W7-AS Contributions to the 24th EPS Conference on Controlled Fusion and Plasma Physics
(9 - 13 June 1997, Berchtesgaden, Germany)

Joint Conference of 11th Internat. Stellarator Conference & 8th Internat. Toki Conference (ITC-8)
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Analysis of W7-AS Mirnov data using SVD and correlation techniques

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Introduction: The modular stellarator W7-AS is equipped with three poloidal arrays of Mirnov probes, two with eight, one with sixteen coils measuring the rate of change of the poloidal magnetic field. Data are acquired at a rate of 333kHz and 250kHz, respectively. Analysis comprises SVD, correlation and Fourier techniques. The aim is to identify MHD instabilities, the influence of which on stellarator confinement is often unclear.

Fig. 1: Schematic of mode analysis: circles indicate input data/parameters, boxes represent entities of different decompositions. Triangles describe ways to analyse the decomposition entities, e.g. power spectral density or phase analysis of SVD chronos. Output data are represented by 'eggs'. Further explanations see text.

Data Analysis: Raw data are FFT'd and band-passed to eliminate parasitic signals such as pick-up from the thyristors of W7-AS' power supply.

Singular value decomposition (SVD) or biorthogonal decomposition (BD) [1] splits a matrix $X (M \times N, M > N)$, containing $N$ time series of $M$ samples of $N$ probes into three matrices: $U(M \times N)$, $V(N \times N)$, and a diagonal $S$.

$$X = U \ast S \ast V^\dagger$$  \hspace{1cm} (1)
The columns of $V$ are spatial “eigenvectors” or topos (figs. 2c, d), the columns of $U$ are temporal “eigenvectors” or chronos (figs. 2e, f), ordered with respect to variance (“importance”) which is reflected by the monotonously decreasing singular values contained in $S$ (fig. 2b).

If two of the singular values are approximately equal, the corresponding topos and chronos describe one single, but rotating perturbation. In that case, the topos reveal the dominant $m$ (figs. 2c, d), as does the relative phase of the two neighbouring topos (figs. 2g, h), which is analysed in the same manner as phases obtained from FFT [3] The relative phase of the corresponding chronos yields a time-resolved frequency [1] (figs. 2i, j).

In our example, SVD components with $k \geq 4$ will contain mostly noise (see the discussion in [1]). A non-Fourier noise-filtering may thus be obtained if the higher order singular values are set to zero when reconstructing $X$ according to eq. (1).

The normalised cross correlation function (NCC) [2] is calculated from SVD-filtered data using selected topos and chronos $k_1 \ldots k_2$. Probe 1 (outboard midplane, fig. 3a) serves as a reference channel and $j = 1 \ldots N$:

$$c_{ij}(p) = \frac{1}{2M + 1} \sum_{i=-M}^{M} x_1(i)x_j(i-p)$$

(2)

The correlation diagram, i.e. a plot of $c_{ij}(p)$ vs. the time lag $p$ and the probe position reveals the frequency, the sense of rotation and the dominant poloidal harmonic $m$ (figs. 2k, l, o, p).

For the frequencies of maximum power spectral density yielded by the coherence spectra and the spectra of the relevant chronos, $m$ is obtained by FFT phase analysis** [3].

A comparison with calculated phases or correlation diagrams which are obtained using the assumed $m$ and the straight field line angle $\theta^*$ from the vacuum configuration is often extremely helpful. Such calculations are shown in figs. 2g, h, o, p as dashed curves and in figs. 2k, l as overlayed solid curves.

Result of the analysis example: We state the presence of two different modes. The first one rotating in the ion diamagnetic drift direction (fig. 2k) is possibly a beam-driven GAE mode. The power spectral density of its two chronos is very sharply peaked at $37 kHz$. The poloidal structure appears to be more consistent with $m = 4$ (see figs 2g, h).

*An FFT band-pass filter may be used additionally.

**This is not shown here, but looks essentially the same as figs. 2g, h
k, o), although the correlation diagram exhibits some distortions (fig. 2k). The second mode with \( m = 3 \) rotates in the electron diamagnetic drift direction (fig. 2l). The power spectrum reveals a rather broad peak about 25kHz.

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**Fig. 2:** Mironov analysis of W7-AS discharge # 39181, a purely NBI heated D plasma: \( P_{NBI} = 0.5MW \), line density \( \approx 4 \cdot 10^{19}m^{-2} \), central electron temperature \( \approx 700eV \), extraordinary good energy confinement of \( \tau_E \approx 40ms \) [4] in spite of comparatively strong MHD activity. Data were taken with the 16-probe array (compare fig. 3a) at a sampling rate of 250kHz. 1000 samples centered about 370ms were selected. Further explanations see text.
Reliability: Any difficulties arising from noisy data are presumably circumvented by the use of the analysis procedure described above. Persisting problems are essentially caused by geometry, i.e. by a too high $m$, a too great distance between the probes and the source of the magnetic signal or by inadequate poloidal positions of the probes.

In order to quantify these limitations, time series of simulated Mirnov signals have been produced for the 16-probe array (fig. 3a). Based on theoretical equilibria, modes with $m$—numbers between 2 and 6 at different radial positions $2cm \leq r_{eff} \leq 15cm$ have been simulated. The degree of agreement of the analysis result with the actual $m$ is displayed in fig. 3b. Identification of modes with $m \geq 6$ is impossible, due to the short decay lengths of high multipoles. It is feasible to identify $m = 5$ if the perturbation is located at $r_{eff} \geq 12cm$. If the mode is located inside $r_{eff} \approx 10cm$ which is frequently the case for GAE modes [5,6], $m = 3$ and $m = 4$ may already be hard to distinguish.

Fig. 3a: array of 16 $B_0$-probes (MIR-1). Contours indicate flux surfaces of a sample vacuum configuration. The dashed contour corresponds to $r_{eff} \approx 17cm$.

Fig. 3b: fraction of correctly determined poloidal harmonics $m$ as a function of $r_{eff}$ and $m$. Contours represent levels of 10% to 90% agreement.

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The role of the radial electric field and plasma rotation for the W7-AS stellarator confinement

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Experimental principles and neoclassical calculations

In the advanced stellarator W7-AS [1], the toroidal and poloidal plasma rotation velocity is measured by charge exchange recombination spectroscopy CXRS [2] on impurity ions, mainly on Helium. For that sake, two different spectroscopic systems are available [3], the first one with observation chords in one poloidal plane, the second one with viewing lines with an angle of roughly 50 degrees to the magnetic field lines. The CX light intensity is excited in the beam of a modulated diagnostic neutral beam injector.

From the spectral CX line intensity, Doppler line broadening and line shift, the impurity density \( n_{i+1}(r) \), the impurity ion temperature \( T_1(r) \) and the rotation velocities \( V_\phi \) and \( V_\theta \) are determined, respectively. Here, \( \phi \) and \( \theta \) are the toroidal and poloidal angle co-ordinates.

The CX reaction changes the impurity charge state \( Z \) from \( I+1 \) to \( I \). It is assumed that \( T_1(r) = T_1(r) = T_{i+1}(r) \). The radial electric field profile \( E_r(r) \) is calculated from the simplified radial force balance equation [4]:

\[
E_r = \frac{\partial(n_{i+1}(r) \cdot T_{i+1}(r))}{\partial r} + \frac{1}{e Z_{i+1} n_{i+1}(r)} \cdot \frac{T_1(r)}{e Z_1} \cdot \frac{\partial \zeta_1}{\partial r} + (B_0 V_\phi - B_\phi V_\theta)
\]  

(1)

Here, \( B \) stands for the magnetic field, \( e \) for the elementary charge, \( \zeta_1 \) is the excitation probability for the CX spectral line under consideration. The lifetime \( \tau_1 \) of the excited state and the ion gyro frequency \( \omega_1 \) are taken into account for the factor \( \gamma \):

\[
\gamma = \frac{(\omega_1 \tau_1)^2}{1 + (\omega_1 \tau_1)^2}
\]

(2)

In eq. (1), the first term on the right hand side gives the impurity ion pressure gradient contribution to \( E_r \). The second term arises from finite lifetime of the excited electronic state during the ion gyro motion after CX, for the case that the spectroscopic line of sight \( \bar{I} \) is perpendicular to the magnetic field, and a gradient of the excitation probability exists, perpendicular to \( \bar{I} \) and \( B \). In W7-AS this direction is parallel to the minor radius \( r \), therefore the derivatives are with respect to \( r \). Eq. (1) holds for CX approximately for \( \gamma \ll 1 \), as this is given in our case for \( \text{He}^+ \), and therefore that second term is small. The third term is the \( \bar{J} \times \bar{B} \) force. In stellarators, the toroidal and poloidal rotation are de-coupled, in contrast to tokamaks [14]. It is found experimentally for W7-AS that the main contribution to \( E_r \) comes from the poloidal rotation, the toroidal rotation is strongly damped because of
the missing axi-symmetry of the magnetic stellarator field. A fast toroidal plasma motion can be provoked only by non balanced neutral beam heating NBI, with counteracting toroidal viscosity which is found to be in good agreement to neoclassical calculations [9]. But even then its contribution to \( E_r \) is negligible (typically < 5%), because \( B_\theta << B_\phi \).

In non-axisymmetric devices like stellarators a strong dependence of transport from \( E_r \) is expected from neoclassical theory. The formation of \( E_r \) is determined by the radial particle fluxes \( \Gamma \) from the ambipolarity condition:

\[
\Gamma_e(r, E_r) + \Gamma_i(r, E_r) + Z_i \Gamma_i(r, E_r) = 0
\]  

(3)

Especially in the long mean free path regime LMFP, the differences between tokamaks and stellarators are essential, because in LMFP the ripple of the magnetic field along the field lines produces large populations of trapped particles in stellarators, with different trapping types (toroidally, helically or combinations) with an enhanced radial drift. With respect to the development of a stellarator-reactor this unfavourable effect has to be minimized, as this is done by the so-called optimization concept for the advanced stellarators W7-AS and W7-X [12]. These drifts are also strongly reduced by the ambipolar \( E_r \), an effect which is for the ions much more pronounced than for the electrons.

The neoclassical definitions for the particle fluxes \( \Gamma \) and heat fluxes \( Q \) are given by:

\[
\Gamma_\alpha = -n_\alpha \left\{ D_{11}^\alpha \left( \frac{n'_\alpha}{n_\alpha} - \frac{q_\alpha E_r}{T_\alpha} \right) + D_{12}^\alpha T'_\alpha \frac{T_\alpha}{T_\alpha} \right\}
\]

(4)

\[
Q_\alpha = -n_\alpha T_\alpha \left\{ D_{21}^\alpha \left( \frac{n'_\alpha}{n_\alpha} - \frac{q_\alpha E_r}{T_\alpha} \right) + D_{22}^\alpha T'_\alpha \frac{T_\alpha}{T_\alpha} \right\}
\]

The subscript \( \alpha \) stands for ions or electrons, respectively, \( q \) for their electric charge, \( D \) are the coefficients of the transport matrix. For the definition of \( \Gamma \), a possible Ware's pinch term proportional to \( D_{13}^\alpha \) is neglected. Besides the linear dependence of \( \Gamma \) and \( Q \) from \( E_r \), also the \( D \)'s are functions of \( E_r \), for the ions much stronger than for the electrons. For \( E_r \), an odd number of solutions is expected from neoclassical theory [13]. Two of them should be stable, the others are unstable.

For the neoclassical calculation of \( E_r \) the numerical DKES code [5] is used. Starting from the specific magnetic configuration, represented by Fourier modes of the magnetic field, DKES calculates the mono-energetic transport coefficients by solution of the drift-kinetic equation, as function of the effective minor radius, the collisionality and \( E_r \). By energy convolution, the \( 3 \times 3 \) thermal transport matrix \( D_{ij} \) is then obtained, which is used to calculate the radial particle fluxes following eq. (4). The solution of the ambipolarity condition (3) provides then the solutions for the neoclassical \( E_r \), which are finally compared to the measured one.
**E\(_r\) and transport in W7-AS**

For the case of low electron collisionality (low electron density \(n_e(0)\) below \(3 \cdot 10^{19} \text{ m}^{-3}\), high electron temperature \(T_e(0)\) above 2 keV) only one strong positive solution is expected in W7-AS near the plasma center, the "electron root", with \(E_r = + 400 \text{ V/cm} \). For that particular case, the central electron heat transport is considerably reduced [10], allowing experimentally for maximum \(T_e(0)\) up to 4 keV. This situation can be obtained only in conjunction with rather low \(T_i\), typically below 400 eV.

For the case of higher \(n_e(0)\) (between 4 - 12 \cdot 10^{19} \text{ m}^{-3}\), only one negative solution is expected near the plasma edge, the "ion-root". For that type of discharge, with combined NBI and ECRH heating and high power > 1 MW, maximum negative \(E_r = -1000 \text{ V/cm} \) are obtained in the gradient region which act as a potential barrier, together with maximum \(T_i(0)\) up to 1.5 keV. The global energy confinement time exceeds the prediction from the ISS95 regression database [11] by more than a factor of two.

Multiple field solutions are expected for medium \(n_e(0)\) between 2 - 5 \cdot 10^{19} \text{ m}^{-3} and lower heating power, typically with negative \(E_r\) values ("ion-root") near the plasma edge but smaller values than described above for \(E_r = -100 \text{ to } -400 \text{ V/cm} \), and either strong positive \(E_r\) ("electron-root") near the plasma center for sufficiently high \(T_e > 2 \text{ keV} \), or small positive \(E_r = + 20 \text{ V/cm} \) for lower \(T_e\) (positive "ion-root").

The comparison between the measured and calculated \(E_r\) is a sensitive means to investigate the mutual interference between \(E_r\) and the particle transport in detail [6]. It is found that the measured and the neoclassically calculated \(E_r\) are in general consistent to each other [7]. Thus, the validity of the neoclassical particle transport model for the prediction of \(E_r\) in W7-AS is demonstrated, at least for the central part of the plasma up to \(r < 0.7-a\). The most striking feature appears when investigating the impact of \(E_r\) on the shape of impurity density profiles because of the higher \(Z\). This point is of particular interest for a future stellarator reactor, because impurity accumulation has to be prevented, and efficient Helium exhaust is desired. For the calculation of the impurity density profiles, the SITAR code [8] is used which employs a tokamak axisymmetric magnetic field model. The flux ansatz \(\Gamma_i = -D_i \cdot n_i^\prime - V_i \cdot n_i\) with diffusion and convection is used in SITAR in the form:

\[
\Gamma_i = \frac{\rho^2}{Z_i \tau} (0.5+q^2) \left( B \frac{\partial n_i}{\partial r} - \frac{n_i}{Z_i n_i} \frac{\partial n_i}{\partial r} - A \frac{n_i}{T_i} \frac{\partial T_i}{\partial r} - \frac{n_i}{T_i} eZ_i E_r \right)
\]

where \(\rho\) stands for the mean path length between two collisions, \(\tau\) is a mean collision time, \(q\) is the tokamak safety factor, \(A\) and \(B\) are collisionality dependent factors [15]. SITAR fulfils internally the ambipolarity condition, thus taking into account the \(E_r\) influence implicitly. The detailed information on the stellarator magnetic field as required for the explicit calculation of the transport coefficients for W7-AS, however, is taken into account only by DKES. For some discharges, which in principle allow for multiple solutions as described above, small positive \(E_r = + 10 \text{ V/cm} \) (positive "ion-root") are predicted by DKES.
in W7-AS near the plasma center. In the gradient region, the "ion-root" is negative with values $E_r = -100 \text{ V/cm}$. The spectroscopic $E_r$ measurement bars in this case are, however, too large to confirm the small central positive $E_r$ from the evaluation following eq. (1) alone, see fig. 1 on the left side. The spectroscopic measurement by CXRS in fact shows a hollow He$^{++}$ density profile for that type of discharge, as plotted in fig. 1 below on the right side (dots and broken line). With the small positive $E_r = +10 \text{ V/cm}$ taken into account in SITAR in eq. (5), the calculated He$^{++}$ density profile also shows a hollow shape. Thus, an outward convection for He$^{++}$ is confirmed, as a consequence of the positive $E_r$.

![Diagrams showing radial electric field and He$^{++}$ density profiles](image)

Fig. 1: Left plot: $E_r$ profile: CXRS measurement (circles), DKES calculation for the ion-root (solid line) and the electron-root (broken line) which is not realized. The ion-root solution is positive for $r < 6 \text{ cm}$. Right plot: He$^{++}$ density profile measured by CXRS (dots), least squares fit of a generalized Gaussian function to the measured points (thin broken line), result of the SITAR transport calculation without the small positive $E_r$ taken into account (solid line), the SITAR result with $E_r$ reproduces the outcome of the Gaussian fit (thick dotted line).

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Structure of the edge fluctuations in the W7-AS stellarator

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1. Introduction. The turbulent fluctuations in the SOL and the outer confinement zone of the W7-AS stellarator are spatially and temporally resolved using Langmuir probe arrays. The main topic is the investigation of the structure of these fluctuations both perpendicular to and along the magnetic field.

Perpendicular to the magnetic field we apply two different techniques to obtain a 2d-representation of the fluctuations’ structure. A right angled probe array, with legs in the radial and the poloidal direction, allows the calculation of a 2d-correlation function under the assumption of poloidal homogeneity of the turbulence. A 2d-representation is also derived using one poloidal probe array and one single probe toroidally separated for conditions where a high correlation along the connecting field line exists as in our earlier experiments [1]. Generally, a 2d-analysis is necessary for the investigation of radial properties of the fluctuations like the radial size or the radial propagation velocity.

Formerly, a very high correlation of about 90% for floating potential and ion saturation current fluctuations over a distance of about 6 m along a magnetic field line [1] was found. In these experiments two toroidally separated Langmuir probes were connected by a field line that stayed on the outside of the torus, where fluctuations are known to be driven unstable, over the whole connection length. In this paper, measurements along a field line that passes the inside, where fluctuations are partly suppressed, are presented. In general the maximum correlation is observed along the wavefront which is nearly parallel to the magnetic field since \( k_\| \) is usually very small. The inclination between wavefront and magnetic field can be observed in our experiments and therefore \( k_\| \) can be calculated.

2. Structure in the radial-poloidal plane. The derivation of a 2d-correlation function from a measurement using an angled probe array is described in [1]. The second method uses an experimental setup with two toroidally separated probes connected by a field line of 6 m connection length. One of them is a single probe, while the other is a poloidally resolving probe array. Since for this configuration a correlation of about 90% was observed along the field line [1] the fluctuations undergo nearly no change along the field. Thus the signal observed by the single probe is nearly the same as it would be in the toroidal plane of the probe array at the point where the projection of the single probe along the field line is situated. This point is used as a reference point and therefore defines the
origin of the 2d-correlation function. During the shot the probe array is moved radially about 6 cm into the plasma. This movement is relatively slow (400 ms) so that the probe movement is negligible within time windows of 10 ms, which are necessary to obtain a statistically stable correlation function. All the correlation functions between the single probe and the probe array that can be calculated for the different radial positions of the array can be combined into a 2d-correlation function depending on time-delay, the radial distance and - due to the poloidal resolution of the array - also on the poloidal distance from the reference point. The experimental flexibility of this method is quite restricted since it needs a specific magnetic field geometry to connect both probes. Nevertheless it is possible to measure 2d - correlation functions that are extended across the LCFS in the radial direction.

![Diagram](image)

**Figure 1:** 2d-correlation function that extends over the LCFS for constant time-delay $\tau=0$. Radial-poloidal plane.

In the case shown in fig. 1 the origin is very close to the LCFS indicated both by the LCFS from a vacuum-field calculation which is typically 1 cm inside the LCFS corrected for finite $\beta$ and by the velocity-shear-layer which is typically 0.5 cm outwards respectively. The fluctuating structure of fig. 1 is clearly extended both into the confinement zone and the SOL.

The 2d correlation analysis in the SOL by use of an angled probe array [1] shows quite similar structures for both floating-potential- and ion-saturation-current-fluctuations. Both
have the well known oblique structure in the radial-poloidal plane as it has been observed before in case of floating-potential-fluctuations [1]. Setting up the probe array such that the radial probe tips measure floating-potential and the poloidal ones ion-saturation-current, the 2d cross-correlation between potential and density fluctuations can be obtained. The maximum correlation is hereby shifted in time and space indicating a phase shift between potential and density which causes transport due to these fluctuations. The obliqueness in the radial-poloidal plane means that the particle flux by fluctuations is not purely radial but that is has a significant poloidal component.

Figure 2: 2d-correlation functions ($\tau = 0$) obtained by an angled probe array for floating-potential (left) and ion-saturation-current (middle). The cross-correlation between floating-potential and ion-saturation-current is shown right.

3. Correlation along field lines passing the torus inside. The experimental setup at W7-AS allows connection of two toroidally separated probes by a magnetic field line that passes the inside of the torus and has a connection length of 32 m. Due to the suppression of fluctuations in the regions of good magnetic curvature and due to the long connection length the correlation observed is only around 40% for both floating-potential and ion-saturation-current. In these experiments the connection between both probes is
only possible within the confinement zone of W7-AS. In this case the correlation is no longer along the computed magnetic field line but there is a shift in the poloidal direction observed. This shift can be explained by a nonzero $k_{||}$ since the correlation is observed along the wavefront which is no longer along the magnetic field, if $k_{||}$ is different from zero.

![Schematic view of the experimental setup for the measurement of $k_{||}$](image)

Figure 3: Schematic view of the experimental setup for the measurement of $k_{||}$.

The experimental results give $k_{||} = 0.0098 \frac{1}{m}$, which corresponds to a parallel wavelength of $\lambda_{||} \approx 320 m$.

4. Conclusion. Analysing the structure of edge fluctuations in the W7-AS stellarator we found that in the radial-poloidal plane these fluctuations look like deformed convective cells with a longer half width in the poloidal than in the radial direction. As mentionnd earlier [1], these fluctuations are oblique in the radial-poloidal plane. The potential-density cross-correlation showed in addition that the fluctuation-induced particle flux is not purely in the radial direction but that is has a significant poloidal component. Along the magnetic field we observed very long wavelengths of the fluctuations in the order of some 300 meters. Good curvature does not suppress fluctuations completely.

5. References

Confinement in W7-AS and the role of radial electric field and magnetic shear


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Abstract

Improved neoclassical electron confinement in the centre of low density ECRH plasmas has been observed in the presence of a strong positive radial electric field, which resembles the electron root solution of the neoclassical ambipolarity condition but is obviously driven by the loss of ECRH-generated suprathermal electrons. At higher densities and with NBI heating, a high confinement regime substantially above the ISS95-scaling and different from the H-mode is established with a strongly sheared negative radial electric field at the boundary. The application of plasma current induced magnetic shear reveals that confinement in W7-AS is essentially determined by perturbations at high order rational surfaces. For optimum confinement, these resonances have either to be avoided in the boundary region or magnetic shear must be sufficiently large. Independent of its sign, magnetic shear can reduce transport which is enhanced in the presence of such resonances to the neoclassical level.

1. Introduction

Shear in the magnetic field and flow field, the latter driven by the $E \times B$-drift, have each attracted much attention in the recent discussion on strategies for how to reduce anomalous transport. Both are expected to reduce anomalous transport by radial decorrelation of turbulent structures. For example, advanced tokamak concepts rely on reversed magnetic shear (in the tokamak sense, $dq/dr < 0$), which provides ballooning stability and access to both high $\beta$ and high bootstrap current fraction [1]. With reversed shear transport can be reduced to the neoclassical level [2]. (In the following we define magnetic shear as the radial variation $\partial \psi/\partial r$ of the flux surface averaged rotational transform $\tau$ rather than of the safety factor $q = 1/\psi$). Sheared
poloidal flow for example is supposed to be fundamental to the development of the edge transport barrier in the L-H transition of both tokamak [3] and stellarator [4].

In stellarators, the radial electric field $E_r$ plays an additional important role: it may help to overcome the enhanced neoclassical transport losses in the reactor-relevant long-mean-free-path (LMFP) low collisionality regime. In this so called $1/v$-regime (if $E_r = 0$) the transport coefficients scale with temperature $\sim T^{\eta/2}$, typically exceed the anomalous ones, and limit access to higher temperatures. The poloidal $E\times B$-drift in a sufficiently strong radial electric field averages the radial $VB$-drift motion and effectively reduces transport [5]. The radial electric field adjusts itself by the ambipolarity of the particle fluxes. Different to tokamaks, multiple roots of the ambipolarity condition are predicted for stellarators due to the nonlinear dependence of the neoclassical transport coefficients on $E_r$ [6, 7]. Up to recently, only the so called ion root, with usually moderate negative or small positive $E_r$, had been observed under stationary conditions. The electron root with strong positive $E_r$ appears more attractive: it can theoretically be sufficiently strong to reduce the electron LMFP transport, while the electric field is directed such as to prevent impurity accumulation by driving impurity ions outward.

Magnetic shear sensitively determines the quality of flux surfaces and their susceptibility to magnetic perturbations. Owing to the plasma current density profile, shear is inherent to tokamaks; in stellarators, it is a design parameter since the basic rotational transform is provided externally, although the vacuum field properties are modified by pressure driven plasma currents. In low magnetic shear devices, e.g. Wendelstein stellarators, magnetic islands related to resonant perturbations at low order rational values of the rotational transform should be excluded from the region of strong plasma gradients, where their effect on confinement is most detrimental. At moderate/high shear, e.g. Torsatrons/Heliotrons, the radial island width is kept small by sufficient shear. However, in the case of strong shear a large number of radially overlapping islands at high order rationals may exist, the overlap parameter scaling with $(dt/dr)^{1/2}$ [8]. With low shear $(|\Delta \theta|/\theta) \sim |t_a - t_c|/\theta < 0.04$ for the vacuum field of Wendelstein 7-AS (W7-AS) [9]), optimum confinement is usually found close to the narrow "resonance-free" zones [10] which exist in the direct vicinity of major resonances ($\ell = 1/3, 1/2, \ldots$). At moderate $\beta$, pressure-driven currents (bootstrap and, in W7-AS reduced but still significant, Pfirsch-Schlüetter (PS) currents) modify the rotational transform, introduce moderate shear and may drive a discharge into perturbed configurations. Therefore, the boundary value $t_a$ of the rotational transform is usually controlled by inductive compensation of the pressure-driven net toroidal plasma current. With moderate shear and control of $t_a$, good confinement is found even in the presence of a major resonance [11]. Finally, at high $\beta$, the sensitivity of confinement to the major resonances is strongly reduced due to pressure-induced shear [12]. For the maximum $\langle \beta \rangle = 1.8 \%$ achieved so far, $|\Delta \theta_0| = 0.35$ is estimated [13]. For comparison, the vacuum field of the optimized Wendelstein 7-X (W7-X) stellarator has moderate magnetic shear $(|\Delta \theta_0| \approx 0.1)$. As additionally PS and bootstrap current are both strongly reduced, the vacuum rotational transform should be little affected by plasma pressure [14].
This paper presents, after a short overview of the device (Section 2), recent results on the influence of both radial electric field (section 3) and magnetic shear (section 4) on confinement in W7-AS. The high flexibility of the magnetic configuration can be used to control to a certain extent the Fourier spectrum of $B$ and thus to influence neoclassical transport, which, in turn, affects the radial electric field by the ambipolarity of the particle fluxes and vice versa. The low shear vacuum rotational transform provides the possibility to study the effect on confinement of both positive and negative high shear by external current drive.

In the presence of a strong positive central $E_r$, which is indicative of the neoclassical electron root, central electron temperatures exceeding $T_e = 4$ keV have been achieved in low density ECRH discharges. These temperatures are the highest obtained in W7-AS so far (see also Ref. 15). At higher densities, using NBI heating alone or combined with ECRH, high ion temperatures $T_i = T_e \leq 1.5$ keV are achieved under low recycling conditions. These discharges show excellent confinement, which substantially exceeds the ISS95-scaling [16]. They are associated with a strongly sheared negative $E$, forming an edge transport barrier. They do, however, not show the characteristic features of an H-mode (see also Ref. 17).

With respect to magnetic shear, an increasing amount of shear has been applied by inductive current drive over a wide range of boundary rotational transform values in ECR heated discharges at moderate $\beta$ (< 1% in the centre). In the presence of high order rational values of the rotational transform, confinement is found to depend strongly on the magnetic shear. The effect of the sign of the shear will be addressed.

2. Overview of Wendelstein 7-AS

Despite its stability and inherent steady state capability the classical stellarator concept may be limited by some major drawbacks: (i) a rather low equilibrium $\beta$-limit because of the Shafranov shift, (ii) sensitivity of the flux surfaces to perturbations which may lead to island formation or even ergodization, and (iii) poor neoclassical confinement in the LMFP collisionality regime. To overcome these difficulties simultaneously has been the subject of a long stellarator optimization which has resulted in the design of the W7-X stellarator now under construction [14]. The present W7-AS advanced stellarator is partly optimized: the PS currents are reduced by a factor of two compared to a classical equivalent, leading to improved equilibrium properties [18], MHD-stability [19], and neoclassical plateau transport [20].

W7-AS (5 toroidal field periods, pentagon-shaped magnetic axis, $R = 2$ m, $a = 0.18$ m, $B_0 \leq 3$ T [21]) is designed for high experimental flexibility. The confining magnetic field is produced by 45 non-planar, poloidally closed modular coils providing alone a rotational transform of 0.4, which can be changed within the range $0.25 < \tau < 0.7$ by the use of 10 planar toroidal field coils. Vertical field coils can be used to shift the plasma column radially. The magnetic mirror ratio, i.e. the toroidal variation of $B$, can be varied by feeding the currents in the modular coils positioned at the corners of the pentagon separately. Here lies the region of
strong toroidal curvature and large $\nabla B$-drift. The mirror ratio in particular affects the effective magnetic ripple and thus the neoclassical LMFP transport. A loop voltage can be applied for inductive current drive, either to compensate the bootstrap current or to study, e.g., the effect of current-induced magnetic shear.

Plasma heating is presently achieved by up to 0.8 MW ECRH at 140 GHz (2nd harmonic X-mode at 2.5 T), 0.4 MW ECRH at 70 GHz (1st (2nd) harmonic O(X)-mode at 2.5 T (1.25 T)), and 3.2 MW NBI at up to 50 keV (power through port). Successful ICR-heating has been demonstrated recently [22]. ECR-heating is located in the region of strong toroidal curvature (one of the pentagon corners) so that heating of "ripple-trapped" particles can be selected by the specific choice of the magnetic configuration.

3. The role of the radial electric field

At low collisionality, particles trapped in local helical mirrors may be radially lost by the $\nabla B$-drift which leads, in the absence of a radial electric field $E_r$, to the formation of the unfavourable $1/\nu$-regime. In the presence of a radial electric field, the particle orbits are additionally subjected to the poloidal $E \times B$-drift which reduces the excursions of the trapped particles from the flux surfaces. For moderate $E_r$, a $\nu^{\alpha}$-regime develops. With sufficiently strong $E_r$, the poloidal drift motion dominates over the $\nabla B$-drift resulting in the $\nu$-regime with dramatically reduced transport coefficients. The beneficial effect of the radial electric field scales with $E_r/\nu_{th}$ and is accordingly much stronger for ions than for electrons.

The radial electric field is self-consistently calculated for each flux surface from the local ambipolarity $\Gamma_e = \Gamma_i$ of the neoclassical electron and bulk ion fluxes (impurity ions not being considered) where

$$\Gamma_a = -n_a(D_{11} \partial / \partial n_a - q_a E_r / T_a) + D_{12} \partial \nabla T_a / T_a).$$

$D_{11}$ is the particle diffusion coefficient, $D_{12}$ accounts for temperature gradient driven particle fluxes, and $q_a$ is the charge of particle species $a$. Additional non-ambipolar fluxes are neglected. Because the LMFP transport coefficients depend on $E_r$, multiple solutions of the ambipolarity condition may exist, even with the ions in the plateau regime. For example, at high electron temperatures $T_e >> T_i$, as is typical in the central region of low density ECRH discharges with electron heating and low collisional energy transfer to the ions, a strong positive solution - the so called **electron root** - is predicted, roughly given by $\Gamma_e = 0$. For these conditions the $\partial \nabla T_e$-driven flux dominates over the $\nabla n_e$-driven one, i.e. $eE_r < (D_{12} / D_{11}) \nabla T_i > 0$. (A weakly positive $E_r$ (ion root) may exist as well for $T_e >> T_i$). For the condition $T_e = T_i$, as is found at higher density with NBI heating or generally at the plasma boundary, a smaller and negative $E_r$ - the so called **ion root** - is expected, roughly given by $\Gamma_i = 0$, i.e. $q_i E_r = T_i \nabla n_i / n_i + (D_{12} / D_{11}) \nabla T_i < 0$. For intermediate regions, both solutions may be possible with an additional unstable root. The radial transition from the inner electron to the outer ion root, which
corresponds to a poloidal rotation shear layer, can be predicted by means of a thermodynamic energy principle [20].

Experimentally, the radial electric field is determined from the poloidal plasma rotation by both passive as well as active charge exchange recombination spectroscopy (CXRS) [23]. Active CXRS provides radial resolution by use of a diagnostic neutral beam. In the following, a detailed transport analysis compares for various discharge scenarios the measured $E_r$-profiles and experimental heat conductivities (from power balance analysis) with the neoclassical predictions from the DKES code [24, 20], which uses the measured $n_e$, $T_e$, and $T_i$-profiles as input (from Thomson-scattering, electron cyclotron emission (ECE), active charge exchange neutral particle analysis (CXA) and CXRS). To determine the experimental particle and energy fluxes for the respective balance equations, particle and energy sources have been used as follows: recycling and gas puffing particle sources as well as thermal CX-losses from the EIRENE code [25], NBI particle and energy deposition sources as well as fast CX-losses from the FAFNER code [26], ECRH deposition from ray-tracing. Radiation losses are modeled according to bolometry data.

3.1. The neoclassical electron root feature

Very recently, central electron temperatures of $T_e \approx 4$ keV have been observed in low density ($n_e \leq 2.0 \times 10^{19}$ m$^{-3}$) discharges heated by up to 800 kW ECRH at 140 GHz. Figure 1 shows the radial profiles for this type of discharge. The ions are energetically decoupled from the electrons and have flat temperature profiles at low values of $T_i = 0.35$ keV. In the central region ($r < 6$ cm), an additional peaking of the $T_e$-profile is observed. A strong positive radial electric field up to $E_r = 600$ V/cm is measured in this range. In the outer region $E_r$ is only weakly positive and becomes negative at the edge. In addition to the measured $E_r$, the neoclassically predicted value based on thermal electron fluxes is also given. At $r < 7$ cm, only a positive electron root is predicted. At intermediate radii ($7$ cm $< r < 13$ cm), both the electron root and a weakly positive ion root can in principle exist. The third unstable root is shown dotted. From the energy principle, only the ion root should be realized, with the transition layer from the inner electron to the outer ion root being located just at the onset of the ion root at $r = 7$ cm (the solid line in Fig. 1). Finally, at the edge ($r > 13$ cm), only a negative ion root exists. The experimental $E_r$-profile is in reasonable agreement with these predictions. In the centre, the experimental electron heat conductivity $\chi_e$ from power balance analysis agrees within a factor of about 2 with the neoclassical value calculated for the electron root. For comparison, with $E_r = 0$ the neoclassical estimate by far exceeds the experimental value. At outer radii ($r > 12$ cm), i.e. in the region of strong density gradient and small $E_r$, energy transport becomes strongly anomalous, by far exceeding the neoclassical prediction.
Fig. 1: Profiles of $T_e$ (dots: from Thomson scattering, diamonds: from ECE) and $T_i$, $n_e$, $E_r$ (diamonds: from CXRS) and $x_e$ (dashed dotted: from power balance) for a 770 kW ECRH-discharge (140 GHz) at $B = 2.5$ T. For $E_r$ and $x_e$ the neoclassical predictions from the various roots of the ambipolarity condition are also given (solid line for $E_r$; predicted transition from electron to ion root).

So far the experimental findings seem to be consistent with the prediction of a central electron root driven by thermal neoclassical particle fluxes. However, why does this feature not appear in other cases where it is also neoclassically predicted? The fundamental difference from previous experiments with otherwise comparable parameters is the heating method: the present experiments use 2nd-harmonic X-mode ECRH at 140 GHz, whereas previous experiments used 1st-harmonic O-mode at 70 GHz. O- and X-mode absorption are significantly different. In particular for perpendicular launch and $v_{\parallel} = 0$ (deeply trapped particles) there is no O-mode absorption whereas X-mode absorption is maximum. This may result in a more pronounced suprathermal tail in the ripple-trapped electron distribution for X-mode heating [27]. This difference suggests that the strong positive electric field in W7-AS is related to the heating of electrons trapped in a local magnetic mirror at the ECRH launching position. Indeed, in the W7-AS standard configuration, there is a local minimum of $B$ at this position. To clarify this conjecture, experiments have been conducted which systematically change the fraction of trapped suprathermal particles either by collisional detrapping (density scan), variation of the
ECRH power and its radial position, and by changing the depth of the local magnetic mirror at the ECRH position. With X-mode heating the electron root feature should be less pronounced, or even vanish, with increasing density, decreasing power, off-axis heating and reduced magnetic mirror depth.

Fig. 2 (left) shows for the standard configuration and for ECRH powers of 770, 460 and 220 kW that the central peaking of the $T_e$ profile, which is attributed to the existence of the electron root, is lost at the lowest power. In this case the measured $E_r$ (not shown) is small in the centre, i.e. similar to the ion root, whereas the electron root is still predicted from the thermal fluxes. The electron heat conductivities are nearly identical for the higher powers, leading to a significant increase of $T_e$ with heating power. For the profiles in Fig. 2 (right) the magnetic mirror ratio has been reduced by increasing the ratio of currents in the corner coils and the modular field coils from 1.00 to 1.18. The local mirror depth of $B$ at the ECRH position is thereby reduced from $\Delta B/B = 5.6\%$ to 1.2\% and the effective ripple on axis decreases from $\langle \epsilon_r \rangle = 1.4 \%$ to 0.4 \%, i.e. the configuration is improved with respect to neoclassical transport. Nevertheless, the $T_e$-profiles show neither the high central values nor the steep gradients seen in the high ripple case. With increasing heating power, only a slight gain in $T_e$ is observed. This is also reflected in the rather strong increase of the experimental heat conductivity with $T_e$, which is typical for the 1/\nu-regime with small $E_r$. The central electron temperature (from ECE) decays differently for the two configurations after switching off ECRH. With high ripple, it decays along two time-scales: a fast one with a time-constant characteristic for the decay of suprathermal electrons and a slow one along the confinement time-scale. Thus the radial electric field which is responsible for the high central $T_e$ seems to vanish on the same time-scale as the
suprathermal electrons. In the neoclassically improved low ripple case, no electron root feature is found and the central $T_e$ decays on the thermal time-scale, which however is longer, in agreement with the neoclassical prediction. All these observations consistently lead to the conclusion that the strong positive $E_r$ is driven by the $\nabla B$-drift loss of fast trapped electrons produced in the local magnetic field ripple by ECR-heating [15].

3.2 High confinement with strong negative electric fields

High confinement associated with a strongly sheared radial electric field close to the boundary has been achieved at higher density ($5 \times 10^{19} \text{ m}^{-3}$). With moderate power NBI heating (370 kW), energy confinement times up to $\tau_e = 55$ ms, the highest in W7-AS so far, and temperatures $T_e = T_i$ up to 1 keV have been achieved [17]. At higher NBI power (850 kW) and with additional ECR-heating (350 kW), the highest ion temperatures of 1.5 keV are observed [13]. The characteristics of these discharges are similar, in particular with respect to low boundary density and narrow density profiles. With respect to global confinement, both exceed the ISS95 scaling by more than a factor of two. Low recycling obtained by boronization and He glow discharge conditioning is a prerequisite to obtain this regime. Otherwise density control is lost, broad density profiles develop and temperatures remain below 0.5 keV.

Figure 3 illustrates for a discharge heated by 370 kW NBI that the high confinement is established by a gradual improvement rather than a fast transition. The diamagnetic energy, $T_e$ and $T_i$ continuously increase up to 0.4 s at constant heating power and line density. Simultaneously, a reduction in the $H_\alpha$ light indicates improving particle confinement. A slow increase of the impurity content in the core is indicated from $Z_{eff}$ (from soft-x radiation); it continues up to the end of the discharge and is consistent with impurity transport studies which yield a very low impurity diffusivity of

Fig. 3: Waveforms for a high confinement NBI-discharge: diamagnetic energy, line averaged density and gas puff and $Z_{eff}$ from soft-x, $T_e$ from ECE and $T_i$ from CXA, $H_\alpha$ at a limiter, radial electric field from passive CXRS.
≈0.07 m²/s and inward convection of ≈5r/a m/s [28]. A decrease of density fluctuations is indicated by reflectometry data. Along with improving confinement, the value of the negative radial electric field increases at the boundary (from passive CXRS). During this phase, the edge density decreases, leading to rather narrow density profiles. Simultaneously the temperature increases over the whole cross section and broad profiles with strong gradients close to the edge develop. The process of improvement stops abruptly at about 0.4 s.

![Graph showing profiles of ne, Te (Thomson scattering), Ti (CXA and CXRS), Er (CXRS and neoclassical prediction) and χ (one-fluid power balance and neoclassical predictions) for a high confinement discharge. The T_i profile predicted from power balance is shown dashed.](image)

Fig. 4: Profiles of $n_e$, $T_e$ (Thomson scattering), $T_i$ (CXA and CXRS), $E_r$ (CXRS and neoclassical prediction) and $\chi$ (one-fluid power balance and neoclassical predictions) for a high confinement discharge. The $T_i$ profile predicted from power balance is shown dashed.

Profiles of density, temperatures and radial electric field are shown in Fig. 4 for the steady state phase along with results of a power balance analysis. The measured $E_r$ is in reasonable agreement with the neoclassical "ion root" prediction. The power balance analysis was carried out with the same heat diffusivity $\chi = \chi_e = \chi_i$ for electrons and ions. $n_e$ and $T_e$ profiles from Thomson scattering serve as input and $\chi$ as well as the $T_i$ profile follow from power balance analysis. They are compared with measurements and neoclassical predictions from the DKES code. The calculated $T_i$ profile agrees with the experimental one. Although suggested by the one-fluid analysis, it cannot be concluded that $\chi_i$ is definitely below the neoclassical value for the central region, $r < 0.1$ m: a two-fluid power balance shows that using the neoclassical prediction for $\chi_i$ is also consistent with the central $T_i$, although the two-fluid model fails to predict the ion temperature in the gradient region. Here, in the region of strong
shear in the radial electric field, the heat diffusivity is strongly reduced and drops to the neoclassical level of the electrons and below the neoclassical ion diffusivity. Hence this region satisfies the criterion of a transport barrier in which anomalous transport seems to be suppressed.

The fact that confinement improves continuously rather than in a fast transition could point to a causality loop, where anomalous transport is reduced by sheared flow generated by the neoclassical electric field. Reduced transport steepens the gradients, which again increase $|E|$, etc. This loop would stop when anomalous transport becomes suppressed to the neoclassical level. The time constants of this process would be in the order of the confinement time. At present it cannot be completely ruled out that changes in the magnetic shear profile on the resistive time-scale play a role. However, by imposing positive and negative net plasma currents and keeping the total boundary rotational transform fixed ($\tau_s = 0.35$), it was shown that magnetic shear does not seem to be the key parameter.

As compared to the H-mode, the differences of this high confinement regime can be summarized as follows: (i) the improvement of confinement is gradual, not instantaneous, (ii) relaxation phenomena such as ELMs are not observed, (iii) the edge density is low and the density gradient is located well inside the last closed flux surface rather than at the edge, (iv) there is no pedestal in the temperature.

4. The role of magnetic shear

The dependence of energy confinement in W7-AS on the boundary value $\tau_s$ of the rotational transform with optima close to $\tau_s = 1/3$ and $1/2$ cannot be explained for net current-free discharges by the external perturbations of the vacuum field (in particular "natural" 5/m-components arising from the 5-fold toroidal symmetry and 1/3- and 1/2-components from field errors) [9]. Therefore higher order perturbations at $\tau = n/m$ resonances may be important. The impact of such perturbations will depend on the magnetic shear. At moderate $\beta$ and small net currents (i.e. up to the order of the bootstrap current), a mutual dependence of confinement and shear is established by the balance between pressure-driven and inductive currents, resulting in low or moderate shear. Large inductive currents (i.e. significantly larger than the bootstrap current) introduce high shear and provide to a certain extent experimental control of shear.

The effect of shear on energy confinement has been studied for ECR heated plasmas ($\beta_0 < 1\%$). Plasma currents up to $I_p = \pm 30$ kA are inductively driven, corresponding to $q(a) = 5$. The Ohmic heating power is negligible ($< 10$ kW) as compared to 450 kW ECRH power at 140 GHz. The toroidal net current contributes $\Delta\tau_s = (\mu_0 R I_p)/(2\pi a^2 B_0)$ to the boundary transform, which is $0.007/kA$ for the typical parameters $a = 0.15$ m, $R = 2$ m and $B_0 = 2.5$ T (by definition, positive currents increase $\tau$; in particular, the bootstrap current is positive). The plasma aperture was reduced by having the two main limiters moved inward, in order to weaken the effect of the 5/m boundary islands and to quench the occurrence of the H-mode [4].
For fixed limiter aperture, the effective plasma radius slightly changes with current by typically \((\Delta a/a)/I_p = 0.004/kA\), as has been confirmed by equilibrium calculations. The interpretation of the experiments is based on finite-\(\beta\) equilibria from the NEMEC code [29] which account for the Ohmic and bootstrap current density profiles from the DKES code [24].

Fig. 5:  Left: Dependence of the plasma energy on boundary rotational transform for plasma currents of 0 (diamonds), 5 (squares), 15 (triangles) and 25 kA (circles). For three discharges at 5 kA, density control was lost and the measured energy (crossed squares) has been scaled by \(n^{1/2}\) (squares). Right: Dependence of the electron kinetic energy on plasma current for \(\tau_a = 0.51\) and \(\tau_a = 0.42\). \((B = 2.5\ T, P_{ECRH} = 450\ kW, n_e = 4\times10^{19}\ \text{m}^{-3}\)\)

Figure 5 (left) shows for various plasma currents the variation of the plasma energy with \(\tau_a\) in the vicinity of \(\tau_a = 1/2\). Limiter aperture and line-integrated densities are identical for all discharges, i.e. the plasma radius increases from 0.15 m at \(I_p = 0\) to 0.165 m at \(I_p = 25\ kA\) and the central densities decrease from 4 to \(3.5 \times 10^{19}\ \text{m}^{-3}\). The plasma energy has been derived from the diamagnetic loop signal \(W_{\text{dia}}\), which is not fully compensated for magnetic flux from large plasma currents, by adding for each current a correction \(\Delta W_{\text{dia},lp} = W_{\text{dia},0} (W_{e,lp} / W_{e,0}) - W_{\text{dia},lp}\), which is assumed to be independent of \(\tau\) and is determined from \(W_{\text{dia}}\) and the electron kinetic energy \(W_e\) (from profile integration) at \(\tau_a = 0.51\). The electrons dominate the total energy, ion temperatures being below 450 eV. With small net current \((I_p \leq 5\ kA)\), confinement exhibits a pronounced maximum close to \(\tau_a = 1/2\). By further increase of the current (to \(+15\ kA\) the dependence on the rotational transform is smoothed and finally lost (at \(+25\ kA\)). Here, the value of the optimum confinement in the current-free case is reached.
The variation of the electron kinetic energy with plasma current is shown in Fig. 5 (right) at fixed boundary rotational transform. Here, the variation of plasma radius is compensated by proper adjustment of the limiter aperture. At $t_a = 0.51$ (optimum for $I_p = 0$) a strong degradation of confinement is observed at intermediate currents (7.5 kA). In contrast, at $t_a = 0.42$ (degraded for $I_p = 0$) confinement improves continuously with increasing current up to the optimum. This improvement is independent on the sign of the current (the stagnation at low density with high negative currents is attributed to the approach of very small central $t$-values [30].)

4.1. Confinement at moderate shear

Confinement at small net currents has been analysed for $t_a = 0.495$ at $I_p = 0$ (optimum), $t_a = 0.48$ at $I_p = 0$ (degraded), and $t_a = 0.48$ at $I_p = 5$ kA (optimum). In the latter case, the rather small current improves confinement to the optimum. The density profiles are similar for all cases. Degraded confinement is associated with a flat $T_e$ profile at larger radii (Fig. 6). Also given in Fig. 6 are the calculated profiles of the rotational transform, including the vacuum transform, the PS-current contribution, the bootstrap and the inductive plasma current. For reference, the vacuum transform plus PS-contribution is shown for one discharge.

Fig. 6: Experimental profiles of $n_e$ and $T_e$ for $t_a = 0.495$ and $t_a = 0.48$ at $I_p = 0$ and for $t_a = 0.48$ at $I_p = 5$ kA (left). Calculated $t$-profiles for the same discharges (right). Plasma current and bootstrap current are indicated. The vacuum transform plus PS current contribution for 0 kA and $t_a = 0.495$ (dashed-dotted line) and the rational values $n/m, m \leq 30$ (horizontal lines) are given for reference.
Additionally, the rational values $t = n/m$ up to $m = 30$ are indicated. The vicinity of $t = 1/2$ is free from such resonances.

In the low confinement discharge, the $t$-profile is nearly shearless and close to the vacuum transform. Due to the flat $T_e$-profiles at the boundary, both the bootstrap current (+2 kA) and the compensating inductive current (-2 kA) are low, and concentrate in the central region. Over the whole plasma cross section the rotational transform is within the "resonance" region of densely spaced high order rational values $m/n$. With a slight increase of $t_b$ into the "resonance-free" region, optimum confinement is established for $I_p = 0$. The bootstrap current increases to +7 kA with a correspondingly large inductive current of -7 kA. In combination, these currents provide positive shear in the "resonance-region". In the 5 kA discharge, the total current is governed by the bootstrap current (+7 kA) with only a small inductive contribution (-2 kA). The bootstrap current raises the rotational transform into the "resonance-free" region with very low shear down to about half the plasma radius. The isolated major resonance $t = 1/2$ is crossed with negative shear. Common to all discharges is positive shear in the centre, due to the negative inductive current. The analysis demonstrates that confinement in W7-AS essentially depends on magnetic shear and resonant $t$-values in the outer plasma region: for optimum confinement $t(r)$ has to be in the "resonance-free" region or shear has to be sufficiently large in the presence of high order $n/m$-resonances of the rotational transform. In contrast, degraded confinement is caused by low shear in the presence of such resonances.

The general trends in Fig. 5 can be understood from this analysis as well. With a further decrease of $t_b$, the boundary transform remains in the "resonance" region. For $I_p = 0$, shear remains low since bootstrap and inductive current cannot develop at low confinement and vice versa. Confinement improves slightly, when the next widest "resonance-free" zone (close to $t_b = 2/5$) is reached. The net current of 5 kA always tends to produce some negative shear close to the boundary. Therefore the decrease of confinement is weaker. Confinement degradation above the "resonance-free" zone at $t_b > 0.53$ behaves in a complementary fashion: for $I_p = 0$ the decrease is weaker than for $I_p = 5$ kA, because the upper "resonance" region is entered with moderate positive shear, but nearly shearless for 5 kA. However, in this $t$-range the rather strong natural 5/9 perturbation will be important as well.

The strong decrease of confinement when the boundary transform is kept fixed at $t_b = 0.51$ and the current is increased from 5 to 7.5 kA, is probably due to the circumstance that the upper "resonance" region is entered with low shear. This situation is cured by providing increasing negative shear at higher currents: the energy recovers to the optimum.

4.2. Proceeding to high shear

The previous results indicate that the effect of magnetic shear on local transport can at best be assessed at $t_b = 0.42$, being free from effects related to the "resonance-free" zones. Fig. 7 shows the profiles of electron temperature, heat conductivity and rotational transform for
discharges with plasma currents of 0, ±10, and +25 kA. With increasing shear, a continuous steepening of the $T_e$-gradient is observed at the boundary independent of the sign of the shear. In this region (8 to 14 cm), $\chi_e$ decreases by a factor of up to 5. As compared to neoclassical transport, the experimental $\chi_e$ is anomalous over the whole plasma cross-section for $I_p = 0$, neoclassical in the very center for ±10 kA, and neoclassical up to $r/a \leq 0.7$ for +25 kA. Only at the very plasma edge does the confinement remain strongly anomalous. Thus, with increasing shear, the region dominated by neoclassical transport continuously expands towards the boundary due to the increase of neoclassical transport with temperature and a simultaneous strong reduction of anomalous transport.

These results agree with those of a previous study [30] (the low density points in Fig. 5, $T_e \leq 4$ keV). There, $\chi_e$ in the outer plasma region was decreased by a factor of up to 2 with current-induced shear. These high temperature discharges at lower density showed for the highest currents (±30 kA) neoclassical confinement even up to $r/a \leq 0.8$. At the higher density and thus lower temperature of the present study, neoclassical transport is significantly smaller, and the reduction of anomalous transport with increasing shear becomes more pronounced.

Fig. 7: Profiles of the rotational transform, $T_e$ and $\chi_e$ for discharges at $r_a = 0.42$ and plasma currents of 0 kA (dotted), +10 kA (dashed-dotted), -10 kA (dashed) and +25 kA (solid lines). The $\tau$-profile of W7-X is given for reference.
Summary

The radial electric field in W7-AS plays an important role with respect to both neoclassical and anomalous transport. High central electron confinement is achieved by a reduction of neoclassical LMFP transport in a strong positive $E_r$, which is similar to the neoclassical electron root. However, the observed field seems to be driven by the loss of trapped suprathermal electrons produced with X-mode ECR-heating rather than by thermal fluxes. So far it is not conclusively understood why the thermal fluxes do not suffice to sustain the electron root and an additional drive is required. The predicted expulsion of impurities in the presence of a strong positive $E$, remains to be confirmed. The electron root feature in W7-AS is similar to the high ion confinement in W7-A, which was due to a reduction of the ion plateau transport by a strong negative $E_r$, driven by orbit losses of the perpendicular NBI [31, 32].

High confinement with tangential NBI in W7-AS is related to large shear of the plasma flow in the boundary. The observed strong negative radial electric field is consistent with the neoclassical ion root derived from thermal fluxes. A boundary transport barrier with reduced anomalous transport develops, which is distinctly different from the H-mode by the long timescale for improvement, the absence of ELMs and a very low edge density.

In the standard mode of net current-free operation, optimum confinement in W7-AS is established with the boundary value of the rotational transform in the "resonance-free" zones close to the low order rationals 1/2 and 1/3. With $\tau_a$ in the region of densely spaced high order rationals, confinement degrades significantly. So far it is not clear which specific perturbations at the high order rational surfaces are responsible for the enhanced transport. Increasing the magnetic shear for such degraded situations by an inductively driven plasma current continuously improves confinement back towards the optimum by a reduction of anomalous transport. The effect is independent of the sign of the shear.
References

Extension and optimization of lithium beam diagnostic methods

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Introduction
Lithium beam diagnostics are a multi-faceted technique for investigating fusion edge plasmas. While the determination of electron densities by lithium impact excitation spectroscopy (Li-IXS) has already reached a satisfying standard on both large fusion experiments at IPP Garching [1,2], a neutral lithium beam can also be used to determine local concentrations as well as temperatures of impurity ions by charge exchange spectroscopy (Li-CXS) [3,4]. In order to achieve simultaneous Li-IXS and Li-CXS measurements, both existing IPP setups for electron density measurements have been extended. First results prove the feasibility of Li-CXS in W7-AS plasmas.

In order to check the modelling of the lithium beam we have also investigated the population of higher Li\textsuperscript{+} energy levels.

Experimental setup on W7-AS
The existing lithium beam diagnostic on the stellarator W7-AS has been extended by a 14 channel CXS observation system. Two achromats (Ω/4π ≈ 2.9\times10^{-4}sr) are used to image the light onto 14 bundles of quartz fibers, each bundle consisting of 8 single fibers (400μm). The detection region covers a radial range of about 13cm from the plasma center to the last closed flux surface of standard plasma configurations. For these investigations the bundles are coupled one by one to the entrance slit of a monochromator (ACTON, Czerny-Turner, f=0.75m). A two-dimensional detector (Proscan CCD camera, 512x512 pixels, each 19x19μm\textsuperscript{2}) is directly connected to the monochromator exit. The spectral resolution is up to 0.018nm/pixel using a 1800g/mm holographic grating. An additional system of photomultipliers equipped with interference filters (λ=529.0nm) has been introduced, which can be coupled to the same light guides for simultaneous measurements at 14 radial locations. We used the improved extraction geometry of the ion source [1,2] to inject a beam (I_{Li}=53mA) in the energy range of 20-66keV.

Experimental setup on ASDEX Upgrade
On the ASDEX Upgrade tokamak the lithium beam diagnostic has a new location 33cm above midplane. It is equipped with a completely rebuilt 35 channel electron density optics, a 16 channel charge exchange optics and a 3 channel neutral density monitoring system (fig. 1). At
the new position of the ion gun, which is similar to the one on W7-AS, the magnetic field is weaker. There is also additional μ-metal shielding around the ion beam. Thus, we expect magnetic field effects on the lithium beam (E=35keV, I=1.2mA) to be reduced considerably.

Electron density measurement setup
In the former optical system each spatial channel consisted of three 400μm fibers. The light collection/transmission efficiency varied drastically from fiber to fiber and the system was sensitive to the beam position [5]. To overcome these difficulties, a new head for the fiberguides was built with each spatial channel corresponding to a bundle of 35 quartz fibers (100/130μm). The light from each bundle is coupled to a 48m long, 600μm monofiber (N.A.=0.37) that carries the light to the detection units. About seven percent of the incoming light can be detected this way. A new BK7 optical system -a wedge and two lenses- was built to gather light into 35 channels from about 16cm along the beam path. The observation region can be radially shifted by a couple of centimeters, allowing it to be adjusted to different plasma scenarios. Each detection unit is equipped with an interference filter (T>40%, FWHM≤0.7nm) followed by a photomultiplier. We expect 5 to 10 times more signal with the new detection system which will permit increased temporal resolution and investigations of density fluctuations outside the separatrix.

CXS measurement setup
Two quartz lenses (Ω/4π = 1.5x10^{-3}sr) focus the light onto an array of 16 quartz fibers (400μm) enabling the radial distribution of the emitted light in a 16cm region around the separatrix to be measured with 1cm spatial resolution. The fibers are coupled to the entrance slit of the same CXS spectroscopic system as described for W7-AS. Temporal resolution is limited by the readout time of the CCD detector for 16 channels (13 ms).

Neutral density measurement setup
Two BK7 lenses (Ω/4π = 1.8x10^{-3}sr) focus the light onto an array of 21 quartz fibers (400μm) that are grouped into three bundles. The light is gathered from a region far from the separatrix where lithium excitation by collision with neutrals is the dominant process [2]. Each channel is equipped with an interference filter (T>40%, FWHM≤0.7nm) to select the Li(2p→2s) line.
Results

Lithium beam composition

Because of highly different cross sections for charge exchange processes, depending on the excitation state of the donor atom, the composition of the lithium beam is an important parameter for evaluating CXS data. We have therefore investigated several LiI spectral lines (2p→2s, 3d→2p, 4s→2p, 4d→2p) in W7-AS discharges. While the measurements of the most relevant line Li(2p→2s) at λ=670.8 nm were performed for calibrating the CXS setup relatively to the Li-IXS photomultiplier setup, all other LiII lines were investigated to check the attenuation model of the lithium beam [6]. Measured intensities of emission from higher levels were found to differ considerably (30-60%) from corresponding theoretical values. We observed no dependence on magnetic field strength and beam energy, whereas changing plasma densities had strong effects on the conformity of experimental and theoretical values with the deviation becoming larger at higher densities.

Inadequate scaling relations for excitation and ionization processes involving protons and impurity ions in the underlying database [7] have been identified as the major reason for these disagreements. These are now being recalculated by more advanced means. First results for the Li(3l→2l) transitions show a significant decrease of theoretical values of Li(3d)/Li(2p) ratios (fig.2) when using the improved cross section data. However, since the relative population of the Li(3d) level in the lithium beam is in the range of 1% only, with populations of all other Li(nl) levels (n>2) even smaller, the influences on electron density calculations remain below 10%.

Furthermore, recent simulations have suggested that the population of higher excited states depends on Z_{eff}. Thus, the measurement of only one additional line besides the resonance line offers a possibility to determine an estimate for Z_{eff} - under the assumption of a reasonable radial charge state distribution of the impurities and if the present disagreement for n=3 populations between simulation and experiment further diminishes as a consequence of more accurate cross sections.

Fig.2: Population ratios of the Li(3d) and Li(2p) energy levels. Comparison of experimental data and beam simulation results for a series of similar W7-AS discharges. Plasma center at z_L=26 cm.
CXs investigations

Measurements of spectral lines of impurity ions have proven the feasibility of charge exchange spectroscopy on W7-AS plasmas with the extended setup. The photomultiplier system has been introduced for simultaneous measurements of C6+ concentrations at 14 radial locations, while the spectrometer system has been used for the determination of temperature profiles of C6+ ions (fig. 3).

Fig. 3: a) Radial temperature profiles of C6+ compared to e- [8] and D+[9]. All data were obtained during a series of similar W7-AS discharges. b) Raw photomultiplier signals for concentration evaluation at Reff=12.3 cm.

Injection of a carbon pellet at t=0.8s increases the charge exchange induced signal significantly.Injected lithium current Ili≈2mA at Ei=48V.

Temperature values, which were corrected for line broadening effects such as Zeeman splitting and l-level mixing [10], are similar to proton/deuteron as well as electron temperature profiles. Evaluation of impurity ion concentration values requires extension of the charge exchange cross section database, which is presently being carried out.

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References
Investigation of impurity tracer transport in high density plasmas at the stellarator Wendelstein 7-AS

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Introduction - For the decay time of injected tracer impurities (aluminum laser blow-off) in W7-AS, the electron density is, among others, an important scaling parameter [1], indicating improved confinement of impurities towards high electron density. Previous impurity injection experiments (H2S gas oscillation and Al laser blow-off) [1] at medium density (ne0=2.5-10^{19} m^{-3}) could partially be simulated within the errors by the one-dimensional radiation and transport code SITAR [2], based on neoclassical and Pfirsch-Schlüter transport for axisymmetric devices. At low density, the transport was found to be significantly higher than predicted, whereas at high density (ne0=6.5-10^{19} m^{-3}), the neoclassical fluxes had to be reduced to fit the experimental data. Analysis of further discharges at different electron densities, using a simple transport model (diffusion coefficient D(r)=const, inward velocity v(r)=(r/a)v(r=a), a: plasma radius), supports the trend of decreased diffusion coefficients towards higher density, which cannot be attributed simply to a decrease in Z_{eff}. A similar dependence of D on electron density was already observed in ECF heated Heliotron E plasmas [3] and supposed to be caused by changes in D rather than in the flow velocity, the latter being close to the classical expectations. In order to elucidate the density dependence in W7-AS, discharges with densities varied by a factor of 2 (ne0=3.5/7.10^{19} m^{-3}) are analyzed in more detail (fig.1), together with fluctuation- and MHD-diagnostics and measurement of the radial electric field. Especially in non-axisymmetric devices like stellarators, the latter can play an important role for impurity transport [4], but is not yet included in SITAR.

Simulations with simple assumptions about D(r)- and v(r)-profiles might mask possible local changes in transport. Therefore, the radial transport coefficients were tried to be directly derived from the temporal and radial behaviour of spectral radiation, detected by the SX-camera during the penetration process of injected aluminum by laser blow-off. Indications were expected, whether the modification of the transport coefficients happens in the core plasma or somewhere in the plasma boundary.

Transport analysis - Because of its good radial and temporal resolution, the SX-camera is a proper diagnostic tool at W7-AS for transport investigation. In spite of its energy-integrated information, the use of a 25μm Be-filter in front of the camera offers the possibility to restrict the number of ionization states contributing to the measured intensity and simplifies the reconstruction of the total impurity density profiles. For total impurity density reconstruction during the penetration process of injected aluminum, the radial intensity profile at each time-step was abel-inverted (figs.1b,e) and converted to a total impurity density profile assuming coronal equilibrium (quasi-stationary condition) in a first step. This is considered to be applicable in
Fig.1a-f: Comparison between medium density (left) and high density (right):
(b,e) : Temporal evolution of radiation detected by SX-camera during the first 50ms after aluminium laser blow-off:
(c,f) : Graphical derivation of transport parameters $D$ and $v$ from reconstructed total density profile evolution;
(a,d) : Simulation of SX-radiation with the derived $D$ and $v$.

Fig.1g: Experimentally derived transport parameters $D, v$ for medium density (thin curve) and high density (thick curve) (dotted lines: extrapolations).
low-transport discharges at high electron density. The transport coefficients at a certain radial position can then be derived from the local temporal evolution of the total impurity density profiles n(r,t) [5]. With the ansatz $\Gamma = D \text{ grad}(n) + \nu n$, $D$ and $\nu$ can be determined by fitting a straight line, when plotting the normalized total impurity fluxes $\Gamma/n$ vs. the normalized total density gradients grad$(n)/n$ for all time points at this radial position (figs.1c,f). The flux $\Gamma$ can be estimated from the density profile evolution using the continuity equation $\text{dn}/\text{dt} = \text{div}(\Gamma)$ with restriction to radial regions where external sources and sinks can be neglected.

In cases where the assumption of quasi-stationarity might not hold, the reconstruction with coronal equilibrium can cause errors. Therefore, the radial density profiles for each time step were reconstructed again, now using a reconstruction factor $\beta(r,t) = n(r,t)/P(r,t)$ (P: total local emissivity contributing to SX-camera) obtained from a transport and radiation calculation with SITAR, using the $D$- and $\nu$-values derived in the first step as input. A repeated derivation procedure for $D$ and $\nu$, but now with the corrected density profile evolution as described above provides new transport coefficients (fig.1g) which, in fact, fit the SX-camera better in most cases (fig.1a,d). The accuracy of this iterative method strongly relies on the quality of Abel inversion and atomic data base. Possible errors have to be discussed in this context.

**Results and discussion** - In order to study the density dependence of impurity confinement at W7-AS, ECF-heated discharges ($P_{ECF}=400kW$) at two different densities (figs.1a,b,c: $n_{e0}=3.5 \times 10^{19} \text{ m}^{-3}$; figs.1d,e,f: $n_{e0}=7 \times 10^{19} \text{ m}^{-3}$) were compared. The transport coefficients were derived according to the procedure described above and are plotted in fig.1g for both densities.

For the outer plasma region, the last reliably determined value of the diffusion coefficients at $r=10-11\text{ cm}$ were kept constant up to the plasma edge. They represent average diffusion coefficients for this region, determining essentially the central time behaviour during the inflow phase. The extrapolated average diffusion coefficient for medium density approaches values consistent with neoclassically predicted ones ($0.09 m^2/s \leq D \leq 0.04 m^2/s$, $10 cm \leq r \leq 17 cm$), but falls below the predictions ($0.2 m^2/s \leq D \leq 0.07 m^2/s$, $10 cm \leq r \leq 17 cm$) for the high density case. This trend is even more pronounced at higher electron densities of $n_{e0}=1.2 \times 10^{20} \text{ m}^{-3}$, e.g. in high confinement neutral-beam heated discharges [6], where transport coefficients were derived ($D(r)=0.07 m^2/s$, $\nu(r)=5 m/s/(r/a)$, #38551), being clearly smaller than predicted by SITAR, in which the fluxes are already reduced by 50% to account for the W7-AS transport optimization.

The convection velocity was extrapolated to vanish in the plasma center and was adjusted in the outer part to fit better the temporal decay of spectral line intensities from different ionization states of aluminum, observed by central-line-of-sight crystal- and VUV-spectrometers (fig.2). For the high density discharge, only a slight correction in the derived $\nu$ (fig.1g) was necessary to excellently fit all experimental data. In the case of medium density, $D(r=10\text{ cm})$ was used for extrapolation and $\nu$ has to be reduced by a factor of 2 for a good compromise in fitting the experimental data radially as well as temporally. A reduction of $D$ in the outer region down to the value $D(r=11\text{ cm})$ would fit the time traces of the spectrometers only for the case of vanishing convection velocity, but will lead to a misfit of the radial profiles.

In the two ECF-heated discharges under investigation, no substantial changes as well in the radial electric field $E_r$ (within the error) as in MHD-activity and electron density profile shape
could be observed from which the difference in transport can be deduced. However, previous measurements of the density fluctuation level as a function of electron density [7] show some inverse dependence, but an effect on the impurity transport cannot simply be concluded.

In the case of high electron density, D is overall lower by a factor of 2-3 compared to the medium density discharge, \( v \) being quite similar. The resulting difference in confinement can be illustrated quite impressively by the change in decay time for the injected aluminium (fig.2). Consequently, also the time evolution of intrinsic impurity radiation and \( Z_{\text{eff}} \) is remarkably different during the flat-top phase for this two densities (fig.3). However, assuming a constant impurity influx of, e.g. intrinsic chlorine from the walls and using just the derived set of transport parameters, the difference in the time traces of Cl-XIV can be qualitatively well described (fig 3, fits): compared to medium electron density, where stationary conditions were achieved well within the pulse length (fig.3 curve b, arrow) due to higher transport, the reduced transport in the high density case causes longer times to establish stationarity (fig.3 curve a).

![Fig.2: Time traces of spectral lines from different ionization states of aluminium after injection by laser blow-off: (a) \( n_{\text{E0}}=7 \times 10^{19} \text{ m}^{-3} \), (b) \( n_{\text{E0}}=3.5 \times 10^{19} \text{ m}^{-3} \).](image1)

![Fig.3: Temporal evolution of impurity radiation and \( Z_{\text{eff}}(\text{SX})\) for (a) \( n_{\text{E0}}=7 \times 10^{19} \text{ m}^{-3} \), (b) \( n_{\text{E0}}=3.5 \times 10^{19} \text{ m}^{-3} \) and simulations for Cl-XIV.](image2)

Similar behaviour for other intrinsic impurity species might explain the signals of the SX-camera, the bolometer and \( Z_{\text{eff}} \). Nevertheless, at high density (\( n_{\text{E0}}=7 \times 10^{19} \text{ m}^{-3} \)) the extrapolated total radiation (bolometer) at the time when the Cl-XIV radiation should reach 90% of its stationary level (approx. at 1.7s, from simulation), stays well below the critical value of 60% of the heating power [8], where the plasma is severely affected by radiation.

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Review of 3-D equilibrium calculations and reconstructions for W7-AS

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Abstract
Knowledge of the 3-dimensional structure of the plasma equilibrium is a prerequisite for experiments in stellarators and for the interpretation of the results. Therefore, since the calculation of equilibria consistent with the experimental data is an important task, we review the calculations done for W7-AS with applications to high β and large toroidal currents for stellarator-tokamak hybrid operation. We also present a novel method for fast equilibrium reconstruction for stellarators based on function parameterization.

Introduction
Although the geometry of the flux surfaces of a stellarator is largely determined by the external coil system, the plasma current densities induced by finite β, internal or external current drive (bootstrap-, Ohkawa-, ohmic- and ECCD currents) may lead to considerable changes in the equilibrium fields.
To calculate 3-D MHD equilibria with free boundary, we use the NEMEC-code [1] assuming nested flux surfaces. The input consists of profiles for pressure and toroidal current, an estimate of the magnetic axis position and the plasma boundary, and the vacuum magnetic field. Based on an energy principle the equilibrium is determined iteratively using a steepest gradient method. The equilibrium quantities like flux surface geometry and magnetic field are given in Fourier series with respect to the cylindrical toroidal angle and a poloidal angle coordinate on a radially discretized grid.
An equilibrium reconstruction of a discharge at a given time point clearly implies an iterative process of adjusting the input parameters such that the resulting equilibrium data best match the experimental ones.

Equilibrium Calculations(1) : β effects
Usually, the pressure induced current densities dominate the changes in the magnetic configuration in net toroidal current free discharges. Although the reduced average toroidal curvature of W7-AS leads to smaller Pfirsch-Schlüter (PS) currents (a factor $\sqrt{2}$ compared to a conventional stellarator) they may give rise to appreciable Shafranov shifts, changes in the rotational transform ($\iota$) profile and a displacement of the plasma as a whole in the accessible $\iota$-range of W7-AS ($0.26 < \iota_{\text{v}} < 0.56$). Very good agreement of the NEMEC calculations with experimental data was shown in Refs 2 and 3 for the high β cases. For W7-X the higher reduction of the average toroidal curvature as well as the higher operational range in $\iota$ ($\iota_{\text{v}} \approx 1$) leads to much more stable magnetic configurations in the sense of finite-β equilibrium changes.

Equilibrium Calculations(2) : Toroidal current effects
Internally or externally driven net toroidal currents also affect the $\iota$-profile and the geometry of the equilibrium flux surfaces, depending on their magnitude and distribution. At W7-AS, the usual operation is net current free, which means that the toroidal plasma current is kept at zero by inductive compensation. However, the toroidal current ($I_\phi$)
Figure 1: Comparison of Soft-Xray emissivity with NEMEC flux surfaces for $I_p = 30$ kA (right) and $-30$ kA (left) shows good agreement in the gradient regions. A very flat central emission profile dominates the central contours for $I_p = 30$ kA.

within each flux surface does not vanish and thus alters the $\psi$-profile. A large aspect ratio estimate for the change in $\psi$ is given by $\Delta \psi(r_m) = I_p(r_m)R/(2\Phi(r_m))$ with $r_m$ and $R$ the minor and major radii, respectively, and $\Phi$ the toroidal magnetic flux. Nevertheless, for vanishing net current there is generally no influence on the flux surface geometry. Therefore, to evaluate $\psi$-profiles in net-current free discharges, it is sufficient to know the finite $\beta$ changes on $\psi$ and then correct the profile using the toroidal current profile deduced from neoclassical theory and deposition profiles.

Introducing a net toroidal current also changes the flux surface geometry depending on the current’s magnitude. Calculations show that the excursion of the flux surfaces in $R$ increases with positive net current (the shape tends to be more oblate) and decreases with negative net current (higher vertical elongation). Good agreement of the calculated flux surfaces in 2 cases with large toroidal currents ($\pm 30$ kA at $B=2.5$ T, $r_{out} \approx 0.4$) with the tomographically reconstructed soft X-ray emissivity measured by the new MiniSoX camera system [4] is shown in Fig. 1.

Discussion of NEMEC calculations

Up to high $\beta$ and rather high currents, the equilibrium calculations for W7-AS with NEMEC are in good agreement with the experiment. However, the computational effort is orders of magnitude higher compared to the calculation of tokamak equilibria. For fast reconstructions accompanying the experiment on a shot to shot basis or even online, NEMEC is not suitable with present-day computing power. Therefore, other approaches like function parameterization, which is discussed below, have to be explored.

Despite the good agreement with the experiment, there are limitations in the NEMEC applications. The assumption of nested flux surfaces excludes the detection of ergodic regions or islands. For this, more advanced equilibrium solvers have to be applied, like HINT or PIES. However, their computational requirements are orders of magnitude higher again than that of NEMEC. Furthermore, the Fourier representation of the equilibrium quantities together with the energy minimization method applied limits the resolution of boundary structures which require high Fourier harmonics and contain comparably small energies. Such $5/m$ resonant structures ($m \geq 12$) have been seen by video observations of visible light in high $\beta$ discharges. Nevertheless, NEMEC shows a smooth plasma boundary without the actual indentation.
A hybrid FP/interpretive method for equilibrium recovery

Interpretive methods for determining plasma equilibria are widely used in tokamak analysis. Input parameters to an equilibrium code are iteratively adjusted such that simulated diagnostic signals from the resulting equilibrium best match experimental data. These methods are generally unsuitable for stellarators since each iteration involves the full solution of a 3D equilibrium code, a task requiring roughly one hour of CPU time on the Cray J-90 at IPP for standard W7-AS equilibrium calculations using NEMEC. The application of function parameterization[5] (FP) to W7-AS is under development. FP seeks simple functional relationships between plasma parameters and diagnostic measurements over a database of simulated equilibria. This facilitates rapid equilibrium reconstruction, here in terms of magnetic data and a prescribed pressure profile. We have developed a novel interpretive method for equilibrium identification based on FP reconstructions that can be performed in the order of a few tens of seconds on a workstation.

FP database

Here, the database consists of circa 400 NEMEC equilibria with zero net toroidal current. They are chosen by randomly varying 8 input parameters over ranges appropriate to W7-AS, namely 3 ratios of the 4 field coil currents, a limiter position and a 4-parameter pressure profile chosen from the following family (s is normalized toroidal flux):

\[ p(s) = p_0 (1 - s)^2 \exp(as + bs^2 + cs^3) \]

Interpretive scheme

The Thomson scattering diagnostic on W7-AS gives electron temperature and density (and thus the electron pressure \( p_e \)) on \( \leq 20 \) channels along a horizontal line-of-sight through the magnetic axis in a symmetry plane (\( \phi = 0 \)) at a single timepoint during a discharge, i.e. \( p_e = p_e(R_i) \), \( R_i \) being the major radii of the Thomson channels. A smoothing polynomial in \( R \) is fitted to the \( p_e \) data, allowing evaluation between channels. The iterative procedure attempts to reproduce the spatial to flux transformation \( s(R) \) in \( p_e(R) \) since \( p_e \) is constant on flux surfaces. Starting from an initial guess, the equilibrium pressure profile \( p_{eq}(s) \) is varied to minimize the quantity

\[ \int (p_e[R_{in}(s)] - p_e[R_{out}(s)])^2 ds \]

subject to the constraints that \( p_{eq}(s) \) is non-hollow and the boundary flux surface coincides with the zeroes of \( p_e(r) \). An additional restriction that the kinetic energy content match the diamagnetic energy from the experiment can be optionally enforced. The interpreted fit thus depends only on the topology of the \( p_e(R) \) profile and not on its magnitude. In fact, it could be done separately for the Thomson temperature and density profile. Clearly, mismatches between the theoretical and physical magnetic configurations could falsify both the present scheme and the standard calculations. Consistency checks with additional spatially resolved or global diagnostic data would reveal such discrepancies, if present.

Comparison with standard calculations

Results from the interpretive method and the standard NEMEC simulation for shot 31909 are illustrated. The upper plots show the Thomson \( p_e(R) \) with smoothing polynomial and the interpreted \( p_{eq}(s) \) together with the standard NEMEC profile. The lower plots show the equilibrium flux surfaces in the Thomson (\( \phi = 0 \)) plane and the final asymmetry error in \( p_e(s) \).
Though the two fits are of comparable quality, the interpreted fit is better than the standard calculation in the inner half of the plasma. This is not unexpected since the standard calculation uses the same optimization criterion with manual intervention required between iterations, rather than an automated least-squares fit. The $p_e$ symmetry can be further improved by varying additional pressure parameters during iterations. Only three of a possible six were varied here to prevent overfitting. The interpretation for this case took roughly 20 seconds on a workstation, whereas the standard simulation required several NEMEC calculations.

Work in progress includes investigation of the effects of signal noise on the recovery. Moreover, although scalar equilibrium parameters and flux geometry are well recovered, a model that reproduces the $\varepsilon$-profile well has so far proved elusive. Once this is achieved, the procedure will be applied to equilibria with finite net toroidal current.

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RADIATIVE INSTABILITIES IN W7-AS PLASMAS WITH HIGHLY RADIATING BOUNDARIES

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Introduction

In low radiation ATF (Advanced Toroidal Facility [1]) stellarator plasmas it was concluded that the lack of efficiency in the radial energy transfer to sustain the local energy balance in regions with enhanced losses leads to plasma collapse [2]. The Proctr code [3] was used to analyse NB heated ATF collapsing discharges with moderate radiation levels (Prad/Pin = 30 %) finding that radiative instabilities near the plasma periphery may be the trigger mechanism. On the other hand, low to intermediate Z impurity injection is being considered as a suitable procedure to cool the plasma edge region and, therefore, protect divertor plates in high power fusion experiments from thermal overloads. Then, the dominant plasma contaminant must be the adequate in every device for radiative edge cooling purposes, since cooling rate profiles determine the local energy lost by radiation.

Nitrogen is widely considered as an appropriate impurity for those purposes because of its radiative properties and its capability to be sufficiently pumped by the wall, especially when no active pumping is available [4]. In particular, it is used in tokamaks [5] to get the so called detached plasmas.

Experimental

In a previous work [6], nitrogen was confirmed to be an adequate impurity to effectively cool the edge region of Wendelstein 7-AS stellarator (R = 2.0 m, a = 0.18 m, no active pumping capability). However, steady state complete detachment from the limiters could not be established due to the onset of a radiative instability at a radiation level of about 60 % of the heating power. Further nitrogen injection experiments have been performed in W7-AS stellarator with feedback control of the radiation levels, via VUV line emission (N IV, 765 Å), to maintain the impurity concentration under the limit of radiative collapse [7].

For the present studies a series of net current free ECRH discharges are considered, with flat-top of 1.5 s, at B=2.5 T, I_0/2π=0.34, injected power about 430 kW and electron density n_e = 10^{19} m^{-3}. Transport analysis have been done with the predictive transport code Proctr. Here we are presenting the results obtained for three representative discharges, with no N_2 injection, 1 % and 2.5 % N_2 concentrations, shots # 36770, 36780 and 36783, respectively. In Figure 1 time evolution of line density, plasma stored energy and radiated
power are plotted for the three mentioned discharges. A fraction of about 60% radiated power (measured from bolometry) is reached in discharge 36783 and the attempts to exceed this limit lead to a radiative instability and to feedback induced oscillations of the discharge parameters rather than to a collapse. Decreasing the nitrogen influx and, thus, the radiation level within the discharge duration leads to re-establishment of stationary conditions.

W7-AS is operated with two carbon limiters and so a moderate concentration of carbon is considered to be present in discharge 36770. The experimentally measured density profile is flat up to \( r/a = 0.5 \), whereas the electron temperature profile is rather peaked for the three discharges.

**Transport Analysis**

To simulate the discharge dynamics the following experimental parameters have been used as inputs for Procr: total radiated power, stored energy, heating power, line density, and the dominant species. Density is feed-back maintained with gas puffing plus limiter recycling. Power deposition is notably localized, according to [8]. The experimental electron temperature and density profiles have been used to benchmark the transport models chosen. The modelled temperature and density profiles can be seen in Figure 2, marked with the nitrogen concentration of the discharge.

In Figure 3 the radial profiles of loacal and integrated power electron balances are plotted for the three discharges in steady state regime, at time \( t = 700 \) ms. Again every plot is marked with the nitrogen concentration. It can be seen the very different behaviours of the radiated power profiles, depending on the nature and concentration of the impurity. In discharge 36770 carbon impurities are accumulated mainly at \( r/a = 0.6 \), and the radiation profile is rather flat. When about 1% of nitrogen is injected (discharge 36780) the global radiated power increases in the whole plasma column, mainly due to a small increase of \( Z_{\text{eff}} \), but this increment is more pronounced at the edge, where the radiation profile becomes clearly peaked. For this nitrogen concentration only slight changes appear in plasma transport. Further nitrogen injection notably modifies the power balance at the outermost region of the plasma. Due to the strongly localized ECR absorption in these plasmas, almost no changes in the power deposition profile are seen because 100% single pass absorption is reached, since the temperature is still high enough. Differences are found mainly in particle convection and in the power interchanged with ions by collisions. The edge radiation peak propagates inwards and increases in intensity. This nitrogen concentration somehow represents the maximum admissible limit beyond which a steady state discharge cannot be held.

Surpassing this limit leads to a sudden plasma contraction together with an accumulation of impurities in the plasma core. That is, a thermal collapse occurs. An artificial discharge has been modelled based on the data and transport parameters considered in
discharge number 36783. Concentration of nitrogen, radiated power, and stored energy have been proportionally modified to provoke the appearance of radiative instabilities.

Discussion

Nitrogen injection has been demonstrated as a good procedure for edge cooling in W7-AS, provided that feedback is used to maintain the plasma under the limit of radiative collapse. In spite of the high radiated power, energy confinement time is only slightly degraded in discharge 36783 as compared with the reference one. This fact is confirmed by the diminishing of the estimated heat conductivity in the bulk plasma. This result agrees with the one presented in [7].

When comparing these results with the ATF NB-heated discharges [2] a difference is found in the plasma behaviour. As was concluded there, the detailed radial profile of the electron thermal conductivity plays a key role to supply power to the place where is needed. In the plasmas studied here this term is not able to cure the local increase of power losses and so the plasma can only reacts by modifying convection and power transmission to ions, i.e., particle transport must be mainly involved. One possible reason for these diverse behaviours can be found in the plasma heating method, since NBI power deposition profile is wider than ECRH (for high electron density and temperature) and changes according to the evolution of plasma parameters, such as density and temperature, higher power density is available near the maximum radiative losses region. Another likely cause may be that the magnetic configuration, strong sheared in ATF torsatron and almost shearless in W7-AS helias, is the responsible of the different transport regimes that appear in the two devices.

To clarify this problem simulations of ECR heated ATF discharges are underway. In particular the formation of transport barriers is planned to be studied.

References

The effects of field reversal on the W7-AS island divertor at low densities

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1. Introduction

In the last years, considerable effort has been made on W7-AS to explore the diversion properties of the boundary magnetic islands in high $t$ ($\geq 0.5$) configurations for control and exhaust of the edge plasma [1]. Besides investigating basic aspects of divertor operation, such as high recycling, detachment and impurity control, understanding the drift effects on the particle and energy balance between the contacting plates is essential to improve the divertor performance.

Poloidal asymmetry of the plasma distribution in the island SOL of low density discharges has been observed for W7-AS divertor configurations at $t=5/9$, with the inboard plates intersecting the islands through the O-point. Unlike in most single-null tokamak divertors where, generally, higher density is found on one divertor plate while higher power flux on the other, the W7-AS divertor experiments show that the density, the power flow and the particle flux distributions have the same phase shifts in the island SOL. After reversing the magnetic field, the observed asymmetry is reversed. This asymmetry can be explained by an $E_T \times B$ drift in the island SOL as follows. The radial temperature gradient within the open island surfaces, which are intersected and radially linked by the electrically conducting plates, leads to a radial gradient of the plasma electrostatic potential and hence to a radial electric field inside the islands. The $E_T \times B$ drift delivers additional particles and energy in the island SOL along the upper or lower island fans to the targets, depending on the direction of the toroidal field. This picture is in good agreement with the experimental observations. In order to understand the drift effects quantitatively, in a first step the 3D Monte Carlo transport code EMC3/EIRENE [2] has been extended to allow the treatment of the poloidal drift. Calculations show the same phase shift of the density contours as observed.

2. Experimental observations

At high $t$ ($\geq 0.5$), the edge magnetic structure of W7-AS is governed by inherent magnetic islands. Having a considerable size and an appropriate internal rotational transform, the $t=5/9$ islands have been chosen for the divertor experiments. Ten symmetric divertor plates are installed on the inboard side of the wall with a toroidal location which is symmetric to the triangular cross sections (Fig. 1). Each plate intersects poloidally two islands and the radial intersection position can be easily changed by application of a vertical field shifting the magnetic flux surface configuration horizontally with respect to the plates. In this work, however, only a configuration with the plates cutting the islands through the O-point is considered, which allows easy estimation of the radial electric field throughout the islands. Each divertor plate is poloidally segmented into 8 tiles and calorimetry measurements on each tile give a poloidal
distribution of the power load. In addition, a two-dimensional Hα diode array is fixed horizontally at the outside of the torus, looking at a plate poloidally.

Fig. 1 Right: poloidal cross section of the \(\tau=5/9\) configuration at the toroidal plane of an inboard plate. Left: calorimetry and Hα data for two discharges with B-field reversal.

The discharges investigated are at extremely low density (\(\langle n_e \rangle_{\text{line}} < 10^{19} \text{ m}^{-3}\)) with ECR heating power of about 200 kW. The results for two discharges with B-field reversal are compared in Fig. 1, in which the interaction of the two islands with the plate is well reflected by the diagnostics. It should be mentioned that the present inboard plates are not yet the optimized ones for W7-AS divertor operation. This explains why the power load and recycling are strongly inhomogeneous between the two islands. However, this does not affect the investigation of the drift problem addressed in this paper. We pay our attention only to the lower island, which, due to the longer connection length, carries higher particle and power flows to the plate. This island is cut by the two lowest tiles of the plate. Therefore, the two lowest diagnostic channels of the calorimeter and the Hα array are best suited to verify the differences in particle and energy flows between the upper and lower island fans. In the positive B-field case, higher energy and particle

Fig. 2. Left: Arrangement of the Langmuir probe array and the vacuum structure of the touching islands. Right: Measured density contours for two low density discharges with B-field reversal.
outflows are found at the lower island fan, shifting to the upper one as the B-field is reversed.

A Langmuir probe array consisting of 16 probes is used to measure the density distribution of the edge plasma. The probe array is poloidally shaped to follow the main edge structure of the configuration and toroidally placed on a triangular cross section (Fig. 2). It can be shifted horizontally inward and outward by a few cm, thus providing 2D plasma density distribution in the region of interest. The resulting density contours from the probe array for the two field cases are illustrated in Fig. 2, in which the poloidal structure of the island chain is clearly identified. However, up/down asymmetries in the density contours emerge, showing a strong dependence on the field direction. Comparing density contours for the positive B-field with the vacuum island structure, we find that the plasma density is higher on the lower island fan than on the upper one. As the B-field is reversed, the higher plasma density shifts to the upper island fan, showing a simple correlation with the B-field direction.

3. Theoretical considerations and 3D simulations

As the observed asymmetries change with B-field reversal, the classical drifts are reasonably considered to play an important role because of their close dependence on the B-field direction. We examine the classical drifts, in order to isolate the candidates which can explain our experimental observations. The classical drift terms contributing to the fluid equations of Braginskii [3] are the electric drift $\mathbf{E} \times \mathbf{B}$ and the diamagnetic drift $\mathbf{B} \times \nabla \mathbf{B}$. As has been shown by Chankin [4], the diamagnetic drift and the corresponding diamagnetic energy flux are dominated by divergence-free parts which do not deliver particles and energy to the target, nor give rise to any particle or energy accumulation or sink. The non-divergence-free components are due to spatial variations of the magnetic field, i.e. $\mathbf{B} \times \nabla \mathbf{B}$ and $\nabla \times \mathbf{B}$. In a net-current-free stellarator such as W7-AS, the latter term is directly related to the plasma diamagnetic current and, in practice, can be neglected for the low density discharges involved in this paper. The $\mathbf{B} \times \nabla \mathbf{B}$ term is a vertical drift and, therefore, its effect on the asymmetries in the islands is strongly reduced by the helical path of the islands around the torus.

According to these considerations, we can restrict our attention to the electric field drifts $E_r \times \mathbf{B}$ and $E_\theta \times \mathbf{B}$. The former is related to the radial gradient of the sheath potential due to the temperature drop towards the O-point. The latter is due to the potential drop along the B-field line from up- to downstream, associated with the parallel E-field arising from the thermal force. Following the particle flux estimation made by Chankin for tokamak divertors [5], we have the total radial and poloidal particle drift fluxes:

$$\Gamma_r = N_2 \pi R (k T_{e,up} - k T_{e,down}) n_0 e B$$

and

$$\Gamma_\theta = N_2 \pi R 3k T_{e,down} n_0 e B,$$

where $N$ is the poloidal periodicity of the configuration. Since $N=1$ for single-null divertor tokamaks, the drift effects in $N=1$ island divertors like W7-AS ($N=9$) are more pronounced than in tokamaks. Comparing the radial and poloidal drift fluxes, we find that the condition for $\Gamma_\theta$ dominating over $\Gamma_r$ is $T_{e,down} > (T_{e,up} - T_{e,down})/3$. This is the case for the low density discharges discussed here. Heat transport simulations with the EMC3/EIRENE code showed that the temperature is about 90 eV at upstream position, dropping to 60 eV at the target, which agrees with the downstream temperature deduced from the Langmuir probe array data. A radial drop in temperature of about 20 eV was estimated from the Monte Carlo code. The resulting positive and negative poloidal drift velocities were inserted into the code, leading to the same
phase shift of the density contours as observed (Fig. 3), which can be briefly explained as follows. In the picture assumed here, the particles created by low recycling in the main plasma and diffusing outside across the LCFS experience a poloidal drift in addition to the parallel motion along the island fans. A particle accumulation results in the upper or lower island fan, depending on the B-field direction.

The convective power flux driven by the poloidal drift is estimated to be $-90$ kW, which is about half of the total heating power. This implies stronger asymmetry of the power fluxes between the upper and lower island fans than the one observed. One reason for the discrepancy is that the power fluxes along the upper and lower island fans cannot be sharply resolved by the two lowest channels of the calorimeter. Secondly, due to the large ratio of connection length to island size, a significant fraction of the drift power flux can circulate around the island between the discontinuous plates without reaching them. A quantitative assessment of these effects for island divertors needs a selfconsistent 3D treatment of the drifts, which is beyond the present capability of the EMC3/EIRENE code.

4. Conclusions

For low density, low recycling discharges, poloidal asymmetries in the island SOL have been observed for the W7-AS divertor configurations at $t=5/9$, with the islands deeply cut by the present inboard plates through the O-point. Higher density, power flow and particle flux are measured for the lower or upper island fans, depending on the direction of the toroidal magnetic field. The asymmetries are considered to be driven by an $E_T \times B$ drift resulting from the radial temperature gradient in the island. The $E_T \times B$ drift delivers additional particle and energy fluxes along the island fans to the divertor plates, leading to the same phase shift of the density, energy and particle flux distributions in the island. 3D Monte Carlo calculations including the poloidal drift can well reproduce the density asymmetries measured by the 2D Langmuir probe array.

5. References

Study of density turbulence and coherent mode activity in W7-AS by microwave reflectometry


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Abstract - In the context of transport investigations different heating scenarios have been studied in W7-AS. When the plasma is heated by NBI (pure as well as combined with ECRH) the signals from the microwave reflectometer show coherent mode activity in the frequency range of 20 to 90 kHz depending on the iota profile. These modes have been observed for all the radial positions probed by the reflectometer and are identified as global Alfvén modes.

Experimental system - At W7-AS (R = 2 m, a < 0.17 m, 5 field periods, modular stellarator coil system) a reflectometer is installed [1,2] for density profile and fluctuation measurements. The system uses X-mode propagation in the W-band (75-110 GHz) for probing radial positions corresponding to densities between 1 to 6 \(10^{19}\) m\(^{-3}\). At the toroidal position where the reflectometer is installed the plasma has a nearly elliptical shape with the minor axis along the equatorial plane of the torus. Corrugated horns in combination with elliptical mirrors are used as emitting and receiving antennas resulting in a focused Gaussian beam of 2 cm diameter at the reflecting layer. Heterodyne detection allows to measure phase and amplitude fluctuations separately. An Amplitude Modulation (AM) system [2] integrated into the reflectometer provides a time delay measurement which is used to obtain electron density profile information. This time delay of the signal depends on the cutoff position and on the local density gradient being, in most of the cases, more sensitive to the latter one. Therefore the time delay signal can be considered as a monitor of the local density gradient.

Experimental results - In the context of transport investigations the microwave reflectometer was operated in discharges with different heating scenarios: ECRH, NBI, and combined NBI+ECRH. Two types of discharges are studied in this paper. (1) Plasmas with combined, ECRH (500 kW) and NBI (1.5 MW), heating which show a very high ion temperature, reaching values up to 1.5 keV [3]. (2) Discharges with pure NBI heating with lower injected power (0.5 MW), which reach energy confinement times about twice as large as predicted by neoclassical theories [3]. Both improved confinement scenarios are achieved under low wall-recycling conditions and are characterized by a low edge density and high density peaking factor [4]. Electron density profiles from Thomson scattering and AM-reflectometer for these two types of discharges are shown in Figure 1. The radial positions probed by the microwave reflectometer are between 11 cm and 15 cm. Shots with pure NBI heating and very
high energy confinement time show an electron density profile steeper than shots with combined heating (ECRH+NBI) and very high ion temperature.

Figure 1: Electron density profiles from Thomson scattering (circles) and from microwave reflectometer (lines) in the discharges of interest: combined ECRH+NBI heating and pure NBI.

In general, if the plasma is heated by NBI (pure as well as in combination with ECRH) the reflectometer signals show coherent mode activity. Other diagnostics (X-ray, ECE, Mirnov coils, etc.) also observe these modes depending on the iota profile, and have been identified as global Alfvén modes [5].

The interpretation of reflectometry phase measurements is complicated by an asymmetry in the phase fluctuation spectra, called "phase runaway" [6]. Under both plasma conditions discussed in this paper this effect is several times lower than in pure ECRH plasmas, therefore the measured phase fluctuations mainly can be interpreted in terms of density fluctuations.

Figure 2-a: Phase fluctuation spectra during a purely NBI heated plasma for different radial positions.

Figure 2-b: Time delay spectra during a purely NBI heated plasma for different radial positions.

During pure NBI heating the coherent modes observed by reflectometry appear in the frequency range of 10-40 kHz. As an example Figure 2-a shows the power spectra of the reflectometer phase. Two different modes with frequencies of 27 kHz and 37 kHz are present.
simultaneously in the spectra. The frequency of these modes is independent of the radial position probed by the reflectometer ($r_{\text{eff}} = 11-15$ cm) and their amplitude increases towards the plasma centre.

The time delay signal has been measured simultaneously. The dispersive effects, associated with the local density gradient, represent the largest contribution to the measured time delay. Therefore, fluctuations in the time delay signal can be considered as a monitor of the local density gradient. The spectral analysis of this signal, shown in Figure 2-b, does not always show the coherent modes. Only in the most inner radial positions probed with the reflectometer, the highest frequency ($f = 37$ kHz) coherent mode is seen. This absence of coherent modes in the time delay signal can be interpreted as a movement of the density profile.

During discharges with combined heating, $P_{\text{ECRH}} = 400$ kW and $P_{\text{NBI}} = 1.5$ MW, where the ion temperature reaches values of about 1.45 keV, the frequency of the observed modes is higher than during pure NBI heating, appearing modes up to 90 kHz. Similarly to purely NBI heated discharges, the amplitude of these modes also increases towards the plasma centre and their frequency is constant along all the probed radii. As an example Figure 3-a shows the spectra of the phase signal where several global modes are observed at different frequencies. The coherent mode with the highest frequency appears at 85 kHz in all the probed positions. The spectra of the time delay signal, Figure 3-b, shows this highest frequency coherent mode in most of the radii tested while the lower frequency modes are not observed.

**Figure 3-a:** Phase fluctuation spectra during a combined NBI + ECRH heating for different radial positions.

**Figure 3-b:** Time delay spectra during a combined NBI + ECRH heating for different radial positions.

Microwave reflectometer results were also obtained during discharges in which the heating was switched from pure NBI (0.5 MW) to pure ECRH (0.4 MW). In contrast to the purely NBI heated phase in the purely ECRH heated plasma these coherent modes disappear in the reflectometer signal in all of the probed radii, indicating that the heating and not the magnetic configuration is related with the origin of the modes. Similar radial positions were probed by
the reflectometer in the two heating scenarios. In the NBI heating regime (Fig. 4-a) a coherent mode with a frequency of about 45 kHz can be observed in all the probed radii while with ECRH (Fig. 4-b) this mode does not appear. The absence of the 45 kHz coherent mode in the ECRH regime has been also observed by other diagnostics, e.g. Mirnov coils as shown in the figure 5.

**Figure 4-a:** Phase fluctuation spectra during a purely NBI heated plasma for different radial positions.

**Figure 4-b:** Phase fluctuations spectra during a purely ECRH heated plasma for different radial positions.

![Figure 5: Temporal spectrum evolution of Mirnov coil signal in a shot where the heating was switched from pure NBI to pure ECRH. The 45 kHz coherent mode disappears when the heating scenario is changed at 0.5 s.](image)

In both heating regimes the low frequency broadband incoherent density fluctuations ($f < 40$ kHz) decrease as positions further inside are probed. In contrast the power of the 45 kHz coherent mode observed with NBI is nearly independent of the probed positions.

**Conclusions** - In NBI heated plasmas coherent mode activity is observed which shows up as coherent oscillations in the reflectometer signal if the phase runaway is low. The spectral analysis of the phase signals shows very sharp frequency lines. The modes decay with about 200 $\mu$s after NBI is switched-off and their frequencies scale roughly with the Alfvén velocity [5]. They are observed in all the probed radial positions indicating a global mode structure. These observations are in agreement with an interpretation as global Alfvén eigenmodes.

ENERGY and DENSITY INHOMOGENEITIES DRIVEN by TOROIDALLY LOCALIZED ECRH in W7-AS

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W7-AS is a modular stellarator with 5 field periods. Close to the plasma axis the poloidal Fourier components of the 3-D magnetic field tend to vanish and the configuration becomes similar to five “toroidally linked mirrors”. During on-axis ECRH, the highly focused injected beam interacts, besides passing particles, with only the trapped particle population confined in the mirror where ECRH is launched. In low density ECRH discharges in magnetic configurations with high fractions of trapped particles, strong toroidal anisotropy in the electron distribution function may arise. Aim of this paper is the quantitative evaluation of the collisional redistribution of the deposited power between the different classes of trapped particles and an investigation of the possibility to detect experimentally the ECRH driven toroidal inhomogeneities by means of the ECE diagnostic.

**ECRH DRIVEN TOROIDAL ANISOTROPIES**

The problem has been investigated by means of the bounce-averaged Fokker-Planck code FPTM /1/ suitable for the study of ECRH in periodic magnetic fields. The analyzed scenarios refer to second harmonic on-axis heating by a collimated beam of X-mode polarized waves injected perpendicularly from low-field-side in the “minimum B” and the “maximum B” magnetic configurations characterized by a local minimum or a local maximum, respectively, in the toroidal plane where the ECRH beam is injected /2/. While in the case of “maximum B”, the power is mainly deposited to passing particles and redistributed toroidally in a rather uniform way, in the case of “minimum B” strong local perturbations of the distribution function can arise for the population of electrons trapped in the local toroidal mirror where ECRH is injected, as shown in Fig. 1 (for a target plasma with density \(n_0 = 1 \cdot 10^{13} \text{ cm}^{-3}\), temperature \(T_0 \simeq 1400 \text{ eV}\), ECRH power density 5 \(\text{W/cm}^2\)). In this mirror, the deviations from the Maxwellian distribution are considerably stronger than those found in the other mirrors where heating appears only due to collisional energy transfer from passing particles.

In steady-state, the redistribution of density and energy between the different populations of particles is mainly determined by the balance between ECRH and collisions. The increase of the perpendicular energy of the electrons resonant with the ECRH beam, has two main consequences: the power absorbed by “barely” passing particles can cause their trapping, while the power deposited in the trapped particle region of velocity space tends to create an enhancement of trapped particles at higher energy and a depletion at lower energy. Due to the velocity dependence of pitch angle scattering, the anisotropy at low velocity is more efficiently counteracted than at high velocities, causing a net flux through the boundary between trapped and passing particles. On the other side, as trapped particles of higher perpendicular energy are poorly confined in Stellarators, losses will also play a role in the redistribution.

The ECRH driven density perturbation is shown in Fig. 2. In the “minimum B” configuration, the density of the population of directly heated trapped particles is increased mainly due to the collisional redistribution driven by ECRH. In the other mirrors loss effects are dominant and the number of trapped electrons is slightly reduced.

In the “maximum B”-scenario, the deviations are considerably smaller in amplitude (different scales are used for the ordinates in Fig. 2), as most of the power is now absorbed by passing particles and homogeneously redistributed toroidally. A qualitative similar toroidal dependence is obtained for the ECRH driven energy perturbation (the relative energy perturbation being three times greater in magnitude for the case under examination).
Fig. 1 Energy spectra, along the line $v_\parallel = 0$, of the bounce-averaged distribution function. Solid line: toroidal mirror with ECRH launch; dashed line: other mirrors; dotted line: initial Maxwellian (for reference). Fig. 2 Toroidal dependence of the relative density perturbation driven by ECRH for the "minimum B" (top) and for the "maximum B"-scenario (center). The small rectangle in these figures indicates the toroidal position of the launching system. On the bottom the normalized magnetic field used for the bounce-averaging is shown.

The inhomogeneity in the electron density generates toroidal electric fields. The density perturbation under stationary conditions is determined by the balance between the particle fluxes driven by ECRH and by these electric fields. As ions remain unperturbed, any perturbation of the electron density results in the generation of an electric field (proportional to the density gradient), that, in a self-consistent treatment, would tend to reduce the ECRH driven density perturbation.

**DETECTION BY ECE**

The ECE intensity, $I$, can be evaluated from the "radiative transfer equation" /3/:

$$N_T^2 \frac{d}{ds} \left( \frac{I}{N_T^2} \right) = \eta - \alpha I \quad \Rightarrow \quad \frac{I(s)}{N_T^2} = \int_0^s ds' \frac{\eta(s')}{N_T^2(s')} e^{-\int_s^{s'} \alpha(s'') ds''}$$

(1)

where $N_T$ is the ray refractive index, $\eta$ and $\alpha$, the emission and absorption coefficients, respectively, and $s$ the space coordinate along the ray trajectory.

For a Maxwellian distribution function, $\eta$ and $\alpha$ are related by a simple multiplicative constant, in agreement with Kirchhoff’s law:

$$\frac{\eta(k, \omega)}{\alpha(k, \omega)} = N_T^2 \frac{\omega^2 T}{8\pi^3 c^2}.$$  

(2)

In the general case the evaluation of $\eta$ and $\alpha$ requires the determination of the (warm) polarization of the wave and an integration along the resonance curve $\omega - k_\parallel v_\parallel - \frac{m_e e}{2} = 0$ in velocity space /4/. While for the emission coefficient the integrand is proportional to the electron distribution function, in the expression for the absorption coefficient a differential operator appears (corresponding to the derivative of the distribution function along the “diffusion path” in velocity space) /4/. In this paper the wave polarization has been evaluated from the Maxwellian bulk, while for the integration in velocity space the deviation of the electron distribution function predicted by the FPTM-code has been taken consistently into account.
Fig. 3 left (normalized) emission coefficient, centre absorption coefficient and right "radiative temperature" for the scenario of Fig.1. The same line notation as in Fig.1 is used and the vertical lines corresponds to resonant electrons with energy 15,10,5 keV, respectively.

The normalized emission \( \tilde{\eta}(k, \omega) = \eta(k, \omega) \cdot (8\pi^2 c^2)/(N^2 \omega^2) \) and the absorption coefficients are shown in Fig.3, together with the "radiative temperature", \( \tilde{\eta}(k, \omega)/\alpha(k, \omega) \), for the distribution function in Fig.1.

In the ECRH injection plane ("elliptical plane") the magnetic field topology is rather similar to that found in a Tokamak \((B \propto 1/R)\). Two viewing geometries have been investigated for the ECE-antenna. The first refers to perpendicular viewing from the low-field-side and the second to vertical viewing. Low-field-side observation is characterized by a rather strong gradient of the magnetic field, while for vertical observation the changes in B are weaker and a local maximum appears close to the crossing of the viewing chord with the plasma axis /5/. In order to integrate Eq.(1), the distribution function, determined by the Fokker-Planck code only close to the plasma axis, has to be extended radially. The dominant broadening mechanism is related to \( \nabla B \)-drift of ripple trapped particles. This kind of "convective losses" were examined in /2/. In this paper we use a simple analytical model, \( f(r, v) = f_{Max}(r, v) + [f_{FP}(0, v) - f_{Max}(0, v)] \cdot \exp(-r^2/\sigma^2) \), where \( f_{Max} \) is the Maxwellian, \( f_{FP} \) the Fokker-Planck distribution and \( \sigma \) a "broadening parameter" \((\sigma = 3 \text{ cm, in this paper})\). The deviation of the distribution function from the Maxwellian is therefore distributed on the magnetic surfaces in a poloidal uniform way, only the amplitude of the ECRH-driven perturbation decreases with increasing radii.

The emission spectra obtained integrating Eq.(1) along these viewing chords, for the case of 2nd harmonic X-mode emission, are shown in Fig.4. Together with the emission of a Maxwellian plasma the spectra for an antenna localized in the ECRH injection mirror and in a mirror where no ECRH is injected are shown too. As the deviations expected by Fokker-Planck simulations are found at relatively low energies (see Fig.1), the emission detected by the low-field-side antenna, sensitive to the bulk properties of the distribution function, can be strongly perturbed, especially for the mirror where ECRH is injected. On the other side, the vertical observation is generally "insensitive" to perturbations in the bulk of the distribution close to the plasma axis, as this emission tends to be reabsorbed by thermal optically thick plasma at outer radii, during its propagation towards the ECE antenna. Only emission at sufficiently down-shifted frequencies can reach the antenna. No ECRH power is expected to be directly absorbed at these resonant energies and collision and radial losses mainly determine the tail population /2/. The vertically observed ECE-spectrum is therefore determined by the balance between collisions and radial losses. The response to sudden changes in the injected ECRH power (switch-on, -off, modulation) can in principle be used to separate the two mechanisms, due to the different time scales and energy dependencies /6/.

In case of ECE at optically thin harmonics, reabsorption becomes negligible and the emission from the central region, where the deviation from Maxwellian are expected to be stronger, can reach the ECE antenna under any viewing geometry. Fig.5 refers to the case of emission at
The (optically thin) 4th harmonic (the magnetic field has been reduced from 2.5 to 1.25 T). The radiation temperature is relatively weak for these frequencies. The interpretation of the signals related to optically thin harmonics can be made critical due to possible contributions from wall reflection, making the presence of a “beam dump” mandatory.

**CONCLUSIONS**

Strong ECRH driven toroidal anisotropies are predicted by Fokker-Planck modelling. The deviations in the bulk of the distribution function in the toroidal mirror where ECRH is injected are expected to be strong enough to affect the ECE detected from low-field-side. The detected “radiation temperature” is expected to be substantially higher than the one detected in the other toroidal mirrors. Vertical ECE, along B \( \approx \) const, due to the relation between emitting frequency and energy (\( \omega - \frac{\nu_{th}}{\gamma} \approx 0 \)) allows a more direct interpretation of the observed ECE-spectrum in terms of the “line averaged” distribution function. However, while emission at optically thin harmonics and/or high enough resonant energies can reach the vertical antenna, at optically thick harmonics, the emission related to low energy electrons located close to the plasma axis, i.e., where Fokker-Planck simulations predict the strongest deviations from Maxwellian, is generally reabsorbed during its propagation towards the antenna.

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Bolometer measurements and transport simulations of the density limit on the W7-AS stellarator


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1. Introduction

A theoretical basis for understanding the density limit is of importance for a fusion reactor as high edge densities will be needed for power handling. Basic aspects of the density limit can be described by a simple two point power balance of the scrape off layer (SOL) [1,2]. This model assumes pressure constancy and classical parallel electron thermal conductivity along field lines and stable thermal equilibrium in front of divertor/limiter plates with detachment at the density limit when power balance in front of the limiter can no longer be supported. The radiation term in the power balance equation assumes radiation from carbon (coronal equilibrium) in the SOL. The electron density and temperature measurements at the limiter and last closed flux surface, together with the bulk density, temperature and radiation profiles are the experimental quantities of interest. Using the 4 equations of the two point model for a limiter machine [2], the edge density, $n_e$, as a function of electron temperature at the limiter, $T_{\text{lim}}$, for a given temperature fall off length, $\Delta_T$, and net power flux at the last closed flux surface, $q_{\perp}$, can be predicted.

In W7-AS a discharge with rising line integrated electron density reaching a density limit is shown in Fig.1. The sudden degradation in confinement leading to a decrease in the diamagnetic energy indicates the density limit. In contrast to tokamaks [3,4], current free stellarators at the density limit do not suffer from MHD disruptions and can even recover to a steady state.

This paper is a first attempt to apply the simple two point model to density limit experiments in the W7-AS stellarator. The time dependent ASTRA code [5] was coupled to the impurity transport code STRAHL [6] to facilitate modelling of the radiation profile measured by bolometry. The addition of the simple two point model to ASTRA allowed time dependent simulations of the measured density and temperature at the limiter.

2. Experiments and comparison with the two point model

A magnetic field scan from 0.6 T to 2.5 T in NBI discharges with 1100 kW absorbed power, a rotational transform, $\tau$, of 0.43 and an upward density ramp was performed. At the lowest and highest magnetic field the total radiated power is 35% and 70% of the input power respectively. This procedure was repeated in discharges at 1.25 T and 2.5 T for absorbed powers from 380 kW to 1440 kW with $\tau = 0.34$. The edge densities were taken to be 0.25 of the line averaged density and these values are in good agreement with those from the lithium beam and
Thomson scattering diagnostics.

The two point model was adapted for W7-AS and compared with data from these power and magnetic field scans. From the difference between the input power and bolometer measurements of the bulk radiated power, $q_L$ can be calculated. The experimental observation of rising radiated power and decreasing $q_L$ with rising line integrated density means that a critical edge density will be reached where $q_L$ is no longer sufficient to maintain power balance. Shown in Fig.2, for a typical case in W7-AS with $\Delta_T = 1.2$ or 3 cm, and a connection length to the limiter plate, $L_c$, of 18 m, this maximum $n_S$ at the density limit is plotted. It can be seen from the magnetic field scan that for low magnetic field $\Delta_T > 3$ cm is suggested, well above the maximum indicated by scaling studies.

3. Scaling

The scaling of $n_S$ with $q_L$ and $B_0$ found by regression analysis, as shown in Fig. 3, is $n_S = 0.52 \ q_L \ 0.5^{\pm 0.2} \ B_0 \ 0.8^{\pm 0.15}$. This scaling can be compared to that found in density limit studies for JET limiter discharges with coefficients of 0.66 and 0.33 respectively [2]. The JET scaling is close to that expected when Bohm diffusion is assumed for the scaling of $\Delta_T$. In the case of the $B_0$ scan on W7-AS, the plasma is pushed up against the inboard limiters. To first order it is expected that $L_c$ remains constant as a function of $B_0$ so that the geometrical heat flux enhancement factor, $L_c/\Delta_T$, on the left hand side of the power balance equation is not varying significantly with $B_0$. In the case of the power scan series, obtaining probe measurements at the limiter involves the compromise of using the up/down limiters which are tangential to the last closed flux surface. This introduces the complication that $L_c$ will be a function of $\Delta_T$ and $B_0$. It may be possible to explain the stronger $n_S$ scaling with $B_0$ in W7-AS and the apparently large $\Delta_T$ inferred in the above section by a detailed consideration of the $L_c$ dependence of $\Delta_T$ and $B_0$.

Compared to ASDEX and ASDEX Upgrade density limit discharges, the Greenwald limit for W7-AS in the large aspect ratio, low beta circular approximation appears to be at least a factor of 2 greater but because of the discussed configuration considerations in W7-AS, the usual scaling may not be valid.

4. Transport simulations

The values of $n_S$ from the lithium beam and the measured radiation profiles to calculate $q_L$ were taken as inputs while a constant $\Delta_T = 1.5$ cm was assumed. The time evolution of temperature and density measured by Langmuir probes in front of the limiter in the density limit discharge at 2.5 T with 380 kW NBI input power can then be predicted and are compared to the Langmuir probe measurements in Fig.4. The general features of a rising electron density and falling electron temperature at the limiter can be reproduced and the collapse of the discharge begins as the temperature at the limiter reaches 10 to 20 eV in both the simulation and experiment. From the lithium beam measurements it is clear that the edge density rises continuously to the point of the collapse in diamagnetic energy and that a sudden increase in the density fall off length
occurs in the collapsing phase of the discharge, while Langmuir probe measurements indicate that the edge density remains constant for a further 50 ms in the collapsing phase of the discharge.

Impurity puffing experiments were conducted to increase the radiated power and reduce the power flux so that a density limit could be reached in a similar discharge at a lower edge density. In NBI discharges with 380 kW of deposited power at 2.5 T, nitrogen gas was introduced. For this impurity, the radiation rate coefficient is strongest at 10 eV and with a second maximum at 100 eV and total radiation is proportional to the impurity ion and electron density. The impurity ion density profile for a given impurity flux rate assuming a diffusion coefficient of 0.2 m²/s and an inward pinch of 5r/a m/s as calculated by STRAHL and the electron density profile measured by Thomson scattering or an 8 channel microwave interferometer allows the radiated power to be calculated. In discharges without nitrogen gas puffing, carbon was assumed to be the dominant impurity. Impurity concentrations of 15% and 5% in discharges with and without impurity puffing were necessary to reproduce the total radiated power preceding the collapse of the discharge at the density limit. In discharges with impurity puffing, toroidally separated bolometer cameras indicate that impurity gas injection produces significantly more radiation in the vicinity of the gas inlet valve and MARFE formation was observed. In discharges where the impurity gas puffing was systematically increased the density limit decreased.

5. Conclusions

Density limit discharges in power and magnetic field scans in W7-AS have been compared to the predictions of the two point model. At low magnetic field, a Δr greater than 3 cm is inferred, well above the maximum indicated by scaling studies. A stronger ns scaling with B₀ compared to JET is found. It is speculated that these two observations can be accounted for by variations in Lc as a function of Δr and B₀. Transport simulations of the time evolution of the density and temperature in front of the limiter according to the two point model were found to reproduce the experimental features of a rising density and falling temperature preceding the density limit. This implies that the limit to the maximum allowable edge density for a given q⊥ necessary to satisfy power balance, with the collapse occurring for a temperature at the limiter of 10-20 eV, is the relevant mechanism determining the density limit in W7-AS.

Helpful discussions with K.Borrass and R.Schneider are gratefully acknowledged.

6. References
**Fig 1.** Diamagnetic energy, line integrated density, radiated power and net heat flux at the plasma edge for a density limit shot. Discharge reaches density limit at 0.5 s with a sudden fall in the diamagnetic energy.

**Fig 2.** Predicted maximum edge density as a function of heat flux value at the last closed flux surface for various values of the temperature fall off length, $\Delta_T$, and $L_e = 18$ m with experimental point from the power and B scan.

**Fig 3.** Edge density at the density limit scaling with respect to the magnetic field and perpendicular heat flux with errors of +/-0.15 and +/-0.2 respectively in the power law coefficients.

**Fig 4.** Time evolution of limiter temperature and density compared to the calculated values from the two point model. At 10-20eV, the maximum allowable edge density, $n_S$, is reached and the discharge collapses.
Tomographic Reconstruction of Plasma Equilibria and MHD-Modes at WENDELSTEIN 7-AS

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I. Introduction

The usage of soft-X radiation for detection of MHD-activities in a plasma and for mapping of plasma equilibrium flux surfaces is a widely applied method on fusion devices. The specific magnetic configuration of W7-AS is the result of the optimization procedure for the stellarator design and, compared with the axial-symmetric tokamak configuration, the flux surfaces are three-dimensional and are characteristically shaped in a poloidal plane, such that a tomographic system needs an enhanced poloidal resolution. To achieve proper plasma imaging for the soft-X radiation, a small sized camera system was installed inside the vacuum vessel to overcome the technical restriction of access from the inboard side of the torus. The system consists of 10 cameras with a total sum of 320 channels and therefore a good spatial resolution in both the radial and poloidal direction is obtained. Together with a well balanced distribution over the poloidal plane, tomographic reconstruction of the soft-X emissivity distribution (6μm Be filter) can be done without additional assumptions about the topology.

In this paper we introduce the camera system and the numerical algorithm, applied for the tomographic reconstruction. Additional data processing is also presented. With the whole tomographic system, comprising the hard- and software, we show some applications to

Fig. 1: 10 camera tomographic system at W7-AS. All 320 lines of sight are drawn. Typical χ = 1/3 flux surfaces are additional plotted. The surrounding ellipse is the W7-AS vessel.
measurements of the equilibrium flux contours and MHD-activities, seen in the stellarator W7-AS.

II. The 10-camera tomographic system at W7-AS

Figure 1 shows an overview of the hardware of the tomographic system together with the idealized lines of sight of the X-ray detectors. The spatial resolution is comparable with the separation between two neighboring lines. Data acquisition is performed with a sampling rate of up to 200kHz in time windows with a typical duration of about 50ms (defined by a maximum total data amount of 8MByte per shot).

Three different algorithms for the 10-camera tomography and additional techniques to extract the relevant physical data are used, because different reconstruction problems may require an appropriate algorithm. A choice can be made between algorithms using first order regularisation [1], which tries to avoid steep gradients, a minimum Fisher-information regularisation scheme [2] and the maximum entropy [3, 4] solution.

\[ \Lambda_{\text{LinReg}} = \frac{1}{2} \chi^2 + \alpha_R R \quad R = \| \nabla \rho \|^2 \quad \text{1. order regularisation} \]

\[ \Lambda_{\text{MinFisher}} = \frac{1}{2} \chi^2 + \alpha_R I_F \quad I_F = \int \frac{g'(x)^2}{g(x)} \, dx \quad \text{min. Fisher-information} \]

\[ \Lambda_{\text{MaxEnt}} = \frac{1}{2} \chi^2 - \alpha_R S \quad S = \sum_{i=1}^{N_p} \left[ \rho_i - m_i \ln \frac{\rho_i}{m_i} \right] \quad \text{max. entropy} \]

Maximum Entropy is not a widely used algorithm, because of the long numerical calculations. On the other side, it is the only algorithm, implemented such, that the result is a real mathematical result without the need of guessing the regularisation parameter. Additionally, error bars for the reconstructed distribution are obtained by this method.

For visualisation and quantitative analysis, additional processing is necessary. This includes the implementation of different meshes and expansion functions of the emissivity to be reconstructed, singular value decomposition (SVD), filtering in the frequency-space and error calculations for the case of maximum entropy reconstructions. SVD thereby is the most important method, to extract the fluctuating part of the emissivity attributed to MHD-modes from a sequence of reconstructions. In order to check the reliability, simulated plasma radiation distributions were used.

III. Equilibrium reconstructions

Concerning equilibrium effects, quantitative measurements of the magnetic surface struc-
ture were made. Two effects have been investigated in particular: the change of the plasma shape due to toroidal currents and the effect of the Shafranov-shift, depending on the plasma beta. The minimization of the Shafranov-shift is one of the most important optimization criteria of W7-AS. The comparison between reconstructions and equilibrium calculations [5] was made for a mean plasma radius where SX-profiles have steep gradients. For both cases very good agreement was observed. Only minor differences in the reconstructions was observed in comparing the three tomographic algorithms mentioned above.

IV. MHD-activities

The physics issues addressed in this paper comprise, firstly, different kinds of MHD-instabilities including pressure- and current-driven modes and fast particle driven global Alfvén eigenmodes (GAE). An example is shown figure 2, where different modes are plotted for the case of \( m = 3 \). GAE-modes usually extend over a large part of the plasma cross section, whereas tearing modes and pressure driven modes are localized around a rational surface.

In the case of GAE modes, the reconstruction of the radial and poloidal mode structures, i.e. the determination of the poloidal mode number \((m)\) and the radial eigenfunction, a variety of peaks found in the Alfvén spectrum were analysed. Most of the peaks correspond to the lowest \((m,n)\)-modes expected from theory, but also higher \(m\)-numbers and modes with different radial structure were observed. At Wendelstein 7-AS, GAE-modes are commonly seen with neutral beam heating, and the main features are compatible with theoretical predictions, see [6]. With the new tomographic system,
more complex mode structures can be assessed. For example, figure 3 shows an \((m=9)\)-GAE mode that shows up together with a \((m=3)\) and \((m=6)\) mode (not plotted). Another example, Figure 4, shows two \((m=3)\)-GAEs: one with a radial node at about \(R_{\text{eff}}=10\text{cm}\) and the second without node are observed simultaneously. In both discharges, \(m = (3,6,9)\) and \(m = (3,3\) with radial node), the modes propagate at different frequencies, that seem to develop independently during the live time of the modes.

References


Feedback controlled radiative edge cooling experiments in the Wendelstein 7-AS stellarator


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Introduction. Radiative edge cooling by seeded impurities with appropriate radiation characteristics is widely considered as an option to protect targets in fusion experiments from thermal overloading. In devices with active pumping capability (divertors, pump limiters), noble gases (e.g. Ne) are preferred as edge radiators due to both, their favourable radiation characteristics as well as their recycling properties [1, 2]. They are not pumped by plasma facing materials and, consequently, do not build up long-term reservoirs leading to uncontrollable release in particular during long-pulse or steady state discharges. In devices without active pumping capability as W7-AS, on the other hand, the seeded impurity has to be sufficiently pumped by the walls to allow control of the concentration. As was shown in Ref. [3], nitrogen is relatively well suited. It has favourable radiation characteristics, and its capability to be pumped by the walls enables feedback control of the radiation level over typical discharge durations (=1-2s) in W7-AS. A shot-to-shot build-up of an intrinsic nitrogen reservoir was found to settle at a very low, stationary background radiation level which could be completely removed by ECRH discharges without nitrogen puff. In continuing a previous study [3], this paper reports results from a nitrogen concentration scan with improved feedback control of the nitrogen radiation levels and with extended target diagnostics.

Experimental. The study in W7-AS (R = 2 m, a = 0.18 m) was performed at B = 2.5 T and t = 0.34 with the configuration bounded by two horizontal graphite limiters at the top and bottom of an elliptic cross section. Nitrogen was injected into net current free ECRH (140 GHz, 430 kW) discharges with flat-top phases of 1.5 s at a line-averaged density \( \bar{n}_e = 4 \times 10^{19} \text{m}^{-3} \). Feedback control of the radiation levels was performed via VUV line emission (N IV, 765 Å). Compared to Ref. [3], the control spectrometer was now positioned further away from the nitrogen inlet thus actually enabling quasi-stationary radiation levels. In addition to the diagnostics mentioned in [3], limiter-integrated, poloidal Langmuir probe arrays allowed to study in particular downstream parameters in more detail. The limiters are poloidally segmented by ten tiles per limiter, each equipped with thermocouples allowing poloidally resolving target calorimetry.

Results and discussion. Stationarity within the injection phase of 0.7 s could be obtained up to a central nitrogen concentration of about 2.5% (estimated from CXRS) corresponding to a radiated power fraction of about 60% (from bolometer), Figs.1, 2. Attempts to exceed this limit lead to radiative instability and to feedback induced oscillations of the discharge parameters rather than to a complete collapse. CCD camera observations covering three of the five torus modules indicated strong shrinking of the hot plasma cross section, but did not give any evidence for MARFE formation during the radiative excursions. Decreasing the prescribed nitrogen radiation level and thus the nitrogen influx within the discharge duration to below this stability limit leads to re-establishment of stationary conditions. It was found that relatively
small nitrogen concentrations (=1%) already effectively suppress medium-Z radiation from the core which is primarily ascribed to lower impurity release from stainless steel components due to lower edge temperature.

Fig. 1: Stored energy, line-averaged density, nitrogen radiation intensity, total radiative power (bolometer) and N₂ valve control voltage versus time for a reference discharge without nitrogen and for nitrogen concentrations marginally below (= 2.5%) and above the thermal stability limit, respectively.

Under quasi-stationary conditions, the stored plasma energy degrades approximately linearly with the radiated power fraction (Fig. 2) because, due to the small size of W7-AS, the nitrogen radiation zone extends inward to r/\alpha = 0.4 (Fig. 3). The electron temperature T_e is decreased only at the profile wings, whereas the value at the centre is not affected by the nitrogen radiation. A transport analysis (from power balance, Fig. 3c) indicates slightly improved confinement rather than a degradation. The density profiles (not shown) are flat inside r/\alpha = 0.7 with steep gradients at the outside (from multi-channel interferometer and Li-beam). They are, within the error limits, not altered by the nitrogen radiation which means that the improvement of the central confinement does no coincide with density profile peaking as is often observed in tokamaks.

Fig. 2: Stored energy from diamagnetic signal W_d, total radiated power P_d from bolometer, power flow P_{ECRH} + P_d across the LCMS, and total power flow P_{limiter} onto both limiters from calorimetry versus the NIV radiative intensity from SPRED spectrometer.

Fig. 3: Radial profiles of a) the electron temperature T_e, b) the soft-x emissivity and c) the electron heat diffusivity \chi_e (from power balance) for discharges with nitrogen concentrations from zero (reference) to slightly below the stability limit.
Calorimetric measurements at the two limiters (corrected for the injection phases being shorter than the pulse lengths) show much stronger reductions of the total power onto limiters than expected from bolometry, Fig. 2. The power accountability referred to \( P_{\text{ECRH}} - P_{\text{bol}} \) (with \( P_{\text{ECRH}} \) and \( P_{\text{bol}} \) being the heating and radiated power, respectively) is about 70 - 80% without nitrogen and decreases towards higher nitrogen content. This seems to indicate increasing toroidal asymmetry of the radiation shell not registered by the bolometers. CCD camera observations support this conjecture, but due to the lack of toroidally distributed bolometer cameras there is not yet direct quantitative evidence.

Data from the limiter-integrated Langmuir probe arrays show that the downstream electron temperature \( T_{\text{ed}} \) (measured at about 2 mm outside the last closed magnetic surface, LCMS) nearly linearly decreases from about 90 eV without nitrogen to 20 eV slightly below the stability limit, whereas the downstream density \( n_{\text{ed}} \) stays approximately constant, Fig. 4. The upstream temperature \( T_{\text{eu}} \) at the LCMS is decreased from about 90 to 40 eV (from a fast reciprocating Langmuir probe close to the stagnation plane). In order to check the consistency with calorimetric data, a simple two-point model of the scrape-off layer (SOL) power balance [4, 5, 6] was applied to estimate \( T_{\text{eu}} \), \( T_{\text{ed}} \) and \( n_{\text{ed}} \) from the upstream density \( n_{\text{eu}} \) and the calorimetric power onto the central limiter tile which determines the LCMS. It includes parallel heat transport by classical parallel electron heat conduction, pressure constancy along field lines and the sheath boundary condition for the heat transfer to the target:

\[
T_{\text{eu}} = \left( \frac{T_{\text{ed}}^{7/2} + \frac{7}{4}\gamma L_c q_{\parallel} n_{\text{eu}}}{4k_0} \right)^{2/7}; \quad n_{\text{ed}} = n_{\text{eu}} T_{\text{eu}} / 2 T_{\text{ed}}; \quad q_{\parallel} = n_{\text{ed}} c_s \gamma k T_{\text{ed}}
\]

\( k, c_s, \gamma, L_c \) and \( k_0 \) are the Boltzmann factor, ion sound speed, heat transfer factor (=8), connection length and parallel heat conductivity coefficient, respectively. The parallel power flux was derived from the fitted calorimetric power onto the central limiter tile, \( q_{\parallel} = P_{\text{target}} / \lambda_{\text{avg}} w \). For the power flux decay length \( \lambda_{\parallel} \), an average value was inferred from the poloidal (representing also radial) power deposition profiles on the flat limiters, and \( w \) is the poloidal tile width (2.7 cm). The upstream density \( n_{\text{eu}} \) at the LCMS was kept fixed at \( 10^{19} \, \text{m}^{-3} \) (from upstream probe and Li-beam). The results (lines in Fig. 4) satisfactorily agree with the probe data thus indicating basic consistency.

In order to check the conditions near the stability limit towards smaller \( T_{\text{ed}} \) in somewhat more detail, the sheath boundary condition was extended to include energy losses to hydrogen, \( q_{\parallel} = n_{\text{ed}} c_s \lambda T_{\text{ed}} \) with the temperature and density dependence of \( \lambda \) fitted according to Ref. [7]. These losses could be neglected for calculating the results in Fig. 4, but become increasingly important towards smaller \( T_{\text{ed}} \). The above model was then applied to calculate densities as functions of \( T_{\text{ed}} \) for calorimetric \( q_{\parallel} \) values of the discharges without nitrogen and with the highest nitrogen.
content below the stability limit (upper and lower curve set, respectively, in Fig. 5). Fig. 5 suggests that the marginally stable discharge closely approaches to the SOL thermal instability associated with the existence of a local maximum of $n_{\text{ed}}$ versus $T_{\text{ed}}$ [6]. Considering nitrogen radiative losses at coronal equilibrium from the SOL does not significantly affect this result. The deviation of the measured $T_{\text{ed}}$ from the critical value at the maximum of the calculated upstream density is not too much of a concern because probe data generally tend to over-estimate $T_{\text{ed}}$ in the presence of strong $T_e$ parallel gradients. This SOL instability may destabilize the core radiative mantle with respect to radial shrinking as is seen in the experiments (see also Ref. [8]). However, to definitely confirm this interplay, a more refined analysis is needed which will be given elsewhere.

### Summary and conclusions

Radiative edge cooling experiments have been performed in W7-AS by feedback controlled nitrogen injection into limiter-bounded, net current-free ECRH discharges. Within injection phases of 0.7s, quasi-stationarity was obtained up to central nitrogen concentrations of about 2.5%. Bolometer data indicate radiated power fractions up to about 60% whereas calorimetric data, together with downstream probe results, indicate target load reductions by up to factors of five to six. This discrepancy may indicate asymmetries of the radiation shell not registered by the bolometers. Though the plasma stored energy degrades with increasing nitrogen content due to the small machine size, the central $T_e$ is not affected. A transport analysis shows improved confinement at $n/a < 0.3$. Exceeding the nitrogen level mentioned leads to radiative instability. Downstream data combined with model estimates suggest that this instability coincides with a SOL thermal instability. More detailed studies in particular on this latter issue including B2/EIRENE code analysis with selfconsistent treatment of impurity radiation are in preparation.

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ICRF Experiments on the Stellarator W7-AS

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INTRODUCTION

The first successful ICRF plasma heating experiments in the W-7AS stellarator [1] have demonstrated effective plasma heating and plasma sustainment, for 2nd harmonic and two-ion hybrid heating regimes without significant increase of impurity radiation. The results presented in this paper establish ICRH as an attractive method to heat and sustain the plasma in stellarators under steady-state conditions.

TECHNICAL SETUP

The RF system of W7-AS consists of two RF generators with a nominal power of 1.5 MW that are connected to two antennas. For the experiments to be described the RF power was launched with a broad four-port antenna exciting a narrow \( k_\parallel \)-spectrum around \( k_\parallel = 6 \text{m}^{-1} \) for \( \pi \)-phasing [2,3] After a faster arc protection system (full power switch-off within 20\( \mu \text{s} \) of an arc) and a section of lossy transmission line to avoid generator self-oscillations [4] had been installed more effective conditioning of the antenna became possible. Thus RF pulse lengths longer than one second with voltages in the transmission lines up to 70 kV were reached.

EXPERIMENTAL RESULTS

Plasma heating, sustainment and startup from a seed discharge were investigated for the two-ion hybrid heating scenario for hydrogen minority in a deuterium and \(^4\)helium plasma. Good absorption of the RF power radiated from the antenna was observed only if the hydrogen concentration was less or about 10%. Such low hydrogen concentrations were obtained after boronization of the vessel interior with \( B_2D_6 \) or after helium glow discharges with about 5% deuterium. The hydrogen concentration was inferred from comparison of the intensity of the \( H_\alpha \) and \( D_\alpha \) lines and comparison of the CX hydrogen and deuterium fluxes. This ratio was uncontrolled and only determined by wall recycling. Good wall conditions (obtainable only after many hundreds of plasma discharges) were imperative for keeping the hydrogen concentration sufficiently low during the ICRF experiments.

ECRH generated deuterium and helium plasmas \( (P_{ECRH} = 500 \text{ kW}, n_e(0) = 3. \times 10^{19} \text{m}^{-3}, T_e(0) = 2.5 \text{ keV}, T_D(0) = 450 \text{eV}, \iota = 0.34, W_{diam} = 5 \text{kJ}, \text{central heating}) \) were used as targets. Antenna loading approximately doubled with plasma compared to vacuum
loading. Thus half of the generator power was radiated into the plasma, the remainder was ohmically dissipated in the antenna.

Fig. 1 shows the temporal evolution of an ICRF-heated ECRH plasma. During the ICRF pulse an increase of about 25% of the diamagnetic energy was observed. The dominant part of this increase is due to an increase in deuterium temperature. Fig. 2a shows the deuterium temperature profiles with and without ICRF. An overall increase from 400 eV to 550 eV is observed. Fig. 2b shows the electron temperature profiles with and without ICRF. The strongest increase in electron temperature was observed off-axis at the approximate location of the two-ion hybrid resonance and was mainly due to direct electron heating.

Fig. 3 shows the radial profile of the RF power density absorbed by the electrons. This was evaluated from the change in slope of the ECE electron temperature and the Thomson electron density profile at the turn-on time of the RF power. The total RF power absorbed by the electrons was around 30 kW for this case.

The transient drop of the central electron temperature within the first 50 μs is not understood. It cannot be attributed to a central density increase or an increase in impurity radiation.

For RF powers up to about 400 kW the increase in diamagnetic energy followed the stellarator confinement time scaling [5]. From this one can conclude that most of the radiated power is absorbed in the plasma. For higher RF powers the diamagnetic energy increased less than expected as can be seen in Fig. 4. There the increase in diamagnetic energy is shown for different ICRF powers radiated from the antenna. The solid line is the prediction based on the confinement time scaling for W7-AS [5]. For RF powers larger than 400 kW, which corresponds to $P_{ICRF} / P_{ECRH} \approx 0.6$, a saturation in the relative increase of the diamagnetic energy is observed. Data at higher ICRF power that should be available now are necessary to clarify this point.

Bulk hydrogen temperature up to 1.2 keV was measured at the plasma center with active CX whereas the deuterium temperature was about 400 eV. Hydrogen fluxes with energies up to 45 keV were observed at an angle of about 45° to the magnetic field lines. At RF powers less than about 400 kW an increase of the tail temperature with RF power was observed. However, the tail temperature saturated at about 7 keV for RF powers greater than 400 kW.

It was possible to sustain the plasma solely with ICRF. Pulse lengths up to 1 sec (>50 energy confinement times) at RF powers of 500 kW were achieved. Fig. 5 shows the time evolution of such an ICRF plasma. The achieved plasma parameters were: $W_{diam} \leq 4.2$ kJ, $T_e(0) \leq 800$ eV, $T_D(0) \leq 500$ eV, $n_e(0) \leq 4.5 \times 10^{19}$ m⁻³, $P_{RF} \leq 700$ kW. Good wall conditions facilitated density control over the length of the discharge. After some slightly transient behavior steady-state was reached where all measured plasma parameters were constant. In particular, no increase of impurity radiation was observed.
The total radiation as determined by bolometry remained comparable to ECRH plasmas. VUV measurements showed a slow increase of iron and chromium radiation that saturated after a few hundred milliseconds in agreement with the long particle confinement times expected at these densities and power levels.

It was also possible to generate the plasma with ICRF starting from a seed plasma given by the afterglow of a discharge. The time evolution of such a discharge is shown in Fig. 6. At the start of the ICRF pulse the density of the seed plasma was less than the cut-off density for the fast wave and the antenna loading was equal to the vacuum loading. Within 100 msec of the start of the RF pulse typical ICRF-plasma parameters were recovered. The stub tuners of the system had to be set to a compromise setting to ensure sufficiently low VSWR during the whole discharge.

CONCLUSIONS

Plasma heating, sustainment and generation from a seed plasma is possible with ICRF for the two-ion hybrid heating scenario in an advanced stellarator. ICRH-sustained plasma parameters were comparable to those achieved with ECRH. The experiments showed, however, that the heating efficiency degraded rapidly if the hydrogen concentration increased beyond 20-30 %. Therefore, control of the minority concentration is of crucial importance for this heating scenario in a large aspect ratio device as the W7-AS stellarator. A viable alternative to hydrogen minority could be $^3$He, whose concentration can easily be controlled.

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FIGURE 1. Shot 37298. ICRF heating of ECRH plasma.

FIGURE 2. Electron and deuterium temperature profile for shot 37298.

FIGURE 3. Relative increase in diamagnetic energy versus the ratio of ICRF power radiated from the antenna to ECRH power.

FIGURE 4. Radial profile of RF power density absorbed by electrons.

FIGURE 5. Shot 39387. ICRF sustained plasma.

FIGURE 6. Shot 39389. ICRF generated plasma.
Role of recycling to achieve high $nT \tau_E$ in W7-AS

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Introduction. High $nT \tau_E$ at W7-AS in pure or combined NBI and ECRH discharges is only obtained for low edge density. In the case of combined heating, steep temperature gradients ($T_e'$ as well as $T_i'$) in the region of low densities in connection with strongly negative radial electric fields are observed [1]. In the bulk plasma, both energy balances (electrons and ions) and particle balance are in good agreement with the neoclassical fluxes. In the gradient region the particle fluxes are still anomalous, but are much lower than in discharges with higher edge densities and lower confinement. Central values of high confinement plasmas with e. g. 825 kW NBI and 370 kW ECRH absorbed power are $n_e = 1 \cdot 10^{20} \text{ m}^{-3}$, $T_i = 1 \text{ keV}$ and $\tau_E \approx 22 \text{ ms}$. This is by a factor of about 2 larger than the one from the International Stellarator Scaling [2].

The same relation between confinement and profile characteristics appears if one considers the dependence of confinement on rotational transform $\iota$ or density in W7-AS. The variation of the latter leads partly to remarkable transitions in confinement. This points to a crucial dependence of global confinement quality, edge profiles and recycling.

Particle Transport. The influence of recycling and particle transport is studied in ECRH deuterium discharges with moderate heating power ($\approx 450 \text{ kW}$) where either $\iota$ or the averaged line density was varied between shots. In addition, discharges with a density ramp were performed and these are discussed in the next section. The plasmas were limiter bounded ($B_s = 23 \text{ mT}$) by the symmetric inboard limiters [3], which allows particle transport analysis with defined sources. $\iota$ was in the range around $1/3$, which shows a resonance-like dependence of confinement and $\iota$ in W7-AS (upper part of Fig. 1) [4]. At first sight this behaviour might indicate a direct influence of magnetic islands on transport, but neither this nor MHD activity seem to explain the observations. Possible explanations are given on the basis of anomalous transport effects induced by rational $\iota$-surfaces and/or $\iota$-shear [5-7]. The $H_\alpha$-signal (lower part of Fig. 1), which roughly corresponds to the particle flux at the edge, shows a strong anti-correlation to $W_{\text{dia}}$. This means that the particle transport is very important in this resonant phenomenon.

Fig. 1: Diamagnetic energy $W_{\text{dia}}$ and $H_\alpha$-signal at the inboard limiter vs. central $\iota_0$. The shadowed lines mark the $\iota_0$-values of the discussed discharges with low or high confinement and density ramps.
Fig. 2: Plasma profiles in the gradient region for low and high confinement. Below (#29925, solid line, ◦) and above (#29927, dotted, △) density threshold at $\alpha = 0.338$. At $\alpha = 0.335$ (#29928, dashed, □). High $nT_{e}$ NBI discharge (#34609, dash dot, ×). Explanation see text.

Besides $\epsilon$ the density has an equivalent influence on confinement. This can be shown as follows: if at $\alpha = 0.338$ the line density is increased the confinement suddenly decreases at a certain density. Fig. 2 shows the profiles near the confinement transition. Above the density threshold the density at the limiter increases and the density profile broadens; also the particle edge flux ($H_{e}$) increases. Simultaneously the temperature significantly lowers over the entire plasma radius and as a consequence the energy content decreases by almost 20%.

If then $\epsilon$ is decreased to a high confinement value ($\alpha = 0.335$, same line density), the former state with a narrow density profile and good confinement is re-established. Therefore at this point a slight change in $\epsilon$ or $n_{e}$ have the identical effect. The comparison with the high confinement NBI heated discharge shows a nearly identical density profile at the edge ($r_{eff} \geq 12cm$) although the central density is almost twice as high. So low edge densities seem to be a necessary condition for obtaining high confinement.

The particle transport is investigated with the help of the 3D Monte-Carlo code EIRENE [8], which provides the particle sources. The adaption to the experimental conditions was done with a similar procedure as formerly [9]. From a radial 1D diffusion equation ($n_{e}$-distribution in cylindrical symmetry) radially resolved particle fluxes and an effective diffusion coefficient are obtained. In fig. 3 the diffusion coefficients $D$ of the discharges of fig. 2 are shown. It clearly can be seen that the relative density dependence is the same for all types of confinement. The $n_{e}^{-2}$ dependence found here is in contrast to earlier observations, which scale with $n_{e}^{-1}$ [10].
but the discrepancy is not understood so far. However the absolute value of \( D \) differs according to the quality of confinement: better confinement means lower \( D \). The same applies to the edge particle flux, which in the high \( nT \tau E \) NBI discharges with \( \sim 10^{21} \) part/s is very low compared to \( \sim 5 \cdot 10^{21} \) part/s in the low confinement case. A more astonishing point is that the differences in \( D \) can be resolved, if the local \( T_e \) is also taken into account. Then the diffusivities can be described by \( D \sim (n_e T_e)^{-1} \). However the implication is not clear: there could e. g. exist a mechanism, which acts both on \( n_e \) and \( T_e \) profiles, or there is some direct relation between energy and particle transport. Therefore it should carefully be interpreted as some additional parameter whose value effects \( D \).

**Density ramps** In order to investigate the dynamic behaviour of the confinement transition, discharges with positive and negative density ramps were performed. In fig. 4 the difference of a density ramp discharge with low and high confinement is shown. While in high confinement the energy content increases linearly with density, in the other case the diamagnetic energy deteriorates above the density threshold again. Both energy loss and edge flux are increased by a factor of 2 compared to high confinement. Due to the excellent wall conditions the typical time scale of a density ramp-up (\( \sim 0.5s \)) could also be obtained in ramp-down discharges. Since in both cases the transition occurs at the same line density, the density can be clearly identified as the transition parameter. Fig. 5 shows a discharge with a density ramp-down at low confinement \( t_0 \approx 0.354 \) (see fig. 1). It exhibits a very sharp transition to high confinement as the density decreases. This can be seen especially in the edge flux \( (H_0) \), which drops on a ms-scale, but also in \( n_e \) and \( T_e \)-profiles. After the transition the profiles evolve over 150 ms until a new quasi-stationary state is formed. The energy content increases (due to \( T_e \)) up to the value, which is observed in an equivalent discharge at high confinement \( \epsilon \). At the same time the density profile at the edge stays nearly constant, but the line density still decreases. This indicates the change of the density profile shape. The transition is not always as fast as in Fig. 5. A higher wall recycling e. g. smooths the transition, because it influences the edge density profile.

From the data, no clear causality between the influence of \( n_e \) and \( T_e \) can be drawn. At the transition the edge \( n_e \)-profile (by Li-beam) decreases, which marks the lower edge density with better confinement. Simultaneously the particle flux is reduced, which corresponds to the lower \( D \) in high confinement (see fig. 3). Also the sudden increase of \( T_e \) as well as the increased negative radial \( E \)-field \( (v_{pol}) \) indicate the better confinement.
To elucidate the relation between particle flux and the edge density, the radial diffusion equation in cylindrical coordinates has been solved. For the diffusion coefficient typical "experimental" diffusivities as shown in fig. 3 were used. The simple $\frac{1}{r^2}$-dependence can reproduce broad and narrow density profiles similar to the experimental ones. However, the edge fluxes were about 10% higher in the case of a narrow profile. It was not possible without additional assumptions to bring down the fluxes as found in the experiment. This corresponds to the observation of the last section, that a $D = D(n_e)$ description is not sufficient to explain the different particle confinement.

**Conclusions** A low edge density seems to be a prerequisite for high confinement in W7-AS. This applies also to the resonance-like dependence of confinement with $t$. The strong anti-correlation of edge fluxes and $W_{dia}$ shows the close connection of particle transport and this resonant phenomenon. There exists a density threshold, where a fast transition of confinement occurs. The particle transport analysis of different confinement states gives to the known $n_e$-dependence an additional $T_e$-dependence for the "experimental" diffusivities. However this should be interpreted as the influence of a parameter, which could be the temperature or a related quantity. A causal influence of $n_e$ or $T_e$ cannot be distinguished.

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Plasma Radiation with Local Impurity Injection into a Magnetic Island of W7-AS Stellarator and at the separatrix of AUG Tokamak

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Introduction

Energy dissipation by impurity radiation at the plasma edge has become the preferred solution for reducing the heat loads of divertor plates in tokamaks and stellarators. However the magnetic configuration at the plasma boundary is quite different for the proposed ITER-tokamak and the W7-X-stellarator [1]. Whereas in ITER open magnetic field lines outside the separatrix are guided to a divertor chamber being locally separated from the main plasma magnetic islands at the plasma edge are intersected by target plates in the main chamber with relatively small radial separation from the separatrix in W7-X. While impurity radiation cooling in the tokamak edge configuration has been studied already extensively, the influence of magnetic islands on the properties of the edge plasma is still quite unknown. First investigations on the stellarator W7-AS are reported in this paper. For this purpose nitrogen pulses of a few ms duration have been injected into natural islands at the plasma edge of W7-AS discharges by a reciprocating erosion probe with a radial precision of injection of about 5 mm [2]. The cylindrical probe head consists of a boronitrid casing of 5 mm diameter and length and a small electrode of titanium or carbon of 2 mm diameter and length on the front side. For comparison similar experiments with nitrogen injection were also made in the tokamak ASDEX-Upgrade. In this case a boronitrid probehead of 30 mm diameter was exposed near the separatrix on the outboard side for 10 ms several times during a discharge using the midplane manipulator.

Modelling of Nitrogen radiation from a magnetic island

The power which can be stationarily radiated by nitrogen from a magnetic island region has been studied using a simple particle transport model. A continous impurity source at the centre of a circular plasma region, representing the island with density n_e, electron temperature T_e and impurity residence time of $\tau = r^2 / D$ is considered where D is the cross field diffusion coefficient. The temporal evolution of the impurity charge states in this region is calculated from the equation
\[
d\frac{n_q}{dt} = n_e (S_q - 1 \cdot n_{q-1} - (S_q + R_q) n_q + R_q + 1 \cdot n_{q+1}) - n_q / \tau_q
\]
with \(n_q\) being the density of the impurity species, \(q\) the ionization state; \(S\) and \(R\) are the rate coefficients for ionization and recombination, respectively. The radiated power density has been calculated from the relation \(p = \Sigma L_q(T_e) n_q n_e\) in which \(L_q\) is the radiation rate function. Assuming \(n_e = 1 \times 10^{19} \text{m}^{-3}\), \(\tau = 2 \text{ mm}\) and an atomic nitrogen source strength of \(1 \times 10^{22} \text{ m}^{-3} \text{s}^{-1}\) the results shown in fig.1 are obtained. Considerable nitrogen radiation from a magnetic island can only be expected if \(T_e\) is about 15 eV and if furthermore the diffusion coefficient within the island is smaller than 0.2 m²/s. Using the values of fig.1 a radiated power from the magnetic island of a 5/9-configuration up to 400 kW are predicted in case of W7-AS.

Experiments in the stellarator W7-AS
The experiments were performed in NBI-heated discharges with a 5/9- island configuration and limiters on the inboard side [2]. The nitrogen was injected with the reciprocating probe in upstream position (bottom side). Due to the injection the electron temperature within the island of W7-AS measured by the probe itself decreased from values up to 80 eV to about 10 eV. The total radiation raised up to 90 % of the input power, however, particular enhancement of radiation from island region could not be observed. The spatial distribution of the radiation after injecting nitrogen was found to be only slightly dependent on the location of the injection (X- or O-point). Enhanced radiation mainly occurred on the inboard side. When the density was increased, plasma shrinking induced by impurity injection was observed to start at densities of about \(1 \times 10^{20} \text{m}^{-3}\). This shrinking was accompanied by partial transient plasma detachment from the inboard limiters. However, there are strong indications for poloidal asymmetries in connection with plasma shrinking. In fact, at plasma densities of \(1.6 \times 10^{20} \text{m}^{-3}\), the ion current measured by probes near the O-point decreased by more than an order of magnitude on the outboard side (upstream) and less than a factor of two in the equatorial plane on the inboard side (downstream). Strongly localized radiation zones at the inboard side (MARFEs) were induced by the impurity injection at this density. In contrast to tokamaks the radiating region is not toroidally symmetric but seems to follow a helical line within a modular section similar to the helical edge observed earlier on the outboard side of W7-AS [3]. The power load to the inboard limiters measured by thermography was reduced by less than 25 % when a Marfe was formed in front of these limiters. The close proximity of the radiating region to the limiters possibly prevent a stronger reduction of the limiter load.

Besides the radiation from nitrogen a considerable enhancement of radiation from intrinsic impurities occurred. This observation is explained by the increase of radiation with decreasing plasma edge temperature and a deeper penetration of intrinsic impurity atoms due to the low edge plasma density and temperature.
Experiments in the tokamak ASDEX-Upgrade

Nitrogen was repeatedly injected just inside the separatrix in a 2 MW neutral beam heated L-mode discharge with programmed increasing plasma density from 4 to $8 \times 10^{19} \text{m}^{-3}$. During the injections the electron temperature at the separatrix measured by ECE on the outboard side decreased to values between 30 eV and 10 eV depending on the density (see fig.3) while the line averaged plasma density was transiently affected by less than 10%. Radiation measured by the bolometer at a toroidal distance of 2.5 m from the impurity source was enhanced mainly above the X-point and on the inboard side. The fraction of the total radiated power in the main chamber to the input power reached about 0.8. Divertor detachment is indicated by the CII-signal of the divertor spectrometer at all plasma densities. The particle and the power fluxes to the target plates measured by Langmuir probes and thermography were also found to be strongly reduced during the injection (not shown here). Most remarkable is the following observation: Despite the fact that the carbon flux released from the target plates of the divertor was strongly reduced the concentration of carbon and oxygen in the main chamber increased as evidenced by the signals from the monitor for these impurities.

Modelling of the plasma boundary of AUG

In order to model the AUG-experiment described above an artificial carbon or boron ion source inside the separatrix on the outboard side was introduced in the B2/Eirene-code for a discharge with a plasma density of $5 \times 10^{19} \text{m}^{-3}$ at the separatrix and 3 MW heating power. The diffusion coefficient was assumed to be 0.2 m$^2$s$^{-1}$. With increasing strength of the artificial impurity source and decreasing edge temperature the carbon radiation has been found to shift continuously from the divertor to the X-point. Generally, when a MARFE is formed with an electron temperature of 1-2 eV finite temperature gradients exist along the separatrix. Thus the electron temperature varies between 1 and 20 eV on the inboard side whereas the value on the outboard side is still 30 eV (see fig.4). This is consistent with the observed radiation pattern. Evidently, the intrinsic impurity fluxes of carbon and oxygen produced at the inner wall shield can penetrate much deeper into the plasma at temperatures below 10 eV; i.e. the screening of the SOL-plasma due to ionization is reduced. This reduced screening may explain the experimental observation of enhanced C and O concentration in case of strong nitrogen injection into the edge plasma of AUG.

Summary:

Despite the quite different magnetic configurations in the stellarator W7-AS and the tokamak AUG the radiation pattern is similar when the electron temperature at the separatrix is decreased to values lower than 30 eV. A particular radiation from island regions at the plasma edge is not observed even after injecting nitrogen into a magnetic island of the stellarator W7-AS. The experiments in AUG have revealed that a too extensive radiation cooling should be avoided in order to prevent an increase of the intrinsic impurities caused by loss of screening.
References:


fig. 1: Calculated power density radiated during nitrogen injection into a magnetic island in dependence on the electron temperature with the diffusion coefficient as parameter.

fig. 2: 2D-plot of radiated power (bolometer) during nitrogen injection showing a MARFE on the inboard side below the equatorial plane of W7-AS.

fig. 3: Temporal traces of diagnostic signals of the AUG-discharge # 7738. From top to bottom: electron temperature at the separatrix, C-II-signal of the divertor spectrometer, C- and O-concentration inside the separatrix.

fig. 4: Electron temperature and plasma density at the separatrix in dependence on the poloidal position for a discharge with a MARFE near the X-point as calculated from B2/Eirene.
Dynamic behaviour of the H-mode edge transport barrier
in the W7-AS stellarator

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Introduction

In the stellarator W7-AS the H-mode is characterized by an edge transport barrier which
is associated with a strong reduction of magnetic and density turbulence and a sheared poloidal
rotation in the electron diamagnetic direction [1,2,3]. A specific feature of W7-AS is the
restriction of the H-mode operational range to narrow windows of the edge rotational transform
where the H-mode is obtained even at the lowest available heating power. The quiescent ELM-
free H-mode is always obtained through a phase with quasi periodic ELMs [3]. Once the
quiescent state is reached, edge profile gradients and the spectroscopically measured poloidal
plasma rotation develop on a timescale of about 20 ms while the energy confinement improves
by ΔW/W≈30%.

Influence of edge parameters on the H-mode operational range

In the low-shear stellarator W7-AS H-mode operation is obtained within narrow
windows of the edge rotational transform t_a (e.g. at t_a = 0.525±0.005 and t_a = 0.48 ) at the
lowest available heating power, 200 kW of ECRH (one gyrotron) or 340 kW of NBI (one
source), respectively. The actual power threshold might even be lower. At the onset of an ELM-
free quiescent H-mode the energy flux density across the separatrix is comparable or less to that
found for the L-H transition in tokamaks [4].

The magnetic field topology within the operational windows is characterized by a
comparatively large plasma minor radius and a plasma boundary determined by the inner
separatrix of a natural island chain (e.g. 5/10 and 5/9 ). The limiter does not disturb this
LCFS. Under these conditions two mechanisms are believed to contribute to the easy access
into the H-mode: (1) outside the LCFS the connection lengths decrease to a value of some
meters within a radial distance Δr ≈ 1 cm. This allows for the development of a strong radial
variation of the radial electric field and a corresponding velocity shear layer already before a
fully developed H-mode is achieved [5]. (2) the poloidal viscosity is lower than for other values
of t_a, as the island structures which create strongly corrugated flux surfaces increasing the
magnetic pumping are shifted out of the confinement region [2,6].
The quiescent H-mode is always reached through a phase with quasi periodic ELMs [2] with a typical repetition frequency $f_{\text{ELM}} > 1$ kHz. Between the ELMs turbulence is strongly reduced and indistinguishable from the quiescent H-mode phase. Global confinement in this ELMy H-mode is close to the L-state which only exists outside the H-mode operational range. Within the H-mode operational windows no stationary turbulent L-mode is found to precede the transition. Instead short (<1 ms) quiescent phases and sequences of periodic ELMs are observed intermittently even in the early phase of the discharge. For the onset of an ELM-free quiescent H-mode a threshold density $n_{\text{th}}$ is required which depends on the edge rotational transform as an important parameter. In Fig. 1a the operational window around $t_a = 0.525$ is marked by the hatched area and the observed values of $n_a$ are given by the squares. The numbers inserted in the figure indicate the repetition frequencies $f_{\text{ELM}}$ observed in ELMy H-modes. As the density approaches the threshold $f_{\text{ELM}}$ decreases to a minimum of 1 kHz. Stationary ELMy H-modes have been obtained below the threshold for up to 700 ms if the line averaged density was kept constant. In Fig. 1b the same range of the total edge rotational transform $t_a$ is scanned continuously within 130 ms in a single discharge by adding a small ohmic current ($0 \text{ kA} < I_{\text{tor}} < 5 \text{ kA}$). The average density is kept constant at $<n> = 4 \times 10^{19}$ m$^{-3}$. ELMs appear as $t_a$ crosses the operational range and $f_{\text{ELM}}$ approaches a minimum as the chosen density is close to the corresponding threshold density for a quiescent H-mode.

**Fig. 1a:** Operational range (hatched area) and threshold density (squares) for the quiescent H-mode in the $t_a$ window around $t_a = 0.525$. The numbers inserted indicate ELM repetition frequencies of ELMy H-modes in kHz.

**Fig. 1b:** Edge electron temperature from ECE (Fig. 1b) and $H_\alpha$ signal (Fig. 1c) as the total edge rotational transform is tuned over the operational window shown in Fig 1a by adding an increasing ohmic current. $<n>$ is kept constant at $4 \times 10^{19}$ m$^{-3}$. The time window shown lasts 130 ms.
The influence of magnetic shear

The H-mode can also be achieved if positive or negative magnetic shear is induced in addition to the vacuum configuration by positive and negative ohmic currents. An improvement of the edge confinement can be obtained by magnetic shear even outside the H-mode operational windows [7]. With increased shear the fluctuation spectra observed with reflectometry are found to broaden and shift further into the electron diamagnetic drift direction.

Dynamic behaviour of edge parameters

The ELMs within the ELMy H-mode display characteristics similar to type III ELMs found in tokamaks [8]. They show up as bursts of magnetic and density turbulence with a repetition frequency >1kHz and a typical length of 200μs followed by an interval with strongly reduced turbulence identical to that in the quiescent H-mode phase. In most cases at the onset of the ELM magnetic coils measure a quasi coherent precursor activity with a frequency around 400 kHz. Edge profiles of $T_e$ and $n_e$ are obtained with high time resolution from EC- and SX-emission, reflectometry and a 10 channel mm-wave interferometer, respectively. As the broadband magnetic and density turbulence level starts to grow the edge profiles of $T_e$ and $n_e$ flatten over the first 3 cm inside the separatrix, emphasising the edge localized character of the phenomenon and the associated loss of confinement. As soon as the level of turbulence begins to decrease edge gradients again begin to steepen.

The radial propagation of the change of density perturbations is followed with reflectometry on a shot to shot basis using the onset of the H_α-burst as reference. Over a distance from 3 cm inside to 1 cm outside the separatrix the increase of the density fluctuation level occurs within a time interval as short as 20 μs.

The poloidal propagation of density perturbations can also be measured with the reflectometer system since the antenna beams are tilted with respect to the normal of the magnetic surfaces. The measured Doppler-shift of the reflected mm-wave results from the selected poloidal wavevector component and the poloidal propagation velocity of the density perturbation [9]. As an example Fig.2 shows a frequency power spectrum of the reflected wave during an ELMy H-mode. The spectrum is scanned over 20 ms, therefore about 30 ELMs and intermitted quiescent phases are covered. Selecting only time intervals during ELMs Fig.3 shows the radially resolved Doppler-shift on a shot to shot basis. The observed Doppler-shift towards negative frequencies is due to the poloidal propagation of the density perturbations in the electron diamagnetic drift direction. The poloidal propagation velocity of fluctuations derived from the Doppler-shift depends critically on the resulting tilt angle i.e. on the details of the complex edge topology. For positions inside the separatrix region a maximum poloidal velocity of about 20 km/s is estimated. Note that the observed velocity of the turbulence structures $v_{pol}$ is the sum of poloidal plasma rotation and the intrinsic phase velocity of the turbulence itself.
Termination of the quiescent H-mode

After a quiescent H-mode of more than about 50 to 100 ms confinement tends to degrade which is accompanied by an increase of bolometry and SX radiation. During this phase strong isolated ELMs can appear as a distinct (duration less than 300μs) burst of magnetic and density turbulence associated with a huge spike in Hα-emission and a reduction of nₑ and Tₑ gradients at the edge (i.e. less than 3 cm inside the separatrix). In contrast to the quasi periodic ELMs found in the ELMy H-mode the edge gradients do not recover immediately after the event but remain flat for a typical period of several ms. During that time the decrease of nₑ and Tₑ propagates to the inner part of the plasma. In comparison to the quiescent H-phase the fluctuation level at the plasma edge in most cases remains significantly higher during this period.

References

RESONANT ELECTRON BERNSTEIN WAVE HEATING VIA
MODE CONVERSION IN W7-AS

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Abstract
Electron cyclotron resonance heating (ECRH) above the plasma cutoff density with electron Bernstein waves (EBW) was investigated and successfully demonstrated at W7-AS stellarator. The EBW's were generated via O-X-B mode conversion heating from an O-wave to an X-wave and, finally, to an electron Bernstein wave. A narrow power deposition profile could be determined from a careful analysis of the soft-X emission at the power switch-off. The deposition could be shifted by changing the position of the cyclotron resonance layer.

Introduction
ECRH with electromagnetic waves is a very efficient method to heat magnetically confined fusion plasmas. However, the accessible plasma density is limited by a critical density (cutoff density). On the other hand, the prospected large stellarator W7-X will have operational regimes above the cutoff density of the proposed ECRH heating system. A possibility to overcome the density limit is the O-X-B mode conversion process proposed by J.Preinhaelter and V.Kopecký [1] in 1973. This process is a general physics phenomenon of EC-waves propagating in hot magnetised plasmas, such as ionospheric or fusion plasmas. Here O, X, and B represent the ordinary, extraordinary and electrostatic mode, the so called electron Bernstein mode. The essential part of this scheme is the conversion of an O-wave launched by an antenna from the low field side into an X-wave at the O-wave cutoff layer. This mode conversion requires an O-wave oblique launch near an optimal angle. The transverse refractive indices $N_x$ of the O-wave and X-wave along a wave trajectory in a density gradient are connected at the optimal launch angle with a corresponding longitudinal (parallel $B_0$) index $N_{z, opt}^2 = (Y/(Y+1))$ with $Y = \omega_{ce}/\omega$ ($\omega$ is the wave frequency, $\omega_{ce}$ is the electron cyclotron frequency) without passing a region of evanescence ($N_x^2 < 0$). Therefore power can be transmitted through the plasma cutoff and a fast X-wave is generated. At the upper hybrid resonance (UHR), where the X-mode branch is connected to the electron Bernstein branch, EBW's are generated and propagate, since for EBW's no density limit exists, toward the dense plasma center. A detailed description of the O-X-B mode conversion experiments at W7-AS is found in [2]. In this paper the propagation and cyclotron absorption of the EBW's in W7-AS was investigated and successfully demonstrated.
Ray tracing

Ray tracing calculations were performed in order to get a more detailed insight into the O-X-B-scheme and to show the propagation and absorption of the EBW's. Density, temperature and magnetic field profiles similar to that of a typical neutral beam sustained W7-AS plasma were used for model calculations in a torus. We use the nonrelativistic hot dielectric tensor with a correction for electron ion collisions given by [3] and an isotropic electron temperature. The ray trajectories in the equatorial plane are shown in Fig. 1. The beam is launched from the low field side and propagates through the cutoff, where it is converted into an X-mode. Then it moves back to the UHR-layer, where the X-B-conversion takes place. The EBW’s are absorbed near the cyclotron resonance at the plasma centre. A small fraction of the beam power is lost at the UHR due to finite plasma conductivity. The power deposition zone for resonant heating strongly depends on the magnetic field and electron temperature as shown in Fig.1.

![Ray tracing results](image)

*Fig.1* Results of ray tracing calculations.

- **Left picture:** Ray trajectories of 70 GHz EBW’s in a plasma torus for different magnetic fields. The central density was $1.5 \times 10^{20} \text{ m}^{-3}$ and the central temperature was 500 eV. In the resonant case the trajectory was calculated until 99% of the power was absorbed.

- **Right picture:** Ray trajectories for different central plasma temperatures at a central magnetic field of 2.2 T and a density of $1.5 \times 10^{20} \text{ m}^{-3}$.

EBW’s experience a cutoff layer ($N \rightarrow 0$) at the UHR surface, which in the nonresonant or higher harmonic ($\omega_{ce}/\omega<1$) field totally encloses the inner plasma. The radiation is then trapped inside the plasma like in a hohlraum. The EBW’s are either reflected at the UHR surface in the case of an oblique angle of incidence or are back converted to X-waves which are converted again to the EBW’s at their next contact with the UHR. Radiation can only escape through the small angular window for O-X- and X-O-conversion, respectively. In the absence of an electron cyclotron resonance in the plasma the EBW’s may be absorbed due to finite plasma conductivity after some reflections at the UHR-layer. In the resonant case due to the nonvanishing parallel
component $N_{\parallel} = N_{z,\text{opt}}$ of the refractive index in the oblique launch the cyclotron absorption is strongly Doppler shifted.

**Experimental set-up**

The experiments were performed at the W7-AS stellarator (major radius $R=2.0\text{m}$, minor radius $a=0.18\text{m}$) with two 70 GHz gyrotrons with 170 KW power each. A detailed description of W7-AS and its 70 GHz ECRH system can be found in [4]. The central magnetic field was set between 2.0 and 2.5T and the edge rotational transform $\nu$, taken from the magnetic reconstruction, near 0.35 according to the experimental requirements. The central density of the neutral beam injection (NBI) sustained target plasma was up to $1.5 \times 10^{20} \text{m}^{-3}$, which is more than twice the 70 GHz O-mode cutoff density. The NBI power was 360 KW in co-injection. At the ECRH launch position the stellarator plasma has an elliptical shape similar to D-shape tokamak. In the equatorial plane the magnetic field as a function of the effective radius $r_{\text{eff}}$ is approximately given by the following relation: $B(r_{\text{eff}}) = B_0 A / (A - r_{\text{eff}} / a)$ with $A=10.5$. The power deposition was estimated from the change of the temperature profile at the power switch-off. Since the density was far above the ECE cutoff, the temperature profiles were calculated from the soft-X emission and the Thomson scattering diagnostic. The central temperature was 500 eV. The X-ray emission was monitored by an array of 36 silicon detectors with a 25 $\mu\text{m}$ beryllium filter. To obtain the radial X-ray emission profile the signals were inverted to the magnetic flux co-ordinates. The time resolution was 0.1 ms and the radial resolution was about 1cm.

**Experimental results**

The soft-X intensity $I_{sx}$ at the detector is approximately given by the following formula:

$$ I_{sx}(T_e) = c_{\text{fit}} n_e^2 Z_{\text{eff}}^2 \sqrt{T_e} \exp\left(-\frac{E_{\text{filter}}}{T_e}\right), $$

where $n_e$ is the electron density, which we take from the Thomson scattering diagnostic, $T_e$ is the temperature, and $E_{\text{filter}}$ is the cutoff energy of the filter (1KeV for 25 $\mu\text{m}$ Be). $Z_{\text{eff}}$ is assumed to be constant over the radius. At the time of Thomson scattering the linearity of this relation was checked with the Thomson temperature over the full temperature range and a proportionality factor $c_{\text{fit}}$ was estimated. Assuming that $n_e$ and $Z_{\text{eff}}$ do not change during the ECRH power switch-off the temperature difference is then

$$ \Delta T_e = \left(\frac{dI_{sx}}{dT_e}\right)^{-1} \Delta I_{sx}. $$

Since for small time scales the radial heat transport is low, the difference of two soft X-ray profiles with $\Delta t=3\text{ms}$ immediately after the ECRH switch-off represents the upper limit of the relative ECRH absorption profile. The local power deposition was estimated from the heat wave at the power switch-off taken into account the radial symmetry of the temperature profile. In Fig.2 the absorption profiles for different magnetic fields are shown. The absorption is strongly Doppler shifted due to the oblique launch and moves from the high field side at 2.0T.
through the center (2.2T) to the low field side at 2.3T with increasing magnetic field, which clearly demonstrates the propagation and the local cyclotron absorption of the EBW's for the first time. In comparison with the ray tracing results the experimental absorption is more Doppler shifted than the calculated one, which may be related to an only rough modelling of the stellarator magnetic field geometry in the ray tracing code.

![Graph showing changes of temperature 3ms after O-X-B heating switch-off and the related ECRH absorption profiles different central magnetic fields.](image)

**Fig.2**

**Conclusions**
ECRH of an overdense plasma with 70 GHz electron Bernstein waves was clearly demonstrated in W7-AS. The EBW's were generated via mode conversion in the O-X-B process. The position of the narrow absorption profile, estimated from the soft-X emission, could be changed by a shift of the cyclotron resonance layer. Thus the propagation and the local resonant cyclotron absorption of EBW's was shown, which is an excellent test of hot plasma wave theory.

**References**
Analysis of D Pellet Injection Experiments in the W7-AS Stellarator

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A centrifugal injector was used to inject deuterium pellets (with $3-5 \times 10^{19}$ atoms) at =600 m/s into currentless, nearly shearless plasmas in the Wendelstein 7-AS (W7-AS) stellarator. The D pellet was injected horizontally at a location where the noncircular and nonaxisymmetric plasma cross section is nearly triangular (Fig. 1.) Visible-light TV pictures usually showed the pellet as a single ablating mass in the plasma, although the pellet occasionally broke in two or splintered into a cloud of small particles.

Density Evolution Following Pellet Injection in W7-AS

Figure 2 shows Thomson scattering profiles for the electron density $n_e$ and temperature $T_e$ immediately before and at times shortly after pellet injection (at 0.4 ms, 0.5 ms, 0.65 ms, 5.7 ms, 7.6 ms, and 18.5 ms) for a neutral-beam-heated plasma. Pellet injection leads to a rapid ($< 400$-$\mu$s) rise in $n_e$ in the main part of the plasma and to a slow increase in the plasma edge on a longer (>20-ms) timescale. Microwave reflectometer measurements of the time delay (which reflects the density gradient) and FIR interferometer measurements of the line-average density also indicate that the density rises in $<500$ $\mu$s (Fig. 3) and then remains unchanged. The same behavior is seen in the soft X-ray intensity profiles in Fig. 4 (with 1.4 ms between curves); this signal is more influenced by density than by temperature because there was no
Fig. 3. Time delay and fluctuation level from a microwave reflectometer, MHD oscillations from a Mirnov coil, and line-average density around the time of pellet injection.

Fig. 4. Soft X-ray intensity profiles from 6.5 ms before pellet injection to 8.4 ms after a pellet filter in front of the soft X-ray diode array. Neutral lithium beam measurements of \( n_e \) at radii \( r > 12.5 \text{ cm} \) (Fig. 5) show that the outer density rises slowly after pellet injection (on a 60-ms timescale), presumably due to recycling from the wall. The increase in \( n_e \) decreases with distance into the plasma.

Pellet ablation calculations using the standard neutral gas shielding model and the \( n_e(r) \) and \( T_e(r) \) profiles before injection indicate that the increment in density should be peaked. However, this is not seen in W7-AS on the timescale of any of the measurements. The change in the \( n_e(r) \) profile in Fig. 2 is not diffusive. If it were, then the particle diffusivity would have to increase by a factor of >100 for a relatively short time (<400 \( \mu \text{s} \), about the transit time for the pellet through the plasma). There was not a significant loss of injected particles during the fast density change: comparing the number of particles in the pellet with the increase in the number in the plasma after pellet injection indicates that >60% of the pellet is retained in the plasma for times >20 ms. This behavior, where the density profile after pellet injection assumes the shape of the preexisting profile without significant loss, is not seen in tokamaks.

Either the ablation model is not correct and the deposition profile is the same as the initial density profile for different initial densities and field values, or a fast density redistribution occurs in these experiments in W7-AS.

Fig. 5. Increase in the plasma density in the edge region from 30 ms before to 70 ms after a pellet.
Effect of Pellet Injection on Energy Confinement and Fluctuations in W7-AS

Figure 2 shows a rapid drop in $T_e$ after pellet injection because the stored energy $W$ does not change immediately. Although the $n_e(r)$ profile remains at the new (higher) value for the duration of the discharge following pellet injection, $T_e(r)$ gradually returns toward its preinjection value, indicating an increase in the energy confinement time $\tau_E$. This is illustrated in Fig. 6 where $W$ and $\tau_E$ increased by a factor of 2.7 for a factor of 1.7 increase in $n_e$. From the W7-AS $\tau_E$ scaling expression [1] $W = P \tau_E^{W7-AS} \propto n_e^{0.5}p^{0.46}$, the increase in $n_e$ should lead to a factor of $\approx 1.3$ increase in the stored energy because $P = P_{\text{absorbed}}$ is approximately constant at the densities in these experiments. The additional factor of 2 improvement in $\tau_E$ is due to an increase in confinement: $\tau_E/\tau_E^{W7-AS}$ is 0.7 before the pellet and 1.4 afterward. In Fig. 2, the confinement improvement (1.35) is less: $\tau_E/\tau_E^{W7-AS}$ was 0.8 before the pellet and 1.1 afterward.

A few-ms quiescent period is sometimes seen in the MHD signals on Mirnov coils following pellet injection when the stored energy starts to increase, as seen in Fig. 3. In all cases, the Mirnov loop signal (often associated with degraded confinement) is reduced, and $\tau_E/\tau_E^{W7-AS}$ is $\approx 1$ after a pellet. Low density fluctuation levels, followed by a decaying high-frequency burst in several discharges, were seen with the microwave reflectometer after the rapid density rise (Fig. 3). These oscillations are correlated with the Mirnov coil signal. An example of pellet injection suppressing MHD activity for a longer period with a resultant increase in $\tau_E$ is shown in Fig. 7. Without the pellet, the MHD activity and the associated enhanced energy loss persist for a longer time. It is thought that the pellet rapidly depletes the fast particle population that drives the instability [2].

![Fig. 6. Response of the plasma to pellet injection at $t = 300$ ms.]

![Fig. 7. MHD activity that retards transition to a high-$\beta$ phase (2 top traces) is suppressed by pellet injection (2 bottom traces).]
The soft X-ray (and MHD loop) signals in Fig. 8 show short-duration oscillations (for \(\approx 300 \mu s\)) immediately after pellet injection, followed by slow (650-Hz) oscillations that damp out after a few ms and are in phase across the plasma \((m = 0)\). The phase velocity of the slow oscillations indicates a disturbance propagating outward at 350 m/s. The initial burst is coincident with the fast change in \(n_e\) in the bulk of the plasma seen in Fig. 2. If the rapid change in \(n_e(r)\) is due to a fast density relaxation rather than an anomalous particle deposition profile, then this burst may be evidence of an instability that causes the fast radial spreading of the high-\(\beta\) pellet blob. The magnetic field has low shear, a relatively weak radial gradient along the pellet path, and stronger field regions some distance toroidally on either side, which may allow a radial mixing to occur. The process may be different than in tokamaks because of the low \(B \times \nabla B\) radial drift in W7-AS. In earlier experiments in W7-AS [3], the density profile became more peaked, and soft X-ray tomographic reconstruction showed small-amplitude (=2-mm) internal oscillations in the location of the peak of the soft X-ray emission following pellet injection. In that case, the soft X-ray oscillations showed a different \((m = 1)\) mode structure. Experiments with higher shear and the opposite toroidal variation of the magnetic field strength are needed to understand these observations.

The next steps involve use of a one-dimensional transport code to model the evolution of the density profile and experiments with a more versatile gas gun that allows injection with a larger range of pellet masses. Improved diagnostics — Thomson scattering triggered by the pellet, fast multichannel ECE and microwave and FIR interferometers, a Cotton-Mouton-effect polarimeter, an extensive soft X-ray array that surrounds the plasma poloidally, improved MHD loop arrays — should allow better study of the evolution of \(n_e(r)\) and the associated transport.

References


The Neoclassical “Electron-Root” Feature in W7-AS

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Introduction: The confinement properties of stellarators in the long mean free path regime are determined by the radial component of the $\nabla B$-drift of particles trapped in local ripples. In the worst case, the unfavourable neoclassical temperature dependence of the heat flux, $q \propto T^{3/2}$, leads to rather poor confinement properties. Optimisation of the magnetic configuration - e.g., as performed for W7-X - can significantly reduce the averaged radial drift of these localized particles leading to an essential confinement improvement. A sufficiently strong radial electric field (the poloidal $E \times B$-drift) will force the trapped particle orbits close to the flux surfaces also leading to significantly reduced transport. With the neoclassical transport coefficients in the $Imfp$ regime depending on $E_r$, multiple roots of the ambipolarity condition, $\Gamma_e = Z_i\Gamma_i$, may exist: the “ion root” solution with weak $E_r$ for which only the ion transport coefficients are reduced, and the “electron root” solution at strongly positive $E_r$ with additionally improved electron confinement. For classical stellarators (without optimisation of the magnetic configuration) the “electron root” scenarios are mandatory to obtain acceptable confinement properties. Furthermore, operation with the positive $E_r$ may turn out to be essential for preventing neoclassical impurity accumulation.

Experimental Findings: Recently, strongly positive radial electric fields have been measured at W7-AS in low density discharges at high ECRH power level ($\geq$ 400 kW) with 2nd harmonic X-mode (140 GHz). The electron temperature profiles are highly peaked (with $T_e(0)$ up to 4 keV), and the ion temperatures (with $T_i$ of several 100 eV) fairly flat, see Fig. 1. The density profiles are flat or even slightly hollow. The finding of the strongly positive $E_r$ is related to an additional peaking of the central $T_e$ profile indicating improved electron energy confinement. The corresponding experimental heat diffusivity, $\chi_e$, from the power balance is much lower than the neoclassical one for $E_r \simeq 0$. $E_r$ simulations based on the neoclassical ambipolarity condition with only thermal fluxes taken into account predict only the “electron root” in the inner plasma region. The predicted neoclassical $\chi_e$ with these $E_r$ are, however, smaller than the experimental ones.

This “electron root” feature at sufficient ECRH power is only found for W7-AS configurations where a significant fraction of the ECRH power at 2nd harmonic X-mode is absorbed by ripple trapped electrons close to the magnetic axis. For 70 GHz O-mode launching, an “electron root” feature was not observed, so far [1]. Equivalent experiments in a configuration without trapped electrons in the ECRH launching plane neither show these strongly positive $E_r$ nor the additional peaking of the $T_e$ profile. In spite of the fact that this specific configuration is a neoclassically improved one, the central $T_e$ are lower than the ones of the “electron root” feature at higher power levels. A strong indication, that the ECRH driven electron flux (related to the generation of suprathermal electrons as shown by bounce-averaged Fokker-Planck calculations [2]) is responsible for the “electron root” feature, is found from the ECE temperature measurements after the ECRH is switched off, see Fig. 2. For the configuration with significant trapped particles in the launching plane (even more than in the “standard” configuration of Fig. 1), the decay of the central $T_e$ is characterized by two different time scales. After a very fast decay (within less than 1 ms immediately after switch-off), the central $T_e$ relaxes on a time scale similar to that slightly outside of the “electron root” region. For the “neoclassically improved” configuration, the initial fast decay is not found, and the central $T_e$ is lower although the confinement (reflected by the time scale of $T_e(r,t)$) is clearly higher. Furthermore, the confinement
Fig. 1. Temperature and density profiles for an ECRH discharge (X-mode at 140 GHz with 400 kW) in the W7-AS "standard" configuration. The $E_r$ profile is simulated by means of the ambipolarity condition with only the thermal neoclassical fluxes taken into account (experimental $E_r$ data from active CXRS, lower left plot). The electron heat diffusivities (lower right plot): the neoclassical ones with the simulated $E_r$ ($\times$), and with $E_r = 0$ (dotted line); the $\chi_e$ from power balance (dot-dashed line) is given for reference.

properties of high energetic electrons found in the strongly down-shifted ECE channels (with small reabsorption by the plasma at outer radii on low-field side) also show the fast response, depending on the magnetic configuration.

Monte Carlo Simulations: The "convective" contribution of the ECRH driven electron flux is estimated by Monte-Carlo simulations in 5D phase space [3]. The quasi-linear diffusion term describing the ECRH in the Fokker-Planck equation is approximated by an explicit source term, $\nabla_v \cdot (D \cdot \nabla_v f_{\text{Max}})$. The quasi-linear diffusion coefficient, $D_{\perp \perp}(v_\|, v_\perp)$, is obtained from ray-tracing calculations for the different heating scenarios. In this linear approach, the driven electron flux is proportional to the heating power. Quasi-linear degradation effects at higher ECRH power can only be treated by means of the bounce-averaged Fokker-Planck code [2]. The effect of the different magnetic configurations, however, is completely taken into account.

Fig. 2. ECE electron temperatures after the ECRH is switched off for ECRH discharges (400 kW, X-mode at 140 GHz) with $n_e \simeq 1.5 \times 10^{19} \text{ m}^{-3}$ for configurations with slightly increased (solid lines) and nearly without ripple (neoclassically improved, dotted lines) at the ECRH launching: central, intermediate (from high field side) and outer channels (from low field side).

The electron fluxes driven directly by the ECRH are shown in Fig. 3 in comparison with the neoclassical ambipolar fluxes for the discharge parameters of Fig. 1. These ECRH driven fluxes are significantly decreased by $E_r$, but less than the neoclassical ones at the transition region to the "electron root". Consequently, these "convective" electron fluxes dominate at high $E_r$ in the ambipolarity condition. Within the traditional neoclassical transport theory (which is
the basis for the DKES code [4] used at W7-AS), this discrepancy cannot be resolved since all fluxes are decreased with increasing $E_r$ (e.g., the electron transport coefficients in the $\sqrt{r}$ regime scale with $E_r^{-3/2}$ which is the case in the “electron root” region). Due to the higher absorption for the X-mode scenarios, the $D_{\perp\perp}(v_{||},v_\perp)$ has a maximum closer to the thermal bulk than for the O-mode case. In case of X-mode, however, the power is mainly absorbed by deeply trapped electrons whereas in the O-mode case mainly by barely trapped electrons. In addition to the $\nu v_B$ drift being proportional to energy, the energy dependence of the collisional detrapping leads to the broadening of the electron fluxes in the O-mode case. These results are well in agreement with the findings from the “effective power deposition” profiles by the electron heat wave propagation analysis [2].

![Graph showing ECRH driven fluxes from Monte Carlo simulation for $n_e \approx 2 \times 10^{19} \text{ m}^{-3}$, $T_e \approx 3.5 \text{ keV}$, and 400 kW heating power in the W7-AS “standard” configuration: with $E_r = 0$ (X-mode: solid, and O-mode: dot-dashed line), and with $E_r$ corresponding to Fig. 1 (X-mode: dashed, and O-mode: dotted line). The ambipolar neoclassical fluxes (purely thermal) are given for reference (open circles: unstable root).]

Effects of strong poloidal rotation: The strongly positive radial electric field within the region of the “electron root” feature leads to complex effects on the flux surfaces. In the continuity equation, the inhomogeneity of $B$ with $\nabla \cdot \mathbf{E}_{\times B} \approx -2 \mathbf{E}_{\times B} \cdot \nabla \ln B$ drives a parallel flow (which is the equivalent to the Pfirsch-Schlüter current) as well as a (1st order) density inhomogeneity, $n_1$. In the ion force balance equation, the $(\nu \cdot \nabla) \nu$ term is approximately given by the gradient of the kinetic energy related to the poloidal rotation as well as the parallel flow. Viscous damping of this parallel flow by, e.g., ripple trapped particles leads to an enhanced density disturbance $n_1$ which, in turn, has to be compensated by a (1st order) potential, $\Phi_1$, in the electron force balance equation. These effects can play an essential role mainly in the poloidal component of the force balance, and an external energy flux seems to be necessary to drive a strong poloidal rotation, i.e., the “electron root” feature.

On the other hand, the strong $E_r$ leads also to modifications in the neoclassical transport. In the “usual” ordering, the $E \times B$ drift is assumed to be constant on flux surfaces (and terms of the order of $(B \times \nabla) \cdot E_r$ are neglected to allow for a mono-energetic solution). The only mode coupling in the Fourier expansion of the drift-kinetic equations appears due to the mirror term, $B \cdot \nabla B \partial f_1/\partial p$ (with $f_1$ the 1st order distribution function, and $p = v_{||}/\nu$), leading to all the neoclassical transport in the $lmfp$ regimes. In addition to the $B \cdot \nabla \Phi_1 \partial f_1/\partial p$ in the self-consistent approach [5], also the $\mathbf{E}_{\times B} \cdot \nabla f_1$ term leads to mode coupling, and, as a consequence, to a 1st order density and potential, $n_1$ and $\Phi_1$, respectively. As this $\mathbf{E}_{\times B}$ is partly counteracted by the $v_{||}$ term, this effect will be dominant in the ion drift-kinetic equation, whereas the 1st order potential affects both equations. Although no particle fluxes are driven by the $\Phi_1$ with respect to (the 0th order) $f_0$, the energy flux is directly affected. With respect to $f_1$ (which leads to the usual neoclassical transport), the strong $E_r$ will modify the particle and energy transport both for electron and ions in this self-consistent approach.

The bifurcation in the radial electric field: After flux-surface averaging of the poloidal component of the combined electron and ion force balance equation, the contributions of the pressure as well as of the kinetic energy (from rotation and parallel return flow) vanish. In this context, only the neoclassical prediction of an “ion/electron root” transition based on the local ambipolarity condition is analyzed. With the shear viscosity in the divergence of the pressure
tensor taken into account, the transition between both roots is smoothed out leading to a local violation of the neoclassical ambipolarity condition, i.e., to a radial current, \( e(Z_i \Gamma_i - \Gamma_e) \), of different sign on both sides of the poloidal rotation shear layer. With the \( \partial \nu / \partial t \) term from ion inertia, the (time dependent) diffusion equation equivalent to Ref. [1] is obtained (for simplicity, \( \partial \eta / \partial E_r = 0 \) is assumed here for the normalized shear viscosity coefficient \( \eta \))

\[
\frac{nm}{B^2} \frac{\partial}{\partial t} E_r = \frac{2}{r^2} \frac{\partial}{\partial r} r^2 \eta \left( \frac{\partial E_r}{\partial r} - \frac{E_r}{r} \right) - e(Z_i \Gamma_i - \Gamma_e).
\]

Under the assumption of a stationary \( E_r \), this diffusion equation leads for small \( \eta \) to the shear layer position, \( r_{SL} \), within the region of several roots, see the solid line in the \( E_r \) plot in Fig. 1. With the current density, \( e(\Gamma_i - \Gamma_e) \), roughly of the form \( \propto (E_r - E_r^i)(E_r - E_r^{un})(E_r - E_r^e) \) where the index refers to the “electron”, an unstable, and the “ion root”, the diffusion eq. for \( E_r \) leads to a “bifurcation” problem which is quite similar to 1st order phase transitions in non-linear thermodynamics, e.g., see Refs. [6,7]. Within the poloidal rotation shear layer, the rotation energy is dissipated by viscosity leading to a damping mechanism adding to the one discussed in the previous section. Here, however, only the radial motion of \( r_{SL} \) can be analyzed (i.e., the flux-surface averaged force balance has one degree of freedom).

For the special scenarios with \( T_i \ll T_e \) and the pronounced \( lmfp \) transport in W7-AS at the inner radii (here, the “effective” helical ripple is mainly determined by the toroidal mirror terms) the region with multiple roots of the ambipolarity condition is radially restricted. The thermodynamic arguments can only be applied in this region for estimating \( r_{SL} \), i.e., in this picture the viscous damping of the poloidal rotation within the shear layer cannot be the main reason for the total disappearance of the “electron root”. Then, in both the transition and the “electron root” region, the crucial point is related to the “current dependence” on \( E_r \) within the “electron root” region, i.e., the neoclassical \( Z_i \Gamma_i(E_r) - \Gamma_e(E_r) \) dependence.

Conclusions: The “electron root” feature found at W7-AS with strongly positive radial electric fields is driven by the ripple-trapped suprathermal electrons generated by the ECRH. After switching-off the heating, the “electron root” feature disappears nearly immediately, i.e., on the same time scale as this ECRH driven flux. Monte Carlo simulations in 5D phase space clearly indicate that the additional “convective” electron fluxes are roughly of the same order of the ambipolar neoclassical prediction for the “ion root” at much lower \( E_r \). For the predicted “electron root” \( E_r \), the ambipolar fluxes predicted by the traditional neoclassical ordering are much too small. These findings indicate strongly, that the traditional neoclassical theory (at least for the ions) has to be checked in case of very strong radial electric fields.

Due to the variation of the magnetic field strength, a strong poloidal plasma rotation drives also a Pfirsch-Schlüter-like parallel flow on the flux surfaces, and, as a consequence, density and potential variations. This parallel flow velocities can give a significant contribution to the inertia term in the ion force balance. Although experimental indications of instabilities affecting the central confinement properties are missing so far, the “free energy” related to the strong poloidal plasma rotation as well as the parallel return flow may also be responsible for an “electrostatic” instability which superposes the neoclassical transport in the central region. All these features of the “electron root” at W7-AS, however, cannot be extrapolated to the high density scenarios with \( T_i \cong T_e \) in next generation stellarators. One essential aspect is the questionable reliability of the traditional neoclassical transport predictions for very strong radial electric fields as obtained in the “electron root” scenarios.

References
Simulation and analysis of neutral particle spectra from W7-AS in combination with neutron activation measurements

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For the investigation of confinement of fast ions we present results of an experimental and theoretical study at W7-AS. It consists of the analysis of non-thermal neutral particle spectra and absolute neutron yield measurements combined with theoretical calculations.

The fast ion distribution function is measured with the neutral particle diagnostics. It consists of four neutral particle energy analysers for spatially resolved measurements. The lines-of-sight of the analysers cross a neutral particle diagnostic beam which is passing vertically through the plasma centre. For absolute neutron yield measurement indium activation samples are used which are irradiated inside the vacuum vessel near the plasma. Their γ-activation is measured with a calibrated germanium detector. Using a fluence factor which is calculated with the Monte Carlo code MCNP, the volume integrated absolute neutron yield of the plasma discharge is determined. For the numerical simulation of the neutral particle spectra and the neutron yield the time-dependent 2D Fokker-Planck code NRFPS is used.

The local charge-exchange particle flux \( S(v, \mu, t) \) from the plasma is given by

\[
S(v, \mu, t) = g n_i n_0 f(v, \mu, t) \langle \sigma v \rangle_{\text{CX}} \cdot \exp \left[ - \int_0^1 \left( n_i \langle \sigma v \rangle_{\text{CX}} + n_e \langle \sigma v \rangle_e + n_i \langle \sigma v \rangle_i \right) \right] .
\]  \( (1) \)

Here, \( f(v, \mu, t) \) is the ion velocity distribution, \( v \) the particle velocity, \( \mu = v_i/v_\perp \) the pitch-angle, \( g \) a geometrical factor, \( n_i \) and \( n_0 \) are the ion and neutral particle densities, \( \langle \sigma v \rangle_{\text{CX}} \) and \( \langle \sigma v \rangle_e \) are the rate coefficients for charge-exchange and for electron impact ionisation. The exponential factor describes the absorption of the particles on their trajectory through the plasma and is calculated with respect to the real stellarator geometry.

The local neutron rate \( Q \) from the plasma is given by

\[
Q(t) = \frac{n_i n_0}{1 + \delta_v} \int f_i(v, t) f_j(v', t) \sigma(|v - v'|) |v - v'| d^3v d^3v',
\]  \( (2) \)
where \( f_i(v,t) \) and \( f_j(v',t) \) are the velocity distributions for the ion species \( i \) and \( j \), \( |v - v'| \) the relative velocity and \( \sigma \) the fusion cross section, respectively.

Measurements and calculations have been carried out for a variety of W7-AS discharges with \( \text{H}^0 \) and \( \text{D}^0 \)-injection with average electron densities ranging from \( 2.8 \times 10^{19} \) to \( 7 \times 10^{19} \text{ m}^{-3} \), electron temperatures from 0.5 to 2.8 keV and ion temperatures of about 0.32 to 1 keV. Electron temperatures and densities were taken from the Thomson scattering and ion temperatures from the neutral particle diagnostics. The deposition profile for the injected particles was calculated with the 3D FAFNER code.

**Fig. 1:** Comparison of calculated neutral particle flux with experimental data.

1a: \( n_e = 7.3 \times 10^{19} \text{ m}^{-3} \)

\[ T_e = 0.5 \text{ keV} \]

\[ T_i = 0.44 \text{ keV} \]

\[ \tau_w = 3.5 \text{ ms} \]

1b: \( n_e = 3.7 \times 10^{19} \text{ m}^{-3} \)

\[ T_e = 2.8 \text{ keV} \]

\[ T_i = 0.83 \text{ keV} \]

\[ \tau_w = 52 \text{ ms} \]

Fig. 1 shows as an example the calculated and measured neutral particle flux for two different plasma discharges. For the high density discharge the measurement and the calculation agree rather well. However, discrepancies occur for lower densities. The smaller the density becomes in a discharge, the larger becomes the neutral gas density and the higher the electron temperature. Increasing electron temperature and decreasing density result both in an enlargement of the energy relaxation time of the fast particles. As shown in fig.2 the discrepancies increase strongly with increasing classical energy relaxation time \( \tau_w \).

In principle, the observed discrepancies can be caused by three different processes, namely by reduced particle deposition owing to the high neutral gas density, and by particle or energy loss mechanisms with time scales comparable to the classical relaxation times.
Fig. 2: Comparison between calculated and measured fast particle density and between calculated and measured neutron yield in dependence of $\tau_w$.

These three processes will influence the time behaviour of the plasma signals in a different way. First of all it is of interest to investigate the initial rise of the plasma signals at the onset of the injection. To that end it is sufficient to use the relaxation time model which takes into account only the relaxation of energy $dW/dt = -W/\tau_w^*$ and leads to the simple approximation

$$ f(W) = \frac{dn(W)}{dW} = \frac{\tau_w^*}{W} \left( S - \frac{n(W)}{\tau_n} \right) $$

for the energy distribution function of fast particles. Here, $\tau_w^*$ is the total energy relaxation time, resulting from classical and anomalous energy losses ($1/\tau_w^* = 1/\tau_w^c + 1/\tau_w^a$). $S$ is the density of injected particles per second and $\tau_n$ the time constant of direct particle losses. From this expression follows that the initial rises of the plasma energy $E_I$ and of the neutron-rate $Q$ at the onset of the neutral beam injection are solely determined by the particle injection properties:

$$ \left. \frac{dE_I(t)}{dt} \right|_{t=0} = SW_0, \quad \left. \frac{dQ(t)}{dt} \right|_{t=0} = S n_e \sigma(W_0) \sqrt{\frac{2W_0}{m}}. $$

Here, $W_0$ is the injection energy, $m$ the mass of the injected ions. An example for the initial rise of the plasma energy and the neutron signal is shown in fig. 3. The dotted lines indicate the theoretical rise according to eq. 4. Agreement with the measured signals is obvious for the high density discharge but - owing to the statistical fluctuations in the signals - not as clear for the low density one.
Fig. 3: Initial rise of the plasma energy and the neutron signal.

3a: $n_n=5.7 \times 10^{19}$ m$^{-3}$
$T_e=1.2$ keV
$\tau_W = 21$ ms

Fig. 3b $n_n=3.9 \times 10^{19}$ m$^{-3}$
$T_e=2.7$ keV
$\tau_W = 50$ ms

Although further investigations are required this may be an indication that the observed discrepancies are not caused by a reduced deposition but by a loss mechanism. Moreover, from the FAFNER calculations it follows that the neutral gas density at the plasma edge has to be at the level of $2.5 \times 10^{16}$ cm$^{-3}$ in order to explain the observed effects. This seems to be a rather high value. On the other hand, a direct particle loss owing to radial diffusion would require a diffusion coefficient in the order of 1 m$^2$/s, this may be a reasonable value. Finally, up to now we have not detected plasma mode activities which correlate with the observed discrepancies and thus could lead to an anomalous energy relaxation by ion-plasmon interactions. Thus, in conclusion, though we can presently not exclude definitely the other two processes, it seems that enhanced radial diffusion of the fast ions may be the reason for the observed reduction in fast particle density and neutron emission in low density discharges.
The Shear Alfvén Continuum of an Ideal MHD Equilibrium without Spatial Symmetry

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The shear Alfvén continuum in an asymmetric plasma configuration that is an exact solution of the equilibrium equations of ideal magnetohydrodynamics is treated. It is found that the Alfvén continuum has two components: a continuous component that is characterized by modes defined over the entire magnetic surface, and a discrete component characterized by modes localized with a finite decay length on specific magnetic field lines. The localized modes, the nonsymmetry induced Alfvén eigenmodes (NAE), do not occur in symmetric plasmas.

The shear Alfvén continuum of ideal magnetohydrodynamics has received considerable interest in the past because of its potential importance for plasma heating and current drive, and because of its possible effects on plasma stability. The key feature of the shear Alfvén wave in the continuum is its spatial singularity about surfaces of constant pressure. What is known about the Alfvén continuum has arisen principally from studies of Alfvén wave phenomena in MHD equilibria with spatial symmetries, such as the one-dimensional cylindrical screw pinch and the two dimensional toroidal tokamak. Equilibria without symmetries have received minimal attention, in part because the existence of asymmetric MHD equilibria has not been clearly established. Recently, however, several classes of asymmetric ideal MHD equilibria have been found [1]. These particular equilibria are parallel to a straight, infinite, magnetic axis (z axis). The magnetic lines of force twist about the axis and form closed magnetic surfaces. In the limit of plasma $\beta \ll 1$ the Cartesian components of the magnetic field $\mathbf{B}$ are

$$B_x = -B_0 e^{\kappa z} \sin(\kappa x) / 2, \quad B_y = B_0 e^{-\kappa z} \sin(\kappa y) / 2,$$

$$B_z = B_0 \left[ e^{\kappa z} \cos(\kappa x) + e^{-\kappa z} \cos(\kappa y) \right] / 2, \quad (1)$$

where $B_0$ is the value of magnetic field at $z = 0$ on the axis $x = y = 0$, and $\kappa$ is a length scale parameter. The magnetic field lines are embedded in the surfaces $F = \text{const}$, i.e. $\mathbf{B} \cdot \nabla F = 0$, with

$$F(x, y, z) \equiv e^{2\kappa z} \sin^2(\kappa x) + 2[1 - \cos(\kappa x) \cos(\kappa y)] + e^{-2\kappa z} \sin^2(\kappa y). \quad (2)$$

see Fig. 1. In the limit $z \to \pm \infty$ the cross section condenses into an infinitely thin sheet.
The Alfvén continuum, i.e., its frequencies and modes, in the limit of small plasma \( \beta \), is governed by a single second-order ordinary differential equation along the magnetic field lines [2]. If \( z \) is used as a label along the field line chosen the mode equation becomes

\[
\frac{d}{dz} \left[ p(z) \frac{du}{dz} \right] + \Lambda s(z) u(z) = 0 , \tag{3}
\]

where \( p = \dot{B}_z/Q, \ s = 4/(\dot{B}_z Q), \ Q = \dot{B}^2 / (\nabla F)^2 \) and \( \Lambda = \mu_0 \rho \omega^2 / B_0^2 \). In these definitions, \( \mu_0 \) and \( \rho \) are, respectively, the vacuum permeability and the plasma mass density, while \( \dot{B}_z \) and \( \dot{B} \) are, respectively, the \( z \) component and the magnitude of \( \mathbf{B} \) normalized with respect to \( B_0/2 \). For the time dependence of the mode the ansatz \( u(z,t) = u(z) \exp(i \omega t) \) was made. \( \Lambda \) equals the frequency squared divided by the square of the Alfvén speed on the magnetic axis. The coefficients \( p(z) = p(x(z), y(z), z) \) and \( s(z) \) in Eq. (3) have to be taken along the magnetic field line \( x = x(z), \ y = y(z) \). The field line orbit is determined by the coupled system of differential equations

\[
\frac{dx}{dz} = B_x(x, y, z)/B_z(x, y, z) , \quad \frac{dy}{dz} = B_y(x, y, z)/B_z(x, y, z) . \tag{4}
\]

If \( \Lambda \sim \omega^2 \) is real, Eq. (3) is self-adjoint, provided \( R \equiv \lim_{a \to \infty} p u^* \frac{du}{dz} |^+_a \) is bounded. The latter condition in practice reduces to the requirement that \( u(z) \) be bounded asymptotically, as \( z \to \pm \infty \).

Within these premises, the nature of the Alfvén spectrum still depends on the details of the asymptotic behavior of \( u(z) \). Two general classes of solutions and therefore spectral components can be identified according to whether the integral \( I \equiv \int_{-\infty}^{+\infty} u^2(z) \, dz \) is bounded or not. If \( I \) is bounded, \( \Lambda \) is a discrete eigenvalue. The associated mode \( u(z) \) is a discrete or localized eigenfunction which falls off to zero asymptotically for \( |z| \to \infty \). If \( I \) is unbounded, \( \Lambda \) is in the continuous spectrum and the associated mode \( u(z) \) is an extended eigenfunction which stretches along the whole infinite length of the field line.

On magnetic surfaces close to the axis, \( |\kappa z|, |\kappa y| \ll 1 \), Eqs. (4) can be solved analytically. There results

\[
x(z) = x_0 \left[ 2 e^{-\kappa z} / (e^{\kappa x} + e^{-\kappa x}) \right]^{1/2} , \quad y(z) = y_0 \left[ 2 e^{\kappa y} / (e^{\kappa x} + e^{-\kappa y}) \right]^{1/2} ,
\]

where \( x_0 \) and \( y_0 \) are the transverse coordinates of the field line in the plane \( z = 0 \). Equation (3) becomes

\[
\frac{d}{dz} \left[ \left( e^{\kappa z} x_0^2 + e^{-\kappa z} y_0^2 \right) \frac{du}{dz} \right] + \frac{e^{\kappa z} x_0^2 + e^{-\kappa z} y_0^2}{\cosh^2(\kappa z)} \Lambda u(z) = 0 . \tag{5}
\]

For \( \alpha \equiv y_0^2 / x_0^2 = 1 \), Eq. (5) can be solved by substituting \( \kappa \tau(z) = \arctan[\sinh(\kappa z)] \). This yields the equation, \( d^2 u / d\tau^2 + \Lambda u = 0 \), which, for \( \Lambda \) real, has the solutions,

\[
u(z) = d_1 \sin \left[ \sqrt{\Lambda} \arctan(\sinh z) \right] + d_2 \cos \left[ \sqrt{\Lambda} \arctan(\sinh z) \right] , \quad \text{for } \Lambda > 0 \ , \tag{6}
\]

\[
u(z) = d_3 \exp \left[ \sqrt{-\Lambda} \arctan(\sinh z) \right] + d_4 \exp \left[ -\sqrt{-\Lambda} \arctan(\sinh z) \right] , \quad \text{for } \Lambda < 0 , \tag{7}
\]

where \( (d_1, d_2, d_3, d_4) \) are integrating constants, and \( \kappa \) has been set equal to 1.
On the stable side of the spectrum, where $\Lambda > 0$, localized eigenmodes do exist. The boundary conditions $u(\pm \infty) = 0$ for Eq. (6) are satisfied with $\Lambda = \Lambda_m = m^2$, $m = 1, 2, 3, \ldots$ with the corresponding eigenfunctions $u_1(z) = 1/\cosh z$, $u_2(z) = \sinh z/\cosh^2 z$, $u_3(z) = (1 - 3\sinh^2 z)/\cosh^3 z$, etc. In Fig. 2, the eigenfunctions $u_1(z)$ and $u_{10}(z)$ are shown. Note that the analogue of Eq. (3) does not have localized solutions in axisymmetric configurations as follows from the theory of differential equations with periodic coefficients. In nonaxisymmetric toroidal MHD equilibria, however, provided such equilibria exist, localization seems to be possible again [3].

In addition to discrete modes, the $\alpha = 1$ solutions also yield continuum modes. For $\Lambda \neq \Lambda_m$ all solutions (6) and (7) asymptotically tend towards a nonvanishing constant, at least at one end of the configuration. This implies the divergence of the integral $I$. Consequently, the continuous spectrum consists of the entire complex $\omega$ plane. If Im($\omega$) < 0, the continuum modes are unstable. Symmetric and antisymmetric stable continuum modes are shown in Fig. 3, while examples of unstable continuum modes are shown in Fig. 4. The occurrence of an unstable Alfvén spectrum characterized by Im($\omega$) < 0 in a static equilibrium contradicts conventional results. However, modes in the unstable continuum generally will not contribute to the wave dynamics of the plasma equilibrium because they are tied to specific inhomogeneous boundary conditions that must be specified at the singular ends of the equilibrium, $z \to \pm \infty$. If these conditions are not satisfied, the unstable continuum modes do not appear.

For field lines with $0 < \alpha < 1$ the spectrum and the modes have to be determined numerically. The dependence of the eigenvalues $\Lambda_1 - \Lambda_3$ on $\alpha$ is shown in Fig. 5. For a fixed value of $\Lambda$ a sequence of continuum modes is shown in Fig. 6, displaying a smooth transition as $\alpha$ varies from zero to one. For $1 < \alpha < \infty$ the spectrum and the modes can be obtained by replacing $\alpha$ with $1/\alpha$ and $z$ with $-z$.

Concerning the global behaviour of modes on the whole magnetic surface the following picture emerges from Figs. 5 and 6. There are frequency domains ($\omega^2 \sim \Lambda$) in which a single Alfvén continuum mode, i.e. with fixed frequency, is able to cover smoothly the whole surface, $0 \leq \alpha \leq \infty$, just as in, for example, axisymmetric equilibria. A phase on each field line and a poloidal mode number at one arbitrary position $z = \text{const}$ are free. Interspersed between these continuum domains there are frequency domains in which on at least one field line the mode extension shrinks to a finite value. It is conceivable that the interaction of the plasma with e.g. externally applied waves or fast particles is concentrated in or excluded from this region.

Fig. 1: Magnetic surface $F = 1.0$ for $-2.3 < z < 2.3$. With four field lines.

Fig. 2: Eigenfunctions $u_1(z)$ (dashed) and $u_{10}(z)$ (solid).

Fig. 3: Stable continuum modes $u(z)$ for eigenvalue parameter $\Lambda = 1/\sqrt{3}$ (dashed) and $\Lambda = 100/\sqrt{3}$ (solid).

Fig. 4: Unstable continuum modes $u(z)$ for eigenvalue parameter $\Lambda = -1/\sqrt{3}$ (dashed) and $\Lambda = -100/\sqrt{3}$ (solid).

Fig. 5: Dependence of $\Lambda_1$ (lower curve), $\Lambda_2$ (middle curve) and $\Lambda_3$ (upper curve) on field line position $\alpha$.

Fig. 6: Dependence of continuum modes $u(z)$ on field line position $\alpha$: $\alpha = 1$ (upper thick curve), $\alpha = 10^{-1}$, $10^{-2}$, $10^{-3}$ (curves in descending order) and $\alpha = 0$ (lower thick curve). With $\Lambda = 1/\sqrt{3}$ fixed.
Time–resolved Transport in W7–X as predicted by Neoclassical Theory

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In next–generation stellarators such as W7–X, recycling and gas puffing affect only the edge of the plasma. Hence, an active density profile control via central particle sources, e.g., NBI and/or pellets, is needed. To reach a certain experimental plasma density, the central particle source is not only determined by diffusive losses due to the density gradient as for instance in tokamaks. The off–diagonal elements of the neoclassical transport matrix become essential in the stellarator long mean free path (LMFP) regime which is mandatory for the central region of W7–X plasmas. Especially the off–diagonal term which is related to the particle balance predicts additional particle losses due to the temperature gradient.[1] Hence, strong temperature gradients together with too small a central gas refuelling rate will drive the density hollow, i.e. a positive density gradient has to compensate the off–diagonal losses to fulfil the particle balance. On the other hand, too large a filling rate may give rise to increasing density and decreasing temperature.

Neoclassical Transport Matrix

The neoclassical particle and energy flux densities, $\Gamma_\alpha$ and $Q_\alpha$, with $\alpha = e, i$, are given by

$$\Gamma_\alpha = -n_\alpha \{ D_{11}^\alpha \left( \frac{n_\alpha'}{n_\alpha} - \frac{q_\alpha E_r}{T_\alpha} \right) + D_{12}^\alpha \frac{T_\alpha'}{T_\alpha} \}$$

$$Q_\alpha = -n_\alpha T_\alpha \{ D_{21}^\alpha \left( \frac{n_\alpha'}{n_\alpha} - \frac{q_\alpha E_r}{T_\alpha} \right) + D_{22}^\alpha \frac{T_\alpha'}{T_\alpha} \}$$

with $q_\alpha$ being the particle charge and with the "convective term" $(3/2)T_\alpha \Gamma_\alpha$ included in $Q_\alpha$. The radial electric field $E_r$ is determined by the roots of $Z_i \Gamma_i = \Gamma_e$. The neoclassical transport coefficients $D_{jk}^\alpha$ (with $j,k = 1,2$) are determined from the mono–energetic transport coefficients, e.g., $D_{1/1/0} \sim 1/\nu$, $D_{\sqrt{\nu}} \sim \sqrt{\nu}$ and $D_\nu \sim \nu$ with $\nu$ being the collision frequency [2,3].

In order to show the time evolution of the density and the temperature profile, the ASTRA code [4] is used in a stellarator specific version [5] with the full neoclassical transport matrix for the W7–X high–mirror configuration [6] (see Fig. 1). The ambipolar $E_r$ is calculated by direct iteration yielding the so–called “ion root” which is expected for the plasmas of interest here. For simplicity, $n_e = n_i = n$ is used in all cases. To yield quantitative estimates about the central refuelling rate which is needed to reach a scenario with non–hollow density and pressure profiles calculations are performed at various heating powers of up to 20 MW (electron heating only) and for different initial conditions of densities and temperatures.

Need for Central Particle Refuelling

As can be seen from Fig. 2, too large a central heating and too small a central refuelling rate produces a hollow density profile as expected from theory [1]. In this case, the start conditions are $n = 1 \times 10^{20} \text{ m}^{-3}$, $T = 1 \text{ keV}$, $P = 10 \text{ MW}$. The filling rate is $1 \times 10^{20}$
s\(^{-1}\). Here, a strong temperature gradient driven by the heating cannot be balanced or exceeded by the external particle source. While the density gradient becomes positive, the negative electron temperature gradient is increased by the increased value of heating power per particle, driving the density gradient more and more positive. Hence, no stationary situation is reached here and the whole central density may be lost within a few hundred milliseconds. In this case also an inverted pressure profile occurs with \(p' > 0\), which may drive MHD instabilities.

Enlarging the central refuellng rate to over-compensate the particle losses which are caused by the strong off-diagonal term of the transport matrix coupled to the temperature gradient yields an almost contrary situation. The hollow density profile is observed only at the beginning of the discharge owing to the start conditions. At later times, when the influence of the start conditions is more and more lost, the high refuellng rate is able to drive a steep density profile. Close to the plasma edge, however, the high temperature together with an unfavourable scaling of the transport coefficients result in local particle fluxes which generally are less than the central refuellng rate. A transport barrier occurs. Due to this barrier, the central density increases and becomes more and more peaked. Both temperatures decrease slowly and become identical, as can be seen from Fig. 3 (refuellng: \(15 \times 10^{20} \text{ s}^{-1}\); \(n: 2 \times 10^{20} \text{ m}^{-3}\); \(T: 1 \text{ keV}\); \(P: 10 \text{ MW}\)). As far as an active control of the central particle source is possible, e.g. via pulsed pellet scenario, this behaviour is less critical than the case shown in Fig. 2. This situation does not differ much by changing the density and temperature start conditions, as can be seen from Fig. 4. This figure shows also that independent from the start situations three different regimes can be distinguished: The low rate regime, where the whole central density may be lost within less than a few hundred milliseconds; the high rate regime with highly peaked and increasing densities and lifetimes of a few seconds before the pressure profile becomes hollow due to almost vanishing central temperatures; and the intermediate regime with almost constant plasmas. The rate needed for this intermediate regime is almost proportional to the heating power, i.e. the filling rate must be close to \(1 \times 10^{20} \text{ m}^{-3}\) per second per MW heating power according to the value determined in [1] with stationary calculations.

Further, simulations show that not only the total particle source is linked to the total heating power. To reach non-hollow pressure profiles, the radial profile of the particle deposition must be strongly correlated to the power deposition profile. In all cases with almost stable conditions, both profiles are nearly identical. With a too narrow or too broad particle source, either the density or the temperatures are driven hollow, depending on the total refuellng rate. The reason for this behaviour can be found from the strong coupling of the transport equations in the central region. Assuming sufficiently high temperatures, the radial electric field can be neglected. Solving the stationary equations for the filling rate yields an ambipolar particle flux which is roughly proportional to the energy fluxes of the electrons and the ions. Here the radial dependences of the transport coefficients cancel each other. Hence, the particle refuellng profile is similar to the power deposition profile.

In W7-X, this strong coupling is necessary for the high-mirror configuration. Lowering the toroidal mirror term will change the local behaviour of the transport coefficients with the effect of broader refuellng profiles which are less coupled to the heating profile. Central refuellng, however, will be mandatory.

**Transport Barrier at outer Radii**

At outer radii, the neoclassical particle flux must exceed at least the central particle source. Otherwise the outer density increases and the global density control is lost. In
this region, however, the neoclassical transport coefficients of the electrons are much smaller than those of the ions. Solving the transport equations for the particle flux at outer radii shows that the ambipolar partial flux can simply be described as being proportional to $D_{ii}^e$ and the normalized temperature gradient. For the same reason, the ion energy flux is close to the whole energy flux, which should be equal to the central heating power in stationary plasmas. As far as both the central heating and the central particle source have to be strongly coupled, a strong constraint between the particle flux at the edge and the central heating power can be found which mainly depends on the ratio $D_{ii}^e/D_{ii}^i$. In terms of temperatures, a critical temperature for the edge region can roughly be estimated (mainly proportional to $D_{ii}^e/D_{ii}^i$ and to the central temperature) which defines a transport barrier. For $T$ exceeding this critical value, the outer particle flux is less than the central refuelling rate, and the outer density increases. Additional moderate anomalous heat conductivity (up to $1 \text{ m}^2/\text{s}$) or radiative losses (up to $3 \text{ MW}$) will amplify the problem.

In Fig. 5, the time evolution of density and particle flux is shown for such a plasma (refuelling: $15 \times 10^{20} \text{ s}^{-1}$; $n$: $10^{20} \text{ m}^{-3}$; $T$: $1 \text{ keV}$; $P$: $10 \text{ MW}$). At the beginning, the whole refuelling rate is balanced by particle losses. The transport barrier which occurs at about 100 milliseconds reduces the particle losses to nearly two third of the central filling rate.

The ratio of both transport coefficients depends, in principle, on the magnetic configuration. Hence, in W7–X this problem can be analyzed by varying the toroidal mirror term over a wide range.

References


**Fig. 1:** Magnetic field topology of the W7–X high mirror configuration [6] at two radii with $\zeta$ and $\Theta$ being the toroidal and poloidal angle.
Fig. 2: Density and temperatures of a discharge with low particle refuelling rate ($P = 10$ MW).

Fig. 3: Density and temperatures of a discharge with high particle refuelling rate ($P = 10$ MW).

Fig. 4: Operational regimes for different heating powers and refuelling rates.

Fig. 5: Time evolution of density (left), particle source (middle) and particle flux (right) for a discharge with occurring transport barrier.
High-Confinement NBI Discharges in W7-AS


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Abstract: In W7-AS, the longest energy confinement times were achieved in NBI-heated discharges under low wall-recycling conditions. Low recycling is needed to control the density at line averaged values of $\bar{n}_e \approx 10^{20}$ m$^{-3}$. Under these conditions, electron and ion temperatures of up to 1 keV and confinement times of 55 ms were obtained with an absorbed heating power of $P \approx 0.35$ MW. In NBI discharges without density control, the density rose up to typically $2 \times 10^{20}$ m$^{-3}$ and the temperatures remained at 0.3 keV only.

From the H-mode, where after a fast transition the profiles broaden, these discharges can be distinguished by low edge densities and a rather gradual improvement of energy confinement. The reduction of transport is concentrated to a layer at about 2/3 of the plasma radius. In this region steep temperature gradients and a strong gradient in the radial electric field develop. In this respect the discharges are similar to discharges with stronger NBI and additional electron cyclotron resonance heating (ECH), where the highest ion temperatures were achieved [1].

What is specific for the discharges here is the gradual transition. Since the measured electric field is consistent with the neoclassical ambipolar field, this high-confinement mode could be an example where the neoclassical electric field leads to a suppression of anomalous transport, this in turn produces steeper gradients which lead according to neoclassical theory again to a stronger field. This loop is interrupted when transport reaches the neoclassical level. The result is a transport barrier similar as in high performance tokamak discharges like in JET [2].

Global Energy Confinement Time: In Fig. 1, we compare the energy confinement time $\tau_E$ of high-confinement discharges with the W7-AS dataset contributed to the ISS95 scaling [3], which covers our previously accessible parameter range. In the ISS95 dataset a trend was already observed of the kinetic confinement times, as deduced from profiles, being about 15% higher than the diamagnetic ones. For the discharges here, the kinetic confinement time is about 35% higher than the diamagnetic one. A 15% discrepancy can be understood in terms of the diamagnetic loop picking up a component of the vertical field generated by Pfirsch-Schlüter currents. The 35% difference remains unclear at the moment. Since the profiles are well documented for

Fig. 1: Diamagnetic confinement times of high-confinement NBI, high ion temperature and H-mode discharges compared with the ISS95 scaling (W7-AS pre-factor) and W7-AS data contributed to the ISS95 database (open symbols). Crosses: kinetic values.
these discharges, we have added the kinetic confinement times to the plot. The figure underlines that the confinement time of these discharges is a factor of 2 above our standard confinement level.

**Characteristics:** In Fig. 2, the main characteristics of a high-confinement NBI discharge are depicted. The distinct feature is the continuous increase from 0.25 to 0.4 s of the energy content at constant heating power and line density. During this phase, which corresponds to three energy confinement times, electron and ion temperatures increase. Simultaneously, a reduction in the H₉ light indicates improving particle confinement. The slow increase of the impurity content in the core is consistent with impurity transport studies [4] yielding a very low particle diffusivity of 0.07 m²/s and inward convection of 5r/a (m/s).

The improvement is accompanied by a reduction in the density fluctuation amplitude, which is partly due to the fact that the reflecting layer moves 1.5 cm inward during the measurement. We observe the reduction primarily in the frequency range below 60 kHz. The radial electric field at a location close to the minimum of the thermal diffusivity becomes continuously stronger while confinement is improving.

In Fig. 3, the temporal evolution of electron density and temperature profiles of this discharge can be studied. The increase of the electron temperature occurs over the entire cross-section and stops at around 0.4 s first in the outer region and later in the core. The increase of the energy content can be partly attributed to the broadening of the profile. The evolution of the density profile continues up to 0.55 s. After an initial evolution from a broad to a more peaked profile, it is primarily the edge density which decreases. We see that a continuing reduction of the edge density does not lead to a further improvement of energy confinement. The possibility of controlling the density seems to be the crucial element rather than the low edge density. Radial derivatives of the data in Fig. 3 show a pronounced steepening of both temperature and density gradients in the radial zone between 10 and 14 cm during the initial phase where confinement improves. The gradients remain unchanged after 0.4 s. The bolometric
Fig. 3: Temporal evolution of the electron temperature profile from fits to 24 ECE radiometer channels (left) and the density profile unfolded from 8 interferometer channels (right).

radiation profiles are hollow with a maximum at about 0.14 m and start slowly to fill up in the center. At the end of the discharge, this process is not yet completed.

**Power Balance Analyses:** The transport properties of a high-confinement NBI discharge are compared with those of an ECH discharge at similar density and a heating power of 0.5 MW. The ECH discharge serves as "normal" confinement case. Although the absorbed power is higher, the kinetic energy content is only 13 kJ compared to 17 kJ of the NBI discharge. The power balance analyses shown in Fig. 4 were carried out with the same heat diffusivity $\chi \equiv \chi_e = \chi_i$ for electrons and ions. Electron temperature and density profiles from Thomson scattering serve as input and $\chi$ as well as the ion temperature profile are derived from the power balance. They are compared with measurements and neoclassical predictions from the DKES code [5]. For both discharges, the calculated ion temperature agrees with the measurements.

The neoclassical diffusivities in Fig. 4 are total energy fluxes divided by the temperature gradients, hence the same quantity which is extracted from the power balance. In case of the ECH discharge, $\chi$ agrees well with the neoclassical estimate for the ions and is well above the electron value. Hence, if the ions behave
neoclassically, the electron heat diffusivity must be strongly anomalous.

In case of the NBI discharge, there is an overall reduction of the diffusivity. \( \chi \) is below the neoclassical ion diffusivity and touches at its minimum the neoclassical electron diffusivity. From this result for \( \chi \) for \( r < 10 \) cm, we cannot conclude that \( \chi_i \) is below the neoclassical value. A two fluid power balance showed that using the neoclassical prediction for \( \chi_i \) is also consistent with the central ion temperature measurements. In such an analysis in the core, \( \chi_e \) drops only by 20%. At the transport barrier (\( r \approx 12 \) cm), however, the two fluid analysis fails to reproduce the data. Either the ion temperature is predicted as too high or, if in this region additional anomalous losses are added to the ion channel, \( \chi_e \) drops well below its neoclassical value. The result at the barrier is that either \( \chi_e \) or \( \chi_i \) or both are below the neoclassical prediction. Hence this region satisfies the criteria of a transport barrier in which anomalous transport seems to be suppressed.

\textbf{E \times B and Magnetic Shear:} The slow temporal evolution of confinement could also point to current penetration and magnetic shear as an important ingredient. The \( t \) profile of the high-confinement NBI discharge has a low magnetic shear zone in the outer half of the plasma. By adding positive and negative currents to the discharge it was shown, that the magnetic shear was not the key parameter. It can be ruled out too that the existence of negative shear is, as in tokamaks, a pre-requisite of achieving the high confinement times.

In comparison with the transport barriers observed in the H-mode and in high performance tokamak discharges the role of the radial electric field is of special interest. In W7-AS in general, there is agreement between the electric field as deduced from spectroscopy and from the ambipolarity of the neoclassical fluxes [6]. Fig. 5 shows that this also holds for the high-confinement NBI discharges. The radial zone of low transport is characterized by strong gradients in the radial electric field and hence strong shear in the perpendicular plasma flow.

In Fig. 2, it was shown that the electric field increases together with electron and ion temperatures. The fact that confinement improves gradually rather than through a fast transition could point to a causality loop where transport is reduced by sheared flow generated by the neoclassical electric field. The reduction of transport steepens the gradients, which increase again the neoclassical electric field etc. This loop would stop when anomalous transport is suppressed and transport is on the neoclassical level. The time constants of the process would be in the order of the confinement time.

\textbf{References}


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Collective Thomson scattering at W7-AS

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Abstract. Collective Thomson scattering (CTS) of electromagnetic radiation from thermal plasma fluctuations in principle allows the velocity distribution of plasma ions and its composition in the plasma to be measured. The use of powerful microwave radiation from gyrotrons opens new perspectives for the application of CTS, which is considered to be a promising candidate for alpha-particle diagnostics in reactor-size tokamaks with D/T operation. We have performed the first experiments at W7-AS with different scattering geometries to prove the applicability of gyrotrons for CTS. The experiments were performed with a 140 GHz gyrotron which is routinely used for ECRH, delivering a power of 0.45 MW. The receiver antenna and detection system for the registration of CTS spectra were especially designed for the scattering experiment. In backscattering experiments, which have inherently no spatial resolution, we have measured a transversely propagating, non-thermal lower-hybrid turbulence, which is driven by perpendicularly injected fast particles from a diagnostic neutral beam. The instability is excited by the beam ions under double-resonance conditions, where the LH frequency coincides with some harmonic of the beam ion gyrofrequency. For scattering geometries with the scattering wavevector not perpendicular to the magnetic field, thermal density fluctuations in the plasma were experimentally detected. The ion temperatures derived from these thermal spectra agree well with other diagnostics.

A modified scattering geometry (90° scattering) allows local measurements of the ion temperature and is considered a prototype for the design of a routine diagnostic for ion-temperature measurements.

1. Introduction

Powerful gyrotrons are well suited for application in collective Thomson scattering (CTS) experiments designed to investigate the ion velocity distribution function. Their output in the millimetre wavelength range allows the use of a 90° scattering geometry with good spatial resolution as compared to, for example, small-angle forward scattering with CO₂ lasers. D₂O lasers were used for ion-temperature measurements in tokamaks in the FIR wavelength range [1, 2], but these sources have a short pulse duration of a few microseconds which strongly limits the sensitivity of the CTS technique. High power gyrotrons developed for electron cyclotron resonance heating (ECRH) at present reach pulse lengths of the order of a few seconds and are now approaching CW operation at a power level of a few 100 kW.
This allows the sensitivity of CTS systems to be considerably increased by the so-called radiometric gain $G = \sqrt{\Delta f \tau}$, where $\Delta f$ is the bandwidth of the spectral channel and $\tau$ is the averaging time. With a typical bandwidth of the order of 100 MHz a gain of two to three orders of magnitude or more can be reached. This is the reason why CTS experiments with gyrotrons as the radiation source were planned and realized in a number of fusion installations [3–5].

We discuss the problems related to the use of gyrotrons as the probing source for CTS and the interpretation of the scattering spectra measured at the W7-AS stellarator.

2. Scattering geometry at W7-AS

W7-AS is a modular stellarator experiment with major radius of about 2 m and an effective minor radius of about 0.2 m. Since ECRH usually provides plasma start-up and operates at 70 and 140 GHz, this implies that the confining magnetic field is around 1.25 T (second cyclotron harmonic for 70 GHz) or 2.5 T (first harmonic for 70 GHz and second harmonic for 140 GHz). There is also the possibility of operation in regimes with a non-resonant field; in this case the start-up is achieved with a 900 MHz source after which the plasma is sustained with neutral-beam injection (NBI).

Microwave power is injected in toroidal sectors where the plasma column has an elliptical cross section with strong elongation in the vertical direction. The magnetic-field configuration has low shear and radial variation with a scale length of about 1.3 m. Typical plasma parameters for the results to be presented are: $0.2 \times 10^{20} \text{ m}^{-3} < n_e < 1.2 \times 10^{20} \text{ m}^{-3}$, $0.5 \text{ keV} < T_e < 2.0 \text{ keV}$, $0.2 \text{ keV} < T_i < 0.7 \text{ keV}$.

Two antenna blocks for CTS at 140 GHz are installed in two ECRH sectors; both of them employ existing launching antenna systems for ECRH at 140 GHz. The antenna installed in sector I is an off-axis parabolic reflector antenna with a circular (monomode) waveguide feed to produce a nearly Gaussian beam in the plasma volume (see figure 1). The antenna, monomode section and overmoded transmission line to the receiver introduce an attenuation of 25–30 dB. This antenna is rigidly attached to the movable mirror of the launching antenna thus defining a fixed geometry close to backscattering which can be scanned both in the vertical and horizontal directions. Two receiving antennas installed in sector II consist of a scalar horn with two reflecting mirrors (Cassegrainian configuration). Total losses from receiving antenna to receiver input are significantly reduced ($\approx 10$ dB). The vertical receiving antenna can be radially shifted by $\pm 3$ cm (see figure 2). Two 140 GHz gyrotrons are installed in this sector; both of them provide scanning of horizontally launched beams in the vertical and horizontal directions and can be used for CTS in a scattering geometry at approximately 90°.

In sector I, the incident beam is vertically polarized corresponding mainly to the X-mode; the scattered radiation can be analysed in two orthogonal linear polarizations (X- or O-mode). Within the scanning range used the mode purity is $> 75\%$.

In sector II, the polarization of radiation launched into the plasma is controlled by a universal polarizer in each microwave beam. With two antennas, linear polarization corresponding to O- or X-mode propagating transverse to the magnetic field can be received; so, in principle, in sector II all possible schemes of collective scattering can be investigated ($X \Rightarrow X$, $X \Rightarrow O$, $O \Rightarrow X$ and $O \Rightarrow O$ mode scattering). In all cases both probing and receiving beams have circular cross sections with beam diameters of about 4 cm in the plasma. This means that in sector I (backscattering) there is practically no spatial resolution along the major radius (cf figure 1), while in sector II (90° scattering) it is possible to perform local measurements with a spatial resolution of about 4 cm (cf figure 2).
3. CTS relevant characteristics of gyrotrons

At present gyrotrons used in collective scattering experiments are designed as powerful microwave sources for plasma heating. However, there are some essential differences in the gyrotron parameters required for plasma heating in contrast to collective scattering. For ECRH it is necessary to provide high power in a long pulse at the frequency corresponding to the EC-resonance condition. For collective scattering it is usually sufficient to work with shorter pulses with 10–100 ms duration. In order to reduce the level of ECE noise polluting the CTS spectrum it is highly desirable that the frequency of probing radiation lies between electron cyclotron harmonics. Additional requirements are: spectral purity, sufficiently small frequency drift during the scattering time interval and high shot-to-shot reproducibility of gyrotron operation. With respect to frequency drift the gyrotrons installed at W7-AS are not the optimum candidates for CTS experiments.

A serious problem for CTS is the stray radiation which becomes more pronounced with decreasing size of the plasma chamber. The reason for this is the low level of scattered power which lies typically 15 orders of magnitude below the gyrotron power. The stray
radiation has several consequences: (i) It is necessary to protect the receiver system. To provide this the insertion of notch-filters at the gyrotron frequency is mandatory. (ii) The intrinsic noise generated by the gyrotron outside the central frequency appears as a parasitic signal which cannot be eliminated by the notch filter. The only way to minimize this noise is to provide maximum decoupling between emitting and receiving antennas. (iii) If several gyrotrons with slightly different frequencies (\(\Delta f/f \approx 10^{-3}\)) are operating at the same time as can be the case during ECRH it was found that it is impossible to perform CTS measurements. Usually the frequencies of the other heating gyrotrons lie within the receiver bandwidth but outside the suppression band of the notch-filter tuned to the gyrotron line. Even stray radiation at the second harmonic of the 70 GHz ECRH system will overload the detection system operating around 140 GHz.

The optimum solution to this problem is a specially designed tube operating at a frequency different from that of the heating gyrotrons. A compromise could be the construction of a step-tuneable tube which in the heating mode provides powerful long pulse radiation and in a second operating regime provides output power at a different (by 10–15%) frequency with the power and pulse duration ensuring a sufficient signal-to-noise ratio for CTS.

All three 140 GHz gyrotrons installed for ECRH at W7-AS with a nominal power of 500 kW each may be used in CTS experiments. For all of them a study of frequency drift and of shot-to-shot stability of operation regimes was performed. The frequency variation during rf pulse for these tubes is presented in figure 3. Gyrotron A (used in the backscattering geometry of sector I) with relatively short pulse duration (\(\leq 1\) s) shows a practically linear frequency decrease during a 1 s pulse with a rather low constant rate of \(\approx 100\) MHz s\(^{-1}\) (see figure 3(a)) and very high shot-to-shot stability (frequency deviation within a few MHz). The long-pulse tubes C and E designed for pulse durations up to 3 s need a rather long time to establish quasi-steady-state regime of microwave power generation; during this transient
stage a strong frequency drift of the central line occurs. The transient regimes are strongly different for the two tubes investigated: in a pulse of gyrotron C during the transient time of \( \approx 100 \) ms the decrease of the central line frequency is \( \approx 600 \) MHz; the corresponding values for gyrotron E are \( \approx 70 \) ms and \( \approx 70 \) MHz (figure 3(b)). Shot-to-shot reproducibility in the transient regime for gyrotron C is not satisfactory; so it is practically impossible to use it for short pulse CTS measurements. With the gyrotron E this is possible due to a considerably better shot-to-shot reproducibility in the transient regime. In the quasi-steady-state stage of the pulse the frequency drift is small and the repetition stability for both tubes is adequate for CTS requirements. This is illustrated in figure 3(c) where the frequency evolution in the rf pulse of gyrotron C is shown after the first 200 ms of gyrotron operation.

Figure 3. Frequency drift during the gyrotron pulse for three tubes used in CTS experiments: (a) gyrotron A (in sector I), (b) gyrotron E (in sector II), (c) gyrotron C (in sector II) after the first 200 ms of operation.
4. Detection system

The detection system uses a superheterodyne receiver with a balanced mixer and a BWO as a local oscillator tuneable within a frequency range 110–170 GHz. The noise temperature of the receiver is $\approx 0.2$ eV. The LO phase locking system provides a relative frequency stability better than $10^{-8}$. At the receiver input a notch filter for the suppression of stray radiation is inserted. The tuneable notch-filter is based on a monomode rectangular waveguide with coupled cylindrical cavities [6] which provides attenuation of more than 40 dB in the $\pm 20$ MHz band around the line centre. Attenuation outside the $\pm 100$ MHz band from the line centre is $\leq 3$ dB (see figure 4). The voltage controlled attenuator at the receiver input increases the dynamic range by 50 dB. It also provides protection of the mixer (by means of programmed gating with 50 dB attenuation) from the stray radiation which may be present outside the notch-filter suppression band especially during the chirping phase of gyrotron operation.

![Graph](image)

**Figure 4.** Typical frequency characteristics of the mechanically tuneable notch-filter [6].

In general the detection system operates with the LO frequency coinciding with the gyrotron frequency which results in the superposition of the upper and lower sidebands with respect to the gyrotron line. This scheme assumes that the scattered radiation spectrum is symmetric. The scattered spectrum is analysed in the IF frequency range using two different filter banks. The main frequency analyser covers the range 50 to 1200 MHz and consists of 32 closely packed channels with a relative bandwidth of 10% and a relative frequency difference between adjacent channels of 10%. It is designed to cover the thermal ion feature over a wide range of plasma parameters (D and H plasmas, ion temperatures from 0.1 to 2 keV, scattering angles from 60 to 180°). The second filter bank can ‘zoom’ a frequency window of the scattering spectrum covering a total bandwidth of about 100 MHz. It consists of 20 identical channels with a bandwidth and channel separation of about 5 MHz. The ‘zoom’ system is tuneable over the full range of the main frequency analysing system and allows the possible fine structures in the spectrum of the scattered radiation to be investigated.

The standard integration time of the detection system is 1 ms which provides a radiometric gain ranging from $\approx 50$ (in narrow channels) to $\approx 200$ (in wide channels). The analogue outputs from all 52 channels are stored during a plasma shot for up to 1 s duration.
Numerical routines allow the processing of stored scattering data. In particular, it is possible to increase the radiometric gain for CTS further by increasing the integration time.

5. Calibration of the detection system

The calibration of the detection system was performed using a gas discharge noise source. In addition, for the scattering geometry in sector I, absolute calibration of the whole detection system (together with the receiving antenna and transmission line) is possible due to the fact that the receiving antenna looks practically in the direction of the major radius and thus black-body plasma radiation at the second ECE harmonic in the X-mode (absolute temperature known from other diagnostics) can be used as a calibration source. In order to avoid all disturbances from the gyrotron the calibration was performed with the plasma only heated by NBI. From both calibration procedures we obtain a signal attenuation in the transmission system of 25–30 dB.

The test of the detection system was performed by measuring the ECE spectrum at the fourth harmonic (figure 5). It can be seen that a flat ECE spectrum is registered at the low level of \( \approx 2 \) eV which demonstrates the high sensitivity of the registration system in spite of strong attenuation in the transmission. The absolute value of ECE radiation temperature at the fourth harmonic is confirmed by independent measurements with the ECE diagnostic system operating routinely at W7-AS.

![Figure 5. ECE spectrum in O-mode at the fourth cyclotron harmonic measured with the CTS detection system.](image)

For the scattering geometry in sector II, an absolute calibration is not possible because of the vertical viewing direction. However, a relative calibration can be performed there using either the gas discharge noise source or ECE at the second harmonic which is 'grey-body' for O-mode and black-body for X-mode; in both cases the absolute value is not well determined because radiation originates from the plasma with radially non-uniform temperature. Nevertheless, the latter results can be related to the data from the gas discharge noise source and yield approximate values for the absolute calibration by estimating total losses from the receiving antenna to the receiver input. There is also the possibility to use the horizontal access through the ECRH windows in sector II to perform calibration by ECE at the second harmonic.
6. CTS experimental results and their interpretation

The main aim of CTS experiments at W7-AS is the detection of thermal ion features in the scattered signal. The signal power in terms of the equivalent noise temperature may be estimated as (see also [7])

$$T_s(\text{eV}) = 350G_pG_bP_0(\text{MW}) \left( \frac{L}{20 \text{ cm}} \right) \left( \frac{4 \text{ cm}}{d} \right)^2 \left( \frac{n_t}{10^{14}} \right) \frac{m_i}{m_p} \frac{1}{\sqrt{T_i(\text{keV})}}$$

where $G_p$ and $G_b$ are the polarization factor and the emitting–receiving antenna beam crossing factor, respectively, $P_0$ is the probing radiation power in MW, $L$ is the length of the scattering volume in the direction of the probing beam propagation, $d$ is the diameter of the emitted and receiving beams. For the CTS experiment on W7-AS, the signal level is estimated to lie in the range from ten to a few tens of eV (the uncertainty stems from the $G_p$ and $G_b$ factors). This formula does not take into account ‘fine’ effects like the difference between the electron and ion temperatures or the deviation of the plasma refractive indices from the vacuum value for both the probing and scattered radiation in the scattering volume.

6.1. Thermal CTS spectra from backscattering

For the backscattering geometry of sector I, stray radiation turned out to be a major problem. In order to avoid strong ECE background radiation the CTS measurements were performed with a magnetic field of 1.25 T which is the fourth harmonic for the probing frequency and is optically thin. This resulted, however, in an enhanced level of stray radiation which is practically not absorbed by the plasma, a circumstance which also limited the pulse duration of the probing radiation to avoid damage to the vacuum vessel. The pulse length was restricted therefore to typically 50 ms.

An optimal position of the launching–receiving antenna block was found which minimized the stray radiation level at the receiving antenna input to less than 100 W which corresponds to an antenna decoupling of up to 40 dB. Shots without plasma were also used for the measurement of the spectrum of the intrinsic gyrotron noise. At frequencies above 200 MHz, the noise was low enough to allow measurements of thermal spectra.

Measurements of CTS thermal spectra were performed in alternating sequences of gyrotron shots without plasma and with NBI supported plasmas at 1.25 T. The spectra due to scattering of the probing gyrotron radiation by the plasma were obtained as the difference between signals measured with and without plasma. An example of CTS spectra obtained with this subtraction procedure [8] is presented in figure 6 for two different plasma densities. The signal level increases with density as expected. The decrease in the total bandwidth of the measured spectrum with the density increase together with the more than proportional increase in the signal power may be qualitatively explained by changes in the plasma refractive index.

For the CTS spectra taken at low density a good fit to theoretical curves was obtained by taking into account the radial dependence of plasma density as well as electron and ion temperature [4]. Good agreement of the CTS ion temperature with independent diagnostics is found.

6.2. Thermal CTS spectra from 90° scattering

For the 90° scattering geometry (sector II), the decoupling between emitting and receiving antennas is at least two orders of magnitude higher than in the backscattering geometry.
This diminishes the level of stray radiation to a very low level and allows to register CTS spectra without any interference from the gyrotron noise.

To obtain CTS spectra from a finite-sized plasma volume the probing radiation should experience considerable absorption after the first passage through the scattering volume. This was achieved by launching the probing radiation in the X-mode and by placing the absorbing resonant layer close to the inner chamber wall (a toroidal magnetic field of about 2.35 T instead of 2.5 T which corresponds to absorption on-axis). The example of a CTS spectrum thus obtained in a $X \rightarrow O$ scattering scheme is shown in figure 7. It should be noted that the scattered signal (with the equivalent noise temperature of about 15 eV) is easily separated from the ECE background of $\approx 50$ eV. The spectrum of this background is flat with little deviation over a wide frequency range. The localization of the measurements was checked by toroidal scanning of the probing beam from the position of its optimal crossing with the receiving beam; as expected, the signal drops according to the angular width of the antenna radiation pattern (which corresponds to 4 cm spatial resolution).

The experimental CTS spectra are fitted well by theoretical curves for thermal density fluctuation spectra with plasma parameters in agreement with data from independent diagnostics. In the fitting procedure density and electron temperature are taken from independent diagnostics; plasma ion composition is fixed and it is assumed that all ion components have equal temperatures. Two values are fitted by a least-square procedure, namely the ion temperature and the vertical scale of the spectra. An example of fits of the same thermal spectrum with different ion compositions is presented in figure 8. The values of the derived ion temperature are dependent on the H/D mixing ratio; naturally the value of the ion temperature providing the best fit is decreasing with the increase of hydrogen content. The fitting procedure is also sensitive to the contents of heavy impurities (this is discussed in detail in [9] in the context of ion-temperature measurements in TCA using D$_2$O laser radiation). The results show that good agreement for the value of $T_i$ can be obtained provided that the H/D mixing ratio is known. In addition, an estimate for $Z_{e_{\text{eff}}}$ in the plasma should be feasible.
6.3. Non-thermal CTS spectra from lower-hybrid turbulence

In the backscattering geometry (sector I) a very intense narrow-band feature was detected in ECR heated plasmas (X-mode, second harmonic) simultaneously with a hydrogen neutral-beam injected for the active charge-exchange diagnostics [4]. The signal was registered in both of CTS sectors (sector I and sector II) toroidally separated from the port of NB injection by $\approx 2$ m and $\approx 5$ m, respectively. Spectral measurements with the 'zoom' analyser show a relative bandwidth of about 2–3% of the mean frequency of the registered feature and a signal level up to a few MeV of equivalent noise temperature (see figure 9). Note that this is at least five orders of magnitude higher than the thermal CTS signal. A scan in the toroidal direction with the antenna block shows that this enhanced signal exists only in a very narrow angular range (comparable with the radiation pattern of emitting and receiving antennas $\approx 1$–$2^\circ$) around the direction normal to the magnetic field. This intense (certainly of non-thermal origin) scattered signal is clearly correlated with the diagnostic NB injection. From the dependence of the frequency on the plasma density [4] the narrow-band feature was identified as being related to LH waves, propagating almost transversely to the magnetic field.

7. Theoretical modelling of the LH instability

The unique feature of the detected LH turbulence—the very narrow bandwidth of the scattering signal presented above—suggests the proposal of a theoretical model for its origin. Keeping in mind that the scattered signal is due to line integration throughout most of the minor plasma cross section along the major radius, the only possibility to explain the well-defined LH frequency is to assume that the instability is generated at some intrinsic plasma frequency (e.g. some definite number of ion cyclotron harmonic or LH frequency) in a very narrow region along the major radius. Otherwise the bandwidth should be considerably larger due to the natural inhomogeneity of plasma parameters (the characteristic scales along the major radius are of the order of the integration length of the scattering volume and are comparable for both intrinsic frequencies).

The LH turbulence is triggered by a weak diagnostic neutral beam (beam energy $\leq 20$ keV, equivalent beam current $\approx 2$ A) launched transversely to the magnetic field. As a result of the charge-exchange processes an ion beam is produced in the plasma with a velocity distribution corresponding closely to the injected neutral beam, i.e. having a
Figure 8. Fitting of a CTS spectrum by theoretical curves for plasmas with different ion compositions and fixed electron density and temperature.

well-defined energy distribution function and moving transversely to the magnetic field. Beam ions escape from the plasma due to vertical drift in the inhomogeneous magnetic field almost without energy and momentum losses. These ions fill a considerable part of the plasma volume (3–5 m in the toroidal direction, from the CX diagnostic port) with a density $n_b \approx (2-4) \times 10^{-4} n_0$.

To explain the experimental findings an instability mechanism for LH waves was proposed [4] which takes place in the presence of a transverse ion beam under the double-resonance condition. This strong instability of hydrodynamic type with a growth-rate proportional to the square root of the beam density (see also [10]) can exist in those regions where the LH frequency coincides closely with some harmonic of the beam ion gyrofrequency. An example of the normalized growth-rate of unstable LH waves as a function of the wavenumber $k$ for a typical set of plasma parameters is shown in figure 10.
The instability is predicted to exist over a broad range of wavenumbers and is therefore easily detectable with the CTS at 140 GHz. The narrow angular range of detected turbulence is explained by the propagation effect of LH waves in the non-uniform magnetized plasma. It is proposed that the instability saturates due to a mechanism similar to stochastic ion heating by LH waves which is consistent with the measured turbulence level. More detailed information about the theoretical and experimental results on LH turbulence is presented in a separate paper [11].

8. Anomalous CTS spectra from 90° geometry

A considerable amount of CTS spectra in 90° scattering geometry was obtained in operating regimes without the cyclotron absorption of probing radiation. Under these conditions, the dependence of the signal amplitude on the toroidal launching angle of the probing radiation was much less pronounced. This means that multiple reflections from the chamber walls (of probing or scattered radiation or both) definitely contribute to the resulting CTS spectra. Therefore, it is concluded that these spectra do not only originate from the volume defined by the crossing of the probing and receiving antenna beams.

Some of these spectra may be interpreted as the result of scattering from thermal fluctuations like those obtained by the absorption of probing radiation; but there is another reproducible type of CTS spectra which was obtained in regimes with a toroidal magnetic field of 1.27 T (fourth-EC harmonic) and 1.4 T (non-resonant). It possesses a very distinguished feature which looks like a well pronounced ion acoustic peak (figure 11). In all cases this peak has a frequency of about 200 MHz and demonstrates no dependence on plasma density, besides the fact that it was not observed in low-density plasma ($N_e < 2 \times 10^{13}$ cm$^{-3}$). These spectra cannot be fitted at all with a realistic temperature of all the ion components. All these ‘abnormal’ CTS spectra still require more understanding and a theoretical explanation. It is planned to repeat these measurements with a beam dump for the receiving antenna in order to exclude the possible effects from multiple reflections.
Figure 10. Instability growth-rate normalized to LH frequency as a function of $k$-number for typical plasma parameters in the ECRH regime ($B_0 = 2.5$ T, $N_e = 5 \times 10^{13}$ cm$^{-3}$, $T_e = 2$ keV, $T_i = 0.4$ keV) and a transverse ion beam with normalized density $N_b/N_e = 10^{-4}$ composed of three monoenergetic fractions corresponding to those of the injected neutral beam.

Figure 11. An example of the spectrum obtained without absorption of the probing microwave beam in plasma: the stars represent the experimental points and the full curve the expected thermal CTS spectrum from independently measured plasma parameters.

9. Conclusions and outlook

Both thermal and non-thermal CTS spectra were obtained at W7-AS in a proof-of-principle experiment using gyrotrons as the sources of microwave probing radiation and heterodyne...
detection. The spectra of scattered radiation are characterized by high-frequency resolution and small error bars. Two different scattering geometries were examined: backscattering and 90° scattering.

The ion temperature was determined from the thermal spectra in both scattering geometries yielding results in good agreement with independent diagnostics. The measurements performed in the 90° scattering geometry provide a spatial resolution of about 4 cm which allows the measurement of radial ion-temperature profiles even in a machine with the small dimensions of W7-AS.

In the backscattering geometry a very narrow-band non-thermal LH-turbulence was detected and carefully investigated experimentally with an appropriate theoretical model. It originates from LH wave instability which is driven by a fast transverse ion beam generated by a diagnostic neutral beam.

It has become evident that moderately sized installations such as W7-AS require special precautions in order to reduce the level of undesired stray radiation. We also conclude that for CTS experiments a separate gyrotron is desirable with its frequency differing from EC resonances. To detect thermal-ion features in CTS spectra at W7-AS high power and long pulse duration are not needed. It is, however, necessary to achieve high mode purity, low-frequency drift and high shot-to-shot stability. A compromise could be a step-tunable gyrotron which can be used at different frequencies for plasma heating or for CTS experiments.

In reactor-scale installations gyrotrons in the frequency range 140–200 GHz with long pulse operation are attractive for α-particle diagnostics, and the present experiments can be regarded as a proof-of-principle for this application. Future CTS experiments must be optimized with respect to the probing radiation frequency, scattering geometry and types of normal waves in the probing and scattering radiation. Such a system should provide a low enough ECE background at the antenna input as well as low cyclotron absorption of both the probing and scattered beams. Corresponding theoretical investigations have been published which analyse the requirements for an α-particle scattering experiment for JET [12] and ITER [13].

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Correlation between MHD-activity, energetic particle behaviour and anomalous transport phenomena in WENDELSTEIN 7-AS

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Introduction - Energy and particle transport in W7-AS exhibits a resonance-like dependence on the edge rotational transform (iota) as long as the magnetic shear is relatively weak (low $\beta$, no significant net toroidal currents). MHD modes at resonant surfaces may cause enhanced radial transport depending on the magnitude and radial extent of the magnetic perturbations. In many cases discharges in W7-AS are very quiescent, or in case of mode activity, often no influence on energy and particle confinement is found. In the high beta regime ($\langle \dot{\beta} \rangle \leq 1.8 \%$) shear is increased due to the effect of the Shafranov shift leading to the formation of low order rational surfaces inside the plasma. Pressure driven mode activity appears at corresponding resonant surfaces. These modes could be resistive interchange instabilities since the respective stability criterion can be violated at least in the outer part of the plasma. Only around the highest beta values and in cases, where the magnetic well of the configuration was reduced, relaxations of the plasma energy are observed, indicating the vicinity of a soft beta-limit. In most cases, however, the maximum achievable beta is determined by the available heating power.

Effect of NBI driven global Alfvén Eigenmodes on plasma confinement - Most of the MHD activity observed occurs during neutral beam injection (NBI) in the lower beta regime, where low order rational surfaces can be avoided due to the very low shear of the configuration. Away from but close to resonant surfaces gaps in the shear Alfvén spectrum are present, where global Alfvén eigenmodes (GAE) can be excited by resonant fast ions of the beam distribution. In NBI heated plasmas GAE activity often coincides with the iota-range of degraded confinement. However, this is also the range where formation of low (m,n) Alfvén gaps becomes possible. Since the resonant confinement degradation also occurs during ECRH, where no fast ion population is present and only very weak MHD-activity is observed, it is conjectured, that magnetic turbulence phenomena around the dense high (m,n) resonant surfaces are causing this effect. From the analysis of iota-profiles it is concluded, that the main low (m,n) resonances are less unfavourable because their neighbourhood is free of resonances. In the period after switch on of NBI very pronounced GAE activity causing enhanced transport can be excited due to transiently unstable fast ion velocity distributions. In particular, when the target plasma has very low density or the discharge is initiated by the
neutral beams, even more unfavourable fast ion distributions can be formed due to charge exchange effects. In this case in addition to the low frequency (20-40 kHz), low (m,n) GAE-activity a broad frequency range of activity extending up to 500 kHz is found by magnetics and various fluctuation diagnostics. The origin of this activity can be explained by Alfvén modes of higher mode numbers. GAEs, but also global Alfvén modes in gaps induced by ellipticity or higher non-symmetry (EAE, NAE) could play a role. Fig. 1 shows the change of the frequency spectrum (Mirov coil) with the change of the Alfvén speed. Simulations of a (virtual) antenna loading spectra with the CASTOR code [1] reveal a number of resonances in the frequency range of the experimental data. The high frequency activity has set in later than the low (m,n) GAE’s, which can already be excited by sub-Alfvénic beam ions via toroidal sideband excitation. For a number of cases also MHD calculations of Alfvén instabilities were performed with a gyrofluid model for the fast particles [2], which are consistent with the observed MHD activity. The analysis of mode structures, which is important for comparisons with theory, has been improved by a 10-camera soft X-ray system with 320 channels [3] and new analysis techniques for the magnetic probes [4].

Fig. 1: Global Alfvén frequency contours (top) during NBI density ramp. The dashed lines indicate the temporal evolution of the Alfvén speed. Bottom: antenna loading spectrum calculated with the CASTOR code for n=1,2,3.

Neutron Flux (exp. - code) DD-n-bar-Y/p

Fig. 2: The discrepancy between measured and calculated neutron rates increases at low collisionalities with increasing slowing down times (selected shots)
Another regime where strong effects on transport are caused by beam driven Alfven modes is at low electron collisionality achieved with combined ECRH and NBI heating. Under these conditions the classical slowing down time can be up to 50 ms exceeding clearly the energy confinement time. A transition from continuous GAE modes to bursting GAE activity occurs, which is accompanied by losses of fast particles and thermal plasma energy. This has been inferred from data obtained with deuterium injection using the neutron rate as a measure of the fast particle density since beam-target reactions are the dominating neutron production process. The magnitude of the stationary neutron flux was found to be significantly lower than predicted values (fig. 2), and in addition, relaxations of the D-D neutron rate and soft X-ray signals in correlation with the MHD bursts are observed (fig. 3). The loss rate of energetic particles, can be estimated roughly from \( \tau_{\text{fast}} = \Delta t \cdot (\Delta \phi / \phi)^{-1} \), where \( \Delta t \) and \( \Delta \phi / \phi \) are the average time between bursts and the average relative drop of the neutron flux, respectively. The fast particle confinement times derived in single cases are of the order of the slowing down time. Since resonances often occur with ions of relatively low velocities (< 1/3 injection velocity), which do not contribute to the neutron production, this analysis may underestimate the total loss rate.

**Alfven modes in the absence of fast particles** - MHD activity presumably due to GAE modes is also observed in ECRH plasmas without NBI. These modes are weaker as compared with NBI driven GAE modes and are preferentially seen under degraded confinement conditions without being the direct cause of enhanced thermal transport. Whereas during NBI the propagation is in the (fast) ion diamagnetic drift direction as expected from the ion drift excitation process, it is opposite during ECRH only. The frequencies are consistent with GAE modes inside the lowest Alfven continuum gap (fig. 4). The observations are similar to results of TFTR [5] and ASDEX Upgrade [6] where TAE modes were found in the OH phase. A possible excitation mechanism considered recently is the coupling of \( \mathbf{E} \times \mathbf{B} \) turbulence to
electromagnetic drift \( Alfvén \) turbulence cascading into low \( k_\parallel \) \( Alfvén \) waves [7,6]. The local \( \beta \)-values at the location of the modes is in the range where such coupling is expected (fig. 4, bottom). A correlation between MHD-activity, broadband turbulence and enhanced transport is found in many cases including NBI discharges.

**Conclusions** - The dependence of the confinement on the magnetic configuration cannot be explained simply by low \( (m,n) \) magnetic field resonances ("natural" islands and other static field perturbations) or mode activity. Global \( Alfvén \) modes emerge from the multiformity of MHD activity in W7-AS as most prominent instabilities. Since their propagation requires finite \( k_\parallel \) (or a gap in the continuous \( Alfvén \) spectrum), GAE modes with mode numbers \( (m,n) \) cannot be easily suppressed by avoiding the corresponding rational surface \( \iota = n/m \) in the confinement region. The effect of \( Alfvén \) modes on energy and particle transport is significant only during transient phases of bursting modes and in the low collisionality regime with NBI. Drift \( Alfvén \) turbulence is considered to be an important process to explain both, the resonant anomalous transport and the excitation of GAE's in ECRH plasmas.

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Viscous Damping and Plasma Rotation in Stellarators

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Poloidal rotation with shear flow is one of the key elements in the theory of H-mode confinement in toroidal systems. There are various driving forces which may excite poloidal and toroidal rotation: Stringer spin-up, turbulent Reynolds stresses and lost orbits. Viscous damping is the main candidate to retard the rotation. In a collision dominated plasma viscous damping is provided by the magnetic pumping effect which arises from the variation of the magnetic field strength along the stream lines of rotation. It only depends on the Fourier spectrum of $B$ on magnetic surfaces and not on the details of particle orbits. A comparison of viscous damping rates of various toroidal configurations has been given in ¹. The non-axisymmetric components in the Fourier spectrum of $B$ lead to extra maxima in the poloidal force which complicates the bifurcation problem as compared to axisymmetric tokamaks. Enhanced poloidal damping also occurs if magnetic islands exist in the confinement region. In stellarators islands occur on rational magnetic surfaces if there exists a resonant field perturbation. Such perturbations always exist on "natural surfaces with the rotational transform $\tau = M/k$, $k=1,2,3,..M$ is the number of field periods. The denominator indicates the number of islands. Symmetry breaking field perturbations introduce another class of islands which can easily dominate over the natural islands.

In general there is a combination of natural islands and symmetry breaking islands. Examples are given by the Wendelstein stellarators W 7-A and W 7-AS ² Islands lead to enhanced radial transport and therefore also to enhanced poloidal damping. In the neighbourhood of islands magnetic surfaces are modified and the Fourier spectrum of $B$ exhibits harmonics with the periodicity of the island. Due to toroidal curvature the field variation on magnetic surfaces is roughly $\delta B/B = \tau/R$.

Even if no islands are present in the plasma region magnetic surfaces at the plasma edge are modified by resonant Fourier harmonics. This in particular occurs in the boundary region of Wendelstein 7-AS, where the last magnetic surfaces exhibit the $M/k$ structure of the natural islands.


Fig. 1: Magnetic islands in the boundary region of a stellarator.

In the following we start from a magnetic field configuration with nested magnetic surfaces which satisfies the condition of ideal equilibrium. On these nested magnetic surfaces the Hamada coordinate system is introduced which is characterised by straight magnetic field and a Jacobian equal to unity. In this coordinate system the base vectors are defined by \( e_p = \nabla s \times \nabla \varphi \), \( e_t = -\nabla s \times \nabla \theta \). \( e_p \) is the poloidal base vector and \( e_t \) the toroidal base vector. \( s \) is the volume of the magnetic surface. The lowest order rotation \( \nu_0 = -E(\psi)e_p + \Lambda(\psi)B \) stays on magnetic surfaces and satisfies the equation \( \nabla \nu_0 = 0 \). The two flux functions \( E \) and \( \Lambda \) describe the poloidal and the parallel motion on magnetic surfaces. \( E \) is the radial electric field.

In a toroidal plasma the forces by magnetic pumping inhibit the rotation of the plasma, these forces are

\[
< e_p \cdot \nabla \varpi > = < (p_t - p_i) e_p \cdot \frac{\nabla B}{B} > ; \quad < B \cdot \nabla \varpi > = < (p_t - p_i) B \cdot \frac{\nabla B}{B} >
\]

In a collisional plasma these equations reduce to (see ref. 1)

\[
\begin{pmatrix}
-< e_p \cdot \nabla \varpi > \\
< B \cdot \nabla \varpi >
\end{pmatrix}
= 3\tau P \begin{pmatrix}
C_p & C_b \\
C_b & C_t
\end{pmatrix}
\begin{pmatrix}
E \\
\Lambda
\end{pmatrix}
\]

\( P \) = plasma pressure, \( \tau \) = collision time. The coefficients are

\[
C_p = < (e_p \cdot \frac{\nabla B}{B})^2 > ; \quad C_t = < (B \cdot \frac{\nabla B}{B})^2 > ; \quad C_b = < (e_p \cdot \frac{\nabla B}{B}) (B \cdot \frac{\nabla B}{B}) >
\]

Thus the damping forces are the product of plasma parameters and geometrical coefficients. In the following numerical computations of these coefficients \( C_p, C_t, \) and \( C_b \) are shown. The poloidal coefficient \( C_p \) is of particular interest, since a small poloidal damping facilitates poloidal rotation and poloidal shear flow. The results of numerical calculations show large damping coefficients in standard l=2 stellarators and torsatrons in contrast to optimised stellarators of the Helias type.
Fig. 2: Poloidal damping $C_p$ coefficient in $l=2$ stellarators and the optimised stellarator W 7-X. In W 7-X the coefficient is roughly a factor 2 smaller than in W 7-AS and standard stellarators. A similar result exists for the threshold $D=C_p - C_b \cdot C_p / C_t$ of poloidal spin-up. These coefficients increase towards the plasma boundary. However, since these coefficients are multiplied by $\tau P$ which decreases strongly towards the boundary the viscous forces are small in the boundary region.

This may explain why plasma rotation is mainly observed in the boundary regions.

Magnetic islands may arise in the boundary region of some stellarators like Wendelstein 7-AS. These islands distort the magnetic surfaces in the neighbourhood and lead to an enhanced viscous damping (or magnetic pumping). Inside the island region the plasma equilibrium is not represented by the ideal MHD-theory, therefore the coefficients given here are only relevant in the region outside magnetic islands.

Fig. 3: Poloidal viscous damping rate $C_p$ (collisional regime) in Wendelstein 7-AS. The regime of the rotational transform is between 5/10 and 5/9. The abscissa is the average radius of the magnetic surface. Close to the island the damping rate is enhanced by a factor of 2. The coefficient $C_p$ exhibits a similar enhancement. Little effect of magnetic islands is seen in the toroidal coefficients $C_t$. In conclusion, mainly poloidal rotation will be inhibited by magnetic islands.

The standard case of Wendelstein 7-X has 5 islands at the boundary. These islands are the base of the divertor concept and the issue arises whether these islands suppress plasma rotation in this region. The following figure shows that only in a small neighbourhood of these islands enhancement of the poloidal damping coefficient occurs.

Fig. 4: Comparison with tokamak. Poloidal viscous coefficient. The gray rectangular region is the island at the boundary of W 7-X. Enhancement of viscous damping in W 7-X
occurs in a few cm distance from the islands. The two curves of W 7-AS differ by the rotational transform. Poloidal damping in W 7-X and a tokamak (Tore Supra) are nearly equal at equal plasma radii.

**Discussion**

The viscous damping of poloidal and toroidal plasma rotation depends on two factors, one factor is $\tau P$ - collision time times plasma pressure - and the other factors are geometrical factors $C_P$, $C_t$ and $C_b$. With respect to the geometrical factor significant differences among the various stellarators exist. $C_p$ is the relevant geometrical factor of poloidal rotation. In axisymmetric devices only the toroidal curvature effect gives rise to magnetic pumping and a finite coefficient $C_p$. In stellarators, however, additional helical harmonics lead to increased poloidal viscous damping. Reducing the poloidal variation of $B$ - as has been done in optimised configurations of the Helias type - also reduces the poloidal viscous damping. In the neighbourhood of magnetic islands enhanced viscous damping arises due to the distortion of magnetic surfaces. This effect may be of importance in Wendelstein 7-AS, where magnetic islands in the boundary region strongly corrugate the surfaces. The viscous damping coefficients $C_p$ and $C_b$ are enhanced in the neighbourhood of the islands. In comparison to axisymmetric configurations standard stellarators exhibit a more complex Fourier spectrum of $B$ and therefore magnetic pumping effect is larger. The optimisation scheme realised in Wendelstein 7-X, however, reduces the poloidal damping rate to the level of axisymmetric configurations. As seen in Fig. 4 the coefficient $C_p$ is nearly the same in both configurations. The present analysis is based on a collisional plasma model. Such a model is applicable to a plasma boundary, where the temperatures are below 100 eV and the densities above $10^{19}$ m$^{-3}$. At higher temperatures the plateau regime is reached and the viscous damping must be computed from a kinetic equation. This will be done in a subsequent paper.
Debye Length in a Neutral-beam-heated Plasma

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1. Introduction

One of the major auxiliary heating methods for tokamak fusion plasmas is neutral beam injection (NBI). The fast injected neutrals become ionized and slow down to thermal energy through Coulomb collisions with the background plasma. As a consequence of the injection, the velocity distribution of the injected ion species consists of a Maxwellian part and a non-Maxwellian high-energy part. The evolution of the velocity distribution \( f \) of the neutral-beam-injected ions can be described by the following kinetic equation:

\[
\frac{\partial f}{\partial t} + v \cdot \nabla_r f + \frac{eZ}{m} \nabla \Phi \cdot \nabla_v f = C(f) + S - L,
\]

(1)

where \( \nabla_r \) and \( \nabla_v \) are the gradients with respect to coordinate and velocity, respectively, \( C(f) \) is the collision operator, \( S \) is the fast ion source term and \( L \) is a loss term for ions. Furthermore, \( eZ\nabla \Phi/m \) is the acceleration of a particle due to electrostatic collisions in a quasi-neutral multi-species plasma, where \( Z \) and \( m \) are the charge and the mass of the particle and \( \Phi \) is the resulting potential of the macroscopic field.

If a test charge \( Q \) is placed at the origin in the plasma, its effect is to attract particles of opposite sign and to repel charges of the same sign. The resulting modifications of the ion and electron velocity distributions are simply positive and negative shifts of energy \( E_\Phi = eZ\Phi \). The potential \( \Phi \) follows from the Poisson equation

\[
\nabla^2 \Phi = -\frac{e}{\epsilon_0} \left[ \sum_i Z_i \int f_i(v, \Phi) \, d^3v - \int f_e(v, \Phi) \, d^3v + Q\delta(r) \right],
\]

(2)

where \( Z_i \) and \( n_i \) are the charge and density of the \( i \)th ion species and \( f_e \) and \( f_i \) are the electron and ion velocity distributions normalized to the density \( n_{e,i} \) as \( \int f_{e,i}(v) \, d^3v = n_{e,i} \).

The Poisson equation in the general form given above is impracticable to solve analytically. In deriving the expression for the Debye length in a plasma it is therefore commonly assumed that the velocity distributions of the plasma ions and electrons are isotropic Maxwellsians with temperatures \( T_i \) and \( T_e \), respectively. Thus, effects of plasma heating due to fast particle injection (and other mechanisms leading to highly non-Maxwellian velocity distributions) are neglected. The resulting Poisson equation

\[
\nabla^2 \Phi = -\frac{e}{\epsilon_0} \left[ \sum_i n_i Z_i \exp \left( -\frac{eZ_i\Phi}{kT_i} \right) - n_e \exp \left( \frac{e\Phi}{kT_e} \right) + Q\delta(r) \right],
\]

(3)

where \( k \) is the Boltzmann constant, can be linearized if the perturbation is small, i.e. if \( eZ_i\Phi \ll kT_i \) and \( e\Phi \ll kT_e \). From the spherically symmetric solution of equation (3), one obtains for the Debye length: \( \lambda_D = (\Sigma_i k_i^2 + k_e^2)^{-1/2} \), where \( k_i^2 = e^2 Z_i^2 n_i/\epsilon_0 kT_i \) and \( k_e^2 = e^2 n_e/\epsilon_0 kT_e \) are the Debye constants for Maxwellian ions and electrons, respectively.
2. Debye Length for a Plasma with Neutral-Beam-Injected Ions

In deriving an expression for the effective Debye length in a plasma with NBI slowing-down distributions, it is reasonable to assume that the electron and background ion velocity distributions are Maxwellians. Since the main effect in the calculation of the effective Debye length is the existence of a high-energy tail and the result does not markedly depend on whether the velocity distribution is isotropic or anisotropic, the non-Maxwellian velocity distribution of the injected ions is assumed to be isotropic. Furthermore, for simplicity, the effective screening length will be derived for steady-state conditions. If the subscript '0' refers to the injected ion species and \( f_0 \) is the corresponding isotropic non-Maxwellian velocity distribution, the Poisson equation can be written as:

\[
\nabla^2 \Phi = -\frac{e}{\epsilon_0} \left[ Z_0 \int f_0(v, \Phi) \, d^3v + \sum_{i \neq 0} n_i Z_i \exp \left( -\frac{e Z_i \Phi}{kT_i} \right) - n_e \exp \left( \frac{e \Phi}{kT_e} \right) + Q \delta(r) \right].
\]

(4)

An appropriate non-thermal velocity distribution of the injected ion species can approximately be obtained from the steady-state kinetic equation with the Fokker-Planck collision operator in its high-energy approximation neglecting diffusive terms and loss terms [1]. The complete solution for the velocity distribution of the injected ion species (which is normalized to the density \( n_0 \)) is the sum of the Maxwellian homogenous solution \( f_{th} \) of the kinetic equation and the approximate particular (non-Maxwellian) solution \( f_0 \):

\[
f_0(v) = f_{th}(v) + f_i(v) \approx \frac{n_0 - n_b}{\pi^{3/2} v_{th}^3} \exp(-v^2/v_{th}^2) + \frac{s_0 \tau_e \sigma (v_0 - v)}{4\pi} \frac{m_0 C_e}{m_e} \left( \frac{2kT_e}{m_e} \right)^{3/2},
\]

(5)

where the non-thermal density \( n_b \) and the slowing-down time \( \tau_e \) of ions on electrons are given by:

\[
n_b = s_0 \tau_e \frac{1}{3} \log \left( 1 + \frac{v_0^3}{v_{th}^3} \right), \quad \tau_e = \frac{3}{2} \pi^{1/2} \frac{m_e}{m_0 C_e} \left( \frac{2kT_e}{m_e} \right)^{3/2}.
\]

(6)

with \( C_e = 8\pi n_e Ze^2 e \) and \( v_0 \) is given by:

\[
E_{\omega} = \frac{1}{2} m_0 v_0^2 = 14.8kT_e \left[ \frac{A_0^{3/2}}{n_e} \sum_{i} \frac{n_i Z_i^2}{A_i} \right]^{2/3}.
\]

(7)

Furthermore, \( v_0 \) and \( s_0 \) are the injection velocity and source rate, respectively, \( v_{th} \) is the thermal velocity, \( \sigma \) is the unit step function, \( \log \Lambda \) is the Coulomb logarithm and \( A_i \) denotes atomic masses. Again, the effect of the test charge on the velocity distribution of the injected ion species results in a positive shift of energy \( E_{\Phi} = e Z_0 \Phi = \frac{1}{2} m_0 v_0^2 \):

\[
f_0(v, \Phi) = \frac{n_0 - n_b}{\pi^{3/2} v_{th}^3} \exp(-v^2/v_{th}^2) \exp(-\frac{e Z_0 \Phi}{kT_0}) + \frac{s_0 \tau_e \sigma (v_0 - v - \Phi - v)}{4\pi} \frac{m_0 C_e}{m_e} \left( \frac{2kT_e}{m_e} \right)^{3/2} + v_{th}^2.
\]

(8)

For small perturbations, \( (v^2 + v_{th}^2)^{3/2} \approx v^2 + \frac{3}{2} v_{th}^2 \) and the denominator can be expanded. One obtains an approximate expression for \( f_0 \) which can be introduced into the Poisson equation (4). After linearization and integration over the velocity domains one obtains a
simplified Poisson equation. Neglecting the term $3^{-1/2} \tan^{-1} 3^{-1/2}$ in that equation, the Debye constant $k_b^2$ which accounts for the fast particle contribution due to NBI heating can be written as:

$$k_b^2 = \frac{e^2 Z_0 s_0 T_e}{e_0 2 m_0 v_a^2} \left[ \frac{2 Z_0 m_0 v_a^2}{3 k T_0} \log \left( 1 + \frac{v_0^3}{v_a^3} \right) + \frac{v_0 v_a^2}{v_0^3 + v_a^3} - \frac{1}{6} \log \left( \frac{v_0 + v_a}{v_0^3 + v_a^3} \right) \frac{1}{\sqrt{3}} \tan^{-1} \frac{2 v_0 - v_a}{v_a \sqrt{3}} \right].$$

(9)

The Debye screening length for a neutral-beam heated plasma can, in parallel with the above, be written as:

$$\lambda_{\text{NBI}} = \left( \sum_i k_i^2 + k_e^2 - k_b^2 \right)^{-1/2}.$$  

(10)

This equation is, however, only valid for completely static screening. In reality, the ions are moving very slowly. Therefore they are completely screened by the fast moving electrons, but they are only partially screened by surrounding ions. In order to discuss effects of dynamical screening one can use the results given in [2] for a Maxwellian plasma. Generalizing the expression for the effective screening length given in [2] to NBI heated plasmas, one obtains:

$$\lambda_{\text{NBI}}^{\text{eff}} = \left[ k_e^2 \exp \left( \sum_i k_i^2 + k_e^2 - k_b^2 \log \left( \frac{\sum k_i^2 + k_b^2 - k_b^2}{k_e^2} \right) - 1 \right) \right]^{-1/2},$$

(11)

where the summation is over all ion species. This effective Debye length is in between the total Debye length for static screening in a Maxwellian plasma $\lambda_D$ and the one for static screening of Maxwellian electrons $\lambda_e = 1/k_e$.

3. Results and Conclusions

In order to illustrate the effects of a non-Maxwellian NBI slowing-down distribution on the screening distance, a pure deuterium plasma with temperatures $T_i = T_e$ ranging from 5 keV to 8 keV and $D^0$-injection with varying source rates and injection velocities is considered. To be independent of the densities, only the ratios of the Debye lengths, viz $\lambda_{\text{NBI}}/\lambda_D$ and $\lambda_{\text{NBI}}^{\text{eff}}/\lambda_D$, are calculated as functions of the fast particle fraction. The results are shown in figure 1 for a plasma with 80 keV injection energy and temperatures of 5 keV and 8 keV, respectively. It can be seen that for reasonable plasma parameters with fast particle fractions below about 40%, the Debye lengths for completely static screening differ by up to about 12%. For more extreme plasma parameters differences by up to 20% can occur. In addition, the effect of dynamical screening is also shown in figure 1. For the same plasma parameters and fractional fast particle densities below 40%, the differences to the Maxwellian case with static screening are about 17-25% while for more extreme plasma differences around 30% can occur. Generally, the differences are somewhat increasing with increasing temperatures. The dependence on the injection energy is logarithmically weak and with increasing injection velocity, the ratios of
the Debye lengths are slightly decreasing. (This is because the fast particle fraction is kept constant and therefore with increasing injection velocity, the Maxwellian part of the velocity distribution extends to higher velocities.) As described in [3], the Coulomb logarithm for the interaction of a test particle with a background particle is in the classical limit \( \log \Lambda = \log \lambda_D - \log \rho_\perp \), where \( \rho_\perp \) is the impact parameter. Therefore, although the improved Debye length differs by about 20% from the Debye length for static screening in a Maxwellian plasma, the final effect on the Coulomb logarithm is small. For parameters relevant for fusion plasmas, where the Coulomb logarithm is in the order of 20, the modifications are only in the order of 1%. In conclusion, the effects of non-Maxwellian slowing-down velocity distributions on the Debye length and, in particular, on the Coulomb logarithm may be neglected in many cases without substantial loss of accuracy.

Figure 1. Ratios of the NBI Debye lengths for completely static (s) and for dynamical screening (d) to the Maxwellian Debye lengths (\( \Lambda_{NBI}/\lambda_D \) and \( \Lambda_{NBI}/\lambda_D \)) as functions of the fractional fast particle density for a deuterium plasma with 80 keV deuterium injection and temperatures of 5 keV (—) and 8 keV (---).

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RECENT OBSERVATIONS
OF MHD INSTABILITIES ON W7-AS

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Abstract
Magnetic fluctuations are observed in ECRH plasmas on the W7-AS stellarator which can not be identified with any known type of MHD mode. We investigate this type of fluctuation using cross correlation and Fourier methods as well as phase space analysis from dynamical systems theory. In spite of its weakly turbulent character, a spatial structure can be identified. The development of the spatial structure coincides with a decrease of the energy confinement time.

Keywords: Mirnov probes, magnetic fluctuations, MHD instabilities, weakly developed turbulence, phase space analysis

1 Introduction
Magnetic fluctuations on W7-AS are measured using poloidal arrays of 8 or 16 Mirnov coils at different toroidal locations. Sampling rates are 250 – 330 kHz. The spatio-temporal behaviour of the Mirnov data is routinely analysed using Fourier techniques, singular value decomposition and cross correlation methods [1]. In addition, we use Lomb’s normalized periodogram [2] in order to determine the poloidal structure. Instead of the poloidal mode number m which is well-defined only for rational values of the rotational transform \( \alpha \)
we use the notation \( k_\theta r_f \), where \( k_\theta \) is the poloidal wavevector and \( r_f \) the radial position of the perturbation.

This is motivated by the facts that a) it is an open question whether the signals observed with magnetic diagnostics in electron cyclotron resonance (ECR) heated plasmas are due to coherent mode activity or to fluctuations in a narrower sense and b) there are usually no low \( m, n \) rational values of \( t \) inside the last closed magnetic surface.

If we assume a current perturbation along magnetic field lines, we obtain a heuristic formula for the frequency of the magnetic fluctuations \( \hat{B}_\theta \) in the laboratory frame which reads

\[
\omega_{lab} = \frac{v_{\text{dia}}(r_f) + v_{E\times B}(r_f)}{2\pi r_f} \cdot k_\theta r_f
\]

where \( v_{\text{dia}} \) is the electron diamagnetic drift velocity and \( v_{E\times B} \) is the \( \vec{E} \times \vec{B} \) drift speed of the electrons. Experimentally, \( v_{\text{dia}} \) is obtained from temperature and density profiles measured by Thomson scattering whereas \( v_{E\times B} \) is determined using charge exchange recombination spectroscopy [3].

## 2 Observations

Series of ECR heated discharges have been produced in order to investigate the dependence of the energy confinement time \( \tau_E \) on the rotational transform at the plasma boundary \( \tau \). Parameters of these discharges are \( 0.33 < \tau < 0.36, B_\phi = 2.5 T, 0 < B_z < 22 mT, P_{\text{ECRH}} \approx 480 kW \) and \( n_e \approx 4 \cdot 10^{19} m^{-2} \).

Fig. 1 displays the normalized cross correlation function [2] of Mirnov signals as a function of the time lag \( \tau \) and the poloidal position of the probes. The energy confinement time of discharge \# 38806 (\( \tau = 0.339, B_z = 0 \)) is twice larger than \( \tau_E \) of \# 38808 (\( \tau = 0.345, B_z = 0 \)). For \# 38808, a pronounced poloidal structure with \( k_\theta r_f \approx 3 \) rotating in the electron diamagnetic drift direction is observed. The power spectrum of \# 38806 is broader compared to \# 38808 (see fig.5) with a peak between 10 – 20 kHz.
In fig. 2, we compare $f_{lab}$ eqn. (1) to Mirnov spectra for two discharges with $\tau_a = 0.349$ (# 39218) and $\tau_a = 0.361$ (# 39235, both with $B_z = 220 mT$). As was the case for # 38806 and # 38808, the spectrum of the discharge with the lower $\tau_E$ exhibits more structure and a lower average frequency (# 39235, darker curves in fig.2). Note the low “shear” of $f_{lab}$ for # 39235, which possibly explains the relatively narrow peak in the Mirnov spectrum at $f \approx 20 kHz$.

In a similar discharge (# 39911, $B_0 = 2.5 T, \tau_a = 0.345, P_{ECRH} = 350 kW, n_t \approx 2 \cdot 10^{19} m^{-2}$) temperature fluctuations were observed using electron cyclotron emission (ECE) diagnostics [4]. Magnetic and temperature fluctuations at $f \approx 40 kHz$ are correlated. The maximum amplitude of these fluctuations is localised at an effective radius of $r_{eff} \approx 7 cm$. Frequency and radial position are consistent with eqn.(1) as seen from fig.3. The poloidal structure inferred from Mirnov data is $k_\phi r_f \approx 3$.

In discharge #40365 ($\tau_a = 0.523, P_{ECRH} = 450 kW$) density fluctuations with a frequency of $\approx 7 kHz$ were observed using a Li-beam probe [5]. They extend $1 cm$ inside and $3 - 4 cm$ outside the last closed magnetic surface (fig. 4). The poloidal structure of the magnetic perturbations in this frequency range is $k_\phi r_f \approx 2$, again consistent with eq.(1).

3 Phase Space Analysis

Broadened but structured (i.e. non-power law) spectra are considered as an indicator for weakly developed turbulence [6]. Dynamical systems theory provides numerous instruments to characterize temporal behaviour [7].

We have applied an embedding procedure the basic idea of which is to switch from time evolution to geometry in phase space, where the time enters as a parameter only. The basic assumption is that the dynamics is governed by a set of autonomous ordinary differential equations, where the full set of time-dependent variables span the phase space. This allows to distinguish between noise and strong
tubulence (almost infinite dimensional), weakly developed turbulence (high dimensional), and regular dynamics (low dimensional).

The phase space geometry (the attractor) is reconstructed from the time evolution of $X(t)$,

$$X(t) = [x(t), x(t + \tau), x(t + 2\tau), \ldots, x(t + (d - 1)\tau)]$$  \hspace{1cm} (2)

where $x(t)$ is the time series of one single Mirnov coil signal. To perform an optimal embedding, we apply the fill-factor criterion described in Ref. [8] which yields the time lag $\tau$. The embedding dimension $d$ is estimated by calculating the correlation integral [9].

We have investigated time series of 32k length each from discharges # 38806, # 38808 (compare fig. 1) and # 39727 ($B_\phi = 2.5 T, a = 0.334, P_{ECRH} = 350 kW, n_t = 5 \cdot 10^{18} m^{-2}$). The power spectra (signal and noise) and the visualization of the attractors are displayed in fig. 5.

The phase space 5 (a) shows just a cloud of scattered points hinting stochastic behavior. We note however in 5 (b) a certain distortion of the phase space structure. Both data sets (a) and (b) are unlikely to correspond to a coherent mode structure, it rather appears to be a weakly turbulent regime.

In fig. 5(c) fluctuation data of a state where a coherent mode is established are shown. Correspondingly, the power spectrum is sharply peaked. We obtain a phase space reconstruction that clearly shows a limit cycle, the phase space attractor of a regular oscillation.

The conclusion that both data sets (a) and (b) are not due to a coherent mode structure, in contrast to (c), is confirmed by a detailed analysis of the corresponding fill factors and correlation integrals.

4 Summary

MHD activity in ECR heated plasmas of W7-AS appears to be governed by weakly turbulent behaviour. This turbulence can exhibit a pronounced poloidal structure and can be correlated with temper-
ature and density fluctuations. Changes of frequency or poloidal structure of the magnetic perturbations which coincide with a change of the energy confinement have been observed. Whether the magnetic fluctuations are relevant for transport remains the subject of future work.

References


Figure 1: Left: # 38806, right: #38808. First row: cross correlation as a function of time lag $\tau$ (1 sample = 5$\mu$s) and poloidal position, second row: $cc$ of two neighbouring probes, third row: Lomb periodogram of $cc(\tau = 0, \theta)$

Figure 2: $f_{sa}$ eqn (1) (left) compared to Mirnov spectra (right) for # 39218 (light curves) and # 39235 (dark curves)

Figure 3: like 2, but for # 39911, bars indicating the radial position and frequency of $T_e$--fluctuations correlated with $\dot{B}_\phi$

Figure 4: like 3, but for # 40365, bars indicating radial position and frequency of $n_e$--fluctuations correlated with $\dot{B}_\phi$

Figure 5: Power spectra of Mirnov signals and noise (first row) and phase space diagrams for three discharges. (a) left: # 38806 , (b) middle: # 38808, (c) right: # 39727
Fig. 1

Fig. 2

Fig. 3

Fig. 4

The analysis and optimization of drift losses by the so-called optimization procedure (1) was performed in such a manner that these drifts were minimized. Another scenario of radial losses in the LMF2 occurs as a consequence of a change in the magnetic field $B$, which develops in accordance to the...
fig. 5
Overview on the radial electric field, plasma rotation and transport in the stellarator W7-AS

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Abstract
In the advanced stellarator W7-AS the radial electric field $E_r$ is measured by active charge exchange recombination spectroscopy CXRS. In parallel, it is calculated by using the neoclassical DKES code. A comparison of calculated and measured solutions reveals in how far the neoclassical model is valid for the description of the radial particle transport and the formation of $E_r$. In general good consistency is found, even for the outer radii where the neoclassical fluxes become much smaller than the experimental ones. In this paper the interplay between the particular $E_r$ roots and transport is considered. For strongly positive $E_r$ a reduction of $\chi_e$ is observed in the vicinity of the magnetic axis. The typically negative ion-root in the gradient region strongly influences the local $\triangledown T_i$, thus determining the maximum attainable $T_i(0)$.

Keywords: Radial electric field, spectroscopy, neoclassical transport, W7-AS, ion root, electron root.

I. Introduction
Stellarators suffer, in comparison to tokamaks, from enhanced neoclassical transport especially in the long mean free path LMFP regime. There, large populations of trapped particles are subject to drift losses. This unfavourable effect is partly compensated by the so-called optimization procedure [1], where the magnetic configuration is tailored in such a manner that those drifts are minimized. Another reduction of radial losses in the LMFP occurs as a consequence of a radial electric field $E_r$ which develops in accordance to the ambipolarity condition:
\[ \Gamma_i(r, E_r) + \Gamma_i(r, E_r) + Z_i \Gamma_i(r, E_r) = 0 \] (1)

Following neoclassical theory, the radial particle fluxes \( \Gamma \) and heat fluxes \( Q \) are functions of \( E_r \) [2]:

\[ \Gamma_{\alpha} = -n_{\alpha} \left\{ D_{11}^{\alpha} \left( \frac{n_{\alpha} - Z_{\alpha} E_r}{T_{\alpha}} \right) + D_{12}^{\alpha} \frac{T_{\alpha}'}{T_{\alpha}} \right\} \] (2.1)

\[ Q_{\alpha} = -n_{\alpha} T_{\alpha} \left\{ D_{21}^{\alpha} \left( \frac{n_{\alpha} - Z_{\alpha} E_r}{T_{\alpha}} \right) + D_{22}^{\alpha} \frac{T_{\alpha}'}{T_{\alpha}} \right\} \] (2.2)

In LMFP, the transport coefficients \( D_{ij}^{\alpha}(E_r) \) for electrons \( \alpha = e \) and ions \( \alpha = i \) are reduced, in the plateau regime only those of the ions. \( Z \) is the charge, \( n \) the density, \( T \) the temperature. A possible Ware-pinich contribution in eq. (2.1) proportional to \( D_{13} \) is neglected, because \( E_{II} \) is very small as a consequence of the net current free stellarator operation. To calculate the neoclassical \( E_r \), the DKES code [3] is used which takes into account the complex magnetic configuration of W7-AS. Starting from a Fourier representation of the magnetic field, DKES calculates the monoenergetic transport coefficients \( D_{ij}^{\alpha} \) to solve eq. (1). These solutions are the neoclassical roots of the ambipolar \( E_r \).

The most positive is referred to as electron-root, the most negative as ion-root, solutions between them are not stable. The total number of all roots is always odd [4]. DKES, however, might fail in the estimation of \( \Gamma_i \) for very strong \( E_r \), because in that case variations of \( n_i \) and the electrostatic potential on flux surfaces might appear which are not taken into account. The strongly positive electron-root should be accessible for the case that the electrons are in the LMFP, and the condition \( \Gamma_e(E_r=0) > \Gamma_i(E_r=0) \) holds. This is typically the case in W7-AS for low \( n_e(0) < 3 \cdot 10^{19} \ m^{-3} \) and high \( T_e \). For higher \( n_e(0) \) only the ion-root is expected with \( \Gamma_e(E_r=0) < \Gamma_i(E_r=0) \). It is characterized by small negative values near the magnetic axis and strongly negative values in the gradient region.

Charge exchange recombination spectroscopy CXRS [6] in a modulated diagnostic neutral beam is used in W7-AS [7] to determine radial profiles of the toroidal and poloidal rotation velocity \( V \) [5]. From the simplified radial force balance [8]:

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\[ E_r = \frac{\partial(n_{I+1}(r) \cdot T_{I+1}(r))}{\partial r} \cdot \frac{1}{eZ_{I+1}n_{I+1}(r)} + \frac{T_I(r)\gamma}{eZ_I} \frac{\partial \zeta_I}{\zeta_I} + (B_\theta V_\phi - B_\phi V_\theta) \] (3)

the radial electric field profile \( E_r(r) \) is determined. CXRS also provides the impurity density \( n_I(r) \) and ion temperature \( T_I(r) \) after CX which changes \( Z \) from \( I+1 \) to \( I \). \( B \) is the magnetic field, \( \zeta \) the electronic state excitation probability, \( e \) the electron charge. With the ion gyro-frequency \( \omega \) and the mean excited state lifetime \( \tau \) we define \( \gamma = (\omega \tau)^2 / (1 + (\omega \tau)^2) \). In general it is found that the major contribution to \( E_r \) comes from the poloidal rotation (50% - 80%). \( V_\theta \) dominates because of the relatively small magnetic pumping due to the large aspect ratio, and because of the small mean VB as a consequence of the optimization procedure. Only a minor contribution comes from the ion diamagnetic pressure term (15% - 45%). \( V_\psi \) typically contributes to less than 5%, because \( B_\theta V_\psi << B_\phi V_\theta \). Furthermore, the toroidal viscous damping caused by collisions between passing and trapped particles is very large due to the missing axi-symmetry. A fast toroidal rotation is found only in discharges with non-balanced neutral beam heating NBI, but even then its contribution to \( E_r \) is negligible. The value of \( V_\phi \) is described well by neoclassical toroidal viscosity [9].

The comparison of neoclassical and measured \( E_r \) is also a sensitive test in how far possible anomalous fluxes near the plasma edge are intrinsically ambipolar.

II. Electron root
For a stellarator reactor, a starting scenario in the electron-root is considered [10]. In particular, for this case in eq. (2.1) the positive linear \( E_r \) term dominates [2], thus leading to an enhanced outward drift of impurities. Accumulation of impurities, especially of the Helium ash, could thus be prevented. In W7-AS in fact hollow He\(^{++}\) density profiles are measured [11] for positive \( E_r \). Fig. 1 shows W7-AS profiles for a discharge with an \( E_r \) profile resembling the neoclassical electron-root. This is a low density discharge at high on-axis electron cyclotron heating ECRH power (770 kW 2nd harmonic
140 GHz X-mode). A strong positive $E_r$ develops near the plasma axis, leading to a considerable reduction of the central $\chi_e$, as revealed by the experimental power balance analysis. Thus a maximum $T_e(0)$ up to 4 keV is obtained, the largest measured so far in W7-AS. The appearance of the strongly positive $E_r$ is accompanied by a local increase of $\nabla T_e(r \geq 3\text{cm})$, directly indicating the locally reduced electron heat transport. That local transport reduction can be seen in the plot on the lower right, where the experimental $\chi_e$ near the plasma center comes closer to the neoclassically expected one from DKES for the $E_r$ electron-root solution than for $E_r = 0$. The measured $E_r$ from CXRS and the neoclassically calculated agree well. For higher $n_e(0)$ or reduced heating power the central positive $E_r$ disappears. It also disappears for the case that the local B-field minimum in the ECRH launching plane is removed by modifications of the magnetic configuration, or by shifting the ECRH power deposition from on-axis away towards off-axis [14]. It is assumed that drifts of ripple trapped suprathermal electrons in the ECRH launching plane contribute strongly to the radial electrons fluxes, additionally to the thermal ones. This is supported by the much faster temporal decay of the central $T_e$ in comparison to the outer $T_e$ after switching off the ECRH power at the end of the discharge.

III. Ion root
The second discharge, shown in fig. 2, is characterized by combined ECRH (750 kW at 140 GHz) + neutral beam NBI (680 kW absorbed) high heating power at medium $n_e(0) = 6 \cdot 10^{19} \text{ m}^{-3}$. Typically, only the ion-root solution establishes for this type of discharge. The values for $E_r$ and $\nabla T_i(r \geq 14\text{cm})$ shown in this example are the largest observed in W7-AS so far. As a result of the steep $\nabla T_i$, $T_i(0)$ reaches maximum values. Up to 1.5 keV are found [16]. The strong $E_r$ in the gradient region might act as a potential barrier for the ions, coupling the ion heat fluxes strongly to the electron ones. The transport analysis reveals that both heat and particle fluxes are described well by neoclassical theory, at least up to about 12 cm. The central particle fluxes are well consistent with the values obtained from FAFNER.
Further outside, recycling neutral gas from the vessel wall provides an additional particle source increasing locally the anomalous particle fluxes. These fluxes contribute strongly to the convective heat fluxes, therefore higher particle sources in the gradient region could limit the maximum attainable $\nabla T_i (r \equiv 14 \text{ cm})$ and $T_i(0)$ as well as the maximum $E_r$ for the total heating power available. That critical limit, however, is not yet reached for the discharge shown. Sufficiently low neutral particle sources in the gradient region seem to be the pre-requisite to obtain these strong $E_r$ in conjunction with maximum $T_i(0)$. Furthermore, the experimental energy confinement time $\tau_E$ exceeds the prediction from the ISS95 regression database [13] for W7-AS by roughly a factor of two. That might, to some extent, also be due to the strongly sheared poloidal rotation in the gradient region, leading to a reduction of the correlation length of electron density fluctuations [17], thus reducing the turbulent anomalous transport. The reason for the deviation between measured and calculated $E_r$ near the plasma periphery is subject to further investigations. Several mechanisms as fast ion orbit losses, helical resonances, non-ambipolar CX ion fluxes are quantitatively considered but found not to be sufficient as explanation [15]. The DKES ordering, however, might fail for these strong $E_r$. It is emphasized that the deviation shown is not at all typical for W7-AS: about 40 discharges have been investigated so far and good agreement is found (also in that radial range). That holds even when the neoclassical particle fluxes are much smaller than the experimental ones. Therefore we conclude that any anomalous fluxes seem to be intrinsically ambipolar.

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Figure Captions

Fig. 1 - Upper left: electron (solid) + ion (broken) temperature; upper right: electron density; lower left: \( E_r \) from CXRS (dots), from DKES: stable roots (crosses), instable solution (dots); lower right: experimental \( \chi_e \) (dot-dashed), neoclassical for stable roots (crosses), for \( E_r = 0 \) (dots).

Fig. 2 - Upper left: electron (solid) + ion (broken) temperature; upper middle: electron density; upper right: \( E_r \) from CXRS (dots), from DKES (crosses); lower left: particle fluxes from DKES (solid) and experimental (dot-dashed); lower middle: neoclassical ion heat fluxes (dotted + dashed), from power balance (dot-dashed); lower right: neoclassical electron heat fluxes (solid + dotted), from power balance (dot-dashed).
The role of magnetic shear
in the confinement of W7-AS plasmas

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Abstract

Magnetic shear has systematically been varied by inductive current
drive in moderate $\beta$, ECR heated W7-AS discharges. Confinement is
found to depend strongly on magnetic shear in the presence of high
order rational values of the rotational transform. For optimum
confinement, high order resonances have to be avoided in the boundary
region, or shear must be sufficiently large. The observed improvement
of degraded confinement with increasing shear is associated with a
reduction of transport to the neoclassical level.

Keywords

W7-AS, stellarator, magnetic shear, confinement, transport

1. Magnetic shear in W7-AS confinement

Shear in the magnetic field is generally expected to reduce
anomalous transport by radial decorrelation of turbulent structures. In
tokamaks, a synergy of improved stability at low or reversed (in the
tokamak sense, i.e. $dq/dr < 0$) magnetic shear and $E \times B$ flow shear is
supposed to give rise to the recently observed internal transport barriers
[1]. In particular in stellarators, the quality of flux surfaces and their
susceptibility to magnetic perturbations sensitively depends on the
magnetic shear.
In Wendelstein 7-AS (W7-AS) the rotational transform τ of the vacuum field has very low shear which, at finite β, is modified by pressure driven (bootstrap- and, in W7-AS reduced but still significant, Pfirsch-Schlüter (PS)-current) and, optionally, by externally driven plasma currents. Due to the small shear, magnetic islands of considerable size may form at rational τ-values being in resonance with the major external field perturbations ("natural" 5/m-components from the 5-fold toroidal symmetry, and 1/3- and 1/2-components from field errors) [2]. With increasing β, the sensitivity of confinement to these resonances diminishes, giving evidence of a beneficial effect of pressure induced shear [3]. Although the major resonances play a role, the detailed dependence of W7-AS confinement on the boundary value τ_a of the rotational transform cannot be explained by the external field perturbations [2]. Therefore perturbations at the higher order τ = n/m resonances may be important. Their impact will depend on the magnetic shear.

In the present study, shear has been applied by inductive current drive to ECR heated discharges at moderate β (β_0 < 1%). The experiments are interpreted on the basis of τ-profiles calculated by the NEMEC equilibrium code [4], with inductive and bootstrap current profiles from the DKES code [5]. Plasma currents up to I_p = ±30 kA have been applied. They contribute Δτ_a = (µ_0 RI_p)/(2πα^2 B_o) to the boundary transform, i.e. 0.007/kA for the typical parameters α = 0.15 m, R = 2 m and B_o = 2.5 T of this study (positive currents increase the rotational transform). The NEMEC calculations show that for fixed limiter aperture the effective plasma radius changes with plasma current as Δa/a = 0.004I_p/kA. Ohmic heating is negligible (< 10 kW) as compared to ECRH (450 kW at 140 GHz). The pulse length of 1.1 s is above the current diffusion time scale.

Figure 1 shows the dependence of energy confinement on both τ_a and I_p. The plasma energy has been derived from the diamagnetic loop signal, with a correction accounting for residual magnetic flux from large plasma currents, which is determined from the kinetic electron energy content at τ_a = 0.51 (from profile integration). The electrons dominate the total energy: ion temperatures are below 450 eV compared
to electron temperatures up to 2 keV. Limiter aperture and line-integrated densities are identical for all discharges, i.e. the plasma radius increases from 0.15 m at \( I_p = 0 \) to 0.165 m at \( I_p = 25 \) kA and the central densities decrease from 4 to \( 3.5 \times 10^{19} \) m\(^{-3}\) (optimum confinement in W7-AS scales \( \propto d^2 n^{0.5} \) [6]). With small net current (\( I_p \leq 5 \) kA) a strong dependence of confinement on the rotational transform is observed which is smoothed by increasing the current (to +15 kA) and disappears at the highest current (+25 kA). Here, the level of optimum confinement in the current free case is reached.

2. Confinement at moderate shear

For the discharges of Fig. 1, degraded confinement is associated with a very flat electron temperature profile at larger radii (\( r > 0.05 \) m), whereas for good confinement, a strong \( T_e \)-gradient extends to the edge. The density profiles are similar for all discharges. Figure 2 shows the calculated profiles of the rotational transform for low current discharges (\( I_p = 0 \) and 5 kA) at selected \( t_* \)-values close to 1/2, where the sensitivity of confinement to \( t_* \) is strongest. The rational values \( t = n/m \) up to \( m = 30 \) are given for reference. The close vicinity of \( t = 1/2 \) is free from such resonances (see also Ref. [7]). The calculated bootstrap current reaches up to 7 kA in case of optimum confinement, but only up to 2 kA for degraded confinement. For low confinement with small bootstrap current (dashed lines in Fig. 2), shear is basically determined by the inductive current, i.e. remains low if no net current is imposed. For good confinement (full lines in Fig. 2), both low or moderate shear may result from the combination of inductive and bootstrap current density profiles.

The analysis gives evidence that confinement in W7-AS essentially depends on magnetic shear and resonant \( t_* \)-values in the outer plasma region: for optimum confinement \( t(r) \) has to be in the "resonance-free" region (e.g. \( t_* = 0.48 \) at \( I_p = 5 \) kA), or shear has to be sufficiently large in the presence of high order \( n/m \)-resonances of the rotational transform (e.g. \( t_* = 0.495 \) at \( I_p = 0 \)). In contrast, degraded confinement is associated with low shear in the presence of such resonances (e.g. \( t_* = 0.48 \) at \( I_p = 0 \)). With respect to the maximum \( m \)-number which
determines the width of the "resonance-free" region, the experimental $t_a$-values for the onset of confinement degradation below ($t_a = 0.49$) and above ($t_a = 0.53$) $t_a = 1/2$ are not consistent. However, since at present we do not know which specific perturbations at the high order rational surfaces enhance transport, there is no a priori reason for such a symmetry. Furthermore, the rather strong natural 5/9 perturbation will be important as well for $t_a \geq 0.53$. It gives rise to magnetic boundary islands which cannot be accounted for by the NEMEC code.

3. Proceeding to high shear

If, for $t_a = 0.51$, the net current is increased from 5 to 7.5 kA, optimum confinement strongly degrades, probably because $f(r)$ is raised with low shear into the upper "resonance" region (Fig.3; here, the variation of plasma radius with current is compensated by proper adjustment of the limiter aperture). Confinement recovers, when shear is increased by higher currents. For $t_a = 0.42$, confinement improves continuously with currents of both sign. In the low density case included in Fig. 3 from a previous study, the stagnation with high negative currents is attributed to the approach of very small central $t$-values [8].

It is obvious, that the effect of magnetic shear on local transport can at best be assessed from the degraded situation at $t_a = 0.42$, which is free from effects related to the "resonance-free" zones. Fig. 4 shows the radial profiles of electron temperature and heat conductivity (from power balance analysis) for discharges with plasma currents of 0, ±10, and +25 kA. With increasing shear a continuous steepening of the $T_e$-gradient is observed at the boundary independent on the sign of shear. In this region (0.08 to 0.14 m) the heat conductivity decreases by a factor of up to 5. As compared to neoclassical transport, the experimental $\chi_e$ is anomalous over the whole plasma cross section for $I_p = 0$, neoclassical in the very center for ±10 kA, and neoclassical up to $r/a \leq 0.7$ for +25 kA. Only at the very plasma edge transport remains strongly anomalous. Thus, with increasing shear, the region dominated by neoclassical transport continuously expands towards the boundary due to the increase of neoclassical transport with temperature and a
simultaneous strong reduction of anomalous transport. The results are consistent with those of Ref. [8] (low density in Fig. 3, $T_e \leq 4$ keV) where $\chi_e$ decreased by a factor of up to 2. At the higher collisionality of the present study neoclassical transport is significantly smaller, and the observed improvement by reduction of anomalous transport with increasing shear is more pronounced.

Reflectometry spectra indicate significant systematic changes of plasma turbulence with increasing magnetic shear (Fig. 5). A broadening and a large red-shift in the spectra is observed. In addition, the coherent peak of the unshifted mm-wave carrier, which cannot be observed for reflection in the low-shear plasma, reappears. These features gradually strengthen with the magnitude of shear and are independent of its sign. The quantitative interpretation of reflectometry spectra is difficult, since the propagation and $k$-spectra of the turbulence and the specific antenna-plasma geometry influence the spectrum of the reflected and backscattered mm-wave. However, the observed broadening and frequency shift would be consistent with a decrease of the turbulence scale length and an increasing plasma rotation in the electron diamagnetic drift direction, respectively. Unfortunately, direct spectroscopic measurements of the plasma rotation which could support this hypothesis are not available for these discharges.

4. Summary

Confinement in W7-AS at moderate $\beta$ is found to be strongly related to high order rational values of the rotational transform and to the magnetic shear. For optimum confinement, these resonances have to be avoided in the boundary region or shear must be sufficiently large if they are present. Therefore, with moderate plasma current (up to the order of the bootstrap current) the boundary rotational transform has to be adjusted to the narrow "resonance-free" region close to the major resonances 1/2 and 1/3. With $t_i$ in the "resonance" region, confinement degrades strongly since shear remains low. So far it is not clear which specific perturbations at the high order rational surfaces enhance transport and are stabilized or decorrelated by the magnetic shear. Increasing the shear with higher plasma current for a degraded situation
continuously improves confinement back towards the optimum. Transport analysis indicates a strong reduction of transport to the neoclassical level. The effect is independent of the sign of shear.

References

![Graph](image)

Fig. 1: Dependence of plasma energy on the boundary rotational transform for various plasma currents. For three discharges at $I_p = 5$ kA density control was lost and the measured energy (crossed squares) has been scaled by $n^{12}$ (squares)
Fig. 2: Calculated $\omega$-profiles for discharges with net plasma current as indicated (full lines: optimum confinement; dashed lines: degraded confinement). Rationals $n/m, m \leq 30$, are given for reference (horizontal lines; the shaded area covers the range of rationals with $m \leq 15$).

Fig. 3: Dependence of the electron kinetic energy content on plasma current for $t_a = 0.51$ and $t_a = 0.42$. 
Fig. 4: Measured $T_e$-profiles and $\chi_e$-profiles from power balance for discharges at $t_o = 0.42$ for various plasma currents.

Fig. 5: Reflectometry spectra from discharges with plasma currents of 5 and 25 kA at $t_o = 0.42$. The spectrum from wall reflection without plasma is given for reference.
IMPURITY TRANSPORT INVESTIGATIONS
AT W7-AS

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Abstract

In order to study the density dependence of impurity confinement in
ECF-heated W7-AS plasmas, impurity transport coefficients were
derived from SX-camera radiation during aluminum injection by laser
blow-off at two different electron densities ($n_{eo}=3.5\times10^{19}$ m$^{-3}$). At
high density, the diffusion coefficient turned out to be smaller by a
factor of 2-3. This result is qualitatively compatible with the temporal
increase of intrinsic impurity radiation during high density discharges
due to reduced transport. The helium concentration did not show a
temporal evolution in both cases. No striking differences in MHD-
activity, radial electric field or density fluctuations were observed.

Keywords:
stellarator, W7-AS, impurity transport, laser blow-off

1. Introduction

For the decay time of injected tracer impurities (aluminum laser
blow-off) in W7-AS, the electron density is, among others, an
important scaling parameter [1], indicating improved confinement of
impurities towards high electron density. This observation could not
simply be attributed to a decrease in $Z_{eff}$. Similar dependence of
impurity confinement within a certain density range was already observed in ECF heated Heliotron E plasmas [2], where the change in transport was supposed to be caused by changes in the diffusion coefficient $D$ rather than in the flow velocity $v$, the latter being close to classical expectations. Previous impurity injection experiments at W7-AS (H$_2$S gas oscillation and Al laser blow-off) [1] at medium density ($n_{eo}=2.5\cdot10^{19}$ m$^{-3}$) could partially be simulated within the errors by the one-dimensional radiation and transport code SITAR [3], based on neoclassical and Pfirsch-Schlüter transport for axisymmetric devices. At low density, the transport was found to be significantly higher than predicted, whereas at high density ($n_{eo}=6.5\cdot10^{19}$ m$^{-3}$), the neoclassical fluxes had to be reduced to fit the experimental data. In order to elucidate this density dependence, discharges with densities varied by a factor of 2 ($n_{eo}=3.5/7\cdot10^{19}$ m$^{-3}$) were analyzed in more detail, together with fluctuation- and MHD-diagnostics and measurement of the radial electric field. Especially in non-axisymmetric devices like stellarators, the latter can play an important role for impurity transport [4], but is not yet included in SITAR.

Simulations with simple assumptions about $D(r)$- and $v(r)$-profiles might mask possible local changes in transport. Therefore, the radial transport coefficients were tried to be directly derived from the temporal and radial behaviour of integrated spectral radiation, detected by the SX-camera during the penetration process of injected aluminum by laser blow-off. Indications about the radial location of the transport modification at different densities were expected.

2. Transport analysis

Because of its good radial and temporal resolution, the SX-camera is a proper diagnostic tool at W7-AS for transport investigation. In the case of tracer injection, the pure contribution of tracer radiation was obtained by subtraction of a discharge without injection. In spite of its energy-integrated information, the use of a 25μm Be-filter in front of the camera offers the possibility to restrict the number of ionization states contributing to the measured intensity and simplifies the reconstruction of the total tracer density profiles. For total aluminum
density reconstruction during the penetration process of injected aluminum, the radial intensity profile at each time-step was abel-inverted (Fig.1b: injection time t=0.5s) and converted to a total aluminum density profile assuming coronal equilibrium (quasi-stationary condition) in a first step. This is considered to be applicable in low-transport discharges at high electron density. The transport coefficients at a certain radial position can then be derived from the local temporal evolution of the total aluminum density profiles n(r,t) [5]. With the ansatz Γ=-D grad(n)+vn, D and v can be determined by fitting a straight line, when plotting the normalized total aluminum fluxes Γ/n vs. the normalized total aluminum density gradients grad(n)/n for all times at this radial position (Fig.1d). The flux Γ can be estimated from the density profile evolution using the continuity equation dn/dt=-div(Γ) with restriction to radial regions where external sources and sinks can be neglected.

In cases where the assumption of quasi-stationarity might not hold, the reconstruction with coronal equilibrium can cause errors. Therefore, the radial density profiles for each time step were reconstructed again, now using a reconstruction factor β(r,t)=n(r,t)/P(r,t) (P: total local emissivity contributing to SX-camera) obtained from a transport and radiation calculation with SITAR, using the D- and v-values derived in the first step as input. A repeated derivation procedure for D and v, but now with the corrected density profile evolution as described above, provides new transport coefficients (Fig.2) which, in fact, fit the SX-camera better in most cases (Fig.1b,c). The accuracy of this iterative method strongly relies on the quality of abel-inversion and atomic data base. Bremsstrahlung and radiation from dielectronic recombination and inner-shell processes were not yet implemented in SITAR, but were checked to contribute not more than 20% to the SX-radiation in the case of stationary conditions. Possible errors have to be discussed in this context.

3. Results and Discussion

In order to study the density dependence of impurity confinement at W7-AS, ECF-heated discharges (P_{ECF}=480kW, t=0.338, B=2.56T,
a=0.173m) at two different densities \((n_{eo}=3.5/7×10^{19} \text{m}^{-3})\) were compared. The transport coefficients were derived according to the procedure described above (Fig.1) and are plotted in Fig.2 for both densities.

For the outer plasma region, the last reliably determined value of the diffusion coefficients at \(r=0.10-0.11\text{m}\) were kept constant up to the plasma edge. They represent average diffusion coefficients for this region, determining essentially the central time behaviour during the inflow phase. The extrapolated average diffusion coefficient for medium density approaches the neoclassically predicted one (Fig.2: light shaded area), but falls below the predictions (Fig.2: dark shaded area) for the high density case (fluxes in SITAR are already reduced by 50% to account for the W7-AS transport optimization; 1.7% C as additional impurity background concentration was assumed in both cases, compatible with \(Z_{eff}\) measurements). This trend is even more pronounced at higher electron densities of \(n_{eo}=1.2×10^{20} \text{m}^{-3}\), e.g. in high confinement neutral-beam heated discharges [6], where transport coefficients were derived (\(D(r)=0.07 \text{m}^2/\text{s}\), \(v(r)=5 \text{m/s} \times (r/a), \#38551\)), being clearly smaller than predicted by SITAR.

The convection velocity was extrapolated to vanish in the plasma center and had to be adjusted in the outer part to fit better the temporal decay of spectral line intensities from different ionization states of aluminum, observed by central-line-of-sight crystal- and VUV-spectrometers (Fig.3: left). For the high density discharge, only a slight correction in the derived \(v\) (Fig.2: see arrows) was necessary to excellently fit all experimental data. In the case of medium density, \(D(r=0.10\text{m})\) was used for extrapolation and \(v\), however, has to be reduced by a factor of 2 for a good compromise in fitting the experimental data radially as well as temporally. A reduction of \(D\) in the outer region down to the value \(D(r=0.11\text{m})\) would fit the time traces of the spectrometers only for the case of vanishing convection velocity, but will lead to a misfit of the radial profiles.

In the two ECF-heated discharges under investigation, no clear changes as well in the radial electric field \(E_r\) (within the errors) as in MHD-activity, density fluctuation level and radial density and
temperature profile shape (Fig.4) could be observed from which the difference in transport can be deduced reliably.

In the case of high electron density, D is overall lower by a factor of 2-3 compared to the medium density discharge, v being quite similar. The resulting difference in confinement can be illustrated quite impressively by the change in decay time for the injected aluminium (Fig.3: left). Consequently, also the time evolution of intrinsic impurity radiation and \( Z_{\text{eff}} \) is remarkably different during the flat-top phase for this two densities (Fig.3: right). However, assuming a constant impurity influx of, e.g. intrinsic chlorine from the walls and using just the derived set of transport parameters, the difference in the time traces of CI-XIV can be qualitatively well described (Fig.3: see fits): compared to medium electron density, where stationary conditions were achieved well within the pulse length, the reduced transport in the high density case causes longer times to establish stationarity. Similar behaviour for other intrinsic impurity species might explain the signals of SX-camera, bolometer and \( Z_{\text{eff}} \) (Fig.3 right, Fig.1a). Nevertheless, at high density \( n_{\text{eO}}=7\cdot10^{19}\text{m}^{-3} \) the extrapolated total radiation (bolometer) at the time when the CI-XIV radiation should reach 90% of its stationary level (from simulation: approx. at 1.7s), stays well below the critical value of 60% of the heating power [7], where the plasma is severely affected by radiation. The helium inventory (introduced by a gas-puff and kept in the machine by recycling) as measured by CXRS (He-II at \( r=6\text{cm} \)) shows no temporal increase in the high density case.

References


Figure Captions

Fig. 1: Evolution of radiation at medium and high density: (a) total radiation, (b) SX-camera radiation during impurity injection, (c) simulation of SX-radiation with derived D,ν, (d) graphical derivation of D,ν from aluminium density profile evolution.

Fig. 2: Derived transport coefficients D,ν at medium (thin curve) and high (thick curve) electron density (dotted lines: extrapolations, arrows: corrections, shaded areas: predictions by SITAR).

Fig. 3: Left: time traces and simulations of spectral lines from different ionization states of aluminum after injection by laser blow-off: (1) n_{e0}=7·10^{19} m^{-3}, (2) n_{e0}=3.5·10^{19} m^{-3}. Right: time evolution of impurity radiation and Z_{eff}(SX), and simulations for Cl-XIV.

Fig. 4: Electron density, temperature and ion temperature profiles of the investigated ECF-heated discharges: (●) medium, (○) high density.
IMPURITY INVESTIGATIONS BY MEANS OF LI-BEAM INDUCED CHARGE EXCHANGE SPECTROSCOPY ON W7-AS

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Abstract
Knowledge of the impurity concentration and temperature in the core plasma gradient region as well as the SOL is a vital element on the road to documenting and understanding the physics of L- or H-mode transport and the transport barrier itself. To this end, the Li-beam diagnostic capabilities on W7-AS have been expanded to include the measurement of radial profiles of ion impurity density and temperature via charge-exchange- spectroscopy (Li-CXS).

After briefly discussing the method of Li-CXS this paper describes the experimental setup on W7-AS and presents first results. Measurements of C\textsuperscript{6+}, C\textsuperscript{5+} and B\textsuperscript{5+} spectral lines prove the viability of Li-CXS. In the plasma edge region (r\textsubscript{eff} > 8 cm) a C\textsuperscript{6+} concentration of about 0.5\% could be measured. Temperature values found for C\textsuperscript{6+} are similar to proton/deuterium temperatures.

The intensities of several Li II spectral lines (2p\rightarrow2s, 3d\rightarrow2p, 4s\rightarrow2p, 4d\rightarrow2p) have been measured and are used to critically check the underlying cross section database employed within the collisional-excitation Li-beam model, especially for collision processes involving higher Li-states (n\geq3). It was found that the ratio of the 3d\rightarrow2p and 2p\rightarrow2s spectral lines is overestimated by the model but is within uncertainties in the cross section database.

Keywords: diagnostics, atomic beam, Lithium, charge exchange spectroscopy, impurity density, electron density, impurity temperature
Introduction
Li-beam diagnostics are a multi-purpose technique for investigating fusion edge plasmas. While the determination of electron densities by lithium impact excitation spectroscopy (Li-IXS) has already reached a satisfactory standard on both large fusion experiments at IPP Garching [1,2], neutral Li-beams can also be used to determine local concentrations as well as temperatures of impurity ions via charge exchange spectroscopy (Li-CXS) [3,4]. This method has been proposed by Winter [5] and was applied for the first time at the TEXTOR tokamak at KFA Jülich [6,7]. In order to achieve simultaneous Li-IXS and Li-CXS measurements, the present setup for electron density measurements has been extended. First results prove the feasibility of Li-CXS with the improved Li-injector [1] in W7-AS plasmas.

Principles
From observation of the resonant line radiation profile from the injected Li-beam it is possible to deduce the electron density profile as well as the local Li(nl) state distribution along the injected Li-beam by modeling the Li-beam-plasma interaction. In addition to the impact excitation process the weakly bound outer electron of the Li-atom can also be captured by impurity ions. This charge exchange process populates highly excited states of the impurity ions, giving subsequently rise to characteristic impurity line radiation. Observation of this line emission in conjunction with the calculated Li(nl) state distribution allows evaluation of the impurity density profile along the injected Li-beam. Finally the temperature profile of the impurity ions is determined from the spectral line shape.

Experimental setup on W7-AS
The existing Li-beam diagnostic layout [1] has been supplemented by a 14 channel observation system with a radial resolution of $\delta r \approx 6\text{mm}$ for a range of $\sim 13\text{cm}$ along the beam, corresponding to an effective plasma radius from 3 to 17cm, cf. fig.1. Two glass lenses ($\Omega/4\pi \sim 2.9\times10^{-4}\text{ sr}$) image the light onto 14 bundles, each consisting of a $2\times4$ array of
400μm quartz fibers. The bundles are coupled one by one to the entrance slit of a monochromator (ACTON, Czerny-Turner, f=0.75m) to permit spectral resolution for Li-CXS and LiII radiation. A two-dimensional detector (Proscan CCD camera, 512x512 pixels, each 19x19μm²) is directly connected to the monochromator exit. The spectral resolution can reach up to 0.018nm/pixel, using a 1800g/mm holographic grating (blazing for 500nm). An additional system of R928 or R3896 Hamamatsu photomultipliers in conjunction with interference filters (λ=529.0nm for C6+, δλ1/2 ~ 5Å) can be coupled to the same light guides for simultaneous measurements at the 14 radial locations.

The extraction geometry of the Li-beam injector was changed to increase the Li-beam current delivered by the gun [1]. These experiments were carried out at 50keV injection energy with an equivalent neutral Li-beam current of 1.9mA and a Li-beam diameter of D_{FWHM}≤1cm.

**Lithium beam composition**

Since highly different cross sections for charge exchange processes follow from different excitation states of the donor atom, the composition of the Li-beam is of great importance for evaluating CXS data. We have therefore investigated several LiII spectral lines (2p→2s, 3d→2p, 4s→2p, 4d→2p) in W7-AS discharges. While the measurements of the most relevant line Li(2p→2s) at λ=670.8 nm were performed to calibrate the CXS setup relative to the Li-IXS photomultiplier setup (see below), all other LiII lines were investigated to check the attenuation model of the Li-beam [8]. Measured intensities of emission from higher levels were found to differ considerably (30-60%) from corresponding theoretical values. We observed no dependence on magnetic field strength and beam energy (20-66keV). The plasma density had a strong influence on the conformity of experimental and theoretical values, with the deviation becoming larger at higher densities.

As the major reason for these disagreements, inadequate scaling relations for excitation and ionization processes involving protons and impurity ions in the underlying database [9] have been identified. These
are now being recalculated by more advanced means. However, since
the relative population of the Li(3d) level in the Li-beam is in the range
of 1% only, and populations of all other Li(nl) levels (n>2) are even
smaller, the influences of cross section discrepancies on electron
density calculations remain below 10%.
Furthermore, recent simulations have suggested that the population of
higher excited states depends on Z_{eff}. Thus, the measurement of only
one additional line besides the resonance line offers a possibility to
determine an estimate for Z_{eff} under the assumptions that a reasonable
radial charge state distribution of the impurities is given and that the
present disagreement for n=3 populations between simulation and
experiment diminishes as a consequence of more accurate cross
sections.

CXS investigations
Impurity density profiles
To determine the absolute concentration of C^{6+} impurity ions, both Li-
beam diagnostic systems are necessary (IXS and CXS). While the IXS
system (28 PM) records the Li(2p) light emission, the CXS system
(2CCD-camera) delivers the CVI line radiation at 529.0nm. Since both
systems use different observation optics the detection efficiency has to
be cross-calibrated. This is done by measuring the Li(2p) light with
both systems, the IXS-PM- (U^{PM}_{2p}(i)) and the CXS-CCD-system
(U^{CCD}_{2p}(i)), respectively in a calibration discharge.

\[ U^{PM}_{2p}(i) = k^{PM}_{671}(i) \cdot n_{Li} \cdot A_{2p \rightarrow 2s} \cdot N(2p) \]  (1a/b)

\[ U^{CCD}_{2p}(i) = k^{CCD}_{671}(i) \cdot n_{Li} \cdot A_{2p \rightarrow 2s} \cdot N(2p) \]

k denotes the detection efficiency for the two systems at \( \lambda = 671.0 \text{nm} \), i the corresponding radial channel, \( A_{2p \rightarrow 2s} \) the transmission probability, \( n_{Li} \) the particle density of the Li-beam and N(2p) the relative occupation
number of the 2p-state. The CXS-signal \( (U^{CCD}_{529}(i)) \) can be expressed by

\[ U^{CCD}_{529}(i) = k^{CCD}_{529}(i) \cdot v_{Li} \cdot n_{C^{6+}} \cdot n_{Li} \cdot \sum_{n,l} \sigma_{529}(nl) \cdot N(nl) \]  (2)

Calculating the ratio of equ.1b and 2 for the discharges in question and
using the ratio of equ. 1a and 1b for the calibration discharge, the C^{6+}
density can be expressed by

\[ n_{6^+}(z_{Li}(i)) = \frac{U_{529}^{(i)} \cdot k_{721}^{PM}(i) \cdot k_{821}^{CCD}(i) \cdot A_{2p+2s} \cdot N(2p)}{U_{2p}^{(i)} \cdot k_{671}^{CCD}(i) \cdot k_{529}^{CCD}(i) \cdot v_{Li} \cdot \sum_{n} \sigma_{529}(nl) \cdot N(nl)} \]  \( (3) \)

In equ. 3 the first ratio describes the measured signal ratio (IXS and CXS systems), the second ratio is determined by the calibration procedure (see above) and the third ratio expresses the different detection probabilities of the CCD camera for the two wavelengths. \( v_{Li} \) denotes the Li-beam particle velocity and \( \sigma_{529}(nl) \) the cross section for charge exchange from the (nl)-level, giving rise to subsequent line radiation at \( \lambda = 529.0 \text{nm} \). The relative occupation numbers N(nl) of the Li-beam atoms are calculated in the reconstruction process for the electron density (Li-IXS). The described algorithm was applied in a series of equivalent discharges \( (P_{\text{ECHI}}=400 \text{kW}, D_2, t=0.34, B_z=20.0 \text{mT}, \text{up/down limiters attached}, I_{n_c} dl(\text{HCN})=2 \cdot 10^{19} \text{ m}^{-2} \) to determine the C6+ impurity ion density for different radial positions (different \( z_{Li} \) in equ. 3). The result is shown in fig. 2. The C6+ impurity ion concentration increases from about 0.41% at \( r_{eff}=16.7 \text{cm} \) to 0.63% at \( r_{eff}=8.3 \text{cm} \). In the gradient region of the probe good agreement is found with the result from the H-CXS diagnostic [10].

For a single light guide in the radial range of the maximum of the Li(2p) profile we also investigated the signal to background ratio for other impurity ions. While the ratio was about 0.25 for C6+ we found a ratio of 0.2 for C5+ and B5+ respectively, clearly demonstrating the applicability of Li-CXS also for these impurity ions. The exposure time of the CCD camera was about 5.7ms with a readout time of 3ms. Each 10ms a CCD-picture was recorded. The Li-beam was chopped electronically with a beam on and off time of 40ms each. Thus 4 CCD-pictures could be used to determine Li-CXS and background signals.

For the determination of impurity densities the Li-CXS signals were summed over the Li-beam on-time interval (40ms). No impurity ion concentrations giving rise to CXS-line radiation in the UV-spectral range could be investigated, due to insufficient transmission of the observation optic (glass).
Impurity temperature profiles

Temperature values are obtained by fitting a Gaussian profile to the CCD-camera data. To reduce the scatter of the calculated spectral line width, all CCD-pictures taken in the flat top phase of one discharge had to be summed. This typically implies a integration time of about 300-400ms. The time resolution will be improved by a better coupling of the light guide bundles to the spectrometer entrance slit, where now an important fraction of the light signal is being lost.

For the fitting procedure line broadening effects such as Zeeman splitting and l-level mixing are taken into account [7]. C^{6+} temperature values in the plasma edge obtained in the same series of discharges (see above) are shown in fig. 3. For $r_{\text{eff}} < 12\text{cm}$ C^{6+} temperature values are similar to proton/deuteron temperatures measured via active neutral particle analysis [11].

Conclusions

With the improved Li-injector now available we have demonstrated the applicability of Li-CXS, both for impurity density and impurity temperature measurements, on the W7-AS stellarator. Concentrations found for C^{6+} ions are in the order of 0.5%. The radial profile in the gradient region coincides with values from the H^{a}-CXS-diagnostic. Temperature values in the plasma edge can be considered to be equal to those for deuterium ions. Radial and temporal resolutions of 0.5cm and 40ms, respectively (400ms for temperature measurements) could be achieved. By improving the detection efficiency, mainly by new fibers with smaller aperture and a new construction of the fibers-spectrometer coupling, the temporal resolution can probably be increased by more than a factor of 5.

Acknowledgments

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References


Figure Captions

Fig. 1: Experimental set-up on the W7-AS stellarator. The observation geometry for Li-IXS (28 channels), Li-CXS (14 channels) and neutral density $N_e$ is indicated.

Fig. 2: Radial impurity density profile of $C^6+$ and electron density profiles as a function of effective radius $r_{ef}$ for discharges #39360-39368. The last closed magnetic surface (LCMS) is indicated by a dotted line.

Fig. 3: Radial impurity temperature profiles of $C^6+$, electron temperature $T_e$ and deuterium temperature $T_D$, measured by neutral particle analysis (NPA).
Fig. 1: Experimental set-up on the W7-AS stellarator. The observation geometry for Li-IXS (28 channels), Li-CXS (14 channels) and neutral density $N_o$ is indicated.
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Fig. 3: Radial impurity temperature profiles of $\text{C}^{6+}$, electron temperature $T_e$ and deuterium temperature $T_D$, measured by neutral particle analysis (NPA).
Density Limit Study on the W7-AS Stellarator


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Abstract

Data from currentless NBI discharges in W7-AS strongly indicate that the maximum density for quasi-stationary operation is limited by detachment from limiters. The threshold density at the edge scales with $P^0.5 B^{0.8}$ (with $P$ being the net power flow across the LCMS) which is consistent with an edge based analytic estimation presuming constant threshold downstream temperatures.

Keywords:
stellarator, W7-AS, plasma edge, density limit, limiter

1. Introduction

In currentless stellarator discharges, the maximum achievable plasma density is basically limited by an impurity radiation induced thermal instability (density limit, DL) [1, 2] which, however, does not lead to MHD instability with subsequent disruption as is typical for tokamak DL discharges [3-5]. The sequence of events leading to thermal instability is governed by the same basic physics as in tokamaks and should exhibit some similarity, at least in limiter machines. This paper continues previous DL studies on W7-AS [1, 2] and concentrates on the starting event of the DL sequence (onset condition) which can be described in terms of edge quantities alone, and which practically
determines the actual operational window for quasi-stationary limiter discharges. The aim is to give a status report rather than a comprehensive description.

2. Experimental

The analysis was made for net current-free NBI discharges in W7-AS \( (R = 2 \text{ m}, \ a = 0.18 \text{ m}) \) at \( B = 1.25 \text{ T} \) and \( t = 0.34 \text{ T} \) and \( t = 0.34 \text{ T} \). The configuration was bounded by two tangential graphite limiters at the top and bottom of an elliptical cross section. The absorbed NBI power was varied between 0.35 and 1.2 MW. The power deposited onto the limiters (from limiter-integrated thermocouples) in discharges with flat-top density below the DL was typically about 80% of the non-radiated power fraction. For the DL study, the density was slowly ramped up (typically for about 0.4 s) until the discharges collapsed. Particle confinement times are estimated to be less than 50 ms.

3. Results and discussion

Core parameters, radiation. With the line-averaged density \( \bar{n}_e \) ramped up, the stored energy slightly increases up to a more or less pronounced maximum and subsequently rolls over, at first along a time scale of some ten milliseconds and then much faster (see example in Fig. 1). This fast decay coincides with a rapid decrease of the density, even though the feedback controlled gas valves fully open in this phase. At \( B = 2.5 \text{ T} \), the total radiated power \( P_r \) (from bolometer cameras) is found to increase approximately as \( P_r \propto \bar{n}_e^2 (Z_{\text{eff}}^{-1}) \) (\( Z_{\text{eff}} \) from bremsstrahlung) up to a critical density \( \bar{n}_e^{\text{crit}} \). Above this density the increase becomes much steeper indicating radiative instability. Discharges at \( B = 1.25 \text{ T} \), where \( Z_{\text{eff}} \) values were not available, show in general a quite similar evolution of \( P_r \) with time. Abel-inverted radiation profiles (from 12-channel bolometer camera) are typically hollow below \( \bar{n}_e^{\text{crit}} \), then increase first at the edge, and become strongly peaked during the collapsing phase. Data from bolometer cameras at three different toroidal positions do not indicate significant radiation asymmetries within the last closed magnetic surface (LCMS); MARFE formation was not observed. Whether the phase after having attained \( \bar{n}_e^{\text{crit}} \) leads to
a complete collapse following the example in Fig.1, seems to depend primarily on the external gas programme. In contrast to feedback control, with constant or slowly increasing gas feed the discharges could often be further sustained at a degraded energy level or even oscillated between low and high energy contents. Within the low energy phases, \( n_e, T_e \) (from Thomson scattering) and \( P_r \) radial profiles indicate a radial shrinking of the plasma column.

**Scrape-off layer (SOL) parameters.** The upstream densities \( n_{es} \) at the LCMS (from Li-beam) increase \( \propto \bar{n}_e^{1.140.1} \) at \( B = 1.25 \) T and \( \propto \bar{n}_e^{1.590.1} \) at \( B = 2.5 \) T up to a maximum which coincides with the stored energy maximum, and then decrease at higher \( \bar{n}_e \) (see example in Fig. 2). LCMS upstream temperatures \( T_{es} \) are typically above 60 eV up to \( n_{es}^{max} \) and then drop rapidly. The high \( T_{es} \) values are consistent with the small wetted limiter areas, short power flux decay lengths \( \lambda_q \) (\( \leq 1 \) cm), and long field line connection lengths \( L_e \) (\( \approx 30 \) m) in W7-AS. Downstream strike-point temperatures \( T_{ed} \) (from limiter-integrated Langmuir probes) decrease from slightly below \( T_{es} \) at low density to a more or less pronounced pedestal of about 20-25 eV at about \( n_{es}^{max} \) and then drop further. The downstream ion saturation currents \( I_{sat} \) and densities \( n_{ed} \) derived from it increase with \( n_{es} \) up to a certain density \( n_{es}^{det} < n_{es}^{max} \) and then fall typically by factors of about 1.5 to 5. This indicates reductions of the thermal plasma pressure along field lines by more than a factor of eight and detachment from limiters. Detached phases with high stored energy as shown in Fig.1 do not exhibit \( n_e, T_e \) and \( P_r \) radial profile shrinking. This is in contrast to limiter tokamak DL scenarios which alternatively show detachment by radial shrinking or MARFE formation. The difference could be due to the radial flux expansion by about a factor of two at the limiter positions in W7-AS.

The consistency of the measured SOL parameters with basic SOL physics is checked by a simple two-point estimation including the power balance along field lines, momentum balance and the sheath boundary condition [5, 6]:

\[
T_{es} = \left( T_{ed}^{2/7} + \frac{7}{4} \frac{q_{ls} L_e}{\kappa_0} \right)^{2/7} \quad \text{with} \quad q_{ls} = \frac{P_s}{\lambda_q w} \quad (1)
\]
\[ 2n_{ed} T_{ed} = (1-f_m)n_{es} T_{es} \quad (2) \]

\[ n_{ed} \sqrt{\frac{2k}{m_i}} T_{ed}^{1/2} (\xi + \gamma T_{ed}) = (1-f_r)q_{ls} \quad (3). \]

\( P_s, q_{ls}, w, f_r, f_m \) and \( \xi \) are the net power crossing the LCMS, the upstream parallel power flux density, the wetted limiter width, the SOL energy and momentum loss fractions and the potential energy transferred per electron-ion pair (\( \approx 25 \) eV), respectively. Further symbols have the usual meaning [5].

To visualize the role of momentum losses, the measured \( P_s \) and \( T_{ed} \) values (beside geometrical parameters and \( \lambda_q \)) and \( f_r = f_m = 0 \) are used as input in a first step. Then a high recycling scenario with \( n_{ed} > n_{es} \) is predicted which clearly was not observed (see the discrepancy between calculated and measured \( n_{ed} \) values in Fig. 2). Taking \( P_s, n_{es} \) and \( I_{sat} = en_{ed} c A_p \) (with \( A_p \) being the probe area) as input and varying the remaining free parameter \( f_r \) for \( f_m \) to match the \( f_m(T_{ed}) \) values for CX momentum losses from Fig. 24 in Ref. [7] results in a detached scenario consistent with the measured data with the exception of \( T_{ed} \). The latter should be \( < 5 \) eV near the \( I_{sat} \) minimum, see Fig. 3. The example further indicates that both strong momentum and volumetric energy losses from the SOL are required for consistency within this estimate. Though \( f_r \) in the SOL could not yet be measured, this explanation seems to be somewhat realistic because the measured \( n_{es} \) and \( I_{sat} \) values are much less prone to error than the \( T_{ed} \) probe values. In the presence of strong \( T_e \) parallel gradients, the latter are very likely overestimated due to contributions from fast tail electrons from regions further upstream [8]. On the other hand, we cannot exclude that momentum losses due to radial transport [9] contribute to detachment in this geometry. The further analysis is not affected by this uncertainty.

**Operational limit scaling.** Detached scenarios could not yet be quasi-stationarily sustained in W7-AS. We therefore focus on the upstream limit density \( n_{es, det} \) for attached discharges which presently determines the practical operational limit for quasi-stationary limiter discharges. From equations (1)-(3) we get
\[ n_{\text{es}} = \left( \frac{4K_0}{7} \right)^{2/7} \frac{(1-f_g)}{(1-f_m)} \frac{q_{\text{LS}}^{5/7}}{L_{\text{c}}^{2/7}} F(T_{\text{ed}}); \quad F(T_{\text{ed}}) = \frac{2T_{\text{ed}}^{1/2}}{\sqrt{\frac{2k}{m_i} k (\xi + n_{\text{ed}})}} \]  

(4)

Considering carbon (from limiters) as dominating SOL impurity, it can be assumed that detachment starts at a threshold \( T_{\text{ed}} \) value between about five and ten eV where both radiative and momentum losses from the SOL become efficient. The exact value is not critical because \( F(T_{\text{ed}}) \) is only a weak function of \( T_{\text{ed}} \) in this range and can be taken as constant. The above described calculation to reproduce the measured \( n_{\text{es}} \) and \( I_{\text{sat}} \) values for a detached state yields \( f_r, f_m \ll 1 \) and \( (1-f_g)/(1-f_m) = \text{const.} \) when applied to the attached phases. With this approach we get the onset condition

\[ n_{\text{es}}^{\text{det}} \propto \frac{q_{\text{LS}}^{5/7}}{L_{\text{c}}^{2/7}} (\lambda_{qw})^{5/7} = \frac{P_{qw}^{5/7}}{L_{\text{c}}^{2/7}} \]  

or, with \( (\lambda_{qw}) \propto P_{qw}^{1/7}(n_{\text{es}}^{\text{det}}B^2) \) (see below),

\[ n_{\text{es}}^{\text{det}} \propto \frac{P_{qw}^{1/7}B^2}{L_{\text{c}}^{5}} , \quad \alpha = \frac{5/7(1-x)}{1-5/7y} , \quad \beta = \frac{5/7z}{1-5/7y} , \quad \delta = \frac{2/7}{1-5/7y} \]  

(6)

\( \lambda_{qw} \) values derived from SOL \( n_e \) and \( T_{\text{c}} \) radial profiles (from fast reciprocating Langmuir probes) are found to scale as \( P_{qw}^{0.45} n_{\text{es}}^{0.45} B^{-0.5} \) (with the \( B \) dependence inferred from profiles at only two different \( B \) values, 1.25 and 2.5 T). With tangential limiters, the wetted limiter width slightly increases with \( \lambda_{qw} \), \( w \propto \lambda_{qw}^{0.4} \). Inserting these values into eq. (6) yields \( \alpha = 0.53\pm0.12 \), \( \beta = 0.83 \) and \( \delta = 0.48\pm0.05 \) which is consistent with the experimental findings shown in Figs. 4 and 5 \( (L_{\text{c}} \) was not varied). Nevertheless, the scaling with \( B \) needs further confirmation as high power discharges at low magnetic field in general were less close to quasi-stationarity (deteriorated density control) than those at high field. With respect to the dependence of the limit density on the absorbed NBI heating power we found \( n_{\text{es}}^{\text{det}} \propto P_{qw}^{0.5} B^{0.8} \) This, together with the similar scaling with \( P_{qw} \), indicates that the limit should occur at an approximately constant core radiated power fraction which was actually observed \( (f_{r,c}^{\text{core}} = 0.55\pm0.1) \). It has to be noted that any long-term effects, for example possible impurity accumulation, are not considered in this study.
4. Summary and conclusions

Data from currentless NBI discharges in W7-AS show that the maximum achievable density in quasi-stationary discharges is actually limited by detachment from limiters indicated by a drop of the downstream $I_{sat}$. Corresponding threshold densities at the edge scale as $n_{det} \propto P_s^{0.5} B^{0.8}$ which is consistent with an edge based, analytical estimation presuming constant threshold downstream temperatures. Model improvement (B2/EIRENE code) and further experiments with respect to the detached phase and possible long-term effects are planned.

References

Figure Captions

Fig. 1: Typical evolution of core parameters for the line-averaged density $\bar{n}_e$ ramped up until collapse. $W_{\text{dia}}$ is the stored energy, $T_{\text{es}}$ the electron temperature at the centre; for other symbols see text.

Fig. 2: Temporal evolution of LCMS upstream and downstream parameters. The $I_{\text{sat}}$ and $H_{\text{d}}$ line intensity drops at 0.44 s indicate detachment from limiters. A two-point estimation (see text) based on the less reliable probe $T_{\text{ed}}$ and neglecting momentum losses predicts high recycling with $n_{\text{ed}} > n_{\text{es}}$ (lines in the bottom figure) which was not observed.

Fig. 3: SOL momentum and energy loss fractions $f_{\text{m}}$ and $f_{\text{r}}$, respectively, required to explain the observed reduction of the downstream/upstream plasma pressure at 0.53 s in Fig. 2 by friction with neutrals. The free parameter $f_{\text{r}}$ was varied for $f_{\text{m}}$ to match the $f_{\text{m}}(T_{\text{ed}})$ data for CX momentum losses in Ref. [7]. The dotted lines indicate the scatter range of data in [7].

Fig. 4: Threshold densities for the onset of detachment from limiters indicated by the $I_{\text{sat}}$ drop. Errors are estimated to be about ±10% for the densities and ±30 kW for the power. $\bar{n}_{\text{es}}^{\text{max}}$ indicates maximum densities which were transiently achieved during detachment.

Fig. 5: Summary of the edge threshold densities in Fig. 4. The scaling pre-factor is valid for $P_s$ in MW and $B$ in T.
Fig. 5

- B = 2.5 T
- B = 1.25 T
Core Fluctuations and Non-Thermal Electron Distributions at W7-AS

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Abstract
Two special applications of ECE radiometry at W7-AS, for measurement of temperature fluctuations and non-thermal electron distributions, are presented. In the first case correlation-radiometry is applied. The fluctuation spectra are found to be composed of two components, a diffusive and a turbulent one, with relative levels of 0.1–2%. The latter one is discussed in the frame of a turbulent mixing process. To investigate non-thermal electron energy distributions a vertical ECE observation geometry allowing for separation of both thermal and suprathermal populations and their time behaviour is used. Experimental results by variation of plasma parameters are shown and compared with simulated spectra.

Keywords
ECE diagnostic, temperature fluctuations, turbulence, non-thermal electron distributions, suprathermal emission

1 Core Temperature Fluctuations

1.1 Characterisation
Temperature fluctuation measurements are carried out in the context of turbulent transport, which is assumed to be driven at least in parts by fluctuations [9]. The sensitivity limit set by the inherent noise of an ECE-radiometer can be overcome by correlation of two separate radiometers focusing a common plasma volume [8]. Due to the demands for stationarity of the discharge during the correlation analysis only ECR heated plasmas are investigated. Typical parameters are: 2.5 T on-axis magnetic field,
rotational transform $\kappa = 1/3$, $P_{ECRH} = 300–800\,\text{kW}$ (on-axis), peaked temperature profile ($T_{\text{max}} \sim 2.5\,\text{keV}$) and flat density profile ($n_{\text{max}} = 2 \cdot 8 \cdot 10^{19}\,\text{m}^{-3}$). Six radial positions on the high-field side of the core plasma ($r_{\text{eff}} = 5–10\,\text{cm}$, i.e. $r/a = 0.3–0.6$) are observed simultaneously. The relative fluctuation spectra extend up to $50–150\,\text{kHz}$, with radially increasing fluctuation power. Two components which reveal completely different physical behaviour are separated by a minimum around $15\,\text{kHz}$ [4].

### 1.1.1 Radial Properties

The integrated relative power of the fluctuation component below $15\,\text{kHz}$ is about $\tilde{T}/T = 2\%$ ($150\,\text{Hz–15\,kHz}$). Its radial amplitude profile and radial propagation properties (fig. 1), determined from the correlation of different radial channels, point to a diffusive propagating temperature disturbance. Its origin is located outside the observed radial range and it is propagating radially inward according to a heat diffusivity of $\chi_{\text{mc}} = 0.8\,\text{m}^2/\text{s}$. A radial decay length of $2–3\,\text{cm}$ (at $450\,\text{Hz}$) and a radial coherence length $> 20\,\text{cm}$ (i.e. infinite with respect to the radial range diagnosed) confirm its diffusive character. This is in perfect agreement with active temperature perturbation experiments where ECRH is modulated. The radial propagation of the fluctuation component around $40\,\text{kHz}$ is completely different. Its phase velocity $2\pi f/k_r$ is directed outward (fig. 1) with a dispersion relation which cannot be explained diffusively. Its coherence length ($1.5–3\,\text{cm}$) is smaller than the observed radial region, pointing to a radial extended turbulence-like feature. Comparison of high-field-side and low-field-side measurements show an about 2 times larger fluctuation level on the low-field-side whereas the profile of the diffusive component is identical on both sides, i.e. it is constant on a flux-surfaces.

### 1.1.2 Correlation with Mirnov-coil signals

The integrated relative power of the component around $40\,\text{kHz}$ which is believed to be the one relevant to turbulent driven transport, is about $0.6\%$ on the high-field side. A clear coherence of
10–60% with a Mirnov-coil signal indicates corresponding magnetic field variations which are estimated to be in the order of some T/s. However, no mode-like spatio-temporal structure can be identified [1]. Correlation of density and temperature fluctuations have been found in a combined reflectometry-ECE experiment [3] and will be discussed in detail in a forthcoming paper.

1.1.3 Vanishing Temperature Gradient

As expected in a turbulent, convective picture, the temperature fluctuation component above 15 kHz disappears completely in the flat region of the temperature profile. Measurements with reflectometry show significant density fluctuations even in the $\nabla T = 0$ region, indicating persistent plasma turbulence. According to the convective picture, the $\nabla T$-dependence is separated and the "turbulence" $\dot{\Lambda} = \dot{T}/\partial_T T$ is investigated. The fluctuation component below 15 kHz is found in the $\nabla T = 0$ region with a similar amplitude like in the finite-gradient region, as expected due to its diffusive nature.

1.1.4 Parameter Scans

The dependence of the integral turbulence $\dot{\Lambda}$ of the component above 15 kHz on parameters like ECR heating power, plasma density, rotational transform $\ell$ and plasma current has been investigated. Earlier results [5] are correspondingly found looking at $\dot{\Lambda}$. In the $\ell$-scans discharges with poorest confinement show highest fluctuation or turbulence levels. An empirical attempt to assign local plasma parameters to the $\dot{\Lambda}$-profiles points to an universal decrease of $\dot{\Lambda} \propto T^\alpha$ with increasing temperature. In the case of the ECRH-powerscan (fig. 2), an exponent $\alpha \approx -1.5$ has been found, whereas a similar analysis for the density scan yields different exponents ($\alpha = -2 \ldots -3.5$) for different densities (fig. 3).
2 Non-Thermal Electron Distribution

2.1 Introduction

During ECRH heating at high power density, distortion of the Maxwellian electron distribution function is expected near the resonance zone with the creation of a suprathermal tail. In the simplest model a bi-Maxwellian distribution can be assumed consisting of a thermal and a suprathermal part. Whereas low field side ECE observation does not allow for unique interpretation (ambiguity between position and energy), by vertical observation the thermal emission is reduced to a relatively narrow frequency range because of the nearly constant magnetic field along the line of sight. The down-shifted radiation emitted by the suprathermal electrons is much less reabsorbed along the vertical path with a vertically viewing antenna system with Gaussian beam optics. It is arranged in a toroidal plane close ($\Delta \Phi \approx 5^\circ$) to that where the ECRH is deposited. It is movable in radial direction, so it allows to localize the suprathermal electrons as the only diagnostic in the ECRH deposition plane. The non-thermal part of the emitted spectrum can be identified by comparison with numerical calculations and by investigation of the different time behaviour of the two components.

2.2 ECE spectra and Numerical Simulations

Low density plasmas heated by high power ECRH (400 kW and 800 kW, respectively, X-mode polarization) have been investigated. Fig. 4 shows a measured spectrum for the case of $P = 800$ kW, $n_{e0} = 1 \cdot 10^{19}$ m$^{-3}$ and vertical observation (circles). A ray-tracing code is used to calculate ECE spectra including a suprathermal component. The ECE-spectrum is obtained by solving the radiative transfer equation along the ray path. Due to the presence of a local maximum of B along the viewing chord (close to the plasma centre) two peaks appear in the calculated ECE-spectrum (full line in fig. 4) [2]. The maximum at higher frequencies comes from radiation emitted by a region close to the plasma centre. Decreasing the frequency the thermal resonance moves to the plasma edge where
the temperature is much lower. At frequencies below the minimum at 138-139 GHz the thermal resonance is shifted outside the plasma and only down-shifted emission is measured. Assuming a bi-Maxwellian distribution a good fit of the spectrum is possible (dotted line in fig. 4). Spectra obtained by variation of plasma parameters are given in Fig. 5. Doubling the density leads to a strong decrease of the detected radiation temperature, although the central plasma temperature remains nearly unaffected. If the ECRH power is changed from 400 kW to 800 kW the low frequency peak increases and shifts to lower frequencies, as expected because of higher central temperature. By varying the radial position of the vertically viewing antenna it was found out that the radial extent of the spatial distribution of the suprathermals is only about 2 - 3 cm, coinciding with the power deposition zone [6]. Calculations result in a strong dependence of the central density of the suprathermals on \( n_{e0}^{th} \) and \( P^{ECRH}, 0.5\%n_{e0}...n_{e0}^{th}...3\%n_{e0} \), whereas there is only a weak variation of \( T_{e1}^{th} \) between 6keV and 8.5keV. Switch-off experiments give a further possibility to distinguish between the thermal and the suprathermal component of the spectrum [7, 6].
The suprathermal radiation disappears on a time scale of about 1 ms which is attributed to the loss mechanism or at least to the collisional slowing-down. That shows the direct connection between ECRH power deposition and creation of suprathermal electrons.

### 2.3 Different magnetic configurations

At W7-AS the depth of the toroidal mirror of the magnetic field configuration can be changed. Measurements have been made in the so-called "minimum B" and "maximum B" configurations. In the latter a significantly smaller number of electrons in the ECRH launching plane is trapped. Fig. 6 shows the response of the ECE-emission after the switch-off of the ECRH for the two configurations. It is not possible to decide exactly the maximal radiation temperature in case of "minimum B" configuration because of the technically determined measurement gap between 138 and 141 GHz. But the "minimum B" configuration shows a larger amount of
suprathermal electrons under otherwise same conditions (integration of suprathermal emission over the hole frequency range). This gives a hint that suprathermals are trapped in the local mirror, which leads to a higher particle density in the dynamic equilibrium of production, slowing down and drift. Besides fig. 6 shows higher energies of the suprathermal in the "minimum B" configuration.

References


Figure 1: Spectrum of the radial Wavenumber $k_r$. 

Figure 2: ECRH power scan.

Figure 3: Density scan.
Figure 4: Measurements vs. simulation.

Figure 5: ECE spectra measured with the vertically viewing antenna. (×) \( n_{e0} = 1 \times 10^{19} \text{ m}^{-3} \), \( P = 800 \text{ kW} \); (○) \( n_{e0} = 2 \times 10^{19} \text{ m}^{-3} \), \( P = 800 \text{ kW} \); (□) \( n_{e0} = 1 \times 10^{19} \text{ m}^{-3} \), \( P = 400 \text{ kW} \).

Figure 6: Time behaviour of the ECE-spectrum after ECRH power switch-off. Left: "minimum B"; Right: "maximum B". No difference can be determined in the decay of the suprathermal radiation within the frame of measurement accuracy.
W7-AS, Programme and Recent Results

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Abstract
The programme of W7-AS concentrates mainly on topics related to the optimisation principles of the WENDELSTEIN stellarator line such as improved equilibria with reduced Shafranov shift, neoclassical transport in improved magnetic configurations including the effects of the ambipolar radial electric field, confinement of high energetic particles, radial diffusion of trapped electrons generated by ECR power deposition. But also topics related to stellarator physics, more generally, are intensely treated, such as the radiative limit at high densities, particle transport, GAE-modes, fluctuations and turbulence, the influence of shear on confinement, and the H-mode. Furthermore new diagnostics with respect to long pulse operation of W7-X are developed. Optimum (neoclassical) confinement discharges with energy confinement times (up to 50 ms) of more than 2.5 times larger than predicted by the International Stellarator Scaling, ISSN95, with $T_i \leq 1.5$ keV are presented as well as discharges with $T_e = 4$ keV showing a strongly positive ECRH driven "electron root" feature of the radial electric field and high density high power NBI discharges leading to $\langle \beta \rangle = 1.8\%$. A 3D Edge code is developed to model edge transport and to optimise energy and particle deposition in a future island divertor. The role of $E \times B$ drifts in islands is considered in detail. Heating scenarios with relevance to W7-X are finally addressed which are independent of cut-off densities and do not necessarily require magnetic resonance conditions, such as the O-X-B heating scenario of ECR waves. The prospects of ICRH for long pulse operation are explored.

Keywords
W7-AS, stellarator, optimum confinement, transport, ISSN95, electric field, ion root, electron root, equilibrium, heating scenarios

1.1 W7-AS
WENDELSTEIN 7-AS, W7-AS, is an Advanced Stellarator [1] (torus radius 2 m, effective plasma radius ≤ 0.18 m) with 5 toroidal periods. The magnetic configuration, with $B_0 = 2.5$ T in the standard case, is optimised by a reduction of the Pfirsch-Schlüter currents. It is generated by a system of 9 non-planar coils in each period with 10 additional planar coils allowing a variation of the rotational transform, $\rho$, from 0.25 to 0.6. It has low shear and stability is given by a moderate magnetic well [2]. Different heating scenarios, electron cyclotron resonance heating, ECRH, with $P \leq 0.4$ MW at 70 GHz and
P ≤ 1.2 MW at 140 GHz, ion cyclotron resonance heating, ICRH, with P ≤ 0.5 MW at the antenna as well as neutral beam injection, NBI, (P ≤ 3.5 MW, 50 keV) give access to net current free plasmas in a wide range: electron densities, \( n_e \leq 3 \cdot 10^{20} \text{ m}^{-3} \) [3], electron temperatures, \( T_e \leq 4 \text{ keV} \) [4] [5] [6], ion temperatures, \( T_i \leq 1.5 \text{ keV} \) [7] [8] [9], and energy confinement times, \( \tau_E \leq 50 \text{ ms} \) [10]. All above results have been derived with a magnetic field on axis of \( B_0 \leq 2.5 \text{ T} \). \( \beta \leq 1.8\% \) has been achieved at \( B_0 \leq 1.25 \text{ T} \) (see chapter 1.3.3). The theoretically predicted reduction of the Shafranov shift by a factor of 2 has been proven experimentally already in an early stage of operation [3]. The bootstrap current of several kA can be compensated integrally either by a small ohmic current or more locally by current drive by ECR [11]. Current drive also can be used to introduce shear for better confinement [12] [4]. With 0.4 MW ECRH at 140 GHz at high densities H-mode transitions with the characteristics as known from tokamaks could be achieved [13].

Plasma build-up can be carried out not only by ECRH but also by non-resonant Rf at 900 MHz leading to a low density target for NBI thus allowing operation independently of the "electron cyclotron resonance fields" of 2.5 and 1.25 T, respectively. Also independent on resonant magnetic fields is the O-X-B heating scenario [14] [15] which recently has been successfully applied at W7-AS. ICRH has been operated successfully with long pulses of up to 1s [16]

### 1.2 Programme

The major element of the optimisation is to reduce the relative Pfirsch-Schlüter current \( \langle |j| / j_\text{total} \rangle^2 \) and - for W7-X - to reduce the bootstrap current. Low PS currents lead to improved equilibria with small Shafranov shift, high stability and strongly reduced neoclassical losses; the minimisation of the bootstrap current avoids a strong disparity between high pressure vacuum field equilibria so that the characteristics of the field design are maintained towards high beta. The optimisation rests on the reduction of average curvature by proper plasma shaping, the repellence of trapped particles from zones of high curvature by the mirror effect and from a proper match of toroidal and helical curvature. Despite of the optimisation neoclassical transport displays the deleterious \( 1/\nu \) collisionality scaling in the long mean free pass, \( \ln \nu \), regime. The presence of an ambipolar field will further reduce the transport coefficients to a tolerable level. It is mandatory to measure the ambipolar field and to assess the balance of
radial species flow. If the ambipolar electric field is determined by radial neoclassical fluxes - because the superimposed turbulent contributions are low or intrinsically ambipolar - various transport equilibria (roots) might develop. Both negative and positive electric fields reduce the neoclassical ion losses. The electron root which develops at high electron temperature with thermally decoupled ions has the advantage of good electron and low impurity confinement.

Stellarator optimisation is expected to improve also the confinement of energetic particles (in case of a reactor that of the α-particles). W7-AS is partly optimised in its field properties and can therefore address some of the above issues. The improvement in equilibrium is well documented [3] [8]; the maximal beta is in the range \( \langle \beta \rangle = 1.8\% \) and is well above that of a classical stellarator [3] [9]. The bootstrap current is dominated by toroidal curvature in W7-AS and increases the rotational transform just like in a tokamak. The measured bootstrap current agrees with the expected one in a large range of collisionalities [11].

The ambipolar electric field and its impact on transport will be discussed in detail below.

The confinement of NBI ions is complete in case of slowing down times \( \tau_s \) smaller than the energy confinement time \( \tau_E \). For cases with high \( T_e \) when \( \tau_s \) can be much larger than the thermal \( \tau_E \), the measured slowing down spectra cannot be modelled satisfactorily with the assumption of complete ion confinement. The discrepancy between measurement and the expectation from good confinement is also borne out by neutron flux measurements in case of D into D -injection. The results are still preliminary and the loss mechanisms are not yet identified [17]. With respect to transport of energetic electrons the fast radial diffusion of trapped electrons, heated by the interaction with ECRH, is observed in power deposition measurements (fast ECRH modulation). The deposition profile deviates from that expected from ray-tracing and absorption and shows wings in the plasma periphery [18]. The widening of the deposition profile is well modelled [19].

There are important topics in the development of stellarators which are, at present, beyond the reach of a rigorous optimisation like the high density operational range [20], turbulent transport [21], [22], aspects of MHD stability like the beam induced Global Alfvén waves, GAE, [22], [23], and the exhaust conditions. Present stellarators can operate at high density (see chapter 1.1) which will ease the conditions for plasma exhaust. In addition, stellarators can safely operate at high density. The operational limit is not connected with a virulent MHD
phenomenon but is a radiative loss of plasma energy. Therefore, the collapse time is slow along the energy confinement time, the plasma energy is homogeneously distributed by radiation and there is time to interfere and to restore a stable plasma [20]. The confinement of W7-AS is determined by the superposition of neoclassical transport in the plasma core and turbulent transport in the gradient region towards the edge. The global confinement is determined by the turbulent processes. The pattern of confinement of the low-shear field system of W7-AS is rather capricious. Depending on the choice of rotational transform $\phi$ confinement can be good or low. The energy content $W$ shows a strong variability with $\phi$. The following picture emerges: The confinement is good, when the iota profile fits into a iota range (edge to core) with few low-order resonances as it is the case around the major resonances $1/3$ and $1/2$. Outside these intervals, shear is required for good confinement [12] [4]. Sufficient shear is provided by PS- and bootstrap current - if the plasma beta is large enough. Outside the resonances free intervals, the development of good confinement depends on a cyclic process because improved confinement enhances the pressure-driven currents which feedback positively on confinement. If the pressure driven currents are enforced by operation at high beta or low collisionality, the interplay between pressure driven currents and confinement is avoided and the sensitivity of $W$ on $\phi$ is lost. Because of the complexity of the dependence of energy content and field distribution not yet all details are satisfactorily understood.

The base-line confinement of W7-AS is provided in the good confinement intervals. It resembles the L-mode confinement of tokamaks. In these $\phi$-ranges and possibly with further restrictions, improved confinement regimes can develop. It is important that the conditions of external confinement as provided in stellarators are also capable of allowing states with improved confinement. An example which is well summarised in the literature is the H-mode of W7-AS [13]. Below, another regime is described.

1.3 Recent Results
Selected topics from the W7-AS programme are briefly discussed below.

1.3.1 Optimum Confinement Discharges
At W7-AS, purely NBI heated and combined NBI/ECRH discharges at medium and high electron densities, $n_e$ between $1.1 \cdot 10^{20}$ and $0.5 \cdot$
$10^{20}$ m$^{-3}$, with $T_e \geq T_i = 0.8 - 1.5$ keV lead to high performance if good wall conditioning and low recycling are provided. Energy confinement times up to $\tau_E = 50$ ms are determined. The experimental transport analyses in the plasma core are consistent with the neoclassical predictions from DKES code [24], [25]. In purely NBI heated discharges with a power of $P_{\text{NBI}} \leq 400$ kW, $\tau_E = 50$ ms has been determined [10], exceeding the ISS95 scaling [26] by more than a factor of 2.5, see fig. 1. During the gradual transition to good confinement the density profile becomes narrower, whereas the temperature profile broadens. In both profiles the gradients steepen and the radial electric field $E_r$ decreases in the gradient region to $E_r \leq 300$ V/cm [10]. Both, density and temperature show very low values at the edge. Also in high power combined NBI/ECRH discharges, narrow density profiles, steep temperature gradients and large negative $E_r$ (up to -700 V/cm) close to the plasma edge are found where locally $n_e$ becomes very small (fig. 2). For this type of discharges with $T_e > T_i \leq 1.5$ keV, optimum confinement properties are found. $\tau_E$ exceeds the ISS95 scaling by a factor of about 2.5 as in the purely NBI heated case (fig. 1). The experimental particle fluxes as well as the ion and electron energy fluxes are in good agreement with the neoclassical predictions up to $r_{\text{eff}} = 12$ cm, see Refs. [7] - [9]. The radial electric field, $E_r$, obtained from the ambipolarity condition of the neoclassical fluxes is consistent with the experimental one deduced from the poloidal rotation of tracers measured by active CXRS in the plasma core (up to 12 cm) as well as by electron impact spectroscopy [27] and probe measurements [28] at the edge. The observation of neoclassical ambipolar radial electric fields up to the very edge where the turbulent fluxes dominate the neoclassical ones indicate that the anomalous particle fluxes are intrinsically ambipolar. Deviations in the intermediate range may be due to the violation of the condition of small electric fields for neoclassical theory [18]. In all discharges with optimum confinement the measured density profiles at outer radii are very similar for all central densities from 0.5 to $1.1 \cdot 10^{20}$ m$^{-3}$ and heating powers between 0.4 and 1.3 MW. Also, the $T_e$ and $T_i$ profiles are very similar in this region [18]. Prerequisites for this type of discharges are good wall conditioning and very low recycling in order to obtain narrow density profiles and to provide global density control even for high NBI power levels with the related strong particle sources of up to $2.5 \cdot 10^{20}$ s$^{-1}$. All discharges have been carried out at $t = 1/3$ at $B_0 = 2.5$ T with a vertical field, $B_z = 0.01 \cdot B_0$. As a consequence the plasma is shifted towards the inboard limiters thus influencing the recycling behaviour but also the magnetic
configuration is improved with respect to neoclassical transport, because the magnetic ripple at the location of strong (grad B) is reduced [29]. Once the above described conditions are fulfilled optimum confinement is reproducibly attained. During the next experimental campaign the question of attaining optimum confinement under separatrix conditions around $+ = 1/2$ will be a major topic of investigation.

1.3.2 ECRH driven "electron root" feature

In low density discharges, heated by ECR ($P_{ECRH} = 400$ kW, 2nd harmonic X-mode at 140 GHz), very strong positive $E_r$ of up to $+600$ V/cm close to the plasma centre with peaked electron temperature profiles and $T_e(0) = 4$ keV, $T_i(0) = 0.3$ keV and flat density profiles ($n_e(0) = 0.2 \cdot 10^{20}$ m$^{-3}$) have been measured [4] - [6], see fig. 3. $\chi_e$ in the intermediate radial range agrees with the neoclassical prediction for the slightly positive $E_r$ (ion root), in the centre it is larger than calculated for the neoclassical (thermal) electron root but much lower than neoclassically predicted if $E_r$ is assumed to be zero, see fig. 3. These large positive electric fields, correlated with strongly peaked electron temperature profiles have only been found at X-mode ECRH with powers above 400 kW in magnetic configurations of W7-AS where a significant part of the ECRH power is absorbed by ripple trapped electrons close to the axis [4] - [6] [18]. These features are lost if the ECRH power is decreased or if the configuration is changed such that the toroidal ripple is reduced. The evidence, that ripple-trapped suprathermal electrons induced by ECRH generate the strong positive $E_r$ in the centre, is supported by Monte Carlo simulations (in 5D phase space) [19].

1.3.3 Equilibrium and Stability

For a discharge heated by about 2.2 MW of NBI at $B_0 = 1.25$ T with $B_z = 0.026 B_0$ with $n_e(0) = 2 \cdot 10^{20}$ m$^{-3}$ and $T_i = T_e = 0.35$ keV the free boundary equilibrium code NEMEC [30] predicts a central $\beta_0$ of 4% and an averaged $\langle \beta \rangle$ of 1.8%. This is in excellent agreement with the $\beta$ profile evaluation from Thomson scattering data, see fig. 4, where the ion contribution is estimated neglecting the effect of $Z_{eff}$ which is expected to be small for this discharge. Furthermore the shift of the magnetic axis as deduced from the maximum of the $\beta$ profile agrees very well with the NEMEC code result. Though the Shafranov
shift of W7-AS is reduced by a factor of 2 compared to classical stellarators [3] it is still considerable. Nevertheless, the equilibrium $\beta$ limit has not been reached, so far. In the upper part of fig. 4 the surfaces of constant soft X-ray emission intensity are shown as are obtained by tomographic reconstruction using 2 soft X-ray cameras [8]. They are in good agreement with the NEMEC results. Though in high $\beta$ configurations with $\beta_0 = 4\%$ the magnetic well is reduced and NEMEC predicts resistive interchange unstable regions at the plasma edge no instabilities have been observed, so far.

1.3.4 First Steps to an Island Divertor

The W7-AS ,,high iota" configurations (iota > 0.5), bounded by ,,natural" islands of considerable size, have the basic edge topology required for an island divertor. The presently installed target arrangement (10 inboard target plates, preserving both the fivefold periodicity and up/down symmetry of the configuration) is the first step on the way to an optimised divertor, which will start operation within the next two years.

Extensive measurements of the island edge structures by plasma spectroscopy and target calorimetry are in excellent agreement with predicted vacuum and equilibrium configurations, which are available up to central $\beta$ values of $\approx 1\%$. Analysis of high density NBI discharges for $\iota = 5/9$ gives strong indications of stable high recycling conditions for $<n_e> \geq 10^{20}$ m$^{-3}$ [31]. The observations are reproduced by the 3D plasma edge transport code EMC3 [32] which has been coupled selfconsistently to the EIRENE [34] neutral gas code.

For low density ECRH discharges, poloidal asymmetries in the island SOL have been observed for $\iota = 5/9$ configurations, with the islands deeply cut by the inboard plates through the O-point (fig. 5a) [33]. Higher densities from Langmuir probe array data were measured for the lower or upper island fans, depending on the B field direction (fig. 5b). The asymmetries can be explained by the effects of an $E_r \times B$ drift driven by the radial temperature gradient in the island which leads to a radial gradient of the electrostatic potential. The poloidal density gradient also implies an imbalance of the power and particle fluxes to the plates, which is consistent with calorimetric and $H_\alpha$ measurements. The drift velocities resulting from estimated radial temperature gradients were inserted into the EMC3/EIRENE code, leading to the same phase shift of the density contours as experimentally observed (fig. 5b).
1.3.5 Heating Scenarios with relevance for W7-X

The operation of W7-AS using ECR waves for heating (pure or in combination with NBI) as well as for start up of the plasma desires a proper adjustment of the magnetic field to meet the resonance condition: \( B_0 = 2.5 \, T \) for 70 GHz fundamental O-mode and 140 GHz second harmonic X-mode and \( B_0 = 1.25 \, T \) for 70 GHz second harmonic X-mode waves. Furthermore ECRH is restricted in density because of the cut-off condition. Non resonant plasma build-up has been successfully carried out with turbulent Rf heating at a frequency of 900 MHz in combination with NBI at various magnetic fields between 0.6 and 2.5T (thus "non resonant") starting from a thin cold plasma \( n_\text{e} < 10^{18} \, \text{m}^{-3}, \, T_\text{e} < 100 \, \text{eV} \). Another scenario which is neither restricted to a certain fixed magnetic field nor by an upper limit of the density is heating by electron Bernstein waves, EBW. Since these waves cannot be excited from outside the plasma they have to be generated by mode conversion from electromagnetic waves, e.g. in a so called O-X-B process. An O-mode wave launched with an optimum angle oblique to the magnetic field vector is converted into an X-mode at the O-mode cut-off density. This X-wave propagates towards the upper hybrid resonance layer where it is converted into an EBW. This heating scenario has been successfully applied to W7-AS [14], [15].

Fig. 6 shows an example where at a non resonant magnetic field of \( B_0 = 2.0 \, T \) an additional ECH pulse of \( P_{\text{ECH}} = 220 \, \text{kW} \) has been applied to a NBI heated plasma with \( P_{\text{NBI}} = 720 \, \text{kW} \) at a density of \( n_\text{e}(0) = 1.6 \cdot 10^{20} \, \text{m}^{-3} \) which is well above the 70 GHz cut off density of \( n_\text{e} = 0.6 \cdot 10^{20} \, \text{m}^{-3} \). The O-X-B heating leads to a clear increase of the plasma energy content compared to the purely NBI heated case. Increases in the energy content up to 30\% or 2.5 kJ (\( T_\text{e}^{\text{NBI}} = 360 \, \text{eV}, \, \Delta T_\text{e}^{\text{OXB}} = 110 \, \text{eV} \)) with \( P_{\text{ECH}} = 330 \, \text{kW} \) and \( P_{\text{NBI}} = 360 \, \text{kW} \) have been achieved, so far.

Long pulse operation in W7-X will need long pulse heating systems. Besides ECRH in cw operation ion cyclotron resonance heating, ICRH, is foreseen. Since matching of the outer magnetic surfaces with appropriate antenna systems turned out to be much more difficult for stellarators than for tokamaks, ICRH in stellarators is still in an experimental status. Nevertheless ICRH has been successfully operated at W7-AS for a pulse length up to 1 s [16]. Feeding 500 kW to a "broad antenna" [35] led to a deposition of about 200 kW into the
plasma which had been built up by ECRH and taken over by ICRH via minority heating of hydrogen in a deuterium bulk plasma. The purely ICR heated plasma could be sustained up to 1 second. The electron temperature after dropping from about 2.5 keV during the phase of combined ECR and ICR heating to about 400 eV some time after the beginning of the purely ICR heated phase recovers to about 800 eV at the end of the ICRH pulse whereas the line integrated electron density remained constant at about 0.16 \( \cdot 10^{20} \) m\(^{-3}\). During the whole ICRH pulse the impurity radiation remained low. In order to increase the power deposited in the plasma a double strap antenna will be installed during the next experimental campaign.

References


[28] GRIGULL, P. private communication
Figure captions

Fig. 1: Energy confinement times in W7-AS compared to ISS95-W7. Dashed line: ISS95 for all helical devices; solid line: ISS95-W7 for W7-AS data up to 1995; open squares: pure ECRH; full triangles: ECRH H-mode; full circles: low power NBI, leading to $\tau_E = 50$ ms, open circles: NBI with higher power; full (light) diamonds: combined NBI and ECRH, leading to $T_i = 1.5$ keV. Crosses for the same discharges taking the kinetic energy instead of the diamagnetic one. Both, circles and diamonds, show the increase of $\tau_E$ by about a factor of 2.5 compared to ISS95.

Fig. 2: Optimum confinement. Discharge # 34313, combined ECRH and NBI, $P_{ECRH} = 750$ kW, $P_{NBI}^{abs} = 900$ kW, at medium density. $T_e$ from Thomson scattering and ECE (dark symbols) and $T_i$ (light symbols) from active CX-NPA and CXRS (upper left); $n_e$ from Thomson scattering (upper middle), radial electric field (upper right) as deduced from active CXRS, (dots) and passive spectroscopy (circles) compared to the neoclassically calculated ambipolar electric field (crosses). In the lower part the experimental particle and energy fluxes for ions and electrons, respectively, are compared with neoclassical predictions by DKES code taking the ambipolar electric field (upper right) into account.

Fig. 3: "Electron root" feature. Low density ECR heated, $P_{ECRH} = 750$ kW, discharge with strongly peaked $T_e$ (upper left). The strongly positive electric field from CXRS is compared to the neoclassical electron root close to the centre and slightly positive ion root at intermediate radii (lower left). The lower right picture shows the experimental heat conductivity (dash-dotted line) compared to neoclassical predictions if $E_r = 0$ is assumed (dotted line) and if electron and ion roots of $E_r$ are taken into account (crosses).

Fig. 4: Tomographic reconstruction of soft X-ray emission compared to magnetic flux surfaces calculated by NEMEC code for high $\beta$ discharge heated by about 2.2 MW of NBI at $B_0 = 1.25$ T with $B_Z = 0.026$ B, $n_e(0) = 2 \times 10^{20}$ m$^{-3}$ and $T_i = T_e = 0.35$ keV. The shift of the magnetic axis agrees well with that resulting from the $\beta$ profile calculated from Thomson scattering.

Fig. 5: a) left: Poloidal cross section of the $t = 5/9$ configuration at the poloidal plane of an inboard plate, right: Flow to target plate resulting from the $E_r \times B$ drift (schematic). b) Density contour from probe array (curves) compared to EMC3/EIRENE results for positive and negative B-field. For
positive (negative) B the deposition maximum is shifted downward (upward).

**Fig. 6:** Comparison of energy contents of a discharge heated by combined NBI and O-X-B (heating scenario with conversion from O-mode to X-mode to electron Bernstein waves) (upper curve) compared to a purely NBI heated one (lower curve) at $B_0 = 2.0$ T with $P_{ECH} = 220$ kW at 70 GHz and $P_{NBI} = 720$ kW at a density of $n_e(0) = 1.6 \cdot 10^{20}$ m$^{-3}$. 
Fig. 1:

Fig. 2:

Fig. 3:
Fig. 4:
Fig 5a:

Fig 5b:

Fig. 6:
ELECTRON CYCLOTRON HEATING BEYOND THE CUTOFF DENSITY BY O-X-B MODE CONVERSION IN W7-AS

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Abstract
Electron cyclotron heating (ECH) above the plasma cutoff density with electron Bernstein waves (EBW) was successfully demonstrated at W7-AS stellarator. The EBW's were generated via O-X-B mode conversion from O-waves to X-waves and, finally, to electron Bernstein waves. Clear evidences for both mode conversion steps were detected and resonant absorption with a narrow profile was demonstrated.

Keywords:
electron cyclotron heating, electron bernstein waves, mode conversion, o-x-b process

Introduction
The accessible plasma density for electron cyclotron heating (ECH) with electromagnetic waves is limited by the plasma cutoff. For the electrostatic electron Bernstein wave (EBW), the third EC-mode which is able to propagate in a hot plasma, no such limit exists. However, since EBW's could not be exited from the outside, they have to be generated via mode conversion from the electromagnetic waves. This can be performed by the O-X-B process which was proposed by J.Preinhaelter and V.Kopecký [1] in 1973. Here O, X, and B represent the ordinary, extraordinary and the electron Bernstein mode. In a first step an O-wave launched by an antenna is converted into a slow X-
wave at the O-wave cutoff layer. This mode conversion requires an oblique launch near an optimal angle, a plasma density above the O-wave cutoff and a frequency above the first cyclotron harmonic in the plasma. In a second mode conversion an EBW is generated from the slow X-wave at the upper hybrid resonance (UHR), where the X-mode branch of the solution of the hot plasma dispersion relation is connected to the electron Bernstein branch. Since for EBW's no density limit exists, they propagate toward the dense plasma center where they are absorbed by cyclotron damping or in the nonresonant case by collisional multiple pass damping. A detailed description of the O-X-B mode conversion process is found in [2],[3]. In this paper the experimental results of the O-X-B mode conversion and EBW heating will be presented.

**Experimental results**

The experiments were performed at the W7-AS stellarator (major radius R=2.0m, minor radius a=0.18m) with two 70 GHz gyrotrons with 110 kW power each. A detailed description of W7-AS and its 70 GHz ECH system can be found in [4]. The central magnetic field was set between 1.25T and 2.0T and the edge rotational transform \( \tau \), taken from the magnetic reconstruction, near 0.35 according to the experimental requirements. The central density of the neutral beam injection (NBI) sustained target plasma was up to \( 1.6 \times 10^{20} \) m\(^{-3} \), which is more than twice the 70 GHz O-mode cutoff density. Co- and counter NBI with 360 kW power each were used to compensate the momentum transfer to the plasma.

1. **Variation of the launch angle.**

The launch angle of the incident O-mode polarised wave was varied at fixed heating power (220 kW) at a nonresonant magnetic field. The increase of the total stored plasma energy (from the diamagnetic signal) depends strongly on the launch angle (see Fig. 1), which is typical for the O-X-conversion process, and fits well to the calculation. Here the power transmission function given by [5] was normalised to the maximum energy increase. The central density was \( 1.5 \times 10^{20} \) m\(^{-3} \) and the central electron temperature was 500 eV. Heating at the plasma edge could be excluded since at the nonresonant magnetic field of 1.75T no electron
cyclotron resonance existed inside the plasma. Due to technical limitation of the maximum launch angle, only the left part of the transmission function could be proved experimentally. The plasma energy content increased by about 1.5 kJ compared to a similar discharge with neutral beam injection (NBI) only as shown in Fig. 2. Here two 70 GHz beams in O-mode polarisation (110 kW power each) were launched with an angle of 40° with respect to the perpendicular launch into a NBI (720 kW) sustained target plasma with a central density of 1.6 10^20 m^-3 and a central temperature of 560 eV. More than 70% of the heating power was found in the plasma if the power scaling of the energy confinement (P^0.6) was taken into account. Thus O-X-B-heating turned out to be efficient.

2. Density variation and parametric instability (Pl).

In these experiments, it should be demonstrated, that a density threshold (O-cutoff) for the O-X-conversion exists and that the parametric decay process which is a footprint of the X-B conversion takes place. For this the plasma was build up by one 70 GHz gyrotron in X-polarisation in a resonant central magnetic field of 1.25T. Then the density was slowly ramped up to density above the O-cutoff. In parallel as shown in Fig. 3 a second 70 GHz beam O-mode polarised with the optimal launch angle and modulated with 20% amplitude was launched into the plasma. During the plasma build up thermal EC emission (ECE) was detected. As soon as the cutoff density is reached ECE vanished and O-X-B heating started, which caused an increase of the plasma energy and central soft-x emission shown in Fig. 3. Simultaneously the PI at the X-B-conversion process generated a decay spectrum, whose high frequency part could be measured with the ECE-detector. The modulation amplitude strongly exceeded that of the pump wave, what clearly demonstrated the nonlinear character (power threshold) of the PI. Fig. 4 shows the high frequency decay spectrum. Two red shifted and one blue shifted lines can be recognised. Their spectral distances to the 70 GHz pump wave, which was suppressed by a Notch filter, are multiples of the lower hybrid resonance (LHR) frequency (~900MHz).

The spectrum of the LHR oscillation itself could be detected by a loop
antenna. The LHR oscillation shows a high degree of correlation with the high frequency decay waves.

3. Resonant cyclotron absorption.
Here the central magnetic field was varied between 2.0T and 2.5T to show resonant absorption of the EBWs. In the equatorial plane the magnetic field as a function of the effective radius \( r_{\text{eff}} \) is approximately given by the following relation: 

\[
B(r_{\text{eff}}) = B_0 A \left(\frac{A}{A + r_{\text{eff}} / a}\right)
\]

with \( A = 10.5 \). The power deposition was estimated from the change of the temperature profile at the power switch-off. Since the density was far above the ECE cutoff, the temperature profiles were calculated from the soft-X emission and the Thomson scattering diagnostic. The central temperature was 500 eV. The X-ray emission was monitored by an array of 36 silicon detectors with a 25 µm beryllium filter. To obtain the radial X-ray emission profile the signals were inverted to the magnetic flux co-ordinates. The time resolution was 0.1 ms and the radial resolution was about 1 cm. In Fig. 5 the absorption profiles for different magnetic fields are shown. The absorption is strongly Doppler shifted due to the oblique launch and moves from the high field side at 2.0T through the center (2.2T) to the low field side at 2.3T with increasing magnetic field, which clearly demonstrates the propagation and the local cyclotron absorption of the EBW's for the first time.

**Conclusions**
ECH of an overdense plasma with 70 GHz electron Bernstein waves was clearly demonstrated at W7-AS. The EBW's were generated via mode conversion in the O-X-B process. Both, the angular dependence of the O-X-conversion and the parametric instability which is typical for X-B-conversion could be experimentally verified. The position of the narrow absorption profile, estimated from the soft-X emission, could be changed by a shift of the cyclotron resonance layer. Thus generation, propagation and local resonant cyclotron absorption of EBW's was shown, which is an excellent test of hot plasma wave theory.
References

Figures

Fig. 1: Increase of the plasma energy content by O-X-B-heating versus the longitudinal vacuum refractive index $N_Z = \cos \varphi$ of the incident O-wave ($\varphi$: launch angle). The solid line is the calculated transmission function multiplied by the maximum energy increase.

Fig. 2: Energy content (diamagnetic signal) of a NBI-discharge with and without O-X-B-heating.
Fig. 3: Temporal development of some plasma parameter during a O-X-B heated discharge. From the top: plasma energy estimated from the diamagnetic signal, average density from the interferometric measurement, heating power, intensity of ECE and PI, central soft X signal. The markers show the O-X-B heating interval.

Fig. 4: High frequency spectrum of the parametric decay waves generated in the X-B-process. The incident wave frequency is 70 GHz and the LH frequency is about 900 MHz.
References


Figures

Fig. 1: Increase of the plasma energy content by O-X-B-heating versus the longitudinal vacuum refractive index $N_z = \cos \phi$ of the incident O-wave ($\phi$: launch angle). The solid line is the calculated transmission function multiplied by the maximum energy increase.

Fig. 2: Energy content (diamagnetic signal) of a NBI-discharge with and without O-X-B-heating.
Fig. 5 Changes of temperature 3ms after O-X-B heating switch-off and the related ECRH absorption profiles different central magnetic fields.
Density Control Problems in Large Stellarators with Neoclassical Particle Transport

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Abstract
With respect to the particle flux, the off-diagonal term in the neoclassical transport matrix becomes crucial in the stellarator long-mean-free-path regime. Central heating with peaked temperature profiles can make an active density profile control by central particle refuelling mandatory. The neoclassical particle confinement can exceed significantly the energy confinement at the outer radii. As a consequence, the required central refuelling may be larger than the neoclassical particle fluxes at outer radii leading to the loss of the global density control.

1 Introduction
Within neoclassical theory, the particle and energy fluxes are linked together. The energy dependence of the $\nabla B$ drift in the stellarator $l m f p$ regime, where the neoclassical transport is dominated by the ripple trapped particles [1], results in significant non-diagonal terms in the transport matrix. Central power deposition and, as a consequence, peaked temperature profiles will lead to significant outward particle fluxes. To avoid very hollow density profiles or even pressure profiles (a positive pressure gradient will lead to strong MHD instability), central particle sources (e.g., central refuelling by pellets) are needed. In larger stellarators, gas puffing as well as wall recycling will only affect the region close to the plasma edge. Outside of the power deposition region at lower temperature, the neoclassical particle fluxes may be in conflict with the central refuelling rate needed to control the inner density profile. The appearance of a particle transport barrier can lead to the loss of the global density control. In particular, the effect of different stellarator magnetic-field configurations on the expected density control problems is analyzed in this paper.

2 Mono-energetic transport coefficients
In the simplest approximation, a classical stellarator can be described by the superposition of a toroidal and a helical magnetic field, with magnitude $B/B_0 = 1 - \epsilon \cos \theta - \epsilon_h \cos (m \theta - N_p \phi)$, where $\theta$ [$\phi$] is the poloidal [toroidal] angle, $\epsilon$ [$\epsilon_h$] is the poloidal [helical]
modulation of $B$ (need not be equal to $\varepsilon_t = r/R$, the inverse aspect ratio) and $m$ [$N_p$] is the helical field multipolarity [period number]. In the imfp regime particles localized in the helical ripple may experience significant radial drift and are generally expected to provide the dominant contribution to the overall neoclassical transport. Within the traditional analytic theory (e.g., see Refs. [1, 2]) the form this contribution takes is determined by the relative magnitudes of the effective collision frequency, $\nu_{eff} = \nu/2\varepsilon_h$, and the $E \times B$ precessional frequency, $\Omega_E = E_r/rB_0$. For the simple model field of a classical stellarator, the results may be summarized by the “mono-energetic” diffusion coefficients in the $1/\nu$, the $\sqrt{\nu}$ and the $\nu$ collisionality regimes

\[
D_{1/\nu} = \frac{4}{9\pi} \left( \frac{v_d}{\varepsilon_t} \right)^2 \left( \frac{2\varepsilon_h}{\nu} \right)^{3/2}
\]

(1)

\[
D_{\sqrt{\nu}} = \frac{4\sqrt{2}}{9\pi} \left( \frac{v_d}{\varepsilon_t} \right)^2 \frac{\nu^{1/2}}{|\Omega_E|^{3/2}}
\]

(2)

\[
D_\nu = \left( \frac{v_d}{\Omega_E \varepsilon_t} \right)^2 \frac{\nu}{2F_{\text{bl}}}
\]

(3)

which are predicted analytically for $\nu_{eff} \gg \Omega_E$, $\nu_{eff} \ll \Omega_E$ and $\nu_{eff} \ll \Omega_E$, respectively. In these expressions, $v_d = mv^2 / 2qB_0R$ is the $\nabla B$-drift velocity of mono-energetic particles (the averaged radial component of the $\nabla B$-drift is decreased by the factor $\varepsilon/\varepsilon_t$ corresponding to the reduction of the averaged toroidal curvature in an advanced stellarator), and $F_{\text{bl}} = \sqrt{\varepsilon + 2\varepsilon_h} - \sqrt{2\varepsilon_h}$.

For classical stellarator configurations, which can be reasonably well represented by the simple model field, the predictions of the bounce-averaged analytic theory in eqs. 1 to 3 have been compared with the numerical solution of the mono-energetic drift kinetic equation by using the DKES code [3, 4]; good agreement was found [5]. Additionally, the analytic approach has also been confirmed by Monte-Carlo simulations [2, 6].

For a more complex magnetic field geometry, however, the traditional analytic approach cannot be used. The magnetic field strength on flux surfaces is represented by the $m, n$ Fourier modes, $\beta_{mn}(r)$, with respect to the poloidal, $\theta$, and toroidal angle, $\phi$, in magnetic (Boozer) co-ordinates. The coefficients of the simple model field for classical stellarators are given by $\varepsilon \equiv -\beta_{10}$ and $\varepsilon_h \equiv \beta_{m1}$; see above. For a given Fourier spectrum, $\beta_{mn}(r)$, the DKES code estimates the mono-energetic transport matrix as a function of $\nu/\nu$ and $E_r/\nu$. Neglecting the Ware pinch term, the particle and energy transport is determined completely by the mono-energetic transport coefficient.

DKES computations of this mono-energetic transport coefficient
in the \textit{Imfp} regime for quite different stellarator configurations (even in the case of a fairly broad \(\beta_{mn}\) Fourier spectrum as is necessary to describe the W7-AS configurations) showed that the \(1/\nu\), the \(\sqrt{\nu}\) and the \(\nu\) collisionality regimes are clearly present [5]. Furthermore, the different dependencies on the radial electric field for each regime could also be identified corresponding to the traditional analytic theory; see eqs. 1 to 3. Consequently, these analytical transport coefficients were used to fit the DKES results for the different regimes, i.e., numerical parameters only depending on radius have been obtained to fit the “analytical” transport coefficients to the DKES data. Additionally, the different regimes are smoothly connected (using \(D^{-1} = D_{1/\nu}^{-1} + D_{\sqrt{\nu}}^{-1} + D_\nu^{-1}\)).

3 Thermal neoclassical transport matrix

For each particle species \(\alpha\), the “thermal” transport matrix, \(D^\alpha_{jk}\), is obtained by convolution with respect to the Maxwellian based on the mono-energetic transport coefficient. In the energy convolution both \(\nu(v)/v\) and \(E_r/v\) vary significantly leading to a mixing of the mono-energetic transport regimes. Neglecting the Ware pinch term, the neoclassical particle and energy flux densities, \(\Gamma_\alpha\) and \(Q_\alpha\), are given by

\[
\Gamma_\alpha = -n_\alpha \cdot \left\{ D_{11}^\alpha \left( \frac{n'_\alpha}{n_\alpha} - \frac{q_\alpha E_r}{T_\alpha} \right) + D_{12}^\alpha \frac{T'_\alpha}{T_\alpha} \right\} \tag{4}
\]

\[
Q_\alpha = -n_\alpha T_\alpha \cdot \left\{ D_{21}^\alpha \left( \frac{n'_\alpha}{n_\alpha} - \frac{q_\alpha E_r}{T_\alpha} \right) + D_{22}^\alpha \frac{T'_\alpha}{T_\alpha} \right\} \tag{5}
\]

with \(\alpha = e, i\) (impurity fluxes are omitted here), and \(q_\alpha\) being the particle charge. The so-called “convective term”, \(\frac{3}{2} T_\alpha \Gamma_\alpha\), is included in this neoclassical definition of \(Q_\alpha\). For the total energy flux density, an additional \(T_\alpha \Gamma_\alpha\) contribution is obtained from the term \(\nabla \cdot \Pi \cdot v\) (with the approximation \(\Pi \simeq n T_1 \frac{1}{2}\) for the pressure tensor, and \(\Gamma = n u_r\)), and taken into account in the energy balance.

The radial electric field is determined by the roots of \(Z_i \Gamma_i = \Gamma_e\), additional non-ambipolar particle fluxes are disregarded. As the \(D^\alpha_{jk}\) depend on \(E_r\), multiple roots of the ambipolarity condition may exist, e.g., see Refs. [7, 8]. Only in the very \textit{Imfp} regime, i.e., at low density, may the additional “electron root” at fairly strong \(E_r > 0\) (roughly given by \(\Gamma_e = 0\)) be expected. The “ion root” with negative \(E_r\) for \(T_e \sim T_i\) is predicted for all collisionalities. As at higher densities only the “ion root” can exist, the following analysis considers only this case.

In general, the ambipolar electric field must be determined numerically (roots of \(Z_i \Gamma_i = \Gamma_e\)). Under the assumption of very “extended” mono-energetic transport regimes, the mixing due to the
energy convolution can be neglected. For these “pure” regimes, the thermal transport matrix can be easily estimated. With the normalisation \( \delta_{jk} = D_{jk} / D_{11} \) and \( \Delta(\tau) \) describing the additional radial dependence as, e.g., obtained from numerical fits of the “analytical” coefficients given by eqs. 1 to 3 to the DKES results, the thermal transport matrix is given for the different regimes by

\[
D_{1/\nu} = \Delta_{1/\nu} \frac{1}{n} T^{7/2}; \quad \delta_{12} = \frac{7}{2}, \quad \delta_{21} = 5, \quad \delta_{22} = \frac{45}{2} \tag{6}
\]

\[
D_{\sqrt{\nu}} = \Delta_{\sqrt{\nu}} \sqrt{n} T^{3/4} \left( \frac{r}{E_r} \right)^{3/2}; \quad \delta_{12} = \frac{5}{4}, \quad \delta_{21} = \frac{11}{4}, \quad \delta_{22} = \frac{99}{16} \tag{7}
\]

\[
D_{\nu} = \Delta_{\nu} n T^{1/2} \left( \frac{r}{E_r} \right)^2; \quad \delta_{12} = \frac{1}{2}, \quad \delta_{21} = 2, \quad \delta_{22} = 3 \tag{8}
\]

\[
D_{pl} = \Delta_{pl} T^{3/2}; \quad \delta_{12} = \frac{3}{2}, \quad \delta_{21} = 3, \quad \delta_{22} = \frac{15}{2} \tag{9}
\]

Here, the plateau regime scaling is included, the other axisymmetric transport regimes can be disregarded. Note in this context, that the transport matrix in the \( \nu \) regime scales differently from the tokamak “banana” regime \( (\delta_{12} = -\frac{1}{3}, \delta_{21} = 1, \delta_{22} = \frac{1}{3}) \). Furthermore, the “effective” helical ripple, \( \langle \epsilon_h \rangle \), the generalisation of \( \epsilon_h \) for non-classical stellarator configurations, is defined by \( \Delta_{1/\nu} \propto \langle \epsilon_h \rangle^{3/2} \).

4 Link of Particle and Energy Fluxes

The central particle sources needed to control the density profile (e.g., refuelling by pellets) are related to the power deposition by the neoclassical link of particle and energy fluxes. With the assumption of “pure” transport regimes, this link will be estimated in the following. For simplicity, \( T \equiv T_e \approx T_i \) is assumed (which is the case at higher density), and only the combined energy balance is considered. For this case, the collisional power transfer as well as the \( \Gamma_\alpha E_r \) terms in the electron and ion energy balances cancel. Furthermore, \( n \equiv n_e = Z_i n_i = \text{const} \), i.e., external density profile control by central particle sources is assumed. Then, the ambipolarity condition is given by

\[
(D_{11}^e + Z_i D_{11}^i) \frac{e E_r}{T} = (D_{12}^e - D_{12}^i) \frac{T'}{T}. \tag{10}
\]

In order to analyze the neoclassical link of the particle and energy fluxes, the radial electric field from the ambipolarity condition is discussed for two limits: \( D_{jk}^e \ll D_{jk}^i \) and \( D_{jk}^e \approx D_{jk}^i \). The first approach is valid at the outer radii with lower temperature independent of the specific magnetic configuration, while the second one becomes important only in the central region for configurations with a significant \( 1/\nu \) regime, e.g., for the W7-X “high-mirror” configuration.
The limit $D_{jk}^e \ll D_{jk}^i$:

The ambipolar $E_r$ can be estimated from $\Gamma_i \simeq 0$, and the ambipolar particle flux density, $\Gamma$, is approximately given by the electron flux density with this $E_r$ taken into account,

$$\Gamma = \Gamma_e \simeq -n D_{11}^e \left( \frac{1}{Z_i} \frac{\delta_{12}^i}{\delta_{12}^e} \right) \frac{T'}{T}.$$  \hspace{1cm} (11)

The total neoclassical energy flux density, $Q_{\text{tot}} = Q_e + Q_i + 2\Gamma T$ (with the additional $\Gamma T$ from the pressure tensor term taken into account), is related to the particle flux density by

$$\Xi = \frac{Q_{\text{tot}}}{\Gamma T} = 2 + \xi_e + \xi_i \frac{D_{11}^i}{D_{11}^e}.$$  \hspace{1cm} (12)

with $\xi_e = \frac{\delta_{22}^i + \frac{1}{2} \delta_{12}^i \delta_{21}^i}{\frac{1}{2} \delta_{12}^e + \delta_{12}^i}$ and $\xi_i = \frac{\delta_{22}^i - \frac{1}{2} \delta_{12}^i \delta_{21}^i}{\frac{1}{2} \delta_{12}^e + \delta_{12}^i}$.

The electrons are assumed to be in the $1/\nu$ regime, and $Z_i = 1$. Then, $\xi_e = 6.25$ and $\xi_i = 0.5$ for the ion $\nu$ regime, $\xi_e = 6.05$ and $\xi_i = 0.58$ for the ion $\sqrt{\nu}$ regime, and $\xi_e = 6$ and $\xi_i = 0.6$ for the ion plateau regime. Consequently, these ratios turn out to be fairly equivalent for the different ion transport regimes. Most important is the ratio of the ion to the electron transport coefficient, $D_{11}^i / D_{11}^e < \sqrt{m_i/m_e}$, which determines the fraction of the ion energy flux. For the electrons at the beginning of the $1/\nu$ regime, the ion energy flux dominates, and $\Gamma T$ is much less than $Q_{\text{tot}}$.

The limit $D_{jk}^e \approx D_{jk}^i$:

The electron and ion transport coefficients can be comparable in magnitude if the temperature is sufficiently high, and if the electron $1/\nu$ regime is significant. These conditions can be realized in the central region in configurations with a fairly large toroidal mirror term. Then, the ambipolar radial electric field in eq. 10 becomes small in magnitude, and is implicitly given by $D_{12}^i \approx D_{12}^e$. Neglecting $E_r/T$ as a driving term for these conditions, the ambipolar particle flux density, $\Gamma$, is linked to the total neoclassical energy flux density

$$\Xi = \frac{Q_{\text{tot}}}{\Gamma T} = 2 + \frac{\delta_{22}^e}{\delta_{12}^e} + \frac{\delta_{22}^i}{\delta_{12}^i}.$$  \hspace{1cm} (13)

Since the ratios $\delta_{22}^e / \delta_{12}^e$ corresponding to the different transport regimes are rather similar ($1/\nu$: 6.43, $\nu$: 6, $\sqrt{\nu}$: 4.95, and plateau regime: 5), the electron and ion energy fluxes are also comparable. For example, for 10 MW central heating at $T(0) = 5$ keV, a particle source of about $10^{21}$ /s is needed.
5 Need for central refueling

The link of the neoclassical energy and particle fluxes in the bulk part of the plasma leads to a requirement for particle sources to control the central density profile. The assumed purely neoclassical particle flux density defines the particle source profile, \( S_p(r) = r^{-1} d(r\Gamma) / dr \). In the combined energy balance, additional losses are taken into account, \( P^*(r) = r^{-1} d(r Q_{\text{tot}}) / dr \), with \( P^* = P_h - P_l - r^{-1} d(r Q_{\text{an}}) / dr \) where \( P_h(r) \) is the heating power density profile, \( P_l(r) \) describes radiative losses, and \( Q_{\text{an}} = -n \chi_{\text{an}} T' \) is an additional energy flux density due to an “anomalous” heat diffusivity, \( \chi_{\text{an}} \). Then, the particle source profile needed to control the density profile (\( n = \text{const.} \)) is given by

\[
S_p = \frac{P^*}{T^2} - \left( \frac{T'}{T} + \frac{\Xi'}{\Xi} \right) \Gamma .
\]  

(14)

For conditions where \( D_{jk}^i \approx D_{jk}^\xi \) holds, \( \Xi' = 0 \), see eq. 13. Neglecting the additional \( T' \) term the profile of the particle sources, \( S_p(r) \), must be quite similar to the power deposition profile, \( P_h(r) \), if the additional energy losses are negligible which is reasonable for peaked central heating. These estimates are independent of the specific magnetic configuration as long as the condition \( D_{jk}^\xi \approx D_{jk}^i \) holds (which is, however, violated at outer \( r \) with lower \( T \)).

The other limit, \( D_{jk}^\xi \ll D_{jk}^i \), is rather similar, but the additional \( r \) and \( T \) dependence of the ratio \( D_{11}^i / D_{11}^\xi \) in eq. 12 must be analyzed. With \( \Gamma_i \approx 0 \), \( eE_r \approx 1.25 T' \) is obtained from eq. 10 for the ion \( \sqrt{\nu_e} \) regime leading to \( D_{11}^i \propto T^{5/4} (r/T')^{3/2} \) where an additional \( r \) dependence in \( \Delta_{\sqrt{\nu}} \) is neglected. \( D_{11}^\xi \) in the \( 1/\nu \) regime reflects the radial dependence of the helical ripple. Corresponding to eqs. 1 and 6, the electron transport coefficients scale with \( T^{7/2} (\langle \nu_e \rangle / \langle \nu_h \rangle)^{3/2} \) where \( \langle \nu_h \rangle \) is the “effective” helical ripple. For high-mirror advanced stellarators such as W7-X, \( \langle \nu_h \rangle \) is mainly determined by the toroidal mirror term, i.e., \( \langle \nu_h \rangle \) is roughly constant in the inner region leading to a dominant \( 1/\nu \) regime (i.e., eq. 13 is fulfilled). For configurations with only a small toroidal mirror term, \( \langle \nu_h \rangle \) has the same dependence on \( r \) as the leading helical Fourier mode: \( \langle \nu_h \rangle \propto \beta_{11} \propto r \) for \( m = 1 \) configurations (e.g., see Ref. [9] for the W7X “low-mirror” configuration), and \( \langle \nu_h \rangle \propto \beta_{21} \propto r^2 \) for \( m = 2 \). Then,

\[
\frac{D_{11}^i}{D_{11}^\xi} \propto \left( \frac{r}{T' \langle \nu_h \rangle} \right)^{3/2} T^{-6/5}
\]  

(15)

is obtained. For configurations with \( m \geq 1 \), \( \Xi \to \infty \) for \( r \to 0 \) leads to smaller \( S_p \) close to the axis compared to a “high-mirror” configuration.

Consequently, for a peaked central power deposition profile a central particle refuelling is also needed in the \( D_{jk}^\xi \ll D_{jk}^i \) case. The
required particle source strength is comparable with that needed for $D_{i1}^j \approx D_{jk}^j$. Furthermore, the particle source profile needed to control the central densities turns out to be broader for $m \geq 1$ configurations (due to $\Xi' < 0$).

6 Neoclassical particle transport barrier

The particle flux at outer radii (i.e., outside of the power deposition region) must be at least as large as the central particle source. Otherwise the global density control may be lost, i.e., the outer density increases. In this context, the outermost particle sources from recycling are disregarded; shielding is assumed to dominate at the outermost radii where the strong density gradient is established.

An outer particle transport barrier may be formally defined by $S_p \leq 0$. Additional energy losses or an enhanced energy flux tends to establish such a scenario. In principle, the temperature profile, which differs significantly for the magnetic configurations under investigation, must be estimated. In eq. 14, however, the main effect of the magnetic configuration on the appearance of this crucial particle transport barrier is directly included in the $\Xi'$ term. With the $\sqrt{\nu}$ regime scaling of $D_{i1}^i$ and the $1/\nu$ regime scaling of $D_{i1}^i$, and neglecting the radial dependence of $T'$ (in eq. 15) leads to

$$
\Xi' = -\xi_i \frac{D_{i1}^i}{D_{i1}^f} \{ \frac{9}{4} \frac{T'}{T} + \frac{3}{2} \left( \frac{\langle \epsilon_h \rangle}{r} \right)' \}.
$$

(16)

Here, $\Xi' > 0$ amplifies the transport barrier problem, i.e., $S_p$ is decreased. Since $\xi_i D_{i1}^i / D_{i1}^f > 0.5 \Xi$ (the relative fraction of the ion energy flux; compare eq. 12), the $T'/T$ term in eq. 14 is overcompensated by the $T'/T$ contribution from $\Xi'/\Xi$ in eq. 16 resulting in a negative contribution to $S_p$. The other term $\langle \ln(\langle \epsilon_h \rangle/r) \rangle'$ contains the magnetic configuration effect in $\langle \epsilon_h \rangle$ which can lead to a negative or a positive contribution to $S_p$: $\langle \ln(\langle \epsilon_h \rangle/r) \rangle' = -1/r$ for a high-mirror configuration ($\langle \epsilon_h \rangle \approx \text{const.}$); = 0 for an $m = 1$ configuration without toroidal mirror ($\langle \epsilon_h \rangle \propto r$); and = 1/r for an $m = 2$ configuration ($\langle \epsilon_h \rangle \propto r^2$). Within this simple analysis, an $m = 2$ or even an $m = 1$ configuration with only small toroidal mirror terms are less susceptible to loss of the global density control than an $m = 0$ configuration. In the "high-mirror" W7X configuration, $\langle \epsilon_h \rangle$ even decreases with $r$; see Ref. [9]. This kind of "over-optimisation", i.e., the neoclassical confinement is mainly improved at the outer radii, leads to the most dangerous scenario in this context.

Consequently, $S_p < 0$ outside of the power deposition region is in contradiction to the assumption of a stationary and flat density profile, and the formal "particle sink" needed is equivalent to an
increase of the density \((\partial n / \partial t > 0)\) at these outer radii. With increasing density, the temperature is reduced and, for example, the radiative losses are increased, leading to even more negative \(S_p\). The \(D^H_{1i}\) coefficient related to the positive density gradient is too small to compensate the non-diagonal temperature gradient driven flux. In this way, the neoclassical particle transport barrier can lead to the loss of the global density control.

7 Conclusions

In large stellarators central particle refuelling will be mandatory. To avoid hollow density and pressure profiles or rapidly increasing density, the particle source strength must be nearly proportional to the heating power. Especially for the W7X “high mirror” configuration, the particle and power deposition profiles must be quite similar. A particle transport barrier must be avoided, otherwise the global density control is lost. In W7X, these effects can be studied by varying the toroidal mirror term. For example, lowering the mirror term slightly degrades the particle confinement at outer radii which may help to prevent the loss of the global density control. Especially for the most dangerous W7X “high-mirror” scenario, a “density limit” for stationary operation is predicted. At higher densities, the global density control is lost. This “density limit” increases with heating power, but decreases with radiative losses and additional “anomalous” energy fluxes.

References


Transport Analysis in
Low-collisionality W7-AS Plasmas

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Abstract

At W7-AS, the confinement properties are analyzed and compared mainly with neoclassical predictions for quite different conditions. Low density ECRH discharges allow access to the very long mean free path regime for electrons \( T_e \) up to 4 keV at \( n_e \approx 1 - 2 \cdot 10^{19} \text{ m}^{-3} \) whereas combined NBI/ECRH discharges at high density \( T_e \approx T_i > 1 \text{ keV at} \ n_e \approx 10^{20} \text{ m}^{-3} \) lead to high performance. Depending on the achieved temperatures, the experimental transport analysis in the plasma core is consistent with the neoclassical predictions. The experimentally observed "electron root" feature with strong \( E_r > 0 \) is driven by the ECRH at high power levels. The neoclassical prediction of a purely thermal "electron root" is not supported experimentally. The neoclassic ordering scheme is violated for ions in case of very strong \( E_r \).

1 Introduction

High temperatures are most important for the transport analysis in the stellarator long-mean-free-path (lmfp) regime. For these conditions, neoclassical theory predicts an unfavourable dependence of the transport coefficients on temperature. Thus, the transport analysis in low-collisionality plasmas is best suited for the examination of the neoclassical predictions. With respect to future large stellarators, neoclassical theory seems to be a fairly reliable tool which may be used for predictive transport codes.

Close to the plasma edge at low temperatures, neoclassical theory fails. The confinement in this region is dominated by anomalous transport (i.e., the physics are not yet understood). Especially in W7-AS with the fairly low vacuum shear, the confinement properties depend sensitively both on the value of the rotational transform (low order rational values of \( \epsilon \) can lead to confinement degradation)
and on the shear, \( dt/dr \) (the shear in the \( E \times B \) rotation may also play a role). However, this paper concentrates on the confinement properties in the bulk plasma for optimum conditions.

2 Optimum confinement discharges

In high power NBI/ECRH discharges, a narrow density profile allows steep temperature gradients (and large \( E_r < 0 \)) close to the plasma edge where \( n_e \) becomes very small. For this type of discharges, optimum confinement properties are found (\( T_e > T_i \simeq 1.5 \) keV) with \( \tau_E \) exceeding the ISS95 scaling [1] by at least a factor of 2. The experimental particle fluxes as well as the ion and electron energy fluxes are in good agreement with the neoclassical predictions up to 70% of the plasma radius; see Refs. [2]. Furthermore, the predicted \( E_r \) obtained from the ambipolarity condition of the neoclassical fluxes is consistent with the experimental findings [3] also at the outer radii where the ambipolar neoclassical fluxes become very small. These findings indicate that the additional anomalous particle fluxes may be intrinsically ambipolar.

Good wall conditioning and very low recycling are mandatory to obtain the narrow density profiles and to provide global density control even for high NBI power levels (with a particle source strength of up to \( 2.5 \cdot 10^{20} \) /s). Contrary to NBI discharges in the "early phase" of W7-AS operation, the \( n_e \) profiles are fairly narrow, and the edge density is very low. Moreover, the outer density profile is - within the experimental errors - independent of both the central density and the heating power. This very surprising result is shown in Fig. 1. Here, the \( T_e \) gradient at the outer radii reflects the heating power. Furthermore, \( T_i \simeq T_e \) is found in this region [3]. The steep temperature gradients flatten in the bulk region due to the strong temperature dependence of the neoclassical transport coefficients leading to the optimum confinement in this type of discharges.

Also in ECRH discharges at moderate density, the achieved electron temperatures are limited by neoclassical transport in the bulk region. An example with optimized edge confinement at high magnetic shear (\( \epsilon' \)) by a plasma current of 25 kA is shown in Fig. 2; see [4]. The evaluation of the electron heat diffusivity - both the
experimental one from power balance and the neoclassical one - are based on the $T_e$ data from Thomson scattering, the ECE data indicate even better edge confinement (the high plasma current was taken into account in the equilibrium calculations of the magnetic coordinate transformation). With this uncertainty, the electron energy flux is in agreement with the neoclassical prediction within 70% of the plasma radius.

3 ECRH driven "electron root" feature

Strongly positive radial electric fields have been measured at W7-AS in low density discharges at high ECRH power level. The electron temperature profiles are highly peaked (with $T_e(0)$ up to 4 keV), whereas the ion temperature and the density profiles are flat. The finding of the strongly positive $E_r$ is related to an additional peaking of the central $T_e$ profile indicating improved electron energy confinement. The corresponding experimental $\chi_e$ is much lower than the neoclassical one for $E_r \approx 0$. This "electron root" feature at sufficient ECRH power is only found for W7-AS configurations where a significant fraction of the ECRH power at 2nd harmonic X-mode is absorbed by ripple-trapped electrons close to the magnetic axis. Equivalent experiments in a configuration without trapped electrons in the ECRH launching plane show neither these strongly positive $E_r$ nor the additional peaking of the $T_e$ profile. The transient response of the central $T_e$ in case of ECRH power modulation or switch-off strongly indicates that the ECRH driven electron flux is responsible for this "electron root" feature; see Fig. 3. For the configuration with an enhanced minimum of the magnetic field strength in the ECRH launching plane, the central $T_e$ is higher than in the neoclassically improved configuration (with a maximum of B), however, this additional peaking in the central $T_e$ disappears within less than 1 ms after the ECRH is switched off. Furthermore, the temperature decay is significantly faster than is measured for the improved configuration confirming directly the neoclassical prediction.

The evidence, that ripple-trapped suprathermal electrons generated by the ECRH contribute to the anipolarity condition of the particle fluxes, is supported by Monte Carlo simulations [5] (in 5D
phase space). At intermediate radii, slightly positive \( E_r \) ("ion root" for \( T_e \gg T_i \)) are predicted for "purely" thermal particle fluxes which are consistent with the experimental data. Also the experimental heat diffusivity (from power balance) is in good agreement with the neoclassical one. In the region of the "electron root" feature, however, the strongly positive \( E_r \) reduce the thermal neoclassical electron fluxes, but also the ion ones are so strongly decreased that the suprathermal ECRH driven fluxes cannot be balanced. In the conventional neoclassical theory, the ion transport coefficients in the \( l m f p \) regime decrease at least with \( E_r^{-3/2} \). Consequently, this approach must be reviewed for the case of strong \( E_r \).

4 Need of self-consistent neoclassical theory

The strongly positive radial electric field within the region of the "electron root" feature leads to complex effects on the flux surfaces. In the continuity equation, the inhomogeneity of \( B \) with \( \nabla \cdot \mathbf{v}_{\times B} \approx -2 \mathbf{v}_{\times B} \cdot \nabla \ln B \) drives a parallel flow (which is the equivalent of the Pfirsch-Schlüter current) as well as a (1st order) density inhomogeneity, \( n_1 \), which is linked to a (1st order) potential varying on flux surfaces. With respect to the "usual" neoclassical theory, a strong \( E_r \) leads also to modifications. For example, in the DKE [6] the \( E \times B \) drift is assumed to be constant on flux surfaces, and density inhomogeneities are neglected.

In the conventional neoclassical approach, the 0th order distribution function, \( f_0 \), must satisfy \( C(f_0) = 0 \) with \( C \) being the Coulomb collision operator, i.e., \( f_0 \) corresponds to the Maxwellian, \( f_M(r,v^2) \), and the radial derivative drives the neoclassical transport in 1st order. This approach can be generalized to include the transport along the magnetic field lines due to small poloidal and toroidal electric fields on the flux surfaces, \(-\nabla \Phi_1\), and the 0th order DKE is given by

\[
\mathbf{v}_\| \cdot \nabla f_0 + v \frac{\partial f_0}{\partial v} = C(f_0) \quad \text{which yields} \quad f_0 = e^{-\frac{w_1}{E_r}} f_M(r,v^2).
\]

Here, \( \dot{v} \) is only the parallel acceleration term (\( \propto p B \cdot \nabla \Phi_1 \) with the pitch, \( p = v_\| /v \)). The other contribution to \( \dot{v} \) originating from the radial \( \nabla B \)-drift in a strong radial electric field (\( \propto (1 + p^2) (B \times \nabla B) \cdot \nabla \Phi_0 \)) violates the isotropic ansatz for \( f_0 \), i.e., \( C(f_0) \neq 0 \), and
must be omitted in 0th order. It follows, that the poloidal \( B \times \nabla \Phi_0 \)
can be treated only in 1st order.

In the 1st order DKE, the terms related to \( \partial f_1/\partial r \) and \( \partial f_1/\partial v \)
are neglected with respect to the equivalent terms in \( f_0 \), and the Coulomb term, \( C \), is approximated by the simple pitch-angle collision operator. Then, \( r \) and \( v \) are only parameters in the 1st order “mono-energetic” DKE. To maintain consistency with the 0th order ansatz, also the term \( \propto (B \times \nabla B) \cdot \nabla \Phi_0 \) must be omitted in the \( \dot{p} \) expression. In this way, \( (B \times \nabla B) \cdot \nabla \Phi_0 \) disappears completely in both orders, i.e., the radial electric field must be sufficiently small. Strong \( \nabla \Phi_0 \) are in conflict with this conventional neoclassical ordering.

\[
\begin{align*}
\left( \frac{B \cdot v}{B} B + \nabla \Phi + \frac{1}{B^2} B \times \nabla \Phi_0 \right) \cdot \nabla f_1 \\
- v \frac{1-p^2}{2} \left( \frac{q}{B} \frac{2}{mv^2} B \cdot \nabla \Phi_1 + \frac{1}{B^2} B \cdot \nabla B \right) \frac{\partial f_1}{\partial p} \\
- C(f_1) = -(\nabla \Phi + \nabla \Phi_{EB} \times B) \dot{f}_0
\end{align*}
\]

Here, \( f_0' \) is the radial derivative (under the constraint of invariant total energy) which drives all the neoclassical transport in combination with the radial components of \( v_{EB} = (mv^2/2q B^3) (1+p^2) B \times \nabla B \) and \( v_{EB} \times B = B \times \nabla \Phi_1 / B^2 \), respectively.

The Fourier expansion (with respect to the toroidal and poloidal angles) of the 1st order DKE leads to mode coupling. In DKEs only the “mirror term” \( B \cdot \nabla B \partial f_1/\partial p \) is taken into account leading to the contribution of the ripple trapped particles which dominates in the \( Lmfp \) regime. In addition to the \( B \cdot \nabla \Phi_1 \partial f_1/\partial p \), also the \( \nabla \Phi_1 \times \nabla f_1 \) term leads to mode coupling (the poloidal component of \( \nabla \Phi \) is of minor importance in this context), and, as a consequence, to a 1st order density and potential, \( n_1 \) and \( \Phi_1 \), respectively. However, \( \nabla \Phi_{EB} \) has to be sufficiently small in this neoclassical ordering. As this \( \nabla \Phi_{EB} \) is partly counteracted by the \( v_{v_\parallel} \) term, this effect will be much stronger in the ion drift-kinetic equation, whereas the 1st order potential affects both equations. Finally, the radial electric field is determined from the ambipolarity condition of the particle fluxes and \( \Phi_1 \) from quasi-neutrality condition \( (n^e = n^i) \) on the flux surfaces.

In principle, this coupled system of drift-kinetic equations can
be solved iteratively. The 1st order mono-energetic DKE is solved with Fourier expansion of $B$, $f_1$ and $\Phi_1$ both for electrons and ions. By energy convolution, the 1st order densities are calculated, and $\Phi_1$ from the Boltzmann factor in 0th order is estimated to satisfy the quasi-neutrality condition. A first attempt of integrating self-consistently this system of DKE’s has already been performed [7].

Although no “convective” particle fluxes are driven by $\Phi_1$ in 0th order (i.e., with $f_0$), the energy flux is directly affected. With respect to $f_1$ (which leads to the conventional neoclassical transport), a strong $E_r$ will modify the particle and energy transport both for electron and ions. In particular, the nearly vanishing ion transport coefficients for strong $E_r$ as obtained from DKES may be substantially affected. Consequently, the prediction of a “purely neoclassical electron root” (typically obtained for low density ECRH discharges in W7-AS) without the drive of fast ripple trapped electrons generated by the ECRH seems to be unreliable.

Finally, this self-consistent approach seems to be essential for describing the impurity transport. The relative effect of the radial electric field scales with $\sqrt{m_i}$, and the impurity transport coefficients in the conventional approach can become much too small. Even in a “tracer” modelling, the 1st order potentials generated by the bulk ions may dominate the impurity transport, and the restriction of sufficiently small $\nabla \Phi_0$ may be satisfied only for the bulk ions. A simple impurity transport modelling was found to be inappropriate with respect to the W7-AS findings [8].

5 Conclusions

Neoclassical theory is confirmed by the experimental transport analysis in low-collisionality W7-AS plasmas, i.e., in the bulk part at sufficiently high temperatures. This conclusion holds for the ion and electron heat conduction as well as for the particle transport. Also the predicted ambipolar electric field is consistent with experimental findings except when the strongly positive “electron root” is predicted. Based on purely thermal neoclassical particle fluxes, the ambipolarity condition typically predicts these strongly positive $E_r$ (“electron root”) in low density ECRH plasmas with highly peaked $T_e$ profiles which is not consistent with the experimental
findings. Only for special conditions, where a significant amount of the ECRH power is absorbed by ripple-trapped electrons, do the ECRH generated suprathermal electron fluxes lead to the strong \( E_r > 0 \). In this sense, an ECRH driven "electron root" feature was found at W7-AS (with \( T_e(0) \) up to 4 keV). However, a basic assumption of the conventional neoclassical theory is violated for the ions in case of these very large \( E_r \). In the case of higher density, the radial electric fields are typically slightly negative allowing for the conventional neoclassical approach.

In a next step, however, a self-consistent neoclassical code, which takes also the density variations on flux surfaces into account, should be developed. Such a code seems to be mandatory to extend the neoclassical theory to the impurity transport in stellarators. So far, the neoclassical predictions for the impurity transport show an extremely sensitive dependence on the radial electric field; poloidal and toroidal electric fields (which are also responsible for the transport in the Pfrisch-Schlüter regime) are not taken into account. Since the impurity transport will likely be a crucial topic for future large stellarators, one should start with a self-consistent neoclassical description.

References


F. Wagner, M. Kick and W7-AS Team, this conference.


Figure Captions

Fig. 1 Electron density and temperature profiles for optimum confinement conditions at different density and heating power with purely NBI and combined NBI/ECRH (solid line: 450 kW; dashed line: 1650 kW; dot-dashed line: 2100 kW).

Fig. 2 The electron temperature profile (upper plot) and the electron heat diffusivity for an ECRH discharge (∼400 kW at 140 GHz 2nd harmonic X-mode) with optimum confinement (25 kA plasma current) at moderate density ($n_e(0) \approx 4 \cdot 10^{19}$ t$^{-3}$): • from Thomson scattering and ◼ from ECE; The experimental $\chi_e$ (from power balance, dot-dashed line) and the neoclassical predictions (lower plot, the dotted line is for $E_r = 0$).

Fig. 3 Time traces of central ECE temperatures after the ECRH is switched off for equivalent discharges in different magnetic configurations with ripple-trapped particles in the ECRH launching plane (solid lines) and without (dashed lines).
Fig. 1

Fig. 2

Fig. 3
Alfvén Instabilities in WENDELSTEIN 7-AS

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Abstract

Global Alfvén eigenmodes excited by resonant neutral beam injected fast ions cause a pronounced MHD activity in W7-AS. The characteristics of these modes in different parameter regimes are described. In particular, improved reconstructions of mode structures have been obtained by X-ray tomography. The effect of increased magnetic shear on the Alfvén spectrum and fast particle effects will be discussed.

Keywords:
Alfvén Eigenmodes, neutral beam injection, fast particles, X-ray tomography, MHD-activity, magnetic shear

1 Introduction

The excitation of Alfvén Eigenmodes (AE) by energetic particles is an important issue because the associated magnetic perturbations may cause increased losses of resonant particles, in particular of fusion born alpha particles in a fusion reactor. Such losses can lead to a reduced margin for ignition and also to excessive heat loads on the vessel wall. First observations of alpha-particle-driven toroidal Alfvén eigenmodes (TAE) were made in the TFTR tokamak [1]. Most of the experimental investigations, however, refer to plasmas with energetic particle populations, which originate from the heating by neutral beam
injection (NBI) or ion cyclotron resonance heating (ICRH). The experimental determination of AE frequencies and damping rates has been successfully achieved in the JET tokamak by using antenna excitation [2]. Considerable progress has been made to describe the present experiments by theoretical models thus providing a more reliable basis for predictions of the AE stability in future machines such as ITER. The magnetic configuration in W7-AS is characterized by non-axisymmetry, very low shear and large aspect ratio, and therefore, the Alfvén spectrum is expected to differ with respect to a typical tokamak case. The observations in W7-AS can potentially contribute to a broader understanding of Alfvén instabilities in toroidal systems.

During neutral beam injection in W7-AS pronounced MHD activity is observed, which is attributed to global Alfvén eigenmodes resonating with ions of the slowing down distribution [3]. In most cases net-current-free plasmas are investigated. Two almost tangential beamlines injecting in co-and counter- direction, each with a power of up to about 1.5 MW at 50-55 kV, into target plasmas generated by electron cyclotron resonance heating (ECRH) at magnetic fields of 1.25 T and 2.5 T or by a 900 MHz generator at arbitrary field, respectively. Usually hydrogen is injected into a deuterium target plasma (H → D), but D → D injection has been used occasionnally in order to study fast particle effects by monitoring the beam-target neutron rates.

2 Low Frequency GAE Modes

The rotational transform at the plasma edge is typically slightly above the main resonances at $\epsilon = 1/3, 1/2$, where the confinement is good due to the reduced number of possible resonances [4]. Therefore, the shear Alfvén continuous spectra defined by the simple dispersion relation $\omega^2 = (k_\parallel \cdot v_A)^2$, with $k_\parallel = (m+n)/R$, $v_A$ Alfvén speed, $\epsilon = 1/q$ rotational transform, poloidal and toroidal mode numbers m,n do not extend to the zero frequency limit since $k_\parallel$ remains finite. In the gaps below the Alfvén continua weakly damped global Alfvén eigenmodes (GAE) of both helicities, $n/m > 0$ and $n/m < 0$, can exist. In the first case the mode helicity is in the same direction as of the
equilibrium field, and very low frequency GAEs in the range of 10...40 kHz are typically found. The mode numbers correspond to low order values of m and n with the ratio n/m closest to the value of the rotational transform at the location of the mode. Toroidal coupling between modes of same toroidal mode number n and adjacent poloidal mode numbers m and m±1 does not play a role due to low shear, and therefore, TAE gaps and TAE modes as in tokamaks are not present at least in case of low m and n. Low frequency GAEs are mostly relatively coherent and appear as continuous oscillations on signals of various diagnostics. The particle drive is inferred from transient mode behaviour at the time, when a neutral injector is switched off. The decay of the mode activity corresponds to the slowing down of fast particles below the resonance velocity, and is therefore much faster than the decay of the plasma pressure. The GAE frequencies are still much larger than those of pressure driven modes or tearing modes in case of ohmic current drive (OH). The GAE modes show the characteristics of waves propagating with a real frequency in the plasma rest system. The propagation is in the direction of the fast ion diamagnetic drift (opposite to the case of the other modes), and this is in agreement with the excitation mechanism, where free particle energy is tapped from the spatial fast ion gradient. The condition for this to overcome the fast particle velocity space damping is approximately given by $\omega \approx \omega_{\text{fast}}/\omega_{\text{GAE}} > 1$ [5].

New data on the spatial structure of GAE modes have been obtained with a 10-camera miniature soft X-ray system (MiniSoX) [6]. This system with a total number of 320 channels is mounted inside the vacuum vessel and provides tomographic reconstructions of the soft X-ray emissivity with good radial and poloidal resolution. Different regularisation schemes including an advanced maximum entropy method [7] and additional analysis methods such as singular value decomposition (SVD) were applied to reconstruct equilibrium profiles and mode structures. The mode structures usually extend over a large fraction of the plasma radius. In case of single peak spectra, the mode structures are consistent with the lowest mode numbers to be expected. Frequently, however, two or more peaks appear in the Alfvén
spectrum. Two effects have been found to cause multiple frequency peaks: - firstly, modes of different toroidal mode numbers but same pitch n/m eg. \((m,n) = (3,1), \ (6,2), \ (9,3)\), - secondly, modes of same poloidal and toroidal mode numbers but different numbers of nodes in their radial eigenfunction. The second effect shows the global waveguide mode structure of the GAEs.

In many cases, particularly at high magnetic field, the velocity of the injected ions does not reach the Alfvén speed. However, excitation of GAEs can occur through \(m \pm 1\) sideband resonances because of toroidal coupling and particle drift effects. The sideband resonance, which is at \(v_A/3\) in the TAE case, can be as low as \(v_A/10\) for GAEs because \(k_{||}\) can be different by this factor for \(m\) and \(m \pm 1\).

In quasi-stationary discharges coherent GAE activity does not seem to cause significant losses. Typical saturation amplitudes of the magnetic perturbations are about \(\dot{B}/B \leq 10^{-4}\), which seem to be in a subcritical range. Only in transient phases at the start of NBI, where the velocity distribution can be more unstable, plasma energy losses in combination with larger mode amplitudes have been found.

### 3 GAE Activity at higher Frequencies

Pronounced but less coherent and sometimes bursting GAE mode activity in the higher frequency range 100 - 500 kHz has been occasionally observed. The frequency spectrum in this range is more complex and typically contains several peaks and also broader features. The frequencies are consistent with GAE modes of higher poloidal mode numbers with typically \(m = 5 - 8\) for \(n = 1\) or even higher in case of \(n > 1\). Additional gap modes such as toroidicity or ellipticity induced AEs (TAE, EAE) can also be present in case of higher mode numbers. Individual mode structures could not be experimentally identified so far, but the frequencies, which scale with the Alfvén speed, are in the range, where strong AE resonances are predicted by the MHD code CASTOR [8-10]. There is some evidence for enhanced transport due to this activity, probably because of multiple resonances. The appearance of the high frequency modes is correlated with the
condition, that the Alfvén speed is below the full energy fast particle velocity. Only under this condition GAEs with higher mode numbers can resonate with the fast ions, since in this case the sideband resonance velocities do not differ very much from \( v_A \). The transition from sideband to fundamental resonance dominated excitation is clearly seen during density ramps, when \( v_A \) drops below \( v_{\text{beam}} \). Simultaneously with the onset of the high frequency AEs the activity at low frequencies becomes weaker indicating that the available free energy is redistributed.

3 Transition from GAE to TAE Modes

In order to investigate the effect of shear on the GAE stability and to show the common physics of GAEs and TAEs shear variation experiments were performed by driving toroidal currents in both directions with the OH transformer. With respect to stability no conclusive result has emerged, since confinement, and therefore, plasma parameters depend on the iota profile. However, evidence for TAE modes at higher shear was found by soft X-ray tomography (fig. 1) in spite of the narrow TAE gap due to large aspect ratio. The observed frequency is consistent with the Alfvén gap structure for \( n = 2 \), and the spatial structure is dominated by \( m = 5 \) and \( m = 6 \) as expected for the \( n = 2 \) TAE. A MHD code with a gyrofluid model for the fast particles [11], which explains the main features of the GAE modes, also gives consistent results for the TAE case (fig. 2).

4 Conclusions

The basic observations of beam driven modes are consistent with global Alfvén eigenmodes. New information about the internal mode structure was obtained from X-ray tomography analysis. As soon as the Alfvén velocity becomes low enough, modes in the higher frequency range are destabilized. With increasing shear a transition of the GAE to TAE type occurs. The effect of Alfvén instabilities on the fast particle confinement is not yet clear. Global Alfvén eigenmodes cannot be suppressed by avoiding rational surfaces, since their eigenfunctions do
not peak there. Therefore they are of potential danger in plasmas with a large fraction of ions in the super-alfvénic range.

References

Figure Captions

Fig. 1 Iota-profiles (left), Alfvén continuum gap structures (middle) and Alfvén mode structures obtained from X-ray tomography (left). In the case of low shear (zero toroidal current) the dominant mode is a \((m,n) = (3,1)\) GAE in the gap below the continuum (top). Increased shear due to OH current drive causes a TAE gap for \(n = 2\), and the mode structure is dominated by \(m = 5\) (inner part) and \(m = 6\) (outer part) consistent with the \(n = 2\) TAE mode (bottom).

Fig. 2 Poloidal Fourier analysis of SX tomogram (top) for \#39042 (fig. 1) yields \(m = 5\) in the inner part and \(m = 6\) further out in qualitative agreement with gyrofluid model calculations (bottom).
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