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Abstract: The local electron and ion heat transport as well as the particle and impurity transport properties in stellarators are reviewed. In this context, neoclassical theory is used as a guideline for the comparison of the experimental results of the quite different confinement concepts. At sufficiently high temperatures depending on the specific magnetic configuration, neoclassical predictions are confirmed by experimental findings. The confinement properties in the LMFP collisionality regime are discussed with respect to the next stellarator generation, for which at higher temperatures the neoclassical transport is expected to become more important.

1. Introduction

The particle and energy transport properties of stellarators and tokamaks in the bulk plasma show some common features, but also significant differences. For the family of quite different stellarator configurations, the main experimental results are briefly reviewed without claim for completeness. The most critical problems of transport in stellarators are discussed with respect to next generation experiments: the Large Helical Device (LHD) [1] being under construction, and W7-X [2, 3], an optimized Helias configuration. These future stellarators, comparable to operating medium size tokamaks, have to demonstrate favourable confinement properties (several keV ion and electron temperatures at densities of up to $10^{20} \text{ m}^{-3}$) essential for demonstrating the reactor potential of the stellarator line.

Present-day stellarators (with major and minor radii up to 2 m and 0.25 m, respectively) have a rather broad spectrum of magnetic configurations [4]. Within the torusatron line, Heliotron-E (Japan) with an aspect ratio of $A \approx 10$ and $N_p = 19$ field periods, the Advanced Toroidal Facility (ATF, USA) with $A \approx 8$ and $N_p = 12$, the Compact Helical System (CHS, Japan) with $A \approx 5$ and $N_p = 8$, and, finally, Uragan-2M (Ukraine, starting operation) with $A \approx 10$ and $N_p = 4$ have moderate or high shear (with the edge value of the rotational transform $\epsilon_\alpha \lesssim 1$ except Heliotron-E with $\epsilon_\alpha \gtrsim 2$). In all these torsatron configurations, the outer part of the plasma confinement region is characterized by rather strong shear and by a magnetic hill. The classical L–2 stellarator (Russia, $A \approx 8$, $N_p = 14$) with moderate shear, but very large helical ripple is complementary to the nearly shearless W7-A (Germany, $A \approx 20$, $N_p = 5$, the predecessor of W7–AS) with very small field ripples. The W7–AS stellarator (Germany, $A \approx 11$, $N_p = 5$) with a full modular coil system is partly optimized with respect to reduced Pfirsch–Schlüter currents (reduction in the Shafranov-shift). W7–AS is also nearly shearless, but has a very complex magnetic field structure (a broad Fourier spectrum of $B$). Finally, the Helias (stellarators with a helical magnetic axis) II–1 (Australia, started operation) and TJ-III (Spain, under construction) with $A \approx 5$, $N_p = 3$ and $A \approx 7$, $N_p = 4$, respectively, complete the stellarator family.

In all stellarators, the magnetic field strength, $B$, shows typically a rather strong modulation along field lines. The radial motion ($\nabla B$-drift) of particles trapped in these ripples can
affect significantly the confinement properties in the long mean free path (LMFP) regime. In an axisymmetric tokamak, the average radial motion of the “banana” particles cancels, and the neoclassical transport coefficients in the LMFP regime decrease with the collisionality, $\nu^*$, and become negligible. In stellarators, however, the neoclassical transport coefficients increase with $\frac{1}{\nu}$ for decreasing collisionality. A sufficiently large poloidal $E \times B$-rotation can average out the radial $\nabla B$-drift of ripple trapped particles. In the so called $\frac{1}{\nu}$-regime, the neoclassical transport coefficients decrease significantly with the radial electric field, $E_r$, but this reduction is much stronger for the ions (in tokamaks, there is no $E_r$ dependence in the neoclassical transport coefficients as long as the “banana” orbits exist). The dependence of the neoclassical transport coefficients on $E_r$ can lead to multiple roots of the ambipolarity condition of the particle fluxes [5, 6]. Deep in the LMFP regime, the stellarator-specific “electron root” with large $E_r > 0$ reduces the neoclassical transport coefficients much stronger than the “ion root” with typically $E_r < 0$ (for $T_i \ll T_e$, also $E_r > 0$ is possible). Consequently, the radial electric field plays a very important role for the neoclassical transport in stellarators.

With respect to the reactor potential of stellarator configurations, the following aspects are considered to be the most important ones. Firstly, the neoclassical ion heat flux, $Q^\text{neo}_i$, scales for $E_r = 0$ as $Q^\text{neo}_i \propto T_i^{9/2}$ in the LMFP regime. As the worst case, this unfavourable $T_i$-dependence may limit the achievable ion temperatures. Secondly, the temperature gradient drives an outward convective flux which is related to the off-diagonal term in the neoclassical transport matrix. The balancing diffusive inward term may lead to fairly hollow density profiles in the inner region with negligible particle sources. Finally, accumulation of impurities is neoclassically expected for the “ion root” with $E_r < 0$ which would not allow steady state reactor conditions, the principal advantage of the stellarator line.

2. Global Energy Confinement

As usual, the description of confinement properties has to start with the analysis of global energy confinement. Two types of energy confinement time scalings are available: theory based scalings as, e.g., the Lackner-Gottardi scaling ($\tau_E^{L-G}$) [7] which is based on the neoclassical ion plateau diffusivities or the gyro-reduced Bohm scaling ($\tau_E^{gB}$, e.g., [8]) where a local turbulence model is assumed, and power laws obtained from statistical regression of experimental data as, e.g., the so-called LHD-scaling ($\tau_E^{LHD}$) [9]. The $\tau_E^{L-G}$, $\tau_E^{gB}$ as well as the $\tau_E^{LHD}$ scalings are rather similar, and agree with the experimental $\tau_E$ of both stellarators and tokamaks as is shown in Figure 1 for the Lackner-Gottardi scaling. For the different stellarator lines, the global energy confinement is quite similar. The common features of these stellarator-specific $\tau_E$-scalings are the improvement of confinement with volume (roughly with $a^2R$), with $B$ and with averaged density, $\bar{n_e}$, and the degradation with heating power, $P$. There is no indication of a dependence on the heating method (e.g., [10]).

Replacing the positive $I_p$-dependence (plasma current) in the regressive tokamak $\tau_E$-scalings by $B$ and by the edge value of the rotational transform, $\epsilon_a$, yields much less agreement with the stellarator data. The clear and favourable $\bar{n_e}$-dependence in the stellarator data is not observed in the regressive $\tau_E$-scalings (with strong $I_p$) of tokamaks, e.g. in the “Goldston L-mode” scaling. No clear indication for an isotope effect is found in stellarators [11].

The nearly shearless stellarators W7–A as well as W7–AS show improving confinement with $\epsilon_a$ as a general trend. Close to low order rational values of $\epsilon_a$ optimum confinement is obtained [43]. In the intermediate $\epsilon_a$ regions, rather strong degradations in $\tau_E$ are found. This effect, attributed to the magnetic configuration (see the next Section) is not observed
Figure 1: Experimental energy confinement times of stellarators [3; data from ATF, CHS, Heliotron-E, W7-A and W7-AS] and tokamaks (Y; from ITER L-mode database [12]) versus the Lackner-Gottardi scaling [7]. For this comparison, the plasma current in tokamaks is replaced by $B$ and $\epsilon(a)$, and $\epsilon(\frac{3}{2}a)$ is used for the stellarators.

In stellarators with moderate or high shear. Due to the small plasma size in the present stellarators, the global confinement can be affected by edge effects (e.g. at low order rational values of $\epsilon_a$).

3. Electron Heat Transport

ECRH discharges at low or moderate densities, for which the collisional power transfer, $P_{ei}$, allows the separation of the electron and ion power balances, are best candidates for the analysis of the local electron heat transport. The ECRH power deposition is highly localized, and the experimental estimates of the deposition profile (from $T_e(r,t)$ analysis after the ECRH is switched off) are in reasonable agreement with ray-tracing calculations (but, typically, only 50% to 80% of the RF input power is found by this analysis). With the purely diffusive ansatz for the radial electron heat flux density, $q_e = -n_e \chi_e^{\exp} T_e'$, the experimental electron heat diffusivity, $\chi_e^{\exp}$, is derived from the local power balance analysis. These $\chi_e^{\exp}$ values are in reasonable agreement with those obtained from analysis of the heat wave propagation stimulated by ECRH power modulation [14]. In the first part of this section, the so-called “anomalous” electron heat transport in stellarators is summarized which is typically found at lower and moderate $T_e$. The second part deals with the neoclassical effects, found experimentally in case of higher $T_e$ values (within the very LMFP regime).

“Anomalous” electron heat transport:

In stellarators, the $T_e$ profiles reflect the power deposition profiles. “Profile resilience” effects as found e.g. in the D III-D tokamak for off-axis ECRH [15] are not observed. An example for off-axis ECRH at moderate density measured at W7-AS is shown in Figure 2. The hollow $T_e$ profile (due to collisional power loss to the ions, $P_{\text{rad}}$ is of minor importance) leads to a $\chi_e^{\exp}$ profile which is compatible, at least within the errors, to that obtained for the equivalent discharge with central ECRH power deposition showing peaked $T_e(r)$ and slightly hollow $n_e(r)$. However, rather small uncertainties in the off-axis power deposition profile result in large errors for the central $\chi_e^{\exp}$. In the example of Fig. 2, $\chi_e^{\exp}$ is much higher than the neoclassical prediction. Contrary to the D III-D results, where a strong inward heat pinch is postulated [15, 16], no indication for a convective term in the electron heat transport is found. This result is supported by analysis of heat waves for off-axis power modulation [17].

In tokamaks, the electron heat transport seems to be linked to the current density profile (e.g., [18]). In stellarators with only rather small toroidal net currents (a few kA), the effect
of the internal shear on local transport is also not fully clear. In W7-AS at low order rational values in the $\epsilon(r)$ profile for low internal shear, confinement degradation is indicated by a local flattening in the $T_e$ profile [19]. At higher internal shear (generated by ECCD, bootstrap and ohmic currents), however, this local confinement degradation disappears. In general, localized mode activity and fluctuations are found at low order rational values in $\epsilon(r)$ [19, 20].

The effect of external field errors on the magnetic flux surfaces [21] may lead to island formation [22, 24], or, at high shear, to field ergodization. This picture reflects the optimum confinement properties of W7-A as well as of W7-AS close to $\epsilon_a \simeq \frac{1}{3}$ and $\frac{1}{2}$, where only high order rational values of $\epsilon$ are present. The hypothesis, that the magnetic field in a high shear region can be ergodized (e.g. due to small external error fields), and that parallel electron heat transport is responsible for degraded confinement properties, is not directly confirmed in the experiments. From perturbation field experiments in CHS, Heliotron-E and ATF, no conclusive result related to the ergodization effect is obtained (see, e.g., [23]). Furthermore, in the high shear region in torsatrons close to the plasma edge, where the magnetic hill is present, the electron heat transport may also be attributed to unstable resistive interchange modes [11, 25].

As in tokamaks, the “anomalous” electron heat transport is usually attributed to mode activity and turbulence, but a consistent (at least to some extent) theoretical understanding is missing so far, see, e.g., Ref. 26. Regarding the full 3D and rather complex field topology in stellarators, the theoretical modelling of fluctuations, turbulence and related transport is still at its beginnings. Nevertheless, some general statements seem to be justified. In the torsatrons, an external outward shift of the plasma column or the Shafranov shift of the magnetic axis deepen the magnetic well and increase the radius where the magnetic hill appears. The more favourable MHD properties of these magnetic configurations with respect to global mode activity have been confirmed by Heliotron-E, CHS as well as ATF (e.g., [28]). The better confinement, however, is typically found for the opposite case, i.e. for an inward shift of the magnetic axis. Consequently, the observed low $n/m$ modes do not determine the local electron heat transport. For ATF [11], the dissipative trapped electron mode (DTEM) is discussed to be responsible for the observed $\chi_e^{\text{exp}}$. The fraction of trapped particles (controlled externally by the dipole and quadrupole moments of the configuration), however, affects only weakly the global confinement properties [8]. Furthermore, no indication for the DTEM is found in CHS [29]. Ballooning modes are not important for stellarator confinement.
The typically "anomalous" $\chi_e^{\exp}$ obtained from power balance analysis decreases with increasing $n_e$ and $B$, respectively. A slight improvement with $\epsilon$ for optimum confinement is found in W7-AS [30]. $\chi_e^{\exp}$ increases with heating power, $P$. An ECRH power scan (within a factor of 5) shows only a rather small increase in the electron temperature gradient [19]. So far, no regression of $\chi_e^{\exp}(r)$ with the power dependence being replaced by the local plasma parameters $T_e$ and $T_e^\ast$ is available. The dependence of $\chi_e^{\exp}(r)$ on $T_e(r)$ as well as $T_e^\ast(r)$, which are most relevant to get a feeling for the "anomalous" transport drive, are not yet clear. An other approach is the comparison of dimensionally similar discharges at different $B$. In stellarators, the $\chi_e^{\exp}$ profiles are found to be more close to "gyro-Bohm-like" scaling rather than to "Bohm-like" scaling [31]. These findings are different to the JET [32] and TFTR [33] results. The "gyro-Bohm-like" transport scaling may be attributed to radically extended turbulence with correlation lengths of the order of the ion gyro-radius. E.g., the neoclassical transport coefficients, however, are also "gyro-Bohm-like". So far, only preliminary conclusions can be drawn from all these kinds of analysis.

**Neoclassical electron heat transport:**

For analyzing the neoclassical contribution in the electron heat transport, rather solid theoretical tools (compared to the previous part) are available. In low density ECRH discharges with sufficiently high $T_e$, the neoclassical heat diffusivity, $\chi_e^{\text{neo}} \propto T_e^{-7/2}$ in the $1/\nu$-regime, exceeds the "anomalous" one in the inner part of the plasma. By comparing equivalent discharges for W7-A, L-2 and W7-AS, the configurational dependence of $\chi_e^{\exp} \approx \chi_e^{\text{neo}}$ could be confirmed [34]. The rather poor confinement in L-2 is related to the very large helical ripple, whereas optimum confinement was obtained for W7-A, with the very small field ripples. For electron collisionalities full in the $1/\nu$-regime and with $T_e \gg T_i$, rather strong radial electric fields $E_r > 0$ ("electron root") are predicted for the inner radii. The predicted improved confinement for this "electron root", however, is in contradiction to the experimental power and particle balance. Also at Heliotron-E [35] and CHS [29], neoclassical heat diffusivities are found for equivalent conditions. The $E_r$ values obtained from the measured poloidal plasma rotation are found to be positive as expected, but by a factor of 2 or 3 less than the predicted values of the "electron root" [36]. All these predictions of $E_r$ are mainly based on the ambipolarity condition of the particle fluxes, even if viscosity effects close to the shear layer in the poloidal rotation (transition from "electron" to "ion root", comp. Fig. 5) are taken into account [37]. Possible non-ambipolar contributions (e.g. due to magnetic fluctuations or turbulence) have not been considered.

These results are very significant with respect to the next generation stellarators where at higher $T_e$, but equivalent collisionalities, sufficiently low electron heat transport has to be demonstrated. As a consequence, attention must be payed to the reduction of the effective helical ripple by optimization of the magnetic configuration. As long as the strong "electron root", which reduces also the neoclassical electron transport coefficients, is not confirmed experimentally, the unfavourable neoclassical electron heat flux ($q_e \propto T_e^{7/2}$ in the $1/\nu$-regime) is expected to exceed significantly the "anomalous" heat transport found in present day tokamaks. However, the neoclassical confinement itself can be essentially improved for an optimized magnetic field configuration [3, 38].

4. **Ion Heat Transport**

At high densities, where the collisional power transfer, $P_{ei} \propto n_e^2(T_e - T_i)/T_e^{7/2}$, causing a strong coupling, the electron and ion heat loss channels cannot be separated based on mea-
ured $T_e$ and $T_i$ profiles. Nevertheless, the ion channel seems to play a main role in the power balance. At low densities, $P_{ei}$ becomes small, and only direct ion heating can lead to higher $T_i$. For NBI sustained discharges at low densities (shortly after the ECRH target phase), rather high ion temperatures ($>1$ keV) have been observed [10, 11]. In ATF (at $n_e \leq 10^{19}$ m$^{-3}$), this high $T_i$-phase immediately after switch-on of the NBI [39] is much shorter than the slowing-down time. These findings may be related to an unstable fast ion distribution leading to preferential ion heating [21] and to rather strong radial electric fields (due to fast ion orbit losses) improving significantly the ion confinement. For this kind of discharges, both the estimation of the degradation in the NBI heating efficiency due to CX losses (which may be essential at low $n_e$) and an ion power balance analysis are not available.

In NBI discharges at moderate $n_e$ ($<10^{20}$ m$^{-3}$), high ion temperatures ($T_e < T_i$ up to 1 keV) were obtained at W7-A [40, 21]. This very good ion energy confinement is attributed to strong $E_r < 0$ generated by fast ion orbit losses of the nearly perpendicular NBI [41, 42]. Although the thermal ions are within the plateau collisionality regime, their neoclassical transport coefficients are significantly reduced by these very large $E_r$ [43, 41]. For this special case, the fast ion orbit losses affected significantly the particle balance (ambipolarity condition). For increased plasma radius as well as for mainly parallel injection, however, fast ion orbit losses are of minor importance. As a consequence, this method to obtain significantly improved thermal ion heat confinement has no prospects.

Figure 3: Density, electron and ion temperature as well as the radial electric field profiles of NBI heated discharges at CHS [44]. The two analytical predictions for $E_r$: solid lines for high $n_e$ (•) and dashed lines for low $n_e$ (O).

For the comparison of the neoclassical ion heat transport with the experimental results, the radial electric field plays an essential role. However, as the $E_r$ profiles are not available in general, neoclassical predictions based on the ambipolarity condition of the particle fluxes are used. Experimental $E_r$ profiles derived from the poloidal plasma rotation measured at CHS [44] are shown in Figure 3. These $E_r$ profiles agree rather well (except for the outer radii) with the neoclassical predictions based on analytical models. For the rather low ion and electron temperatures of this example, the predicted "ion root" is basically confirmed. Nevertheless, the possible effect of non-ambipolar contributions to the particle fluxes can play a role close to the edge (e.g., direct ion loss cone effects).

At W7-AS, the ion power balance is analyzed [45] for both ECRH (with 140 GHz) and NBI discharges (and for combined heating) at moderate densities ($n_e < 10^{20}$ m$^{-3}$). For both types of discharges, the ion heat flux, $Q_i^{\exp}$ (estimated from the power balance within 70% of the plasma radius, CX-losses are here negligible), is found to exceed the neoclassical prediction, $Q_i^{\text{neo}}$, with the ambipolar $E_r$ included by a factor of 2 to 3, but to stay well below $Q_i^{\text{neo}}$ with $E_r = 0$. As an example, the ion heat flux and the neoclassical prediction are shown in Figure 4 for an ECRH (140 GHz) discharge at optimum confinement. The full
Figure 4: Profiles of electron and ion temperature (on the left) and density (in the center) of an on-axis ECRH (140 GHz, 2nd harmonic) discharge at W7-AS (after an "H-mode transition"). The additional \( T_i(r) \) (dotted line) is the neoclassical prediction (based on DKES code): the self-consistent calculation of both \( T_i \) and \( E_r \). In the region where CX-losses are negligible (\( r < 12 \text{ cm} \)), \( Q_{i,\text{exp}} \) from power balance (dashed line, based on the measured \( T_i \)) exceeds \( Q_{i,\text{neo}} \) (dotted line) by a factor of 2 to 3 (on the right).

\( T_i \)-prediction based on neoclassical transport coefficients calculated by the DKES code [46, 47] with the self-consistent \( E_r \) (here, only the "ion root" is found) is also given. At lower \( n_e \) and higher \( T_i \) (lower \( \beta_i^* \)), the effect of \( E_r \) on the neoclassical \( T_i(r) \) and \( Q_i(r) \) is more pronounced.

The detailed analysis shows clearly that "convective" terms (non-diagonal elements in the transport matrix being related to both the density gradient and \( E_r \)) are important: the comparison of only \( \chi_{i,\text{exp}} \) (a purely diffusive ansatz for the ion heat transport) with \( \chi_{i,\text{neo}} \) is not sufficient.

In high power NBI discharges in W7-AS, \( \tau_p \gg \tau_E \) is indicated, and \( n_e \) increases with time. Without efficient density control, the ion collisionality increases within the plateau regime \( (Q_{i,\text{neo}} \propto n_i T_i^{5/2}) \). With efficient density control (low recycling and additional ECRH, see Sec. 5), stationary conditions are obtained in the transition to the LMFP regime \( (n_i \simeq 6 \cdot 10^{19} \text{ m}^{-3} \text{ and the central } T_i \leq 800 \text{ eV}) \). Here, the ion heat loss is a main part in the total power balance. At lower collisionality, good neoclassical confinement at higher \( T_i \) is only predicted for the "electron root" (at rather large \( E_r > 0 \)) with \( Q_{i,\text{neo}} \propto n_i^2 T_i^{3/2} \) whereas for the "ion root" \( Q_{i,\text{neo}} \propto T_i^{5/2} \) is expected. Consequently, access to significantly higher ion temperatures should depend on the realization of the "electron root". However, the transition from the "ion" to the "electron root" is predicted only for sufficiently low collisionalities deep in the LMFP regime. \( E_r > 0 \) was found experimentally in Heliotron-E [35] and CHS [48] only in low density ECRH discharges. As a consequence, the existence of the "electron root" and the associated confinement improvement has still to be experimentally confirmed for higher temperatures and densities.

5. Particle Transport

In general, the density profiles in stellarators and tokamaks are quite different. In stellarators, the density peaking due to an inward convective term is typically not observed. The inner part of the \( n_e \) profiles (about 60% to 80% of the plasma radius) is only slightly peaked (typically for NBI) or flat, and for ECRH with central deposition even slightly hollow. An example for an ECRH discharge at W7-AS is shown in Figure 5 (see also [35]). In this discharge, most of the ECRH power was deposited off-axis (at about 8 cm, only 30% highly peaked central deposition). Contrary to the purely off-axis case in Figure 2, the ECRH off-axis
Figure 5: Profiles of electron and ion (neoclassical prediction) temperature (on the left) as well as density (in the center) of an ECRH (70 GHz, fundamental o-mode) discharge at W7-AS with on- and off-axis power deposition (profile from ray-tracing is indicated). In the inner region, the particle flux, $\Gamma^{\text{exp}}$ (on the right), estimated by using DEGAS code (dot-dashed line) is in reasonable agreement with the neoclassical (mainly ambipolar) prediction (solid line). At the position of the shear layer in the poloidal plasma rotation ($r \approx 5$ cm), the transition from the inner “electron” to “ion root” at outer radii is related to a small (numerical) viscous flux. 

deposition is not visible in the $T_e$ profile. The 3D distribution of the neutrals is calculated by the DEGAS code [49, 50] (calibrated to $H_\alpha$-measurements). The averaged particle fluxes, $\Gamma^{\text{exp}}$, are estimated and compared with the ambipolar neoclassical predictions, $\Gamma^{\text{neo}}$, obtained by using DKES code [46, 47] (s. Fig. 5, on the right). In the region of the hollow density profile, both estimates agree quite well. $\Gamma^{\text{neo}}$ is mainly driven by the “convective” term (the non-diagonal term in the transport matrix $\propto T_e$). Reasonable agreement of $\Gamma^{\text{exp}}$ and $\Gamma^{\text{neo}}$ is also obtained for low $n_e$ ECRH as well as for high $n_e$ NBI heated discharges in W7-AS [50]. As a consequence, the particle fluxes in the inner part of the plasma are consistent with the neoclassical prediction if the ambipolar radial electric field is taken into account.

At the outer radii, the particle transport analysis based on the neutral gas distribution calculated with DEGAS code leads to “anomalous” particle fluxes (being much larger than $\Gamma^{\text{neo}}$). Here the main particle sources (from recycling and from gas puffing) and rather strong density gradients are located. Contrary to tokamaks where an inward particle pinch is required for peaking of the $n_e$ profiles, a purely diffusive ansatz for the particle flux, $\Gamma^{\text{exp}} = -D_p n_e^\prime$, seems to be sufficient within the density gradient region [50]. In this region, both the electron heat diffusivity $\chi_e^{\text{exp}}$ and $D_p$ show a strong increase towards the plasma edge. Analogous to $\chi_e^{\text{exp}}$, $D_p$ decreases with increasing $B$, $n_e$ and $\epsilon_a$, and increases with heating power [19]. For W7-AS, Figure 6 shows $D_p$ and $\chi_e^{\text{exp}}$ from power balance (both at 80% of the plasma radius) as well as the impurity diffusivities, $D_I$ (from AI ablation), versus density. All these transport coefficients scale roughly with $n_e^{-1}$. The rather large difference in $\chi_e^{\text{exp}}$ for the two $\epsilon_a$ values is mainly attributed to the different heating power ($P$ is 3 times larger for $\epsilon_a = 0.345$ than for $\epsilon_a = 0.5$), but only a weak effect of $P$ on $D_I$ is indicated. Discharges at equivalent $P$ show clearly a moderate reduction in $D_p$, $D_I$ and $\chi_e^{\text{exp}}$ with $\epsilon_a$.

With respect to the transport analysis in the stellarator scrape-off layer (SOL), the experimental data base is relatively small compared to tokamaks. Furthermore, due to the very complex stellarator field geometry, theoretical transport modelling in the plasma edge is still in the beginnings. The general findings indicate “anomalous” transport with basic similarities to tokamaks. The observed edge fluctuations (electrostatic turbulence) measured by Langmuir probes are related to particle fluxes, $\Gamma \propto \langle \dot{n}_e \phi \rangle$, which are roughly consistent (ignoring possible poloidal and toroidal asymmetries in the SOL) with those from the global
particle balance. In W7-AS, SOL Langmuir probe data are analyzed by a 1D radial “flux tube” approach [52] using particle sources (by DEGAS code) and parallel averaging. The $D_p$ obtained within the SOL decrease towards inner radii matching quite well to the inner $D_p(r)$ [53] as obtained by DEGAS code simulations. A scaling of $D_p$ in the SOL inversely to $n_e(r)$ is indicated corresponding to the $D_p$ in the bulk plasma, see Fig. 6. In W7-AS, there is no indication for different transport mechanisms in the density gradient region and in the SOL.

The plasma turbulence at the plasma edge may be related (see, e.g., [54]) to atomic physics processes (radiation, ionization and charge exchange) in combination with non-linear mode coupling [55]. This rather complex drive mechanism is a common feature in stellarators and tokamaks. Inside the plasma, the fluctuation level decreases monotonically. By ECRH power scans, the $H_\alpha$-signals as well as the global particle confinement clearly indicate a power dependence. With the assumptions that the main turbulence drive is located close to the SOL, and that the turbulence level decays smoothly within the $n_e$-gradient region, the suggestion of $D_p$ as well as $D_f$ being dependent on $P$ is equivalent to the assumption that the heating power affects the turbulence driving mechanism. It seems to be unlikely that this link is associated with the variation of the edge $T_e$ values [56]. Furthermore, within the $n_e^*$ region at ATF, the density fluctuations are attributed to resistive interchange turbulence related to the magnetic hill [57] (opposite to the magnetic well in W7-AS). This mechanism is mainly related to the local plasma parameters.

In high power NBI heated discharges at optimum confinement, $n_e$ increases with time until the discharges are terminated by radiative collapse. A MHD related “density limit” is unknown in stellarators (up to $n_e \approx 3 \cdot 10^{20}$ m$^{-3}$ has been observed in W7-AS). Applying additional ECRH, both density and impurity control can be obtained [10, 51, 19]. So far, a consistent picture for this so called “ECRH pump-out” is missing. In the additional ECRH phase at W7-A [51], the $T_e$ profile broadened whereas the $n_e$ profile was nearly unchanged. Contrary to the remarkable degradation of the global particle confinement, the energy confinement was mainly unaffected. Al ablation experiments (see next Sec.) clearly indicate an increase of $D_f$ for the combined phase (a factor of about 3 at Heliotron-E [10]) whereas the inward velocity is only weakly affected. With the assumption of electrostatic turbulence ($\Gamma \propto \langle n \phi \rangle$), these findings suggest an equivalent increase of $D_p$ at least within the $n_e^*$ region. In combined ECR and NBI heated discharges in W7-AS, on- or off-axis ECRH deposition can suppress the density increase leading to rather flat $n_e$ and highly or slightly peaked $T_e$ profiles, respectively. After switching-off the ECRH, the $T_e$ profile flattens, and the density starts to increase. These observations show the effect of the $T_e$ profile shaping on the particle
transport properties in the inner part (via a non-diagonal term) and, consequently, on the central density evolution.

The "H-mode transition" found in W7-AS [58, 59] shows a steepening of the density gradient only at the outer radii which corresponds to a shrinking of the $n_e^c$ region (an increase of $T_e'$ is also observed in this outer region). At the inner radii where $n_e(r)$ is flat, the electron temperature gradient, $T_e'$, is nearly unchanged (but at higher $T_e$) leading to power balance $\chi_{\text{exp}}^e$ which are the same within the error bars. Within the "H-mode" phase (see the profiles of Fig. 4), the reduced gas feed as well as the decreased $H_{\alpha}$ signals indicate significantly improved particle confinement within the region of steep $n_e'$ just inside the separatrix. First reflectometry measurements [60] show a significant reduction in the $n_e$ fluctuations being most pronounced at low frequencies (< 10 kHz, corresponding to the long wave lengths) which contribute more efficiently to electrostatic turbulent transport. Assuming $D_p \sim (\tilde{n}_e^2)$, $D_p$ decreases within the $n_e'$ region by a factor of at least 4 which is consistent with the rough estimate from particle balance (reduced gas feed and recycling ($H_{\alpha}$) as well as steepening of $n_e'$; no DEGAS simulation being available).

A partly similar transition is also found at W7-AS under limiter conditions at $\varepsilon_a = 0.32$ (and at lower $\tilde{n}_e$) [19]: $D_p$ decreases by about 50% in the region of the strong density gradient. The "H-mode-like" NBI heated discharge found at CHS [61] where the rotational transform was increased by an additional ohmic current, is also characterized by a decrease of the $H_{\alpha}$ signals and by a steepening of the density as well as both temperature gradients close to the edge. All these findings suggest that improved confinement in the steep $n_e'$ region might be due to similar mechanisms in both stellarators and tokamaks.

6. Impurity Transport

The prediction of the impurity fluxes is also closely linked to the radial electric fields. From general arguments, $E_r$ can be identified with a "thermodynamic force" driving "thermodynamic fluxes". For $E_r < 0$, an inward convective term for the ion and, in particular, for the impurity fluxes ($v_I \propto Z_I E_r/T_i$) is expected. Consequently, neoclassical theory predicts impurity accumulation for strongly negative $E_r$. Contrary to tokamaks [63], no fully consistent formulation of the neoclassical impurity transport in stellarators is available, so far. From the experimental point of view, impurity transport analysis is based on two approaches: the time analysis of impurity line radiation (with different ionization stages) by laser ablation technique, and the simulation of the accumulation phase.

The analysis of laser ablation experiments is often based on extensive simplifications: e.g., the assumption of $D_I \propto \text{const.}$ and $v_I \propto r$ both being the same for all ionization stages. $D_I$ is mainly determined from the increasing line radiation of the highest ionization stages during the inflow phase, whereas the ratio of $D_I$ and $v_I$ by the decay phase. The radiation of low ionization stages, being a convolution of impurity transport, atomic processes and, for the lowest stages, inhomogeneity effects (of the ablated material on flux surfaces), is usually more difficult to simulate. Consequently, these estimates of both $D_I$ and $v_I$ should be interpreted carefully in comparison to the local transport coefficients $D_p$ and $\chi_{\text{exp}}^e$. Time-dependent simulations with full radial resolution (e.g., the SITAR code [62]) need more precise experimental data. In general, however, no profile information of the impurity concentrations at the different ionization stages is available. In a predictive version, these codes are used to test theoretical models, e.g., the neoclassical impurity transport. As a consequence, theoretical predictions can only be cross checked with a limited data base.
In low $n_e$ ECRH discharges, no indications for impurity accumulation are found at W7-A, Heliotron-E, ATF and W7-AS. Typically, the average impurity diffusion coefficient, $D_I$, decreases with increasing $\bar{n}_e$ [10] (see also Fig. 6), and an inward convective term ($v_E < 0$) is not found for these low $\bar{n}_e$. For good confinement properties at high $n_e$, however, the decay-times, $\tau_{\text{dec}}$, of the laser ablated impurity lines become very large, and sometimes, no decay is found within the duration of the discharge. For this type of discharges, also the spectral lines related to wall material (e.g. Fe) increase with time indicating “impurity accumulation”.

For the NBI heated discharges at W7-A where the optimum ion confinement was attributed to the strongly negative $E_r$, strong impurity accumulation was found [62]. Fully time-dependent simulations of the impurity transport using a neoclassical, but axisymmetric (tokamak) model (equivalent to [63]) were performed and compared with the experimental findings. Fairly good agreement was obtained. In this axisymmetric model, $E_r$ is implicitly included by “convective” terms related to the background ions, however, these $E_r$ values are less negative than the measured ones (related to the fast ion loss of the nearly perpendicular NBI).

E.g., after the “H-mode transition” impurity accumulation is also indicated at W7-AS. The neoclassical prediction for $E_r$ (“ion root”) close to the separatrix yields more negative $E_r$ after the transition. This result is mainly related to the steepening of the density gradient. Consequently, an increased inward velocity in the region of steep $n_e$ is neoclassically expected. However, both the reduced particle outflux and wall recycling (see Sec. 5) lead to a decreased impurity influx. Strong ELM activity may affect the impurity transport at outer radii (indicated by the outer SX channels). For stationary conditions in the global plasma parameters, the increasing of SX signals and of Fe lines indicate accumulation in the bulk part. The radiation level starts to increase already in the phase of good confinement obtained by increasing $\bar{n}_e$. This effect is consistent with the requirement of an inward convective term at higher $\bar{n}_e$ in the analysis of Si ablation experiments at Heliotron-E [10] (not so clear in W7-AS, so far). The “H-mode transition” in W7-AS improves mainly the particle confinement just inside the separatrix, and the impurity accumulation becomes more pronounced.

![Figure 7: The decay time, $\tau_{\text{dec}}$, of the SX signal after Si laser ablation and the global energy confinement time, $\tau_E$, versus the radial electric field, $E_r$, measured at about 75% of the plasma radius at Heliotron-E [35]. The $E_r$ was controlled at fixed density ($n_e \approx 10^{19} \text{ m}^{-3}$) within an ECRH power scan.](image)

A very important result is obtained at Heliotron-E [35]. At about 75% of the plasma radius, $E_r$ is deduced from poloidal rotation measurements. At moderate and higher $\bar{n}_e$, the predicted “ion root” with $E_r < 0$ is found. In low density ECRH discharges, positive $E_r$ (possibly the “ion root” for $T_i \ll T_e$) being by a factor of 2 to 3 smaller than the prediction for the “electron
root” [36] (comp. also Sec. III) are obtained. Ablation experiments are performed at fixed
density ($\bar{n}_e \simeq 10^{19} \text{ m}^{-3}$) and with different heating power to control $E_r$. The decay-time,
$\tau_{\text{dec}}$, as well as the global energy confinement time are shown in Figure 7 versus $E_r$. This result
demonstrates the strong effect of $E_r$ on the impurity pinch: as $\tau_{\text{dec}} \simeq \tau_E$, an outward pinch is
indicated for $E_r > 0$. Extrapolating these findings to the high temperature LMFP regime at
higher $\bar{n}_e$ being expected for the next generation stellarators, the problem of impurity control
is connected to the control of the radial electric field: the realization of the “electron root”
(the “ion root” feature with $E_r > 0$ cannot be expected for $T_e \sim T_i$) within the bulk part of
the plasma may prevent impurity accumulation.

7. Summary and Conclusions

Transport in stellarators seems to be less sophisticated than in tokamaks. The density and
temperature profiles in stellarators reflect mainly the particle and heat sources. It is sug-
gested that ion heat transport as well as particle and impurity transport in the bulk part of
the plasma are not so far from the neoclassical predictions, this holds also for the electron
heat transport at high electron temperatures. At lower $T_e$, the experimental electron heat
diffusivity is typically “anomalous”, however, no indication for a “heat pinch” is found. Simi-
lar to tokamaks, this electron heat flux is basically not understood. “Anomalous” particle
transport is found in the outer region with steep density gradients where the main particle
sources are located. In this region, electrostatic type turbulence seems to be responsible for
the particle flux. After “H-mode transitions”, the transport coefficients together with the
density fluctuations are significantly reduced in the region of steep density gradients.

For the next generation stellarators, much higher electron and ion temperatures are ex-
pected. For the “anomalous” electron heat diffusivity, no strong and unfavourable temper-
ature dependence is indicated in the present experiments. The neoclassical transport is ex-
pected to become more important which in the stellarator-specific $1/\nu$-regime (LMFP) shows a
strong and unfavourable temperature dependence. A classical stellarator configuration with
a strong helical ripple will be limited to rather low $T_e$ due to this neoclassical electron heat
loss. Consequently, optimization of the magnetic configuration with respect to the reduction
of the effective helical ripple is essential [2].

To suppress the unfavourable $T_e$ dependence of the neoclassical ion heat flux, operation
under the stellarator-specific “electron root” seems to be desirable. Furthermore, this strongly
positive radial electric field could prevent impurity accumulation as it is confirmed experi-
mentally even for rather small $E_r > 0$. However, as the “electron root” solution of the
ambipolarity condition is only predicted in the LMFP regime, sufficiently high temperatures
are required for the realization of the “electron root”. Consequently, this aspect relies on
optimized magnetic field configurations with sufficiently low neoclassical losses. So far, the
existence of the “electron root” is only predicted theoretically.

The off-diagonal terms in the neoclassical transport connect density and temperature
profiles. For the particle transport, no “anomalous” inward convection is experimentally in-
dicated. As the main density gradients (a region of only a few cm, probably not scaling with
the plasma radius) are located close to the edge, the temperature gradient driven contribu-
tion to the particle flux may result in rather hollow density profiles. As a consequence, also
particle sources within the bulk part of the plasma may be necessary in larger stellarators.
On the other hand, temperature profile shaping seems to be a possible tool for the control of
the density profiles in case of particle sources in the bulk plasma (e.g., for high power NBI
heating).

So far, ion temperatures of only up to 1 keV have been obtained in stellarators. It seems to be unlikely, that the rather low $T_i$ values (in comparison to big tokamaks) are only attributed to the smaller plasma radius. Due to the limited temperatures, the full access into the deep LMFP regime with $T_e \approx T_i$ is hardly possible in the present stellarator experiments. For this regime, however, essential improvements of the neoclassical confinement in the LMFP regime are theoretically predicted for optimized stellarator configurations. Consequently, only the next generation stellarators with significantly increased plasma radius should be able to conclude on the reactor potential of the stellarator line.

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References

CXRS-Measurement and Code-Calculation of Impurity Profiles

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INTRODUCTION.
Charge Exchange Resonance Spectroscopy CXRS is performed for impurity lines of several ionization states, especially for carbon and boron, on the advanced stellarator W7-AS. These lines are excited by means of a neutral beam diagnostic injector. The spatially resolved spectroscopic observation in the visible and near ultraviolet region makes it possible to evaluate density profiles of the impurities without a tomographic calculation or an Abel inversion. The measured impurity profiles are compared with the results of calculations of the IONEQ-Code, which take into account either a neoclassical or an anomalous transport model. Thus, diffusion coefficients and convective velocities for the impurities are evaluated. Discharge scenarios before and after H-mode transitions are compared, and impurity influxes before and after boronization of the vessel wall are considered.

PROFILE MEASUREMENT.
The charge exchange CX lines are excited by a pulsed neutral beam injector with particle energies typically at about 30, 15 and 10 keV for the full, half and one third energy components, respectively. Because the beam is pulsed, background radiation can be subtracted from the CX radiation. The light is guided by means of a turnable mirror device towards a 1.26m spectrometer. Intensity can be recorded between 200 nm and 700 nm by a multi-channel photomultiplier array with 15 channels, the spectral resolution for the whole device is about 0.02 nm / channel.

We start from Thomson scattering data for electron temperature and density profiles, an assumed neutral gas background and the geometry of the experimental arrangement for the first loop of an iterative calculation of the impurity profile. A first-step attenuation profile for the neutral beam is calculated by assuming a model for a first-step impurity profile, because the impurity itself contributes to the beam attenuation, too. That impurity profile model is calculated from the electron density profile.

This first-step beam attenuation profile together with the measured CXRS are used as input data for the calculation of a second-step neutral beam attenuation profile, making possible the evaluation of a second-step impurity profile, and so on. A recursion algorithm makes it possible to calculate the n-step impurity profile, which is considered as satisfactory for the case that numerical convergence and stability for the calculated profiles, step by step, are attained for higher n. The whole problem can be condensed to the solution of the following equation:

\[
\frac{N_{\text{phot}}(l)}{n_k(l) \cdot \text{Vol}(l) \cdot (\sigma_{\text{CX}}(V) \cdot V_{\text{coll}})} = \text{Int}(l) = \sum_{i=1}^{3} \text{Int}(l=0)_i \cdot X \cdot \exp \left( -1 \cdot \int_{-Z_0}^{z_0+1} \left( \sigma_{\text{CX}}(V) \cdot n_{H^+}(z) + \sum_{j=1, j \neq k}^{m} \sigma_{\text{CX}}(V) \cdot n_j(z) + \sigma_{\text{CX}}^e(V) \cdot n_k(z) \right) \, dz \right)
\]

Here, \(N_{\text{phot}}(l)\) stands for the CX light intensity as a function of the coordinate \(l\) along the neutral beam trace, \(\text{Int}(l)\) is the neutral beam intensity, \(n_k(l)\) is the impurity density profile for the ionization state \(k\), \(\text{Vol}(l)\) is a geometric factor, \(V_{\text{coll}}\) is the collision velocity, \(\sigma_{\text{CX}}(V)\) is the CX cross section for the excitation of CX radiation, \(\sigma_{\text{CX}}^e(V)\) is the CX cross section for
the neutral beam attenuation by collisions between neutral hydrogen atoms with impurity ions in the ionization state j. The summation in the first line spreads over the mentioned three energy components for the neutral beam, z is the vertical coordinate with respect to the plasma, the summation in the exponent spreads over the assumed impurities present in the plasma. The exponential factor gives the total neutral beam attenuation, the integral corresponds to a line density multiplied by a cross section. Atomic data for cross sections are taken from Barnett, and Tawara et al. (1). Variations of that procedure consist of variations of the cross sections. Thus, we determine the influence of the absolute values of the cross sections on the final results, see Rehker and Speth (2), to be able to estimate errors stemming from possible uncertainties of the input data. Further variations are done by replacing the left sides of the following expressions by the right sides for the equation given above.

\[ \sum_{i=1}^{3} \sigma_{CX}(V) \rightarrow \sum_{i=1}^{3} \frac{[\sigma_{CX}(V) \cdot V_{coll}]}{V_{beam,i}} \]

\[ \text{replace : } \sum_{i=1}^{3} (\sigma V) \rightarrow \sum_{i=1}^{3} \int_{0}^{\infty} \sigma_{CX}(V_{coll}) \cdot V_{coll} \cdot f(V_{plas}) \cdot d^{3}V_{plas}. \]

Here, \( f(V_{plas}) \) stands for a velocity distribution function, assumed for the plasma particles, the brackets \( <.> \) denote the mean defined by the integral (for monoenergetic beam particles):

\[ <\sigma V> = \int_{0}^{\infty} \sigma(V_{coll}) \cdot V_{coll} \cdot f(V_{plas}) \cdot d^{3}V_{plas}. \]

Thus, the role of differing velocity distributions, and variations of it, can be estimated. The impurity density profiles, evaluated in the described way, are then compared to IONEQ code calculations.

CODE CALCULATIONS.
The code IONEQ can, among other things, calculate impurity density profiles starting from electron density, electron temperature and background neutral gas profiles as input parameters, see Weller et al. (3). The calculation includes either an anomalous model, which needs estimated values for diffusion and convective inward velocity for the impurity under consideration, or a neoanomalous ansatz. The calculated result of IONEQ can be varied by a variation of the input anomalous diffusion coefficient \( D \) for an impurity in ionization state \( k \), the inward convective velocity \( V \) and an exponential parameter \( \alpha \), such that the result of the calculation fits best to the measured CXRS impurity density profiles. The inward convective velocity \( V \) for a impurity species at minor radius \( r \) is given by \( V(r) = (-2D/a) \cdot (r/a)^{\alpha} \). Here, \( a \) is the minor plasma radius. First, a ansatz for diffusion and convection is made, then a set of coupled continuity equations is solved by IONEQ for each flux surface, taking into account radiation, ionization and recombination as sources and sinks from the neighbor flux surfaces. For our investigations, we first calculate anomalous transport parameters by comparison with CXRS measurements. From the impurity density profiles obtained in this way we then calculate the corresponding neoanological values.

RESULTS.
First, we present two results for density profiles of C VII, observed by CXRS of C VI at a wavelength of 343.4 nm. The CXRS light intensity is integrated during a time interval of 50 msec, the spatial resolution in direction of the effective minor radius \( r_{eff} \) is about 1 cm. In all plots below, the plasma edge is on the right, the magnetic axis is on the left, \( a_{eff} \) is about 16 cm. The density of C VII is given in arbitrary units, absolute concentrations are for these examples about \( 5 \cdot 10^{17} \text{ m}^{-3} \) in the plasma center, as confirmed by VUV-measurements. Each single dot with an error bar in the plot in figure 1 corresponds to one measured point,
the drawn line shows the result of the IONEQ calculation. The anomalous diffusion coefficient D, the inward velocity V and the exponential factor α are chosen such that the calculated line fits best to the measured data. Note that we do not perform a numerical fit in the usual sense, but much more only an adaptation by hand. For the example in fig. 1, we obtain an anomalous \( D = 1500 \pm 500 \text{ cm}^2 / \text{sec} \), \( V = 2 \pm 1 \text{ m/sec} \) inward velocity at the plasma edge, \( \alpha = 2.0 \pm 0.5 \).

Figure 2 shows some data points for the density of C VII, 100 msec before (lower line) and 100 msec after H-mode transition (upper line). Typical error bars are given, too, the x-axis indicates the position z in the laboratory frame along the neutral beam. Unfortunately, due to a poor raw data basis, only 4 points are available. Apparently, the C VII profiles are much more peaked after the transition occurred. This is due to severely changed transport for the impurities.

![Graph 1](image)

![Graph 2](image)

Figure 3 shows results for the density profiles of B VI, observed by CXRS of B V at a wavelength of 298.1 nm, together with the results of the code calculation, indicated by the drawn line. Once again, the calculated curve is fitted to the measured data points, indicated by single dots. We obtain an anomalous \( D = 1500 \pm 1000 \text{ cm}^2 / \text{sec} \), \( V = 2 \pm 1 \text{ m/sec} \) inward velocity at the plasma edge and \( \alpha = 2.0 \pm 0.5 \) for the anomalous transport calculation.

Figure 4 shows a result for density profiles of B VI, observed by CXRS of B V, 150 msec before (lower line) and 100 msec after H-mode transition (upper line). Once again, severely changed transport can be observed. Unfortunately, a code calculation cannot be performed for that changed profile because of strongly changed profile shape. Some representative error bars are given again. The y-axes are in arbitrary units, the x-axes show the effective minor radius.

![Graph 3](image)

![Graph 4](image)
Figure 5 shows a result for a density profile of B V (H-like), measured by CXRS of B IV at a wavelength of 282.4 nm, together with the result of the IONEQ calculation. The measured data are indicated by single dots, the y-axis is in arbitrary units. Apparently, that line is observable by CXRS due to the fact, that the electron impact excited line from lower l-states is superposed by the charge exchange excited line from higher l-states. After subtraction of the background intensity, CX intensity is observable. As expected, the density profile is concentrated to a shell near the plasma edge. Anomalous values obtained from the IONEQ calculation are D = 2000 ± 1000 cm²/sec, V = 4 ± 2 m/sec, α = 3.0 ± 1.0.

Figure 6 shows results of passive measurements of Fe III at a wavelength of 326.7 nm near the plasma edge. The x-axis here is given in cm in the laboratory frame, effective radii cannot be calculated outside the plasma, a = 30 cm. The y-axis is in arbitrary units, but the same for both lines. The upper line is taken several hundreds of shots after the last boronization, showing a relatively high density of Fe III, the lower one about 100 shots after the last boronization, indicating a relatively low level of Fe III near the plasma edge. Unfortunately, no clear profile shape can be deduced from the measured data.

DISCUSSION.
All values obtained for anomalous D by a comparison of IONEQ calculations with the CXRS measurements are between about 1500 and 2000 cm²/sec. Afterward calculation for neoclassical D from the obtained impurity profiles provide almost the same values, indicating only a minor deviation between anomalous and neoclassical transport for impurities.
Passive measurement of Fe III light intensity, directly after and a long period after boronization, show a clear decrease due to the boronized vessel wall, a result which is supported by VUV measurements.
The H-mode transition, and the changed transport after the transition, are modelled by taking into account the change of the shape of the electron temperature profiles which normally become slightly broader, furthermore the change of the soft-X-radiation profiles which become more peaked, and finally the increase in electron density. The broadening of the electron temperature profile, together with the peaking of the SX-profiles indicate an increase of density of impurities near the plasma center, which can be supported with the measured CXRS profiles as well as with VUV measurements of impurity radiation. Nevertheless, a change of the diffusive or convective parameters cannot be calculated with the method described, because not only the impurity density increases, but also the profile shape changes strongly after the H-mode transition.

REFERENCES.
TEMPERATURE RELAXATIONS DURING CURRENT DRIVE EXPERIMENTS IN THE W7-AS STELLARATOR

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Introduction The W7-AS stellarator has a low-shear vacuum magnetic field configuration. However, internal currents such as the bootstrap current change the rotational transform and shear. In order to keep the boundary value of the rotational transform at its vacuum value the internal currents are usually feedback cancelled by externally driven currents. Contrasting this standard mode of operation we have used the current drive capabilities at W7-AS, inductive (ICD) and electron cyclotron (ECCD) current drive, for the extreme opposite purpose: The rotational transform profile \( r(r) \) \((r=1/q)\) has been changed such that low order rational surfaces are introduced into the plasma and instabilities can occur. While previous work along this line concentrated on tearing modes saturating at finite amplitudes \( /1/\), this paper deals with non-linearly unstable MHD modes. In particular, small sawtooth oscillations and large \( m=2 \) temperature relaxations are discussed. Sawteeth appear in plasmas with strong central ECCD and correspondingly high values of the central iota and shear. The \( m=2 \) relaxations are mainly destabilised in discharges with ICD.

Inductive and electron cyclotron current drive Two methods of current drive are available at W7-AS. The inductively driven current (ICD) follows the resistivity profile of the plasma. Thus the current distribution is rather broad. The 70 GHz ECRH system allows to drive a current in a narrow region (typically FWHM=2cm) close to the resonance position with an efficiency depending on the launch angle of the waves (ECCD) \(/2/\). The radial coordinate of wave absorption can be varied to yield on- or off-axis current drive. A current can be driven in the direction of the bootstrap current (Co-CD) or in the opposite direction (Counter-CD). Due to their strongly differing current density profiles the separate and the combined application of both CD schemes allows to create a large variety of iota and shear profiles and to investigate the impact on MHD stability.

Results on sawteeth Starting from vacuum values of the almost shearless rotational transform around \( r=1/3 \), sawteeth were not destabilised with ICD alone. This is due to the fact that the necessary condition for the generation of sawteeth, \( \psi(0)>1 \) is not fulfilled even in the discharges with the largest plasma currents \( (I_{\text{p, max}}=I_{\text{BS}}+I_{\text{OH}}=19\text{kA}) \). The central current density \( j_{\text{ICD}}(0) \) is too small. Destabilisation of sawteeth can, however, be achieved in experiments
Fig. 1 Examples of discharges in which sawteeth have been destabilised by (a) ICD only, (b) combined ICD and ECCD. Note that the local density in b) is smaller by a factor of 1.5 than in a), although the central line densities shown are almost the same. This is due to measuring at differently shaped plasma cross sections.

The central line-averaged soft X-ray intensity drops by about 10%. The collapse is preceded by m=1 precursor oscillations with a frequency of 15kHz. The inversion radius \( r_{inv} = 3-4 \text{ cm} \) is consistent with estimates of the iota profile, which yield \( r_1 = 3 \text{ cm} \) for the radius of the \( \iota = 1 \)-surface and a large local shear at \( r_1 \). The ECE signals show a collapse in the relativistically down-shifted emission from central suprathermal electrons in the 30-50keV energy range. It is in phase with the soft-X collapse.

The above scenario with a near-central, high-shear \( \iota = 1 \)-surface due to on-axis ECCD has been taken as a starting point for systematic parameter variations to investigate the impact on sawteeth. The main observations are:

- Co-ICD was added to the Co-ECCD (fig. 1b). The transformer action was feedback controlled to limit the total plasma current to the value obtained in the discharges with ICD only. However, only those discharges which reached a constant transformer current and therefore zero inductive current were sawtoothing. This somewhat surprising result shows that the stability against sawteeth depends very sensitive on details of the rotational transform. Since \( j_{ICD}(0) << j_{ECCD}(0) \) and \( r_1/a << 1 \) is true in all cases, the inductive current in the non-sawtoothing discharges cannot change \( \iota(0) \) and the shear inside \( r_1 \) significantly.

- If the inductive current is increased while the ECCD power is kept constant, so that the total plasma current is \( I_p = 19.5 \text{ kA} \), the sawtooth repetition time and the collapse amplitude decrease (fig 2b). This means that the additional free energy of the added inductive current drives the instability more unstable. For the reasons discussed above, this increase in instability is unlikely to be explained by changes in the shear at \( r_1 \).
Due to the very high local shear at $r_1$ in all cases with ECCD, the inversion radius increases only very slightly by adding ICD.
- If the ECCD-current is decreased by one third, which is probably not sufficient to prevent the generation of a $\tau=1$-surface, sawteeth vanish.
- If the density is increased, the frequency of the precursors decreases.
- If the vacuum iota is increased from 0.28 to 0.37 while the total plasma current is kept constant, the discharge with the higher iota value is sawtoothing, while the other one is not. However, a $\tau=1$-surface exists probably in both cases. This is another hint that $\tau=1$ might not be a sufficient condition for sawtoothing.
- If the radius of ECCD power absorption is increased by $\Delta r_{\text{eff}}(0)=1\text{cm}$, sawteeth vanish. It is not clear, if a $\tau=1$-surface exists in this case.

The results show that the destabilisation of small sawteeth in W7-AS is very sensitive to minor changes of the current and the rotational transform profile. Due to uncertainties in the evaluation of the ECCD driven current density profile the existence of a $\tau=1$-surface in non-sawtoothing discharges cannot be proved. However, there are some indications that the generation of a $\tau=1$-surface might not be a sufficient condition for sawtoothing.

**m=2 temperature relaxations** In the above ECCD experiments, $\tau$ rises from $\tau(a)=0.45$ at the plasma boundary to $\tau(a)>1$ in the centre. Accordingly, a $\tau=1/2$-surface exists close to the boundary. In these plasmas $m=2$ modes are generally found to be stable or to saturate at finite amplitude.

However, in ICD experiments large $m=2$ temperature relaxations have been destabilised. Fig 3 shows such an event in a discharge, in which the plasma current has been ramped up to a flat top value of 16.5kA by co-ICD. Analysis of the soft X-ray data shows that the plasma undergoes successive collapses
caused by a growing m=2 mode. The mode rotates in the electron diamagnetic drift direction, the frequency is 1.6kHz. ECE-measurements show the temperature drop to affect almost the whole plasma cross section. The event is likely to be interpreted in terms of a tearing mode, where the bootstrap current plays an important role for destabilisation.

**Concluding remarks** In standard net current-free operation the internal currents in the W7-AS stellarator are generally too small to drive MHD instabilities. The current drive experiments reported here are aimed at investigating the physics of current driven instabilities by controlled changes of the current and iota profile. For future experiments it is planned to go to inductive currents large enough to introduce a t=1-surface without ECCD. This will allow the investigation of sawteeth in plasmas with a smaller local shear at the resonant surface. Interesting questions to answer will be, if the local shear has a similar strong effect on sawtooth stability as in tokamaks /3/ and if stellarator sawteeth can play a role in impurity control.

**REFERENCES**

/1/ R. Jaenicke et al., 14th Int. Conf., IAEA 1992, Würzburg, Paper IAEA-CN-56/C-2-1
DEPENDENCE OF TRANSPORT ON ROTATIONAL TRANSFORM IN THE STELLARATOR W7-AS


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Introduction - The global energy confinement time $\tau_E$ in Wendelstein 7-AS scales with electron density, heating power, magnetic field and minor radius $/1/$. The scaling deviates from that found in tokamaks, where a scaling of $\tau_E$ with plasma current is observed $/2/$. If one expresses the current scaling in terms of toroidal magnetic field, the tokamak and stellarator scaling laws become similar. In order to explore the intrinsic mechanism, giving rise to current scaling, the $t=(1/q)$ scaling of $\tau_E$ is studied in more detail in the currentless stellarator W7-AS. Indications for the existence of a favourable $t$-dependence on transport were already observed $/1/$. From the multiple linear regression analysis of $\tau_E$, a scaling

$$\tau_E \sim a^{1.8\pm0.12} B_0^{0.72\pm0.06} n_e^{0.53\pm0.03} P_{ECF}^{-0.56\pm0.03} t(a)^{0.29\pm0.08}$$

was derived, but with a poor statistical significance for $t(a)$, because of configurational effects.

There is another interesting reason for such a study, because a scaling of the form $\tau_E\sim t^{-0.4}$ is expected from plateau-scaling $/3/$. In W7-AS, also the scaling of $\tau_E$ with other global parameters is found to be close to plateau-scaling. It is of great importance, to elucidate the $t$-dependence as a sensitive element of plateau-scaling, related to the radial step-size of the transport. Furthermore, W7-X, an envisaged next step stellarator will operate at $t=1$ and will profit from such a $\tau_E$-scaling.

The modular stellarator W7-AS (R=200cm, a=18cm, B=1.25/2.5T) can be operated continuously within a quite large range of rotational transform (0.25<t<0.60) at extremely low vacuum magnetic shear and without externally driven currents. Just because of the low shear, it turns out to be a proper tool to investigate the dependence on rotational transform. Only the bootstrap current has to be compensated by small induced ohmic currents to stabilize the edge value of the rotational transform. Various transport mechanisms are studied: the energy ($\tau_E$), the electron ($\tau_{e0}$), the particle ($\tau_{p0}$) and the impurity ($\tau_{i0}$) confinement were compared at different $t$ in ECF-heated plasmas ($B_0$=1.25 T, 70 GHz 2nd harmonic x-mode).

Configurational effects - Transport was studied in ECF-plasmas under optimum confinement conditions, at t-values in the neighbourhood but distinctly off resonance values $t=1/2$ and $1/3$. In these small zones the confinement region is still free from natural resonances, originating from the five-fold symmetry of the torus. Here, the confinement should not be affected by configurational perturbations, as can be introduced in W7-AS owing to its nearly shear-free magnetic topology.

With the two radially movable main limiters at the outermost position (large aperture $r_{lim}$=18cm), increasing $t$ leads to a steady shrinking of the plasma cross-section from a limiter dominated ($t$<0.4) to a separatrix dominated plasma ($t$>0.4), caused by a change of the magnetic topology at the plasma boundary. The reduced plasma-limiter interaction can actually be demonstrated by a decrease of the limiter load. In the case of fixed programmed line density during an t-scan, variation of the plasma size will cause changes of the central electron density and has to be taken into account for the comparison of transport parameters at different values of $t$. Due to the inverse scaling of many transport coefficients with density, differences in confinement might be masked. Even for the case of a limiter dominated plasma, the plasma shape at the toroidal position, where the interferometer - controlling the density feedback - is
localized, changes slightly with $t$. This leads to corrections in the order of 10% for the line averaged density and, consequently, the central electron density.

To avoid these sources of error, measurements in purely limiter dominated plasmas with equal central electron density (measured by Thomson scattering) and plasma aperture for both $t$-values were additionally compared. Taking equal limiter load for both $t$-cases as a signature for clear limiter domination, this condition was only achieved at relatively small plasma aperture ($R_{lim}=12.5cm$).

**Methods** - The total plasma energy $W_{dia}$, measured by the diamagnetic loop, is used for studies of the energy confinement time.

For local information the electron heat conductivity $\chi_{e}^{pb}$ was calculated from a stationary power balance analysis, using $N_e$- and $T_e$-profiles from the Thomson scattering and ECE-system. An additional independent approach to derive $\chi_e$ is the investigation of heat wave propagation in the plasma. For this purpose, the ECF-heating power, which is absorbed inside a small magnetic resonance zone in the center of the plasma (representing a local heat source) was modulated in amplitude with a fixed frequency (92 Hz). From the spatial phase shift of the wave, observed by a radially resolving ECE-system, the electron heat conductivity $\chi_{e}^{hw}$ could be derived.

The particle diffusion coefficient $D_p$ at $r=0.8a$ (10cm) was calculated from 3D-DEGAS code simulations and calibrated $H_{a}$-emissions, measured at both limiters, which, at the smallest aperture of $a=12.5cm$ were the dominant sources of neutral gas. In addition to the Thomson profiles, the Langmuir probe provided the profile information for the plasma edge. By proper adjusting the input heating power and the external gas feed, two density scans were performed for the lower and higher $t$-case, respectively, keeping the local temperature constant in all cases. This offers the possibility to compare particle diffusion coefficients under identical local plasma parameters and to extract the contribution of rotational transform. Finally, the impurity transport coefficient $D_I$ was obtained from a simulation of the time traces of spectroscopic line emission from aluminum laser blow-off experiments by the one-dimensional transport code STRAHL/4/.

**Results** - As a first approach, the total plasma energy $W_{dia}$ was measured as a function of the edge value of the rotational transform $t(a)$. In the case of large plasma aperture, the total energy drops dramatically in extended regions of $t(a)$ (fig.1), indicating a strongly affected confinement by resonant island structures entering the confinement region. Even at the positions of optimum confinement the plasma has different size due to changes in the vacuum magnetic configuration with $t$. Consequently, it seems to be difficult to discover relatively weak effects, as

![Fig.1 : Total energy vs. edge rotational transform for large (dots) and small limiter aperture (open triangles). Maximum values at $t(a)=0.35, 0.5$ are corrected for constant central density (filled triangles)](image1)

![Fig.2 : Normalized simulated total energy vs. edge rotational transform for small limiter aperture. Regions of optimum confinement are marked by grey bars.](image2)
predicted by the t-scaling. The lower curve represents an t-scan with small limiter aperture ($t_{lim} = 12.5\text{cm}$), where the plasma cross section is well defined by the limiter up to $t = 1/2$. Only under this conditions, an improved energy confinement in the surrounding of $t = 1/2$ compared to the region around $t = 1/3$ could actually be observed. The arrow should elucidate the difference in the total energy at the positions of optimum confinement, assorted now with an additional relative density correction (assuming a scaling $\sim N_e^{0.53}$). An expected degradation of confinement is observable at $t(a) = 0.33$, whereas around $t(a) = 0.50$ (see also fig 3), however, a maximum of total energy was found. This might be due to the fact, that for the case of small limiter aperture and small existing shear, the edge value of $t$ is not identical with $t$ at the limiter position. This, and possible deviations from the vacuum magnetic topology can lead to a shift of the resonances on the $t(a)$-scale.

An attempt to simulate the t-scan for small plasma aperture is shown in fig 2. Here, originally undisturbed Ne- and Te-profiles were flattened in radial regions where magnetic islands are calculated, as a result of a heuristic magnetic perturbation spectrum. For that, locally high diffusion coefficients had to be applied across the extension of the islands. In spite of existing discrepancies in comparison with the experimental data, the simulation illustrates some instructive features which are already mentioned before: in the direct vicinity of the main resonances 1/2 and 1/3 distortions by other stronger resonances are not expected and optimum confinement should be achieved.

Fig 3 shows a reproduction of the scan in the vicinity of $t = 1/3$ and $1/2$ to confirm the positions of optimum confinement with small limiter aperture ($t(a) = 0.35, 0.50$). Under these two conditions, but now with the same central electron density ($N_e = 1.9 \times 10^{19} \text{m}^{-3}$) and an ECF-heating power $P_{ECF} = 360 \text{kW}$, the different types of transport coefficients were investigated. As result, the particle diffusion coefficient $D_p$ at $r = 0.8a$ (10cm) as well as the impurity diffusion coefficient $D_i$ and the electron heat conductivity $\chi_e^{pp}$, actually, show a common decrease for increasing rotational transform (fig 4). The heat wave analysis show the same tendency for $\chi_e^{hw}$, too, indicating an overall improvement of confinement at higher rotational transform.

![Image 3: Small t-scan for the total energy (filled symbols) and the al confinement time (open symbols) in the vicinity of the optimum confinement regions.](image3)

![Image 4: Particle diffusion coefficient, impurity diffusion coefficient and electron heat conductivity at the t-positions of optimum confinement.](image4)

At lower central electron densities between $0.4 - 0.9 \times 10^{19} \text{m}^{-3}$ $D_p$ was found to decrease with higher $t$ by a factor of 2 to 3. In the case of high $t$, $D_p$ may be underestimated due to the neglection of the wall source contribution, which should be strongly reduced with small limiter aperture, but is difficult to predict in this configuration.

Simply taking the maximum values for $\tau_E$ at the positions of optimum confinement in the t-scans would lead to a rough dependence of $\tau_E - t(a)^{0.8}$, which appears, indeed, unexpectedly high. The corresponding electron and density profiles for the two t-cases are plotted in fig 5.
together with the calculated $\chi_e^{pb}$ profile, illustrating the smaller $\chi_e^{pb}(r)$ in the high-t-case. Also, the statistical regression analysis, taking into account all W7-AS discharges (fig.6), provides the same global trends ($\chi_e^{pb} = t(a)^{-0.31}, t_e = t(a)^{0.29}$) for the electron heat transport.

Concerning the impurity transport, there exist only a relatively small database up to now. Evaluation of a global expression for the impurity diffusion coefficient by multiple regression analysis, leads therefore to results with extremely low significance for the dependence on rotational transform. But there might exist also another reason for the scattering. Below the t-scan in fig.3, the corresponding measured al-confinement time is plotted, showing a strong variation around t=1/2, uncorrelated to the behaviour of the total energy. However, impurity transport measurements at slightly different t-values in this region can possibly differ rather strongly.

**Discussion** - Under the assumption, that the configuration is still not distorted at the two compared t-positions with maximum energy, there is, actually, evidence for a general improvement of confinement at higher rotational transform. This is additionally supported by global scalings, obtained from statistical analysis. Nevertheless, it is of great importance to investigate the vicinity of the two main resonances in more detail, to exclude remaining distortions of the configuration. The comparison with the simulation show, that the confinement in large regions of t are not yet quite understood and have to be explored. Also, a difference in the electron density and temperature profiles compared to earlier measurement under similar conditions have to be clarified. These problems are subject of the ongoing investigations.

**References**

4/ Behringer,K., JET-R (87)08 (1987)
ICRF HEATING EXPERIMENTS IN THE W7-AS STELLARATOR USING A NARROW $k_\| \! \!$ SPECTRUM ANTENNA

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ABSTRACT

First experimental results using a toroidally broad ICRH antenna on the W7-AS modular stellarator are presented. The antenna, which is located on the high-field side in the elliptical section, has been designed to launch only a narrow $k_\|$ spectrum to the plasma. As a consequence it should be able to radiate the same power into the plasma with a voltage at the antenna surface a factor of 2 ~ 3 lower than an equivalent double strap antenna. This is expected to reduce the production of fast ions in the SOL and the impurities influx. The antenna is not provided with a Faraday screen, allowing the flexibility to launch preferentially the the FW or the IBW by a simple change of the phasing of the feeders currents. Results are presented indicating that although the dominant field component of the antenna should be $b_\|$, an unexpected high loading, possibly attributable to launching of the IBW masks the intended 2nd harmonic FW heating.

INTRODUCTION

Previous ICRH experiments [1] in W7-AS have used a conventional double strap antenna on the low-field side with an optically closed Faraday screen. Second harmonic heating experiments at both 2.5 and 1.25 T showed no clear evidence of heating even at the maximum coupled power of 300 kW. In the D(H) and He(H) minority heating at 2.5 T an initial increase of the total stored energy, with a relatively low efficiency $\frac{dW}{dt}/P_{IC} = 0.2$, is soon counteracted by a large metallic impurity influx and consequent radiative cooling. A large amount of unconfined fast ion population and a RF-induced density and impurities rise were also observed.

In order to minimize these deleterious effects a broad antenna [3] has been designed, which by removing the unwanted part of the $k_\|$ spectrum reduces the antenna near fields. The broad antenna extends 1 m toroidally and 40 cm poloidally. It is earthed to the vessel wall at both ends and is fed independently at 4 positions (A,B,C,D in fig.1) which allows changing both the $k_\|$ spectrum and the polarisation of the electric field radiated by the antenna. The vacuum $b_z$ field profile in fig.2 is notable in the absence of field peaking at the edges of the antenna strap, and around the image currents in the antenna side-limiters, which in a conventional antenna lead to a spectrum rich in high axial wavenumbers. The vacuum $b_z$ field decreases exponentially away from the antenna surface with the expected $|k_z| \approx |k_\|| \approx 6 m^{-1}$, ensuring that the FW launching efficiency has a weak dependence on the antenna-plasma separation. When the current in two toroidally adjacent feeders is $\pi/2$ out of phase this antenna also gives a very directive narrow spectrum which can be used for current drive generation.

Due to stress limitations on the field coils of W7-AS, operation of the machine has been restricted to half-field operation regime at 1.25T, which, given the generator tunability 30 MHz < $f$ < 120 MHz, allows only hydrogen plasma 2nd harmonic ICRF heating. Two typical target plasmas were used:

(i) 400 kW ECRH plasma, characterized by hot ($\approx 1$ keV) electrons, cold ($\approx 200$ eV) ions and relatively low density. The ECRH cutoff at $6 \times 10^{15}$m$^{-3}$ makes density control critical, especially if the density rises during the ICRH pulse.

(ii) 800kW NBI plasma, with high density and ions and electrons strongly coupled, with a temperature $\approx 400$eV. Because of the low electron temperature, these discharges are sensitive to radiation collapse due to enhanced impurities influx or strong gas puffing.
EXPERIMENTAL RESULTS

Initial experiments were performed feeding only the bottom two ports (B & D) of the antenna, and relying on a half wavelength resonant line on each of the top two ports to provide the desired π phasing up/down and left/right. Although the currents in the 4 ports of the antenna were equal, and the phasings correct in vacuum, during the plasma the currents in the top half of the antenna fell to less than 10% of those in the bottom half. This is attributable to the high measured 6 Ω ≤ R ≤ 12Ω plasma loading (much higher than expected for preferential FW coupling) reducing the Q of the resonant circuit below the value needed for correct feeding the upper half of the antenna. Up to 500kW was coupled to an ECRH plasma in this regime, always to be accompanied by a strong density increase. The initial rate of density increase was proportional to the coupled power however no significant impurities influx was observed. The plasma loading was independent of power over 5-500kW.

Following these experiments, the antenna feeding system was revised to actively feed all 4 ports of the antenna, by T-junctions at the voltage maxima, thus also providing a useful impedance transformation. This ensured that the desired antenna phasing was maintained throughout the plasma. The plasma loading was observed to range between 8 Ω ≤ R ≤ 16 Ω and was, for the standard i = .34 limiter plasma, independent of power. However when the limiter was open (31 cm) and i = 5 a power dependence of the loading i.e. decreasing from 16 Ω (P=10 kW) to 10 Ω (P > 100 kW) was observed. During the RF pulse both the plasma density and the total stored energy increased while T_e and T_i remained almost unchanged. The practical power limit during ECRH discharges is due to the RF induced density rise, which eventually increases the density beyond the ECRH cutoff and terminates the discharge. It was observed, however, that below some power limit (P ≤ 200 kW) the density ceased to increase before the end of the RF pulse, and could be decreased to the initial value by the feed-back control of the gas valve fig.3. The maximum power for which the density rise could be controlled was increased by a power ramp at the beginning of the RF pulse. Over about 200 kW, however, this technique became impractical due to the limitation in the power ramp length (t ≤ 100 ms). Although in vacuum the RF power is limited by thermal outgassing and consequent arcing only at 280kW (33kV) and can last for over 600ms, during NBI discharges, the power was always limited to 150kW (12kV) for 100 ms, decreasing with pulse length, due to arcing in the connections to the antenna (not in the antenna itself, which was observed by a video camera). The observation that both the reactive and resistive antenna loading depends on edge and not on line density indicates a coupling mainly to the edge plasma. Langmuir probe measurements made in the SOL toroidally opposite the antenna showed an increase in the plasma potential by 30 V for 60 kW of ICRH power although no change in either the edge density or electron temperature was observed. This suggests that the power is being deposited in the edge region near to the antenna. Some evidence of FW eigenmodes excitation was observed in the antenna loading and in the RF probe signals after ECRH switch-off when the line density was still high but the plasma temperature was decreased fig.4.

DISCUSSION

The central observation in the first experimental results from the broad antenna is the very large antenna loading which remains large even at very low values of the line density and of the total stored energy where FW loading should be negligible.

This loading is much larger than can be theoretically expected for the FW even at the maximum stored energy (R < 2 Ω). The predicted fraction of the RF power launched in the FW should be under 10%.

The absence of power dependent loading indicates against loading due to Langmuir currents to the antenna surface as seen in comparable experiments with unshielded antennas. [2,6]
The strong dependence of the loading from the scrape-off density indicates IBW launching by the parallel currents as the best candidate for the large loading observed. The increase (almost twice) in the loading after changing the antenna feeding scheme can thus be explained by assuming that the mirror currents in the plasma decrease the value of the parallel inductance and hence increase the ratio of toroidal to poloidal currents in the antenna. The upper and lower parts of the antenna can in this case be considered as separate.

Numerical comparison with theory is difficult due to the sensitivity of IBW to the exact values of the density in the region very close to the antenna ($n_e \leq 3 \times 10^{16}\text{m}^{-3}$). However, although the large loading could be consistent with some choice of the unknown density at the antenna, the theoretical model [3] failed to predict the large decrease in the parallel inductance.

The increase in the line density, without evidence of bulk heating is similar to the results of several IBW experiments [4,5].

The presence of modes, both on magnetic probes and above a large background loading of the antenna, when the stored energy is low (low damping), is consistent with only a small part of the RF power going in the FW which is only seen during sharp eigenmodes.

REFERENCES

Figure 2: Profile of $\tilde{b}_x$ vs. $z$ (cm) along the midplane of the antenna for $\pi$ phasing feeding at A & B.

Figure 3: (a) Total stored energy (kJ), (b) Line density ($10^{18}$m$^{-2}$), (c) Gas Flux (a.u.), (d) Soft-X electron temperature (keV), (e) Antenna loading ($\Omega$), 400kW ECRH plasma with 100kW ICRH.

Figure 4: (a) Total stored energy (kJ), (b) Line density ($10^{18}$m$^{-2}$), (c) Antenna loading ($\Omega$), (d) $\tilde{b}_z$ magnetic probe signal (a.u.), 800kW NI plasma with 100kW ICRH.
Experimental and theoretical investigation of density and potential fluctuations in the scrape-off layer of ASDEX

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Introduction
Electrostatic fluctuations (i.e. the magnetic field is assumed constant) are candidates for the explanation of the anomalous transport of particles and energy in both tokamaks and stellarators. While most theoretical effort has been directed to an explanation of the anomalous transport in the bulk plasma, it is now widely being realized that the anomalous radial transport in the scrape-off layer, determining the width of the power flow channel at limiter or divertor plates, may be equally important to a future reactor experiment.

In the divertor tokamak ASDEX density and potential fluctuations in the scrape-off layer were investigated with high temporal and spatial resolution by Langmuir probes and an H₀ diagnostic. Many results of these measurements were reported in [1] and are summarized below. Several of these properties of the fluctuations have also been reported from other experiments.

Basic properties of the observed fluctuations

• The fluctuations can be described by independent individual “events” superimposed randomly rather than by modes following a dispersion relation.

• In poloidal direction, these events show a periodicity with a wavelength of typically 4–6 cm, but the poloidal correlation decays within 2–3 wavelengths (the poloidal half width is about as large as the wavelength).

• Parallel to the magnetic field the fluctuations are in phase and highly correlated over a distance of at least 10 m.

• The typical lifetime of the largest fluctuation events is of the order of a few 10 μs.

• On the high field side, such fluctuations are only observed in “single null” discharges when there is a connection between inboard scrape-off layer and outboard scrape-off layer, whereas on the low field side the fluctuations are present also in “double null” discharges.

• Outside the separatrix the fluctuations propagate into ion diamagnetic drift direction with velocities of up to a few 100 m/s, but at low temperatures the velocity can also be very close to 0 or even reverse its sign; inside the separatrix the fluctuations always propagate into electron diamagnetic drift direction.

• The phase between the fluctuations of ion saturation current and of floating potential yields maximum transport radially outward and suggests the picture of eddy-like events exchanging plasma between the inner and outer parts of the scrape-off layer; the lifetime of an individual “event” just fits to allow half a turn with the radial $E \times B$ drift velocity and radial width of the scrape-off layer measured.
• The radial transport derived from the fluctuation measurements is of the right order of magnitude to explain the global particle confinement.

Modelling the scrape-off layer
In addition to the gradients of density and temperature and the curvature of the magnetic field, in the scrape-off layer the boundary conditions at the intersection of the magnetic flux tubes with the limiter or divertor plates have to be taken into account. At these intersection points electrostatic sheaths are formed and the flows of ions, electrons and energy to the target plate are connected to the density and temperature in front of the sheath and to the potential drop in the sheath.

In [2] it was shown for a fluid model that under these conditions the scrape-off layer will be unstable in the region of unfavourable magnetic curvature, if potential fluctuations are regarded in connection with either density or temperature fluctuations. Our analysis takes into account simultaneous fluctuations of all three quantities in a simplified cylindrical geometry. We assume that fluctuations are in phase along each magnetic flux tube, which is supported by the correlation measurements mentioned above and is consistent with the physical picture to be presented below. Thus, the time-dependent (fluctuating) part of the equations can be treated two-dimensionally, and a non-linear numerical analysis would be desirable, but has not been performed yet.

We performed a linear stability analysis and arrived at the following results within the limits of our model:

• The scrape-off layer is stable against perturbations
  - of sufficiently large poloidal wavelength $\lambda_{pol}$
  - or for slab geometry (no magnetic curvature) and homogeneous magnetic field.

• For unfavourable magnetic curvature (low field side in a tokamak) a stability limit $\lambda_0$ for $\lambda_{pol}$ exists, and perturbations with wavelengths $< \lambda_0$ will grow with a rate $\propto 1/\lambda_{pol}^2$ for small $\lambda_{pol}$.

• The instability mechanism in this case can be illustrated by assuming a fluctuation with poloidal pressure gradients: The resulting radial diamagnetic currents have a non-vanishing divergence due to the cylindrical geometry. They have to be balanced by parallel currents through the sheath forcing changes in the electric potential of the magnetic flux tube and thus leading to poloidal electric fields. The resulting radial $E\times B$ drift leads to a pressure perturbation in phase with the original perturbation, and thus instability results.

• In the region of favourable magnetic curvature the scrape-off layer is stable in the limit of small and of large wavelengths, but if the density gradient is sufficiently large compared to the pressure gradient, an intermediate range of unstable wavelengths can exist. This result is due to the inclusion of both density and temperature fluctuations in the model.

If we use mixing length arguments, trying to describe the convective transport due to the fluctuations by diffusion, a diffusion coefficient results which itself is proportional
to the pressure gradient. We use this diffusion coefficient to add an anomalous radial transport term to the stationary version of our model equations and thus get estimates for a diffusion coefficient and decay lengths for density and pressure. The diffusion coefficient is 

$$D \sim 1 \frac{m^2}{s} \left( \frac{T_0}{10 \text{eV}} \right)^{7/6} \left( \frac{L_c}{1 \text{m}} \right)^{1/3} \left( \frac{B}{1 \text{T}} \right)^{-1/3} \left( \frac{R}{1 \text{m}} \right)^{-2/3}$$

where $T_0$ is the mean electron temperature, $L_c$ is the connection length between the target plates, $B$ is the magnetic field and $R$ is the major radius. The decay lengths for density $L_n$ and pressure $L_p$ are in the order of 1–3 cm and scale as $T_0^{1/3} L_c^{2/3} B^{-2/3} R^{-1/3}$.

These scalings are only valid, however, if the secondary electron emission coefficient at the sheath and the mean net current through the sheath are constant. But there are indications from Langmuir probe measurements in the divertor that the net current through the sheath can vary strongly with different plasma conditions, and for this case our model predicts opposite behaviour of $L_n$ and $L_p$ which in the divertor scrape-off layer indeed has been observed.

**Comparison between model and experiments**

For density ramp discharges at ASDEX, several model results are plotted versus parameters of the correlation functions in figs. 1–3. As the model is based on several crude simplifications up to now, the qualitative agreement is quite satisfactory. The results of the model are not explicitly dependent on density but on temperature, and the temperature in the scrape-off layer scales inversely with density. Fig. 4 indicates that the typical wavelengths of the fluctuations are closely correlated to the radial width of the energy-carrying layer.

[3] gives a dependence of $L_n \propto T_0^{0.26} L_c^{0.45}$ from lithium beam experiments for a set of comparable discharges, which is not so far away from $L_n \propto T_0^{1/3} L_c^{2/3}$ given above (without taking into account variations of the secondary electron emission coefficient and the net current through the sheath), but no dependence of $L_n$ on the magnetic field could be detected in those experiments.


**Fig. 1** Typical growth time from our model versus correlation time from the temporal and spatial correlation functions of the fluctuations in the scrape-off layer as measured by the Hα diagnostic. The data are from three ASDEX discharges with density ramps (different symbols for each discharge).

**Fig. 2** Poloidal phase velocities of the fluctuations from our model versus the measured values from the correlation functions. Same data as in fig. 1.

**Fig. 3** Density decay lengths from mixing lengths estimates in our model versus the values given by the scalings in [3]. The straight line is fitted through the origin. The qualitative behaviour and the order of magnitude agree. Details may be affected by the secondary electron emission coefficient and the mean net current to the target plates.

**Fig. 4** Experimental wavelengths of the fluctuations as determined from the temporal and spatial correlation functions versus equivalent width of the power-carrying layer in the divertor as given by the scaling in [3]. The straight line is fitted through the origin. The data are from the same discharges as in fig. 1.
RECENT RESULTS WITH 140 GHz ECRH AT THE W7-AS STELLARATOR

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INTRODUCTION

Net current free plasma build-up and heating is achieved at W7-AS with up to 0.8 MW ECRH at 70 GHz (\(<3\) s) with \(B_{\text{res}} = 2.5\) T and \(n_{\text{e,crit}} = 6.2 \times 10^{19} \text{ m}^{-3}\) (1\(^{\text{st}}\) harmonic O-mode). A 140 GHz prototype gyrotron with 0.5 MW power (\(<1.1\) s) was set into operation recently. First experiments with 2\(^{\text{nd}}\) harmonic X-mode launch focused on high density operation up to \(1.1 \times 10^{20} \text{ m}^{-3}\), which could be achieved up to now with NBI only.

H-mode transitions, which are well known from tokamaks [1] were observed in W7-AS. The operational window for the H-mode is discussed. The basic features of the Stellarator H-mode were reported in previous papers [2,3]. Here we concentrate on the influence of gas puffing on the H-transitions and heat wave experiments.

Combined heating experiments with NBI were performed, which is inherently related to high density operation. Typical high-power NBI-heated plasmas in W7-AS are non-stationary, because beam particle fuelling and recycling causes a steady density rise [4]. Under combined heating conditions, however, the density can be controlled despite the beam fuelling and a strong impact on the impurity confinement was found.

H-MODE TRANSITIONS

The plasma is generated and heated by a 70 GHz ECRH prepulse, which is followed by 140 GHz ECRH. The density is ramped up to a line integrated density of \(4 \times 10^{19} \text{ m}^{-2}\) \(\left( n_{\text{e0}} = 0.9 \times 10^{20} \text{ m}^{-3}\right)\) as seen from Fig. 1. The H-mode transition to a moderately improved global confinement occurs in the flat top phase. The density starts to increase further after the transition with no external gas feed, which indicates an improved particle confinement. The H-mode transition phenomena are similar for hydrogen and deuterium operation. The most pronounced signatures of the H-mode is a drop of the H\(\alpha\) line emission signals from both limiters (only one signal is displayed) indicating a reduced particle recycling.

The edge value of the rotational transform, the density and the edge condition (limiter position) are leading parameters to achieve the H-mode. Below a line integrated density of \(2 \times 10^{19} \text{ m}^{-2}\) \(\left( n_{\text{e0}} = 4.5 \times 10^{19} \text{ m}^{-3}\right)\) no transitions were observed at 2.5 T operation. In all cases investigated so far (limited available heating power), the H-mode transitions occur only in a narrow parameter window around \(t = 0.52\). The choice of the rotational transform is closely interlinked with the magnetic topology at the plasma boundary. Magnetic islands have significant influence on the edge structure due to the low shear in W7-AS. The transition was suppressed, if the limiter was inserted about two centimetres from the last closed flux surface.
Fig. 1. Time evolution of the stored plasma energy $W$, the line integrated density $\int ndl$ (m$^{-2}$), the H$\alpha$ emission (top limiter) and the ECRH input power $P$. The dotted line marks the transition for a hydrogen discharge.

Fig. 2. Time delay (top) and amplitude decay (bottom) from stimulated heat wave propagation as a function of the effective radius in the pre-transition phase (squares) and after the transition (dots).

An unperturbed and separatrix dominated plasma boundary seems to play a significant role to establish the H-mode with the available heating power.

The smoothness of the transition is dependent on the strength of the gas puffing. Experiments with a slow ($1 \times 10^{20}$ m$^{-3}$ s$^{-1}$) and fast ($2 \times 10^{20}$ m$^{-3}$ s$^{-1}$) density ramp showed, that the transition occurs in both cases at approximately the same density of $n_{eo} = 5 \times 10^{19}$ m$^{-3}$ as a lower limit for the transitions. With a slower density ramp, however, the transition is smoothed out by a series of ELM’s.

Heat waves were exited by a modulation (typically 15%, 92 Hz) of the microwave power and measured by ECE diagnostics [5]. The phase delay of the propagating heat wave is shown in Fig. 2 as a function of the effective radius for a 200 ms quasi steady state time interval before and after the transition, respectively. The heat wave is travelling slower in the H-mode phase of the discharge as compared to the pre-transition phase. The interpretation in terms of heat diffusivity is difficult, because of the strong electron-ion collisional coupling at the high densities of $0.9 \times 10^{20}$ m$^{-3}$. Furthermore, the slight density increase after the transition would already lead to a small reduction of the heat diffusivity and has to be taken into account. A simple analysis in cylindrical geometry and assuming constant $\chi$ in the analyzed radial interval gives a reduction from 1.0 m$^2$/s in the pre-transition phase to 0.7 m$^2$/s in the H-mode phase. A thorough analysis of the experimental results taking into account the density effects and the electron-ion coupling is under way. The change of the heat diffusivity from the steady state power balance analysis is within the experimental error bars.
COMBINATION OF ECRH AND NBI

Experiments with combined heating were performed at densities around \( n_{e0} = 5 \times 10^{19} \, \text{m}^{-3} \) and \( n_{e0} = 1 \times 10^{20} \, \text{m}^{-3} \), respectively, with both, on- and off-axis power deposition of ECRH. An example is given in Fig. 3 with on axis ECRH at \( n_{e0} = 5 \times 10^{19} \, \text{m}^{-3} \). The time sequence of the heating pulses is seen from the bottom trace in Fig. 3, the discharge is sustained in the final phase by NBI only. During the ECR-heated prephase and the combined heating phase density control is maintained by feedback controlled gas puffing, whereas after switch off of ECRH the density increases with the external gas feed turned off. The density increase with pure NBI heating is a well known feature at W7-AS [4], where beam fuelling (typically 1 \( \times \) \( 10^{20} \, \text{s}^{-1} \) for 0.35 MW) and recycling fluxes have a comparable contribution to the overall density rise.

The density could be controlled also with off-axis ECRH power deposition at \( r/a = 0.5 \), at high plasma densities and with combination of 0.75 MW ECRH and 0.7 MW NBI heating. As a rule, the density control could be maintained only, if the ECRH power was approximately equal or larger than the NBI power.

Evidence for degraded particle confinement in the bulk plasma comes from the profile development in the different heating cases as seen from Fig. 4. The density profile for combined heating is flat and the \( T_e \) profile is strongly peaked for on-axis ECRH. On the other hand, the density profile becomes peaked and the \( T_e \) profile is flat in the plasma centre for off-axis ECRH. These results suggest, that the particle confinement is affected mainly by the temperature profile changes inferred by the additional ECRH. This would also explain the empirical finding, that a certain ECRH-power is needed to provide the density control. For the given high plasma densities sufficient ECRH power has to be added to shape the temperature profile. The impurity confinement was investigated by Al laser-blow off experiments, where the central impurity line emission from Al XII shows a decay time of typically 200 ms for the on axis combined heating case, whereas the decay time is much larger (and could not be quantified within the available discharge time) in the off-axis case. For both cases, the decay of the Al XII line emission is very similar in the ECRH prephase and in the combined heating phase. The degradation of the impurity confinement time with combined heating confirm the results from W7-A [6].

The understanding of the measured behaviour is of crucial importance, because combined heating may provide tools for impurity control in steady state plasma operation.

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**Fig. 3.** Time evolution of the stored plasma energy \( W \), the line integrated density \( \int n_{cl} \, (\text{m}^{-2}) \), the central electron temperature (SX) and the input power \( P \) for combined heating with on-axis ECRH and NBI.
5. SUMMARY AND CONCLUSION

A new 140 GHz ECRH system (0.5 MW, >1.1 s) was set into operation at W7-AS. The parameter range for ECR-heated discharges was extended to densities up to $1.1 \times 10^{20}$ m$^{-3}$.

H-mode transitions were observed at W7-AS, which behave similar for hydrogen and deuterium operation. The H-mode transitions where observed at densities above $4.5 \times 10^{19}$ m$^{-3}$ (2.5 T) and under optimum confinement conditions in the vicinity of $\tau = 0.5$, where the plasma is separatrix dominated. The transition occurs independent of the density ramp rate at about the same threshold density. With lower gas puffing, the transition is smoothed out by a series of ELMs. Heat wave experiments showed a reduction of the heat diffusivity in the H-mode phase.

A strong density clamping effect was measured in experiments with combined heating of ECRH and NBI. Whereas the density increases steadily in purely NBI heated discharges due to the beam fuelling and recycling, the density can be controlled in combined heating scenarios.

There is experimental evidence, that the particle and impurity confinement is affected by the profile changes inferred by ECRH.

References:
STUDY of EC CURRENT DRIVE EFFICIENCY
and BOOTSTRAP CURRENT by
POWER MODULATION EXPERIMENTS

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• INTRODUCTION •

In Stellarators, the absence of a strong "obscuring" Ohmic component permits the experimental investigation of small non-inductive currents with a precision difficult to be obtained in an equivalent Tokamak /1/. In the W7-AS Stellarator ( R = 2 m, < a > ≤ 0.18 m, $P_{ECRH} ≤ 800$ kW ) we have performed a systematic study of Electron Cyclotron Current Drive /2,3/. Uncertainties in the measurements of the ECCD efficiency, $\eta_{ECCD}$, are introduced by the unavoidable presence of the bootstrap current ( and, eventually, an Ohmic one ) and by the strong coupling between confinement properties and rotational transform, which is typical for low-shear vacuum magnetic field configurations, as W7-AS /4/. To improve the accuracy, we have applied a new perturbative procedure for the determination of $\eta_{ECCD}$. The method is based on the independent launch of two microwave beams. While one of the two beams is launched with a toroidal angle of injection corresponding to co-current drive ( with respect to the bootstrap current direction ) the second one is injected in the opposite direction ( counter-current drive ). In absence of effects perturbing the symmetry ( e.g., strong $E_\parallel$ as in Tokamaks, or strong anisotropic electron distribution functions as in presence of LHCD ) the two r.f. beams are equivalently absorbed and the two driven currents compensate. Through a modulation of the powers in both beam with same amplitude but opposite phase, the modulated ECCD contribution can be discriminated from the unaffected bootstrap and Ohmic ones. Furthermore, the time delay between the modulated power and the response of the loop voltage contains information about the radial localization of the modulated ECCD-current /5/. The choice of the optimal modulation frequency depends on the electric conductivity of the target plasma and on the radial localization of the modulated current profile. For on-axis ECRH in W7-AS, modulation frequencies $\nu_{mod} ≤ 10$ Hz have to be chosen to avoid substantial inductive screening.

• THEORETICAL DESCRIPTION •

The current profile is described as the superposition of a bootstrap, an Ohmic and an EC-driven component. These three components are evaluated independently in the frame of the linearized neoclassical theory in the long mean free path regime /2/.

The current diffusion is evaluated by a numerical solution of the diffusion equation for the voltage $V_i(r,t)$:

$$\frac{\partial}{\partial t} \left[ \sigma \cdot V_i \right] = \frac{c^2}{4\pi} \frac{1}{r} \frac{\partial}{\partial r} \left[ r \frac{\partial V_i}{\partial r} \right] - 2\pi R_0 \frac{\partial J_{ext}}{\partial t}$$

with boundary conditions:

$$\left. \frac{\partial V_i}{\partial r} \right|_{r=a} = 0 \quad V_i|_{r=0} = \frac{c^2}{4\pi} R_0 \frac{\partial V_i}{\partial r} \bigg|_{r=a} = V_{ext}$$

Here $a$ and $R_0$ are the minor and major radius, $\sigma$ the neoclassical electric conductivity, $L_{ext}$ the external inductance, $J_{ext}$ the external current source ( $J_{ext} = J_{boot} + J_{ECCD}$ ) and $V_{ext}$ the externally applied loop voltage ( $V_{ext} = 0$ for operation with open OH-transformer ). The second boundary condition follows directly from $V_i(a) + L_{ext}dI/dt = V_{ext}$.

During the antiphase power modulation of the two microwave beams, the temporal changes in $J_{ECCD}$ will induce a local $V_i$. The associated $E_\parallel$ could in principle affect $\eta_{ECCD}$ itself.
Due to the perturbative nature of the experiments, this effect has been neglected in the present study, where $\eta_{ECCD}$ has been assumed to be constant during the modulation. It was verified a posteriori that the perturbed quantities are indeed small.

Using the measured $n_e(r)$ and $T_e(r)$-profiles ( $T_{e,1} \approx 1.2$ keV, $n_{e,1} \approx 1.4 \cdot 10^{19}$ m$^{-3}$ ) and assuming a uniform $Z_{eff} = 1.2, 3$, respectively, (which is the uncertainty in the experimental determination of the $Z_{eff}$-profile) neoclassical theory predicts a bootstrap current $I_{boot} \approx 7.3, 5.4, 4.5$ kA and a resistance $R \approx 12, 20, 27$ $\mu$Ω.

Following linear ECCD theory, including trapped particles effect in the long mean free path /2/, $\eta_{ECCD}$ shows a maximum $\eta_{ECCD} \approx 11$ A/kW for a toroidal angle of injection $\varphi_{tor} \approx 10^\circ$ and $Z_{eff} = 1$. At higher $\varphi_{tor}$, $\eta_{ECCD}$ rapidly decreases reaching $\eta_{ECCD} \approx 2.7, -0.3$ A/kW at $\varphi_{tor} \approx 20^\circ, 30^\circ$, respectively. $\eta_{ECCD}$ scales as $1/(5 + Z_{eff})$.

**EXPERIMENTAL CONDITIONS**

The experiments were performed with 2nd-harmonic, (X)-mode injected from low field side. The ± symmetry in the toroidal angle $\varphi_{tor}$ couldn't be fulfilled. A power $P_{cstr} \approx 200$ kW generated by the first gyrotron was injected at $\varphi_{tor} \approx -10^\circ$ (EC driven current opposite to the bootstrap current) while the power generated by the second source, $P_{cso} \approx 170$ kW, was injected at $\varphi_{tor} \approx 35^\circ$. Following the theoretical predictions and the results of previous ECCD experiments /3/, $\eta_{ECCD}$ should become negative at $\varphi_{tor} = 35^\circ$. Therefore, $P_{cso}$ is injected not to drive any relevant ECCD but with the aim of limiting the excursion in $T_e(r)$ (and consequently $I_{boot}$) during an antiphase modulation cycle. Modulating the two beams in antiphase, ($\nu_{mod} = 10$ Hz, $|\Delta P_{cso}| = |\Delta P_{cstr}| \approx 60$ kW) a $\delta T_e/T_e \approx 5\%$ was measured by ECE which is to be compared to $\delta T_e/T_e \approx 30\%$ in case of in phase modulation. Thus, for antiphase modulation, changes in $I_{boot}$ are assumed to be negligible compared to the induced modulation in $I_{ECCD}$. Experiments using the exact launching symmetry, where the previous assumption becomes unnecessary, are in program and will permit more definite conclusions.

**EXPERIMENTS WITHOUT CURRENT CONTROL**

In this first kind of experiments, the total injected power $P_{inj} = P_{cso} + P_{cstr}$ was kept constant (within ±10 kW) but the relative contribution from the two sources ($P_{cso} + \Delta P$ and $P_{cstr} - \Delta P$) was varied from shot to shot. Because the pulse length of 1.1 s was too short to reach a steady state current, see Fig. 1, the asymptotic steady-state current $I_0$, the plasma resistance $R$ and inductance $L$ were determined by means of a least square fitting of the time traces of the net current $I_{pl}$ and the loop voltage $V_{loop}$ with $I_{pl}(t) = I_0 \{1 - \exp[-(t-t_0)/(L/R)]\}$ and $V_{loop} = -L \cdot d I_{pl}/dt$, respectively. An example of the fit for $I_{pl}$ is seen in Fig. 1, which gives $L/R \approx 430$ ms ($L \approx 5.8$ $\mu$H, $R \approx 13.5$ $\mu$Ω, respectively). This experimentally determined $R$ corresponds to the neoclassical value if a uniform $Z_{eff} \approx 1.2$ is assumed. However, allowing for the enhancement in the conductivity related to the r.f. induced deformations of the electron distribution function /6/, similar $R$-values could be obtained with higher $Z_{eff}$. After wall boronization, $\approx Z_{eff} > 2$ are inferred from spectroscopy measurements for this type of discharges.

The extrapolated value, $I_0$, for the steady state current can be interpreted as:

$$I_0 = I_{boot} + \eta_{ECCD}^{ECCD} \cdot (P_{cso} + \Delta P) + \eta_{ECCD}^{ECCD} \cdot (P_{cstr} - \Delta P)$$

Assuming that $I_{boot}$ doesn't change significantly during the $\Delta P$-scan (note that $P_{inj} = const$) the changes in $I_0$ with respect to the reference case with $\Delta P = 0$ will give a measurement of $\eta_{ECCD}^{ECCD} - \eta_{ECCD}^{ECCD}$ through the linear relation $\Delta I_0 = (\eta_{ECCD}^{ECCD} - \eta_{ECCD}^{ECCD}) \cdot \Delta P$.

As expected a linear relationship, corresponding to $\eta_{ECCD}^{ECCD} - \eta_{ECCD}^{ECCD} \approx 8.5$ A/kW, is found between $\Delta P$ and $\Delta I_0$, as seen from Fig. 2. For $|\Delta I_0| \geq 150$ A the diamagnetic energy start to show a response, limiting the $|\Delta P|$-window where this technique can be applied. $|\Delta I_0|$-window of similar width are observed also under current controlled operation where the linear relationship between the increment in $V_{loop}$ and the induced $\Delta I_{pl}$ is used to determine the plasma resistance from $R = \Delta V_{loop}/\Delta I_{pl}$.

The derived efficiency $\eta_{ECCD}^{ECCD} - \eta_{ECCD}^{ECCD} \approx 8.5$ A/kW, is in good agreement with the values obtained by the linear model (once trapped particles effects are taken into account).
Ray-tracing calculations predict in fact $\eta_{ECCD}^{\text{E}} \approx -0.3$ A/kW and $\eta_{ECCD}^{\text{tr}} \approx -11$ A/kW, for $z_{\text{eff}} = 1$.

* EXPERIMENTS WITH CURRENT CONTROL *

The uncertainty introduced by the strong dependence of the confinement properties on the net current itself can be alleviated by keeping $I_{pl} = 0$ by a feedback-induced Ohmic component, i.e., $\tau_{\text{a}} = \text{const}$ during the discharge. The disadvantage of this operational scenario is that the determination of the non-inductive current, will now rely on the knowledge of the plasma resistance $R$. For launching conditions equivalent to the $I_{pl} \neq 0$ case, a $V_{\text{loop}} \approx -38$ mV has to be applied to keep $I_{pl} = 0$. For a consistency check, we assume that the non-inductive current is that observed under $I_{pl} \neq 0$ operations, $I_{0} \approx 2.8$ kA, this should imply $R = V_{\text{loop}}/I_{0} \approx 13.6$ $\mu$O, in excellent agreement with the value inferred in the previous Chapter. Fig. 3 shows the loop voltage for a discharge where at $t = 630$ ms, $\Delta P_{\text{tr}}$ is reduced from 205 to 170 kW while $P_{\text{co}}$ is simultaneously increased from 127 to 157 kW ($\Delta P = 35$ kW) to keep the total power constant. After having relaxed to the new steady state, an higher $V_{\text{loop}} \approx -44$ mV is necessary to counterbalance the increase in the non-inductive component. This corresponds to an increment of the non-inductive current $\Delta I_{0} \approx 400$ A, using the experimentally determined $R \approx 13.6$ $\mu$O. Under the usual assumptions, this should be interpreted as $\Delta I_{0} = (\eta_{ECCD} - \eta_{ECCD}^{\text{tr}}) \cdot \Delta P$ and consequently

$$\eta_{ECCD} - \eta_{ECCD}^{\text{tr}} \approx +11$$ A/kW, again in good agreement with the previous and the theoretical estimations.

* EXPERIMENTS with PHASE and ANTIPHASE power MODULATION *

Fig. 4 compares the response of the loop voltage when the r.f. power injected by the two gyrotrons is square wave modulated in phase or in antiphase, respectively ($\nu_{mod} = 10$ Hz, $|\Delta P_{\text{co}}| = |\Delta P_{\text{tr}}| \approx 50$ kW). For in phase modulation the total injected power $P_{\text{inj}} = P_{\text{co}} + P_{\text{tr}}$ is also modulated so that the temperature profile and consequently the bootstrap current, shows a relatively large excursion ($\Delta T_{e,p} \approx 300$ eV) during the modulation phase. This is reflected by the strong and immediate response of $V_{\text{loop}}$ seen in Fig. 4. In case of antiphase modulation $P_{\text{inj}}$ remains constant and the loop voltage shows a delayed and considerably smaller response. This is consistent with the assumption that $J_{\text{loop}}$ remains nearly unaffected by the modulation while $J_{ECCD}$ reacts in phase with it. As the $J_{ECCD}$ is driven close to the plasma axis (where the power is absorbed) with a narrow radial profile of $\approx 3$ cm, the time delay between the gyrotron modulation and the response of $V_{\text{loop}}$ reflects the current diffusion time. A time dependent simulation is shown in Fig. 5, where the centrally peaked current profile evaluated by the ray-tracing is used. The total current was fixed to $I_{ECCD} = 500$ kA in agreement with the efficiency of $\eta_{ECCD} = \eta_{ECCD}^{\text{co}} + 10$ A/kW determined in the previous experiments.

* CONCLUSIONS *

Preliminary experiments have shown the potential of perturbation experiments for the analysis of electron cyclotron current drive, bootstrap current, plasma resistivity and their radial profiles. Although we have not used symmetric launch a separation of the different current components was possible.

The results from three different perturbative experiments were found to be consistent with the neoclassical linear modelling of ECCD, bootstrap current and electric conductivity. In particular the deduced value for the ECCD-efficiency shows a good agreement with the expectations of linear theory, when trapped particle effects are included.

The time traces of the loop voltage during in phase and antiphase modulations are consistent with a current profile formed by an ECCD-profile peaked on-axis and a broad bootstrap component as expected from the W7-AS $T_{e}$ and $n_{e}$-profiles.

Experiments with off-axis and using symmetric launch are under way.

* REFERENCES *

/2/ U. Gasparino et al., Theory of Fusion Plasmas, Varenna 1990, 195
/5/ U. Gasparino et al., ETC on r.f. Heating and Current Drive, Brussels 1992, Vol. 16E, 301
/6/ N.J. Fisch, Phys. Fluids, 28 (1985), 245
Fig. 1 Time traces of the net toroidal current for a discharge where no external loop voltage is applied. From the fit (full line) the asymptotic steady-state current $I_0$ and the L/R time scale are determined.

Fig. 2 Reaction of the extrapolated net current $I_0$ to a variation in the co- and counter-injected powers. While keeping the total power $P_{co} + P_{ctr}$ constant, the power responsible for co-current is changed to $P_{co} - \Delta P$ simultaneously to a variation $P_{ctr} + \Delta P$ in the counter-injected power.

Fig. 3 Reaction of the loop voltage when, at the time $t = 0.63$ s, the co- and the counter-injected powers are varied by an amount $\Delta P \approx 35$ kW (keeping the total injected power constant). The full line refers to a simulation, using a $\Delta J_{ECCD}$ in agreement with theoretical predictions (see Fig. 5).

Fig. 4 Response of the loop voltage to a square wave modulation ($\nu_{mod} = 10$ Hz) of the co- and counter-injected powers. In case of in-phase modulation a strong and immediate response is seen, while for antiphase modulation $V_{loop}$ shows a delayed and considerably smaller response.

Fig. 5 Analysis of the antiphase modulation of Fig. 4. The full line shows the theoretical expectations. In this case only the EC-driven current is modulated with $J_{ECCD}(r)$ in agreement with ray-tracing calculations. The dotted line is shown as reference and corresponds to the response in $V_{loop}$ if the modulated current had a profile proportional to $J_{boad}(r)$. 
TEMPERATURE, DENSITY AND POTENTIAL FLUCTUATIONS BY A SWEPT LANGMUIR PROBE IN WENDELSTEIN 7-AS

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1. INTRODUCTION

Numerous experiments using a Langmuir probe to investigate the magnitude of temperature fluctuations and their contribution to heat transport in the edge region of tokamak plasmas have been carried out (1,2). Sweeping the voltage applied to a tip fast enough to ensure that the ion saturation current, floating potential and electron temperature may be assumed to be constant during the sweep is experimentally more difficult than alternative schemes but this disadvantage is compensated by the ability to measure all three of these quantities at one spatial location (3,4). Sweep frequencies up to 600kHz have been employed to obtain the current-voltage characteristic. A radial scan in the vicinity of the velocity shear layer on W7-AS stellarator was performed. Inside and outside the shear layer the normalised magnitude of the temperature fluctuations was found to be approximately 30% larger than the magnitude of the electron density fluctuations, approaching a value of 0.12 and 0.09 respectively at a radial position 1cm inside the shear layer. An increase in the coherency of the temperature, floating potential and density fluctuations between tips with a poloidal separation of 2mm was also measured as the shear layer was crossed. Heat conduction produced by correlated temperature and poloidal electric field fluctuations is therefore possible. An increasing coherence of temperature and floating potential fluctuations leads to an increase in the coherence of temperature and plasma potential fluctuations as the shear layer was crossed.

2. EXPERIMENT

The Langmuir probe characteristic was obtained by sweeping the applied voltage between the tip and the vacuum vessel. Transformer coupling of the broadband amplifier (Amplifier Research; 1000W CW from 10kHz to 220MHz) and a dummy BNC cable were used to balance the capacitive currents. The cable length from the transformer to the probe tip was 5m and so for an applied voltage at 400kHz with an amplitude of 100V a capacitive current of 125mA would flow in one of the BNC cables. With compensation this capacitive current could be reduced to 5mA. Low noise broadband buffer amplifiers were used to deliver the applied voltage monitor and current probe output to the data acquisition system (8-bit Nicolet digital oscilloscope; Model Pro 60). For a typical sweep frequency of 400kHz, a data acquisition rate of 50 MHz was used. The 256kB available memory therefore allowed the analysis of 5ms of data consisting of 4000 individual current-voltage characteristics. A non-linear least squares fit of the current-voltage characteristic
yields the ion saturation current, floating potential and temperature. Only those points at a voltage less than $kT_e$ above the floating potential were fitted. The mean fitting error (with $T_e/T_\theta = 0.03$ in this experiment) is required to be as small as possible as this is one of the quantities which determines the lowest level of normalised temperature fluctuations that may be analysed. The testing of recently available ferrite cores, which have low losses at a frequency of 1MHz and so can cope with the typical 100W power being transformed, is being carried out and this will allow the sweep frequency of the voltage used to obtain the current-voltage characteristic to be increased in the near future. Offset biasing of the swept voltage down to -200V avoided perturbation of the plasma encountered when sweeping the probe into the electron saturation regime (5).

3. RADIAL PROFILES

The radial profiles of the mean values and normalised fluctuation magnitudes of density, floating potential and temperature ( $n_e/n_{e0}$, $\tilde{\phi}_{\parallel}/T_e$ and $\tilde{T}_e/T_e$ respectively ) are shown in Fig. 1. The fluctuation magnitudes shown are the RMS values in the frequency range 30kHz to 350kHz. Existing measurements show a wide discrepancy with $\tilde{T}_e/T_e$ greater than, comparable to and less than $n_e/n_{e0}$ being reported (2,4,6). An apparent temperature fluctuation with $\tilde{T}_e/T_e = 0.05$ due to $n_e$ and $\tilde{\phi}_{\parallel}$ during a sweep is estimated (3). In contrast to recently reported measurements, it was found that $\tilde{T}_e$ and $n_e$ were not out of phase as the shear layer was crossed (2,4). As shown in Fig. 2, both the coherency of $\tilde{T}_e$ and $\tilde{n}_e$ up to a frequency of 100kHz and the coherency of $\tilde{T}_e$ and $\tilde{\phi}_{\parallel}$ increased as the shear layer was crossed. The level above which the coherency is statistically significant is indicated.

4. SPATIAL CORRELATION AND TRANSPORT

Three poloidally separated tips were swept simultaneously. The spatial coherency of density, temperature and floating potential fluctuations measured at radial positions inside and outside the shear layer are shown in Fig. 3. In each case, this coherency decreases as the poloidal separation of the probe tips is increased from 2mm to 4mm and increases as the shear layer is crossed. An improvement in the spatial coherency of $\tilde{T}_e$ is expected for measurements at higher sweep frequencies. The plasma potential, $V_p$, is given by:

$$V_p = \alpha_{\parallel} + \alpha T_e$$

where $\alpha = 2$ was used. Fluctuations of plasma potential generate poloidal electric field and radial ExB velocity fluctuations. Coherent velocity and density fluctuations induce convective heat transport, while coherent velocity and temperature fluctuations induce conductive heat transport. The ratio of heat transport by conduction and convection is given by (3):

$$\frac{\tilde{T}_e}{T_e} / \frac{\tilde{n}_e}{n_{e0}}$$

Ultimately, experiments with the swept Langmuir probe are aimed to yield values of this ratio. The observed systematic increase in the power weighted coherence between $\tilde{T}_e$ and $V_p$ as the shear layer was crossed, is based on the increasing coherence of $\tilde{\phi}_{\parallel}$ and $\tilde{T}_e$ as the shear layer was crossed. Also a systematic decrease in the power weighted coherence between $\tilde{n}_e$ and $V_p$ was observed as the shear
layer was crossed. The implications for the ratio of conductive and convective transport is presently under analysis.

The experiments investigating the spatial correlation of temperature fluctuations were motivated by the static Langmuir probe measurements of spatially correlated $\tilde{N}_e$ and $\tilde{\phi}_f$ and the turbulent eddy model (5,7). The extension of these results to include the cross correlation of $\tilde{T}_e$ with $\tilde{N}_e$ and $V_P$ will be presented.

5. H-MODE

The time development of ELM’s in the recently discovered H-mode on W7-AS is another topic able to be studied by the swept Langmuir probe technique. With measurements on the microsecond time scale, the electron temperature and density pulse associated with an ELM can be resolved, enabling estimates of particle transport by coherent density and plasma potential fluctuations during an ELM to be carried out. At a position approximately 5cm away from the separatrix, an ELM is observed typically as a jump in temperature from 30eV to 50eV with a relaxation to the original value on a time scale of 1ms.

6. CONCLUSION

The spatial coherency of $\tilde{N}_e$ and $\tilde{\phi}_f$ found in static probe measurements have been reproduced with the sweep probe technique. The smaller poloidal correlation length of $\tilde{T}_e$ inferred from its smaller spatial coherency could be partly explained by the contributions of an apparent $\tilde{T}_e$ induced by $\tilde{N}_e$ and $\tilde{\phi}_f$ during a sweep and the gaussian noise introduced by fitting. Combined this amounts to $\tilde{T}_e/\tilde{T}_e = 0.06$ which is half the measured fluctuation level at the innermost radius. An increase in the spatial correlation of $\tilde{T}_e$ and in the coherence of $\tilde{T}_e$ and $\tilde{N}_e$ is needed for consistency with the turbulent eddy model (7). Experiments at higher sweep frequencies will decide whether or not the present results remain in contradiction with this model.

Fig. 2 Coherency for a single tip at a radius inside and outside the shear layer between:
(a) temperature and floating potential fluctuations
(b) temperature and density fluctuations
\( r/\alpha_{\text{shear}} = 1.02 \)
\( r/\alpha_{\text{shear}} = 0.94 \)

Fig. 3 Coherency of density, temperature and floating potential fluctuations between two tips separated poloidally by 2mm and 4mm at a radius inside and outside the shear layer.
Edge plasma profile and particle transport study on the

WENDELSTEIN 7-AS stellarator

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1. Introduction. W7-AS has toroidally varying plasma cross sections (five field periods) and a complex magnetic surface topology at the edge which is more or less governed by "natural" islands (at rationals \(5/m\) of the rotational transform \(t = n/m\)) of up to several centimeters radial and poloidal extension [1]. The complexity is further increased by two movable, asymmetric local limiters introducing inhomogeneous connection lengths. Except a narrow \(t\)-range close to 1/3, where magnetic surfaces extend to within the limiter shadow, a sophisticated plasma edge transport analysis requires three-dimensional treatment. In the case with \(t\) just above 1/3 (closest to tokamak limiter configurations) reasonable results were obtained by a 1D (radial) approach basing on singular flux bundles [2, 3]. What we did now is, to explore the ground by Langmuir probes, in what manner the edge plasma parameter profiles are influenced by the resonant boundary structures in the higher-\(t\) case (0.44 - 0.58, including the resonances at 5/11, 5/10 and 5/9). Further we searched for some ordering parameters allowing a comparison of estimated transport coefficients in the SOL at higher \(t\) (at least in ranges free of stronger resonant perturbations) with results found at low \(t\). We focus on particle transport in this first attempt.

2. Boundary magnetic islands and edge plasma profiles. The chosen discharge parameters guarantee a magnetic field configuration close to the vacuum field (flat top ECRH discharges with \(P_{\text{heat}} = 160\) kW and constant line density, \(<n_e> = 1 \times 10^{19}\text{m}^{-3}\), net plasma current compensated to zero). The two main limiters placed at the top and bottom of an elliptical cross section (shifted by one field period) were kept fixed at the outermost position. Two fast reciprocating Langmuir probes (CFC tips, triple mode, temperatures derived by averaging after exchanging positive and floating tip) were placed at the outer tip and more towards the bottom of a triangular plasma cross section (fig. 2a). A condensed overview on the results is given in fig. 1, showing isolines of the ion saturation current density \(j_s^+\) and \(j_s^+T_e\) measured at the triangular tip (probe 1) as functions of the edge rotational transform \(t(a)\). Multiplied by an adequate energy transmission coefficient (= 10) the latter represents the parallel energy flow to a sink. The outstanding feature is a strong radial modulation of the profiles, with flat regions at the resonances 5/11, 5/10 and 5/9 (a shift with respect to \(t(a)\) is explained by the fact that \(t(a)\) refers to a fixed surface at the edge). Concerning the flat regions at the resonances, a detailed connection length analysis shows that the limiters become efficient only outside the islands in these cases. The expansion of the plasma cross section by the closed islands (relative to neighbouring \(t\) ranges where the islands are intersected by the limiters) and long connection lengths in the vicinity of the X-point broaden the profiles along the line-of-sight of the probe. Transport simulations by field line tracing with cross-field diffusion [4] comparing this case (5/9) with a neighbouring-\(t\) case, show good
qualitative agreement (fig. 2c, d). It should be remarked that, though the configurational cross section is increased at the resonances, temperature profiles in the plasma core measured by Thomson scattering generally contract in these ranges indicating short-circuit transport by the islands. Starting from the 5/9 resonance towards smaller \( i \), the 5/9 islands are shifted outward (positive shear) and become increasingly intersected by the limiters (separatrix-dominated configuration). Inside a small \( i \) range between 0.52 and 0.53, connection lengths become short (< 10 toroidal revolutions, \( \approx 120 \) m) close to the X-point and inside the 5/9 islands. This holds for the whole boundary because in the 5/9 case there is only one island which closes after 9 toroidal and 5 poloidal revolutions. That means in this case we have a separatrix-dominated regime with optimum plasma cross section. This is exactly the \( i \) range where we (up to now exclusively) found the H-mode. In contrast to this region the range governed by the 5/10 resonance (\( \langle i \rangle = 0.47 - 0.5 \)) is limiter-dominated within discrete poloidal sectors (5 separated islands), and in between the field lines end at other installations or the wall. Towards smaller \( i \) values the configuration is generally limiter-dominated.

3. Particle transport. For comparison of the particle transport at high \( i \) with the limiter-dominated "standard case" at \( \langle i \rangle \) close to 1/3 (0.345), the "optimum" separatrix-dominated configuration at \( \langle i \rangle = 0.525 \) was selected. Discharge conditions were 350 kW ECRH, maximum limiter aperture and varying densities. Radial density profiles outside the LCMS as obtained by the probes described above were exponentially fitted. Averaged particle diffusion coefficients were estimated then from the e-folding lengths by the zero-dimensional approach \( \lambda_n^2 = 2 D L_c / c_{si} \) with \( D \) the diffusion coefficient, \( 2 L_c \) the connection length and \( c_{si} \) the ion sound speed. Parallel connection lengths \( L_c \) were calculated from the vacuum magnetic field and averaged along the line-of-sight of the probes and over a poloidal distance of \( \pm 2 \) cm from the probes. Finite \( \beta \) effects should be negligible with this respect. \( \lambda_n \) is referred to distances from the LCMS which were averaged along the parallel direction of the flux tubes intersected by the probes. The result is shown in fig. 3, where averaged particle diffusion coefficients obtained by the two probes for the two configurations are plotted versus the electron density at the LCMS. \( \lambda_n \)-values referred to the local distance from the LCMS were found to differ by up to a factor of 3 for the two probes, reflecting the strongly varying local distances of the flux tubes from the LCMS [5]. Nevertheless, after applying the averaging procedures described above, which account for the topological characteristics, a uniform scaling of \( D \) with the edge density was obtained. For the low-\( i \) range, this behaviour was found already earlier [1, 3]. We conclude that, within this approach (perhaps a factor 2 - 3 of \( D \)) and referred to the edge density as ordering parameter, the particle transport does not significantly differ for the regimes and poloidal ranges investigated.

References
[5] F. Sardei et al., this conference
Fig. 1: Isolines of (a) the ion saturation current density $j_s^+$ and (b) $j_s^+ T_e$ measured by probe 1 (see fig. 2a) versus the edge rotational trans-form $\tau(a)$. The dottet line indicates the LCMS defined by the limiters or the separatrix. Symbols L (fig. b) - limiter-dominated, S - separatrix-dominated configuration ("island-divertor"), IS - limiter surface outside closed islands. Rationals at the bottom indicate the type of resonances at the boundary.
Impurity Fluxes and Profiles in Wendelstein 7-AS

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1. Introduction: Wendelstein 7-AS is a modular stellator with 5 magnetic field periods, low vacuum magnetic shear and high aspect ratio (≥10). The minor and major radii are a ≤ 0.18 m, R₀ = 2.0 m. The plasma cross section changes from nearly triangular in the middle of each field period (ϕ = 0°) to nearly elliptical at both ends (ϕ = ±360°).

Impurity influxes and profiles are investigated in a poloidal plane, intersecting the top limiter, close to the elliptical cross section of the plasma at ϕ = -25.4°. Figure 1 shows the poloidal cross section in this plane for a rotational transform of ι = 0.345 together with the port axis (dashed) and the typical range of lines of sight (LoS) of the visible spectroscopy scanning mirror system (arrows). At this low rotational transform the set of nested flux surfaces within the natural separatrix (outermost plotted flux surface) would spread out to the vessel wall and has always to be restricted by a limiter (horizontal bar). Thus, the limiter determines the last closed flux surface (LCFS), indicated as a bold contour (shown here for maximum aperture). For high rotational transform (ι > 0.5) the plasma is limited by a natural separatrix located inside the maximum aperture of the limiters. However, also in this case the movable limiters can be inserted to control the plasma edge. The plasma aperture can thus be varied from 18 cm ≤ a ≤ 12 cm.

In the following we investigate the poloidal distribution of impurity ions in this cross section to determine local sources. Then we concentrate on the localisation of the ionic shells and determine the influence of the plasma aperture on the radial location of the ions. In a next step we deduce the radial pressure profile of BIV and deduce the local diamagnetic impurity drift which is of importance for the measurements of radial electric fields. Finally we investigate the 'limiter peaks', observed for all impurity species and its dependence on limiter position. With most measurements performed on boron species we want to stress the diagnostic potential of this impurity and give a compilation of all boron transitions throughout the UV- and visible spectral range, including molecules, as measured on W7-AS.

2. Poloidal distribution of impurities: Within the geometry of fig.1 the intensity distributions across the whole poloidal cross section are measured for all ionization stages of boron (B I - B V), carbon (C II - C V) and oxygen (O I - O V) throughout the UV- and visible spectral range (200 nm ≤ λ ≤ 800 nm). In order to localize impurity sources, emission from
Fig. 2: Cross sections of the vacuum magnetic field (a, b) at $\phi = 72^\circ$ (probe plane) showing the $5/9$ boundary islands for two values of the boundary rotational transform $t(a)$, and respective Poincare plots taking into account diffusion (c, d). The diffusion is modelled by perpendicular displacements of the particle motion along the magnetic field after each integration step length ($D_{\text{eff}} = 1 \text{m}^2 \text{s}^{-1}$, for details see ref. [4]). The plots illustrate typical constellations discussed in the text: closed islands (profile broadening) and islands intersected by the limiters (separatrix-regime, "island-divertor").

Fig. 3: Averaged edge particle diffusion coefficients $D$ (see text) versus the edge density for two different configurations (limiter-dominated at $t(a) = 0.345$, separatrix-dominated at $t(a) = 0.525$) and two different poloidal probe positions (probe 1: close to an X-point, probe 2: neighboured to an O-point in the separatrix case).
Fig. 1: Isolines of (a) the ion saturation current density $j_s^+$ and (b) $j_s^+ T_e$ measured by probe 1 (see fig. 2a) versus the edge rotational transform $t(a)$. The dotted line indicates the LCMS defined by the limiters or the separatrix. Symbols L (fig. b) - limiter-dominated, S - separatrix-dominated configuration ("island-divertor"), IS - limiter surface outside closed islands. Rationals at the bottom indicate the type of resonances at the boundary.
radial location as the temperature increases (see also fig.5) or the aperture is changed. Fig. 3 shows the location of BII (squares + dots) and BIV (triangles + diamonds) for plasma apertures of a=17.5 cm (open symbols) and 14.4 cm (full symbols). The BII shell essentially becomes broader and the location relative to the LCFS increases while the BIV shell stays about 2 cm inside the LCFS and is shifted together with the aperture.

For more detailed investigations and for the deduction of ion temperatures (from Doppler broadening), drift velocities (from Doppler shift) and radial BIV density profiles as a function of time we performed shot-to-shot measurements on BIV and BII (also during H-mode discharges). Fig.4 shows the measured LoS integrated BIV intensity profile (dots) for a plasma with \( t = 0.563 \) and a separatrix radius of about 14 cm. The profile was taken at a time point of 0.5 s, where it is steepest, as shown in fig.5. Inversion of the intensity profiles yields the radial emissivity and thus the radial impurity density profiles. The BIV density profile (dashed line in fig.4) is located around 11 cm. BII is found to reside in these discharges around 15 cm. From the boron ion temperatures and densities the pressure profiles and the local diamagnetic impurity drifts are calculated. The local diamagnetic drift of BIV, shown in fig.4 as full line, turns out to be very steep at the outside of the profile. The drift is in ion-diamagnetic direction at the outside and changes sign when the maximum of the pressure profile is crossed. Fig. 5 shows the profile contours of the BIV shot-to-shot intensities from which the radial location and shift of the profile during the discharge as a function of time is measured. From 0.4 to 0.5 s the ECRH is increased from 0.17 MW to 0.34 MW resulting in a temperature increase and in a slow outward shift of the profile centre by 1.5 cm. However, the 0.9 intensity level at the inside shows a much more pronounced shift of over 3 cm leading to a strong change of the gradient at the inside.

The measured boron ion temperatures may also be inverted to yield an ion temperature profile across the spread of the corresponding shell. With the knowledge of the diamagnetic drift of the boron ions a radial drift velocity profile due to a radial electric field in the H-mode may be inferred and thus a profile of the radial electric field.

Fig.3: Radial position of BII (squares + dots) and BIV (triangles + diamonds) ionic shells as a function of \( r_{eff} \) for plasma apertures \( a = 17.5 \) cm (open symbols) and 14.5 cm (full symbols), respectively. Note the broadening of BII and the shift of BIV.

Fig.4: Radial profiles of BIV from a series of 9 shots: measured LoS integrated intensity (red dots and spline fit), emissivity (BIV density) from inversion of the intensity profile (dashed).
The location of the BIV shell as measured by visible spectroscopy has been found to be in agreement with transport calculations using the measured Thomson profiles of electron density and temperature and typical transport coefficients.

4. Influence of the limiter: In order to investigate the impurity influx from the limiter and its influence on the measured LoS intensities at $r_{\text{eff}} > 0$ we studied plasmas with smallest aperture of $a = 12$ cm, limited only by the bottom limiter one field period away, but with the top limiter at full aperture ($a=16$ cm). We find that the BIV intensity shell location (as well as the shells of the other high ionization stages) reside around the same radius of about $r_{\text{eff}} = 10$ cm for open and closed bottom limiter. However, the width of the BIV shell is found to be about twice as wide ($\Delta r_{\text{eff}} = 6.5$ cm) in the case of the open bottom limiter than in the closed case. In case of the closed limiter the 'limiter peak' of BIV drops significantly by about a factor of 4 (and becoming $=30\%$ smaller than the 'shell peak') and is shifted inward compared to the open limiter case. For BI the 'limiter peak' is found to be only slightly shifted inward and the peak is much broader compared to the open limiter case. Furthermore, the BI peak intensity from the top limiter drops remarkably, as expected, by about a factor of 10 for the closed bottom limiter case.

5. Conclusions: We presented poloidal distributions and radial profiles of the light impurities boron, carbon and oxygen measured in a poloidal cross section of the W7-AS stellarator. Scans across the whole poloidal cross section of the plasma, inner wall and limiter identified the limiter as the only prominent impurity source. Important differences of the poloidal distribution between low and high ionization states were found. The radial distribution of the different ionization stages of impurity ions was measured in one discharge or on a shot-to-shot basis. While the low ionized species reside outside the separatrix or the LCFS, BIV and other higher ionization stages were found about 2 cm inside. Density profiles of the boron ions were measured yielding the impurity diamagnetic drift. Our measurements allow for simultaneous determination of impurity densities, temperatures and drift velocities (radial electric fields) and will allow to determine radial profiles of all these quantities across the spread of the impurity shells as a function of time. The different species and transitions of boron yield a high diagnostic potential and will always stay at a diagnostic level due to boronization. We will therefore present a compilation of all species and transitions including BH-molecules, measured recently.
ION ENERGY CONFINEMENT AT THE W7-AS STELLARATOR

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Introduction

The neoclassical ion energy confinement in the long mean free path (LMFP) collisionality regime in stellarators depends sensitively on the radial electric field. For $E_r = 0$, the radial drift of ions being trapped in local field ripples leads to the well known $T_i^{9/2}$-scaling of the ion heat flux, $q_i$, in this so called $\frac{9}{2}$-regime. A sufficiently large $E \times B$ drift of the ripple trapped ions reduces the ion transport to the axisymmetric (“tokamak-like”) $\nu$-regime level. For this case, the neoclassical ion heat flux scales with $n_i^2 T_i^{5/2}$, and is sufficiently low to permit high ion temperatures. Within neoclassical theory, the $E_r$ is determined by the condition of ambipolarity of all particle fluxes, the impurity contribution is ignored in this context. Within the plateau collisionality regime, where $q_i$ scales with $n_i T_i^{5/2}$, the ion flux, $\Gamma_i$, is by a factor of $\sqrt{m_i/m_e}$ typically larger than the electron flux, $\Gamma_e$. In the LMFP regime at moderate $E_r$, the electron transport coefficients increase with $\frac{1}{2}$. This increase can lead to ambipolar $\Gamma_e = \Gamma_i$ fluxes and $q_i$ values well above the “tokamak-like” $\nu$-regime level. In the LMFP regime at lower collisionalities, multiple roots of the ambipolarity condition are expected [1]: the “ion root” at rather small $E_r$ values, and the “electron root” at large positive $E_r$ with significantly improved neoclassical confinement properties.

In low density ECRH discharges, the electrons have access to the deep LMFP regime ($T_e(0) \leq 4$ keV obtained in W7-AS). However, ICRH fails to achieve analogous conditions for ions. So far, efficient ion heating is only obtained by NBI and electron collisions. At low densities, however, this ion heating is not efficient which is consistent with the low $T_i$ values measured in ECRH discharges. At higher densities with efficient ion heating, the electron temperature drops to the $T_i$ level well below 1 keV [2]. This observation is a common feature of all stellarator devices. Only in the W7-A stellarator, the nearly perpendicular NBI was connected with significant fast ion orbit losses leading to strong radial electric fields which improved essentially the thermal ion confinement [3]: $T_i > 1$ keV at $n_e \simeq 10^{20}$ m$^{-3}$ were measured. In W7-AS, ion collisionalities well below the plateau regime were obtained in ECRH discharges (in combination with NBI) at moderate densities ($n_e \simeq 5$ to $8 \cdot 10^{19}$ m$^{-3}$). For these discharges the neoclassical ion heat transport with the ambipolar radial electric field is calculated based on measured profiles of $n_e$, $T_e$ (by Thomson scattering) and $T_i$ (by active CX analysis). For the complex magnetic field topology of W7–AS, the full neoclassical transport matrix is calculated by the DKES code [4]. The resultant transport predictions are compared with the experimental ion power balance. As the neoclassical ion heat flux is driven by the density gradient, the radial electric field as well as the ion temperature gradient, only the full flux $Q_i^{\text{neo}}$ is estimated and compared with the ion power balance. An experimental estimate of only the diagonal transport coefficient, $\chi_{i,\exp}$, is rather meaningless.
Table 1: The neoclassical predictions for the integrated ion heat and particle fluxes, $Q_i^{\text{neo}}$ and $\Gamma$, with and without ambipolar $E_{\text{gy}}$, respectively, and the upper limit, $Q_i^{\text{exp}}$, obtained from the experimental ion power balance. The neoclassical prediction for the central ion temperature, $T_{i,\text{neo}}$, is given as reference.

### Discharges at Low Ion Collisionality

Different heating scenarios have been explored in W7-AS: ECRH with on- or off-axis power deposition and in combination with NBI. In the case of central ECRH deposition, no matter if pure ECRH or combined heating, the $T_e$ profiles are peaked, whereas the $n_e$ profiles are flat or slightly hollow [5]. The collisional power transfer, $P_{e,i}$, turns out to be only moderate ($\propto n_e^2 (T_e-T_i)/T_e^{3/2}$). For off-axis ECRH deposition, however, the $T_e$ profiles are hollow with flat or slightly peaked $n_e(r)$ [5], i.e. there is no indication of “profile consistency”. The $T_e > 0$ found experimentally in the innermost region is consistent with the purely diffusive ansatz for the electron heat transport since $P_e^{\text{NBI}} + P_{e,i} + P_{\text{rad}} < 0$.

The discharges listed in Table 1 have been started by 70 GHz ECR (o-mode at $B = 2.5$ T) with the 140 GHz ERCH (2nd harmonic x-mode) added or taken over. In a later phase NBI has been added. For these scenarios the line-averaged density has been feed-back controlled. The purely NBI heated phase after switch off of the ECRH is characterized by increasing density without gas feed. In this phase, strongly peaked $n_e$ profiles have been found starting from the slightly peaked profiles in the off-axis ECRH cases, whereas the flat profile for on-axis ECRH remains flat. With increasing density, the temperatures decrease. The results shown in Table 1 have been obtained under separatrix conditions at $\epsilon = 0.53$ except the $\epsilon = 0.35$ discharge (limiter case, last line) at lower density.

### Ion Power Balance

On the basis of the measured $n_e$, $T_e$ and $T_i$ profiles, the ion power balance is performed. The collisional power transfer, $P_{e,i}$, depends on $T_e - T_i$ which at moderate $n_e$ is reliable within 30% in the main part ($r \leq 11$ cm, no $T_i$ data are available from active CX analysis at outer radii). For higher densities (about $10^{20}$ m$^{-3}$), the temperature difference is typically small, and the separated ion power balance analysis becomes unreliable. DEGAS code simulations of the distribution of neutrals clearly indicate that the thermal CX losses are of minor importance in the bulk part of the plasma, the radial decay length of the neutral density is typically of the order of 1 up to 2 cm
Fig. 1: $n_e$ (on the top), $T_e$ and $T_i$ profiles (in the center, solid and dashed lines, respectively) and the ion power balance (bottom plot, see text) for an off-axis ECRH discharge at 400 kW heating power.

Fig. 2: $n_e$, $T_e$ and $T_i$ profiles and ion power balance (see Fig. 1) for a combined heating discharge: off-axis ECRH (400 kW at 140 GHz, 320 kW at 70 GHz) and NBI (about 510 kW absorbed power).

[6]. In case of NBI, however, fast ion CX losses become more important with increasing slowing down time. Monte Carlo simulations of the NBI power deposition leads to an uncertainty of up to 30% in the ion heating, $P_i^{\text{NBI}}$, at averaged densities of about $5 \times 10^{19}$ m$^{-3}$. At lower $n_e$, the absorbed NBI power decreases substantially depending on the density of neutrals. Consequently, the estimation of the power transferred to the ions in the bulk part of the plasma should be accurate within 40% for the discharges listed in Table 1.

The local ion power balance is shown in Figs. 1 and 2. In both discharges, the ECRH power is deposited at about 7 to 8 cm (highly localized). The collisional ion heating in the pure ECRH case (140 GHz at about 400 kW absorbed power, 2nd line in Table 1) in Figure 1 is rather low (about 150 kW within 11 cm minor radius), the integrated heating power (solid line in the bottom plot) is about a factor of 3 larger than the neoclassical prediction with the ambipolar $E_r$ taken into account (dashed line). With $E_r = 0$, however, the neoclassical prediction (dotted line) exceeds the available power by a factor of more than 2. For the discharge shown in Fig. 2 (see 6th line in
Table 1) with the highest input power, the ion heating is rather efficient (about 410 kW within 11 cm). For these higher $T_i$, the ions are well within the LMFP regime. Here, the neoclassical ion heat flux with the ambipolar $E_r$ is also by a factor of 2 to 3 less than the available power which itself is by nearly one order of magnitude exceeded by the neoclassical prediction with $E_r = 0$. This trend is a common feature of all discharges listed in Table 1. At a minor radius of about 11 cm, the neoclassical ion heat flux is by the factor of 2 to 3 smaller than the value obtained from the ion power balance, the prediction with $E_r = 0$ even exceeds the total input power. This radius indicates roughly the boundary between the central confinement region with small thermal CX losses and the edge region with strong particle sources.

Furthermore, the ion temperature profile was estimated by integration of the ion power balance. Due to the strong non-linearity of the neoclassical ion heat flux with respect to $T_i$, errors induced by the boundary condition as well as by the omitted thermal CX losses at outer radii are strongly screened (e.g., the central $T_i$ is nearly independent on the boundary value). In case of neoclassical transport with the $E_r$ from the roots of the ambipolarity condition included, the neoclassical $T_i$ values exceed the experimental ones (see Table 1). Assuming $E_r = 0$ leads to central $T_i$ values of only 400 to 500 eV for all cases of Table 1. This result is fully consistent with the comparison of the experimental ion heat losses and the neoclassical predictions. The integrated ambipolar particle fluxes (also given in Table 1 for $r \approx 11$ cm) agree rather well with previous findings at lower $\epsilon$ [6] where DEGAS simulations are available. The NBI heating leads to a particle source strength of about $2 \cdot 10^{20}$ s$^{-1}$ per 1 MW input power. The $\Gamma_i$ values for $E_r = 0$ in Table 1, however, are much too large.

Conclusions

The neoclassical ion heat transport with the ambipolar radial electric field included is calculated based on measured profiles. These $Q_{\text{neo}}^i$ values turn out to be a factor of 2 to 3 smaller than those derived from the experimental ion power balance. This result being clearly outside of the experimental uncertainties is independent of the heating method, pure ECRH, pure NBI or combined, and holds for both on- and off-axis ECRH power deposition.

For the discharges with low ion collisionality, the calculated ambipolar $E_r$ corresponds to the “ion root”. So far, no indication for significantly improved ion confinement properties as predicted for the “electron root” in the LMFP regime is found. Ion temperatures up to 800 eV were obtained in these experiments at W7-AS with ion heating powers of up to 500 kW. So far, the high ion heat losses prevent the access to an improved LMFP confinement regime with higher ion temperatures.

The optimization of the magnetic field configuration of the proposed W7-X stellarator leads to an improvement of the neoclassical confinement properties [7]. Within the plateau regime, a reduction of the transport coefficients by a factor of about 6 is obtained, and the effective helical ripple in the LMFP is rather small. From this point of view, the access to an improved LMFP confinement regime should be easier for W7-X than for W7-AS.

References

EFFECTS OF PLASMA CURRENTS AND PRESSURE ON THE
MAGNETIC CONFIGURATION OF THE W7-AS STELLARATOR

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INTRODUCTION

A precise knowledge of the rather complex magnetic configuration of the W7-AS stellarator is an essential prerequisite for the evaluation and interpretation of experimental data. The transformation from real space to magnetic coordinates is based on equilibrium calculations (KW, VMEC). The KW code has been used up to now using pressure profiles as derived from Thomson scattering data. The VMEC code in the free boundary version is used to include the effect of the neoclassical bootstrap current, the ohmic and EC driven currents which are calculated on the basis of the measured temperature and density profiles.

MULTIPOINT THOMSON SCATTERING SYSTEM

The new single shot multipoint Thomson scattering system has gone into operation at the beginning of the current experimental period. With respect to the W7-AS magnetic configuration the laser chord has twofold symmetry, the equatorial plane and the "triangular" poloidal plane. Twenty radial positions along the laser beam are observed with a spatial resolution of 25mm. With open limiter aperture typically 12 channels are viewing the high field side of a peaked temperature profile and 8 channels the low field side so that the position of the plasma can be determined very precisely by means of the gradients. The scattered spectrum is analyzed by four spectral channels. Three different systems are employed accounting for the plasma centre, gradient and edge by appropriate positioning of the wavelength channels. This way the system is quite sensitive to changes in temperature as well as nonthermal contributions. A high precision is also achieved for small electron densities as the ruby laser has a pulse energy of 15J and the overall transmission of the detection system is about 40%. At the toroidal position of the laser beam the Shavranov shift has a maximum which results in an increased asymmetry of the pressure profile in real space. Thus one can sensitively test the transformation to the magnetic coordinate \(\psi\) which must result in symmetric profiles for the pressure \(p(\psi)\) as well as the temperature \(T(\psi)\).

For the present experimental period the old single point scattering system which resides in the "elliptical" poloidal plane is still in operation. So we have the possibility to check the toroidal symmetry of the magnetic configuration, too.

EQUILIBRIUM CODES FOR COORDINATE TRANSFORMATION

Up to now the transformations from real space to magnetic coordinates for the interpretation of experimental data on W7-AS have been taken from a database of equilibrium calculations by the KW-code\(^1\). Using a low-\(\beta\) expansion of the MHD equilibrium equations the KW-code can calculate iteratively free boundary finite-\(\beta\) equilibria. Nevertheless, it does not have the possibility to consider equilibria with a non vanishing toroidal net current density on magnetic flux surfaces like the bootstrap current, ohmic and EC driven currents. In contrast, the VMEC-code\(^1\), which solves the MHD equilibrium equations by an energy principle assuming nested magnetic flux surfaces, is able to take arbitrary toroidal current
profiles into account. Together with the NESTOR-code for the vacuum magnetic field it is possible to calculate free boundary equilibria with non vanishing toroidal current densities which may also result in a toroidal net current.

The input for a free boundary VMEC run consists of the W7-AS coil system with all coil currents and of the pressure and toroidal current density profiles. The pressure profiles taken are derived from the Thomson scattering data with ion contributions taken into account. The Thomson data also enter the neoclassical DKE5-code which is used to calculate the toroidal current density consisting of the electron bootstrap and the ohmic current densities.

**CONFIGURATIONAL STUDIES**

The reliability of the coordinate transformation will be demonstrated by applying it to measured \( T_e \) and \( n_e \) profiles. Since the calibration of the new Thomson scattering system is still under progress the presented results are somewhat preliminary. The systematic errors of \( T_e, n_e \) are still of the order of 5% and 15% respectively. Hence at the moment this analysis has also the character of a consistency check.

**Fig. 1** shows a typical example of electron temperature and density profiles as measured by means of the multipoint scattering system in a nearly net current free discharge with central ECR heating. The beam coordinate has been transformed into magnetic coordinates on the basis of the KW code and the profiles appear fairly symmetric.

**W7-AS profile database**

<table>
<thead>
<tr>
<th>shots</th>
<th>( T_e ) [keV]</th>
<th>( n_e ) ( [10^{20} \text{ m}^{-3}] )</th>
</tr>
</thead>
<tbody>
<tr>
<td>23500</td>
<td>1.6</td>
<td>0.7</td>
</tr>
<tr>
<td>23500</td>
<td>1.5</td>
<td>0.6</td>
</tr>
</tbody>
</table>

**date:** 18-MAR-93  **time:** 0.549 s

**ECRH:** \( P_{in} = 413 \text{ kW} \)

**\( a_{eff} = 16.0 \text{ cm for } z_{lim} = 31.0 \text{ cm} \)**

**Fig. 1:** Electron temperature and density measured by the multipoint Thomson scattering system. The discharge is ECR heated with central power deposition. Transformation to magnetic coordinate on the basis of KW code equilibrium calculation.

**Fig. 2** shows the effect of beta. Discharge no. 23691 represents a neutral beam heated plasma using two beam sources (750kW) where a central beata of \( \beta(0) \approx 1.5\% \) could be
Fig. 2 Two neutral beam heated discharges with a heating power of 0.75MW and 1.5MW for shots 23691 and 17916 and central $\beta$ values of $\approx 1.5\%$ and $\approx 2\%$ respectively. The first one has been measured by the multipoint system, the second one by the single point system. Coordinate transformation based on VMEC code. For shot 17916 the plasma had to shifted about 12mm towards increasing major radius.

Fig. 3 Thomson scattering profiles measured in an ECR heated discharge by means of the multipoint system (open symbols) and the single point system (full symbols). Transformation based on KW code plus an additional radial shift in the second case.
reached. The transformation in this case is obtained from the VMEC code. Both profiles and thus the pressure profile, too, show a remarkable symmetry which means that the beta dependence, i.e. the Shavranov shift, is well reproduced by the transformation. Discharge no. 17916 shows a neutral beam heated discharge using four beam sources (1.5MW) with $\beta(0) \approx 2.0\%$ as measured by means of the single point Thomson scattering system. In this case symmetry can only be achieved by introducing an additional shift of the plasma of $\approx 12\text{mm}$ towards larger major radius.

Fig.3 shows a direct comparison between a well centered profile of an ECR heated discharge as measured by the multipoint scattering system a second one measured by the single point system in the same discharge. Again the latter one had to be shifted by $\approx 12\text{mm}$ towards larger $R$ to make the two profiles coincide.

In the gradient region of the density profiles a small discrepancy is remaining between the two cases. Two reasons can be made responsible for that. On the one hand the gradients in real space are much steeper for the multipoint system than for the single point one so that a radial error occurs due to the averaging across the observation volume. On the other hand more data are necessary to determine the precise value of the outward shift in the case of the single point system.

DISCUSSION

In the position of the multipoint Thomson scattering system (triangular poloidal plane) the coordinate transformation from real space to magnetic coordinates based on the KW and VMEC code results generally in symmetric profiles for $\beta(0)$ values up to 2.5%. In the elliptical poloidal plane where the single point scattering system is installed an additional shift of about 12mm towards increasing major radius seems to be necessary to obtain consistent profiles in both toroidal positions. Possible reasons might in principle be perturbations of the magnetic configuration or the equilibrium codes are not appropriate in some detail. Further investigations are needed to clarify this situation.

CONCLUSIONS

The transformation to magnetic coordinates based on equilibrium codes has in principle been shown to be the proper tool to map diagnostics in different toroidal positions on each other. For a more detailed analysis the systematic errors of the multipoint Thomson scattering system have to be reduced and on this basis a more refined comparison of the two codes has to be done. Equilibrium calculations including strong toroidal net currents are still under investigation.

REFERENCES

KW:

VMEC:

NESTOR:
SIMULATION OF THE POLOIDAL ROTATION SHEAR LAYER FOR STELLARATORS

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Association EURATOM-IPP, D-85748 Garching, Germany
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INTRODUCTION

In the neoclassical theory based on the first order expansion of the distribution function, the radial electric field, $E_r$, is calculated by the roots of the ambipolarity condition of the local particle fluxes:

$$
\Gamma_e = Z_i \Gamma_i \quad \text{with} \quad \Gamma_{\alpha} = -n_\alpha \cdot \left\{ D_{11}^{\alpha} \left( \frac{n_\alpha'}{n_\alpha} - \frac{q_\alpha}{T_\alpha} \frac{E_r}{T_\alpha} \right) + D_{12}^{\alpha} \frac{T_i}{T_\alpha} \right\}
$$

with $\alpha = e, i$ (impurity ion fluxes are disregarded), $q_\alpha$ being the particle charge. In the particle flux densities, $\Gamma_\alpha$, the Ware pinch term ($\propto E_\parallel$) is omitted. In this context, additional "anomalous" contributions are assumed to be intrinsically ambipolar. For given density and temperature profiles, $E_r$ is estimated separately for each flux surface. As the neoclassical particle transport coefficients depend on $E_r$ (and quite differently for ion and electrons in the different regimes of collisionality), multiple roots of the ambipolarity condition can exist. Especially when both the electrons and the ions are in the LMFP regime three roots can appear [1]: the "ion root", $E_{r,i}$, at small $E_r$ values, and the strongly positive "electron root", $E_{r,e}$. An unstable root of the ambipolarity condition exists in between, while both the "ion" and the "electron root" are stable. At outer radii with higher collisionality, typically only the "ion root" with typically $E_{r,i} < 0$ can exist.

For the low density ECRH discharges (an example is shown in Figure 1), the ions with $T_{i0} \ll T_{e0}$ are in the plateau or in the beginning of the $\frac{1}{r}$-regime. For these conditions, the ion transport coefficients are only weakly dependent on $E_r$ if $E_r \ll E_{r,\text{res}} = \nu_{th} B r / R$. Above this resonance value $E_{r,\text{res}}$, where the poloidal component of $\vec{v} + \vec{v}_{EXB}$ vanishes, the "tokamak-like" banana orbits disappear and the particle transport coefficients are strongly decreased. This reduction of the neoclassical ion heat transport within the plateau collisionality regime was experimentally confirmed for NBI heated discharges in W7-A [2]. The fast ion orbit losses due to the nearly perpendicular NBI generated rather large, but negative $E_r$ (up to $-500$ V/cm). In the low density ECRH discharges, the existence of the "electron root" is based on this reduction of the $D_{kl}$ for $E_r > E_{r,\text{res}}$.

THE POLOIDAL ROTATION SHEAR LAYER

In the analysis of low density ECRH discharges with highly peaked $T_e$ profiles (see Fig. 1), only the "electron root" is typically found in the innermost part while only the "ion root" in the outermost part. A rather broad region with both roots exists in between. No
Fig. 1: $n_e$, $T_e$ and $T_i$ profiles of a low density W7-AS discharge with on-axis ECRH power deposition (320 kW, o-mode launch at $B \approx 2.5$ T) based on Thomson scattering, ECE and active CX measurements.

Fig. 2: The $E_r$-integrals of eq. (4) (dashed and dotted lines: for electrons and ions, respectively) within the range of multiple roots of the ambipolarity condition (1) for the discharge of Fig. 1. The limits $E_r$ and $E_r^e$ are given by the dotted line in Fig. 3 (on the left, also in Fig. 4).

A selection criterion with respect to both stable roots is defined so far. In principle, higher order expansions of the distribution function can resolve the problem within neoclassical theory [3, 4]. By these corrections, “non-local” fluxes being related to the shear of the poloidal rotation (i.e. differential terms in $E_r$) add to the purely local fluxes (depending only on $E_r$) leading to a diffusion equation in radius for $E_r$. In this picture, finite orbit effects (mainly of ions) lead to the coupling of neighboring flux surfaces with different roots of the first order ambipolarity condition (1).

A more general approach is based on thermodynamic arguments. The poloidal rotation shear layer, defined as the zone of the transition between the two stable roots, will develop satisfying the constraint of minimizing the dissipation of the poloidal plasma rotation energy. The generalized heat production of both the poloidal sheared rotation and the neoclassical particle transport has to be minimized [5]:

$$
\dot{Q} = \int_0^a \left\{ \dot{\eta} \cdot (E_r' - \frac{1}{r} E_r)^2 + \epsilon \int_0^{E_r} (Z_i \Gamma_i - \Gamma_e) \, dE_r \right\} \, r \, dr.
$$

The 1st term of the integrand is the generalized heat production rate due to the dissipation of the rotation energy, $\dot{q}^{\text{diss}}$, and the 2nd term due to the neoclassical particle fluxes, $\dot{q}^{\text{neo}}$. The detailed form of the factor $\dot{\eta}$ ($B^2 \dot{\eta}$ corresponding to a viscosity coefficient) has to be determined, e.g. by neoclassical theory ($\dot{\eta}$ depends on $r$, the plasma parameters and on $E_r$ itself). This ansatz fulfills the requirement that $\dot{q}^{\text{diss}}$ vanishes for co-rotation ($E_r \propto r$). Omitting the $\dot{q}^{\text{diss}}$ term, the local ambipolarity condition (1) is obtained by variation of $\dot{Q}$. The Euler-Lagrange differential form of this variational problem leads...
directly to the diffusion equation for $E_r$:

$$\frac{2}{r^2} \frac{d}{dr} \left\{ r^2 \hat{n} \cdot (E_r^* - \frac{1}{r} E_r) \right\} - \frac{\partial \hat{n}}{\partial E_r} \cdot (E_r^* - \frac{1}{r} E_r)^2 - \epsilon (Z_i \Gamma_i - \Gamma_e) = 0.$$  \hspace{1cm} (3)

In this form, the differential operator itself introduces a viscous particle flux, $\Gamma^{vis}$, driven by the poloidal viscous force which adds to the neoclassical particle fluxes: $\Gamma_e = Z_i \Gamma_i + \Gamma^{vis}$.

Following further neoclassical arguments, the broadness of this “shear layer” is determined by the deviation of mainly the ion orbits from the flux surfaces. For the radii on the “ion root” side, this size is of the order of a “tokamak-like” banana orbit. On the “electron root” side, however, the ion orbits remain much closer to the flux surfaces due to the very large $E_r$. Consequently, the broadness of the shear layer can be much smaller than the radial region where both stable roots exist.

The radial position, $r_{SL}$, of the “shear layer” defined by the minimum of $\hat{q}$ in eq. (2) can be estimated without solving the nonlinear diffusion equation for $E_r$. Assume that $\hat{q}^{diss}$ vanishes outside of the narrow “shear layer” and that $\hat{q}^{diss}$ depends only weakly on the value of $r_{SL}$. Then, the minimization of $\hat{q}$ with respect to $r_{SL}$ leads directly to $\hat{q}^{neo}(E_r^*) \simeq \hat{q}^{neo}(E_r^*)$. Consequently, the “shear layer” position is determined if the condition is satisfied:

$$\int_{E_r^*}^{E_r^{res}} (Z_i \Gamma_i - \Gamma_e) dE_r = 0.$$  \hspace{1cm} (4)

For the discharge of Fig. 1, both contributions to this integral are shown in Figure 2.

The diffusion equation (3) for $E_r$ is solved numerically with some simplifications. For the transport coefficients in the LMFP regime, an analytical LMFP regime model [1] with the effective ripple ($\epsilon_k$) fitted to DKES code [6] results is used [7]. The axisymmetric banana-plateau contributions of electrons are approximated by the Hinton-Hazeltine model [8], and of ions by the Kovrizhnykh model [9], respectively. The later one is generalized with respect to $E_r$: $D_{ki}^{-1} = const.$ for $E_r \leq E_r^{res}$ and $D_{ki}^{-1} \propto E_r^{-\alpha}$ for $E_r > E_r^{res}$ (with $\alpha = 2$ in Figs. 3 and 4) being used. With this simplified ansatz for the transport coefficients, the roots of the ambipolarity condition (1) show a similar feature as the solution based on DKES code results [7]. In order to simulate the effect of the deviation of the ion orbits from flux surfaces on the viscous particle flux, the $E_r$ dependence of the “viscosity coefficient” $\hat{n}$ was modeled analogous to the axisymmetric contribution to $D_{ki}$ for simplicity: $\hat{n} = const.$ for $E_r \leq E_r^{res}$ and $\hat{n} \propto E_r^{-\beta}$ for $E_r > E_r^{res}$. The solution of eq. (3) shown in Fig. 3 (with $\beta = 2$) shows a broader “shear layer” than the one of Fig. 4 ($\beta = 4$) for which $\hat{n}$ was decreased by a factor of 10. The broadness of the “shear layer” scales with $\hat{n}^{1/2}$. The transition from the “electron” to the “ion root” is in very good agreement to condition (4), $r_{SL}$ is indicated by the vertical line. Only within the shear layer region, a viscous flux is obtained which adds to the neoclassical fluxes of eq. (1).

SUMMARY

For low density ECRH discharges with highly peaked $T_e$ profiles, a radial region with multiple roots of the ambipolarity condition with only the neoclassical fluxes included is found. In this region the formation of a rather narrow shear layer in the poloidal plasma rotation is modeled leading to a unique solution for the radial electric field. The shear in the poloidal plasma rotation results in an additional viscous particle flux only within this region.
Fig. 3: The radial electric field (on the left; solid line: solution of eq. (3), dotted line: solution of the ambipolarity condition with multiple roots) and the particle fluxes (on the right; solid line: viscous flux, dashed line: $\Gamma_e$ and dotted line: $\Gamma_i$) in the shear layer region (for $\beta = 2$). $r_{SL}$ obtained from condition (4) is indicated (see Fig. 2).

Fig. 4: Analogous to Fig. 3, but with $\eta$ decreased by a factor of 10 and $\beta = 4$.

REFERENCES

Particle and Energy Transport in the ASDEX Scrape-Off Layer

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Introduction: The increasing maturity in edge/SOL modeling approaches within the IPP, and the crystallization of clear trends within edge data bases for ASDEX now begin to permit a more systematic exploration of the underlying physics behind edge phenomena observed in the past. This paper first documents the parametrical behavior of a number of edge/SOL/divertor quantities for ohmically- and NI-heated plasmas. Then, the results of parametric calculations with the B2 code /1/ in conjunction with a simple gas recycling model are compared to experiment. Hereby, the goal is to explore the systematic trends predicted for experimentally-measured quantities with the aims of (a) discerning the transport model most appropriate to the data, (b) elucidating the interdependencies of salient code output quantities, and (c) delineation of the direction of future code work of this nature.

Experimental Scalings: In the edge plasma of the main chamber, temperature ($T_{e39.4}$) and density ($n_{e39.4}$) values are available from the YAG Thomson scattering system at a point approximately one cm inside the magnetically-determined separatrix. Regressions involving $T_e$ and $n_e$ are based on this point in order to minimize scatter of the data. Extrapolations of $n_{e39.4}$ and $T_{e39.4}$ to the approximate separatrix position yields values about 2/3 of those at the point of measurement. The lithium beam probe delivers relative $n_e(r)$ profiles in the outer midplane. A reciprocating Langmuir probe in the divertor yields $n_{ed}(r)$, $T_{ed}(r)$ and equivalent power flux $q_{pp}$ profile widths $\Delta p$. Interferometric line densities $n_{eL}$ near the divertor throat as well as the divertor neutral gas density $n_{odiv}$ and $H_\alpha$ intensity are studied.

Table I presents regression fits for the quantities above from an (OH,NI) experimental data base in terms of those parameters expected to be relevant for edge/SOL physics. The "machine parameter" $B_1$ must also be used to obtain reasonable results. Of note is: (a) $q_{D}$ and $B_1$ are of major importance in the fit, and (b) among edge quantities there is no isotope effect.

In the outer midplane $\lambda_n$ is derived by fitting an exponential to the density profile over the first two cm outside the separatrix position $R_S$. $\lambda_n$ is typically given by $\lambda_n = q_{D}^{0.45}$ for $T_{e39.4} > 0.23 - 0.41$, $\lambda_n = 1.6 - 2.6 cm /2.3/$. The smaller $q_{D}$ exponent is more representative. However for $T_{e39.4} < 0.40eV$, corresponding to high-$q_{D}$ ohmic discharges, $\lambda_n$ abruptly changes its behavior: $\lambda_n(T_{e39.4} < 0.40eV) = q_{D}^{0.79} T_{e39.4}^{-0.51}$. Phenomenologically, the density shoulder seen in /2,3,4/ for R-R_S $> 2cm$ in normal regions of operation becomes characteristic of the entire SOL for low $T_e$ and high $q_{D}$.

$\Delta p$ in the divertor scales as $\Delta p_{TEd} = 0.58 \pm 0.013 q_{D}^{0.53 \pm 0.025}$ (R=0.9) for an ohmic deuterium data base ($\Delta p=1-4 cm$). The general trend with decreasing $T_{ed}$ as the density limit is approached is the $T_{ed}$ profile becomes very flat in the range 5-8eV, whereas $n_{ed}(r)$ maintains a somewhat peaked profile near the separatrix. A shoulder in the $q_{pp}$ profiles is also seen, which dominates $\Delta p_{TEd}$ for low $T_{ed}$ and high $q_{D}/5/$.

In previous work /2,3/ an initial effort to deduce the perpendicular diffusion coefficient $D$ in the SOL from the formulation "$D=\lambda_n^2/\pi q_{D}$" under the assumption of a constant mach number led to the result $"D= T_{e39.4}$/ for a wide variety of conditions for ohmically- and NI-heated plasmas. This result is summarized for an OH data base in fig.2.

Braams Code Results: The "$D= T_{e39.4}$/" relation of above is predicated on a primitive model and should be checked against complete B2-EIRENE /1/ calculations. However, the approach discussed here employs a simplified impurity and neutral gas model, with a single-fluid B2 version, in order to gain a feeling for leading terms. Heat transport along field lines is taken to be classical with a flux limit of 0.2eV/$T_e$. For perpendicular transport, three models are examined: (A) $D=\chi_q^2/3=0.5, 1, 1.5 m^2/s$; constant over the SOL, (B) $D=\chi_q^2/3=$
The code input parameters are varied over typical operational ranges; the results are summarized in table II in the form of regressions for the quantities listed in table I, whereby the B2 input quantities \( n_{\text{es}}, P_{\text{sol}}, q_{\text{a}}, \) impurity concentration \( n_{\text{imp}}/n_0 \) and \( D_0 \) are used as the regression parameters. This statistical approach has the advantage that the validity of individual transport models, as well as inadequacies of other model assumptions, can readily be examined.

Classical drifts or an inward pinch are not considered. Only one equilibrium grid is used with a spatial extent of 1.8cm inside \( R_{\text{cs}} \) and 2.6cm outside \( R_{\text{cs}} \) at the outer midplane. The code operates with feedback loops so that \( n_{\text{es}} \) and \( P_{\text{sol}} \) can be specified; power is divided between electrons in the ratio 2:1. \( Z_{\text{eff}}=3 \) is taken for all calculations. The neutral gas model assumes the \( n_{\text{odiv}} \) profile -motivated by DEGAS/EIRENE calculations- is invariant in the divertor (reasonable for the \( n_{\text{L}} \) values encountered in the plasma fan, see table I). The absolute value of \( n_{\text{odiv}} \) is calculated from the divertor particle balance. \( n_{\text{es}} \) is regulated by balancing all ion losses from the SOL and divertor by a poloidally homogeneous gas puff in the main chamber/6/. Other parameters in the code are selected to give close agreement with complete B2-EIRENE runs for a generic situation.

Fig.1 illustrates B2-generated \( n_{\text{e}}, T_{\text{e}} \)-profiles for \( n_{\text{es}}=1\times10^{13} \text{ cm}^{-3} \) and \( P_{\text{sol}}=0.6 \text{MW} \). For \( R_{\text{cs}}\leq2 \text{cm} \) the differences in the three diffusive models are minimal. The \( 1/n_0 \) variation has the pleasing aspect that a shoulder is produced for \( R_{\text{cs}}\geq2 \text{cm} \), as seen in experiment. But, the shoulder is always there, not scaling in the proper way. For \( T_{\text{e}}(r) \) the conv. case is characterized by rather high \( T_{\text{e}} \) at the separatrix with a very steep falloff in the SOL.

Instead of a point-by-point comparison between experimental and code regressions, the overall trends will be reviewed. The inclusion of \( D_0 \) in table II as regression parameter demonstrates the sensitivity of changes in \( D_0 \) to the quantity under question. Quite generally, the experimental \( q_{\text{a}} \)-scalings are not duplicated in the calculations. There is good agreement for \( H_{\text{c}} \) in all cases. The code-predicted \( P_{\text{sol}} \) dependence of \( n_{\text{odiv}} \) is absolutely not found in experiment. The density dependence for \( \Delta_0 \) and \( \lambda_{\text{n}} \) predicted by case B is in strong contrast to experiment. Case C gives an extremely strong relationship between \( T_{\text{es}} \) and \( n_{\text{es}}, P_{\text{sol}} \) again in strong contrast to experiment. Case C does has the correct \( T_{\text{ed}} \) exponent for \( \Delta_0 \). However, as mentioned above, the \( q_{\text{a}} \)-profile for low \( T_{\text{ed}} \) develops a prominent shoulder which surely strongly contributes to the \( T_{\text{ed}}, q_{\text{a}} \)-dependence seen in experiment, and none of the models produce a shoulder in the correct sense. From these considerations, case A is the preferred form of diffusion coefficient simply from the standpoint that it nowhere strongly violates experimental behavior as do B and C for certain quantities.

For case A, the fit for \( \lambda_{\text{n}} \) is obtained only if values \( n_{\text{es}13}>0.8 \) are excluded from the data base. For \( n_{\text{es}} \) below roughly this limit (i.e. in the direction of higher \( T_{\text{es}}, T_{\text{ed}} \)) \( \lambda_{\text{n}} \) begins to decrease with temperature. This behavior is correlated with the Mach number: centered around \( T_{\text{es}}=35 \text{eV} \), \( M \) augments rapidly with increasing or decreasing \( T_{\text{es}} \) such that, heuristically, \( \lambda_{\text{n}}=(D_0/T_{\text{ed}}^{0.5})-(D_0/LMC_{\text{e}}^{0.5})-(D_0/LMT_{\text{e}}^{0.5})^{0.5} \) would be expected to diminish, which is indeed observed in the calculation. Thus, the experimental scaling of \( \lambda_{\text{n}}^{-0.23-0.41} \) comes about not to the diffusion coefficient scaling as \( T_{\text{es}39.4} \), but due to the variation in \( M \) with \( T_{\text{es}} \). Fig. 2 illustrates this point, where \( "D"=\lambda_{\text{n}}^{2/3}C_{\text{ed}} \) is used to derive a "diffusion coefficient" for case A (with \( D_0=0.5,1.1,5 \text{ m/s} \)) in the same way as the experimental data base. For each \( D_0 \) a set of points arises which behave very much as the experimental data set in both a relative and absolute sense. Going further, keeping in mind the very rough \( T_{\text{es}39.4}^{-1.5} \) relationship, the experimental set is compatible with \( D_0=0.5-1 \text{ at high } T_{\text{es}}, \) and around \( T_{\text{es}39.4}^{-40 \text{eV } D_0=1-1.5} \) is more appropriate.

Note that in fig. 3, the clean variation of "D" with \( T_{\text{es}} \) comes about only because \( q_{\text{a}} \) has been held constant, and the data set restricted to \( n_{\text{es}13}>0.8 \). The point scatter arising when these limitations are lifted is not compatible with experiment, where with \( q_{\text{a}}=2.5-5.9 \) and high temperatures, the scaling for "D" continues. Further code work is necessary here.

Overall, in terms of absolute numbers, the B2 results are satisfactorily in agreement with experiment with respect to \( T_{\text{es}} \) and \( n_{\text{es}}T_{\text{es}} \), keeping in mind that \( T_{\text{es}39.4}, n_{\text{es}39.4} \) when extrapolated to the separatrix are roughly \( 2/3 \) of the values given in table I, and that no
attempt was made to model any particular shot in detail. $T_{ed}$ is predicted to be too high, as $T_{ed}>35eV$ is seldom found in experiment. On the other hand, for such high temperatures the impurity content of the plasma will increase, which would exacerbate excessive $T_{ed}$ values if included in the model. $n_{el}$, $n_{odiv}$ and $n_{ed}$ are all about a factor of two too low; increase of recycling in the model will help to alleviate this inconsistency.

**Conclusions and Discussion:** Case A, D = $\chi_e / 3$, is most often closer to experiment, including the fit to $\lambda_n$. A D-1/$n_e$ scaling is of interest, since indications on the stellarator W7-AS are that such behavior prevails.\textsuperscript{[7]} Nonetheless, the $n_{es}$ scalings for $\lambda_n$ and $\Delta_p$ implicit for the model are simply not present. For example, during a density ramp, $n_{es,13}$ might vary from 0.5 to 2.0; the D(1/$n_e$)scaling $\lambda_n-n_{es}$\textsuperscript{-0.51} requires that $\lambda_n$ decrease by a factor of two, i.e. from two to one cm. This is far outside measurement error. The tendency of the convective model is for $T_{es}$ to respond too sharply to changes in $n_{es}$ and $P_{sol}$.

The strong $q_\alpha$, $B_t$, $I_p$ variations of some experimental observables might be an indication that transport is related to these global parameters. It is interesting, for instance, to find for cases where a pure $q_\alpha$-scaling is present (i.e. not in combination with $B_t$), that the assumption D-$q_\alpha$ leads immediately to the correct $q_\alpha$ dependence for $\Delta_p$ and $\lambda_n$. Such a scaling will also presumably allow the code to more correctly duplicate the $q_\alpha$-scaling of the density limit - without this caveat the density limit is largely unaffected by $q_\alpha$ variations.

The code results can identify combinations of experimental parameters suitable to extract D. For case A, D $\sim \lambda_n$\textsuperscript{1.99}$r_{0.075} T_{es}$\textsuperscript{1.65}$0.035$; that is, $\lambda_n$ is directly proportional to D as expected from simple 1D models, but the exponent of $T_{es}$ is exactly opposite to that normally assumed. This arises from the variation of mach number, both with $T_{es}$ and $q_\alpha$.

Finally, the presence of a shoulder in $n_{el}(r)$ and $q_\alpha(r)$ cannot be explained within the code, other than to assume a diffusion coefficient in excess of Bohm in the low-temperature region where the shoulder occurs. If transport is effected by ExB-driven turbulent eddies, of the size of $\lambda_n$ and with correlation times of the order of parallel equilibration times-for which some evidence exists \textsuperscript{[8]}, then during one rotation the plasma may drift outwards such that the effective value of the anomalous diffusion coefficient may exceed Bohm. One expects, then, $B_t$ to be correlated with quantities associated with the SOL. The regressions of table I, where $q_\alpha$ and $B_t$ have been used as principle components, indicate core/edge quantities are roughly proportional to $B_t$, whereas divertor parameters are inversely proportional.

<table>
<thead>
<tr>
<th>$T_{39.4} \text{ const.}$</th>
<th>$n_{39.4}$</th>
<th>$P_{sol}$</th>
<th>$q$</th>
<th>$B_t$</th>
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**Table I:** Regression fits for experimental quantities. The star indicates the fit is with $T_{ed}$ not $n_{39.4}$. See the subtitle of table II for more details. The div. data contains only an $n_{el}$-P$_{ni}$-scan.

**References:**


**Acknowledgement:** One of the authors S.F. is financially supported by a project, headed by Prof. H. Winter, of the Friedrich Schiedel Stiftung für Energieforschung.
Fig. 1 Above: $n_e$ and $T_e$ profiles for the three diffusion models. $D$ = (A) const.,(B) $1/n$, (C) convection only. $n_{34S}=1.0$, $P_{sol}=0.6$ MW.

Fig. 2 Right: $\lambda_n^{2/\tau_{ll}}$ (for mach no. = 0.3) vs. $T_e$ for an experimental OH data base (top), and the D-const. part of the B2 data base (bottom).

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Table II: Fits for the B2 data base. $T_{es}$=$T_{ed}$ for the $\Delta_{p}$ regression, and $T_{es}$=$T_{ed}$ for the $\lambda_{n}$ regression. In both tables, exponents are given for fits of the form: const. $x n_{es}^{a} P_{sol}^{b} T_{es}^{c}$, $n_{es}$, $P_{sol}$ [MW]; $T_{es}$ [eV]; $n_{es}$ [10^{13} cm^{-3}]; $\lambda_{n}$ [cm]. $R$ is the regression coefficient. Range = the min. and max. values of the parameter range.
OPEN MAGNETIC SURFACES AND RESONANT TOPOLOGY
IN THE SEPARATRIX-DOMINATED BOUNDARY REGION
OF THE W7-AS STELLARATOR

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1. Introduction

The boundary of W7-AS for \( \varepsilon \approx 1/3 \) is defined by the contact with two up-down limiters. Smooth flux surfaces extend deep into the SOL, and the limiters map into large-size flux bundles of homogeneous connection lengths. For this topology, a radial 1D transport model has been developed and used to derive radial profiles and density scaling of the diffusion coefficient in the limiter dominated SOL [1,2]. At \( \varepsilon \geq 0.5 \), the boundary topology is totally different and exhibits open, divertor-like field structures which are responsible for highly non-homogeneous recycling and wall load distributions [3]. A comprehensive understanding of the plasma transport and recycling in this region is needed, for example, to optimize passive and active methods of particle and impurity control, to clarify the effects of the boundary conditions on the main plasma performance and to explore the divertor potential of W7-AS. Evaluation and correlation of local experimental data are more difficult in this open topology, as it cannot be parametrized by standard magnetic coordinates.

2. Modelling open magnetic surfaces outside the LCMS

The Poincaré plots without and with a very small cross field diffusion (\( D = 4 \cdot 10^2 \text{cm}^2/\text{s} \), \( v_\| = 3 \cdot 10^7 \text{cm/s} \)) are shown in Figs. 1 and 2 for a typical separatrix-dominated W VII-AS configuration at \( \varepsilon(a) = 0.52 \). A radial sequence of three topologically different zones can be distinguished in Fig. 1:

a) central region with closed flux surfaces bounded by the LMCS,
b) an intermediate region including the \( \varepsilon = 10/19 \) island chain,
c) strongly diverted open field structures ("island region" and "private region") beyond the separatrix of the main \( \varepsilon = 5/9 \) "natural" resonance.

Plasma confinement beyond the LCMS is not likely to be affected by the 10/19 island chain, as the associated radial component of the parallel transport is generally much smaller than the cross field transport. In the open zone outside the X-point the diffusion plots (Fig. 2) show regular patterns which suggest a laminar parallel motion of the outflowing plasma.

This indicates that a radial ordering of the field structures, as needed both for correlating local measurements and for modelling plasma transport, should be feasible. By using the intersection points of a field line with a symmetry plane, smooth, field line spanned surfaces are obtained outside the LCMS by using the following procedure.

The distance between two subsequent intersection points is a measure of the poloidal progression \( \Delta \) of a field line after a poloidal turn, i.e. after \( m \) toroidal field periods (\( \varepsilon = 5/m \)). By mapping along the torus a poloidal line element (PLE) connecting the
two points a field line ribbon is obtained, which intersects the symmetry planes at m distinct poloidal positions. In the case of a closed magnetic surface, these ribbon branches overlap each other after a number of poloidal turns depending on 2. The overlapping constrain makes the PLE unique. In the case of open field structures, however, the branches do not overlap and the PLE has a certain degree of arbitrariness.

The ordering quantity we adopted to obtain unique surfaces from the given intersection points is the poloidal connection length, \( L_p \), which is the characteristic field line length associated with \( \Delta \). The variation of \( L_p \) is smooth along both the radial coordinate and the PLE. Stellarator symmetry and flux conservation imply that \( L_p \) is a symmetric function of \( \psi - 1/2 \), \( L_p(\psi) = L_p(1 - \psi) \), where \( \psi \) is the normalized flux enclosed by two adjacent PLE and is used as poloidal coordinate (Fig. 3). Furthermore, at \( \psi = 0 \) both \( L_p \) and \( \partial L_p / \partial z \) are known. Within these constrains, \( L_p \) is chosen as smooth interpolation between the LCMS and the separatrix of the 5/m resonance. In the open regions, the \( L_p \) are smoothly extrapolated from the limiting values along the separatrix. This procedure can be seen as a coordinate transformation \( R, z \leftrightarrow R_0, \psi \), where \( R_0 \) is the label of the surfaces. The resulting surfaces are not dependent on material structures or transport coefficients.

A sequence of open magnetic surfaces obtained in the "private region" with the described method is shown in Fig. 4. The surfaces are smooth and dense (within numerical resolution) over the whole radial range up to the wall. In this radial position the field line diversion is very strong, increasing by about an order of magnitude over a poloidal turn. The cross sections of three such surfaces in the two planes of symmetry are shown in Fig. 6. They well reproduce the patterns shown by the diffusion plot, Fig. 2. First results covering the open island region also show smooth nested patterns within the range of relevant field line lengths. The maps are in good agreement with the isolines of the ion saturation current of 2D Langmuir probe array measurements at the same \( \varepsilon \) and fully retracted limiter (Fig. 5). (A reduced limiter action is still detectable in a slight up-down asymmetry of the \( I_s \) contours.) Furthermore, the strong poloidal variation of the radial distance between the surfaces within the open island region is confirmed by density decay length measurements from two fast reciprocating Langmuir probes placed at different poloidal positions [4].

3. Resonant topology for modelling plasma transport

Solving parallel transport along the field lines in the described open structures is difficult for lack of symmetry. However, a regularization of the structures can be achieved without grid interpolation by splitting the parallel motion into a resonant component at \( \tau_{res} = 5/m \) and a poloidal component on \( \Phi = \text{const} \) along the given surfaces. This corresponds to splitting the magnetic field into a resonant component and a poloidal perturbation, both lying on the surfaces:

\[
\mathbf{B} = \mathbf{B}_{res} + \mathbf{B}_{pol}; \quad \mathbf{B}_{pol} = \tau_{res} \nabla \Phi \times \nabla \tilde{\Psi}
\]

where \( \tilde{\Psi} \) is the flux enclosed by the ends of all the PLE from the X-point to the given open surface. \( \mathbf{B}_{pol} \) is additive or subtractive depending on the position of the surface with respect to the X-point of the main edge resonance.

The resonant flux tubes enclosed by two adjacent surfaces are obtained at n equidistant toroidal positions by poloidally shifting the original flux tubes back towards resonance
at a constant rate of $\Delta \tilde{\Psi}/n$. The last shift then leads to a poloidal closure of the flux tube.

Fig. 6 shows the intersections of the trivially resonant $X$-point and of three other resonant field lines with the two planes of symmetry. In the special case of a closed island, the shifts correspond to a compensation of the "island $e$".

In a magnetic flux system rotating with a resonant $\star$, the resonant field lines are fixed points at any poloidal cross section. Gradients of the plasma parameters along this lines are expected to be smaller than along the original field lines, the net radial excursion associated with the diverted structures being eliminated by the shifting procedure. (In a tokamak divertor such gradients are trivially zero everywhere by axisymmetry). This property makes resonant flux tubes particularly well suited for correlation of experimental data obtained at different toroidal and poloidal positions.

The described resonant topology can be used to define a comfortable radial coordinate for plasma transport modelling. In fact, any surface enclosing any ensemble of resonant flux tubes is a "resonant flux surface". This property can be used to smoothly extend the flux surfaces beyond the LCMS, shaping them according to the position and form of the target plates. The result is a $X$-point free set of nested surfaces ranging from the core up to the outermost zones of the boundary [5]. Such a system would, of course, be common to both the plasma and neutral gas transport models.

References

[5] F. Sardei et al. (1993), 9th IAEA Int. Workshop on Stellarators, Garching, Germany

Fig. 1: Poincaré plots of the vacuum W7-AS configuration at $\epsilon = 0.52$, showing the closed flux surfaces, the $\epsilon = 10/19$ island chain and the open field structures at the boundary.
Fig. 2: Poincaré plots of the vacuum configuration as in Fig. 1, but with small cross field diffusion.

Fig. 3: Schematic drawing of two poloidal line elements connecting the respective intersection points ($\psi = 0$ and $\psi = 1$) of two field lines covering a full poloidal turn. The poloidal connection length function $L_p(\psi)$ is symmetric about $\psi = 1/2$.

Fig. 4: Sequence of modelled open magnetic surfaces outside the X-point (cross section in the "triangular" plane of symmetry).

Fig. 5: Isolines of ion saturation current from 2D Langmuir probe array measurements.

Fig. 6: Cross sections of 5 open magnetic surfaces covering different regions beyond the LCMS. The labels on the patterns represent resonant field lines.
On the Diffusive Nature of W7-AS Transport

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Introduction

Particle and energy transport in W7-AS have many aspects which are in qualitative agreement with diffusive transport mechanisms. Unlike in tokamaks, the density profiles are flat in the source-free region and for the electron temperature no profile resilience is observed (see Fig. 1). A strong dependence of transport on the temperature gradient could be ruled out since agreement between the transport coefficients from steady state and perturbative studies has been observed both, in absolute value and parameter dependences [1]. Furthermore, W7-AS transport was found to be compatible with a gyro-Bohm-like and, hence, local transport model [2]. Similar to tokamaks, however, there are problems in relating the local transport coefficients to local plasma parameters. Especially the degradation of confinement with heating power cannot easily be connected to a local quantity. This difficulty is the subject of this paper.

Experiment

To enlighten the possible dependencies of the transport coefficient on heating power \( P \) or electron temperature \( T_e \), steady state power scans are compared with the transient behaviour during energy increase and decay after switch-on or -off of a large amount of heating power. The experiments were performed on purely ECR heated (ECH) plasmas at low densities \( n_e(0) = 1.5 \times 10^{19} \text{m}^{-3} \). Since the electrons are well de-coupled from the ions, this allows a clear separation of the electron heat diffusivity \( \chi_e \). The heating power ranges from 0.2 to 0.6 \text{MW}. Detailed analysis of the power balance in transient phases is much easier in ECH plasmas than in beam heated plasmas, where the fast ions continue to heat the plasma for a slowing down time when the beams are already shut off. The slowing down time of the non-thermal electron population is shorter than 0.1 \text{ms}. Hence, we can begin with the analysis of the transient phases already 1 \text{ms} after the switch-on or -off of gyrotrons. The power balance analysis rely on density profiles from the single pulse Thomson scattering system and the electron temperature profiles from a ECE radiometer with a time resolution of 0.1 \text{ms}. The comparison of Thomson and ECE measurements for \( T_e \) are satisfactory and the density profiles, which have also been measured at some time points during the transient phases do not change in the time interval used for the analysis.

The power balance is done with 2 independent methods: (i) we use a time slice analysis transport code to calculate an experimental \( \chi_e \) and include the dynamics in taking the time derivatives of the \( T_e \) profile into account and (ii) we use a time dependent simulation code to predict the temperature profile evolution using different model assumptions for
Figure 1: Electron density and temperature profiles together with the sources of particles and energy, respectively. As one would expect from a diffusive transport model, the gradients appear outside and in the source regions whereas the source-free regions show flat profiles.

The transport coefficient. In Fig. 2 electron temperatures and transport coefficients from the time slice code of stationary and transient phases (from 1 to 3 and back to 1 gyrotron) of a discharge are depicted. The result is surprising: already 1 ms after the change in heating power the transport coefficient switches to the level expected for steady state on the modified power level. It then stays unaltered while the temperature profile still evolves. The conclusion would be that heat transport is not related to $T_e$. It must be stressed, however, that this result would be altered if much less than 100% of the heating power was absorbed in the plasma.

For more time points, the result is shown again in Fig. 3. Temperatures and transport coefficients are now radially averaged in the range $0.35 \leq r/a \leq 0.85$. The transport coefficient changes with the applied heating power while the evolution of $\chi_e$ is clearly decoupled from the evolution in $T_e$.

We have used the time dependent code to investigate 3 model assumptions for the transport coefficient in more detail: (i) we simulate the transient phase with steady state values for $\chi_e$ obtained from the phase prior to the change in heating power, (ii) we use the values which will be achieved when steady state will be recovered in the phase with altered heating power and (iii) we start with $\chi_e$ from the phase prior to the change in $P$ and modify it with a linear $T_e$ dependence. (i) is relevant for comparisons in the early phase, where $T_e$ has not yet much evolved, (ii) is for verifying our conclusions from the time slice code analyses and (iii) is relevant for comparison to theoretical models with $T_e$ dependences. Although there is bench-marking of the codes still to be done, the results for 2 radii ($r/a = 0.4$ and $0.6$) are consistent with our previous conclusions: only model (ii) with $\chi_e$ adjusted to the modified power level describes all phases satisfactorily. Model (i)
Figure 2: ECE electron temperature and transport coefficient profiles from the time slice code before and after (1 and 4 ms) switch-on (left) and switch-off (right) of 2 gyrotrons. \( \Delta t \) gives the time difference to the change in heating power (0.2 and 0.6 MW).

Figure 3: Radially averaged electron temperature \( \Delta T \) and transport coefficient \( \chi_e \) as a function of time relative to the switch-on and switch-off of the 2 additional gyrotrons. The heating power is indicated by a box diagram. It was 0.2 and 0.6 MW in the 1 and 3 gyrotron phases, respectively.
Figure 4: Results from a simulation code compared to the time evolution of $T_e$ after switch-on and -off of 2 additional gyrotrons. For $\chi_e$ used in the code, 3 different assumptions have been made: dashed curve: $\chi_e$ remains unchanged, dotted curve: starting from the steady state value $\chi_e$ is modified as $\sim T_e$; solid curve the steady state value of $\chi_e$ at the new power is used. The steady state $\chi_e$ profiles are taken from the time slice code. The ECH power decreased during the 3 gyrotron phase by 10%.

predicts in both early phases a faster evolution of $T_e$ than is measured. The $T_e$ dependent model (iii) seems to be appropriate for the decay phase but it fails in the simulation of the phase from 1 to 3 gyrotrons.

Conclusions

The detailed investigation of the continues evolution of the electron temperature in transient phases after the change of heating power is a new method to sensibly study the dependence of the transport coefficient on temperature. An important assumption, which is supported by ray-tracing calculations, is that most of the ECH power is absorbed in the plasma. The results would be drastically altered if the absorption was of the order of 50%. The result is that the transport coefficient reacts in times as short as 1 ms to a change in heating power. The change is much faster than the modification in $T_e$ or the temperature gradient. The conclusion would be that the anomalous transport does not depend on the electron temperature. Before drawing this conclusion firmly, however, further experiments at different plasma parameters and with different steps in the heating power have to be performed and a bench-marking of the 2 different codes applied must be done.


Stability of Neutral Beam driven Alfvén Eigenmodes in the WENDELSTEIN W7-AS Stellarator

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Introduction - During neutral beam injection (NBI) pronounced coherent MHD activity occurs at low and medium $\beta$, which is driven by energetic beam ions. These modes are discussed in terms of marginal stable global Alfvén eigenmodes (GAE)$^1$, which are destabilized by beam wave resonances of the fast circulating particles. Global Alfvén waves have been identified experimentally in large tokamak devices as toroidal Alfvén eigenmodes (TAE)$^2-6$, which are excited inside toroidicity induced gaps of the shear Alfvén continua. In the partially optimized stellarator WENDELSTEIN W7-AS $^7$ ($R = 2$ m, $a \approx 0.17$ m, 5 field periods, modular stellarator coil system) low order rational values of the rotational transform $\tau = l/q$ can be avoided because of the weak magnetic shear. Toroidicity induced gaps involving coupling between adjacent poloidal mode numbers $m$ and $m+1$ (with same toroidal mode number $n$), therefore, can not occur. However, gaps below the edge of the low order shear Alfvén continua with $m = 2,3,4,5$ and $n = 1,2,3$ are typically formed. Discrete global Alfvén eigenmodes (GAE) with frequencies $\omega_{\text{GAE}} \leq k_{\|}v_A = \omega_A$ are predicted in the Alfvén spectrum$^8-10$ ($k_{\|} = (m-\tau-n)/R$, $v_A = B/\sqrt{4\pi n_i m_i}$).

Characteristics of NBI driven mode activity - In W7-AS two almost tangential beam lines (co- and counter-direction, $P_{\text{NBI}} < 1.6$ MW, 45 keV hydrogen beam) are used to heat an ECRH target plasma and to achieve densities above the 70 GHz cutoff and higher plasma pressure. At intermediate densities ($n_i \approx 1 \cdot 10^{20}$ m$^{-3}$) coherent mode activity with typical frequencies in the range of 20...40 kHz is detected by the soft X-ray diode arrays, Mirnov coils and electron cyclotron emission (ECE) diagnostic. In fig. 1 the time evolution of global plasma parameters for a typical case is shown. There is an abrupt onset of the modes after the transition to degraded confinement connected with changes of the plasma equilibrium causing eventually peaked plasma profiles. The fast particle drive of the observed modes has been inferred from their fast decay ($< 300$ $\mu$s) compared with the slowing down time ($\approx 5$ ms) of the fast injected ions and the energy replacement time ($\approx 10$ ms) after switch-off of NBI (Fig. 2). The effect must be due to circulating particles, since the fraction of trapped particles is small because of tangential injection. In most cases single sharp lines are observed in the frequency spectra, but frequently 2 or 3 lines are present, which can have a relative shift of up to 30%. The rotation of the modes, as deduced from the X-ray and the Mirnov signals, is in the ion diamagnetic drift direction. This is opposite to the rotation of the usually observed pressure and current driven MHD modes. The rotation therefore changes its direction, if the magnetic field is reversed. The change from co- to counter-injection does not influence the direction of mode rotation. The perturbation propagates with its own real frequency, since the fluid drift velocities are typically slower as deduced from measured pressure profiles and from low frequency MHD activity.
**Relation to global Alfvén modes** - The frequencies depend on the magnetic field and the density and scale roughly with the mean Alfvén velocity. The full energy hydrogen injection velocity $v_b$, marginally reaches the Alfvén velocity $v_A$ at low field (1.25 T), typically one finds $v_b/n_A = 0.35...1.0$. The modes have a global radial structure as derived from X-ray and ECE measurements centered at $r/a = 0.4...0.7$ and extending over about 1/2 of the plasma cross-section. The modulation is even seen in the $H_\theta$ signals from the plasma edge. The radial displacement vectors corresponding to the smaller frequency peaks sometimes show a radial node, indicating higher radial mode numbers.

For a number of different cases the cylindrical shear Alfvén continua $\omega_A = k q v_A$ have been calculated in order to relate the observed frequencies to the upper gap frequencies. An example for the case of a $(m,n) = (3,1)$ mode at 36 kHz is shown in fig. 3. Experimentally deduced profiles of plasma density and rotational transform have been used to calculate $v_A(r)$ and $k_l(r)$. The wave particle resonance condition is approximately given by $v_p = v_b = \omega_A k_{lm}$ where $v_p$ is the fast particle parallel velocity and the subscript m denotes the poloidal mode number used in $k_l$. Due to the poloidal variation of the magnetic drift velocity in a toroidal system also sideband drive occurs for $v_b = \omega_A k_{lm+1}$. This leads to resonance velocities well below the full energy value (fig. 3 top). The condition can be also given in terms of a resonant rotational transform $\tau_{res} = n/(m+1) + \Delta \tau$ (where $\Delta \tau = R/v_b \omega_A/(m+1)$ and $l = 0, \pm 1$). The effect of $\Delta \tau$ is not relevant in a high shear system, but in W7-AS it crucially determines, which part of the fast particle velocity distribution is resonant with the GAE. There is also a different effect for co- and counter-injected ions, since their drift orbits are displaced in opposite radial direction, and hence $\tau$ is changed accordingly. The horizontal bars in fig. 3 at the expected GAE frequencies indicate the estimated radial range, where beam wave resonances can take place with velocities down to 1/3 of the full energy injection speed ($v_b = 1...3 \times 10^8$ cm/s) via sideband resonances.
In fig. 4 a case is shown, where a GAE (30 kHz, m=5, n=3) and a low frequency (6 kHz, m=2, n=1) pressure driven mode are excited at the same time. The latter does not decay immediately after NBI switch-off, and rotates in opposite direction (electron diamagnetic drift). The pressure driven mode indicates the rational surface \( \tau = 1/2 \) inside the plasma. There, the \((m,n) = (2,1)\) GAE cannot exist, because no gap is formed for this particular shear Alfvén continuum \( (k = 0)\). Therefore, the most likely GAE mode appears to be the \((5,3)\) mode, which is observed in the interior region consistent with the calculated continua shown in fig. 5.

The destabilization of the global Alfvén waves has been explained by inverse Landau damping effects due to the magnetic drift and the presence of spatial pressure gradients of passing energetic particles, whereas velocity space gradients lead to damping (for Maxwellian and slowing down distributions). Estimates of stability were obtained by using the analytic expression for the linear growth rate of GAE's (Fu and Van Dam):

\[
\frac{\gamma}{\omega_{\text{GAE}}} = \frac{\beta_b}{2k || R^2} \left( \frac{\omega_{\text{eh}}}{\omega_{\text{GAE}}} - 1 \right) \frac{v_A}{v_b} \frac{1}{\sqrt{F}} \frac{v_A}{v_e}.
\]

(\(\beta_b, \beta_e\) fast and electron beta, \(F\) proportional to the fraction of resonating ions). The effect of the fast ion pressure gradient is expressed in terms of the diamagnetic drift velocity \(\omega_{\text{eh}}\). In accordance with the observations, therefore, only waves propagating in the direction of \(\omega_{\text{eh}}\) can be excited. The electron Landau damping term is also associated with the curvature drift motion. For typical cases in W7-AS, where \(\omega_{\nu} / \omega_{\text{GAE}} = 3...20\), \(v_e / v_A = 0.3...1.0, \beta_b / \beta_e = 0.05...0.15\) is found, positive growth rates are obtained from this relation. However, the damping is likely to be underestimated because of neglecting collisional effects (trapped electron damping), since \(v_e / \omega_{\text{GAE}} \approx 1\) and resistive dissipation has to be taken into account.

**Numerical calculations** - We have analyzed the GAE modes for a few cases with a gyrofluid model with Landau closure for the energetic particles. This approach includes Landau damping/growth effects as well as continuum damping within a fluid description (with reduced MHD) for the background plasma. In the calculations a Maxwellian distribution instead of a beam slowing down distribution for the energetic ions and circular toroidal equilibria were used. The damping effects from the finite ion Larmor radius effects of the bulk plasma are retained in this model; however, finite orbit width effects of the fast ions...
are neglected. The results of first linear and nonlinear GAE calculations are in good agreement with the experimental data. In the well documented case discussed in fig. 3 with an $\tau$ profile ranging from 0.36 (center) to 0.34 (edge) a dominant (3,1) mode with global behaviour similar to the observed structure was found to be dryen unstable at a low threshold ($\beta_0 \geq 2\cdot10^{-4}$) due to weak damping. The resonance velocity peaks well below $v_b/v_A = 1$ (fig. 6) as a consequence of low GAE frequency and sideband excitation. The calculated mode power spectrum and the nonlinear saturation level of $B_0/B = (1-2)\cdot10^{-4}$ fits quite well to the experimental data. The poloidal mode spectrum is very peaked. In agreement with the observations higher toroidal (and poloidal) mode numbers are progressively more stable due to nonlinear (fig. 7) and resistive effects (fig. 8).

Conclusions - Energetic particle driven GAE modes are favoured global modes in the shear Alfvén spectrum of low shear devices. They are weakly damped and have been found the first time in the W7-AS stellarator. The GAE activity can possibly induce enhanced fast particle losses. However, this effect has not yet been observed in W7-AS. In particular, GAE’s are absent in the high $\beta$ regime, presumably due to more significant background Landau and continuum damping and effects of broad profiles.

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References
Power Balance in a Helias Reactor

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Stellarator reactor studies have been undertaken since the beginning of stellarator research. The first studies of modular stellarator reactors mainly concentrated on technical issues neglecting the limitations set by confinement and stability. In contrast to these early studies, the Helias configuration (Nührenberg and Zille, 1986) offers the chance to develop a self-consistent reactor concept where the plasma losses, MHD stability limits and α-particle losses are not prohibitive to ignition.

The dimensions of the Helias reactor are mainly determined by the technical limits of the coil system and the necessary space for blanket and shield. The main parameters are: Major radius 20 m, average plasma radius 1.6 m, magnetic field on axis 5 T, rotational transform on axis 0.81, transform on the boundary ≈ 1.0. Details of the Helias reactor concept are described in (Beidler et al. 1992). The MHD-stability limit in a 5-period Helias configuration is expected to be $\beta = 4\text{-}5\%$ and the neoclassical transport losses can be characterised by an effective helical ripple of 1\text{-}2%. This immediately poses the question of whether ignition is possible and, if so, if the fusion power is in the desired regime of $P = 2.5\text{ - }3.0$ GW.

Since a local power balance needs too many assumptions about transport coefficients and boundary conditions we follow the standard approach with fixed plasma profiles and empirical scaling laws for the confinement time. The plasma profiles are modelled analytically and the power balance is evaluated numerically.

Various scaling laws of anomalous confinement (LHD scaling, Gyro-Bohm scaling, Lackner-Gottardi scaling (Lackner and Gottardi, 1990) have been tested in present-day stellarator experiments, showing no clear distinction among the various scaling laws. In Wendelstein 7-AS all three scaling laws fit the experimental data. Therefore, the pragmatic approach is to extrapolate the empirical scaling laws to a reactor plasma and to test the ignition conditions under various assumptions on anomalous confinement.

A Helias reactor is a steady-state reactor with a low amount of recirculating power. To maintain the auxiliary systems roughly 10% of the electric power is needed. A rough power balance starting from an electric power of 1000 MW shows that on the plasma side a neutron power of 2000 - 2400 MW has to be delivered; this corresponds to an α-particle heating power of 500 - 600 MW. This figure sets a lower limit on the reactor performance and therefore determines strongly the dimensions and the plasma parameters of a Helias reactor.
Power balance. The power balance of a Helias reactor depends on the confinement properties of the plasma which means the confinement of highly energetic alpha particles, the neoclassical transport of the thermal plasma and the anomalous transport caused by plasma turbulence. Alpha-particles in non-symmetric magnetic fields are not absolutely confined, in general. Especially locally trapped particles tend to drift to the wall very rapidly which not only reduces the available heating power to the plasma but also causes damage to the first wall. In the Helias reactor some effort has been made to overcome these difficulties. Firstly, due to the high density and low temperature of the burning plasma the slowing down time is on the order of 0.1 s, and therefore the confinement time of energetic alpha-particles need not be larger than 0.1 s. Although there is no absolute confinement of trapped particles in a configuration without symmetry, improved confinement of these particles in a Helias reactor can be achieved by localizing the trapped particles in each field period and by utilizing the poloidal magnetic drift to avoid the formation of superbanana orbits. That this method is indeed effective has been demonstrated numerically by (Lotz et al. 1991). The result is that nearly all highly energetic alpha-particles are confined for one slowing down time if the plasma $\beta$ is sufficiently large. The poloidal magnetic drift already becomes effective at $\beta = 2\%$; at $\beta = 5\%$ only a few percent of the alpha particles are lost within one slowing-down time. In view of this result, prompt losses of alpha particles have been neglected in the power balance.

Neoclassical losses in a Helias configuration are strongly reduced by the drift optimization of particles orbits. Numerically these neoclassical losses in Helias configurations have been studied by various techniques: Monte Carlo calculations (Lotz and Nührenberg 1988), the DKES code (Maaßberg et al. 1993) as well as other techniques for solving the kinetic equation (Beidler, 1991). In the P.S.-regime and the plateau regime the transport is mainly determined by the $C_{01}$-term in the Fourier spectrum of $B$, where $B/B_{01} = 1 + C_{01} \cos \theta + C_{11} \cos \theta \cos M \varphi + S_{11} \sin \theta \sin M \varphi + C_{10} \cos M \varphi + \ldots$. $\varphi$ is the toroidal angular variable and $\theta$ a poloidal angular variable. $M$ is the number of field periods. Since the term $C_{01}$ is reduced by a factor of two compared with $\epsilon_0$, the averaged deviation of the drift surfaces of passing particles from magnetic surfaces is also reduced by a factor of two. This means that at fixed rotational transform the reduction factor of P.S.-diffusion and plateau diffusion is four. In the long-mean-free-path regime it is mainly the synergistic effect of the leading Fourier harmonics $C_{01}, C_{11}, S_{11}, C_{10}$, which leads to the reduction of radial transport. The 'helical' ripple in Helias configuration can be characterized by $\epsilon_h \approx 0.1$ (Beidler 1991) whereas the effective ripple is $\epsilon_{eff} = 0.01 - 0.02$; therefore the optimization yields a reduction factor of 10 - 20 in the $1/\nu$-regime. For the ions in the limf-regime the radial electric field is more effective in reducing the radial diffusion than the optimized magnetic drift.

The ignition condition of a Helias reactor requires that the effective thermal conductivity must be not higher the 1 m²/s, therefore neoclassical transport coefficients must stay well below this limit. The plasma parameters in the reactor regime are: $n(0) = 2 - 4 \cdot 10^{20}$ m⁻³, $T \leq 16$ keV. The numerical calculations show that the neoclassical transport coefficients in the parameter regime described above are in the range 0.2 - 0.1 m²/s. This result holds for electrons and ions, however the ion thermal conductivity is strongly reduced by the radial electric field. This radial electric field is determined
by the condition of ambipolarity and lies in the ion root regime. Since a Helias reactor operates at relatively high density \((3 \times 3.5 \times 10^{20} \text{ m}^{-3})\) the collisionality is in the range \(\nu^* = R/\lambda_e = 1 - 2 \times 10^{-3}\). These results show that neoclassical transport is small enough and the confinement is good enough to sustain ignition.

Anomalous transport is a phenomenon found in all present-day stellarator experiments. Several empirical scaling laws of the energy confinement time exist which fit to the experimental results. The Lackner-Gottardi scaling, which is favoured in the W 7-AS stellarator, is \(\tau_E = 0.175 R a^2 B^{0.8} \bar{n}^{0.6} P^{-0.64} (A/1.5)^{0.5}\). Besides the Lackner-Gottardi scaling, the LHID scaling and the Gyro-Bohm scaling are employed to predict plasma parameters in ITER. The LHID-scaling is \(\tau_E = 0.17 P^{0.75} a^2 B^{0.84} \bar{n}^{0.6} P^{-0.58} (A/1.5)^{0.5}\) and the Gyro-Bohm scaling \(\tau_E = 0.25 R^{0.6} a^{2.2} B^{0.8} \bar{n}^{0.6} P^{-0.68} (A/1.5)^{0.5}\). These two scaling laws do not show a dependence on the rotational transform, however a positive scaling with \(a\) (\(\tau_E \propto a^{0.28}\)) was found in Wendelstein 7-AS supporting the Lackner-Gottardi type scaling. The units are: time in s, heating power \(P\) in MW, magnetic field \(B\) in T, line average density \(\bar{n}\) in \(10^{20} \text{ m}^{-3}\), lengths in m. The isotope effect, described by the atomic number \(A\), has not yet been verified in stellarator experiments, however, since this effect has been found in many tokamak experiments, it will be included tentatively to test its relevance for reactor conditions. Columns 1-3 in Table I are calculated with an effective atomic mass of \(A=2.5\), columns 4-7 with \(A=1.5\).

**Results**. Since the anomalous transport is larger than the neoclassical one, the power balance is calculated on the basis of the empirical scaling laws of energy confinement. Below the ignition threshold the plasma is maintained by the external heating power \(P_{ex}\). With fixed plasma profiles the power balance \(E/\tau_E = P_a - P_{brem} + P_{ex}\) can be brought into the form \(f(P_{ex}, n(0), T(0)) = 0\). The curves \(P_{ex}\)-constant are plotted in the form of POPCON-plots. Parameter variation were made in the following regime: 1) Two different plasma profiles (type I and type II, where type II is optimized with respect to MHD-stability); 2) Variation of impurity content: abundance of cold alpha particles \(\leq 10\%\), abundance of Oxygen \(\leq 1\%\), abundance of Carbon \(\leq 0.5\%\); 3) Dependence on the isotope factor in \(\tau_E, A = 1.5 - 2.5\); 4) Major radius \(20 - 22\text{m}\); 5) Magnetic field \(5 - 5.5\text{T}\); 6) Confinement time: LHD scaling, Gyro-Bohm scaling and Lackner-Gottardi scaling.

The aim is to find an operating window where the ignited state \(P_{ex} = 0\) stays below the stability limit which is expected between \(\beta = 4\%\) and \(3\%\). Such an operating window does not exist for confinement times following the Gyro-Bohm or the LHD-scaling; an improvement factor of 2 or less is needed. This is similar to tokamak reactors where L-mode confinement is also insufficient to reach ignition. However, the Lackner-Gottardi scaling predicts a higher confinement time and ignition can be reached without the need for further improvements. The plasma parameters of various ignited states close to the MHD-stability limit are listed in Table I.

**Conclusions**. In the envisaged parameter regime of a Helias reactor \((B = 5\text{T}, n(0) = 2 - 4 \times 10^{20} \text{ m}^{-3}, T \leq 17\text{keV})\) the neoclassical transport is smaller than the anomalous transport. A zero-dimensional power balance based on empirical scaling laws of energy confinement shows that a confinement time following the Lackner-Gottardi scaling satisfies the ignition condition. Any reference to H-mode confinement or other improvement
Table I: Plasma Parameters in the Helias Reactor

<table>
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<tr>
<th>Parameter</th>
<th>Value</th>
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<tr>
<td>Major radius [m]</td>
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<td>Average plasma radius [m]</td>
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<td>Magnetic field on axis [T]</td>
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<td>Fusion power [MW]</td>
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<tr>
<td>$\tau_{LHD}$ [s]</td>
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</tr>
<tr>
<td>$\tau_{Cerom}$ [s]</td>
<td>0.90</td>
</tr>
</tbody>
</table>

Mechanisms is not necessary. The reason for this favourable behaviour of the L.G. scaling is the dependence on rotational transform. Accumulation of cold alpha particles is tolerable up to a fraction of 10%. Under these conditions the reactor runs close to the stability limit of $β=5\%$ and delivers a fusion output of 2.5 - 3.8 GW (see Table I). Critical issues of the power balance are a further increase of the impurity content and anomalous loss of fast alpha particles which diminishes the heating power. In this context any improvement of energy confinement, which might arise from the optimized features of the Helias configuration, would be very useful.

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References
ISLAND DIVERTOR CONCEPT FOR THE STELLARATOR WENDELSTEIN 7-X

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Introduction

The proposed Advanced Stellarator Wendelstein 7-X, is a Helias configuration which has been optimized with respect to several criteria concerning plasma performance and coil geometry. The main data are: \( R_o = 5.5 \) m, \( a = 0.52 \) m, \( B_o = 3 \) T. The aim of steady-state operation and a large heating power require a divertor configuration for control of plasma-wall interaction and minimization of impurity inflow. The goal is to optimize the magnetic field in the boundary region without losing any of the other favourable properties, to locate suitable target plates in the outflowing plasma avoiding too excessive power loads and to hold back the re-emitted neutral gas for high recycling probability and pumping efficiency. The concept described in this paper utilizes the existence of large magnetic islands or their remnants in the plasma boundary.

Island divertor

In the standard configuration of Wendelstein 7-X (W7-X), five islands at the boundary exist corresponding to the rational value of the rotational transform \( \epsilon_a = 5/5 \). Increasing \( \epsilon \) to 5/4 at the edge yields four islands which are larger than those at \( \epsilon = 1 \). Usually they exist only in the vicinity of the O-points surrounded by an ergodic layer. The ergodisation can be enhanced by finite \( \beta \) effects or by increased shear at the boundary /1/. The separatrix of the islands provides the diversion of the field lines and takes over the role of the X-line in torsatrons or in tokamaks. The outflowing plasma crosses the separatrix by diffusion and streams along the field lines towards the rear of the island where the target plates are located. The toroidal and poloidal position of the plates is chosen where the radial dimension of the islands is maximal and where the diverted field lines attain the largest distance from the main plasma. The viability of the island divertor concept depends on the stability of the island structure against external and internal magnetic field perturbations, and the capability to guide the outflowing plasma to the target plates, i.e. a certain range of parallel to perpendicular transport is required. Several studies have been made to check this concept and to optimize the geometry of the target plates. These include:

- Studies of the magnetic field structure, localization of the O-points and X-points outside the last magnetic surface, effect of perturbation fields on the islands, and geometry of flux bundles in the vicinity of islands and ergodic regions;
- investigation of particle diffusion and heat conductivity in and around islands by mapping and Monte Carlo techniques;
- optimization of divertor target plates and calculations of heat loads on them;
- design of sweep coils /2/ for control and modification of the edge field structure and for compensation of symmetry breaking error fields by adjusted DC currents, and for reduction of the heat load by AC currents.
Target plates

For a given power flux along the field lines, the intersection angle with the plates determines the power density on the plates. If the power density is held below the technical limit of 10 MW/m², an intersection angle of about 2° – 3° results with an estimated width of the scrape off layer (SOL) of about 2 cm and a heating power of 10 MW. The total length of the plates is then determined by the intersection angle, the requirement that all field lines in the SOL hit the plates, and the requirement that the leading edge of each successive plate is protected by the “shadow” (see Fig. 1) of the previous one. In each of the 5 field periods two target plates are installed with a toroidal length of about 5 m and an average width of 0.5 m, resulting in a total surface of 25 m². These target plates follow roughly the O-point of the $\tau = 1$ islands in a toroidal range of 54 degree (Figs. 2 and 3).

The problem arising from the different configurations created by variation of the rotational transform, $\tau$, is solved by adjusting parts of the plates to each specific case. In the standard case ($\tau_a = 1$) the field lines hit more the center part of the plates and in the high and low iota case ($\tau_a = 5/4$ and $5/6$, resp.) the target plates are loaded more towards one or the other end.

A Monte-Carlo technique is used to obtain the intersection patterns on the target plates. The SOL is simulated by calculating the particle motion parallel to the field together with perpendicularly displacements after each step length (e.g. displacement 0.1 cm, step length 30 cm, for an anomalous thermal conductivity $\chi_\perp = 3 \text{ m}^2/\text{s}$ in a plasma which has a parallel conductivity of $10^7 \text{ m}^2/\text{s}$). The starting points of the field lines are statistically spread on a magnetic surface close to and inside the bounding separatrix; the integration is stopped at the target plates. The intersection points of the separatrix and the adjacent field lines form stripes on the target plates in a width reflecting the width of the SOL (see Fig. 4) separated in the directions parallel and antiparallel to the magnetic field. The intersection pattern differs for the various cases examined:

Fig. 1. Schematic arrangement of segmented target plates. The length of the plates is chosen to protect the leading edge of each successive plate with the “shadow” of the previous one.

Fig. 2. Magnetic surfaces (one field period, just inside the SOL), target plates, and sweep coils of W7-X viewed from the radial outside.

Fig. 3. Cross-section of the “bean-shaped” toroidal plane of W7-X. Shown are the system of magnetic surfaces, the SOL represented by Monte-Carlo calculation, and target plates. The first wall is shown as a dashed line and the radial dimension of the coil winding pack as the hatched area.
- Standard case (HS5V10N)— variation of $\tau_0$ (0.33, 0.86, 0.89), variation of island size using the sweep coils, variation of diffusions coefficient ($\chi_a = 3$ m$^2$/s and 0.6 m$^2$/s).
- Low iota case — variation of island size and radial position
- High iota case — variation of island position and ergodisation

No deterioration of the divertor action occurs in the high $\iota$ case as long as perturbations do not spoil the general field structure within the short length of about 3-5 toroidal transits, which takes the field line from the front to the rear side of the island. The stripes are narrower when the diffusion coefficient is reduced, and because the SOL width depends on the connection length and this length on the island size, the width of the stripes decreases with increasing size of the islands. Consequently, the smallest stripes and largest power densities are found in the high $\iota$ case.

Neutral gas behaviour

The plasma parameters in front of the target plates depend on the plasma surface interaction, most significantly by the interaction with the neutral gas re-emitted by the targets. Low plasma temperatures and high densities in the SOL reduce the impurity source (sputtering) and reduce the impurity inflow if they are ionized close to the target plates. These favourable edge plasma parameters reduce also the power density on the target plates by radiative losses (recombination) and diffusive broadening of the SOL. In this context the aim of W7-X is to establish a high recycling region in front of the target plates. This is attainable by a proper design of target plates, baffles, gas-puff equipment, an efficient pumping system, and the possible operation at a high edge density without a disruptive density limit.

The behaviour of the neutrals is studied with the neutral gas transport code EIRENE /3/. The geometry part of the code is adapted to the complex geometry of the magnetic surfaces and the island structure; see Fig. 5. The intersection pattern on the target plates from the Monte Carlo calculations serve as particle source regions for the neutrals. At present, the code is used as a stand-alone code with fixed plasma parameters. The profile functions for the electron and ion temperatures and the density in the plasma column are: $T_e(r) = 4.6/[1 + (r/20)^9]$ [keV], $T_i(r) = 2.4/[1 + (r/36)^9]$ [keV], $n_e(r) = 1 \cdot 10^{11}/[1 + (r/38)^9]$ [cm$^{-3}$]; $r$ is the average radius of the magnetic surfaces.

Fig. 4. Intersection pattern on one of the lower target plates, standard case ($\iota_a=1$, $\chi_a=3$ m$^2$/s).

Fig. 5. Cross-section of the mesh for the EIRENE-code at the toroidal plane with $\varphi = 0$ (upper part of the up-down symmetric figure). Shown are the radial and poloidal cell boundaries and the positions of the target and baffle plates which form the pumping chamber (first case). The plasma column is surrounded by the islands and an outer region between SOL and the first wall (largest radial boundary).
The parameters are adapted to code predictions on the basis of transport analysis. In the SOL (outside the separatrix) the field lines contact the target plates: a radial decay length of about 1.5 cm is assumed. The plasma pressure is assumed to be constant along the field lines, however the density towards the target plates is increased by a factor of 4 to a value of $8 \cdot 10^{13}$ cm$^{-3}$ in conformity with a temperature drop to values of 15 to 40 eV. These parameters lead to a re-ionisation of 80% to 90% of the neutral particles just in front of the target plates and nearly all are re-ionized in the SOL; see Fig. 6.

The first interest of the calculations was to explore the pumping efficiency. In W7-X, cryo-pumps are planned behind or near the target plates in a chamber separated from the main plasma by baffle plates, which prevent the neutrals from moving away, especially in those region where the plasma has a small radial extension. Two alternatives are considered in the calculations. The first uses transverse slits in the target plates of about 1.5 cm width and a target/slit ratio of 3.4 to get the neutrals into the pumping chamber. In this case about 1% of the source flux on the target plates enters the chamber, composed mainly of charge-exchange particles (Fig. 5). A factor of 3 larger flux into the chamber is obtained in the second case. Here the target plates have no slits, however one side of the chamber is left open (Fig. 6). The neutral particles coming from the target plates and from the edge plasma arrive at the pumping chamber after a few reflections.

Conclusions

The island divertor concept presented in this paper adapts the axisymmetric open divertor to the 3-dimensional stellarator configuration of W7-X. The proposed target plates are fit to various values of $\varepsilon$, permitting a variation of the island size and position; their operation is rather insensitive to magnetic field perturbations. They are located in regions where the islands have a maximum radial extent which implies a large distance from the main plasma. The toroidal length of the target plates is sufficient to keep the power density within technical constraints; a nearly constant intersection angle of the field lines at the highly loaded regions provides a smooth power load and avoids hot spots. Leading edges are avoided by making use of the “shadow effect”. There is a good chance to enter the high-recycling regime in a wide parameter range and to collect a substantial fraction of the neutralized particle with the cryo-pumps.

References

WEENDELSTEIN 7-AS CONFIGURATIONS
AT VARIOUS MIRROR RATIOS AND $\varepsilon$-VALUES

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The coil systems of Wendelstein 7-AS provide various means to modify the vacuum-field topology for this experiment, to change $\varepsilon$ via the fields of the planar outer coils and to shift the magnetic surfaces via the vertical field. An additional variation of the magnetic field strength along the field lines (toroidal direction $\varphi$) is obtained by different currents in the non-planar modular coils. The 'mirror ratio' is defined in this paper as $MR = (B_2 - B_1)/(B_2 + B_1)$, where the fields $B_1$ and $B_2$ are the values on axis in the 'triangular' plane, $\varphi_1 = 0$, and the 'elliptical' plane, $\varphi_2 = 36^\circ$, respectively. An increased (positive or negative) mirror ratio influences the fraction and behaviour of trapped particles and thus the neoclassical properties of the plasma, e.g. neoclassical losses in the limf-regime, the electric conductivity and the bootstrap current. The main effects of a changed mirror ratio are expected near the magnetic axis.

The present paper continues the vacuum field calculations of Ref. /1/ and presents experimental results of ECRH-produced plasmas at an edge value of $\varepsilon_e \approx 0.34$ with $MR = \pm 10\%$. The ECRH power is applied in the 'elliptical' plane, for ‘on-axis’ heating at 1.25 T (resonance field for ECRH). More recent experiments (two values of the ECRH input power and densities, also with changed mirror ratio) are under analysis. These experiments aim to establish the range of attainable plasma parameters. The paper thus contributes to develop a data base for a future discussion of more general aspects of toroidal confinement. A rather small ohmic current is applied to balance the residual plasma currents. Issues concerning these currents, especially the bootstrap current, are addressed in Ref. /2/.

Fig. 1 shows the nested magnetic surfaces for a configuration with $MR = 10\%$ at a value of $\beta_n = 0.7\%$ on axis, obtained by the VMEC code. The corresponding $\beta$-profile is shown in the top part of Fig. 2. This profile is fitted from Thomson scattering data of the shot series AS19980 with 3 gyrotrons at a total power of about 570 kW. $T_i$ is not measured; its contribution to the $\beta$-profile is expected to be small for this type of discharges. The resulting radial profile of the rotational transform is shown in the lower part of the figure: the central $\varphi$-values are reduced compared to those of the vacuum field (dashed curve with positive shear of about 1%). The change in the $\varepsilon$-profile is due to the internal currents, i.e. ohmic, bootstrap and Pfirsch-Schlüter currents and yields a net value of $-100\ A$ in the results of the VMEC code.

In Fig. 3, the experimental profiles of $n_e$, $T_e$ and $\chi_e$ of this series are compared with those of a series with 'reversed mirror', AS21299 with $MR = -10\%$, where slightly higher axis temperatures are seen at 520 kW heating power, as well as with a recent discharge series at standard low $\varepsilon_e = 0.34$, AS22628, heated by 2 gyrotrons at 350 kW. The three experiments are performed at $n_e, o \approx 1.5 \cdot 10^{19} \ m^{-3}$ on axis. Peak electron temperatures
Fig. 1: Finite-β fields obtained by the free boundary code VMEC for the experimental profile of AS19980, (565 kW ECRH heating power) for $\beta_o = 0.67\%$, $(\beta) = 0.18\%$; two toroidal planes of reference, $\varphi = 0$ and $36^\circ$. The Shafranov shift can be seen at $\varphi = 0$ as an offset of the axis position from $R = 200$ cm.

Fig. 2: Characteristic radial profiles for the fields of Fig. 1.

shots:
- 19980-20023
- 21299-21310
- 22628-22654

Fig. 3: Experimental profiles studied in the present paper.
**Fig. 4:** Radial dependence of the effective ripple according to Shaing/Houlberg shown in the left part, and fraction of trapped particles.

**Fig. 5:** Dependence of the effective ripple according to C.D. Beidler on the rotational transform.

**Fig. 6:** Shafranov shift obtained in the VMEC code for the three cases of Fig. 3 with profile factor $\beta_o/(\beta) \approx 3$, (middle part of the figure), along with standard VMEC profiles with profile factor $\approx 2$. 

**Axis Position cm**

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$T_{e,o}$ are between 1.3 and 1.5 keV. The values of $Z_{eff}$ are unknown. Experimental values of $Z_{eff}$ are unknown. Experimental values of the electron heat conduction coefficient range between 5 m$^2$/s near the axis and 3 m$^2$/s at about 12 cm minor radius in the two cases of enhanced mirror ratio. Differences for $r < 7$ cm are marginal. For the series AS2628 the profile of $\chi_e$ is flat at about 3 m$^2$/s. For the three profiles values of $\nu_e$ below 10$^{-2}$ apply near the axis, and the plasma is within the limf-regime in this region.

Fig. 4 shows effective ripple values for various configurations at low $r$ and the associated fractions of trapped particles. The largest neoclassical losses in the limf-regime would apply for the configuration with reversed mirror, MR = $-10\%$. The $r$-dependence of the effective ripple is depicted in Fig. 5. These ripple values depend on $r$ only for the cases with 'reversed' mirror; values are near that of the inverse aspect ratio at low $r$, and decrease towards $r = 0.53$. Results for the standard case might be underestimated.

The experimental profiles of Fig. 3 are analyzed with the DIES code, using Fourier coefficients of the respective vacuum fields. Radial electric fields are not included; discrepancies exist in the resulting power and current balances. For the experiment with MR = 10 % we find an error of about 2 kA in the balance of the OH and bootstrap currents. The latter amounts to about 7 kA, when assuming a uniform $Z_{eff} = 2$. Near the axis the analysis is uncertain. It is performed at $r = 5$ cm. The calculated values of $\chi_{e, neoc}$ are 7, 10, and 2 m$^2$/s for the two cases with enhanced mirror ratio and the standard case, respectively. The largest value applies for MR = $-10\%$. Both experiments with enhanced mirror ratio yield experimental values of about 5 m$^2$/s at this radius, about 2 m$^2$/s are seen for the standard case. So far, no answer is possible on the relative amount of neoclassical to anomalous energy losses for the inner region of these discharges.

The Shafranov shift is analyzed for the three cases of Fig. 3. The quantity $\langle |j_{||}/j_{\perp}| \rangle$ is slightly larger at modified mirror ratio than under standard operation. Consequently the Shafranov shifts are nearly identical, see Fig. 6: for $\beta_o = 1\%$ with a profile factor $\beta_o/(\beta) \approx 2$ they are about 3 cm in the 'elliptical' plane and about 4 cm in the 'triangular' one. The experimental data of Fig. 3 have profile factors $\approx 3$ and smaller $\beta$-values; their smaller Shafranov shifts are shown in the center part of the figure.

Summary and Conclusion:

1) Configurations with modified mirror ratio of the magnetic field are useful to investigate the relative amount of anomalous and neoclassical losses, to analyze details of the electric conductivity, of bootstrap and of ECCD-induced currents in the limf-regime.

2) Experiments at about 1.25 T were initiated for MR = $\pm10\%$ and compared to discharges at standard low-$r$ operation; more recently constant minor radii were approached at different mirror ratio values by adjusting the limiter positions. The analysis of the results is in progress.

3) Further experiments with modified mirror ratio are to be performed, varying the density, heating power, iota, limiter and axis positions.

4) The relative amount of neoclassical to anomalous losses cannot be specified so far for the three experiments discussed in this paper.

/2/ F. Rau et. al.. 9th Int. Workshop on Stell., Garching 1993, (to be published)
EFFECT OF NEUTRAL PARTICLES ON DENSITY LIMITS IN TOKAMAKS

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1 Introduction

The global stability and confinement of a tokamak plasma are significantly influenced by the boundary plasma parameters. The onset of density disruptions, which limit the maximum plasma density, is triggered by impurity radiation in the edge plasma and can be connected with the radiative thermal instability [1]. At the density $n_c$ the total radiative power $P_{\text{rad}}$ is equal to the total input power $P_{\text{in}}$ into the plasma ($S := P_{\text{rad}}/P_{\text{in}} = 1$). Above $n_c$ ($S > 1$) no steady state of the plasma column exists. Contrary to predictions made in [1], where neutral particle kinetics is not taken into consideration, experimental results show that disruptions can occur for $S < 1$ [2], [3]. It was shown in [4] that the carbon impurity radiation cooling of the plasma is strongly affected by charge exchange between carbon ions and hydrogen atoms which penetrate into the plasma as a result of recycling or gas puffing.

We present analytical expressions for the cooling rate $Q_R$ as a function of the plasma temperature $T$, $\xi_N := N/n$ and $\xi_i := n_i/n$, where $N, n_i, n$ are the densities of hydrogen atoms, impurity ions and the plasma, respectively. We investigate the influence of the neutral particles on the critical densities and the stability of the system, taking into account ionization, charge exchange and impurity cooling.

2 Model Equations

Averaged over the magnetic field lines in the SOL, the dynamics of the plasma and neutral particles perpendicular to the magnetic field is described by the following set of coupled equations:

$$\frac{\partial n}{\partial t} + \frac{\partial \Gamma_n}{\partial x} = k_i(T)nN - \theta n/\tau_1(T),$$

(1)

$$\frac{\partial N}{\partial t} + \frac{\partial \Gamma_N}{\partial x} = -k_i(T)nN - \theta Rn/\tau_1(T),$$

(2)

$$\frac{3}{2}n\theta T/\partial t + \frac{\partial \Gamma_T}{\partial x} = \sigma(T)E^2 - Q_R(T, N, n) - \theta nT/\tau_2(T) - \theta \kappa_{||}(T)T/L^2,$$

(3)

$$\Gamma_n = -D_{\perp}(n, T)\partial n/\partial x,$$

(4)

$$\Gamma_N = -(1/2\kappa_{xx} n)\partial (v_i^2 N)/\partial x,$$

(5)

$$\Gamma_T = -\kappa_{\perp}(n, T)\partial T/\partial x.$$

(6)
\( D_{\perp} \) and \( \kappa_{\perp} \) are the coefficients of the perpendicular anomalous plasma diffusion and heat conduction, and \( \kappa_{\parallel} \sim T^{3/2} \) and \( \sigma \sim T^{3/2} \) are the classical heat and electric conduction coefficients, resp.; \( \tau_m := \beta_m L/v_s \) is the lifetime of the plasma \((m = 1)\) in the SOL and its energy \((m = 2)\) due to streaming along the field lines to the limiter \((\beta_m > 1, L = \pi q R \cdot \text{connection length})\), and \( v_s = \sqrt{2T/m_i} \) is the ion sound velocity; \( \kappa_{\text{i}} \) and \( \kappa_{\text{ex}} \) are the rate coefficients for ionization and charge exchange, resp., and \( R \) is the recycling coefficient at the limiter. \( R \) is the separatrix radius, \( \Theta = \Theta(z) \) the Heaviside function, and \( E \) the electric field of the tokamak.

The analytical expression for the cooling rate \( \dot{Q}_R = n^2 \xi_i G(T, \xi_N) \) [W/cm^3] (see Fig. 1) in the temperature range \( 5eV \leq T \leq 50eV \) with

\[
G = A(T - T_1)e^{-B(T - T_1)},
\]

\[
A = 10^{-26}(3.1e^{-0.024(\ln \xi_N + 11)^2} + 0.7), \quad B = 0.6e^{-7.810^{-4}(\ln \xi_N + 14)^4} + 0.055
\]

\((T \text{ in } eV, n \text{ in } \text{cm}^{-3}, T_1 = 4.4eV)\) may be simplified in two steps:

\[
A = 0.78 \times 10^{-26} \xi_N^{-0.18}, \quad B = 0.041 \xi_N^{-0.31}, \quad \xi_N > 10^{-4},
\]

(see [5]) and according to [1]

\[
G = G_0 \Theta(T_w - T),
\]

\[
L_0 = \frac{A}{B^2} \frac{1 + B(T_w - T_1)}{T_L} e^{-B(T_w - T_1)}, \quad T_L = T_1 + 2/B,
\]

where \( T_w \) is the wall temperature.

3 Equilibrium and Critical Densities

The system of eqs. (1)-(6) was solved under the steady-state condition \( \partial/\partial t = 0 \) for \( z \in [x_c, x_w] \) \((x_c \text{ - position in the central plasma, } x_w \text{ - distance from the separatrix to the wall})\). For the fluxes \( \Gamma_A \) with \( A = n, N, T \) it is of the following form:

\[
d\Gamma_A/dz = H_A - L_A
\]

with the corresponding source \( (H_A) \) and loss \( (L_A) \) terms. An equilibrium state exists if the condition

\[
I_A(z) := \Gamma_A(x_c)^2 + 2 \int_{x_c}^{z} dz' \Gamma_A(H_A - L_A) \geq 0 \forall z \in [x_c, x_w]
\]

is fulfilled. Otherwise critical parameters may exist. For \( z = x_w \) the general 6-parameter solution of our problem can be reduced down to a 3-parameter one.

For the case of complete recycling at the limiter \((R = 1, \Gamma + \Gamma_N = 0)\) we obtain with

\[
D_{\perp} = D_0 n^p
\]

\[
N = \frac{1}{v_s^2} \left( C - \frac{2k_{\text{ex}}D_0 n^{1+p}}{2 + p} \right), \quad C = \text{const.}
\]

For \( p = 0, x_c = -\infty, x_w = \infty, n(x_c) = n_0, N(x_c) = 0, n(x_w) = 0, N(x_w) = N_w \) and averaged temperature it follows (cf. [6]) that

\[
N_w \leq N_w^c = 1/k_i \tau_1, \quad n_0^2 = \frac{1}{2} \frac{N_w v_s^2}{k_{\text{ex}} D_0}.
\]
is the upper limit of the neutral density at the wall. Next we investigate the energy balance equation with averaged densities, assuming the radiation term as the dominant loss term in the approximation (10), (11), and consider the two cases of (i) a transparent \( l_N > l_L \) and (ii) a non-transparent \( l_N < l_L \) SOL for the neutrals. \( l_N = \frac{v_s}{n \sqrt{2k_0 k_w}} \) is the mean free path of the neutrals and \( l_L \) the width of the radiation function. The condition \( I_T(x_c) = 0 \) reads

\[
\Gamma_T^2(x_c) = 2\kappa_\perp n^2 \xi_i G_0 \begin{cases} \frac{T_L - T_w}{T_N - T_w}, & l_N > l_L \quad (i) \\ \frac{T_L - T_w}{T_N - T_w}, & l_N < l_L \quad (ii) \end{cases}
\]

with \( T_L \) as given in eq.(11) and \( T_N = T_w + \Gamma_T(x_c) l_N / 2 \kappa_\perp \), which is equivalent to the above-mentioned condition \( S = 1 \) \( (S \equiv Q_R b / \Gamma_T(x_c), b = N, L) \). Using eq.(14), we obtain the following scaling relations \((D_0 \sim q^a, B(T_w - T_1) \ll 1, A, B \text{ given in } (9))\):

\[
n_c \sim \left( \frac{P_{in}}{q^a N_{w0.44} \xi_i} \right)^{1/(2.56 + p)} \quad (i), \quad n_c \sim \left( \frac{P_{in}}{\xi_i q^{0.13 r}} \right)^{1/(1 + 0.13(1 + p))} \quad (ii).
\]

An example of a self-consistent solution of the problem is displayed in Fig. 2. The radiation density localized in the SOL strongly depends on the profile functions.

4 Radiative Thermal Instability

The stability of the equilibrium state with respect to poloidally and toroidally homogeneous perturbations is investigated by assuming all quantities \( \bar{A} \) to be of the form \( \bar{A}(x, t) = \bar{A}_0(x) + \delta \bar{A}(x) \exp(\gamma t) \). \( \bar{A}_0 \) is the equilibrium solution (section 3) and \( \delta \bar{A} \) is determined by a 6-dimensional linear system of equations with coefficients depending on \( \bar{A}_0 \). A numerical solution to the problem is in preparation. Here we give a preliminary analytical estimate for case (i) considering only the neutral gas density variation in eq. (3) which leads to

\[
\kappa_+ \frac{\partial^2 \delta T}{\partial x^2} - 3 n \gamma \delta T = \frac{\partial Q_R}{\partial T} \delta T + \frac{\partial Q_R}{\partial N} \delta N.
\]

Using approximation (10), this results in the dispersion relation

\[
k_+ + k_+ \cot g k_\perp = S(1 + 0.31\eta_N), \quad k_+^2 = \frac{3n \gamma l_{k_\perp}}{2 \kappa_\perp}, \quad k_+^2 = \frac{0.26 S \eta N}{2 - S} - k_2^2, \quad \eta_N = \left( T \frac{\partial N_{w0.44}}{\partial T} \right) l_L.
\]

The threshold for this thermal instability \( S_e (\gamma = 0) \) is shown in Fig. 3. The essential result is that neutral particles can cause the stable equilibrium state to be unstable to thermal perturbations.

Fig. 1: Impurity radiation for different values $\xi_n$: 1 - $\xi_n = 10^4$, 2 - $\xi_n = 10^3$, 3 - $\xi_n = 10^2$, 4 - $\xi_n = 10^1$

Fig. 2: Profiles of 1 - particle flux [$10^4$ s$^{-1}$ cm$^{-1}$], 2 - energy flux [W cm$^{-1}$], 3 - temperature [450 eV], 4 - plasma density [$3.5 \times 10^5$ cm$^{-1}$], 5 - neutral density [$10^8$ cm$^{-1}$], 6 - radiation density [$10^7$ W cm$^{-1}$], ($x_e = 15$ cm, $r = 1$, $\xi_e = 10^3$, $D_e = 10^4$ cm$^{-1}$ s$^{-1}$, $L = 2 \times 10^4$ cm)

Fig. 3: Radiation density threshold as a function of $\eta_n$
Bifurcation of Electron Temperature in the High Recycling Regime

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Introduction. Thermal instability in the boundary region of tokamaks has been recently studied numerically by Capes et al.[1], who considered the one-dimensional thermal conduction between the mid-plane of a tokamak and the divertor target plates. Multiple solutions and bifurcation of the electron temperature are caused by impurity radiation and its non-linear dependence on the temperature. In this paper it will be shown that the non-linearity of heat flux in the high-recycling regime in front of the divertor target plates can also trigger a bifurcation even in absence of impurity radiation. In this paper we consider a magnetic flux tube which is bounded by target plates on two sides. Such a case exists in tokamaks and stellarators in the scrape-off layer outside the last magnetic surface. Here, the high recycling regime is of particular interest since it allows to keep the temperature on the target plates at a low level and to minimize sputtering and other damaging effects on the target plates. To achieve this state, the density in front of the target plates must be sufficiently high so that ionisation of recycling neutrals occurs in a short distance from the target plates. Thus, the heat flow into the high recycling layer must cover the ionisation and radiation losses and the thermal flux onto the wall. Lackner et al.[2] have modelled this effect and proposed the following ansatz for the heat flow into the high recycling layer, or — since this layer is considered as very small — into the target plate.

\[ q_{\|} = (\delta_t k T_i + R \varepsilon_{\text{eff}}) n_t \sqrt{\frac{k T_i}{m_i}} (\gamma_e + \gamma_i) \]  

(1)

\( \delta_t, \gamma_e \) and \( \gamma_i \) are constants, \( n_t, T_i \) density and temperature at the target plate and \( \varepsilon_{\text{eff}} \) the ionisation energy enhanced by the radiated energy. \( R \) is the recycling coefficient.

On the plasma side of the high recycling layer these losses must be covered by the heat conduction into this layer. One of the standard approximations of the plasma boundary is constant pressure along the flux tubes \( p_o = 2 n_s k T_s = n_s k T_i; n_s, T_s \) are the plasma parameters at the stagnation point where the parallel flow velocity is zero. Introducing the pressure as a parameter on the flux tube yields the boundary condition

\[ (\delta_t k T_i + R \varepsilon_{\text{eff}}) \frac{p_o}{2 k T_i} \sqrt{\frac{k T_i}{m_i}} (\gamma_e + \gamma_i) = -n_{\|} \frac{B}{B^2} B \cdot \nabla T; \quad T = T_i \]  

(2)
where the parallel thermal conductivity is $n\chi_p = \kappa_p T^{5/2}$. The essential feature of this boundary condition is its non-linearity and the non-monotonic dependence on the temperature. This will give rise to bifurcation and multiple solutions as ref.[1] by volume radiation.

The heat conduction equation. In our model we consider the flux tube between two target plates outside the high recycling layer. We neglect convective energy transport and perpendicular thermal conduction. Therefore, the temperature is the solution of an one-dimensional conduction equation along the flux tube. Let be $h(x)$ a source term which describes the power input into the flux tube; $x$ is the length coordinate along the flux tube. The temperature is the solution of the following equation $-\nabla \cdot n\chi^{\parallel} B / B^2 \nabla T = h(x)$, which by introducing the magnetic potential $\psi (dy = B dx)$ as a new independent variable and $f = 2/7 T^{7/2}$ instead of the temperature can be transformed to

$$
- \frac{d^2 f}{dy^2} = \frac{h}{\kappa_p B^2}.
$$

(3)

This transformation eliminates the geometry of the flux tube and leads to a standard Poisson equation which can be solved analytically. The solution which satisfies the boundary condition $F(a) = F'(b) = 0$ is

$$
F(y) = \int_a^b G(y, x) \frac{h}{\kappa_p B^2} \, dx; \quad G(y, x) = \begin{cases} 
\frac{(y-a)(b-x)}{(b-a)} & : y \leq x \leq b \\
\frac{(b-y)(x-a)}{(b-a)} & : a \leq x \leq y
\end{cases}
$$

(4)

The general solution of eq.(3) is $f(y) = F(y) + A_1 + A_2 (y-a)/(b-a)$, where $A_1$ and $A_2$ are correlated to the boundary values of $f$ by $A_1 = f_a$; $A_2 = f_b - f_a$. Inserting this ansatz into the boundary conditions (2) yields

$$
f_b = f_a + \frac{g(T(f_a))}{\kappa_p B(a)} - F'(a); \quad f_a = f_b + \frac{g(T(f_b))}{\kappa_p B(b)} + F'(b)
$$

(5)

where the function $g(T)$ is the left hand side of eq.(2) and $T(f) = (7/2f)^{2/7}$. These non-linear equations determine the boundary values $f_a$ and $f_b$. Having found $f_a, f_b$ the temperature profile is calculated using $F(y)$. With fixed values of $\epsilon_{eff}, \delta_1$ and $\gamma$, three control parameters which determine the solution of eq.(5). These are the total input power $Q$, which yields a factor to $F'$, the pressure $p_0$ and the factor $\kappa_p$ in the thermal conductivity. The symmetry of the boundary conditions is disturbed by the magnetic field, which in general is different on both ends of the flux tube. Furthermore, the constants in the function $g$ may also differ on both sides. Adding the two equations in (5) yields

$$
\frac{g(T_a)}{B(a)} + \frac{g(T_b)}{B(b)} = \int_a^b \frac{h(y)}{B^2} \, dy.
$$

(6)

The left hand side of this equation is positive and has a positive minimum since the function $g(T)$ has a positive minimum at $T = T_{eff}$. As a consequence the integrated source term on the right hand side must be larger than a threshold value, otherwise a solution does not exist. The reason is the ionisation in the recycling layer. The second
term in eq.(1) makes the heat flux \( g(T) \) non-monotonic in \( T \). With \( R \to 0 \) the threshold in heating power also shrinks to zero.

**Numerical results.** After introducing the dimensionless variable \( X = T/T_{eff} \) and \( kT_{eff} = R\kappa_{eff}/\delta \), eqs.(5) have been solved numerically. The source term is modelled by \( h(y) = Q_0 \cdot 0.01/(0.01 + (y - y_0)^3) \). \( Q_0 \) is proportional to the total power input into the flux tube and \( y_0 \) describes the location of the maximum power input. Further dimensionless constants are: 
\[
\begin{align*}
C_a &= [7C_1/2\kappa_0B(a)T_{eff}^3]^{2/7}, \\
C_b &= [7C_1/2\kappa_0B(b)T_{eff}^3]^{3/7}, \\
D_a &= F'\kappa_0B(a)/C_1\sqrt{T_{eff}}, \\
D_b &= F'\kappa_0B(b)/C_1\sqrt{T_{eff}}.
\end{align*}
\]

The constant \( C_1 \) is proportional to the pressure in the flux tube \( C_1 = \delta p_0/2(\gamma e + \gamma i)/m_i \). \( F'(a) \) and \( F'(b) \) depend on the source function \( h(y) \), both are proportional to the total power input \( Q_{tot} \) or \( Q_0 \). In these dimensionless variables the boundary conditions are

\[
\begin{align*}
\left( \frac{X_a}{C_a} \right)^{7/2} &= \left( \frac{X_a}{C_a} \right)^{7/2} + \sqrt{X_a} + \frac{1}{\sqrt{X_a}} - D_a, \\
\left( \frac{X_b}{C_b} \right)^{7/2} &= \left( \frac{X_b}{C_b} \right)^{7/2} + \sqrt{X_b} + \frac{1}{\sqrt{X_b}} + D_b.
\end{align*}
\]

The eqs.(7) are two coupled polynomials of 8th order, in general they have more than one solution. In case of symmetry \( B(a) = B(b), C_a = C_b, D_a = D_b \) two solutions are on the symmetry line \( X_a = X_b \), however, there are also two solutions off the symmetry line. Fig.1 shows these 4 solutions depending on the control parameter \( Q_0 \).

![Figure 1: Bifurcation of the temperature on the target plates. Solution of eqs.(7) with \( C_a = C_b = 4.0 \) and \( kT_{eff} = 10 \text{ eV} \). The x-axis is \( Q_0 \propto \) heating power. \( X = T/T_{eff} \) is the norm. temperature. The curves represent the temperature on one of the target plates.](image-url)

Below a threshold in \( Q_0 \) there are no solutions of eqs.(7); above the threshold a small region exists with two solutions followed by a region of four solutions. In this region the
temperatures at the boundaries $y = a$ and $y = b$ are: case 1) $X_a = T_1$ and $X_b = T_1$; case 2) $X_a = T_2$ and $X_b = T_2$; case 3) $X_a = T_3$ and $X_b = T_4$; case 4) $X_a = T_4$ and $X_b = T_3$. The four temperature profiles are shown in Fig. 2a. A remarkable result is the existence of two asymmetric solutions $X_a \neq X_b$ although the boundary conditions and the source term are symmetric to the midplane of the flux tube. As expected, asymmetric boundary conditions lead to asymmetric solutions (Fig. 2b). These asymmetries arise from the magnetic field which is asymmetric along the flux tube or from the asymmetry of the source function $h(y)$.

![Figure 2: Temperature profiles along the flux tube. Fig. 2a (left): Symmetric case $C_a = C_b = 4.05$, $Q_o = 2500$, $y_o = 0.5$. Fig. 2b (right): Asymmetric power input $C_a = C_b = 4.05$, $Q_o = 2500$, $y_o = 0.6$, $a = 0$, $b = 1$.](image)

Summary and conclusions. The model of a high recycling layer in front of two target plates at the ends of a flux tube leads to non-linear boundary conditions of the heat conduction equation. The heat conduction equation can be transformed to a linear Poisson equation and be solved analytically. In general four solutions exist leading to four temperature profiles. Asymmetric solutions with different temperatures on the target plates ($T_a \neq T_b$) can exist although all boundary conditions are symmetric to the midplane of the flux tube. The thermal stability of the solutions has not yet been considered. The bifurcation point of the multiple solutions depends on the power input and the pressure $p_o$ on the flux tube — as in ref. [1]. However, the bifurcation also depends on the magnetic field on the target plates ($B(a) \neq B(b)$) and the asymmetry of the power input.

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References.