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Impurity Dynamics in Stellarator W7-AS Plasmas

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Abstract

Numerical efforts to understand the neoclassical transport of impurities in stellarator plasmas have been undertaken. The new code solves the radial continuity equations for each ionization stage of the impurity ions for given background plasma profiles and magnetic configuration. An analytic description of the neoclassical transport coefficients based on numerical results from the DKES (Drift Kinetic Equation Solver) code and monoenergetic Monte-Carlo calculation (C.D. Beidler *et al*., EPS 1994), is here applied for impurity transport coefficients. The transition between the different charge states due to the ionization and recombination in balance equation is described by using the ADAS (Atomic Data and Analysis Structure) database. The impurity behavior in some typical discharges from W7-AS with moderate (NC) and improved energy confinement (HDH) has been considered.

It is shown that the spatial distribution results from the competition between the radial electric field and the thermal force (which together produce a convective flux), and the diffusive term, which flattens the radial impurity distribution. The impurity ions are localized at the radial position where the convective flux goes through zero. It is also shown that for typical stellarator discharges there is no pronounced temperature screening effect as in tokamak plasmas.

Keywords:

Impurity dynamics, stellarator physics

1. Introduction

Development of the Stellarator Impurity Transport code (SIT) has been started, aiming to have a numerical tool for better understanding of the existing experimental observations of impurities in stellarator plasmas, and ultimately to prove the feasibility of the stellarator as a reactor. The code solves the system of continuity equations (averaged over the magnetic flux surfaces) for impurity ions in each charge stage, coupled due to the ionization and recombination. In the current stage it calculates the evolution in time and space of impurity ions coming from the wall or due to the pellet ablation. In this sense SIT code is similar to tokamak impurity codes MIST [1] or STRAHL [2]. The new code evaluates impurity behavior in the frame of the stellarator-specific neoclassical transport, which unlike the tokamak configuration is strongly dependent on magnetic topology and the radial electric field. In the case of a test impurity approximation, considered here, the radial electric field is evaluated from the ambipolarity condition in the background plasma and depends on neoclassical transport coefficients for electrons and ions. An analytical description of the neoclassical transport coefficient for the background plasmas (based on numerical results from the DKES code [3] and monoenergetic Monte Carlo simulations) was generalized to impurity ions of arbitrary mass and charge state and used in the code as a neoclassical transport model for impurities. The reduction of Pfirsch-Schluter convection due to the radial electric field and its impact on impurity dynamics has been included. Calculations were performed for given plasma density and temperature profiles. This can be justified because the impurity flux density is ~ $\sqrt{m_i/m_e}$ times larger than the plasma diffusion rate and, in the case of very small impurity concentration, when the radiation is negligible in the plasma energy balance.

2. Impurity density equations, transport coefficients and boundary conditions

The SIT code solves the set of coupled time-dependent one-dimensional continuity equations for

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impurity density, n_j^I for each ionization stages, j of a given impurity species, I, averaged over the magnetic surfaces:

$$\frac{\partial n_j^I}{\partial t} + \frac{1}{r}\frac{\partial}{\partial r}(r\Gamma_j^I) = S_j^I, \quad j = 1, 2 \ldots \qquad (1)$$

Here the impurity flux Γ_j^I can be written as $\Gamma_j^I = -D_j^I \nabla n_j^I + n_I V_j^I$. In general the radial diffusion coefficient $D_j^I = D_{1,1}^{I,j} + D_{an}^I$ and the convective velocity, $V_j^I = V_{Ware}^I + V_{an}^I$ consist of a neoclassical and an anomalous term. Since in stellarators the Ware pinch can be neglected, the neoclassical impurity flux Γ_j^I can be written as (with a slightly different definition of the neoclassical transport coefficients $D_{nk}^{I,j}$ compared to Ref. [1]):

$$\Gamma_j^I = -D_{11}^{I,j}\left(n_j^I\right)' + n_{jI}^I \frac{D_{11}^{I,j}}{T}$$
$$\left\{Z_I E_r - \left(\frac{D_{12}^{I,j}}{D_{11}^{I,j}} - \frac{3}{2}\right)T'\right\} \equiv \Gamma^{diff} + \Gamma^{conv}, \qquad (2)$$

The first term in equation (2) describes the impurity diffusion flux, whereas the rest terms constitute the convective flux. Here and further we assume that all impurity ions have the same temperature, $T_I \approx T_i \equiv T$ and the prime denotes the radial derivative of density and temperature. The sign and the value of the neoclassical convective flux (two second terms in (2)) depend on the radial electric field and the thermal force. We assume that functional dependence of the neoclassical transport coefficients $D_{n,k}^{I,j}$ for impurity species on plasma parameters is the same as for the background plasma coefficients. For background plasma species the mono-energetic radial transport coefficients, $D_{n,k}^I$ (with $n, k = 1, 2$ and $I = e, i$) are obtained by energy convolution of the mono-energetic transport coefficients with a Maxwellian distribution function. For various magnetic field configurations, the databases of these coefficients calculated by DKES code are fitted based on traditional analytic theory, with axisymmetric contributions in the plateau collisionality regime taken into account. Finally, the functional dependence of these coefficients on plasma parameters is expressed analytically, whereas the radial dependence for a given magnetic field configuration is calculated by using the least-squares fitting of numerical results [4,5]. This semi-analytic description makes possible a rapid numerical simulation of the impurity transport.

For the collisional case we are using a new expression for the Pfirsch-Schluter monoenergetic coefficients, $D_{PS}^{I,j}(E_r, V, r)$ [6] :

$$D_{PS}^{I,j}(E_r, V, r) \approx D_{PS}^{I,j}\left(\left(1 + \left(\frac{3\nu_I \Omega_E R_0^2}{\iota^2 V^2}\right)^2\right)^{-1} + \left(\frac{\epsilon_t \iota}{b_t}\right)^2\right),$$

$$D_{PS}^{I,j} \approx \frac{4}{3}\frac{VR_0}{\Omega_j}\left(\frac{b_t}{\epsilon_t \iota}\right)^2 \nu_I, \qquad (3)$$

In this coefficient the dependence on the radial electric field, E_r is taken into account (the second term in (3)). Here $\Omega_E = E_r/aB_0, \Omega_j = Z_j eB_0/m_I, x \approx m_I V^2/2T$, a and R_0 are the average minor and major radii, ν_I is the collisional frequency for impurity ions, and b_t is the "effective" toroidal curvature, V is the velocity of the impurity ion of mass m_I, ι is the rational transform value and $\epsilon_t \approx a/R$. The physical meaning of the new term in (2) can be understood as follows. As the collisionality increases, even modest values of the radial electric field result in a poloidal rotation which is sufficient to suppress charge separation, and Pfirsch-Schlüter diffusion is reduced to the classical level (the last term in (3)). Calculation show, that the dependence on electric field is important particularly for high Z impurity ions. The radial electric field can be found from ambipolarity condition, assuming the tracer approximation for impurity ions, $\Gamma^i(E_r) \approx \Gamma^e(E_r)$. Here we are also ignoring the variation of electric field in poloidal direction (which is difficult to estimate), but which can have a noticeable impact. Excluding the term with electric field from ambipolarity conditions (2), one obtains:

$$\Gamma_j^I = -D_{11}^{I,j} n_j^I \left(\frac{(n_j^I)'}{n_j^I} + \beta\frac{n'}{n} + \gamma_T \frac{T'}{T}\right) \qquad (4)$$

where $n = n_i \approx n_e \gg Z_j n_j^I$, and

$$\gamma_T = Z_j \frac{D_{11}^e\left(\frac{D_{12}^e}{D_{11}^e} - \frac{3}{2}\right) - Z_i D_{11}^i\left(\frac{D_{12}^i}{D_{11}^i} - \frac{3}{2}\right)}{D_{11}^e + Z_i^2 D_{11}^i} + \left(\frac{D_{12}^I}{D_{11}^I} - \frac{3}{2}\right)$$

$$\beta = Z_j \frac{D_{11}^e - Z_i D_{11}^i}{D_{11}^e + Z_i^2 D_{11}^i} \qquad (5)$$

In equation (4) the first term is the diffusive flux driven by the impurity density gradient. The two last terms here are proportional to the plasma density and temperature gradients, respectively. These terms represent impurity convection. Whether impurity will localize at the plasma edge or at the center depends on the sign of γ_T and β coefficients in (4). The second term in (4) describes impurity convection due to the electron density gradient (driven by the radial electric field) and is proportional to $\beta \approx -Z_j/Z_i$, since $D_{11}^e \ll D_{11}^i$. Here Z_i is the charge number for background ion, and Z_j for impurity ions. This term is always negative that corresponds to inward flux. The third term in (4) is proportional to the thermal diffusion coefficient, γ_T. When $\gamma_T > 0$, an outward drift (in positive radial direction) happens and the "temperature screening effect" takes place (notice, that D_{1k}^I coefficients are positive).

The source term in equation (2) describes the atomic processes of recombination and ionization between the different charge states:

$$S_j^I = n_e \left\{ S_{j-1} n_{j-1}^I - (S_j + \alpha_j) n_j^I + \alpha_{j+1} n_{j+1}^I \right\} \qquad (6)$$

Here the ionization S_j and recombination α_j rates for impurity ions are taken from ADAS. In the case of trace impurity concentrations, considered below, the radial electric field E_r is determined mainly by background plasma and can be evaluated from the differential equation, which permits a smooth transition in case of different roots of the ambipolarity condition [7,8].

$$\epsilon_\perp \frac{\partial E_r}{\partial t} - \frac{2}{\rho^2} \frac{\partial}{\partial \rho} \left(\rho^2 \frac{\eta \tilde{E}_r}{B^2} \right) + \left(\frac{\tilde{E}_r}{B} \right)^2 \left(\frac{\partial \eta}{\partial E_r} \right)$$
$$= e \left(\sum_i z_i \Gamma(E_r)_i - \Gamma_e(E_r) \right),$$
$$\tilde{E}_r \equiv E_r' - E_r/\rho \qquad (7)$$

Here $\Gamma_{e,i}$ are the radial fluxes of background components, and $\epsilon_\perp$ is the dielectric constant. The E_r dependence of the viscosity coefficient $\eta(E_r)$ has been taken to be in accordance with as $\eta \approx const.$ for $E_r \leq E_{res} \equiv \iota v_{Te} B\rho/R$ and $\eta \propto E_r^{-\beta}$ for $E_r \geq E_{res}, \beta \approx 4$ [7]. The boundary conditions are the symmetry condition at the center $E_r' = 0$ and at the edge either $E_r = 0$ in limiter plasmas, otherwise $\Gamma_e \approx \sum z_i \Gamma_i$ at the last closed magnetic surfaces.

The impurity density equations (1) are integrated from the plasma centre ($\rho = 0, \rho \equiv r/a$) to the edge ($\rho = 1$). The boundary conditions at the edge depend on the thermal velocity of impurity ions, V_T^I, and the recycling coefficient, R_I

$$\Gamma_j^I = 0.25 R_I n_j^I V_T^I \qquad (8)$$

The boundary conditions at the centre are $\Gamma_j^I(0, t) = 0$. The density of neutral atoms entering the plasma edge, $r = a$ decays inside the plasma as,

$$n_0(r) = n_0(a) \frac{a}{r} \exp\left(- \int_r^a \frac{n_e s_0}{V_0} \mathrm{d}r \right), \qquad (9)$$

where V_0 is the thermal velocity and s_0 is the ionization rate of the injected impurity atoms. The impurity influx can last during the all operation time or be terminated after some short time, simulating the case of pellet injection at the edge region.

3. Results and conclusions

Calculations have been performed for several impurity species (Fe, C) for background plasmas profiles typical for the HDH and NC configuration in W7-AS (see

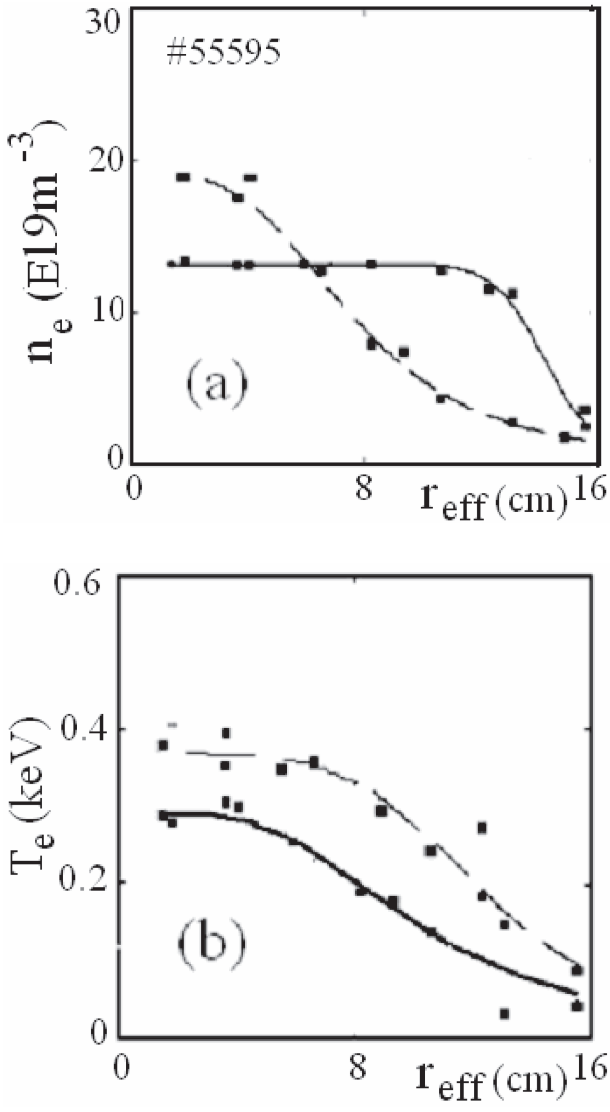

Fig. 1 Background density (a) and temperature (b) profiles in HDH and NC (dashed) modes of W7-AS [9,10].

Fig. 1) [9]. Observation show that the NC discharge is characterized by impurity accumulation at the centre, while in the HDH discharge impurity is mainly radiated from the edge region [9]. A source of neutral impurity atoms was specified at the wall and can be varied in time. The calculated radial electric field (see Fig. 2) is in both cases negative, but for HDH is rather small, particularly at the edge. The thermal diffusion coefficients for Fe impurity ions for the HDH mode are shown in Fig. 3. γ_T for each charge state is negative, which predicts the accumulation at the centre. A steady state impurity distribution is achieved, when the source term is switched off and the convective flux turns to zero. A comparative analysis of the diffusive and the convective fluxes shows that, in steady state impurities are localized radially, where these fluxes are compensate each other. In the HDH mode of operation the flat density profile in the bulk plasma slows down the inward impu-

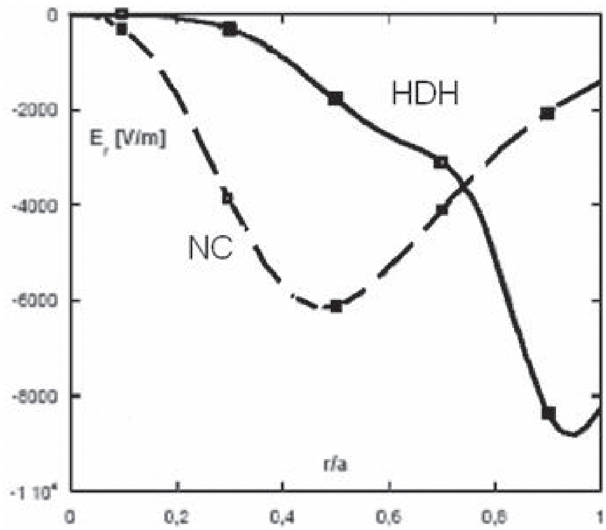

Fig. 2 The radial electric field for NC (dashed) and HDH modes of the W7-AS.

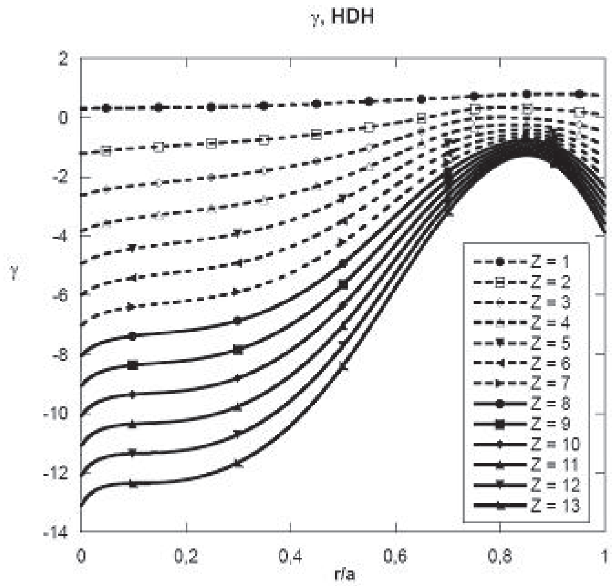

Fig. 3 The thermal diffusion coefficients, γ_T, for Fe impurity ions in HDH mode of W7-AS

rity penetration in this region. In the NC mode, the convective term dominates the transport and brings about a rapid accumulation of impurities in the bulk plasma. In both operational modes impurities are predominantly in the Pfirsch-Schlüter regime, $\nu_I^{eff} \gg 1$, so that the correct Pfirsch-Schlüter description was important. As it was mentioned above, the Pfirsch-Schlüter diffusion should include the dependence on the radial electric field (see Eq. (3)). In this expression the reduction of the Pfirsch-Schluter convection due to the poloidal rotation caused by the radial electric field has been taken into account and, its impact on the impurity dynamics due to the change in $D^I_{n,k}$ coefficients and electric field was investigated. As one can see from equation (3), beyond a critical collisionality of background plasma, $\nu^{eff} > \iota V/3R_0\Omega_E$, the poloidal rotation due to the radial electric field is sufficient to suppress the charge separation responsible for Pfirsch-Schlüter diffusion, and the losses are then reduced to the classical level. So, as the collisionality increases, D^I_{PS} drops to its classical value and then become independent of electric field. This drop in D^I_{PS} value results in a lowering of the electric field value.

The thermal diffusion coefficient is negative, for all considered profiles, thus excluding the temperature screening effect, except for the case with very low and flat density profiles and low charge states. Calculations show that the impurity convection is very sensitive to the value of the electric field and, hence to the plasma profiles. In the HDH mode, convection in the bulk region is smaller than in the NC mode. At 0.05 sec impurities are still at the edge, whereas in the NC mode impurities are already distributed in the centre. In both regimes impurity ions finally accumulate eventually at the centre and the "temperature screening" coefficient is always negative (no "screening"). However, the inward convective drift in the case of the HDH mode is smaller than in the NC mode, because of the flat density gradient and, ultimately, due to the smaller value of the electric field. Figure 4 shows, that in both cases Fe impurity ions are accumulating in the centre. Calculation shows that, in the frame of correct stellarator-relevant neoclassical theory, it is impossible to explain retention of impurities at the plasma edge, seen in the HDH mode of W7-AS. However, the typical timescales needed to reach steady state $\sim$ 0.1 sec in the HDH plasma, and 3 – 5 times less in the case of NC, are of the same order

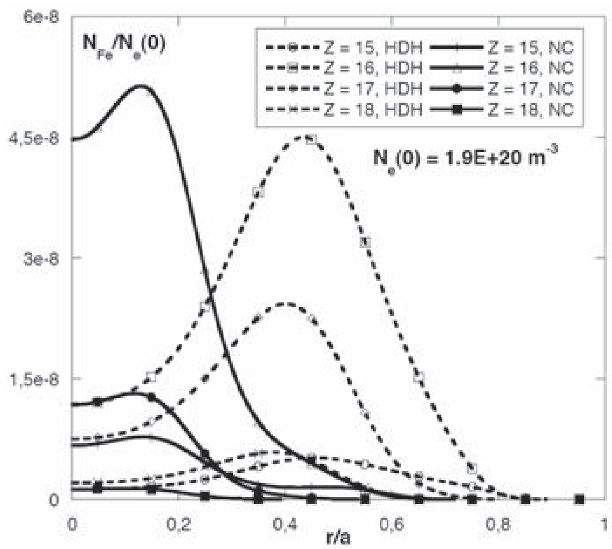

Fig. 4 Evolution of F_e ions profiles in the HDH & NC discharges of W7-AS.

of magnitude as in experiments [9].

Calculations show that the radial electric field plays a dominant role in impurity dynamics and localization. In case of low density and high temperatures (positive radial electric field) one can probably expect impurity retention at the plasma edge.

The different impurity behavior in NC and HDH modes could be probably explained as a result of change in sign of the radial *E*-field at the edge from negative (in NC or H^* mode) to positive, with pronounce screening effect. The main signature of HDH is the increase of average beta value, which causes a strong ergodization of the edge area. At some higher beta value, a stochastic area is setting up in the edge region beyond the pedestal region. Calculations done in [10] for Helias finite-beta equilibrium case confirm, that with increasing beta the width of the macroscopic islands at the edge increases, showing the trend to overlapping of magnetic islands and ergodization. The ergodization could change a sign of the electric field in this area, making it positive the electron losses become dominant. In NC mode beta is small and the E-field is negative in entire area. With increase of average density we are approaching H^* mode, which has higher beta value, than in NC mode, but not sufficient for effective ergodization. When average pressure (beta) is increasing further the edge region is ergodized stronger, and the positive *E*-field prevails over the diffusivity. This model is however not yet included into code.

In conclusion: the main goal of our consideration was to simulate impurity behavior in the case of stellarator W7-AS plasmas in the frame of neoclassical theory. For this purpose an analytical description of the neoclassical transport coefficient for the background plasmas (based on numerical results from DKES code and monoenergetic Monte Carlo simulations) has been applied for the impurity ions of arbitrary mass and charge state. The code describes the evolution in time and space of impurity concentration of different charge states, the corresponding convective and diffusive fluxes, and radiation intensities for given plasma profiles. It was confirmed that, in contrary to tokamak, a temperature screening effect does not appear in stellarator configurations. In the HDH mode the flat density profile in the bulk plasma slows down the inward impurity penetration. In the NC mode, the convective term dominates and brings about a rapid accumulation of impurities in the bulk plasma. However, the impurity screening in the HDH mode can not be understood on the base of the pure neoclassical transport, used here.

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