Neoclassical transport of heavy impurities with poloidally asymmetric density distribution in tokamaks

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Abstract

Heavy impurity ions in tokamaks are not always evenly distributed over flux surfaces. For instance, toroidal plasma rotation can give rise to a centrifugal force large enough to push impurities to the outboard side of the torus, or ion-cyclotron resonance heating of minority ions can cause inboard impurity localization. It is shown that such poloidally uneven distribution of the impurity ions can enhance or reduce their neoclassical transport by one to two orders of magnitude, or even reverse the direction of the neoclassical impurity convection, depending on the level of poloidal asymmetry of the impurity density distribution and on the ratio of the logarithmic ion temperature gradient to the logarithmic density gradient.

I Introduction

Impurity ion accumulation in the core of tokamak and stellarator plasmas has been an issue of concern for several decades. If such accumulation cannot be avoided, the deleterious combination of fuel dilution (mainly by light impurities) and radiation (mainly by heavy impurities) would strongly limit the possibilities of realizing a fusion power plant based on these concepts. Neoclassical transport plays an important role for impurities both in the core [1, 2, 3, 4, 5, 6, 7] and in the pedestal of H–modes [8, 9, 10], where turbulent transport is largely suppressed.

Standard neoclassical transport theory [11, 12], which in particular assumes approximately homogeneous particle densities over the magnetic flux surfaces, predicts inward transport of highly charged impurity ions, driven by the density gradient of the *bulk* ions, in all collisionality regimes in both stellarators and tokamaks [13]. Fortunately, in tokamaks the impurity particle flux driven by the temperature gradient is outward in most regimes, but no such "temperature screening" exists (theoretically) in stellarators [14, 15]. In experiments, turbulence is often observed to expel impurities from the plasma core [5, 6, 16, 17, 18, 19], but central accumulation of heavy impurities is nevertheless common, particularly in transport barriers where the turbulence is suppressed and in plasmas without strong heating in the center.

However, there are effects not accounted for by conventional neoclassical theory, and in the present paper we consider one such effect, which appears to be particularly important. In conventional neoclassical theory, it is usually assumed that the density of all species is approximately constant on flux surfaces. Unless the plasma rotates very rapidly, this is a good approximation for the electrons and bulk ions, but it can easily fail for heavy impurities. Indeed, their emission is frequently observed to be poloidally asymmetric [20, 21, 22]. There are several possible reasons for such an asymmetry. The centrifugal force associated with the toroidal rotation pushes particles to the outboard side of each flux surface [23], so that the out–in asymmetry of the density of species a, with mass m_a and temperature T_a, becomes of order ϵM_a^2, where ϵ denotes the inverse aspect ratio and

$$M_a^2 = \frac{m_a \omega^2 R^2}{2T_a},$$

the squared Mach number of the species in question [24, 25], with ω the plasma toroidal angular velocity, and R the major radius. A mechanism for inboard localization of impurity ions is provided by ion cyclotron resonance heating (ICRH) of minority ions, which tends to increase the perpendicular velocity of these particles, causing the mirror force to push them to the outboard side. A poloidal field will then appear to maintain quasineutrality, and this field drives highly charged impurities to the inboard side [20, 21, 22, 26]. The effect is predicted to be large in the presence of strong ICRH power. It can produce in–out asymmetries of W in plasmas which are at rest or are weakly rotating, or it can completely balance the out–in centrifugal asymmetry of strongly rotating plasmas [27]. Also, neutral beam injection (NBI) can produce temperature anisotropies of the background deuterium which can non–negligibly impact the W density distribution, depending on the beam geometry [27].

In a series of papers published 10-15 years ago, it was established that a nonuniform poloidal distribution of impurities affects the neoclassical transport in tokamaks [24, 28, 29, 30, 31, 32]. In these works, the poloidal distribution of impurities resulting from toroidal rotation and friction against the bulk ions was calculated self-consistently. A very recent work [33] extends those previous studies, and obtains analytical formulae for the complete diffusive and convective components of the impurity flux in a rotating plasma, for arbitrary impurity charge and Mach number in the Pfirsch–Schlüter regime. The analytical formulae are found to be in perfect agreement with the numerical results of the NEO code [34, 35]. Here, instead, we use a complementary approach, and we extend previous studies in order to also compute the effect of other redistribution mechanisms, such as ICRH, on the impurity neoclassical transport in the limit of high impurity charge in the Pfirsch–Schlüter regime. Thereby, in the present paper, we allow for an arbitrary poloidal impurity distribution, not necessarily produced by centrifugal effects only. We find that a poloidally uneven distribution of the impurity ions can enhance or reduce their neoclassical transport by one to two orders of magnitude with respect to its conventional value, or even reverse the direction of the neoclassical impurity convection, depending on the level of poloidal asymmetry and on the parameters of the bulk plasma.

II Orderings and kinetic equation

The mathematics is very similar to that of Ref. [31], where the plasma was taken to consist of low–collisionalty ions (i), electrons (e) (in the banana regime), and highly charged, collisional impurity ions (z) (in the Pfirsch–Schlüter regime), with a charge z. In order to keep our treatment as general as possible, the plasma is allowed to rotate toroidally, but for simplicity the bulk ion Mach number is assumed to be small, $M_i \ll 1$, whilst $M_z = O(1)$. Of course the present results are valid also in the case that the plasma is at rest, that is when $\omega = 0$. In contrast to Ref. [31], however, other particle species may also be present, whose behavior is not further specified. For instance, there may be minority ions interacting with an ICRH operator causing these ions to have a poloidally non-uniform distribution, implying the existence of a poloidally varying electrostatic potential $\tilde{\phi}(\psi, \theta)$. This potential is ordered to be too weak to influence the distribution of the bulk plasma, but strong enough to affect the impurities,

$$\frac{e\tilde{\Phi}}{T_i} \sim \frac{1}{z} \ll 1.$$

As in Ref. [31], a conventional expansion of the drift kinetic equation using these orderings gives for the bulk-ion distribution function

$$f_i = f_{i0} \exp\left(-\frac{e\tilde{\Phi}}{T_i} + M_i^2\right) - \frac{Iv_\parallel}{\Omega_i}\frac{\partial f_{i0}}{\partial \psi} + h_i(H, \mu, \psi, \sigma), \tag{1}$$

$$f_{i0} = N_i(\psi)\left(\frac{m_i}{2\pi T_i}\right)^{3/2} \exp\left(-H/T_i\right), \tag{2}$$

where $\Omega_i = eB/m_i$ is the ion cyclotron frequency, $\mu = m_i v_\perp^2/(2B)$ the magnetic moment, $H = m_i v^2/2 + e_i\bar{\Phi} - m_i\omega^2 R^2/2$ the energy. The indices $\perp$ and $||$ indicate directions perpendicular and parallel to the magnetic field respectively. In Eq. (1), which corresponds to Eq. 6 in Ref. [31], $\tilde{\Phi}$ is the poloidally varying part of the electrostatic potential and the function h_i satisfies

$$v_\parallel \nabla_\parallel h_i = C_i^l(f_i), \tag{3}$$

where C_i^l is the linearized ion collisional operator. The magnetic field is written as $\boldsymbol{B} = I(\psi)\nabla\varphi + \nabla\varphi \times \nabla\psi$, where φ is the toroidal angle and ψ is the poloidal flux function, and a prime denotes differentiation with respect to ψ. The angular frequency of the rotation is the same for all species, $\omega = -d\bar{\Phi}/d\psi$, where $\bar{\Phi}(\psi)$ is the lowest-order electrostatic potential, appearing in the expression of the lower order distribution function f_{i0} in Eq. (2). The velocity $\boldsymbol{v}$ is measured in the rotating frame, and the independent velocity-space variables are H and μ.

From the solution of Eq. (3), one can calculate the ion-impurity friction force, as in Eq. 10 of Ref. [31],

$$R_{zi\|} = -\frac{p_i I}{\Omega_i \tau_{iz}} \left(\frac{p_i'}{p_i} - \frac{3}{2}\frac{T_i'}{T_i} \right) + \frac{m_i n_i}{\tau_{iz}} \left(u - \frac{K_z}{n_z} \right) B, \tag{4}$$

where

$$\tau_{iz} = \frac{3(2\pi)^{3/2} \epsilon_0^2 \sqrt{m_i} T_i^{3/2}}{n_z z^2 e^4 \ln \Lambda},$$

is the ion-impurity collision time, bf p_i and n_i are respectively the pressure and density of the main ions, whereas n_z is the density of the impurity species. Moreover,

$$u = \frac{\tau_{iz}}{n_i B} \int v_\| \nu_{iz} h_i d^3 v$$

is a flux function and the flow speed of the impurities along the magnetic field has been written as [28]

$$V_{z\|} = -\frac{I}{B}\frac{d\bar{\Phi}}{d\psi} + \frac{K_z(\psi) B}{n_z}. \tag{5}$$

The flux function $K_z(\psi)$ appearing in this expression is proportional to the poloidal velocity and can be determined from the parallel momentum equation of the impurities

$$m_z n_z \boldsymbol{V}_z \cdot \nabla \boldsymbol{V}_z \cdot \boldsymbol{b} = -z n_z e \nabla_\| \tilde{\Phi} - T_i \nabla_\| n_z + R_{zi\|}$$

by multiplying this equation by B/n_z and taking the flux-surface average, denoted by angular brackets, giving the solubility constraint

$$\left\langle \frac{B R_{zi\|}}{n_z} \right\rangle = 0, \tag{6}$$

resulting in the relation

$$K_z = \left[u \left\langle B^2 \right\rangle - \frac{T_i I}{e} \left(\frac{p_i'}{p_i} - \frac{3}{2}\frac{T_i'}{T_i} \right) \right] \Big/ \left\langle \frac{B^2}{n_z} \right\rangle. \tag{7}$$

The most difficult of our equations is of course the kinetic equation (3) for h_i, which was solved in various limits in Ref. [31], including in particular the poloidal asymmetry produced by centrifugal effects. The details of this solution are not of any great concern to us since all the information from h_i that we eventually need is the constant u determining the friction force (4). In the simplest case of trace impurities, $n_z z^2/n_i \ll 1$, one obtains

$$u = -0.33 \frac{f_c I}{e \langle B^2 \rangle} \frac{\partial T_i}{\partial \psi}, \tag{8}$$

if the ion-ion collision operator is approximated by a pitch-scattering operator plus a momentum conserving term (Section II.B in Ref. [31]). Here, f_c denotes the usual effective fraction of circulating particles,

$$f_c = 1 - f_t = \frac{3 \langle B^2 \rangle}{4} \int_0^{\lambda_c} \frac{\lambda d\lambda}{\left\langle \sqrt{1 - \lambda B} \right\rangle},$$

which is $f_c \simeq 1 - 1.46\epsilon^{1/2}$ in a large–aspect–ratio torus with circular cross section.

III Impurity dynamics

We now turn our attention to the neoclassical transport of impurities. The particle flux can be related to the parallel friction force by the usual flux-friction relation [12] and is equal to

$$\langle \mathbf{\Gamma}_z^{neo} \cdot \nabla \psi \rangle = -\frac{1}{z} \langle \mathbf{\Gamma}_i^{neo} \cdot \nabla \psi \rangle = -\left\langle \frac{I R_{zi\|}}{zeB} \right\rangle. \tag{9}$$

Substituting the friction force (4) and using Eq. (7) gives

$$\langle \mathbf{\Gamma}_z^{neo} \cdot \nabla r \rangle = \frac{q^2 n_i \langle n_z \rangle T_i}{\epsilon^2 m_i \langle \Omega_i^2 \rangle n_z \tau_{iz} z} \left[\left(\left\langle \frac{n}{b^2} \right\rangle - \left\langle \frac{b^2}{n} \right\rangle^{-1} \right) \left(\frac{d \ln p_i}{dr} - \frac{3}{2} \frac{d \ln T_i}{dr} \right) - \tilde{u} \left(1 - \left\langle \frac{b^2}{n} \right\rangle^{-1} \right) \right], \tag{10}$$

where $r(\psi)$ is an arbitrary flux-surface label, q/ϵ is defined to be equal to $I/(d\psi(r)/dr)$,

$$n = \frac{n_z}{\langle n_z \rangle},$$

$$b^2 = \frac{B^2}{\langle B^2 \rangle},$$

and we have written

$$\tilde{u} = \frac{eu \langle B^2 \rangle}{qT_i} \simeq -0.33 f_c \frac{d \ln T_i}{dr}.$$

The last equality holds in the limit (8). Since the present derivation is not based on a large–aspect–ratio–expansion, the expression of the impurity flux in Eq. (10) is valid for arbitrary aspect ratio. However, the calculation of the flux function u as obtained in Eq. (8) requires the analytical solution of the bulk–ion drift–kinetic equation for h_i, that is Eq. (3) (which is independent of the impurity when $n_z z^2 \ll n_e$). As already mentioned, the analytical solution of the drift–kinetic equation is possible only if the collision operator is approximated by a pitch-scattering operator plus a momentum conserving term. This approximation can be shown to be exact in the limit $\epsilon \ll 1$, but can be expected to work reasonably well also for larger values of ϵ. A more accurate treatment can be found in Ref. [12], but would merely modify the coefficient 0.33 in Eq. (refualfaliten). The main geometric effect in Eq. (8) is encapsulated by the fraction of circulating particles f_c, provided this is appropriately computed (that is not in a large aspect ratio limit). We also note that the product $n_z \tau_{iz}$, at the denominator of Eq. (10), is constant over a flux surface.

The result (10) is identical to Eq. (33) in Ref. [31], but has been derived under slightly more general assumptions. In that article, the distribution $n(\theta)$ of the impurities was calculated self-consistently, i.e., the parallel momentum equation of the impurity fluid was solved to find how the impurities respond to the centrifugal and friction forces. In the present paper, we do not attempt to do so, as we instead allow for the possibility of more particle species such as fast ions produced by NBI or ICRH. These are in general not evenly distributed over the flux surface and therefore cause a poloidal electric field to form in order to maintain quasineutrality. The specific example of minority ions heated by ICRH was considered in Refs. [20, 21] and more recently in [22] (for related studies on turbulent transport, see also Refs. [36, 37], as well as [38], where analytic formulae are applicable also considering more general expressions of the background electrostatic potential, not solely produced by centrifugal effects). The heating acts to increase the perpendicular velocity of the minority ions, thus pushing them into trapped orbits on the outboard side of the torus. These ions are normally so energetic that they are practically collisionless and therefore do not experience much friction from other particle species, so that their poloidal distribution can be taken as given. In order to balance the surplus of energetic ions on the low-field side, an in-out poloidal electric field arises, pulling electrons to the outboard side and thermal ions to the inboard side of the torus. The resulting poloidal density asymmetries are small for most species except highly charged impurity ions, which can in principle be arbitrarily highly localized if $z \gg 1$. We also notice that the ordering $z \gg 1$ implies that terms $O(1/z)$ are not kept. Therefore Eq. (10) does not include any term proportional to the impurity density gradient (that is, in Eq. (10) no diffusion term is present and the impurity flux is solely driven by gradients of the background ions).

In the limit of large aspect ratio and weakly varying impurity density, we take

$$b = 1 - \epsilon\cos\theta + O(\epsilon^2),$$

$$n = 1 + \delta\cos\theta + \Delta\sin\theta + O(\delta^2, \Delta^2, \delta\Delta), \tag{11}$$

where the parameters δ and Δ describe out–in and up–down asymmetries respectively. Although these functions are only known to first-order accuracy in ϵ, δ and Δ, the additional information that $\langle n\rangle = \langle b^2\rangle = 1$ exactly (by the definition of n and b) enables us to calculate the geometric factors in Eq. (10) to second-order accuracy. This is done by writing

$$\left\langle \frac{n}{b^2} \right\rangle = 1 + \left\langle \left(\frac{1}{b^2} - 1\right)\left(n - b^2\right) \right\rangle,$$

$$\left\langle \frac{b^2}{n} \right\rangle = 1 + \left\langle \left(\frac{1}{n} - 1\right)\left(b^2 - n\right) \right\rangle,$$

and noting that the remaining averages can be evaluated to second order in ϵ, δ and Δ, with the result

$$\left\langle \frac{n}{b^2} \right\rangle = 1 + \epsilon\delta + 2\epsilon^2,$$

$$\left\langle \frac{b^2}{n} \right\rangle = 1 + \epsilon\delta + \frac{\delta^2 + \Delta^2}{2},$$

Hence the factors appearing in Eq. (10) are

$$\left\langle \frac{n}{b^2} \right\rangle - \left\langle \frac{b^2}{n} \right\rangle^{-1} = 2\epsilon(\epsilon + \delta) + \frac{\delta^2 + \Delta^2}{2},$$

$$1 - \left\langle \frac{b^2}{n} \right\rangle^{-1} = \epsilon\delta + \frac{\delta^2 + \Delta^2}{2},$$

with an error of third order in ϵ, δ and Δ.

Thus, in the conventional theory, where the impurities are evenly distributed over the flux surface, $\delta = \Delta = 0$, their cross-field flux becomes

$$\langle \mathbf{\Gamma}_z^{neo} \cdot \nabla r\rangle = \frac{2q^2 n_i T_i}{m_i \Omega_i^2 \tau_{iz} z}\left(\frac{d\ln p_i}{dr} - \frac{3}{2}\frac{d\ln T_i}{dr}\right). \tag{12}$$

When the impurities are instead unevenly distributed over the flux surface, this flux is enhanced by a factor $1 + \delta/\epsilon$ (plus terms of second order in δ and Δ). Outboard impurity accumulation enhances the transport [30, 31, 33] and inboard accumulation reduces it. In the analytical expansion (11), the strength of an up–down asymmetry is described by the parameter Δ, which enters only at the second order in the flux surface averages of Eq. (10). Up–down asymmetries are usually very weak in the plasma core, according to the predictions of neoclassical theory. Their impact on the neoclassical impurity transport has been already discussed in detail in Ref. [31].

The opposite limit of strongly localized impurities is more striking. It is instructive to consider the analytical limit of $n(\theta) \propto \delta(\theta - \theta_0)$, where θ_0 is arbitrary, which can be thought to be obtained in the limit of arbitrarily large impurity charge in the presence of a poloidal variation of the electrostatic potential which still fulfils the condition $e\tilde{\Phi}/T_i \ll 1$. Then

$$\left\langle \frac{n}{b^2} \right\rangle = 1 + O(\epsilon),$$

$$\left\langle \frac{b^2}{n} \right\rangle^{-1} = 0,$$

and the transport becomes, to lowest order in ϵ,

$$\langle \mathbf{\Gamma}_z^{neo} \cdot \nabla r \rangle \simeq \frac{q^2 n_i T_i}{\epsilon^2 m_i \Omega_i^2 \tau_{iz} z} \left[\frac{d \ln p_i}{dr} - \left(\frac{3}{2} - 0.33 f_c \right) \frac{d \ln T_i}{dr} \right]. \tag{13}$$

This is an enhancement by a factor $1/(2\epsilon^2)$ over the conventional Pfirsch–Schlüter value (12). Thus, if the impurities are strongly localized anywhere on the flux surface, by whatever mechanism, their neoclassical cross-field transport is greatly enhanced. In the central region of the plasma, if the ratio of the logarithmic temperature gradient to the logarithmic density gradient exceeds a critical value,

$$\eta_i = \frac{d \ln T_i}{d \ln n} \gtrsim \eta_{ic} = \frac{6}{1 + 2 f_t}, \tag{14}$$

this transport is outward and will tend to expel impurities from the core.

The limit just considered, that of strong poloidal impurity localization, can in fact be treated analytically for any impurity collisionality and arbitrary aspect ratio. If $n(\theta) \propto \delta(\theta - \theta_0)$, then it follows that the quantity K_z defined in Eq. (7) must vanish since the parallel flow velocity would otherwise be infinite for $\theta \neq \theta_0$. The friction between the bulk ions and impurities is thus given by Eq. (4) with $K_z = 0$, and the cross-field impurity flux (9) becomes

$$\langle \mathbf{\Gamma}_z^{neo} \cdot \nabla r \rangle = \frac{q^2 n_i \langle n_z \rangle T_i}{\epsilon^2 m_i \langle \Omega_i^2 \rangle n_z \tau_{iz} z} \left[\frac{1}{b_0^2} \left(\frac{d \ln p_i}{dr} - \frac{3}{2} \frac{d \ln T_i}{dr} \right) - \tilde{u} \right],$$

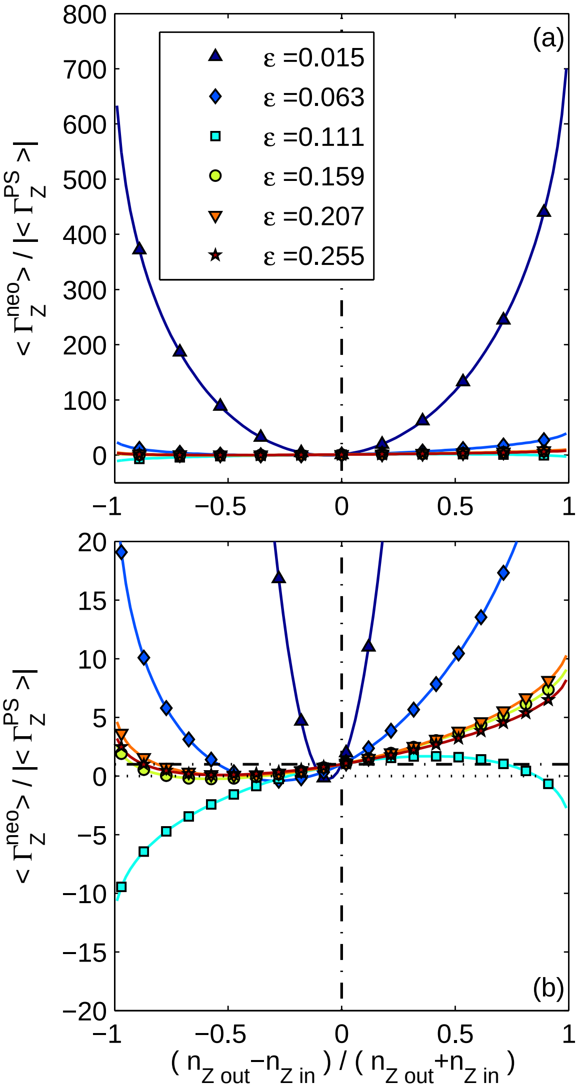

Figure 1: Flux–surface averaged radial flux of heavy impurity, normalized to the absolute value of the Pfirsch–Schlüter flux for poloidally symmetric impurity density, as a function of the level of poloidal asymmetry of the impurity density distribution and for different values of the local inverse aspect ratio ϵ as quoted in the legend (a). The strong enhancements at small values of ϵ with both out–in and in–out asymmetry are illustrated. A zoom around moderate values of the impurity poloidal asymmetry is presented in (b), in linear scale. Positive values correspond to outward flux, negative values to inward flux. The Pfirsch–Schlüter flux is positive (directed outward) at all radii.

where $b_0 = b(\theta_0)$. If we now take the limit used in Eq. (8) and $\epsilon \ll 1$, we again obtain Eq. (13), which is thus seen to hold for all impurity collisionalities in the limit of strong localization.

The implications of Eq. (10) are shown in Fig. 1, where the enhancement factor of the impurity radial flux, defined as $\langle \mathbf{\Gamma}_z^{neo} \cdot \nabla r \rangle / | \left\langle \mathbf{\Gamma}_z^{PS} \cdot \nabla r \right\rangle |$ is plotted as a function of the parameter $(n_{zout} - n_{zin})/(n_{zout} + n_{zin})$, which describes different levels of poloidal asymmetry of the heavy impurity density. Here $\left\langle \mathbf{\Gamma}_z^{PS} \cdot \nabla r \right\rangle$ is the Pfirsch–Schlüter flux in the poloidally symmetric case, and n_{zout} and n_{zin} are the low-field side (LFS) and the high-field side (HFS) impurity densities, respectively, assuming that the impurity density varies according to

$$n = 1 + \delta \cos \theta.$$

Radial profiles of density and temperature, shown in Fig. 2(a), and geometry of a typical H-mode ASDEX Upgrade plasma have been considered (shot ♯27028 at 2.5 s [39]), and the variations of the background plasma profiles and geometry have been consistently taken into account with the variation of the parameter ϵ. The measured η_i profile is compared to the η_{ic} profile in Fig. 2(b), where symbols identify the different radial locations which correspond to the curves presented in Fig. 1. The corresponding values of the inverse aspect ratio ϵ are quoted in the legend of Fig. 1. Consistent with the analytical estimates, we observe that transport is greatly enhanced in the case of very strong poloidal asymmetries, both on the high and the LFS. In contrast, under conditions of weak to intermediate levels of in–out asymmetry, as those which can be obtained by the application of ICRH [22, 26], the neoclassical impurity flux can be significantly reduced.

We also observe that at radial locations at which Eq. (14) is no longer satisfied, but $\eta_i > 2$ (radial domain between $r/a \simeq 0.25$ and $r/a \simeq 0.45$, thereby in particular the curve at $\epsilon = 0.111$ in Fig. 1, corresponding to $r/a = 0.35$ in Fig. 2) the heavy impurity flux can be reversed (in this case from outward to inward) with respect to the poloidally symmetric case.

This example provides a clear demonstration of the fact that the development of poloidal asymmetries (in particular small levels of in–out asymmetries) can strongly affect the size of the neoclassical convection or even reverse its direction, depending on the local value of η_i and on the level of impurity asymmetry. The dependence of

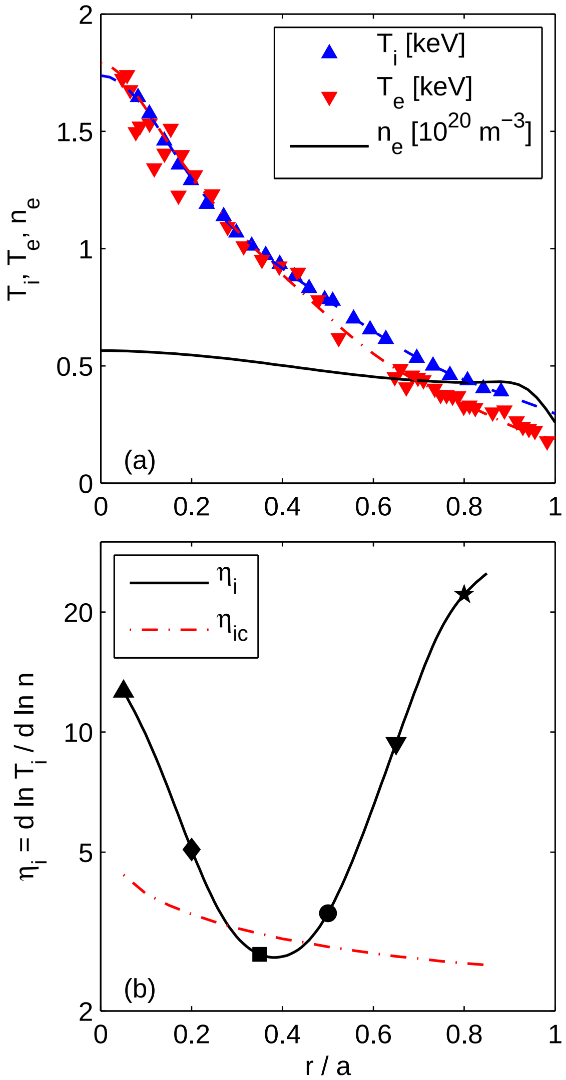

Figure 2: Radial profiles of ion and electron temperatures and density (a) and of η_i and η_{ic} (b) from a typical ASDEX Upgrade H–mode plasma (shot ♯27028 at 2.5 s, $I_p = 1$ MA, $B_T = 2.5$ T, $\bar{n}_e = 8\ 10^{19}$ m^{-3}, $P_{NBI} = 5$ MW, $P_{ECH} = 0.7$ MW, $q_{95} = 4.0$), which has been considered for the computation of the curves in Fig. 1. Symbols identify the radial locations at which the curves of Fig. 1 have been computed.

the impurity radial flux as a function of the poloidal angle θ is shown in Fig. 3 for $\epsilon = 0.159$ and for 4 levels of poloidal asymmetry of the impurity density distribution $(n_{zout} - n_{zin})/(n_{zout} + n_{zin}) = [-0.89,\ -0.3,\ 0,\ 0.89]$. In the poloidally symmetric case (solid line), the flux is outward on the LFS and inward on the HFS, resulting in an almost complete cancellation, as is typical of the Pfirsch–Schlüter regime with densities independent of θ. The reason why the cancellation is imperfect is that B and $R_{zi\|}$ vary by order $O(\epsilon)$ over the flux surface. With the experimental profiles considered in this case, the flux–surface averaged radial flux of the poloidally symmetric case is outward. If the impurity density peaks on the LFS (dash–dotted line, typically produced by centrifugal effects), the outward flux on the LFS is strongly enhanced, resulting in a net outward flux that is much larger than the poloidally symmetric Pfirsch–Schlüter value. In contrast, if the impurities are localized on the HFS (dotted line), the sign of the radial flux as a function of the poloidal angle is reversed, and becomes inward on the LFS and outward on the HFS. Because a density asymmetry with this sign counteracts the "natural" direction of the fluxes (obtained with evenly distributed impurities), the effect is non-monotonic: a very strong level of in–out density asymmetry increases the total flux, but an intermediate asymmetry (as experimentally achievable in the presence of ICRH with LFS resonance, dashed line) can have the effect of making the LHS/HFS cancellation of fluxes *more* complete and thus *reduce* the net flux. This condition can produce a flux–surface averaged radial flux which is very close to zero, for appropriate combinations of parameters.

IV Conclusions

In tokamak plasmas, highly charged impurity ions tend to be unevenly distributed over flux surfaces as a result of centrifugal effects and/or effects caused by auxiliary heating systems like ICRH and NBI, which produce fast ions with anisotropic pressure. These effects create a poloidal electric field that causes a much stronger redistribution of highly charged impurities than that of other species. In the limit of very strong redistribution (of either sign), the neoclassical heavy impurity particle flux increases by a large factor, up to $1/(2\epsilon^2)$ in a large-aspect-ratio tokamak. While this result is already known for out–in asymmetries produced by centrifugal effects [24, 30, 31, 33],

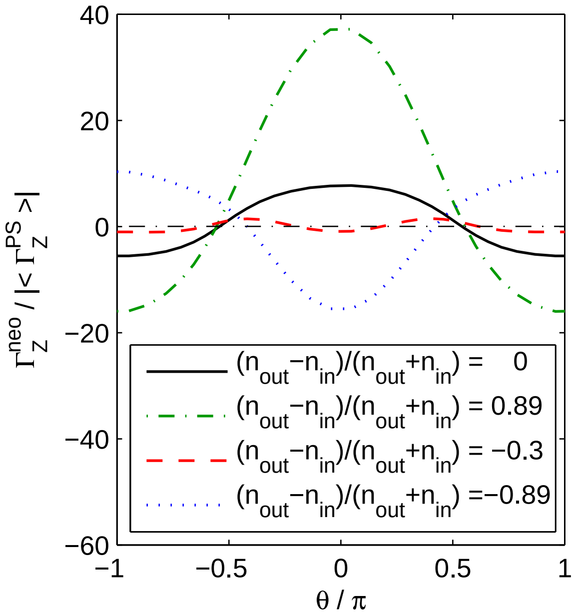

Figure 3: Heavy impurity radial flux as a function of the poloidal angle for different levels of poloidal asymmetry, as shown by the values of the parameter $(n_{zout} - n_{zin})/(n_{zout} + n_{zin})$ quoted in the legend.

this work shows that a similar level of strong enhancement can be also produced by asymmetries created by NBI or ICRH, which tend to have the opposite sign. The enhancement of the neoclassical impurity transport close to the magnetic axis has the unfavorable consequence that turbulent transport (which is often considered as the main transport effect to limit impurity accumulation [40]) can become relatively inefficient in competing against the neoclassical pinch. An additional result is that, in the case of moderate in–out asymmetries in the range reported by experiments with ICRH, the neoclassical transport of heavy impurities can be reduced, thus alleviating the danger of neoclassical accumulation, or can even reverse direction. The strength of the enhancement or reduction, and the direction reversal, of the impurity flux depend sensitively on the magnitude and location of the poloidal asymmetry as well as the flux-surface geometry and the background ion density and temperature profiles. In particular, when the $\eta_i = d \ln T_i / d \ln n$ parameter is above 2 but below a certain critical value, the development of a poloidal asymmetry can imply a reversal from outward to inward of the heavy impurity convection.

These analytical results motivate the development of complete models to accurately compute the strength of the impurity density asymmetries produced by auxiliary heating systems, and the extension of neoclassical transport codes, like NEO [34, 35], to account for their effects on the impurity transport. The result that poloidal asymmetries can produce large modifications of the impurity neoclassical transport should be taken into account not only in the transport modelling, but also in the analysis of the experimental data, when the experimentally inferred levels of transport are compared to the neoclassical levels in order to deduce the relative role of turbulent transport. The strong enhancement of neoclassical impurity transport close to the axis reported here shows that comparisons with neoclassical codes that neglect this effect can lead to incorrect conclusions about the (ir)relevance of neoclassical transport. The result that neoclassical impurity convection can be significantly reduced by appropriate levels of in–out impurity density asymmetry (in the range experimentally achievable by ICRH) may have an interesting potential for the development of external control tools against impurity accumulation.

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