Impurity transport in ITER-like plasmas

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Neoclassical impurity transport is compared with transport calculated from the reactive drift wave model of turbulent transport for an ITER-like [R. Aymar, P. Barabaschi, and Y. Shimomura, Plasma Phys. Controlled Fusion 44, 519 (2002)] scenario. The turbulent transport is inward for both main ions and impurities, but the impurity ion inward transport is much weaker than the main ion inward transport. The neoclassical impurity transport is outward because of temperature screening. The total impurity transport, determined by a balance between turbulent and neoclassical transport, depends sensitively on the charge number of the impurity and the ratio of the ion density and temperature scale lengths, $\eta_i$. © 2006 American Institute of Physics. [DOI: 10.1063/1.2375042]

I. INTRODUCTION

The presence of impurities in the edge region of tokamaks can be beneficial because a strongly radiating boundary distributes plasma power loss. But many tokamak discharges suffer from unwanted and uncontrollable impurity accumulation in the plasma core and this may lead to core radiation, flat radial distribution of temperature, fuel dilution, and sometimes even disruptions. Impurity transport is critical for ITER and the issue of what governs the impurity transport should be given careful consideration.

Neoclassical transport is driven by parallel friction dynamics and is not affected significantly by the fact that the ion cross-field transport is dominated by fluctuations. Thus it is clear that neoclassical and anomalous transport coexist. The assumption that the total impurity transport is a linear sum of turbulence-driven transport and neoclassical transport has been confirmed experimentally. Reference 1 describes low-Z impurity transport in DIII-D, and shows that both types of transport have to be included to explain the observed transport phenomena.

Turbulent transport of ions and impurities is usually much stronger than neoclassical theory predicts. However, it has been found in many tokamak experiments that in advanced scenarios impurity fluxes in the core may be of the order of the predictions of neoclassical theory. Reference 3 investigated the dependence of core plasma impurity transport on the charge number for ASDEX Upgrade H-mode discharges and found that the impurity transport is close to neoclassical predictions for low-Z impurities, when effects related to toroidal rotation can be neglected. In Ref. 5 the impurity behavior in advanced scenarios at ASDEX Upgrade and Joint European Torus (JET) was shown to be consistent with the predictions of neoclassical theory and the accumulation of impurities was connected to the main ion density peaking. In Ref. 7 the transport of impurities was measured in DIII-D in VH-mode discharges having reduced anomalous transport, and it was shown that temperature screening leads to strong outward convection of low-Z impurities. Thus, there is a strong motivation to compare neoclassical and turbulent impurity fluxes to determine which mechanism is stronger in ITER-like experimental situations.

In this work we compute the impurity transport from neoclassical theory and compare it with transport calculated from the reactive drift wave model of turbulent transport for a specific ITER-like scenario. As we will show, the total (neoclassical + turbulent) impurity flux depends sensitively on the charge number of the impurity and the ratio of the ion density and temperature scale lengths, $\eta_i$.

The structure of the paper is as follows: In Sec. II the neoclassical impurity flux is presented, both for collisionless and collisional impurities. Section III describes the reactive drift wave model of turbulent transport for an ITER-like plasma, and the dependence of the transport on $Z$ and $\eta_i$ is discussed. Finally, the results are summarized in Sec. V.

II. NEOCLASSICAL TRANSPORT

Neoclassical ion transport is dominated by collisions with impurities since ion-impurity collisions are more frequent than ion-electron collisions. From the ambipolarity condition it follows that the impurity flux is oppositely directed to the ion flux $I_i = -Z I_e$. In the absence of the thermal force this should lead to impurity accumulation ($n_i \sim n_e^2$, where $n_i$ and $n_e$ are the impurity and background ion densities). The accumulation can be reduced or avoided if the coefficient in front of the thermal force is negative. The outward flow induced by the thermal force, referred to as temperature screening, normally requires that the ions are collisionless and its effectiveness depends on the impurity charge and the fraction of impurities in the plasma.

The collision frequency between the species $a$ and $b$ is defined as

$$v_{ab} = \frac{n_b e_a^2 e_b^2}{m_b} \ln \Lambda / \left(4 \pi e_a^2 m_b^2 v_{T_a} v_{T_b}^2 \right),$$

where $v_{T_a}$ denotes the larger of the thermal velocities $v_{T_a}$ and $v_{T_b}$, and $\ln \Lambda$ is the Coulomb logarithm. The normalized collisionality is

$$\bar{\kappa} = \frac{v_{ab} R}{v_{T_a} e^{3/2}},$$

where $v_{ab} = \Sigma_{a} v_{ab}$ is the collision frequency, $q$ is the safety factor, $R$ is the major radius, and $e = r/R$ is the inverse aspect ratio. Impurity-impurity collisions are more frequent than ion-ion collisions, so that heavy impurities are usually more collisional than the ions: if $T_e > T_i$, then $n_i \sim n_e^2$ due to ambipolar effects,

$$n_i \sim n_e^2 \left(\frac{T_e}{T_i}\right)^{3/2} \sim n_e^2 \left(\frac{T_e}{T_i}\right)^{3/2}.$$
\[ \frac{1}{2} z^2 (1 + \alpha \sqrt{Z}) \], with \( \alpha = n_i Z^2 / n_e \). In today’s tokamak experiments, impurities tend to lie in the collisional regime \((\nu_e/\alpha \gg 1)\). However, for ITER-like parameters (density \( n_i = 10^{20} \text{ m}^{-3} \), temperature \( T_i = 10 \text{ keV} \), \( q = 2 \), \( R = 6 \text{ m} \), \( \alpha = 0.7 \), and \( \epsilon = 0.3 \)), low-Z impurities are collisionless \((\nu_e < 1)\) up to charge number \( Z = 11 \).

In this work we calculate the impurity transport in a plasma containing background ions \( (i) \) and electrons \( (e) \) and one species of impurity \( (z) \). The ion particle flux \( \Gamma_i = n_i V_i \) is given by

\[
\langle \Gamma_i, \nabla \psi \rangle = \sum_{a=0,\ldots,4} L_{ia} A_{ia}^a, \tag{1}
\]

where \( A_{ia}^a = \partial \ln p_a / \partial \psi \) and \( A_{a}^a = \partial \ln T_a / \partial \psi \) are the thermodynamic forces; \( n_a, p_a, \) and \( T_a \) are density, pressure, and temperature; \( \psi = -R A_\psi \) is the poloidal flux function, \( B = I \nabla \phi + \nabla \varphi \times \nabla \psi \) is the magnetic field;

\[
\langle \Gamma_z, \nabla \psi \rangle = f_t \frac{\int f_e \langle \bar{\psi}_{z1}(\ln p_i) + \bar{\psi}_{z2}(\ln T_i) \rangle - T_i \bar{\psi}_{z1}(\ln p_i) + \bar{\psi}_{z2}(\ln T_i) \rangle}{m_z \Omega_{\|}(\mu_{11} + \xi_z \mu_{1z})}, \tag{4}
\]

where \( f_t \) denotes derivative with respect to \( \psi \), \( \bar{\psi}_{z1} = 3 \langle (\nabla \psi)^2 \rangle \tau_{a1} f_e f_i m_i \mu(B_0)^{-1} \bar{\psi}_{a1} \), \( \tau_{a1} = 3 \nabla / 4 \bar{\psi}_{a1} \), \( \xi_z = \tau_{a1} m_i (\mu_{11} m_i)^{-1} \), and \( \Omega_{\|} = eB_0 / m_i \). The normalized viscosity coefficients are \( \bar{\psi}_{z1} = (\nabla \psi)^2 \) and \( \bar{\psi}_{z2} = (\nabla \psi)^2 / \alpha \nabla Z \) and \( m_{1z} \), with

\[
m_{1z} = \sqrt{1 + x_{1z}^2 + x_{2z}^2 \ln \frac{x_{2z}}{1 + x_{1z}^2}}, \tag{5}
\]

\[
m_{2z} = \sqrt{1 + \frac{x_{1z}^2}{x_{2z}^2} - \frac{5}{2} m_{1z}^2},
\]

where \( x_{1z} = v_{pi} / T_{Ti} \). The term proportional to the ion temperature gradient in Eq. (4) leads to outward transport of impurities (temperature screening). The inward flux driven by the impurity temperature gradient can be neglected, since

\[
Z \frac{1 + \bar{\psi}_{z2} / \mu_{1z} T_i / T_z}{1 + \bar{\psi}_{z2} / \mu_{1z} T_z} = 2 Z T_i / T_z \gg 1
\]

is satisfied in the parameter region of interest. Furthermore, if \( |n'_i / n_i| >> (T_i / Z T_z) |n'_z / n_z| \), also the flux driven by the impurity density gradient can be neglected. A criterion for the outward flux of impurities can then be derived to be

\[
n_i'(\eta_i - \eta_i) < 0, \] where \( \eta_i = (n_i T_i / n_i T_z) = L_{ti} / L_{T_i} \) is the ratio between the ion density and temperature scale lengths, and \( \eta_i' = (n_i / n_i + \bar{\psi}_{z2}) \).

In the cylindrical, large aspect ratio limit the impurity flux can be written in simpler form as

\[
\langle \Gamma_z, \nabla r \rangle = \int f_t \frac{\int f_e \langle \bar{\psi}_{z1}(\ln p_i) + \bar{\psi}_{z2}(\ln T_i) \rangle - T_i \bar{\psi}_{z1}(\ln p_i) + \bar{\psi}_{z2}(\ln T_i) \rangle}{m_z \Omega_{\|}(\mu_{11} + \xi_z \mu_{1z})} \times \left[ \frac{d \ln n_i}{dr} + (\mu_{11} + \mu_{12}) \frac{d \ln T_i}{dr} \right], \tag{7}
\]

where we neglected the flux driven by impurity density and temperature gradients and approximated the derivative with respect to \( \psi \) as \( \partial / \partial \psi = 1 / B_{\psi} (\partial / \partial r) \), where \( B_{\psi} \) is the poloidal magnetic field. In this limit \( B_{\psi} / B = e / l q \) and the effective fraction of trapped particles can be approximated as \( f_t = \sqrt{2 e / (1 + e)} \). If the impurity-ion density ratio \( n_i / n_i \) is constant, then \( \xi_z = \alpha \sqrt{m_i m_j T_i / T_z} \). The strong dependence of \( Z \) is strongly dependent on \( Z \) and the flux increases with \( Z \). This behavior has been confirmed experimentally. However, if \( Z_{eff} = (n_Z^2 + n_i) / (n_i Z + n_i) \) is assumed to be constant, then the fraction of impurities \( n_i / n_i \) decreases with increasing \( Z \), and \( \xi_z = \sqrt{Z} \). As a consequence, the neoclassical transport decreases weakly with increasing \( Z \).
B. Mixed-collisionality regime

Even in ITER, impurities with high $Z$ may be collisional, especially at the edge, where the temperature is low. The impurity flux is given by

$$\langle \Gamma_c \cdot \nabla \psi \rangle = \frac{2\eta_z^2}{m_Z k Yak Z \tau_i} \left[ \frac{\partial p_i}{\partial \psi} - \frac{3n_i n_D}{2} \frac{\partial T_i}{\partial \psi} - \frac{n_i \partial p_i}{n_Z Z \partial \psi} \right].$$

III. ANOMALOUS TRANSPORT IN THE REACTIVE DRIFT WAVE MODEL

The reactive model used here\textsuperscript{10} (usually called Weiland model) has been used extensively in describing the present international tokamak database (JET, ASDEX, DIII-D) and also for making ITER predictions.\textsuperscript{11} It uses an “advanced" reactive fluid model where “advanced" here refers to the rule for closure which allows us to use the model close to the fluid resonance in the collisionless case. In the version used here it has two independent ion species with the same physics included for both. The particle transport for the main species was tested successfully on JET discharges in Ref. 12. The particle pinch depends particularly strongly on the magnetic drift frequency. Because of this, species with higher $Z$ have a weaker particle pinch thus giving a favorable net effect on the effective $Z$. The particle pinch is particularly relevant for ITER since central fueling will not be possible there. In fact, the particle pinch has been found to improve ITER performance significantly.

The transport coefficients used in the calculations are derived using quasilinear theory.\textsuperscript{13-15} The saturation level can be estimated by a balance between the linear growth and the dominant $E \times B$ convective nonlinearity as

$$\frac{e\psi}{T_e} = \frac{1}{k_s \rho_s k_y c_s} \gamma,$$

where $x$ and $y$ are slab coordinates corresponding to radial and poloidal coordinates, $k$ is the wavenumber of the unstable wave, $\phi$ is the electrostatic potential, $\rho_s = c_s / \Omega_i$, $c_s$ is the sound speed, $\gamma$ is the linear growth rate of the unstable mode and the components of the wavenumber are assumed to be $k_x \rho_s = k_y \rho_s = \sqrt{0.1}$. The impurity transport is given by

$$\Gamma_c = \delta n_c v_e + \text{complex conjugate},$$

where $v_e = -ik_y \phi / B$ and $\delta n_c$ is the impurity density perturbation due to the turbulence given by\textsuperscript{10}

$$\frac{\partial n_c}{n_c} = \frac{L_{ne} \omega_{ne}}{L_{nc} N_c} \left[ \omega(1 - \epsilon_{nc}) + \left( \frac{\eta_z}{3} + \frac{5}{3} \epsilon_{nc} \right) \omega_{Dr} \right] \frac{e\phi}{T_e},$$

where $N_c = \omega^2 - 10 \omega_{Dr}/3 + 5 \omega_{Dr}/3$, $\omega_{Dr}$ is the diamagnetic frequency, $\omega_{ne} = \omega_{Dr}/\omega_{ci}$, $L_{ne} = -n_{io} / n_{io}$, and $L_{nc} = -T_{pe} / T_{a}$ are the density and temperature scale lengths for particle species $a$ and $\eta_z = L_{nc} / L_{Te}$. The imaginary part of the density perturbation is given by

$$\Im \left[ \frac{\partial n_c}{n_c} \right] = \frac{L_{ne} \epsilon_{nc}}{L_{nc} \epsilon_{nc} N_c} \left[ \left( \frac{\omega_{Dr}}{T_c} - \omega_{ci} \right) (1 - \epsilon_{nc}) \right.$$

$$\left. + \left( \frac{\eta_z}{3} + \frac{5}{3} \epsilon_{nc} \right) \tau_i \omega_{ci} \right] \frac{e\phi}{T_e},$$

where the overbar denotes normalization with respect to the electron magnetic drift frequency $\omega_{Dr} = 2k_y T_e / eBR$, $\tau_i = T_i / (Z T_e)$, and $N_c$ and $N_{ci}$ are the real and imaginary parts of $N_c$. Using Eqs. (9) and (12) in (10), we can calculate the impurity diffusion coefficient $D_c$, defined by $\Gamma_c = -D_c \delta n_c / dr$ to be

$$D_c = 2 \gamma \frac{\rho_c}{R k_y} \left( \Delta_1 - \Delta_2 \right) \frac{\epsilon_{nc}}{2 \Delta_2 L_{nc} / R},$$

where $\Delta_1 = 2 \tau_i (\omega_i + 5 \tau_i / 3) / |N_{ci}|^2$, with

$$\Delta_1 = \left[ \frac{\alpha_i^2}{\gamma^2} + \frac{\gamma^2}{\gamma^2} + 14 \omega_i \tau_i / 3 + 55 \tau_i / 9 \right] |N_{ci}|^2,$$

$$\Delta_2 = - \left[ \frac{\alpha_i^2}{\gamma^2} + \frac{\gamma^2}{\gamma^2} + 10 \omega_i \tau_i / 3 + 35 \tau_i / 9 \right] |N_{ci}|^2,$$

where $|N_{ci}| = N_{ci}^2 + N_{ci}^2$, with $N_{ci} = \omega_i - \gamma^2 + 10 \tau_i \omega_i / 3 + 5 \tau_i / 3$ and $N_{ci} = 2 \gamma (\omega_i + 5 \tau_i / 3)$. These expressions are in agreement with the corresponding ones given in Refs. 15 and 16.

The diffusion coefficient derived above is for the electrostatic case, but electromagnetic effects only influence the unstable wave frequency and have no effect on the expres-
For a specific high-$Q$ ITER-like scenario, with minor radius $a=2.75$ m, major radius $R=6.2$ m, and magnetic field at the axis $B=5.3$ T, the radial temperature and density profiles are given in Fig. 1. For the ITER scenario we selected, the density profile is slightly hollow and is much flatter than the temperature profile. For flat density profiles, the neoclassical impurity transport is outward and the neoclassical transport is outward. The reason for the outward flow is temperature screening, that is effective when $\eta_i > \eta_i^0 = 3.5$, if both the ion density and temperature gradients are negative. Note that even if the density gradient is locally positive $n_i'(\eta_i - \eta_i^0) < 0$ is satisfied, because $\eta_i$ is then negative (assuming that the temperature gradient is still negative). However, it is interesting to note that a more peaked density profile reduces the effect of temperature screening and should reduce the outward transport. In Fig. 2 the influence of the peakedness and the hollowness on the impurity transport is illustrated by modeling the main ion density by \( n_i = n_i(0)(1 - r/a)^3 \) and for various $\sigma$ comparing the impurity transport with the transport in the case of the density profile shown in Fig. 1. The impurity outward transport is higher for the hollow profile shown in Fig. 1 than for the flat profile with $\sigma = 0.1$, especially in the outer part of the plasma. The transport changes sign if the parameter $\sigma > 0.9$, since a strong density gradient overwhelms the temperature screening. Figure 3 shows $\eta_i$ as a function of the normalized radial position for the impurity diffusion coefficient. The frequency and growth rate of the unstable mode, $\tilde{\omega} = \omega_i + i \tilde{\gamma}$, are determined numerically, including electromagnetic effects.

It is well known that the growth rate of the main ion temperature gradient (ITG) mode is reduced by dilution.\(^{13}\) This would tend to reduce the turbulent particle transport. However, for sufficiently large impurity density the impurity ITG takes over. The growth rate of this mode increases with impurity density and, accordingly, also the turbulent particle transport increases. Thus, although the impurity density does not appear explicitly in our formulas, it is an important parameter. In the following section, the results of Eqs. (7), (8), and (13) will be compared with numerical calculations of the turbulent transport for ITER-like parameters and different impurity species. In these calculations we keep $Z_{\text{eff}} \approx 1.7$ constant.

### IV. IMPURITY FLUX IN ITER

For a specific high-$Q$ ITER-like scenario, with minor radius $a=2.75$ m, major radius $R=6.2$ m, and magnetic field at the axis $B=5.3$ T, the radial temperature and density profiles are given in Fig. 1. For the ITER scenario we selected, the density profile is slightly hollow and is much flatter than the temperature profile. For flat density profiles, the neoclassical impurity transport is outward and the neoclassical transport is outward. The reason for the outward flow is temperature screening, that is effective when $\eta_i > \eta_i^0 = 3.5$, if both the ion density and temperature gradients are negative. Note that even if the density gradient is locally positive $n_i'(\eta_i - \eta_i^0) < 0$ is satisfied, because $\eta_i$ is then negative (assuming that the temperature gradient is still negative). However, it is interesting to note that a more peaked density profile reduces the effect of temperature screening and should reduce the outward transport. In Fig. 2 the influence of the peakedness and the hollowness on the impurity transport is illustrated by modeling the main ion density by \( n_i = n_i(0)(1 - r/a)^3 \) and for various $\sigma$ comparing the impurity transport with the transport in the case of the density profile shown in Fig. 1. The impurity outward transport is higher for the hollow profile shown in Fig. 1 than for the flat profile with $\sigma = 0.1$, especially in the outer part of the plasma. The transport changes sign if the parameter $\sigma > 0.9$, since a strong density gradient overwhelms the temperature screening. Figure 3 shows $\eta_i$ as a function of the normalized radial position for the impurity diffusion coefficient. The frequency and growth rate of the unstable mode, $\tilde{\omega} = \omega_i + i \tilde{\gamma}$, are determined numerically, including electromagnetic effects.

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radius for the various density profiles used in Fig. 2, together with the threshold for \( \eta_i \), denoted by \( \eta_i \).

The transport due to the impurity gradients is more than an order of magnitude lower than the transport due to the ion gradients for this specific ITER scenario.

**Comparison of turbulent and neoclassical transport**

The turbulent transport is caused by the ITG mode and trapped electron (TEM) mode. The real part of the frequency and the growth rate for the numerically calculated unstable modes are given in Fig. 4. The growth rate of the ITG mode is considerably larger than the growth rate of the TEM modes, and its real frequency is negative while the real frequencies of the TEM modes are positive. The impurity transport is inward when \( \Delta_i - \Delta_f \eta_i + 2\Delta_f L_{ni}/R < 0 \), which is satisfied by impurities with any \( Z \) in this specific ITER-like scenario. Figure 5 shows the neoclassical and the numerically calculated turbulent impurity particle transport as a function of the normalized radius for impurity charges \( Z = 6, 10, 54 \). It is interesting to note that neoclassical and turbulent transport have opposite signs.

The numerical calculations show that turbulent transport is dominated by the unstable ITG mode. The convective part of the transport, proportional to \( \Delta_i \), dominates, and that gives an inward flux of impurities. The terms proportional to \( \Delta_f \) give rise to outward transport (for the ITG mode), but they are negligible inside the radius \( r/a < 0.7 \). Noting that for \( r/a < 0.7 \), \( \Delta_i = -1/(\bar{\omega}_r^2 + \gamma^2) \), we can rewrite the convective part of the flux as

\[
\Gamma_{i\text{turb}} = -\frac{2\gamma^2}{\bar{\omega}_r^2 + \gamma^2} \frac{T_{i\text{keV}}}{B} \n_i, \tag{14}
\]

where \( T_{i\text{keV}} \) is the ion temperature in keV. The flux given in Eq. (14) is approximately equal to the total turbulent flux inside the radius \( r/a < 0.7 \). Outside this radius, the terms proportional to \( \Delta_i \) and \( \Delta_f \) become comparable to the convective part, and reduce the inward transport, so much that it sometimes even changes sign at the edge.

Assuming flat density profiles, so that the flux driven by the ion density gradient can be neglected, the neoclassical flux in the banana regime can be approximated as

\[
\Gamma_{i\text{neo}} \approx -2 \cdot 10^{19} \frac{f_{\text{i}19} q^2 n_i^2 T_{i\text{keV}}}{f_{\text{e}19} Z B \epsilon^2 T_{\text{keV}}^{3/2} R} \tag{15}
\]

for \( r/a < 0.7 \), where \( n_{i19} \) is the ion density in units of \( 10^{19} \, \text{m}^{-3} \). Using these approximate expressions for the neoclassical and turbulent transport we can determine the charge number for which the outward neoclassical transport is larger than the inward turbulent transport,

\[
Z > Z_{\text{eff}} + 10^2 \frac{Z_{\text{eff}} - 1}{\epsilon^2 f_{\text{i}19} q^2 \bar{\omega}_r^2 + \gamma^2 T_{\text{keV}}^2 R}. \tag{16}
\]

For low-Z impurities, the inward turbulent transport dominates. However, the main ion inward transport is at least an order of magnitude larger.

The growth rate of the unstable ITG mode increases with \( \eta_i \), so that the turbulent transport (which is normally inward) increases. Even if \( \eta_i \) is large, leading to strong outward neoclassical transport, the resulting (turbulent + neoclassical) transport can be inward, if the turbulent transport is large enough to dominate over neoclassical transport. Figure 6 shows the total (turbulent + neoclassical) impurity transport as a function of \( \eta_i \), for impurity charges \( Z = 6, 10, 54 \). For low \( Z \), the impurity transport depends sensitively on \( \eta_i \). The direction of the impurity flux is inward if \( \eta_i \) is large and negative (flat and slightly hollow profiles), and the threshold is at \( \eta_i \approx -10 \). Below \( \eta_i \approx -10 \) turbulent transport dominates and the total impurity transport is inward. For large and positive \( \eta_i \) (outside \( r/a = 0.7 \)), inward transport caused by the convective term \( \Delta_i \) is counteracted by the outward transport.
caused by $\Delta_i$ and $\Delta_T$, and the neoclassical transport, so that the total transport is outward if $\eta_i$ is not too large.

Our expression for the impurity diffusion coefficient (13) agrees with Eqs. (5)–(8) of Ref. 16, and it shows that the magnitude (and sometimes even the direction) of the anomalous flux depends on the sign of the real frequency of the unstable mode, which is different for the ITG and TEM modes. The TEM mode, with a positive real frequency, gives rise to an inward flux, while the ITG mode can give rise to both inward and outward flux depending on the relative magnitude of the terms $\Delta_i$, $\Delta_T$, and $\Delta_2$. In our calculations for the chosen ITER profile, the ITG mode is dominant, and the anomalous flux is inward, except for low-Z impurities in the edge plasma.

Figure 7 shows the turbulent and neoclassical ion particle transport for $Z=6$, 10, 54. Turbulent ion particle transport is not sensitive to the charge number. Both turbulent and neoclassical ion transport are inward. Although the magnitude of the neoclassical transport is $Z$ times larger than the impurity transport, it is still smaller in comparison with turbulent ion transport by at least a factor of 20.

V. CONCLUSIONS

The direction of the total impurity flux depends on the impurity charge number $Z$ and the ratio of ion density and temperature scale lengths $\eta_i$. If $Z$ is high (collisional impurities), then impurity transport is inward for all $\eta_i$. If $Z$ is low (collisionless impurities), then the direction of the total impurity flux depends sensitively on $\eta_i$. If $|\eta_i| \leq 10$ the total flux is outward, but if $|\eta_i| \geq 10$, then turbulent transport dominates over neoclassical and the total transport is inward.

In this work we have only determined the direction and magnitude of the impurity flux for a given ITER scenario, at a specific time, and we did not analyze the evolution of the impurity profiles in time. If the density profile becomes more peaked as the turbulence drives the main ions inward, $\eta_i$ changes, and then both the direction and magnitude of the impurity flux will change.

Our analysis assumes that the effect of the interaction between different impurity species and charge states can be neglected. This assumption is justified in case the impurity density is low, so that collisions among different species may be neglected. Furthermore, we neglect the effect of plasma rotation, since it is not expected to be very strong in ITER. However, we note that toroidal plasma rotation may enhance the neoclassical impurity transport, especially for heavy impurities.

Our conclusions are in qualitative agreement with the conclusions of Ref. 19. In Ref. 19 impurity behavior was numerically calculated for ITER operational scenarios and it was shown that for flat density profiles, temperature screening prohibits impurity accumulation in the core. However, our analysis shows that, depending on the sign and magnitude of $\eta_i$, even for low $Z$, turbulent transport may dominate over neoclassical, and the direction of the impurity particle flux may be inward. However, the turbulent transport of main ions is usually also inward and dominates the effect on the charge balance.

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