A possible mechanism responsible for generating impurity outward flow under radio frequency heating

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Abstract

The effect of poloidal asymmetry of impurities on impurity transport driven by electrostatic turbulence in tokamak plasmas is analyzed. It is found that in the presence of in–out asymmetric impurity populations the zero-flux impurity density gradient (the so-called peaking factor) is significantly reduced. A sign change in the impurity flux may occur if the asymmetry is sufficiently large. This may be a contributing reason for the observed outward convection of impurities in the presence of radio frequency heating. This paper extends a previous work (Fülöp and Moradi 2011 Phys. Plasmas 18 030703), by including the effect of ion parallel compressibility on the peaking factor, which is found to have a significant contribution in the presence of poloidal asymmetry. It is shown here that in the ion temperature gradient mode dominated plasmas the presence of an in–out poloidal asymmetry can lead to a negative impurity peaking factor, and it becomes more negative in regions with larger ion temperature gradients. In the trapped electron mode dominated plasmas an in–out poloidal asymmetry results in a strong reduction of the peaking factor; however, it remains positive for typical experimental parameters. Furthermore, it is shown that an up–down asymmetry reduces the peaking factor while an out–in asymmetry increases it.

(Some figures in this article are in colour only in the electronic version)

1. Introduction

Accumulation of impurities in the core of fusion devices would have a detrimental effect on fusion reactivity due to increased radiation losses and plasma dilution. Significant effort has therefore been devoted in recent years to identify plasma conditions in which accumulation can be avoided. One of the most promising methods for obtaining flat or hollow impurity density profiles is to use radio frequency (RF) heating. This has been shown to work well in various experiments [1–7] but the physical mechanism by which the change in the direction of the impurity convective velocity occurs has not yet been clearly identified in spite of the
various efforts that have been made \cite{8–10}. In particular, the reason for the flat impurity density profiles in ion cyclotron resonance heated (ICRH) discharges in JET has been debated for many years \cite{1,5,6}. In a recent paper \cite{11} it was shown that a poloidal asymmetry could lead to a significant reduction in the impurity zero-flux density gradient (also called the peaking factor), and even a sign change in the impurity flux, if the asymmetry is sufficiently large. Along with other effects, this may be a contributing factor to the avoidance of accumulation of high-Z impurities with ICRH. In this paper, we extend the analysis of \cite{11}, by including the effect of ion parallel compressibility on the peaking factor, and by presenting a numerical investigation of the effect of poloidal asymmetry on impurity transport. In particular, the dependence of the impurity density peaking factor on charge number, inverse ion and electron temperature scale lengths and inverse electron density scale length is analyzed. The results are benchmarked to GYRO \cite{12} in the poloidally symmetric case.

It is found that inboard accumulation gives rise to negative peaking factors (outward impurity convection and hollow impurity profile) in ion temperature gradient (ITG) mode-driven turbulence, if the asymmetry is sufficiently large. Also in the trapped electron (TE) mode-dominated case inboard accumulation results in a strong reduction in the peaking factor; however, it remains positive for typical experimental parameters. The sign and magnitude of the peaking factor will be shown to be sensitive not only to the asymmetry strength but also to the temperature gradient. As the ITG increases in the ITG-mode-dominated case the impurity peaking factor becomes more negative. As noted in previous work \cite{8}, in the poloidally symmetric case, parallel compressibility has a significant effect on the peaking factor. As we will see in this paper this effect is also observed for asymmetric impurity densities, and leads to a peaking factor that is more sensitive to the ITG than in the case without taking into account parallel compressibility.

The rest of the paper is organized as follows. In section 2 we describe the mechanism behind the impurity poloidal asymmetry that arises in the presence of ICRH. In section 3 the model for calculating the quasilinear impurity flux and the peaking factor in the presence of poloidal asymmetry is presented. In section 4 the parametric dependences of the peaking factor are analyzed by presenting scans over relevant parameters such as charge number, and temperature and density scale lengths. Also, the importance of the impurity parallel compressibility is demonstrated. Finally, the results are discussed and summarized in section 5.

2. Poloidal asymmetry

Poloidal impurity asymmetries in tokamaks can arise for various reasons: e.g. difference in impurity source location, toroidal rotation or neoclassical effects. There is a wealth of experimental evidence for poloidal asymmetries \cite{13–18}. In the plasma core in–out asymmetries can arise due to the presence of RF heating (henceforth ‘in–out’ and ‘out–in’ asymmetries will refer to the situations when the maximum of the poloidally varying impurity density is at the inboard and outboard sides of the plasma, respectively). A detailed physical explanation of why ICRH favors inboard accumulation, together with the description of the experimental results, is given in \cite{13}. The RF-heating scheme applied in the experiment described in \cite{13} is hydrogen-minority heating in a deuterium plasma. The underlying principle is that the asymmetry is a result of the increase in the hydrogen-minority density on the outboard side. These particles tend to be trapped on the outside of the torus and the turning points of their orbits drift toward the resonance layer due to the heating. The poloidal asymmetry in the hydrogen-minority density gives rise to an electric field that pushes the other ion species to the inboard side. In the case of highly charged impurities, this effect is amplified by their higher charge $Z$. 

2
The RF-induced accumulation of minority ions on the outboard side leads to a corresponding impurity accumulation on the inboard side by the following mechanism. If the plasma consists of electrons, bulk ions, impurity ions and RF-heated minority ions then it can be expected that all species except the minority ions are Boltzmann distributed (the dynamics of the minority ions is strongly affected by the heating). If a particle species $a$ is Boltzmann distributed, the poloidal variation of the density is

$$\tilde{n}_a/\hat{n}_{a0} \simeq -e_a\phi_E/T_a,$$

where the tilde denotes the variation on the flux surface, $e_a$ is the charge and $T_a$ is the temperature of the species. Quasineutrality requires that

$$n_{e0}(1 + e\phi_E/T_e) - n_{D0}(1 - e\phi_E/T_D) - n_{H0}(1 - e\phi_E/T_H) - Zn_{Z0}(1 - Ze\phi_E/T_Z) = \hat{n}_H,$$

where the subscript zero indicates the density where the equilibrium potential $\phi_E$ vanishes, and $\hat{n}_H$ represents only the fraction of the hydrogen-minority density which feels the ICRH resonance and does not follow the Boltzmann distribution. If $\phi_E$ is normalized so that $n_{D0} + n_{H0} + Zn_{Z0} - n_{e0} = 0$, assuming a similar temperature for the different ion species and $n_{H0} \ll n_{D0}$, the poloidal variation of the impurity density becomes

$$\frac{\delta n_z}{n_{Z0}} = -\frac{Ze\phi_E}{T_z} = -\frac{Z\hat{n}_H/n_{D0}}{1 + (T_i/T_e) + (n_{Z0}Z^2/n_{D0})}.$$

Since the poloidal variation in the impurity density has the opposite sign to that of the minority ions, the accumulation of the latter on the outboard side gives rise to an electric field that pushes the other ion species to the inboard side. Simulations of the hydrogen ion distribution function in the presence of RF heating with the Monte Carlo code FIDO described in [13] show that a considerable out–in asymmetry in the hydrogen ion density can be expected, which is sufficient to account for the observed in–out asymmetry in the impurity density.

In the tokamak edge, where the plasma is sufficiently collisional, steep radial pressure or temperature gradients can also give rise to an in–out asymmetry. These effects have been observed in, e.g., Alcator C-Mod [14], and it has been shown that the observations are in qualitative agreement with neoclassical theory [19–22]. The sign and magnitude of these asymmetries depend sensitively and nonlinearly on magnetic geometry, fraction of impurities in the plasma and rotation. Neoclassical theory also predicts an up–down asymmetry, which is caused by the ion–impurity friction.

3. Impurity flux

Since impurity transport is usually dominated by drift-wave turbulence, in this work we focus on the effect of the impurity poloidal asymmetry on impurity transport driven by microinstabilities. We assume that the processes that cause the asymmetry are not affected significantly by the fact that the cross-field transport is dominated by fluctuations. For simplicity we consider only the collisionless, electrostatic case and due to the low transit frequency of the impurities we neglect electrostatic trapping. The quasilinear impurity particle flux is given by

$$\Gamma_z = -\frac{k_0}{B} \text{Im} [\hat{n}_z, \phi^*],$$

where $\text{Im} [\cdot]$ denotes the imaginary part, $k_0$ is the poloidal wavenumber, $\hat{n}_z$ is the perturbed impurity density, $\phi^*$ is the complex conjugate of the perturbed electrostatic potential $\phi$.

The perturbed impurity density response in an axisymmetric, large aspect ratio torus with circular magnetic surfaces can be obtained from the linearized gyrokinetic equation [23],

$$\frac{\partial g_a}{qR \partial \theta} - i(\omega - \omega_{Da})g_a - C_a(g_a) = -i\frac{e_a f_{a0}}{T_a} (\omega - \omega_{m0}^a) \phi J_0(z_a),$$

where $J_0$ is the Bessel function of the first kind of order zero.
where $g_{\alpha}$ is the non-adiabatic part of the perturbed distribution function, $\theta$ is the extended poloidal angle, $f_{0\alpha} = n_{\alpha}/(\sqrt{\pi} v_{T\alpha})^{3}$ is the equilibrium Maxwellian distribution function, $x_{\alpha} = v_{\alpha}/v_{T\alpha}$ is the velocity normalized to the thermal speed $v_{T\alpha} = (2T_{\alpha}/m_{\alpha})^{1/2}$, $n_{\alpha}$ and $m_{\alpha}$ are the density and mass of species $\alpha$, $\omega_{\alpha} = -k_{B}T_{\alpha}/e_{\alpha}B L_{\alpha}$ is the diamagnetic frequency, $\omega_{\alpha}^{*} = \omega_{\alpha}[1 + (x_{\alpha}^{2} - \frac{3}{2})L_{\alpha}/L_{T\alpha}]$, $L_{\alpha} = -[\partial(\ln n_{\alpha})/\partial r]^{-1}$, are the density and temperature scale lengths, $\omega_{Da} = -k_{B}(v_{T\alpha}^{2}/2 + v_{\alpha}^{2})D(\theta)/\omega_{\alpha}$ is the magnetic drift frequency, $D(\theta) = (\cos \theta + s \theta \sin \theta)$, $\omega_{s\alpha} = e_{\alpha}B/m_{\alpha}$ is the cyclotron frequency, $B$ is the equilibrium magnetic field, $q$ is the safety factor, $s = (r/q)(dq/dr)$ is the magnetic shear, $r$ and $R$ are the minor and major radii, $J_{0}$ is the Bessel function of the first kind and $z_{a} = k_{\perp}v_{\perp}/\omega_{\alpha}$. Here, we assumed that the $E \times B$-drift frequency is negligible compared to the wave frequency.

In the absence of collisions and in the limit $v_{||}/qR(\omega - \omega_{Dz}) \ll 1$, equation (3) can be solved iteratively to find the non-adiabatic part of the perturbed impurity distribution. Including the Boltzmann part of the distribution, the perturbed ion density response becomes $[23, 24]
\begin{align}
\frac{\tilde{n}_{z}}{n_{z}} &= -\frac{Ze\phi}{T_{z}} \int d^{3}v \frac{ZeJ_{0}(zz_\perp)}{n_{z}T_{z}} \left[ 1 - \frac{v_{||}}{qR(\omega - \omega_{Dz})} \frac{\partial}{\partial \theta} \frac{v_{||}}{qR(\omega - \omega_{Dz})} \frac{\partial}{\partial \theta} \right] \\
&\quad \times \frac{\omega - \omega_{\alpha}^{*}}{\omega - \omega_{Dz}} f_{0J_{0}(zz_\perp)} \phi, \tag{4}
\end{align}

where parallel compressibility is represented by the term containing the two $\theta$-derivatives.

The zero-flux impurity density gradient (the peaking factor) can be obtained from (4). Here $\{ \cdot \cdots \cdot \} = (1/2\pi) \int_{-\pi}^{\pi} (\cdot \cdots \cdot \cdot) d\theta$. We model the poloidal asymmetry of the impurity density by the ansatz $n_{\perp} = n_{\perp 0}P(\theta) = n_{\perp 0} \sum n_{\alpha} f_{\alpha} P(\theta, \delta, n)$, where $P(\theta, \delta, n) = [\cos^{2}(\frac{\delta}{2})]^{n}$, with $\delta$ representing the angular position where the impurity density has its maximum, $n$ represents the peakedness of the asymmetry and the weights $f_{\alpha}$ can be chosen to represent populations of impurities with various degrees of peakedness. For evenly distributed impurity density $f_{\alpha} = 1$ and the rest of the weights $f_{\alpha} = 0$ for $n \neq 0$. Figure 1 shows the asymmetry function $P(\theta, \delta, n)$ as a function of normalized $\theta/\pi$, for different values of weights $f_{\alpha}$. We note that, in general, the impurity density cannot be factorized this way, which would mean a possible poloidal variation of the logarithmic density gradients. Nevertheless, to keep the theoretical treatment simple, we adopt the factorization, which is justified for cases where $d(\ln n_{z}/dr)/d\theta \ll d(\ln n_{z}/dr)$. A similar asymmetry function was also used in previous works [11, 15].

The impurity flux can be calculated numerically by solving the velocity-space integrals in the perturbed impurity density (4) in the expression for the impurity flux (2), without the constant energy resonance approximation $[v_{T\perp}^{2} + 2v_{\perp}^{2} \rightarrow 4(v_{T\perp}^{2} + v_{\perp}^{2})/3][23]$ or the assumption on the smallness of the finite Larmor radius parameter. These approximations were used in [11].

The impurity peaking factor $a/L_{nz}^{0}$ can be obtained by setting the impurity flux zero, or $\{ \text{Im}[\tilde{n}_{z}\phi^{*}] \} = 0$. Here, the fluctuating density is given by equation (4), $a$ is the outermost minor radius, and $a/L_{nz}^{0}$ is calculated as

\begin{align}
a/L_{nz}^{0} &= \frac{\{\text{Im}[S_{F}(\theta)\phi^{*}]\}}{\{\text{Im}[S_{n}(\theta)\phi^{*}]\}}, \tag{5}
\end{align}

where

\begin{align}
S_{n}(\theta) = \int d^{3}x_{\perp} e^{-i\vec{x}_{\perp} \cdot \vec{r}} J_{0}(zz_{\parallel}) \left[ 1 - \delta_{p} \mathcal{M}_{0} \right] \frac{\tilde{\omega}_{nz} P(\theta) \phi(\theta) J_{0}(zz_{\parallel})}{\tilde{\omega} - \tilde{\omega}_{Dz}(\theta)}, \tag{6}
\end{align}

and

\begin{align}
S_{F}(\theta) = \int d^{3}x_{\perp} e^{-i\vec{x}_{\perp} \cdot \vec{r}} J_{0}(zz_{\parallel}) \left[ 1 - \delta_{p} \mathcal{M}_{0} \right] \frac{[\tilde{\omega} - \tilde{\omega}_{nz}(x_{\perp}^{2} - 3/2)a/L_{Tz}] P(\theta) \phi(\theta) J_{0}(zz_{\parallel})}{\tilde{\omega} - \tilde{\omega}_{Dz}(\theta)}, \tag{7}
\end{align}
Figure 1. $P(\theta)$ normalized to its maximum, for different values of $f_n$. Solid line (red) represents $f_0 = f_1 = 1$, dashed line (blue) is for $f_0 = 1, f_1 = 0.5$, dashed–dotted line (green) is for $f_0 = 1, f_1 = 0.5, f_2 = 0$. The rest of the $f_n$s are assumed to be zero. $\delta = \pi$ in all cases.

with the operator representing parallel compressibility

$$M_\theta = \frac{x_{c\parallel}}{R(\theta)(\bar{\omega} - \bar{\omega}_Dz(\theta))} \frac{\partial}{\partial \theta} \frac{x_{c\parallel}}{R(\theta)(\bar{\omega} - \bar{\omega}_Dz(\theta))} \frac{\partial}{\partial \theta},$$

where $R(\theta) = R_0(1 + \epsilon \cos \theta)$ with $\epsilon = r/R_0$, $\delta_p = 2a^2m_i/(q^2m_z\tau_z)$, $\tau_z = T_z/T_e$, $b = (\rho_i/Z)\sqrt{2Ae/\tau_e\tau_i}$, $\rho_i$ is the ion sound Larmor radius, $\bar{\omega}_{sc} = -k_0\rho_i/Z\tau_z$, $\bar{\omega}_{Dz} = -2k_0\rho_i(a/R)(x^2_{c\parallel}/2 + x^2_{c\parallel})D(\theta)/Z\tau_z$, $z_{c\parallel}(\theta) = x_{c\perp}b k_{c\perp}(\theta)$ and $k_{c\perp} = k_0\sqrt{1 + s^2\theta^2}$. All frequencies marked with bars are in $c_s/a$ units, where $c_s$ is the ion sound speed, and the main ion and impurity temperature gradients are assumed to be equal, $a/L_T = a/L_{Tz}$.

4. Parametric dependences of the peaking factor

In the calculations presented in this section we have used the following local profile and magnetic geometry parameters: $r/a = 0.3$, $R/a = 3$, $k_0\rho_i = 0.3$, $q = 1.7$, $a/L_{n_e} = 1.5$, $T_i/T_e = 0.85$, $a/L_{Tz} = 2$, $a/L_{Te} = 2.5$, $s = 0.22$ and $\rho_i/a = 0.0035$. This is the baseline case in our study, and these parameters will be used unless otherwise stated. The impurities are assumed to be present in trace quantities, in the sense that $Z_{nz}/n_e \ll 1$ ($n_z/n_e = 2 \times 10^{-3}$ is used in the simulations). To study the effect of the strength of the asymmetry we present results for $n = 0$ to $n = 3$, and assume $f_j = 1$ for the specific asymmetry strength (for instance, if $n = 3$ we use $f_3 = 1$ and $f_j = 0$ for $j \neq 3$).

The perturbed electrostatic potential and eigenvalues (which are practically unaffected by the presence of a poloidally asymmetric trace impurity species) are obtained by linear electrostatic gyrokinetic calculations with GYRO. We note that the electron and main ion densities are assumed to be approximately poloidally symmetric. This is important for our model to be valid, since GYRO assumes poloidally symmetric background plasma parameters. We would like to emphasize that—as can be seen from equations (5)–(7)—the zero flux impurity density gradient does not depend on any constant multiplier of the perturbed potential.
4.1. Charge number dependence

The peaking factor as a function of charge number for various impurity poloidal asymmetries is shown in figure 2. In the poloidally symmetric case (solid line in figure 2(a)), the peaking factor is not sensitive to the charge number, as has been noted before, in both fluid and gyrokinetic simulations of ITG turbulence dominated transport, without taking into account the poloidal impurity asymmetries [8, 25]. The situation is similar also in the case of up–down asymmetric impurity populations or outboard accumulation. Out–in asymmetry gives slightly higher and up–down asymmetry gives lower peaking factors. However, in agreement with the conclusion of [11], impurities experience outward convection (corresponding to negative peaking factor) if the impurity density is accumulated on the inboard side, as shown with the black dotted line in figure 2(a). The change in the peaking factor becomes stronger as the asymmetry is increased, as illustrated in figure 2(b). Note that the strength of the asymmetry is also expected to depend on \( Z \), and usually it is larger for heavy impurities, as was shown in section 2. According to equation (1), in the limit of trace impurities, the poloidal variation of the impurity density is proportional to the charge number.

In order to examine the importance of parallel compressibility in determining the charge dependence of the impurity peaking factor we separate the terms independent of and proportional to \( \delta_p \) (representing the parallel compressibility) in equation (6) as

\[ S^n(\theta) = S^n_1(\theta) + S^n_{pc}(\theta), \]

where

\[ S^n_1(\theta) = \mathcal{P}(\theta)\phi(\theta) \int d^3x e^{-x^2} J_0^2[z_\perp(\theta)] \frac{\tilde{\omega}z}{\tilde{\omega} - \tilde{\omega}_{Dz}(\theta)}, \]

and in equation (7) as

\[ S^T(\theta) = S^T_1(\theta) + S^T_{pc}(\theta), \]

where

\[ S^T_1(\theta) = \mathcal{P}(\theta)\phi(\theta) \int d^3x e^{-x^2} J_0^2[z_\perp(\theta)] \frac{\tilde{\omega} - \tilde{\omega}_{Dz}(x^2 - 3/2)a/Lr}{\tilde{\omega} - \tilde{\omega}_{Dz}(\theta)}, \]
S_{pc}^n T(\theta) = -\delta_p \int dx z e^{-x^2} J_0[z_c(\theta)] M_\theta [\bar{\omega} - \bar{\omega}_{zc}(x_c^2 - 3/2)a/L_T] P(\theta)\phi(\theta) J_0[z_c(\theta)]. \ \ \ \ (12)

Using the above notation the peaking factor can be rewritten as

$$a/L_{nz}^0 = \frac{\langle \text{Im}[S_1^T(\theta)\phi^*] \rangle + \langle \text{Im}[S_{pc}^n T(\theta)\phi^*] \rangle}{\langle \text{Im}[S_{pc}^n T(\theta)\phi^*] \rangle + \langle \text{Im}[S_1^T(\theta)\phi^*] \rangle}.$$ \ \ \ \ (13)

Figure 3 shows these four expressions as functions of impurity charge for the symmetric and in–out asymmetric cases with an asymmetry strength of $n = 3$. In the denominator, the terms independent of $\delta_p$, i.e. $S_1^T(\theta)$, are the dominant contributors in both the symmetric and asymmetric cases. These are proportional to $1/Z$. But in the numerator, the balance between the terms proportional to and independent of $\delta_p$ is very different between the symmetric and asymmetric cases; in the symmetric case $S_1^T(\theta)$ is the dominant term, while in the asymmetric case $S_{pc}^n T(\theta)$ is dominant. These results show that if an in–out asymmetry is present, the parallel compressibility effects become more important than the other terms, and therefore, have to be taken into account.

In the following we will concentrate on the peaking factor for nickel which was one of the impurities studied in [1]. Figure 4 shows the peaking factor for nickel for various accumulation maxima $\delta$ and asymmetry strengths $n$. The peaking factor increases slightly for out–in accumulation but the more dramatic change—including a sign change—is expected only for inboard accumulation. The sign change occurs when $n \simeq 2.5$. In [11] it was shown that without taking into account the effect of parallel compressibility the sign change occurs when $n \simeq 3$. This is in agreement with the results obtained here if parallel compressibility effects are neglected, i.e. by setting $\delta_p = 0$ in equations (6) and (7) (as will be shown later in figure 12).
Figure 4. Peaking factor for nickel as a function of $n$ (a) and $\delta$ (b). (a) symmetric impurity density (solid, red), up–down asymmetry (dashed, blue), and in–out asymmetry (dashed–dotted, green). (b) symmetric impurity density (solid, red), $n = 1$ (dashed, blue), $n = 2$ (dashed–dotted, green), and $n = 3$ (dotted, black).

Figure 5. Real and imaginary parts of $\omega = \omega_r + i\gamma$ as a function of $a/L_{T_i}$ (a) and $a/L_{T_e}$ (b) obtained by GYRO for the baseline case. Red lines (with circle markers) represent the real part, blue lines (triangle markers) correspond to the imaginary part of the eigenvalue. The frequencies are normalized to $c_s/a$.

4.2. Temperature gradient dependence

The eigenvalues and electrostatic potentials as functions of ion and electron temperature gradients are shown in figures 5 and 6. As expected, if we increase the ITG, the turbulence becomes more ITG-dominated (the real part of the mode frequency $\omega_r$ is negative), while if we increase the electron temperature gradient, TE-mode-driven turbulence will dominate ($\omega_r$ is positive). The shape of the imaginary part of the potential $\text{Im}[\phi]$ varies slightly by increasing the temperature gradient for both the ITG- and TE-mode-dominated cases, while the real parts of the potential $\text{Re}[\phi]$ are not modified significantly, see figure 6. The imaginary part of the potential plays an important role in the parallel compressibility terms and therefore the change in the temperature gradient will modify the impurity peaking factor.

Previous works highlighted the difference in peaking factors between ITG- and TE-dominated cases, and concluded that ITG-dominated turbulence will always generate inward pinch of impurities, while in TE-mode-driven turbulence outward convection in the plasma core (for $r/a \approx 0.2$) can be expected. Both linear [7, 8] and non-linear [26] gyrokinetic
simulations have shown that the latter is due to the contribution from the parallel dynamics which can reverse the impurity convection from inward to outward for modes propagating in the electron direction. These results are in agreement, under some conditions, with the experimental observation that the impurity convection changes sign from inward to outward when a strong central peaking of the electron temperature arises as a response to strong localized central electron heating. However, these results cannot explain the outward convection of impurities observed in experiments with RF heating where the ITG mode is the dominant instability.

Our results show that the peaking factor can become negative in the ITG-dominated case if the impurities accumulate on the inboard side, and it is influenced very strongly by increasing temperature gradients.

Figure 7 shows the ITG scaling of the peaking factor and the diamond symbols show these values obtained by GYRO simulations in the poloidally symmetric limit, which present good agreement with our results. The peaking factor increases with the ITG as long as the impurity density is poloidally symmetric or if it is up–down or out–in asymmetric. However, if the impurity density is in–out asymmetric then the sign of the peaking factor is negative for large temperature gradients, if the asymmetry is sufficiently large. The threshold for the sign change for the experimental scenario studied in this paper is \( a/L_T \approx 2 \). As the strength of the asymmetry, \( n \), is increased for an in–out asymmetry the modifications of the peaking factor become stronger, see figure 7(b).

As the electron temperature gradient is increased the turbulence becomes more TE dominated. Figure 8 shows that the absolute value of the peaking factor is strongly reduced in magnitude in the presence of an in–out asymmetry; however, it remains positive. Stronger asymmetries are needed in the TE-dominated case in order to obtain a sign change.
Figure 7. Peaking factor for nickel as a function of ITG for different values of $\delta$ (a) and $n$ (b). In both figures the solid line represents the case of poloidally symmetric impurity distribution; in (a) this case is compared with GYRO simulations (black diamonds). (a) $n = 3$—out–in asymmetry (dashed), up–down asymmetry (dashed–dotted), in–out asymmetry (dotted). (b) in–out asymmetry—$n = 1$ (dashed, blue), $n = 2$ (dashed–dotted, green), $n = 3$ (dotted, black).

Figure 8. Peaking factor for nickel as a function of electron temperature gradient for different values of $\delta$ (a) and $n$ (b). In both figures the solid line represents the case of poloidally symmetric impurity distribution; in (a) this case is compared with GYRO simulations (black diamonds). (a) $n = 3$—out–in asymmetry (dashed, blue), up–down asymmetry (dashed–dotted, green), in–out asymmetry (dotted, black). (b) in–out asymmetry—$n = 1$ (dashed, blue), $n = 2$ (dashed–dotted, green), $n = 3$ (dotted, black).

4.3. Density gradient dependence

The density gradient scaling of the peaking factor is shown in figure 9. As the density gradient is increased the underlying instability changes from ITG to TE mode, see figure 10. The peaking factor is slightly decreasing with electron density gradient in the case of poloidal symmetry and in the cases of up–down or out–in asymmetries. Also here, the in–out asymmetric impurity density leads to a negative peaking factor if the strength of the asymmetry is sufficient. It is interesting to note that in the case of in–out asymmetry the peaking factor is quite sensitive to the density gradient. The difference between the two cases $a/L_n = 0.5$ and $a/L_n = 1.5$ can be understood by comparing the shape of the imaginary part of the electrostatic potential $\text{Im}\left[\phi\right]$ in figure 11(b). It can be noted that by changing the density gradient, $\text{Im}\left[\phi\right]$ changes significantly. Note that it is mainly the part of the potential which is close to $\theta = \pi$ which is important, and it
Figure 9. Peaking factor for nickel as a function of electron density gradient for different values of \( \delta \) (a) and \( n \) (b). In both figures the solid line represents the case of poloidally symmetric impurity distribution; in (a) this case is compared with GYRO simulations (black diamonds). (a) \( n = 3 \)— out–in asymmetry (dashed, blue), up–down asymmetry (dashed–dotted, green), in–out asymmetry (dotted, black). (b) in–out asymmetry— \( n = 1 \) (dashed, blue), \( n = 2 \) (dashed–dotted, green), \( n = 3 \) (dotted, black).

Figure 10. Real and imaginary parts of the eigenvalues as a function of \( a/L_{ne} \) obtained by GYRO. Red lines (with circle markers) represent the real part, blue lines (triangle markers) correspond to the imaginary part of the eigenvalue. The frequencies are normalized to \( c_s/a \).

is considerably different for the two density gradients shown in figure 11(b) and therefore also the result for the peaking factor changes dramatically. Even though for \( a/L_{ne} = 0.5 \) the ITG mode is the dominant instability the peaking factor remains positive even for strong in–out asymmetries. This is different from the ITG-dominated cases in the previous two subsections where the peaking factors were negative in the presence of an in–out asymmetry with the same strength. The difference in the electrostatic potentials is the underlying reason for the difference in the trends seen in figures 7–9.

4.4. Effect of parallel compressibility

It has been shown previously that when the transport is TE-mode-dominated parallel compressibility effects generate an outward contribution to impurity anomalous flux which can, under certain plasma conditions, cancel out the inward contributions, leading to zero or even negative impurity peaking factor [7, 8]. In this subsection, we examine this effect by neglecting the parallel compressibility terms, i.e. terms proportional to \( \delta_p \) in equations (6)
Figure 11. Real (a) and imaginary (b) parts of the electrostatic potentials for two different density gradients.

and (7). Note that in this case only the absolute value of the potential enters in the expressions, while in the case with parallel compressibility both the imaginary and real parts and their derivatives are important.

Figure 12 shows the peaking factor for various \( \delta \) and \( n \). The effect of poloidal asymmetry in this limit is in agreement with our previous work in [11] and is similar to that in the previous sections where the parallel compressibility effects are considered. As seen in figure 12(a), in the absence of parallel compressibility effects an in–out asymmetry can lead to a negative peaking factor (outward impurity flux). An increase in the poloidal in–out asymmetry will increase the outward flux of impurities as shown in figure 12(b). Note that in the case with parallel compressibility, the sign change in the peaking factor occurs for broader range of \( \delta \) and for lower asymmetry strength (compare figures 4 and 12). Also an up–down asymmetry can lead to a slight reduction in the impurity peaking factor. This is similar to the case where the parallel compressibility was taken into account, see figure 4(a).

Figure 13 shows the ITG scan for the peaking factor without parallel compressibility. The diamonds represent the values of the peaking factor obtained by GYRO (without parallel compressibility) which show agreement with our results in the symmetric limit. From figure 13 it is clear that without parallel compressibility the peaking factor is less sensitive to the increase in the ITG, and in the presence of an asymmetry, regardless of the asymmetry sign, when the transport is TE mode dominated the peaking factor remains positive. In the ITG-mode-dominant case, including the effects of parallel compressibility, the in–out asymmetry results in significantly larger negative peaking factors than without, compare figure 13 with figure 7(a).

5. Discussion and conclusions

For tokamak operation the avoidance of central impurity accumulation is a key issue. Conditions under which the convective impurity flux is directed outward are particularly
interesting. In order to find such conditions many experiments have been devoted to explore various techniques. One way to expel the impurities from the plasma core is to maintain sawtooth crashes in a controlled way by applying central ICRH \cite{27,28}. It is observed that this method will indeed remove the impurities from the plasma core. However, it was shown that even though sawtooth crashes hamper the accumulation of the impurities their contribution is less relevant compared with the effect of the ICRH itself \cite{5}. Another technique which has been successful in removing the impurities from the plasma core and is routinely used in tokamak experiments such as in the ASDEX-U tokamak is the application of a very localized central electron cyclotron resonance heating (ECRH) \cite{29}. Simulations with both linear and non-linear gyrokinetic models have shown that under these conditions the electron temperature gradient is strongly peaked and therefore, modes propagating in the electron diamagnetic drift direction are the dominant instability responsible for the turbulence-driven transport \cite{26}. The
interaction between the related electrostatic potential fluctuations and the parallel dynamics of the impurities leads to an outward convection of impurities. These results are in agreement with the experimental observation in ASDEX-U for the very central region \((r/a \simeq 0.2)\) [7]. In other tokamak experiments, for example at JET, the application of the central RF heating has also been explored with success. In these experiments it has been observed that the dominant instabilities are not TE modes but rather modes directed in the ion diamagnetic drift direction, ITG modes, which theoretically should result in an inward impurity flux. It has been debated that in the very central region of plasmas at JET the transport is mostly driven by neoclassical effects. By applying a strong central heating as the temperature gradient peaks the neoclassical temperature screening effects become dominant and result in an outward directed impurity flux. Under some plasma conditions these effects may be the reason for the observed behavior [2]; however, usually the observed transport is an order of magnitude larger than the neoclassical predictions implying that the impurity transport is turbulence driven [26].

In this work we discuss the impact of poloidal dependence of the impurity density on the impurity peaking factor in the core of tokamak plasmas. Various mechanisms may give rise to a poloidal asymmetry of impurity density: difference in impurity source location, toroidal rotation or neoclassical effects. Among these is an in–out poloidal asymmetry observed in plasmas where RF heating is applied. The mechanism responsible for this behavior was explained through RF-induced accumulation of minority ions on the outboard side of the torus giving rise to a corresponding impurity accumulation on the inboard side as was discussed in [13] and in section 2. The strength of the impurity accumulation depends on the impurity charge (higher for heavier impurities) and the plasma parameters such as temperature gradient. However, the exact form of these dependences is not yet known and further analysis is needed in this area. It is out of the scope of this paper to make an extensive analysis of the above-mentioned phenomena; the main objective of this paper is to demonstrate theoretically, assuming a simple ansatz for the poloidal asymmetry of impurity density, what effect it would have on the impurity transport.

In the parameter scans, for simplicity, we only presented the asymmetry function \(P(\theta, \delta, n)\) for a specific asymmetry strength, with weight \(f_j = 1\). The total peaking factor can be estimated from the sum of the various peaking factors weighted in an appropriate way. For instance, for an asymmetric population, such that \(n_z = n_{z0} + n_{z1}\) with \(\langle n_{z0} \rangle / \langle n_{z1} \rangle = f_0 / f_1\) (and \(f_j = 0\) for \(j \neq 0, 1\)) the peaking factor can be estimated to be

\[
\frac{a}{L_{n_z}} \simeq \frac{a}{\langle n_0 + n_1 \rangle} \frac{\partial \langle n_0 + n_1 \rangle}{\partial r} \simeq \frac{f_0}{f_0 + f_1} \frac{a}{L_{n_0}} + \frac{f_1}{f_0 + f_1} \frac{a}{L_{n_1}},
\]

where \(a/L_{n_0}\) is the peaking factor corresponding to the poloidally symmetric part \((n = 0, \text{ with weight } f_0)\), and \(a/L_{n_1}\) is the peaking factor corresponding to the poloidally asymmetric part \((n = 1, \text{ with weight } f_1)\).

We have found that an in–out poloidal asymmetry of the impurity density \((\delta = \pi)\) can lead to an outward impurity flux (negative peaking factor) in both ITG- and TE-mode dominated cases, and this effect becomes stronger as the asymmetry strength \((n)\) increases. However, stronger asymmetries are needed in the TE-mode-dominated case in order to obtain a sign change. The level of asymmetry that is needed to change the sign of the peaking factor is too large in order for this effect alone to explain the observed outward convection. The RF-induced asymmetry can give rise to a poloidal variation of the order of 10–20% and that is much smaller than what is necessary for a sign change. However, together with other effects (e.g. neoclassical temperature screening, sawteeth) poloidal asymmetries can also contribute to the avoidance of central accumulation of impurities. The sign and magnitude of the peaking factor is sensitive not only to the asymmetry strength but also to the temperature gradient.
In the ITG-mode-dominated case the in–out asymmetry results in a more negative peaking factor as the ITG increases while in the TE-mode-dominated case as the electron temperature increases the peaking factor becomes more positive.

Our results indicate that an out–in asymmetry leads to an increase in the peaking factor, as it is observed in the NBI-heated plasmas where the out–in asymmetry is generated by the centrifugal force on the impurity ions. These effects of the background toroidal rotation on the heat and particle transports have been previously studied, see for example [30, 31], where gyrokinetic simulations have shown that in the ITG-dominated case, the convective impurity pinch increases with the rotation (more strongly with increasing impurity mass). It is also shown that an up–down asymmetry can lead to a reduction in the peaking factor. This can be a contributing reason for the observed flat impurity density profiles in plasmas with ECRH, and up–down asymmetries have indeed been observed in EC-heated plasmas [15], although the physical mechanism for these asymmetries is not entirely understood.

The reason for the sign change of the impurity peaking factor in the presence of a poloidal asymmetry is attributed to the interaction between the poloidal variation of the related electrostatic potential and the poloidal dependence of the impurity density. If parallel compressibility effects are taken into account the imaginary part of the electrostatic potential is the determining factor in the sign change of the impurity peaking factor.

In summary, our results suggest that poloidal asymmetries can significantly alter the turbulence-driven impurity transport and therefore, have to be taken into account. These asymmetries may have a significant role in determining the impurity accumulation properties in plasmas with radio frequency heating. Therefore, there is a strong need for development of new tools in order to detect poloidal asymmetries and determine the effect of RF heating in their generation and the dependence of the asymmetry function on various plasma parameters.

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