Interchange transport in electron-positron plasmas with ion impurities

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Interchange drive and cross-field transport of density filaments in quasi-neutral inhomogeneously magnetized electron-positron plasmas is shown to be strongly reduced by the presence of minority ions. Two mechanisms are identified for the reduction in radial propagation and plasma transport: effective mass related inertia, and collisionality dependent Boltzmann spin-up of the filaments. Numerical results are obtained with a three-dimensional full-F multi-species gyrofluid model.

I. INTRODUCTION

Plans and first progress in laboratory confinement of quasi-neutral electron-positron (e-p) plasmas in toroidal magnetic fields [1][2] has also generated renewed theoretical interest in magnetized e-p plasmas [3][4].

It was recently shown that magnetic field-aligned density filaments (in the fusion plasma community often named “blobs” [11][12]) in an inhomogeneously magnetized e-p plasma are interchange unstable for a range of parameters accessible in planned experiments and could lead to crucial transport losses [13].

In the following it is demonstrated that replacing a fraction of the positrons with “impurity” ions effectively reduces the interchange propagation and transport of such e-p density filaments. The relevance of filamentary transport lies in the self-propellation of such elongated density perturbations down any magnetic field gradient, which does not require a background temperature or density gradient for stimulating transport. The initial perturbation may arise out of any plasma density or electric field inhomogeneity.

It is here shown that both the effective mass dependent polarization inertia (which is also active in a 2-d model) and the Boltzmann spin-up (which is a 3-d effect depending on Coulomb collisionality between the species) contribute to the filament propagation reduction.

In sec. II the 3-dimensional full-F gyrofluid model used for the numerical simulations is discussed, and in sec. III it is argued why a delta-f model (evolving only small fluctuations on a constant background plasma) is inappropriate for describing interchange instability in e-p plasmas. The (full-F) computational results are presented in sec. IV, and conclusions are given in sec. V.

II. FULL-F E-P-I GYROFLUID MODEL

We analyse magnetized e-p-i plasmas by means of a nonlinear three-dimensional full-F multi-species gyrofluid model, which is based on the 6-moment derivation of Madsen [10] from a gyrokinetic model that evolves the full distribution function $F(x, v, t)$, including a first order finite Larmor radius (FLR) closure.

In the isothermal electrostatic limit [17][18] the full-F 3-d gyrofluid model consists of normalized continuity and momentum equations for the gyrocenter densities $n_s$ and parallel velocities $v_s$ for all species $s$, which here are given by electrons, positrons and ion with $s \in (e, p, i)$.

$$
\partial_t \tilde{n}_s = \frac{1}{B} [\tilde{n}_s, \phi] - \frac{B}{n_s} \nabla || \left( \frac{n_s v_s}{B} \right) + \mathcal{K}(\tilde{n}_s) \tag{1}
$$

$$
\partial_t v_s = \frac{1}{B} \left[ v_s, \phi \right] - \frac{1}{\mu_s} \nabla || h_s - C_0 J || + \tau_s v_s \mathcal{K}(\tilde{n}_s) + 2\tau_s \mathcal{K}(v_s). \tag{2}
$$

Here $h_s \equiv (\phi + \tau_s \tilde{n}_s)$ is abbreviated. The (species specific) gyro-screened potentials $\phi_s = \Gamma_1 \phi - (\mu_s/2B)(\nabla \phi)^2$ include both FLR and ponderomotive effects.

The total parallel current $J ||$ is given by $J || = \sum_s n_s Z_s e v_s$. We have neglected triple nonlinear terms involving the parallel velocities, and also electromagnetic fluctuations of the vector potential $A ||$.

The gyrocenter densities $n_s$ are normalized to a constant reference density $n_0$, so that the magnitude of the plasma (electron) density $n_e \leftarrow n_e/n_0$ is of order one. Eqs. [1][2] have been divided by the gyrocenter densities $n_s$, so that logarithmic densities $\tilde{n}_s \equiv \ln n_s$ are introduced as the evolving quantity to ensure positivity, with both $\tilde{n}_s$ and $n_s$ appearing in the equations. The potential is normalized to $T_e/e$, perpendicular length scales to the ion drift scale $\rho = \sqrt{T_e/m_i}/(eB)$, and time to $c_s/\rho$ with sound speed $c_s = \sqrt{T_e/m_i}$, $T_e$ and $m_i$ are the temperature and mass of the species $s$, and $B_0$ is a reference magnetic field strength. Parallel derivatives are further scaled as $\nabla || \leftarrow (L_||/L_\perp) \nabla ||$ with the connection length $L_||$, which for toroidal geometry is given by $L_|| = 2\pi qR$ with inverse rotational transform $q$ and major torus radius $R$. The drift parameter $\delta = \rho_s/L_\perp$ is used to set the perpendicular length scale $L_\perp$.

The nonlinear quasi-neutral polarisation equation

$$
\sum_s \left[ Z_s \Gamma_1 n_s + \nabla \cdot \left( n_s \frac{\mu_s}{Z_s B^2} \nabla \phi \right) \right] = 0. \tag{3}
$$
determines the electrostatic potential $\phi$ for given gyrocenter densities $n_s$.

In the (2-d) model used in ref. [12] for studying interchange transport in pure e-p plasmas we had included Debye length effects into the polarisation equation. This had restored the original “Poisson” term from the electrostatic Poisson equation $\sum_s Z_s N_s + \varepsilon \nabla^2 \phi = 0$, in order
to determine the (strongly damping) influence of Debye screening on filament propagation. The Debye parameter \( \varepsilon = (\lambda / \rho)^2 \) represents effects of finite Debye length \( \lambda = \sqrt{\epsilon_0 T_e / (e^2 N_{e0})} \) in relation to the drift scale or Larmor radius \( \rho_s = \sqrt{T_s m_s / (eB)} \) [10]. Here we specifically neglect these Debye effects and set \( \varepsilon = 0 \) (and thus can assume exact quasi-neutrality in the polarization equation), in order to focus only on the influence of ion impurities on e-p filament dynamics and to reduce the number of free parameters in the model.

The particle densities \( N_s \) are linked to the gyrocenter densities \( n_s \) by the relation

\[
N_s = \Gamma_b n_s + \nabla \cdot \left( \frac{n_s \mu_s}{Z_s B^2} \nabla \phi \right). \tag{4}
\]

The charge states are \( Z_e = -1 \) for electrons and \( Z_p = +1 \) for positrons. \( \Gamma_b \) depends on the impurity ion species and ionization degree, but for the expected low temperatures of magnetized e-p laboratory plasma in the range of a few eV we may assume predominantly singly ionized atoms or molecules with \( Z_i = +1 \). Neutral impurities (and their ionization and recombination processes) are here neglected. We also specifically neglect electron-positron annihilation, which can for low e-p densities be expected to occur on much longer time scales compared to the instability growth times [20]. In principle, annihilation (or positronium formation) rates could be easily included as sink terms on the right hand side of eq. (4).

The gyro-averaging operator in Padé approximation is defined by \( \Gamma_{1s} = (1 + (1/2)b_s)^{-1} \) with \( b_s = \tau_s \mu_s \nabla \phi \). The mass ratio is given by \( \mu_s = m_s/(Z_s m_e) \), and the (constant) temperature ratio by \( \tau_s = T_s/(Z_s T_e) \). For electrons, thus \( \tau_e = -1 \), and for positrons we assume an equal constant temperature so that \( \tau_p = +1 \). Ions are assumed to be cold with \( \tau_i = 0 \) so that also \( b_s = 0 \).

Our model in principle can resolve all FLR effects, but in the following we will neglect these also for electrons and positrons, and set \( b_s = b_p = 0 \). Temperature dynamics and gradients could further influence the filament propagation results [21], but we here assume the e-p plasma to be cool and isothermal.

The 2-d advection terms are expressed through Poisson brackets \([f, g] = (\partial_x f)(\partial_y g) - (\partial_y f)(\partial_x g)\) for local coordinates \(x\) and \(y\) perpendicular to \(B\). Normal and geodesic magnetic curvature enter the compressional effect due to field inhomogeneity by \( K = \kappa_y \partial_x + \kappa_x \partial_y \) where the curvature components in toroidal geometry are functions of the poloidal angle \( \theta \) mapped onto the parallel coordinate \( z \). For a circular torus \( \kappa_y \equiv \kappa_0 \cos(z) \) and \( \kappa_x \equiv \kappa_0 \sin(z) \) when \( z = 0 \) is defined at the outboard midplane. The toroidal magnetic field strength is assumed to vary only in parallel direction as \( B(z) = 1 + a \cos(z) \) with inverse aspect ratio \( a \).

The collisionality parameter in the parallel velocity equation is given by \( C = 0.51(\nu_e L/\epsilon_s) \). We note that the collisionality term in the corresponding (electromagnetic) equation for the momentum given in ref. [18] was written as \( C(J \parallel /n_s) \). However, the parameter \( C \sim \nu_e \) also includes a direct density proportionality in the electron/positron-ion collision frequency \( \nu_e \sim n_e \), which cancels the inverse density factor in the collisionality term, so we here use a constant \( C_0 \sim \nu_e(n_0) \), evaluated at a fixed reference density \( n_0 \). The weak density dependence in the Coulomb logarithm is neglected.

For numerical stability, a small perpendicular hyperviscosity term \( -\nu_s \nabla^2 \tilde{n}_s \) is added on the right hand side of eq. (1), and in 3-d computations parallel viscous terms \( \nu \partial_z^2 \tilde{n}_s \) are added to eqs. (1) and (2), respectively. Boundary conditions in y direction are periodic for 2-d simulations, and quasi-periodic (shear-shifted flux tube) for 3-d simulations. The further numerical methods are presented in ref. [18].

### III. INADEQUACY OF A DELTA-F MODEL

The common delta-f isothermal gyrofluid model [22] is regained by splitting \( n_s = n_{s0} + \tilde{n}_s \) into a static constant background density \( n_{s0} \) and the perturbed density \( \tilde{n}_s \). When \( \tilde{n}_s/n_{s0} \ll 1 \), the right hand sides of eqs. (1) and (2) can be linearized by approximating \( n_s \approx n_{s0} \) so that \( \nu_e \approx \tilde{n}_s + (\tilde{n}_s/n_{s0}) \), and neglecting all nonlinear terms except the Poisson bracket:

\[
\partial_t \tilde{n}_s = \frac{1}{B} [\tilde{n}_s, \tilde{\phi}_s] - B \nabla \cdot \left( \frac{\tilde{v}_s}{B} \right) + K(\tilde{\phi}_s) \tag{5}
\]

\[
\partial_t \tilde{v}_s = \frac{1}{\mu_s} [\tilde{v}_s, \tilde{\phi}_s] - \nabla \cdot \left( \frac{\tilde{n}_s}{\mu_s} \right) + 2 \tau_s K(\tilde{v}_s) - \frac{C_0}{\mu_s} \tilde{j}_\parallel \tag{6}
\]

The consistent delta-f polarisation equation in the high-\( k \) limit is \( \sum \alpha_s \Gamma_{1s} \tilde{n}_s + (1/\tau_s)(\Gamma_{0s} - 1)\tilde{\phi}_s \) = 0 with \( \Gamma_{0s} = (1 + b_s)^{-1} \). Linearisation of the low-\( k \) eq. (3) does not include gyro-screening on the potential and gives \( \sum \alpha_s \Gamma_{1s} \tilde{n}_s = (\sum \alpha_s \mu_s) \nabla^2 \tilde{\phi}_s \). The velocities and current are coupled in the electrostatic limit by \( J_\parallel = \sum \alpha_s \tilde{v}_s \). The parameter \( \alpha_s = Z_s n_{s0}/n_{s0} \) denotes the ratio of species reference densities \( n_{s0} \) to \( n_{0} \).

These delta-f gyrofluid equations are a good approximation to the full-F model for example in core and mid-pedestal e-i fusion plasma turbulence simulations [24], where density fluctuations indeed are usually much smaller than the average background plasma density.

The applicability of the delta-f multi-species model, which assumes a quasi infinite background density for all of the species, however specifically fails, when one of the species has a much smaller or vanishing density compared to the others.

Then the term \( K(\tilde{\phi}_s) \approx K(\tilde{\phi}_s + \tau_s \tilde{n}_s) \) inconsistently would generate by \( \partial_t \tilde{n}_s \sim K(\tilde{\phi}_s) \) new density out of any appearing inhomogenous potential fluctuation \( \tilde{\phi}_s \) even if the initial species density (fluctuation) \( \tilde{n}_s \) was zero.

This artefact is not present in the full-F model, which is evident when we do not write eq. (11) in terms of the logarithmic density \( \tilde{n}_s \), but originally as

\[
\partial_t n_s = \frac{1}{B} [n_s, \phi_s] - B \nabla \cdot \left( \frac{n_s v_s}{B} \right) + n_s K(\phi_s). \tag{7}
\]
Here the interchange curvature term \( n_s K(h_s) \) can drive changes in density only in proportion to the locally present species density \( n_s(x,t) \).

**IV. E-P-I FILAMENT PROPAGATION**

Pressure perturbations in magnetized plasmas experience interchange forcing due to an inhomogeneity (gradient and curvature) of the magnetic field, which leads to a “radial” propagation across the magnetic field. Perturbations in toroidal plasmas are mostly flute-like and strongly elongated along the magnetic field direction, and appear as plasma filaments. In the fusion plasma literature, filamentary pressure perturbations at the plasma edge with positive amplitudes are commonly named “blobs”, or “holes” for negative amplitudes.

The basic gradient and curvature drift dynamics, which differ in sign between positive and negative plasma species by their charges, results in an \( E \times B \) drift which radially advects the perturbation. Initially symmetric (e.g. Gaussian) shapes of the perturbation across the field then develop into mushroom shaped plume structures. These effects are already present in 2-d (perpendicular to the magnetic field direction) fluid models.

Along the magnetic field direction the dynamics is usually more wave-like (compared to the fluid-like advection across the field), and pressure perturbations can induce sound waves or Alfvén waves (which are however neglected in the present work, with \( \beta = 0 \)). Collisional coupling between the species can lead to deviations from an adiabatic response on perturbations.

When the initial perturbation filaments are extended with constant amplitude everywhere along the magnetic field \( (z) \) direction, the dynamics again becomes quasi-2-dimensional (except for effects of magnetic shear). Here we rather consider perturbations that are also initially localized in the parallel direction with some parallel width \( \Delta z \) in the maximum ballooning region (which here defines \( z = 0 \)). Then the perturbation will experience spreading along the field direction by pressure driven expansion.

The sound speeds of electrons and positrons are (for equal temperature) identical due to the same mass, but the sound speeds differ substantially between electrons and the much more massive ions. The more rapid excursion of electrons from (initially neutral) pressure perturbations together with ion inertia leads to a positive charging of the perturbation, which again slows down the electrons into an ambipolar parallel diffusion. In e-i plasmas this arising potential perturbation leads to a vertical \( E \times B \) drift around the perturbation, which spins the blob into an eddy and effectively slows down the radial interchange drive of the whole filament. This effect, named “Boltzmann spinning” in the fusion plasma literature [23], is absent in pure mass-symmetric pair plasmas.

In the following, effects of the presence of some fraction of ion impurities in an e-p pair plasma on interchange driven filament transport will be studied. Boltzmann spinning of localized perturbations may be expected to slow the e-p blobs depending on ion concentration. This effect will be addressed with 3-d simulations.

**A. Inertial mass effect through polarization**

But already in a simplified 2-d setup another species mixture effect on blob propagation can be expected by changes of the effective mass of the plasma, which enters into polarisation dynamics mediated by eq. (3).

Linearisation of the polarisation equation (without FLR effects) gives

\[
\sum_s a_s \tilde{n}_s = \left( \sum_s a_s \mu_s \right) \nabla^2 \tilde{\phi} \equiv \tilde{\mu} \Omega. \tag{8}
\]

The development of \( E \times B \) vorticity \( \Omega = \nabla^2 \tilde{\phi} \) out of density perturbations is thus mediated by an effective mass \( \tilde{\mu} = \sum_s a_s \mu_s \). Here \( a_s = Z_s n_{a0}/n_{e0} \) and \( \mu_s = m_e/(m_e Z_s) \), when the electron mass \( m_e \) is used as a reference. For a pure e-p pair plasma (with \( a_i = 0 \)), \( a_e = \mu_e = -1 \) and \( a_p = \mu_p = +1 \), so that \( \tilde{\mu} = 2 \).

The effective mass in an e-p-i plasma is given by \( \tilde{\mu} = a_e \mu_e + a_p \mu_p + a_i \mu_i = (a_e - a_p) \mu_e + a_i \mu_i \). When the electron density is kept constant and a fraction of positrons in an e-p plasma is replaced by ions, then \( a_e = -1 \) and \( a_p = 1 - a_i \), so that \( \tilde{\mu} = 2 + (\mu_i - 1) a_i \approx 2 + \mu_i a_i \).

The effects of variations in the effective mass on interchange driven filaments and turbulence in the edge and scrape-off layer of tokamak fusion plasmas has been recently investigated for the similar ion masses in hydrogen isotope mixtures [26, 28], where relevant changes have been found, so that even stronger effects can be expected for the e-p-i system with large mass differences between the positive species. Significant changes should occur at least for ion density ratios \( a_i > 1/\mu_i \).

In the following numerical examples we assume hydrogen ions as the impurity species, so that \( \mu_i = m_i/(m_e Z_i) \approx 1836 \). When the impurities are generated by e-p plasma-wall interactions or by rest gas contamination in an imperfect vacuum chamber, the ion masses can be larger, depending on the present atomic or molecular species. For thermal ionization the charge state \( Z_i \) of impurity ions will likely be single, but ionization by annihilation photons could lead to stronger degrees of ionization in impurity species.

For simplicity we here thus only consider hydrogen ions (protons) and keep in mind that heavier species would more enhance the reported mass effects. The ion contribution to mass inertia is of order unity and larger when \( a_i \mu_i > 1 \), or when \( a_H > 1/1836 \approx 5 \cdot 10^{-4} \). For comparison, for singly charged iron impurities (from the chamber wall) with \( \mu_{Fe} \approx 56 \mu_H \), the critical concentration, above which inertial mass effects become relevant, would be around \( a_{Fe} \approx 10^{-5} \).

We first investigate the inertial mass effect through \( \tilde{\mu} \) in 2-d simulations. For this we numerically evaluate...
eqs. (1) and (3) for $v_s = 0$ at the location $z = 0$, where the normal curvature is maximum and the geodesic curvature contribution vanishes. As mentioned above, this corresponds to a case of highly elongated filaments.

A Gaussian initial density perturbation with perpendicular width $\sigma = 4\rho$ and amplitude $\Delta n_e = 0.5$ is set on an otherwise homogeneous density background with $n_e(0) = 1.0$. The magnetic curvature is set to $\kappa_0 = 0.01$. The computational domain is $L_x \times L_y = 64 \rho \times 32 \rho$ on a rectangular numerical grid with $n_x \times n_y = 128 \times 64$. Higher grid resolution leads to nicer resolved pictures of the blobs, but does not change the results significantly.

A fraction $a_i$ of positron density is replaced by (hydrogen) ions, and is varied between $a_i = 0$ for a pure e-p pair plasma, up to $a_i = 1$ for a pure e-i plasma.

The average interchange transport by radial blob propagation is determined by $\Gamma_n = \langle n_e v_x \rangle_{x,y}$ with $v_x = \partial_y \phi$. The transport as a function of normalized time is shown in Fig. 1; it increases to a maximum as long as the radial blob propagation velocity accelerates, and then drops again to low levels. Nonlinear breakup of the blob leads to a more unsymmetric decay phase. We observe that the maximum transport (as well as the maximum blob velocity) is strongly reduced with increasing ion fraction by a factor $1/\sqrt{\bar{\mu}}$.

The time scale $T_{max}$ for acceleration, until the maximum velocity and transport level are reached, on the other hand grows with increasing ion fraction: the interchange growth rate and propagation velocity are shifted from electron to ion time scales by the same factor $1/\sqrt{\bar{\mu}}$.

The total integrated transport, which we here (because of the nearly symmetric shape of $\Gamma_n(t)$) approximate as $\int dt \Gamma_n(t) \approx \langle 1/2 \rangle \Gamma_n(0) T_{max}$, is therefore largely independent of the ion impurity fraction, but is only spread over different time scales. Both the maximum transport and the approximate integrated transport are shown as a function of the ion fraction in Fig. 2. The black dots are the numerical values of the maxima from the simulations shown in Fig. 1 and the black line is the analytical function $\Gamma_n(a_i) = \Gamma_n(0)/\sqrt{\bar{\mu}}$, with $\bar{\mu} = 2 + 1836 a_i$.
FIG. 4: Evolution of a density blob (2-d x-y cross sections of 3-d simulations) in an inhomogeneous magnetic field. From top to bottom: (a) initial perturbation at $t = 0$; (b) e-p blob at $t = 50$; (c) e-p blob at $t = 100$; (d): e-p blob with 1% ion fraction, showing Boltzmann spin-up.

B. Boltzmann spinning effect

The inertial polarization mass effect, which effectively scales the blob propagation time, is also still present for 3-d simulations of elongated filaments. For a finite initial filament extension $\Delta z$ along the field line, the time scale of radial advection then competes with the time scales of parallel spreading and charging.

In 3-d the collisionality $C_0$, the filament extension $\Delta z$, and the parallel-to-perpendicular scale ratio $\hat{\varepsilon} = (qR/L_\perp)^2$ enter as additional parameters and control the non-adiabatic electron response.

The difference in electric potential generation by parallel evolution for pure e-p compared to e-i plasmas is shown in Fig. 4 as 2-d x-y cross sections (at $z = 0$) of 3-d simulations for various times. The cross section shows the computational region of $64 \times 32 \rho$ like in the other simulations above, and the blob is again initialized with $\sigma = 4\rho$, now with parallel width $\Delta z = 1/8$. On the top, picture (a) shows the initial density $n_e(x, t = 0)$ with the same Gaussian perturbation for both e-p and e-i blobs. Picture (b) and (c) show the evolution of a pure e-p blob at the times $t = 50$ and $t = 100$, respectively. The radial propagation velocity and the associated outward density transport are maximal at around $t = 50$. Picture (d) shows at $t = 100$ the onset of density spin-up by $E \times B$ drift advection in an e-p blob with 1% ion impurity fraction ($a_i = 0.01$), which effectively suppresses the outward interchange driven propagation of the filament. Note that the density color scale has been adapted in (d), as the blob also loses amplitude at the shown location at $z = 0$ due to parallel spreading.

The combined effect of mass inertia and Boltzmann spinning on filamentary e-p transport is now computed.
for varying ion impurity densities. In Fig. 5 the maximum transport (averaged over the parallel coordinate) is again shown as a function of the ion fraction \( a_i \). Black lines and circle symbols denote the simulation results for \( C_0 = 30 \), while the red lines and diamond symbols denote the results for \( C_0 = 300 \). Further simulation results (not shown here) for a completely adiabatic response with \( C_0 = 0 \) are nearly identical (slightly smaller) compared to the results for \( C_0 = 30 \). This range approximately covers values that may be expected for low-temperature e-p laboratory plasmas of a few eV.

The dotted red line shows the analytical estimate for the inertial mass effect on \( \Gamma_{\infty}(a_i) = \Gamma_{\infty}(0)/\sqrt{2 + 1836/a_i} \) for \( C_0 = 300 \). While in the 2-d case the simulation results were nearly exactly lying on the analytical graph, we here see a systematically lower transport level, although the transport still approximately follows the overall analytical trend of inertia. The further reduction is a combination of density decrease by parallel spreading and of a suppression of radial filament propagation by Boltzmann spinning. While the 2-d integrated transport was independent of the ion impurity fraction, we observe a significant reduction of the values for \((1/2)\Gamma_{\infty}(a_{\text{max}})/\Gamma_{\text{max}}\) with \( a_i \) (depicted by the dashed black and red lines connecting the simulation values) by the 3-d Boltzmann spinning effect. The strongest change of filament transport by \( a_i \) in Fig. 5 still occurs for values around \( \mu_i a_i \sim 1 \), which for the presently assumed hydrogen ions is for \( a_i \sim 1/1836 \approx 0.5 \cdot 10^{-3} \). For more massive impurities the ion effect on e-p interchange transport would accordingly occur already for lower density fractions.

V. CONCLUSIONS AND OUTLOOK

To summarize, we have presented the first computations of interchange transport in inhomogeneously magnetized e-p plasmas with impurity ions. The reduction of transport with increasing ion fraction \( a_i \) roughly follows the inertial mass scaling, and is additionally reduced by Boltzmann spinning which depends on parallel localization of the filament and on the dissipative parallel coupling between leptons and ions.

Is this effect, after all, in any way relevant? Can, for example, a significant impurity density be expected in planned e-p confinement experiments? The parameters of future experiments \( \mathcal{E} \), like achievable e-p densities, temperatures, or radial profiles have large uncertainties. Any reliable predictive theoretical modelling of confinement properties and expectable modes and instabilities is thus not honestly possible. Theory can for now only stake out likely effects and trends.

The higher edge temperatures and more energetic edge localized transport events in magnetized fusion plasmas lead to sputtering and erosion of the plasma-facing wall components, which may enter the confined plasma region as impurity ions. In low-temperature e-p experiments the impurity content may be much lower, but also depends on the purity of the initial vacuum. For iron impurity ions from the vacuum chamber, the critical concentration where interchange mass effects would become noticeable is around \( a_i \sim 10^{-5} \), which is not completely unrealistic. But only the first real experiments will be able to clarify the e-p plasma purity.

On the other hand, the impurity concentration could also be set on purpose to probe the e-p to e-i transition by injecting for example hydrogen ions into a confined e-p plasma. Such dedicated e-p-i experiments would be able to test and validate our theories and models of plasma physics, which would be of general value for other areas like magnetic confinement fusion research.

However, in low density e-p plasmas the mass effect on interchange driving will not appear alone but in context with the Debye screening studied in ref. \[13\]. There we had derived the interchange growth rate (and accordingly the radial propagation velocity and associated transport) to be proportional to \( \gamma \sim 1/\sqrt{\mu + \epsilon} \), where \( \mu = 2 + a_i \mu_i \) was fixedly set to 2 for the pure e-p plasma. The values for the Debye parameter \( \epsilon \) have been estimated to be in the range of 50-300 for planned experiments. This implies that any ion impurity concentration effect will only become inertially relevant for \( a_i \mu_i \) in a similar order of magnitude as \( \epsilon \), or above. For a value of \( \epsilon = 200 \) and hydrogen ions, this would require a concentration of around 10 %, which appears to be unrealistically large for chance wall or rest gas impurities to be of any relevance. If the ion mass effect and any e-p to e-i physics transition should be tested on purpose, then a larger e-p plasma density in the order of the Brillouin density would be required in the experiments to overcome the Debye damping. So it can be concluded that an ion mass effect on interchange transport for vanilla operating conditions is likely to be subdominant.

However, we have so far ignored an additional possible ion impurity mechanism in magnetized e-p plasmas: the presence of ions is expected to be able to trigger the onset of resistive drift wave or drift-Alfven wave instabilities and associated turbulence in e-p plasmas in the presence of a background density gradient. The turbulent transport resulting from e-p-i drift wave turbulence may still turn out to be detrimental for magnetic e-p confinement, if it is not also effectively damped by Debye shielding. The computational investigation of fully developed e-p-i drift wave turbulence is however rather expensive because of the high required resolution to resolve the disparate electron/positron and ion drift scales appropriately. In particular the necessity to use full-F models, which are computationally also much more demanding than delta-f models, presently slows down the acquisition of results. These will therefore have to be reported in a future work.
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