Effects of multiple ion species on two-dimensional trapped-electron modes

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ABSTRACT

Previous calculations for the two-dimensional spatial structure of drift and trapped-electron modes in tokamaks are generalized to include an arbitrary number of ion species. Thus, it is possible to determine the influence of light or heavy impurity ions, and of multiple hydrogen isotopes, on these instabilities. As an example, results are presented for the case of a single light impurity species. Here the effect is found to be either stabilizing or destabilizing, depending on specific parameters, such as the equilibrium density profiles. Moreover, if a sufficiently idealized model is adopted, beam ion effects can also be studied. In the case considered, corresponding to the neutral-beam-heated PLT, these effects are found to be stabilizing.

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I. INTRODUCTION

Previous calculations¹ for the two-dimensional spatial structure of drift and trapped-electron modes in tokamaks are generalized to include an arbitrary number of ion species. Thus it is possible to determine the influence of light or heavy impurity species, and of multiple hydrogen isotopes, on these instabilities. Moreover, if a sufficiently idealized model is adopted, beam ion effects can also be studied.

There is no previous work which is actually comparable to the present nonlocal toroidal calculation. Tsang² investigated the effects of impurities on the trapped-electron mode, but only with a very approximate analytic one-dimensional radial analysis. More recently, Tang et al.³ carried out a numerical one-dimensional radial analysis of the effects of a single light impurity species on collisionless drift modes, but for a sheared-slab geometry. There have been a number of other papers on the effects of impurities on the trapped-ion mode, and also on new branches of drift modes specifically due to the presence of impurities (see Ref. 4 for relevant references). However, these are basically local calculations.

In Sec. II we describe the extensions of the calculation in Ref. 1 necessary to treat multiple ion species. As a specific example, results are presented in Sec. III for the case of a single light impurity species. Here the effect is found to be

either stabilizing or destabilizing, depending on particular parameters, such as the equilibrium density profiles. In Sec. IV, adopting a rather idealized model, the effect of including a beam ion species is studied. The corresponding beam distribution function can be strongly non-Maxwellian in its energy dependence. In the case considered, corresponding to the neutral-beam-heated PLT, the presence of the beam species is found to be stabilizing. Finally, we discuss our conclusions from this investigation in Sec. V.

II. EXTENSION OF TWO-DIMENSIONAL CALCULATION FOR MULTIPLE ION SPECIES

The two-dimensional calculation of Ref. 1 was for a pure plasma, i.e., a single hydrogenic ion species and electrons. Here we describe the extension of that calculation to the case of multiple ion species. These could be different isotopes, like hydrogen (H) and deuterium (D), or impurities like carbon (C). They could even include a beam species (b), as long as the present Maxwellian approximation for the equilibrium distribution function is adequate (when this is not the case, cf. Sec. IV), and as long as $|\omega|$ is larger than $\bar{\omega}_{tb}$, the circulating beam ion transit frequency.

The condition of bulk neutrality,

$$n_e(r) = \sum_j z_j n_j(r) \quad , \quad (1)$$

has to be satisfied at all r , where the summation is over all ionic species. Here $z_j |e|$ is the charge of ionic species j .

Thus

$$1 = \sum_j z_j n_j / n_e \quad , \quad (2)$$

and, differentiating Eq. (1) with respect to r ,

$$1 = \sum_j z_j \frac{n_j}{n_e} \frac{r_{ne}}{r_{nj}} \quad (3)$$

where $r_{nj} \equiv -(d \ln n_j / dr)^{-1}$, and all quantities are evaluated at $r = r_0$.

The perturbed electron density $\hat{n}_e$ is the same as that used in Ref. 1, except for the electron collision frequency. The effective electron collision frequency which is used here is

$$\nu_f(\epsilon, \lambda) = \left(\nu_f^L + \nu_f^H \frac{\nu_e^*}{X^3 Z_f} [Z_f + H_{ee}(X)] \right) \times \left(1 + \frac{\nu_e^*}{X^3 Z_f} [Z_f + H_{ee}(X)] \right)^{-1}, \quad (4)$$

where

$$\nu_f^L = \frac{\nu_e}{X^3} \frac{Z_f + H_{ee}(X)}{|(1 - \epsilon_0) - \lambda|} \quad (5)$$

$$\nu_f^H = 1.08 \nu_e \frac{Z_f + H_{ee}(1)}{1 + H_{ee}(1)}, \quad (6)$$

$$Z_f \equiv \sum_j Z_j^2 n_j / n_e, \quad \nu_e = 1221.8 n_e \lambda_e / T_e^{3/2}, \quad \lambda_e = 15.94 - (1/2)$$

$\ln(n_e / T_e^2)$, n_e is in units of 10^{13} cm^{-3} , T_e is in units of keV,

ν_e is in units of sec^{-1} , $\nu_e^* \equiv \nu_e Z_f / (\epsilon_0^{3/2} \bar{\omega}_{te})$, $H_{ee}(1) = 0.6289$,

and the rest of the notation is as defined in Ref. 1. Equation (4) constitutes the most straightforward generalization of the electron collision frequency of Ref. 1 to the case of multiple ion species, using the results tabulated in Ref. 5, and is consistent with that used in Ref. 6 for the trapped-ion mode.

Generalizing in the same fashion the effective collision frequency that will be used for ionic species i is

$$\nu_i = 1221.8 z_i^2 \frac{(m_e/m_i)^{1/2}}{T_i^{3/2}} \sum_j n_j z_j^2 \lambda_{ij} H_{ee} \left[\left(\frac{m_i T_i}{m_i T_j} \right)^{1/2} \right], \quad (7)$$

where

$$\lambda_{ij} = 18.39 - \ln \left[\frac{z_i z_j (m_i + m_j)}{m_i T_j + m_j T_i} \left(\frac{n_i z_i^2}{T_i} + \frac{n_j z_j^2}{T_j} \right)^{1/2} \right], \quad (8)$$

n_i and n_j are in units of 10^{13} cm^{-3} , T_i and T_j are in units of keV, and ν_i is in units of sec^{-1} . The ionic perturbed density

$\tilde{n}_i = \tilde{n}_i^{\text{nr}} + \tilde{n}_i^{\text{c}} + \tilde{n}_i^{\text{r}}$ is the same as in Ref 1, except with $|e|$ replaced by $z_i |e|$ everywhere, and with the new definitions:

$$\omega_{*i} = -(T_i r_{ne} / z_i T_e r_{ni}) \omega_{*i}, \quad n_i \equiv (d \ln T_i / dr) / (d \ln n_i / dr),$$

$$\bar{\omega}_{di} = (r_{ni} / R_0) \omega_{*i}, \quad \rho_i = \nu_i / \Omega_i, \quad \Omega_i = z_i |e| B / m_i c, \quad \text{and}$$

$\nu_i^* \equiv \nu_i / (\epsilon_0^{3/2} \bar{\omega}_{ti})$. Otherwise all quantities are defined as in Ref. 1.

The mode equation is just the quasi-neutrality equation

$$0 = \tilde{n}_e - \sum_i z_i \tilde{n}_i. \quad (9)$$

This equation is solved to obtain the eigenfrequencies ω and eigenfunctions $\tilde{\phi}(\theta, S)$ just as in Ref. 1.

III. NUMERICAL RESULTS FOR A SINGLE LIGHT IMPURITY

In this section we present numerical results illustrating the effect of a single light impurity species, carbon (C), on the eigenvalue for a deuterium (D)-electron (e) plasma in the banana regime. The parameters used are those of the "basic case" of Ref. 1, namely $\epsilon_0 = 0.1$, $T_e/T_D = 2$, $q = 2.5$, $q'r/q = 1$, $n_1 = n_e = 0$, $r_{ne}/r_0 = 1$, $m_D/m_e = 3672$, and $b_{D0} = 0.2$, (or $k_\theta \rho_D = 0.63$). In addition we specify $T_e = 1$ keV and $R_0 = 132$ cm. The density is chosen to be $n_e = 0.2656 \times 10^{13}$ cm.⁻³, so that $v_e^* = 0.03$ when $n_C = 0$, as in Ref. 1. Then n_e is kept constant as n_C/n_e is varied. Also $T_C = T_D = 0.5$ keV, and $n_C = 0$.

The eigenfunction for $n_C/n_e = 0.03$, or $Z_C n_C/n_e = 0.18$, and $r_{nC}/r_{ne} = 1$ is shown in Fig. 1 for the lowest eigenfunction. It is qualitatively the same as that in Ref. 1 without impurities. Subsequent results in this section will be for this eigenfunction, which has the highest linear growth rate, as different parameters are varied. In the calculation we retain nine poloidal harmonics, nine radial basis functions, and set $\sigma = 0.70 + 0.15i$, as in Ref. 1. The eigenvalue is $\omega/\omega_{*e} = 0.345 + 0.144i$.

In Fig. 2, the variation in ω/ω_{*e} , as n_C/n_e is varied, is shown for $r_{nC}/r_{ne} = 1$, i.e., for radially constant density fractions for the three species. At the left end of Fig. 2, for a

pure deuterium plasma, $v_e^* = 0.03$, while at the right end, for a pure carbon plasma, $v_e^* = 0.1785$. Correspondingly v_D^* increases from $v_D^* = 0.08$ to $v_D^* = 0.60$ and v_C^* increases from $v_C^* = 1.34$ (at $n_C/n_e = 0.01$) to $v_C^* = 12.83$. Thus the electrons and deuterium are always in the banana regime and the carbon is always in the plateau regime. As n_C/n_e varies, n_D/n_e also varies so that the constraint Eq. (2) is always satisfied. Increasing the carbon charge fraction $Z_C n_C/n_e$ is seen in Fig. 2 to increase the real part of the frequency ω_r and to be moderately stabilizing.

In Fig. 3 the variation in ω/ω_{*e} as n_C/n_e is varied is shown for $r_{nC}/r_{ne} = -1$. This could correspond to a situation where the impurity is localized radially at the outside of the discharge. Here r_{nD}/r_{ne} is also varied so that the constraint Eq. (3) is always satisfied. The variation in v_e^* , v_D^* , and v_C^* is the same as for $r_{nC}/r_{ne} = 1$. Increasing the carbon charge fraction $Z_C n_C/n_e$ is seen to decrease ω_r , has little effect on the growth rate γ for $Z_C n_C/n_e < 0.5$, and is stabilizing for $Z_C n_C/n_e > 0.5$. Here several effects are operating simultaneously to give this net effect on γ . The increases in the collisionalities would be stabilizing. The decrease in ω_r/ω_{*e} increases the effect of the magnetic drifts (including their resonances), and can be either stabilizing or destabilizing. The stabilizing

contribution from the ion sound terms and the ion transit resonance terms can change in a complicated fashion as $Z_C n_C / n_e$ and the exact amount of toroidal coupling change.

In Fig. 4 is shown the variation in ω as (r_{ne}/r_{nC}) , (which is proportional to the impurity density gradient), is varied. Increasing (r_{ne}/r_{nC}) causes an increase in ω_r/ω_{*e} . Increasing normal gradients $(r_{ne}/r_{nC} > 0)$ are seen to be slightly stabilizing, and increasing reversed gradients for the carbon $(r_{ne}/r_{nC} < 0)$ are seen to be slightly destabilizing. This sort of behavior was seen previously in a sheared-slab calculation with a light impurity in Ref. 3. In Fig. 4, $n_C/n_e = 0.03$, and n_D/n_e and r_{nD}/r_{ne} are varied so that the constraint Eqs. (2) and (3) are always satisfied.

In Fig. 5 is shown the variation of ω with $\eta_e = \eta_D = \eta_C$, for the case where $r_{nC}/r_{ne} = 1$ and $n_C/n_e = 0.03$. Increasing temperature gradients are seen to be substantially destabilizing. They also decrease ω_r/ω_{*e} , eventually causing a transition from the electron to the ion diamagnetic direction, at $\eta_e = \eta_D = \eta_C = 3.4$.

In Fig. 6 the variation of ω with η_e is shown for the case of a reversed carbon density gradient, $r_{nC}/r_{ne} = -1$, with $r_{Te} = r_{TD} = r_{TC}$, where $r_{Tj} \equiv -(d \ln T_j / dr)^{-1}$. Thus in this case $\eta_e = -\eta_C = 1.44 \eta_D$, with $n_C/n_e = 0.03$ again. Qualitatively the

effects here are the same as in the case of a normal carbon density gradient, as shown in Fig. 5.

In Fig. 7 is shown the variation of ω (in absolute units) with $b_{D0} \equiv k_{\theta}^2 \rho_D^2/2$. Here $n_C/n_e = 0.03$, $r_{nC}/r_{ne} = 1$, and $b_{C0} = b_{D0}/6$. Here ω_x has a maximum for $b_{D0} \approx 0.2$. The growth rate γ increases with increasing b_{D0} up to $b_{D0} \approx 0.8$, and then decreases much more slowly. The circulating ion transit frequencies $\bar{\omega}_{tD}$ and $\bar{\omega}_{tC}$ are indicated in Fig. 7 for comparison.

IV. EFFECT OF BEAM PARTICLES

It is fairly straightforward to include a beam species in the eigenmode calculation in a somewhat approximate fashion. The beam-species equilibrium distribution function is calculated by a multi-group algorithm in the BALDUR transport code.⁷ We consider here a case modeling the neutral-beam-heated PLT, where a deuterium beam is injected into a hydrogen plasma. The beam distribution function F_b , averaged over pitch angle, is shown as a function of the energy E in Fig. 8 for a minor radius of 22 cm. (limiter radius $a = 40$ cm.). The result is shown at the time that the beams turn off, when the ion temperature on the magnetic axis is 7.4 keV.

The beam distribution function F_b is seen to be substantially different from the Maxwellian $N \exp(-E/T_b)$ normalized to have the same number density and average energy as F_b . It

is also substantially different from the distribution⁸ $N'/(v^3 + v_c^3)$ often used to model the beam distribution. The form of $F_b(E)$ can be understood as follows. The beam consists of three components, a full energy component at 37 keV comprising 67% of the beam particles, a half energy component at 18.5 keV comprising 18%, and a third energy component at 12.3 keV comprising 15%. For the current local parameters, as a beam particle scatters down in energy from its injection energy, it is quite likely to be charge-exchanged, and thereby lost. This causes the regions where $\partial F_b / \partial E$ is positive. There is a high-energy tail above the full injection energy due to energy diffusion,⁸ but this tail falls off quickly. In the multi-group calculation, a beam particle whose energy has fallen to (3/2) times the background ion temperature is transferred to a separate (background) thermal deuterium ion population, which is treated like those for the other ion species. Thus $F_b = 0$ for $E \leq 1.76$ keV.

The beam distribution function F_b is shown at three neighboring radii in Fig. 9. It can be seen that the shape of $F_b(E)$ is almost the same at the three radii, with the main difference being in the overall magnitude. Therefore $\partial \ln F_b / \partial r \approx \partial \ln n_b / \partial r$ to a good approximation.

The form of the ion response used in the two-dimensional eigenmode calculation is only slightly different from that described in Ref. 1. For a Maxwellian lowest-order distribution function, Eq. (27) of Ref. 1 for the non-resonant perturbed ion

density contains the factor $f_{mi}(\omega - \omega_{*i}^T)$, where f_{mi} is the Maxwellian equilibrium distribution function. For a non-Maxwellian energy distribution function, this becomes $(-\omega T_b \partial F_b / \partial E - F_b \omega_{*b})$, for $\partial \ln F_b / \partial r = \partial \ln n_b / \partial r$, where T_b is (2/3) times the average energy obtained from $F_b(E)$. A corresponding change is made in the expression for the resonant perturbed ion density.

To model the neutral-beam-heated PLT, it is necessary to include four ion species: hydrogen (H), background deuterium (D), carbon (C), and beam deuterium (b). In the notation of Ref. 1, the equilibrium parameters evaluated at $r = 22$ cm. are: $T_e = 1.047$ keV, $q = 1.375$, $q'r/q = 1.480$, $\epsilon_0 = 0.1689$, $n_e = 2.378 \times 10^{13}$ cm.⁻³, $r_{ne}/r = 1.282$, and $\eta_e = 3.688$. The parameters for the four ion species are given in table I. The values of r_{nC}/r_{ne} and η_C are so large because the ion density profile is almost flat at this radius. For comparison purposes, we will also consider a case with the beam species artificially suppressed. When so doing, n_D/n_e and r_{nD}/r_{ne} are changed so that the constraint Eqs. (2) and (3) will continue to be satisfied.

Results have been calculated for three cases: 1) with no beam species, 2) with the beam species with a Maxwellian distribution function in energy and 3) with the beam species with the non-Maxwellian distribution function shown in Fig. 8. The eigenvalues for the three cases are given in table II, for the

eigenfunction which we presume has the highest growth rate. The eigenfunction is qualitatively similar to that shown in Fig. 1.

All of these cases have $\omega_r/\omega_{*e} < 0$, as is expected for these large values of the η_j 's. The effect of turning on the Maxwellian beam is seen to be substantially stabilizing. Part of this is due to the change in r_{nd}/r_{ne} by more than a factor of two. Another part is due to the replacement of part of the background deuterium density, with $\eta_D \approx 7$, by beam density with $\eta_b = 0$, in a situation where the ion temperature gradient is the strongest destabilizing mechanism. The effect of changing from the Maxwellian to the non-Maxwellian energy distribution is seen to be further stabilizing. This may be due to an enhancement of circulating beam ion transit resonances (Landau damping) by an increase in the number of high energy particles whose transit frequency is comparable to or larger than the mode frequency. In fact, for these parameters, $\bar{\omega}_{tb}/\omega_{*e} = 0.848$ so that $|\omega/\bar{\omega}_{tb}|$ is of order unity, so that circulating beam ion transit resonances have a strong effect.

However, it should be pointed out that the form of the ion response we use is derived in Ref. 1 in the limit $|\omega/\bar{\omega}_{ti}| \gg 1$. Thus this assumption is at the point of breaking down for the beam ions, while it is well satisfied for all of the other ion species. Because of this limitation, the result obtained for the effect of the beam ions should be regarded as more qualitative than quantitative. For this reason, we have not considered

it worthwhile to explore a large number of different cases in parameter space. In future calculations, we hope to employ a more general form of the ion response which will be valid for arbitrary $|\omega/\omega_{tj}|$.

Another limitation of the present calculation is the neglect of the pitch-angle dependence of the beam particle distribution function. This dependence is seen to be quite strong in the multigroup results for PLT referred to earlier. This pitch-angle dependence would introduce new terms into the gyro-kinetic equation solved for the perturbed ion distribution function in Ref. 1. We hope to include this pitch-angle dependence in future calculations.

V. CONCLUSIONS

We have investigated the effects of multiple ion species on trapped-electron modes. As a specific example, we have considered the case of a single light impurity, carbon, in a deuterium plasma for parameters such that the electrons and the deuterium ions are in the banana regime, and the carbon ions are in the plateau regime. For a fixed, moderate carbon density fraction, an increasing carbon density gradient in the same direction as the electron and deuterium density gradients is slightly stabilizing. Conversely, an increasing carbon density gradient in the opposite direction from the electron and deuterium gradients is slightly destabilizing. However, for a fixed carbon

density gradient of either sign, the effect of increasing the carbon density fraction was always slightly stabilizing, due to a combination of effects.

We have also investigated the effects of a beam species, using parameters appropriate to the neutral-beam-heated PLT. The beam ion distribution function is strongly non-Maxwellian in energy in this case. The effect of turning on a Maxwellian beam species is substantially stabilizing. The effect of changing from a Maxwellian to the non-Maxwellian energy distribution for the beam ions is further stabilizing. These results for the effects of the beam species are rather approximate for two reasons. First, the strong pitch-angle dependence of the beam distribution function has been neglected. Second, the form of the ion response used is accurate only in the limit where the mode frequency is large compared to the circulating-ion transit frequency. This condition is well satisfied for the other ion species, but the two frequencies are comparable for the beam species. It is hoped that both of these approximations can be removed in a future calculation.

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Table I: Ionic parameters for the neutral-beam-heated PLT.

j	b_{j0}	n_j/n_e	z_j	m_j/m_H	r_{nj}/r_{ne}	T_j (keV)	η_j
H	0.6	0.245	1	1	0.972	1.246	4.517
D	1.2	0.494 ^a	1	2	1.544 ^a	1.246	7.173
C	0.2	0.032	6	12	-75.78	1.246	-352.15
p	14.7	0.069	1	2	0.160	15.293	0

^a $n_D/n_e = 0.563$ and $r_{nD}/r_{ne} = 0.750$ when the beam species is artificially suppressed.

Table II: Eigenfrequencies for the neutral-beam-heated PLT.

Case	ω/ω_{*e}
No beam	$-1.456 + 0.467i$
Maxwellian beam	$-0.902 + 0.280i$
Non-Maxwellian beam	$-0.827 + 0.188i$

Fig. 1. Real and imaginary parts of $\tilde{\phi}(\theta, S)$, the perturbed electrostatic potential eigenfunction, for $n_0/n_e = 0.03$, $r_{n0}/r_{ne} = 1$, with the other parameters as given in the text. Here $\theta = 0$ corresponds to the outside of the torus and $S \approx (r - r_0)/\Delta r_s$, where r_0 is the radius of the reference mode-rational surface and Δr_s is the separation of mode-rational surfaces.

792400

Fig. 2. Variation of the eigenvalue ω with the carbon impurity charge fraction $Z_C n_C / n_e$ for a normal impurity density gradient. 792402

Fig. 3. Variation of the eigenvalue ω with the carbon impurity charge fraction $Z_C n_C / n_e$ for a reversed impurity density gradient.

792403

Fig. 4. Variation of the eigenvalue ω with the carbon density gradient, for $n_c/n_e = 0.03$. 792404

792401
 Fig. 5. Variation of the eigenvalue ω with the temperature gradient parameter $\eta_j = (d \ln T_j / dr) / (d \ln n_j / dr)$ for a normal carbon density gradient, with $n_c/n_e = 0.03$.

792399
Fig. 6. Variation of the eigenvalue ω with $\eta_e = -\eta_c = 1.44\eta_D$ for a reversed carbon density gradient, for $n_c/n_e = 0.03$ and $r_{Te} = r_{TD} = r_{TC}$, where $r_{Tj} = -(d \ln T_j / dr)^{-1}$.

Fig. 7. Variation of the eigenvalue ω with $b_{D\theta} \equiv k_0^2 \rho_D^2 / 2$, for $n_c/n_e = 0.03$ and $r_{nc}/r_{ne} = 1$. 792398

802050
 Fig. 8. Beam particle distribution function F_b as a function of the energy E . Also shown are a Maxwellian, $N \exp(-E/T_b)$, with the same number density and average energy, and a previous simplified beam distribution function, $N'/(v^3 + v_c^3)$, with the same number density. Here F_b is calculated for the neutral-beam-heated PLT at $r = 22.29$ cm.

802051

Fig. 9. Beam distribution function $F_b(E)$ at three neighboring radii.