Toroidal Microinstability Studies of High Temperature Tokamaks

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ABSTRACT

Results from comprehensive kinetic microinstability calculations are presented showing the effects of toroidicity on the ion temperature gradient mode and its relationship to the trapped-electron mode in high-temperature tokamak plasmas. The corresponding particle and energy fluxes have also been computed. It is found that, although drift-type microinstabilities persist over a wide range of values of the ion temperature gradient parameter $\eta_i \equiv (d \ln T_i / dr) / (d \ln n_i / dr)$, the characteristic features of the dominant mode are those of the η_i -type instability when $\eta_i > \eta_{ic} \sim 1.2$ to 1.4 and of the trapped-electron mode when $\eta_i < \eta_{ic}$.

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I. INTRODUCTION

There have been a number of papers recently concerned with the ion temperature gradient (η_i) mode in the linear and nonlinear regimes.¹⁻⁷ Here, $\eta_j \equiv (d \ln T_j / dr) / (d \ln n_j / dr)$. However, much of this work simplifies the geometry to that of an unsheared or sheared slab, or approximates the electron response as adiabatic, or employs purely fluid equations. Any of these approximations removes the effect of the trapped electrons in toroidal geometry. In the present work, the goals are to show the effect of toroidicity on the usual η_i mode, and to show the relationship between the η_i mode and the usual trapped-electron mode. Toroidal effects enter in a number of different ways. In addition to the trapped-electron effects, the toroidal curvature causes ion and electron magnetic drifts, which can make both resonant and nonresonant contributions to the eigenfrequency. This effect has been studied before in Ref. 8 in an approximate calculation. Also, the ballooning of the eigenfunction to the weak magnetic field side of the tokamak is a consequence of the greater unfavorable curvature effects as well as of the increased number of trapped particles in this region.

In the very simplest analytic calculation of the η_i -mode eigenfrequency ω in unsheared slab geometry, it is assumed that $k_{\parallel} v_i \ll \omega \ll k_{\parallel} v_e$, where $k_{\parallel}$ is the wave number along the magnetic field line and $v_j \equiv (2T_j/m_j)^{1/2}$. Taking $\eta_i \gg 1$ and $k_{\perp} \rho_i \ll 1$, where $k_{\perp}$ is the wave number perpendicular to the magnetic field line and ρ_i is the average ion gyroradius, and assuming the electron response to be adiabatic, the quasineutrality condition implies that $\omega^3 \simeq -k_{\parallel}^2 v_i^2 (T_e/2T_i) \omega_{ci} \eta_i$, where ω_{ci} is the ion diamagnetic drift frequency. The unstable root is the η_i mode while the marginally stable root becomes the electron drift mode, so that the two roots are distinct. However, in toroidal geometry, where $k_{\parallel} v_j$ is replaced by $\bar{\omega}_{tj}$, the average transit frequency for untrapped particles of species j , it will be shown that, in realistic cases, the mode frequency can be such that $|\omega| \lesssim \bar{\omega}_{ti} \ll \bar{\omega}_{te}$. This implies that the untrapped ion transit frequency resonances (ion Landau damping) will be strong, but it also implies that an expansion in $|\omega|/\bar{\omega}_{ti} \ll 1$, yielding a distinct ion sound term, is no longer possible. Thus, the mode equation remains transcendental and the number of roots is

indefinite. This suggests the possibility that the electron drift mode root and the ion sound (η_i) mode root may be combined together into a single, hybrid root. It will be shown in the results here that this can be the case.

A relatively complete computational model for tokamak linear eigenmodes and their associated quasilinear transport is employed here. The equations solved and the solution methods have previously been presented in great detail in Refs. 9-11. This material will not be repeated here, because the motivation behind the present paper is the application of such procedures to address the important transport question regarding the relationship between the η_i -type mode and the trapped-electron instability. The calculation is essentially an exact solution of the linearized gyrokinetic equation for large n (toroidal mode number), employing the ballooning formalism. It allows for arbitrary ratios between the mode frequency, ω , the transit frequency for untrapped particles, ω_t , the bounce frequency for trapped particles, ω_b , and the magnetic (gradient and curvature) drift frequency, ω_d . Trapped and untrapped particles for every species are included, and full finite Larmor radius effects are retained. The model collision operator employed should be relatively accurate in the banana regime and the lower end of the plateau regime, as described in Ref. 10. Input data for the calculation, such as radial density and temperature profiles, are taken from results of BALDUF transport code¹² runs. The corresponding pressure and safety factor profiles are used in calculating the numerical MHD equilibria employed in the instability calculation. For illustrative purposes, all instability results presented here are calculated in the electrostatic limit. As pointed out in Refs. 9 and 11, inclusion of finite- β modifications in these computations is quite straightforward.

The most directly relevant previous calculation is that in Ref. 1, where the effects of an approximate trapped-electron time-average term on an ion temperature gradient mode are considered. However, this calculation makes several significant approximations. In particular, all bounce and transit frequency effects are omitted, so that the ion temperature gradient mode considered is not the one mentioned earlier coming from the ion sound term. Instead,

it is a related instability coming mainly from the toroidal ion magnetic drift frequency, in combination with η_i . As will be shown, for realistic parameters, the relevant instabilities can be in a regime where $|\omega| \lesssim \omega_{Ti}$, so that finite ion transit frequency effects are critical. Also, Ref. 1 neglects the variation of the eigenfunction along the magnetic field line, which is needed for quantitative accuracy. Nevertheless, the results here have some qualitative features in common with those of Ref. 1.

Numerical results are presented in Sec. II, and conclusions are presented in Sec. III.

II. Numerical Results

Results will first be presented for a case¹³ corresponding to a (now superseded) design for the Compact Ignition Tokamak (CIT), with vacuum major radius $R_0 = 1.75$ m, toroidal magnetic field $B_0 = 10$ T, and plasma current $I_p = 9$ MA. This is the same case employed in Ref. 14, but now without any α -particles or helium ash particles, and with the eigenmode calculation performed in the electrostatic limit. Also, the numerical MHD equilibrium is recomputed for decreased equilibrium pressure, but with the same profile shape. The local parameters on the chosen magnetic surface with average minor radius $r \simeq 1/2$ are $\beta = 0.166\%$, $q = 1.17$, $\hat{s} \equiv r(dq/dr)/q = 0.456$, $\epsilon_0 \equiv r/R_0 = 0.154$, $R_0 = 180$ cm, $\nu_e^* \equiv$ effective electron collision frequency/average trapped electron bounce frequency $= 0.0035$, $n_e = n_i = 3.71 \times 10^{14} \text{ cm}^{-3}$, $m_i = 2.5$ amu, $T_e = 18.9$ keV, $T_i = 19.4$ keV, $n = 58$, and $k_\theta \rho_i = 0.400$. Results will be presented as η_i and η_e are varied from their BALDUR values, $\eta_i = 1.34$ and $\eta_e = 1.28$. For this case, as η_i is varied, η_e varies in proportion, so that always $\eta_e/\eta_i = 0.955$. Also, as η_i varies, the total pressure gradient is kept fixed. This means, for instance, that as η_i increases, so that dT_i/dr and dT_e/dr increase, $dn_i/dr = dn_e/dr$ decreases so as to maintain a fixed total pressure gradient. In this way, the parameters remain consistent with the MHD equilibrium, which results from a specific pressure gradient on the chosen magnetic surface.

Results for the growth rate, γ , and the real frequency, ω_r , in units of the electron diamagnetic drift frequency, ω_{ce} , are shown in Fig. 1 for this case as η_i and η_e are varied. It can be seen from the growth rate curve in Fig. 1a for the actual collisionality value, $\nu_e^* = 0.0035$, that the instability has two regimes: for $\eta_i \gtrsim 1.4$, the growth rate increases monotonically, and almost linearly, with η_i . In this regime, the fluid η_i -mechanism, which is responsible for destabilizing the η_i mode in the simple slab geometry limit mentioned earlier, destabilizes this toroidal mode. However, as η_i decreases from 1.4, the growth rate increases again slightly, before falling to zero for $\eta_i \simeq -0.4$. In this regime, the trapped-electron time-average magnetic drift (precession) frequency resonance mechanism is dominant. It is this mechanism which is responsible for the usual collisionless trapped-electron mode. In this case, a single root of the mode equation combines the usual collisionless trapped-electron mode and the usual η_i mode, and will be referred to as the trapped-electron- η_i mode.

The effects of the trapped electrons can be greatly reduced by increasing the collisionality, since the trapped-electron time-average term varies as $1/\nu_e^*$ for ν_e^* sufficiently large. If ν_e^* is artificially increased to 1.0 by increasing $n_e = n_i$ to $1.28 \times 10^{17} \text{ cm}^{-3}$, the results in Fig. 1a show that the regime destabilized by the trapped-electron magnetic drift resonance mechanism is eliminated, and the mode is damped for $\eta_i \lesssim 1.1$. The corresponding real frequencies, shown in Fig. 1b, start out in the electron diamagnetic direction at small η_i and make a transition to the ion diamagnetic direction as η_i increases past 1.1.

The dependence of γ and ω_r on $k_\theta \rho_i \propto n$ for the mode is shown in Fig. 2 for the nominal BALDUR values of $\eta_i = 1.34$ and $\eta_e = 1.28$. The value $k_\theta \rho_i = 0.400$ chosen for Fig. 1 is seen to be close to the maximum of γ for both collisionalities. Increasing $k_\theta \rho_i$ is also seen to push the real frequency in the ion diamagnetic direction.

The results shown in Fig. 1 correspond to a single root resulting from the coupling of the collisionless trapped electron mode and the η_i mode, with an eigenfunction that is essentially a simple gaussian along the magnetic field line, well localized around the magnetic field minimum. Normally, this is the most unstable root. There are other roots as well, some

with lower growth rates, and some that are damped. In fact, since the eigenvalue equation is transcendental in character, there are an indeterminate number of possible roots.

The other case to be considered here corresponds to a chosen time in a BALDUR transport code run¹⁵ modeling a "supershot" (shot 22014) in the Tokamak Fusion Test Reactor¹⁶ (TFTR) with $R_0 = 2.55$ m, $B_0 = 4.5$ T, and $I_p = 0.8$ MA. To make the calculation more realistic, a beam species (b) from neutral beam injection heating is included in the calculation, along with the background deuterium ions (i) and the electrons (e). The equilibrium distribution function $F_b(E, v_{\parallel}/v)$ for the beam species as calculated by the BALDUR code for this case with balanced co- and counter parallel neutral beam injection is shown in Fig. 3. Here, E is the particle energy and $v_{\parallel}/v$ is a dimensionless pitch-angle variable, with $v_{\parallel}$ being the particle velocity component along the magnetic field, evaluated when the particle is at the magnetic field minimum. The results for the trapped-electron- η_i mode for this strongly non-Maxwellian distribution will be compared with those for a Maxwellian distribution with the same number density and the same average energy, and to results with no beam species. In this last case, the density and density gradient of the background ions are adjusted to maintain the condition of bulk neutrality and its radial derivative. The local parameters on the chosen magnetic surface with $r/a \simeq 1/2$ are $\beta = 1.66\%$, $q = 1.15$, $\bar{s} = 0.456$, $\epsilon_0 = 0.068$, $\nu_e^* = 0$, $n_e = 5.54 \times 10^{13} \text{ cm}^{-3}$, $n_i/n_e = 0.901$, $n_b/n_e = 0.099$, $T_i = 12.8$ keV, $T_b = 32.9$ keV for the Maxwellian distribution, $T_e = 4.68$ keV, $\eta_b = 2.65$ for the Maxwellian distribution, $\eta_e = 1.15$, $r_{ni}/r_{ne} = 0.914$, and $r_{nb}/r_{ne} = 6.58$, where $r_{nj} \equiv -(d \ln n_j / dr)^{-1}$. The MHD equilibrium is calculated numerically with a fixed, circular boundary. The instability calculation is again performed in the electrostatic limit, for $n = 13$ or $k_y \rho_i = 0.356$.

Results for the trapped-electron- η_i mode growth rate, γ , and real frequency, ω_r , are shown in Fig. 4. In this case, η_e and η_b and the total pressure gradient are kept fixed as η_i is varied. Looking first at the growth rate curve for the case with no beam species, it can be seen that, for $\eta_i > 1.2$, the growth rate increases monotonically and almost linearly with η_i . Again, this is the regime where the η_i -mechanism is dominant. For η_i decreasing

from 1.2, the growth rate increases again, peaking around $\eta_i = 0$, and falls to zero only for $\eta_i \simeq -0.7$. Again, this is the regime dominated by the trapped-electron time-average magnetic drift (precession) frequency resonance mechanism. Addition of the beam species with the Maxwellian distribution increases the growth rate moderately for all $\eta_i > 0$. With the non-Maxwellian distribution, the increase is almost doubled, but is still moderate. The corresponding real frequencies are shown in Fig. 1b. Again, there is a transition from the electron diamagnetic direction to the ion diamagnetic direction as η_i increases. The presence of the beam species has very little effect on the real frequency. Note that, over the entire range of η_i , $|\omega| \lesssim \bar{\omega}_{hi}$, the average trapped-ion bounce frequency. The common approximation, that $|\omega| \gg \bar{\omega}_{hi}$, is clearly invalid in this realistic case.

The perturbed electrostatic potential eigenfunctions, $\tilde{\Phi}(\theta)$, where θ is the extended poloidal angle in the ballooning representation (ballooning coordinate along the field line), are shown in Fig. 5 for the cases with no beam species for $\eta_i = 0$ and $\eta_i = 3$. For $\eta_i = 0$, the eigenfunction is well localized around $\theta = 0$, the magnetic field minimum, where the trapped electrons are localized and where there is "bad" magnetic curvature. For $\eta_i = 3$, the mode is still well localized, but the width is somewhat greater. This is allowed because the η_i -mechanism is not dependent on trapped particles or magnetic drifts, so the most unstable eigenfunction spreads out somewhat to decrease the stabilizing effects of the ion transit frequency resonances (ion Landau damping) and of the parallel energy convection (that corresponds to the ion sound term in the case of large $|\omega|/\bar{\omega}_{hi}$).

The dependence of the growth rate and real frequency on $k_\theta \rho_i \propto n$ is shown in Fig. 6 for this case with no beam species and with the non-Maxwellian beam species, for $\eta_i = 0.51$ or 0.56 and for $\eta_i = 2$. In this case, the growth rate peaks for $k_\theta \rho_i \simeq 0.6$ when $\eta_i = 2$ and for $k_\theta \rho_i \simeq 1.0$ when $\eta_i = 0.51$ or 0.56. This difference is as expected, since finite Larmor radius effects are strongly stabilizing for the η_i mode in the simple slab analytic limit mentioned earlier, whereas the collisionless trapped-electron mode growth rate typically peaks for $k_\theta \rho_i$

of order unity. Again, the real frequency is pushed in the ion diamagnetic direction by increasing $k_\parallel \rho_i$.

The quasilinear fluxes of ions and electrons, Γ , of electron energy, Q_e , and of ion energy, Q_i , can be calculated as described in Refs. 10 and 11. Here, Q_i is the *total* energy flux, including the convective part. A purely quasilinear calculation does not determine the saturation level of the instability, and hence the absolute level of anomalous transport, but it does specifically determine the ratios of fluxes, *e.g.*, $T_i \Gamma / Q_i$ and Q_e / Q_i . These ratios are shown for this ITFR case with no beam species in Fig. 7. For $\eta_i < 1$, $T_i \Gamma \sim Q_e \sim Q_i$, while for $\eta_i > 1$, $Q_i \gg T_i \Gamma \sim Q_e$. This difference is understandable in terms of the two instability mechanisms involved. If there were no electron dissipation in the eigenmode equation, (*i.e.*, if the electron response were purely adiabatic with perturbed electron distribution function $\tilde{f}_e = (e\tilde{\Phi}/T_e)F_{me}$), then the perturbed electrostatic potential $\tilde{\Phi}$, the perturbed electron density $\tilde{n}_e = (e\tilde{\Phi}/T_e)n_e$, the perturbed ion density $\tilde{n}_i$ ($= \tilde{n}_e$ by the quasineutrality condition), and the perturbed electron pressure $\tilde{p}_e = (e\tilde{\Phi}/T_e)n_e T_e$ would all be exactly in phase. Then $\Gamma = \Gamma_i = \Gamma_e \propto \langle \tilde{v}_{E\tau} \tilde{n}_e \rangle \propto \text{Im}(\tilde{\Phi}^* \tilde{n}_e)$ and $Q_e \propto \langle \tilde{v}_{E\tau} \tilde{p}_e \rangle \propto \text{Im}(\tilde{\Phi}^* \tilde{p}_e)$ would be identically zero. Here, $\tilde{v}_{E\tau}$ is the radial component of the perturbed $\mathbf{E} \times \mathbf{B}$ drift velocity. However, the perturbed ion pressure $\tilde{p}_i$ can be out of phase with the perturbed electrostatic potential, since $\tilde{p}_i \neq \tilde{p}_e$, so $Q_i \propto \langle \tilde{v}_{E\tau} \tilde{p}_i \rangle \propto \text{Im}(\tilde{\Phi}^* \tilde{p}_i)$ can be nonzero, due to the destabilization from the η_i -mechanism. If some small amount of electron dissipation were then added, mainly from the trapped-electron magnetic drift frequency resonance in the present instance, a correspondingly small amount of particle flux and electron energy flux would result. The ion energy flux, however, coming from the η_i -mechanism, would be much larger. This is the situation when $\eta_i > 1$. For $\eta_i < 1$, on the other hand, the trapped-electron magnetic drift frequency resonance mechanism is dominant. Since it gives nonzero contributions to all three fluxes, it is not surprising that they would be roughly comparable.

III. CONCLUSIONS

In both of the realistic cases presented here, the usual η_i mode was seen to connect to the usual collisionless trapped-electron mode as η_i decreased. In an academic sense, this effectively removes the threshold value of η_i for instability, as noted previously in Ref. 1. However, in a more realistic sense, the physics *does indeed change* around the threshold, η_{ic} . Specifically, the dominant destabilization mechanism and the dominant form of anomalous transport do change at $\eta_{ic} \simeq 1.2$ to 1.4 . For $\eta_i > \eta_{ic}$, the usual fluid η_i -mechanism for instability is dominant and the quasilinear ion energy flux greatly exceeds the electron energy flux or the particle flux. On the other hand, for $\eta_i < \eta_{ic}$, the trapped-electron time-average magnetic drift (precession) frequency resonance mechanism for instability is dominant, and all three fluxes are comparable. In future work, it will be of considerable interest to compare the results of this type of linear and quasilinear microinstability study with corresponding experimental results in high-temperature tokamaks.

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FIGURE CAPTIONS

- FIG. 1. a) Growth rate γ and b) real frequency ω_r versus η_i , in units of the electron diamagnetic drift frequency, ω_{*e} , for the CIT case. Here, $\eta_e/\eta_i = 0.955$ and the total pressure gradient are kept fixed as η_i varies.
- FIG. 2. a) Growth rate γ and b) real frequency ω_r versus $k_\theta \rho_i \propto n$, in units of 10^5 sec^{-1} , for the CIT case.
- FIG. 3. Non-Maxwellian equilibrium distribution function F_b for the hot neutral beam species versus energy E and dimensionless pitch-angle variable $v_{||}/v$, for the TFTR case.
- FIG. 4. a) Growth rate γ and b) real frequency ω_r versus η_i , in units of 10^5 sec^{-1} , for the TFTR case. Here, $\eta_e = 1.15$ and the total pressure gradient are kept fixed as η_i varies.
- FIG. 5. Electrostatic potential eigenfunction $\tilde{\Phi}(\theta)$ versus θ , the extended poloidal angle (ballooning coordinate along the magnetic field line) for the TFTR case for a) $\eta_i = 0$ and b) $\eta_i = 3$.
- FIG. 6. a) Growth rate γ and b) real frequency ω_r versus $k_\theta \rho_i \propto n$, in units of 10^5 sec^{-1} , for the TFTR case.
- FIG. 7. Ratios of quasilinear fluxes for the TFTR case with no beam species versus η_i . Here, Γ is the particle flux of electrons and ions, and Q_j is the total energy flux of species j , including the convective part. Also, $\eta_e = 1.15$ and the total pressure gradient are kept fixed as η_i varies.

Fig. 1

Fig. 2

Fig. 4

Fig. 5

Fig. 6

Fig. 7

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