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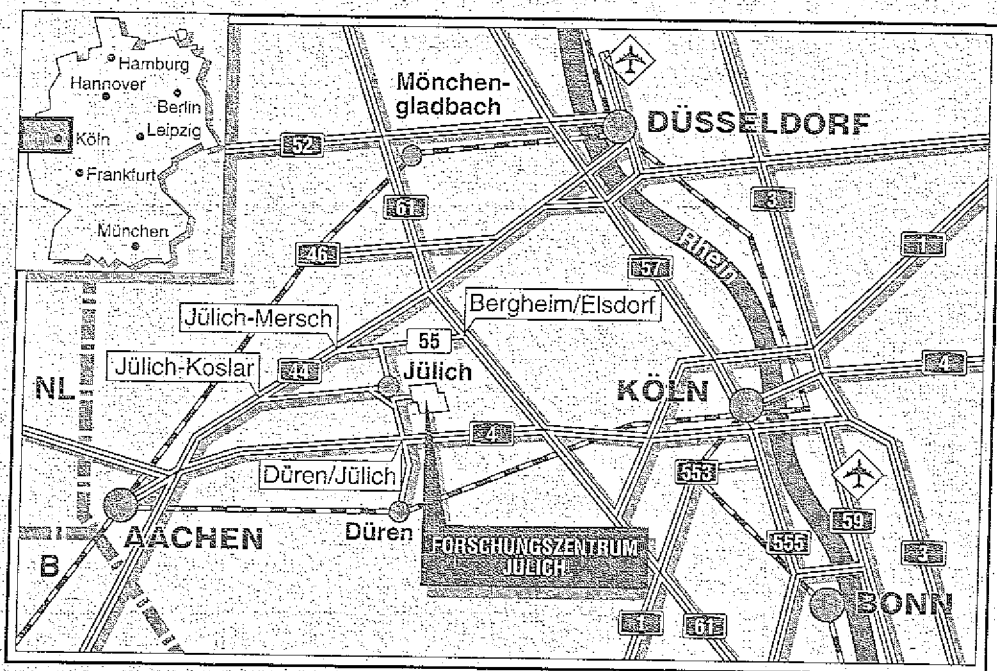

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Effect of Finite-β on the Ion-Temperature-Gradient-Driven Instability in Toroidal Geometry

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Abstract

A set of coupled differential equations is derived to study the finite-β effects on the ion-temperature-gradient driven (ITGD) modes for a collisionless plasma in a toroidal magnetic field. β is assumed low enough to allow ignoring the compressional component of the perturbed magnetic field. The equations will be solved analytically in various limits.

I. Introduction

It is usually considered that low frequency turbulence in magnetic confinement systems, including both electrostatic and electromagnetic modes, is responsible for the enhanced energy and particle losses usually observed in many experiments. To explain such anomalous losses one needs first to study carefully the driving mechanisms responsible for their excitation, and it has been known for many years that the presence of inhomogeneities in plasma equilibrium parameters such as the density and temperature can drive these fluctuations[1]. The well-known electrostatic electron drift wave travels in the direction of electron diamagnetic drift and grows through the resonant interaction of the wave with the electron motion along the field lines. It was shown that the smallest amount of magnetic shear can stabilize the drift wave in a collisionless plasma confined by a cylindrical geometry[2,3]. Stability will still be maintained when β, the ratio of kinetic pressure to magnetic pressure, is increased[4]. The effect of shear can be strongly reduced in a toroidal geometry due to the coupling between the various poloidal Fourier components of the perturbation[5,6,7]. While magnetic shear produces stability in a cylindrical geometry through an outward convection of wave energy, toroidal coupling can result in a radial trapping of the wave energy, thus allowing any destabilizing mechanism to dominate. This applies to the resonance between the wave and circulating electron motion along the magnetic field lines and between the wave and trapped electron magnetic drift in the poloidal direction, in addition to the detrapping electron collisions[8,9,10].

Not only is the presence of an unstable electron drift wave to be expected, but it is also possible to have an unstable lower frequency mode, related to the ion sound wave[11]. Such a wave is electrostatic in the zero-β limit and can be destabilized by an ion temperature gradient, magnetic curvature or by gradients in the parallel or poloidal flow velocity. The shear Alfvén wave, which is electromagnetic in nature ($E_\parallel = 0$) can also be destabilized through magnetic curvature, leading to the well-known β-limit of the magnetohydrodynamic modes which will be localized poloidally in the bad curvature region (to the outside of torus)[12].

It is obvious that the stability analysis of a plasma confined by a helical magnetic field is one-dimensional because of the axial and azimuthal symmetries. The mode frequencies can be obtained by solving a system of integral equations in the general case[13]. Only when the ion Larmor radius is much smaller than the expected effective radial wave length can we approximate these equations by simpler ordinary differential equations. One may still have to resort to a numerical analysis to study the influence of an arbitrary ratio of ion Larmor radius to the azimuthal wave length and of the wave-particle resonance ($\omega = \underline{K} \cdot \overline{\underline{V}}$) where ω is the mode frequency, $\underline{K}$ the wave vector and $\overline{\underline{V}}$ is the Larmor-averaged particle velocity, including both the motion along the field lines and the magnetic drift motion. If the expected frequency is much smaller or much larger that than Doppler shift $\underline{K} \cdot \overline{\underline{V}}$, one can again evaluate the velocity space integrations approximately and the resulting system of ordinary differential equations can sometimes be solved[14]. For a toroidal system, the stability analysis is further complicated by the two-dimensional nature of the problem, the particle magnetic drift which has a component normal to the flux surface and by the presence of a trapped particle component with a general ratio of parallel excursion to parallel wave number. A toroidal analysis can be again simplified by ignoring trapped particles and employing the ballooning formalism, which transforms the two-dimensional eigenvalue problem into a one-dimensional problem along a stretched poloidal angle variable[15]. Further simplification results in the strong coupling limit[16].

The ion-temperature-gradient-driven (ITGD) instability is of particular interest to the study of anomalous ion energy losses in tokamaks. It was discovered in the sixties and remained a subject of interest ever since. This interest was partly due to the fact that it proved not possible to stabilize this mode by using magnetic shear as this results merely in localizing the mode radially and not in the stabilizing effect of outward energy convection[13]. It is also one matter to study the instability in the special cases when the fluid limit of the kinetic analysis is valid, and another matter to study the threshold for instability, close to which the fluid expansion breaks down. Besides, it is not generally allowed to make the simplifying assumption of a long radial wave length and the integral approach is then required and a numerical analysis is usually employed. The differential approach has been frequently used to study the slab limit both analytically (ignoring resonance effects) and numerically when magnetic gradient drifts are ignored[13,17-20]. Some studies also allowed for the presence of a nonuniform plasma flow which has a strong effect on the stability properties[21-23]. As mentioned above, the toroidal case can be analytically treated in the strong toroidal coupling limit when the spacing between the rational surfaces is much shorter than the spacing between the radial turning points in the slab limit solution[7]. The opposite limit of weak toroidicity, where the mode is stretched along the field lines with no significant ballooning, can also be treated analytically, and the effect of ion transit resonance can be accounted for[24,25]. The weak coupling limit requires low magnetic shear or long poloidal wavelength. Various analytic and numerical treatments of the toroidal problem have been produced, where different limits are studied[26-36]. Effects of trapped electrons[34,35], trapped ions[30,36] and an energetic particle component were also considered. Some calculations dealt with flat density H-mode plasmas[28,29,34]. In some calculations, it was preferred to treat the important velocity space wave-particle resonance in detail using the local approximation which ignores the spatial variation of the mode[37-42].

The above discussion applied to the electrostatic zero-β limit. As β is increased gradually, the ITGD mode acquires first a finite radial magnetic field component which can affect the stability properties[43-46]. As β is increased further, one can no longer ignore both the modifications in plasma equilibrium and the compressional magnetic component of the perturbation along the magnetic field lines[47-49]. The two effects are usually opposite in the fluid limit resulting in a significantly reduced effect for the compressional magnetic component[1].

It should finally be mentioned that the stability properties of the ITGD mode are complicated by the large number of parameters which can influence the mode dispersion properties such as degree of steepness of both the density and temperature profiles, the amount of magnetic shear, the amount of toroidicity, the safety factor, the ratio of the ion Larmor ratio of the wave length and the value of β. It is thus not to be expected that simple generalized stability properties can be obtained.

It is our purpose in the present work to treat the effect of a finite but small-β on the ITGD mode in a toroidal geometry using the fluid approximation. An expression for the β required for stability will be obtained.

II. Theoretical Model and Eigenvalue Equation

We will assume that the perturbed electric and magnetic field components can be expressed as:

$$\phi(r,\theta,\xi) = \sum_j \hat{\phi}_j(r) e^{i(m_0\theta + j^\theta - n\xi - \omega t)}$$

where $|j| \ll m_0$. r_0 is taken as the reference surface with $m_0 = nq(r_0)$ with q being the safety factor $= rB_\xi/RB_\theta$ and where circular flux surfaces are assumed. The compressional magnetic perturbation will be ignored. To obtain the perturbed ion distribution function we use the gyrokinetic equation[50],

$$\left\{iv_\parallel \frac{\partial}{\partial s} + (\omega - \omega_{Di})\right\} g_i = \frac{ieF_i}{R_\theta T_i}(\omega - \omega_{*T,i})\{E_\theta + v_\parallel B_r\} J_0 \tag{1}$$

where g_i is the non-adiabatic part of the distribution,

$$f_i = \frac{e_i E_\theta}{iT_i k_\theta} F_j + g_i \exp(iL_i) \tag{2}$$

where F_j is an isotropic Maxwellian distribution of temperature T_i and density N. $L_i = \underline{k} \cdot \underline{V} \times \hat{b}/\Omega_i$

where $\hat{b}$ is the unit vector along the magnetic field lines and Ω_i is the ion gyrofrequency = eB_T/m_i.

$$\omega_{D_j} = k\cdot\hat{b}\times(\mu\nabla B + v_\parallel^2\hat{b}\cdot\nabla\hat{b})/\Omega_j \quad \begin{matrix} < 0 \textit{ for ions} \\ > 0 \textit{ for electrons} \end{matrix}$$

$$= \frac{-k_\theta}{\Omega_j}(v_\parallel^2 + v_\perp^2/2)\left(\cos\theta + \frac{i}{k_\theta}\frac{\partial}{\partial r}\right) \quad \textit{is an operator}$$

$$\textit{and } \omega_{*T,j} = \omega_{*,j}\left[1 - \frac{3}{2}\eta_j + \eta_j\frac{mV^2}{2T}\right] \textit{ and } \eta_j = \frac{d\ln T_j}{d\ln N_j} ,$$

$$\textit{and } \omega_{*,j} = \frac{T_j}{e_j B}k\cdot\hat{b}\times\nabla(\ln N) \quad \begin{matrix} < 0 \textit{ for ions} \\ > 0 \textit{ for electrons} \end{matrix}$$

in the case of radially decreasing density.

$J_0 = J_0\left(\frac{k_\perp v_\perp}{\Omega_i}\right)$ is the zeroth order Bessel function and $k_\perp^2 = k_r^2 + k_\theta^2$.

For the electron component we can use the drift kinetic equation which is similar to the ion equation except that $L_e = 0$ and $J_0(kV_\perp/\Omega_e) = 1$.

The perturbed particle distribution function can then be obtained,

$$f_j = \frac{e_j F_j}{ik_\theta T_j}\left[E_\theta - \frac{\omega - \omega_{*T}}{\omega - K_\parallel V_\parallel - \omega_D}\{E_\theta + V_\parallel B_r\}J_0\right] \tag{3}$$

where $K_\parallel = \frac{K\cdot B}{B} = \frac{j-s}{qR}$ where $s = \frac{r-r_0}{\Delta r_s}$, Δr_s is the radial spacing between two consecutive rational surfaces $= \frac{1}{k_\theta\hat{s}}$, $\hat{s}$ is the shear parameter $\frac{r\partial q/\partial r}{q}$. We may note that the effective shear length is given by $L_s^{-1} = -\hat{s}/qR.$. The quasineutrality condition can be written as

$$\tau\left[NE_\theta - \int d^3v\frac{(\omega - \omega_{*T,i})F_i}{\omega - K_\parallel V_\parallel - \omega_{D,i}}\{E_\theta + V_\parallel B_r\}J_0^2\right]$$

$$= -\left[NE_\theta - \int d^3v \frac{(\omega-\omega_{*T,e})F_e}{\omega - K_\parallel V_\parallel - \omega_{D,e}}\{E_\theta + V_\parallel B_r\}\right] . \tag{4}$$

The parallel component of Ampere's law can be written as

$$B_r'' - \frac{m^2}{r^2}B_r = -\mu_0 i k_\theta J_\parallel = -ik_\theta \mu_0 \sum e_j \int d^3v f_j v_\parallel J_0 \tag{5}$$

Assuming $|\omega_{D,i}/\omega|$ *and* $|K_\parallel V_{+h,i}/\omega| \ll 1$ *and* $|K_\parallel V_{+h,e}/\omega| \gg 1$ we may rewrite Eqs. 4 and 5 in the limit of small ion Larmor radius in the form

$$\left[E_\theta - \frac{\omega-\omega_{*i}}{\omega}E_\theta + \frac{\omega-\omega_{*i}(1+\eta_i)}{\omega}\frac{K_\parallel^2 V_{\parallel,i}^2}{\omega^2}\left(E_\theta + \frac{\omega}{K_\parallel}B_r\right) + \left\{\frac{\omega-\omega_{*i}(1+\eta_i)}{\omega}\frac{\omega_{*i}}{\omega}2\epsilon_n\right.\right.$$

$$\left.\left. - \frac{k_\perp^2\rho_i^2}{2}\frac{\omega-\omega_{*i}(1+\eta_i)}{\omega}\right\}E_\theta\right] = -E_\theta - \frac{\omega-\omega_{*i}}{K_\parallel}B_r \tag{6}$$

$$B_r'' - \frac{m^2}{r^2}B_r = \frac{-\mu_0 Ne^2}{T_e}\left[\frac{(\omega+\omega_{*i})}{K_\parallel}\left(E_\theta + \frac{\omega}{K_\parallel}B_r\right)\right.$$

$$\left. + \frac{\{\omega+(1+\eta_i)\omega_{*i}\}}{K_\parallel}2\epsilon_n\omega_* E_\theta - V_i^2\left\{\frac{\omega-(1+\eta_e)\omega_{*i}}{\omega}\left(B_r + \frac{K_\parallel}{\omega}E_\theta\right)\right\}\right] \tag{7}$$

where we let $\tau=1$ for simplicity and $\omega_{*i} = -\omega_*$ where ω_* denotes the positive electron diamagnetic frequency, we obtain

$$\left(k_\perp^2\rho_s^2 + 2\epsilon_n\frac{\omega_*}{\omega}\right)E_\theta + \left(\frac{\omega-\omega_*}{\omega+\omega_*(1+\eta_i)} - \frac{K_\parallel^2 C_s^2}{\omega^2}\right)\left(E_\theta + \frac{\omega}{K_\parallel}B_r\right) = 0 \tag{8}$$

$$\frac{K_\parallel V_A^2 k_\perp^2 B_r}{\omega+\omega_*(1+\eta_i)} = -\left(k_\perp^2 + \frac{2\epsilon_n\omega_*}{\omega\rho_s^2}\right)E_\theta - \frac{\omega-(1+\eta_e)\omega_*}{\omega+(1+\eta_i)\omega_*}\frac{2\epsilon_n\omega_*}{\rho_s^2}\frac{B_r}{K_\parallel} \tag{9}$$

where $C_s^2 = \frac{T_e}{M}$, $\rho_s^2 = \frac{C_s^2}{\Omega_i^2}$, $\epsilon_n = (R\frac{\partial}{\partial r}\ell n N)^{-1}$ *and* $V_A^2 = \frac{NM\mu_0}{B^2}$.

For small enough β, and assuming $k^2 \rho_s^2 \gg 2\epsilon_n \frac{\omega^*}{\omega}$ *and* $k^2 \gg k_s^2$ as must be verified aposteriori

we obtain from Eq. (9)

$$B_r = -\frac{\omega + \omega_*(1+\eta)}{K_\parallel V_A^2} E_\theta \tag{10}$$

Substituting into Eq. (8) we obtain

$$\left[k_\perp^2 \rho_s^2 + 2\epsilon_n \frac{\omega_*}{\omega} + \frac{\omega - \omega_*}{\omega + \omega_*(1+\eta_i)} - \frac{K_\parallel^2 C_s^2}{\omega^2} - \frac{\omega(\omega - \omega_*)}{K_\parallel^2 V_A^2} + \frac{\omega + \omega_*(1+\eta)}{\omega} \beta_i \right] E_\theta = 0 \tag{11}$$

where $\beta_i = \frac{V_i^2}{V_A^2} = \frac{\mu_0 NT}{B^2}$.

Let $\frac{\omega}{\omega_*} = \Omega$, $k_\theta^2 \rho_s^2 = b_\theta$, $Z = s - j$ $(or\ K_\parallel = \frac{-Z}{qR})$ *and* $u = \frac{\Delta r}{\rho_s}$, noting that

$(\cos\theta + \frac{i}{k_\theta} \sin\theta \frac{\partial}{\partial r})\, \phi_j = \frac{1}{2} [\phi_{j+1} + \phi_{j-1} + \hat{s} \frac{d}{ds} (\phi_{i-1} - \phi_{i+1})]$ and using the strong coupling

and ballooning formalisms we can write:

$$\phi_{j+1} + \phi_{j-1} + \hat{s}\frac{d}{ds}\,(\phi_{j-1} - \phi_{j+1}) =$$

$$\phi_j(Z-1) + \phi_j(Z+1) + \hat{s}\frac{d}{ds}\,[\phi(Z+1) - \phi(Z-1)] =$$

$$2\left[1 + \left(\frac{1}{2} - \hat{s}\right)\frac{d^2}{dZ^2}\right] \phi_j(Z) \ . \tag{12}$$

We may now write the general differential equation which governs the stability of the ITGD mode in a finite-β toroidal plasma in the limit of strong toroidal coupling as

$$\frac{d^2E_\theta}{du^2} = \frac{b_\theta + \frac{\Omega-1}{\Omega+1+\eta} - \frac{\epsilon_n^2\hat{s}^2}{q^2\Omega^2}u^2 + \beta_i[1+\frac{1+\eta_i}{\Omega}] + \frac{2\epsilon_n}{\Omega} + \frac{\Omega(1-\Omega)}{\epsilon_n^2\hat{s}^2}q^2\Omega^2\beta_i}{1-\frac{2\epsilon_n}{\hat{s}^2\Omega b}(\frac{1}{2}-\hat{s})}\phi \tag{13}$$

III. Analysis

It may be easily seen that Eq. (13) agrees with the slab and toroidal equations in the zero-β limit. The analysis of the finite-β toroidal limit will be presented in a later work The slab limit will be treated here where Eq. (13) takes the form

$$\phi'' = \left(b_\theta + \frac{\Omega-1}{\Omega+1+\eta_i} - \frac{\epsilon_n^2\hat{s}^2u^2}{q^2\Omega^2} + \beta(1+\frac{1+\eta}{\Omega}) + \frac{\Omega(1-\Omega)q^2\beta_i}{\epsilon_n^2\hat{s}^2u^2}\right)\phi \tag{14}$$

$$or \quad \phi'' = -\lambda + \mu^2x^2 + \frac{\nu}{x^2} \tag{15}$$

Eq. (15) was studied in Ref. 14.

Let $\delta(\delta-1) = \nu \rightarrow \delta = (1\pm\sqrt{1+4\nu})/2$. For even modes the condition for localized solutions is given by

$$\lambda = \mu\left[4N+2-\sqrt{1+4\nu}\right] \qquad N = 0,1,2,...$$

From Eq. 14 we may then write the eigenvalue condition for even modes as

$$b_\theta + \frac{\Omega-1}{\Omega+1+\eta_i} + \beta + \frac{\Omega+1+\eta_i}{\Omega} = \frac{-iL_n}{\Omega L_s}\left[4N+2\mp\sqrt{1+4\Omega(1-\Omega)\beta\left(\frac{L_s}{L_n}\right)^2}\right] \tag{16}$$

which gives the stability condition

$$\beta^2\frac{L_s^2}{L_n^2} \geq \frac{1+4N}{1(1+\eta_i)}\left(\frac{1}{1+\eta_i} - b_\theta\right) \tag{17}$$

for $(1+\eta_i)\,b_\theta < 1$.

The validity of the assumptions made above can be easily checked. In the same manner we can show that β destabilizes the odd modes. It is interesting to note that the β required for stabilizing the even ITGD mode is of the same order as that obtained numerically in Ref. 46, where β was found to stabilize both odd and even modes. Further work is needed to cover various aspects of the problem.

IV. Conclusions

In the above work we reviewed the previous work made in studying the ion-temperature-gradient driven mode and presented the dispersion relation for the finite-β toroidal case. It was found that small amounts of β can strongly influence the stability properties of the mode.

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