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Electromagnetic electron temperature gradient driven instability in toroidal plasmas

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The fluid theory of a new type of electron temperature gradient instability is proposed. This mode is closely related to the short wavelength Alfvén mode in the regime $k_\perp^2 \rho_i^2 > 1$. Contrary to standard electron temperature gradient modes, which are mostly electrostatic, the considered mode is fundamentally electromagnetic and does not exist in the electrostatic limit. The mechanism of instability relies on gradients in both the electron temperature and magnetic field. It is suggested that this instability may be a destabilizing mechanism for collisionless microtearing modes, which are observed in a number of gyrokinetic simulations. *Published by AIP Publishing.*
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In magnetically confined plasmas, turbulence results in the so-called "anomalous" transport of particles and energy. Despite decades of work, various manifestations of anomalous transport are still not fully understood, particularly those which arise in tokamak physics.[1,2] Significant progress has been made in the understanding of anomalous particle and ion (energy) transport, which is driven largely by ion temperature gradient (ITG) modes.[3] Unfortunately, understanding the sources of anomalous electron (energy) transport has proved to be a more difficult problem.[2,4] Short wavelength electron temperature gradient (ETG) modes and trapped electron modes (TEMs) are being considered as possible sources of turbulence in some experiments.[2,4–8] Still, a number of observations, such as the dependence of electron energy transport on the plasma pressure parameter β, point to other transport mechanisms; particularly, those caused by the magnetic flutter associated with electromagnetic (EM) fluctuations.[2,9] ETG and TEM modes are mostly electrostatic and may in fact be stabilized by electromagnetic effects at higher β.[8]

Microtearing modes[10–13] are inherently electromagnetic and are thus natural candidates for electromagnetic electron transport. For high poloidal mode numbers, $m > 1$, when the linear tearing-mode stability parameter, $\Delta' = -2m/r$, is negative, mirotearing modes can be driven linearly unstable by thermal force effects. These microtearing modes are, however, fundamentally related to collisions, which could make them less relevant for modern tokamaks, where collisions are weak.

Electromagnetic fluctuations could also result from the nonlinear instability of small scale (high m) magnetic islands, which can be excited by diamagnetic and nonlinear diffusion effects.[14–18] These small-scale magnetic islands can be nonlinearly sustained, resulting in island overlap and magnetic stochasticity. In magnetically stochastic regions, the radial diffusion of magnetic field lines can result in significant electron energy transport. A number of electron transport models have been proposed based on this mechanism.[19,20] Magnetic fluctuations, and the subsequent magnetic stochasticity, could also be caused by nonlinear excitation of subdominant (linearly damped) microtearing modes, where electrostatic modes, such as the ITG mode, provide the main source of instability.[9,21,22]

Recent, high performance, gyrokinetic simulations have renewed interest in microtearing turbulence as a potential source of anomalous electron transport. In particular, these simulations show that magnetic fluctuations possessing tearing-mode parity provide a level of anomalous electron energy transport which is consistent with the experimental data, particularly for spherical tokamaks with large magnetic gradients.[23–26] Experimental measurements with HIBP diagnostics in the JIPPT-IIU tokamak also detect magnetic fluctuations with the general characteristics of microtearing dispersion.[27]

Classical microtearing modes,[10–13] destabilized by thermal forces, require finite electron temperature gradients and collisions. Still, a large body of gyrokinetic simulations indicates the presence of an additional, collisionless, destabilization mechanism, likely related to magnetic gradients.[25,28,29] In this paper, we present the local formulation of an instability related to the so-called short wavelength Alfvén mode. This mode is fundamentally electromagnetic and represents the extension of the Alfvén mode in the $k_\perp \rho_i \geq 1$ regime. Density gradients provide the drift corrections to the mode dispersion, which are important for small-scale drift magnetic islands.[16,17,30–33] We show that this mode can be destabilized by gradients in the magnetic field and electron temperature, and thus, we call it the electromagnetic electron temperature gradient (EM-ETG) mode. Although these destabilization mechanisms are similar to the standard toroidal ETG mode, the underlying nature of the short wavelength Alfvén mode is very different, as it is strongly electromagnetic and fundamentally depends on perturbations in the magnetic field.

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The dispersion equation for Alfvén waves, including the effect of finite ion Larmor radius, can be obtained from kinetic theory in the form[34]

$$\omega^2 = z\left(\tau + \frac{1}{1 - I_0(z)\mathrm{e}^{-z}}\right) v_A^2 k_{\|}^2, \tag{1}$$

where $z = k_\perp^2 \rho_i^2 = k_\perp^2 T_i/(m_i \omega_{ci}^2)$, and $\tau = T_e/T_i$. For small ($z < 1$) and large ($z > 1$) ion Larmor radius, (1) reduces to

$$\omega^2 = \left[1 + z\left(\frac{3}{4} + \tau\right)\right] v_A^2 k_{\|}^2 \quad z < 1, \tag{2}$$

$$\omega^2 = z(1+\tau) v_A^2 k_{\|}^2 \quad z > 1. \tag{3}$$

The EM-ETG mode is the generalization of (3) for inhomogeneous plasma with density, temperature, and magnetic field gradients. In this work, we develop an appropriate nonlinear model of this mode using the framework of fluid theory.

In the short wavelength limit ($k_\perp \rho_i > 1$), ions are umagnetized and follow the Boltzmann response

$$\tilde{n}_i = -\frac{e\phi}{T_i} n_0. \tag{4}$$

In the considered regime, the electrons remain strongly magnetized and are described by standard fluid equations. The electron continuity equation takes the form

$$\frac{\partial n_e}{\partial t} + \nabla_\perp \cdot (n_e \mathbf{v}_\perp) + \nabla_{\|}(n_e \mathrm{v}_{\|e}) = 0. \tag{5}$$

We assume quasineutrality in what follows, and thus we have $n_e = n_i = n_0 + \tilde{n}_i = n$. The perpendicular velocity is found from the electron momentum balance equation, where the pressure gradient and gyroviscosity[35] terms are included to account for the diamagnetic drift and finite Larmor radius effects, respectively,

$$m_e n \frac{d\mathbf{v}}{dt} = -en(\mathbf{E} + \mathbf{v}_e \times \mathbf{B}) - \nabla p - \nabla \cdot \mathbf{\Pi}. \tag{6}$$

Considering the low frequency regime $\omega \ll \omega_{ce}$, this yields

$$\begin{aligned}\mathbf{v}_{e\perp} &= \mathbf{v}_E + \mathbf{v}_{Pe} + \mathbf{v}_{Ie} + \mathbf{v}_{\pi e} \\ &= c\frac{\hat{\mathbf{b}} \times \nabla\psi}{B} - c\frac{\hat{\mathbf{b}} \times \nabla p_e}{enB} - \frac{c}{\omega_{ce}}\hat{\mathbf{b}} \times \frac{d}{dt}\mathbf{v}_{e\perp}^{(0)} - c\frac{\hat{\mathbf{b}} \times \nabla \cdot \mathbf{\Pi}}{enB},\end{aligned} \tag{7}$$

where $\mathbf{v}_{e\perp}^{(0)} = \mathbf{v}_E + \mathbf{v}_{Pe}$, $d/dt = \partial/\partial t + (\mathbf{v}_{e\perp}^{(0)} \cdot \nabla)$, and $\mathbf{v}_{e\perp}$ is composed of the $\mathrm{E} \times \mathrm{B}$ and diamagnetic drifts at lowest order, along with the inertial and gyroviscous drifts at higher order. The destabilizing effect is provided through the compressibility of the $\mathrm{E} \times \mathrm{B}$ and diamagnetic drifts in the continuity equation (5)

$$n(\nabla \cdot \mathbf{v}_E) = -2n\mathbf{v}_E \cdot \nabla \ln B = -\frac{n}{\tau}\mathbf{v}_{De} \cdot \nabla\left(\frac{e\phi}{T_i}\right), \tag{8}$$

$$\nabla \cdot (n\mathbf{v}_{Pe}) = -2n\mathbf{v}_{Pe} \cdot \nabla \ln B = n_0 \mathbf{v}_{De} \cdot \nabla\left(\frac{p_e}{p_{e0}}\right), \tag{9}$$

where

$$\mathbf{v}_{De} = -\frac{2cT_e}{eB}\hat{\mathbf{b}} \times \nabla \ln B. \tag{10}$$

The contribution of the inertial and gyroviscous drifts into the continuity equation are simplified by taking into account the gyroviscous cancellation,[35,36] which allows us to use

$$\mathbf{v}_{Ie} + \mathbf{v}_{\pi e} = -\frac{1}{\omega_{ce}}\hat{\mathbf{b}} \times \frac{d_0}{dt}\mathbf{v}_{e\perp}^{(0)}, \tag{11}$$

where $d_0/dt = \partial/\partial t + \mathbf{v}_E \cdot \nabla$. For these higher order terms, gradients and curvature of the magnetic field are neglected in the simplification of the continuity equation, and we find

$$\nabla \cdot [n_e(\mathbf{v}_{Ie} + \mathbf{v}_{\pi e})] = n_0 \rho_e^2 \nabla_\perp \cdot \frac{\mathrm{d}_0}{\mathrm{d}t}\nabla_\perp\left(\frac{1}{\tau}\frac{e\phi}{T_i} - \frac{p_e}{p_{e0}}\right). \tag{12}$$

The parallel velocity term in the continuity equation (5) is simplified using Ampère's law

$$J_{\|} = -en_e \mathrm{v}_{e\|} = \frac{c}{4\pi}\nabla \times \mathbf{B} = \frac{c}{4\pi}\nabla_\perp^2 \psi, \tag{13}$$

where ψ is the z-component of the magnetic vector potential for the perturbed magnetic field, $\tilde{\mathbf{B}} = \hat{\mathbf{z}} \times \nabla\psi$. The expression for the total magnetic field is

$$\mathbf{B} = B_0\hat{\mathbf{z}} + \hat{\mathbf{z}} \times \nabla\psi,$$

and the parallel gradient operator along the total field is

$$\nabla_{\|} = \frac{\mathbf{B} \cdot \nabla}{B} = \frac{\partial}{\partial z} + \frac{1}{B_0}\hat{\mathbf{z}} \cdot \nabla\psi \times \nabla. \tag{14}$$

Using the total electron density from (4) in Equation (8), then linearly expanding the pressure $p_e = nT_e$ in (12), results in the following simplifications:

$$-\frac{n}{\tau}\mathbf{v}_{De} \cdot \nabla\left(\frac{e\phi}{T_i}\right) = -\frac{n_0}{\tau}\mathbf{v}_{De} \cdot \nabla\left[\left(\frac{e\phi}{T_i}\right) - \frac{1}{2}\left(\frac{e\phi}{T_i}\right)^2\right], \tag{15}$$

$$n_0 \rho_e^2 \nabla_\perp \cdot \frac{\mathrm{d}_0}{\mathrm{d}t}\nabla_\perp\left(\frac{1}{\tau}\frac{e\phi}{T_i} - \frac{p_e}{p_{e0}}\right)$$

$$\approx n_0 \rho_e^2 \frac{\mathrm{d}_0}{\mathrm{d}t}\nabla_\perp^2\left[\left(\frac{1}{\tau} + 1\right)\frac{e\phi}{T_i} - \frac{\tilde{T}_e}{T_{e0}}\right]. \tag{16}$$

In Eq. (15), we have retained the scalar nonlinear term due to the density perturbation. Inserting these simplifications, along with (9) and (13), into Eq. (5) then yields the final form of the nonlinear continuity equation

$$\begin{aligned}&\frac{\partial}{\partial t}\left(\frac{e\phi}{T_i}\right) - \frac{\mathrm{v}_{*e}}{\tau}\frac{\partial}{\partial y}\left(\frac{e\phi}{T_i}\right) + \mathbf{v}_{De} \\ &\cdot \nabla\left[\frac{1}{\tau}\left(\frac{e\phi}{T_i}\right) - \frac{1}{2\tau}\left(\frac{e\phi}{T_i}\right)^2 - \frac{p_e}{p_{e0}}\right] \\ &-\rho_e^2 \frac{\mathrm{d}_0}{\mathrm{d}t}\nabla_\perp^2\left(\left(\frac{1}{\tau} + 1\right)\frac{e\phi}{T_i} - \frac{\tilde{T}_e}{T_{e0}}\right) \\ &+\frac{c\lambda_{De}^2}{\tau}\nabla_{\|}\nabla_\perp^2\frac{e\psi}{T_i} = 0,\end{aligned} \tag{17}$$

where

$$\mathrm{v}_{*e} = -\frac{cT_e}{eB}\frac{\partial \ln n_0}{\partial x}. \tag{18}$$

The remaining nonlinear equations are found by considering the parallel electron dynamics. We consider the limit $\omega \ll k_{\|} v_{Te}$, where kinetic theory gives the relation

$$\nabla_{\|} T_e = 0. \tag{19}$$

Note that (19) constrains the electron temperature to be constant along the perturbed magnetic field lines. In the linear limit, this condition becomes

$$ik_{\|}\tilde{T}_e - \frac{1}{B_0}\frac{\partial \psi}{\partial y}\frac{\partial T_{e0}}{\partial x} = 0. \tag{20}$$

The parallel momentum equation under condition (19) takes the form

$$0 = -en_e E_{\|} - T_e \nabla_{\|} n_e, \tag{21}$$

which, for $E_{\|} = -\nabla\phi + c^{-1}\partial\psi/\partial t$, yields the evolution equation for ψ

$$\frac{\partial \psi}{\partial t} = c\nabla_{\|}\phi - \frac{cT_e}{en_e}\nabla_{\|} n_e. \tag{22}$$

Note that the $\nabla_{\|} n_e$ operator involves the gradient of both the perturbed and equilibrium density, the latter of which is associated with perturbations in the magnetic field.

In the linear limit, and using the local approximation for the low field side of the tokamak such that $\mathbf{v}_{De} = \mathrm{v}_{De}\hat{\mathbf{y}}$, from Equations (17), (19), and (22), we arrive at the following dispersion relation:

$$\begin{aligned}\frac{(\omega - \omega_{*e})}{1+\tau}(\omega\tau + \omega_{*e}) = {} & c^2 k_{\|}^2 k_{\perp}^2 \lambda_{De}^2 \\ & + \omega_{De}[\omega - \omega_{*e}(1+\eta_e)] \\ & - k_{\perp}^2 \rho_e^2 [\omega - \omega_{*e}(1+\eta_e)],\end{aligned} \tag{23}$$

where $\lambda_{De}^2 = T_e/(4\pi e^2 n_0)$, $\rho_e^2 = T_e/(m_e \omega_{ce}^2)$, $\eta_e = \partial \ln T_e / \partial \ln n_0$, $\omega_{*e} = \mathrm{v}_{*e} k_y$, and $\omega_{De} = \mathrm{v}_{De} k_y$. Plots of this dispersion relation are shown in Fig. 1. In neglect of the dispersive terms, Eq. (23) is similar to the dispersion equation obtained in Ref. 6 from kinetic theory. Further neglecting temperature and magnetic field gradients, and noting that $c^2\lambda_{De}^2 = \tau v_A^2 \rho_i^2$, it is easy to see that (23) reduces to (3). The first term on the right of (23) thus originates from the Alfvén wave component, and is related to the bending of magnetic field lines. This term is stabilizing, whereas the component $\omega_{De}\omega_{*e}(1+\eta_e)$ of the second term, originating from gradients in the magnetic field and temperature, is destabilizing for sufficiently large η_e. An approximate condition for the instability can be written as

$$\tau v_A^2 \rho_i^2 k_{\|}^2 k_{\perp}^2 < \omega_{De}\omega_{*e}\eta_e. \tag{24}$$

Note that the condition $\omega < k_{\|} v_{Te}$ has to be satisfied for our approximation to remain valid. In terms of plasma parameters, the instability condition (24) can be written as

$$\beta_e > k_{\perp}^2 x^2 \frac{L_T S^2}{Rq^2}, \tag{25}$$

or, with $k_{\perp} \simeq k_x \gg k_y$ and $k_x x \simeq 1$

$$\beta_e > \frac{L_T}{R}\frac{S^2}{q^2}, \tag{26}$$

where x is the distance from the magnetic surface, S is the shear parameter, and where we have used the shear length, $L_s = qR/S$, $k_{\|} = k_y x/L_s$, and

$$\omega_{De} = \frac{2k_y cT_e}{eBR}. \tag{27}$$

Condition (26) is consistent with conclusions from gyrokinetic simulations that the electromagnetic transport increases with plasma pressure, and is more pronounced for spherical tokamaks (smaller R). When the destabilizing term is large,

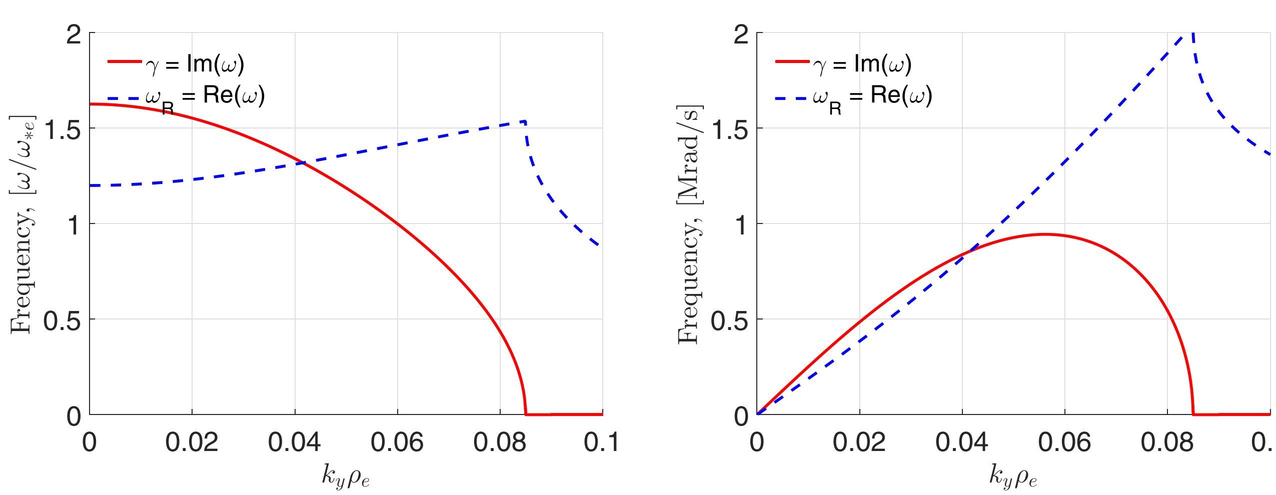

FIG. 1. Real frequency and growth rate of the EM-ETG mode; (left) normalized to the diamagnetic drift frequency, (right) in absolute value; for $1/x = k_x = 10k_y$, $S = 0.5$, and $\eta_e = 3$.

we can use the estimate $\omega^2 \approx \omega_{De}\omega_{*e}$, which makes the condition $\omega < k_{||}v_{Te}$ become

$$(k_y x)^2 > k_y^2 \rho_e^2 \frac{q^2 R}{S^2 L_T}. \quad (28)$$

It is worth discussing the relation of the EM-ETG mode to the standard ETG mode such as in Ref. 8. The fundamental difference is the condition $\omega \ll k_{||}v_{Te}$, which is employed in our work, versus the condition $\omega \gg k_{||}v_{Te}$ which is used for the standard ETG mode. For the standard case, $\omega \gg k_{||}v_{Te}$, the parallel electron motion can be neglected, and the basic equations for the electron density and pressure become

$$\frac{\partial n_e}{\partial t} + \mathbf{v}_E \cdot \nabla n_0 + n_0(\nabla \cdot \mathbf{v}_E) + \nabla_\perp \cdot (n_e \mathbf{v}_{pe}) = 0, \quad (29)$$

$$\frac{3}{2}\left(\frac{\partial p_e}{\partial t} + \mathbf{v}_E \cdot \nabla p_{0e}\right) + \frac{5}{2} p_{0e}(\nabla \cdot \mathbf{v}_E) + \frac{5}{2}\mathbf{v}_{De} \cdot \nabla p_e + \frac{5}{2} n \mathbf{v}_{De} \cdot \nabla T_e = 0. \quad (30)$$

Equations (29), (30), and the ion density equation (4) form the basic standard ETG model.[8] In this limit, the parallel electron current in the continuity equation and the heat flux in the energy equation are small and neglected. These terms are responsible for the electromagnetic corrections. Conversely, in the limit of $\omega \ll k_{||}v_{Te}$, the parallel electron streaming is dominant so that the electron temperature and density are determined by Equations (19) and (21), respectively. In this case, the electron continuity equation (17) must be viewed as an equation for the parallel current, where the density is found from (21). It is easy to see that the eigen-mode described by equations for (17), (19) and (21) is fundamentally dependent on the perturbed magnetic field. If the perturbed magnetic potential is neglected in Equation (22), the simple Boltzmann distribution for the electron density follows and no eigen-mode occurs from (22) and (4).

In this paper, a fluid theory for the EM-ETG instability in toroidal plasmas has been proposed. The closely related short wavelength (drift) Alfvén modes with $k_\perp^2 \rho_i^2 > 1$ have long been considered as a source of electromagnetic fluctuations in magnetized plasmas. Nonlinear structures related to these modes, such as vortices and magnetic islands, have been studied in a number of papers.[17,37–39] Stochastization of the magnetic field related to such structures was a main element of the electron energy transport models in Refs. 19, 20, 31, and 32. Furthermore, this instability may provide a dynamical route to the nonlinear regimes considered in Refs. 19, 20, 31, and 32. We hypothesize that the destabilization mechanism, resulting from electron temperature and magnetic field gradients, is similar to the destabilization caused by magnetic drift effects detected in gyrokinetic simulations.[23–25] Note that in the transition regimes $k_\perp^2 \rho_i^2 \geq 1$, the ion diamagnetic drifts may provide additional destabilization[40–42] to the short wavelength electromagnetic modes.

In order to investigate the EM-ETG mode's relevance to the excitation of small scale (high m) tearing modes, and to the accompanying formation of magnetic islands, two additional aspects must be incorporated into the model. First, the excitation of tearing modes and the subsequent change in magnetic field topology via reconnection requires an additional effect, such as dissipation and/or electron inertia. Neither of these were included in the current model. One can expect that collisionless dissipation due to wave-particle interaction (Landau damping) may be more relevant than collisions, which are weak for modern tokamaks. Landau damping can be incorporated into our model via linear kinetic closures, such as those given in Refs. 43 and 44. A similar model has actually been used by Kadomtsev,[45,46] which is also based on short wavelength Alfvén modes, and includes collisionless dissipation in the form of a kinetic closure for electrical conductivity. The dissipation allows for plasma slipping ("unfreezing") through the magnetic flux surfaces, while the underlying dynamics of the short wavelength Alfvén wave determine the temporal and spatial scales of the structures. Although the linear stability of the model detailed in Refs. 45 and 46 was not investigated, their finding that magnetic field lines pierce surfaces which are constrained to move with the plasma is suggestive of a change in magnetic topology, which is fundamental to the formation of magnetic islands.

Determining if the proposed EM-ETG mode (with the inclusion of dissipative/electron inertia effects) can lead to the formation of magnetic islands will also require the model to be developed in a non-local form. This will involve the explicit inclusion of magnetic shear effects, $k_{||} = k_{||}(x)$, which is the inherent spatial dependence required to create the structure of magnetic islands. The electron inertia may be important on its own and need to be included in the "inner sub-layer," where $\omega > k_{||}(x_i)v_{Te}$. The analysis with all these additional elements is left for future publication.

Other fluid theories of microtearing modes have been proposed in Refs. 47 and 48. Similar to our work, the fluid models in these papers included some terms due to the magnetic gradient drifts, but also electron inertia and equilibrium electron current. The destabilization in Ref. 47 was due to the collisions, similar to the standard theory of microtearing modes.[10,11]

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