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Drift effects on geodesic acoustic modes

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ABSTRACT

We present a fluid model for geodesic acoustic modes including diamagnetic effects due to inhomogeneous plasma density and temperature. Effects of ion parallel viscosity (pressure anisotropy), which allows to recover exactly the adiabatic index obtained in kinetic theory are considered. We show that diamagnetic effects lead to the positive up-shift of the GAM frequency and appearance of the second (lower frequency) branch related to the drift frequency. The latter is a result of modification of the degenerate (zero frequency) zonal flow branch which acquires a finite frequency or becomes unstable in regions of high temperature gradients.

Geodesic Acoustic Modes (GAM) [1], also referred to as a high frequency branch of Zonal Flow (ZF), have been actively investigated due to its role in turbulence and anomalous transport suppression in tokamaks [2]. Various aspect of these modes have been studied in recent years [3–7]. In high safety factor approximation, $q \gg 1$, one can formulate the expression for GAM frequency in the form: $\omega_{gam} = [\Gamma_1 + \Gamma_2/\tau + \mathcal{O}(q^{-2})]^{1/2} v_{Ti}/R_0$, where $v_{Ti} = \sqrt{2T_i/m_i}$ is the ion thermal write speed, $\tau = T_i/T_e$, R_0 is the major radius of tokamak and Γ_1 and Γ_2 are adiabatic indexes that have different values in different models. In one fluid ideal MHD model, the well-known pioneering work [1] yielded $\Gamma_1 = \Gamma_2 = 5/3$. Kinetic models [3,5,8] have predicted $\Gamma_1 = 7/4$ and $\Gamma_2 = 1$. It was shown [9] that pressure perturbations are intrinsically anisotropic in GAMs and that the account of the parallel ion viscosity (or pressure anisotropy) exactly recovers the adiabatic indexes obtained in kinetic models.

Standard GAM involves poloidally symmetric ($M = 0$) perturbation of electrostatic potential coupled with poloidal side-bands ($M = \pm 1$) of plasma pressure, density, and potential. Poloidally inhomogeneous perturbations effectively couple to gradient driven drift wave turbulence. Interaction of geodesic and density and temperature gradients (or diamagnetic effects) result in another type of modes, the so-called Beta Alfvén Eigenmodes (BAE) [10], which can be viewed as a finite $M \neq 0$ generalization of GAMs. BAE dispersion relations were derived and investigated in [11–13] in local approximation and later in [8,14,15] by using the ballooning theory. GAM and BAE have similar dispersion relations [5,16] with the difference residing in the poloidal and toroidal modes ($M = N = 0$ for GAM and $M, N \neq 0$ for BAE). BAE and related modes have been studied in a number of papers. In particular, unstable beta induced temperature gradient (BTG) modes have been found using MHD [12] and kinetic [13] models. Beta-induced Alfvén and Alfvén-Acoustic Eigenmodes (BAE and BAAE, respectively) have been widely investigated due to its role in background turbulence, generation of zonal flows and the possibility of application in MHD spectroscopy to diagnose safety factor profiles, $q(r)$, in tokamaks [17,18,7,19].

Much work in studies of GAM/BAE was done with kinetic theory, particularly, taking into account the resonance damping due to Landau mechanism [20,21]. Yet, the fluid theory remains an attractive alternative to full kinetic studies, especially, for studies of nonlinear effects, e.g. nonlinear generation of GAM due to drift wave turbulence [5,22,23]. Much of such work was done in fluid theory and it is desirable to have a fluid theory that self-consistently includes drift effects in both GAM modes as well in the background drift wave turbulence.

This Letter presents a two fluid MHD theory investigating the diamagnetic effects on GAM. We provide a simple analytical model for GAM in presence of diamagnetic effects. We show that in presence of diamagnetic effects, previously degenerated (zero frequency) mode, normally associated with zonal flow acquires a finite frequency. Moreover, for large gradients of the temperature, the latter mode becomes unstable in the regions of high q.

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In the study of GAM, it is reasonable to assume that ions are in the fluid regime, $\omega \gg v_{Ti}/qR_0$, and the electrons are in the adiabatic regime, $\omega \ll v_{Te}/qR_0$. In this context, electron temperature fluctuations can be neglected in the lowest order and the electron has a Boltzmann response as consequence. On the other hand, for ions, parallel viscosity, $\boldsymbol{\pi}_\parallel = 3\pi_\parallel(\mathbf{bb} - \mathbf{I}/3)/2$, must be consider to account for pressure anisotropy ($\pi_\parallel = 2(p_\parallel - p_\perp)/3$) effects.

In this Letter we use ideal MHD equations for ion and electrons in the form:

$$\frac{\partial n}{\partial t} + \nabla \cdot (n\mathbf{v}) = 0, \tag{1}$$

$$\frac{3}{2}\left(\frac{\partial p}{\partial t} + \mathbf{v}\cdot\nabla p\right) + \frac{5}{2}p\nabla\cdot\mathbf{v} + \nabla\cdot\mathbf{q} = 0, \tag{2}$$

$$\frac{d\pi_\parallel}{dt} + p\left(-2\mathbf{v}\cdot\nabla\ln B - \frac{2}{3}\nabla\cdot\mathbf{v}\right) + \frac{2}{5}\left(-2\mathbf{q}\cdot\nabla\ln B - \frac{2}{3}\nabla\cdot\mathbf{q}\right) = 0, \tag{3}$$

$$mn\frac{d\mathbf{v}}{dt} + \nabla p + \nabla\cdot\boldsymbol{\pi} - e_\alpha n(\mathbf{E} + \mathbf{v}\times\mathbf{B}) = 0. \tag{4}$$

We note that Eqs. (1), (2), (4), representing conservation of mass, energy and momentum respectively, have been widely used in both one and two fluid models. Eq. (3) is the parallel component of Grad type equation for the viscosity tensor, $\boldsymbol{\pi}$, that was derived for general curvilinear magnetic field in Ref. [24]. Parallel velocity, $v_\parallel$, which is responsible for $\mathcal{O}(q^{-2})$ corrections, is neglected in our analysis since GAM is mostly important at the edge of the plasma column.

Taking the cross product of Eq. (4) with $\mathbf{B}/B^2 = \mathbf{b}/B$, one finds to the equation for ion/electron currents

$$\mathbf{j} = \mathbf{j}_I + \mathbf{j}_p + \mathbf{j}_\pi + \mathbf{j}_\parallel + e_\alpha n\mathbf{V}_E, \tag{5}$$

where $\mathbf{j}_I = e_\alpha n\mathbf{b}\times(d\mathbf{v}/dt)/\omega_{c\alpha}$, $\mathbf{j}_p = \mathbf{b}\times\nabla p/B$, $\mathbf{j}_\pi = \mathbf{b}\times\nabla\cdot\boldsymbol{\pi}/B$, $\mathbf{j}_\parallel = J_\parallel\mathbf{b}$, $J_\parallel = \mathbf{j}\cdot\mathbf{b}$, $\mathbf{E} = -\nabla\tilde{\Phi} - \partial\tilde{A}_\parallel\mathbf{b}/\partial\theta$ and $\mathbf{V}_E = \mathbf{b}\times\nabla\tilde{\Phi}/B$.

Eqs. (1)–(5) are closed with quasi-neutrality condition in the forms

$$\nabla\cdot\mathbf{j} = 0, \qquad e(n_i - n_e) = 0. \tag{6}$$

Here, the subscript $\alpha = i, e$, standing for ions and electrons species, were omitted in n, $\mathbf{v}$, p, $\mathbf{q}$, m and $\pi_\parallel$ for simplicity of notation.

As in standard ideal MHD models, frozen field condition,

$$\mathbf{E} + \mathbf{v}\times\mathbf{B} = 0, \tag{7}$$

from Eq. (4), is considered to find the velocity in a first approximation [4]. The velocity, $\mathbf{V}_E$, is then used in the continuity equation for ions,

$$\frac{\partial\tilde{n}_i}{\partial t} + \mathbf{V}_E\cdot\nabla n_0 + n_0\nabla\cdot\mathbf{V}_E = 0, \tag{8}$$

to find the first side-bands components of ion density, $\tilde{n}_{i\pm1}$. Here the second term is responsible for the density drift effects. The following perturbed quantities, $\tilde{\Phi}, \tilde{n}, \tilde{p}, \tilde{\pi}_\parallel$, are assumed to be in the form

$$\tilde{X} = X_{-1}\exp(-i\theta) + X_0 + X_1\exp(i\theta), \tag{9}$$

since for the study of GAM only poloidal modes $M = 0, \pm1$ and the toroidal mode $N = 0$ are important. Then, by using (9) into (8) we obtain

$$\tilde{n}_{i\pm1} = \left(\pm\frac{i}{2}\frac{\omega_{di}}{\omega}\tilde{\Phi}_0 \mp \frac{\omega_*}{\omega}\tilde{\Phi}_{\pm1}\right)\frac{en_0}{T_i}, \tag{10}$$

where $\omega_{di} = k_r\rho_i v_{Ti}/R_0$ is the magnetic drift frequency, $\omega_* = \omega_{*i}$ is the density drift frequency defined by $\omega_{*i} = T_i n_0'/(reB_0n_0) = \rho_i v_{Ti}/2rL_N$ and $L_N = (d\ln n_0/dr)^{-1}$ is the characteristic density scale length. To solve (8) we use $\nabla\cdot\mathbf{V}_E = -2\mathbf{V}_E\cdot\nabla\ln B$ and we consider high aspect ratio approximation, $\epsilon = r/R_0 \ll 1$, and circular magnetic surfaces tokamaks, which gives $B \approx B_0(1 - \epsilon\cos\theta)$ and consequently $\nabla\ln B = (-\cos\theta\mathbf{e}_r + \sin\theta\mathbf{e}_\theta)/R_0$.

For electrons, due to its small mass, the parallel component of the momentum equation can be approximated by

$$T_e\nabla_\parallel\tilde{n}_e - en_0\nabla_\parallel\tilde{\Phi} = 0, \tag{11}$$

which leads to the Boltzmann response as solution for the first side-bands,

$$\tilde{n}_{e\pm1} = \frac{en_0}{T_e}\tilde{\Phi}_{\pm1}. \tag{12}$$

By substituting (10) and (12) in the quasi-neutrality condition, $e(\tilde{n}_i - \tilde{n}_e) = 0$, we obtain the relation between the potentials,

$$\tilde{\Phi}_{\pm1} = \pm\frac{i}{2}\frac{\omega_{di}/\omega}{\tau\pm\omega_*/\omega}\tilde{\Phi}_0, \tag{13}$$

in according with [25,13] (Eqs. (8) and (3.8), respectively) using kinetic model. Here the ratio of temperatures, $\tau = T_i/T_e$, is used from now on to simplify notations.

The reduced energy equation for ions is obtained in the form

$$\frac{3}{2}\left(\frac{\partial\tilde{p}_i}{\partial t} + \mathbf{V}_E\cdot\nabla p_{0i}\right) + \frac{5}{2}p_{0i}\nabla\cdot\mathbf{V}_E = 0. \tag{14}$$

It is solved for the first side-bands of ion pressure,

$$\tilde{p}_{i\pm1} = \left(\pm\frac{5}{3}\frac{i}{2}\frac{\omega_{di}}{\omega}\tilde{\Phi}_0 \mp (1+\eta_i)\frac{\omega_*}{\omega}\tilde{\Phi}_{\pm1}\right)en_0, \tag{15}$$

where $\eta_i = L_n/L_{T_i}$ and $L_{T_i} = d\ln T_i/dr$ is the characteristic ion temperature scale length. It is worth noting here that, in general, the heat flux ($\mathbf{q}$), nor the diamagnetic drift, are not divergence free in toroidal geometry. It can be shown however, that these effects do not contribute to the GAM dispersion in the leading order.

In case of electrons, since fluctuations of temperature are small in the adiabatic regime, electron pressure can be written as

$$\tilde{p}_{e\pm1} = \pm\tilde{\Phi}_{\pm1}en_0. \tag{16}$$

From the reduced ion parallel viscosity equation [4],

$$\frac{\partial\tilde{\pi}_{i\parallel}}{\partial t} - \frac{2}{3}p_{0i}\mathbf{V}_E\cdot\nabla\ln B = 0, \tag{17}$$

in which heat flux contributions can be neglected, we obtain the first side-bands of parallel viscosity,

$$\tilde{\pi}_{i\parallel\pm1} = \pm\frac{1}{3}\frac{i}{2}\frac{\omega_{di}}{\omega}n_0e\tilde{\Phi}_0, \tag{18}$$

where it can be noted that this equation is unchanged by diamagnetic effects.

Considering that electrons are isotropic and so $\tilde{\pi}_e = 0$, it is convenient to compute the $\sin\theta$ component of the combination $\tilde{p} + \tilde{\pi}_\parallel$ from both ions and electrons contribution,

$$(\tilde{p} + \tilde{\pi}_\parallel/4)_s = -\frac{\omega_{di}}{\omega}\left(\frac{7}{4} + \frac{\tau + (1+\eta_i)\omega_*^2/\omega^2}{\tau^2 - \omega_*^2/\omega^2}\right)n_0e\tilde{\Phi}_0. \tag{19}$$

This will be used in the current conservation equation averaged over the magnetic surfaces,

$$\langle\nabla\cdot\mathbf{j}\rangle = 0, \tag{20}$$

where $\mathbf{j}$ is the sum of the ion and electron currents and the average over the magnetic surfaces, $\langle\rangle$, is used to eliminate the parallel current contributions, since $\langle\nabla_\parallel J_\parallel\rangle = 0$. In Eq. (20), since $\omega_{ci} \ll \omega_{ce}$ and the ion and electron currents due to $\mathbf{V}_E$ cancel each other, we need to consider ion inertial current, the total (ion and electron) diamagnetic current, and the viscosity current for ions.

The divergence of these currents averaged over the magnetic surfaces are calculated in the leading order as follows:

$$\begin{aligned}\langle\nabla\cdot\mathbf{j}_I\rangle &\approx \left\langle\frac{en_0}{\omega_{ci}}\nabla\cdot\left(\mathbf{b}\times\frac{\partial\mathbf{V}_E}{\partial t}\right)\right\rangle \\ &\approx \left\langle -i\omega\frac{en_0}{\omega_{ci}}\nabla\cdot\left(\frac{-\nabla_\perp\tilde{\Phi}}{B}\right)\right\rangle \\ &\approx -(2\pi)^2 rR_0 i\omega\frac{en_0}{\omega_{ci}}\frac{k_r^2\tilde{\Phi}_0}{B_0},\end{aligned} \tag{21}$$

$$\begin{aligned}\langle\nabla\cdot\mathbf{j}_p+\nabla\cdot\mathbf{j}_\pi\rangle &= \langle -2\mathbf{j}_p\cdot\nabla\ln B+\mathbf{j}_\pi\cdot\nabla\ln B\rangle \\ &\approx \left\langle -\frac{2ik_r\tilde{p}\sin\theta}{BR_0}-\frac{1}{2}\frac{ik_r\tilde{\pi}_\parallel\sin\theta}{BR_0}\right\rangle \\ &\approx -(2\pi)^2 rR_0\frac{ik_r(\tilde{p}+\tilde{\pi}_\parallel/4)_s}{B_0R_0}.\end{aligned} \tag{22}$$

Finally, the dispersion relation can then be written in the form

$$\left[\omega^2-\frac{v_{Ti}^2}{R_0^2}\left(\frac{7}{4}+\frac{\tau+(1+\eta_i)\omega_*^2/\omega^2}{\tau^2-\omega_*^2/\omega^2}\right)\right]\frac{en_0k_r^2\tilde{\Phi}_0}{\omega\omega_{ci}B_0}=0. \tag{23}$$

This equation has following solutions:

$$\begin{aligned}\Omega_\pm^2 = \frac{1}{2}\Big(&\omega_{gam}^2+\omega_{*e}^2 \\ &\pm\sqrt{\left(\omega_{gam}^2+\omega_{*e}^2\right)^2+(4\eta_i-3)\omega_{*e}^2v_{Ti}^2/R_0^2}\Big),\end{aligned} \tag{24}$$

where $\omega_{*e} = -\omega_{*i}/\tau$ is the electron drift frequency and $\omega_{gam}^2 = (7/4+1/\tau)v_{Ti}^2/R_0^2$ is the GAM frequency in the absence of drift effects.

It can be seen from this equation, that diamagnetic effects due to temperature and density gradients modify the GAM frequency and create a new mode which becomes unstable for larger values of $\eta_i > 3/4$. The new mode occurs as a result of coupling to the degenerate (in absence of drift effects) zero frequency mode normally associated with zonal flow. For weak gradient, $\omega_{*e} \ll \omega_{gam}$ ($\rho_i/L_n \ll r/R_0$), the two branches of the solution can be approximated by

$$\Omega_+^2 = \omega_{gam}^2+\frac{1+\eta_i+1/\tau}{7/4+1/\tau}\omega_{*e}^2, \qquad \Omega_-^2 = \frac{3/4-\eta_i}{7/4+1/\tau}\omega_{*e}^2. \tag{25}$$

The Ω_+^2 branch is related to the GAM mode, and the Ω_-^2 branch is related to the ZF. The zonal flow branch becomes unstable for large temperature gradient, $\eta_i > 3/4$.

GAM results from the balance between the ion polarization current (Eq. (21)) and the geodesic current (Eq. (22)), which is produced by the $M = \pm 1$ side-bands of the perturbed pressure (together with parallel viscosity [4]). The pressure side-bands are in turn created by the flow due to $M = 0$ mode of the electrostatic potential [1]. As pointed out by [5], GAM are degenerated in the absence of diamagnetic effects in the $q \gg 1$ approximation. Drift effects directly affect the diamagnetic current, (Eq. (19)), breaking the symmetry of the dispersion relation (Eqs. (23) and (24)) removing the degeneracy of the zero frequency zonal flow mode. The instability occurs for high temperature gradients $\eta_i > 3/4$. The grow rate of this instability, which is associated with the zonal flow branch, as well as the modification of the GAM frequency are typically small for density/temperature gradients length scale of the order of the minor radius, i.e., $\gamma \sim (\rho_i R_0/rL_n)\omega_{gam}$. However, in the region of large density/temperature gradients, the drift corrections may become of the similar order as the GAM frequency. To our knowledge this instability of the ZF/GAM type mode has not been studied before, though we have shown that our final dispersion relation is a special case of a more general fishbone-like dispersion relation obtained by Zonca et al. [26]. The fluid model presented in this Letter provides a transparent treatment of the diamagnetic drift corrections and thus a shorter path to the understanding of the physical mechanism. The fluid treatment is also more easily generalized to the nonlinear case. One has to note that technically, the fluid treatment requires the condition $\omega > v_{Ti}/qR$, therefore our results are valid either in high q regimes, or in case of large plasma gradients, so that $\omega_* > v_{Ti}/qR$.

In stable case, $\eta_i < 3/4$, the dispersion relation (24), produces two GAM type modes, which, in the region of high density gradient, can have comparable frequencies. It is of interest to note that in experimentally, there are two mode with close frequencies observed [27,28]. One can conjecture that this mode splitting is due to plasma gradients effects.

Previously, the effect of equilibrium (poloidal and toroidal) rotation on GAM were investigated within the framework of one fluid MHD theory in [29,30]. It was found that the equilibrium plasma rotation induces not only the $\sin\theta$, but also $\cos\theta$ component of the perturbed density (and pressure); θ is the poloidal angle in toroidal configuration with equilibrium magnetic field $B = B_0(1-\varepsilon\cos\theta)$, $\varepsilon = r/R$, where r is the minor and R are major radii. In this respect, density gradients have similar effect as equilibrium rotation [29,30], causing the perturbed density to depend also on $\cos\theta$ (Eq. (10)) and $\tilde{\Phi}_c \neq 0$ (Eq. (13)), i.e., $\tilde{n}_i = \mathcal{O}(k_r^2\rho_i^2) + \mathcal{O}(k_r\rho_i)\sin\theta + \mathcal{O}(k_r\rho_i)\cos\theta$ and $\tilde{\Phi} = \tilde{\Phi}_0 + \tilde{\Phi}_s\sin\theta + \tilde{\Phi}_c\cos\theta$.

Higher order of dispersive terms, $\mathcal{O}(k_r^4\rho_i^4)$, which are important to understand the radial structure of GAM, although not included in this Letter, can be obtained by accounting for $\tilde{\Phi}_{\pm 2}$. Similar dispersion relation for electromagnetic modes including dispersive terms were derived in [12,13] (see also [8,5]). It is expected that, generally, $\tilde{\Phi}_k \sim (k_r\rho_i)^k\tilde{\Phi}_0$ (Eq. (13)), $\tilde{n}_{i0} \sim k_r^2\rho_i^2 e\tilde{\Phi}_0/T_i$, $\tilde{n}_{i\pm1} \sim k_r\rho_i e\tilde{\Phi}_0/T_i$ (Eq. (10)) [9], and, hence, as confirmed previously [1] the main quantities involving in the basic GAM study are $\tilde{\Phi}_0$ and $\rho_s\sin\theta$ (mass density).

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