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Fast excitation of geodesic acoustic mode by energetic particle beams

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A new mechanism for geodesic acoustic mode (GAM) excitation by a not fully slowed down energetic particle (EP) beam is analyzed to explain experimental observations in Large Helical Device. It is shown that the positive velocity space gradient near the lower-energy end of the EP distribution function can strongly drive the GAM unstable. The new features of this EP-induced GAM (EGAM) are: (1) no instability threshold in the pitch angle; (2) the EGAM frequency can be higher than the local GAM frequency; and (3) the instability growth rate is much larger than that driven by a fully slowed down EP beam. © 2015 AIP Publishing LLC.

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Geodesic acoustic modes (GAMs)1,2 are finite frequency components of zonal structures (ZS), with a predominantly antisymmetric density perturbation. GAMs have been intensively studied in the past two decades due to their potential roles in regulating symmetry-breaking microturbulence3–5 and the associated wave-induced transport. Further to this wave-wave coupling, resonant wave-particle interactions with energetic particles (EP), i.e., the so called energetic-particle-induced GAM (EGAM), also attracted significant attention.6,7 Due to their “zonal” mode structures, GAM cannot be excited by expansion free energy associated with pressure gradient, and velocity space anisotropy is needed for EGAM instability. In previous works, a slowing-down EP distribution with a localized pitch angle7–9 was adopted to study the EGAM excitation by neutral beam injection (NBI), considering that EP slowing down due to collisions of electrons are much faster than pitch angle scattering due to ion collisions. It was found that, for the fully slowed down EPs, only EPs with pitch angle $A_{B0} > 2/5$ are destabilizing,8 as later confirmed by numerical simulation.10 Here, $A = \mu/E$ is the pitch angle, $\mu = v^2/(2B)$ is the magnetic moment, and $B_0$ is the magnetic field at the magnetic axis and $E = v^2/2$. Berk and Zhou,11 meanwhile, proposed that the sharp gradient in pitch angle induced by NBI prompt loss could lead to fast excitation of GAM. EGAM excitation by a bump-on-tail ion distribution function is addressed in Refs. 12 and 13, for the case where EP source is the high-energy tail in ion distribution function induced by radio frequency (RF) heating; and possible interaction between EGAM and drift wave turbulence is discussed in Refs. 12, 14, and 15.

A recent paper by Ido et al.16 presented experimental evidence in the Large Helical Device (LHD), where EGAM is observed during tangential neutral beam injection, and some peculiarities appear in the comparison with theoretical predictions. The EGAM is excited before the injected NBI is fully slowed down, and the EP birth energy ($\sim 170$ KeV) is much larger than that predicted for wave-particle resonance;8 thus, earlier theories based on a fully slowed down NBI7,8 cannot be directly applied here. In this work, we explain why EGAM is excited by a not fully slowed down EP beam, with the EP distribution function being a function of time. As we will show later, the instability drive comes from the positive velocity space gradient in the low-energy end of the EP distribution function, similar to that of a bump-on-tail distribution. The EGAM excited by such an EP beam can be applied to explain the experimental observations in LHD.16 For the sake of clarity, we focus only on the local dispersion relation of EGAM, while neglecting system nonuniformity and higher order effects, such as finite Larmor radius. Modulation of EP distribution function due to the excitation of EGAM is not taken into account either; thus, this work presents the linear theory of EGAM excitation by EP distribution function peculiar to the LHD scenarios. For the sake of simplicity, we ignore the helicity of the device and assume large aspect ratio, consistent with the experiment observation in the center of the device using heavy ion beam probe.

The dispersion relation of the $n = 0$ electrostatic EGAM can be derived from quasi-neutrality condition

$$\sum_s \frac{Q}{m} \frac{\partial F_{0s}}{\partial \phi} \frac{\partial \phi}{\partial s} \left. \frac{\partial \phi}{\partial s} + J_0 (k_s n_L) \delta H_s \right|_s = 0,$$

(1)

with $\delta H_s$ being the nonadiabatic component of the perturbed particle distribution function

$$\delta f_s = \frac{Q_s}{m_s} \frac{\partial F_{0s}}{\partial \phi} \delta \phi + \exp \left\{ \frac{i m_c}{QB^2} k \times B \cdot \mathbf{v} \right\} \delta H_s,$$

d$\delta H_s$ derived from linear gyrokinetic equation17–19

$$\left( \omega - \omega_d + i \omega_d \frac{d}{dt} \right) \delta H_s = - \frac{Q_s}{m_s} \frac{\partial F_{0s}}{\partial \phi} \omega J_0 (k_s n_L) \delta \phi.$$

(2)

Here, the subscripts $s = e, i, h$ represent electrons, ions, and EPs, respectively; and $Q$ is the electric charge. Furthermore, $J_0 (k_s n_L)$ is Bessel function of the zero-index,
with $\rho_L = mc v_L / (Bq)\eta_L$ being the Larmor radius, $\omega_{tr} = v_L / (q\rho_L)$ is the transit frequency, $\omega_d = -k_r (v_L^2 + 2v_r^2) / (2QR_0)$ sin $\theta$ is the magnetic drift frequency; and other notations are standard.

The thermal plasma response to GAM has already been derived in several earlier works, and we shall not provide the details here. Electron response to EGAM is adiabatic to $m \neq 0$ poloidal harmonics of the scalar potential; while the response of the ions is derived noting $k_r \rho_i \ll 1$ and $q \gg 1$. EP response can be derived by transforming into drift orbit central coordinates. Since, in the scenario considered here, EGAM is excited before the EPs are fully slowed down, we assume the following distribution for EPs:

$$F_{0h} = \frac{c_0 H(E_b - E)H(E - E_L)\delta(A - \Lambda_0)}{E^{3/2} + E_{crit}^{3/2}}$$

This distribution function is derived by solving Vlasov equation as a function of time, keeping only the slowing down term due to collisions with electrons in the collision operator. Here, $c_0 = \Gamma / (4\pi v_e)$ with $\Gamma$ being the NBI particle flux, $\gamma_i$ is the slowing-down rate, $H$ is the Heaviside step function, $E_b$ is the EP birth energy, $E_L \approx E_h \exp(-\gamma_i t)$ is the time dependent lower end of the distribution function, and $E_{crit}$ is the critical energy, where the EP pitch angle scattering rate of thermal ions is the same as the EP slowing down rate on thermal electrons. We typically have $E_b \gg E_L \gg E_{crit}$. Assuming small EP drift orbit normalized to EGAM wave length and keeping only $l = \pm 1$ transit harmonics, one then obtains the following dispersion relation from the surface averaged quasi-neutrality condition:

$$-\frac{e}{m} n_i k_r^2 \frac{1}{Q_i^2} \left(1 - \frac{\omega_{Q_i}^2}{\omega^2}\right) \delta(\phi + \delta\phi_h) = 0,$$

where $\omega_{Q_i}$ is the GAM frequency and the perturbed EP density

$$\delta\rho_{hi} \approx 2\pi B \sum_{\sigma = \pm 1} \int \frac{\delta Ed\Delta E}{|v_i|} \delta H_h$$

$$= 2A \int \frac{(2 - \Lambda_0)^2}{(1 - \Lambda_0)^1/2} \frac{B dE \Delta E^{5/2} \delta E F_{0h}}{2E(1 - \Lambda_0) - \omega^2 q^2 R_0^2}.$$  

Here, $A = \pi c e^2 k_r^2 \delta(\phi_{hi} - \sqrt{2B_0 Q_i})$.

The dispersion relation can then be derived as

$$-1 + \frac{\omega_{Q_i}^2}{\omega^2} + \frac{\pi}{\sqrt{2}} B_0 q^2 \frac{c_0}{n_0}$$

$$\times \left[ C \left( \ln \left( 1 - \frac{\omega_{Q_i}^2}{\omega^2} \right) - \ln \left( 1 - \frac{\omega_{E_i}^2}{\omega^2} \right) \right) \right]$$

$$+ D \left( \frac{1}{1 - \omega_{Q_i}^2/\omega^2} - \frac{1}{1 - \omega_{E_i}^2/\omega^2} \right) = 0.$$  

Here, $C = (2 - \Lambda_0 B_0) / (5\Lambda_0 B_0 - 2) / (2(1 - \Lambda_0 B_0)^{5/2})$, $D = \Lambda_0 B_0 / (2 - \Lambda_0 B_0) / (1 - \Lambda_0 B_0)^{5/2}$, and $\omega_{Q_i}$ and $\omega_{E_i}$ are the transit frequencies defined at $E_b$ and $E_L$, respectively. Note that $\omega_{Q_i}$ is a slowly varying function of $t$.

Note that there are two different singularities in the dispersion relation; i.e., the logarithmic terms (the first ones in the square bracket) typical of a slowing down distribution function, and the simple poles (the second ones in the square bracket) from the integration limits. As noted in Ref. 8, the simple pole at $1 - \omega_{Q_i}^2/\omega^2 = 0$ only modulate the EGAM real frequency, but it does not contribute to mode excitation. The logarithmic terms, on the other hand, are destabilizing for $\Lambda_0 B_0 > 2/5$; thus, there is a threshold, $\Lambda_0 B_0 > 2/5$, for EGAM excitation by the fully slowed down EP beam case. Meanwhile, for the not fully slowed down distribution function considered here, the simple pole at $1 - \omega_{Q_i}^2/\omega^2 = 0$ is also destabilizing; thus, there is no threshold in the pitch angle $\Lambda_0 B_0$ for EGAM excitation by the EPs.

The dispersion relation can be solved numerically as a function of $\tau = \gamma_i t$, and yields the slow temporal evolution of the excited EGAM due to the slowing down of the EP beam. There are three branches: A GAM branch with $\omega_0 \approx \omega_{Q_i}$; a lower beam branch (LBB) with $\omega_0 \approx \omega_{E_i}(t)$; and an upper beam branch (UBB) with $\omega_0 \approx \omega_{Q_i}$. Here, we take $\omega_{Q_i} = 1 - 0.05\gamma_i$ with the imaginary part being the Landau damping rate of GAM due to resonance with thermal ions, $\omega_{Q_i} = 3 \gg |\omega_{E_i}|$, and $\Lambda_0 B_0 > 2/5$ such that the coefficient of the logarithmic term (i.e., $C$) is positive. Note that GAM Landau damping rate typically increases with decreasing GAM frequency; however, Landau damping is much smaller than the strong resonant drive of EPs. Thus, our results are insensitive to the actual value of GAM Landau damping, which is assumed constant for the sake of simplicity. The real frequencies and growth rates are shown, respectively, in Figs. 1 and 2. The horizontal axis is time in units of $\gamma_i^{-1}$. We may see that only LBB with the frequency $\omega_0 \approx \omega_{Q_i}$ can be unstable, consistent with the fact that, in the present case of experimental interest for LHD experimental observations, the UBB frequency is larger than the GAM frequency. The LBB is stable when $\omega_0 \gg \omega_{Q_i}$ as $\omega_0$ approaches $\omega_{Q_i}$ with $t$, the growth rate of LBB increases significantly ($t = 0.8-1.2$), and then decays quickly as $E_0$ further decreases ($t = 1.2-1.5$). Later, ($t > 1.5$), the growth rate decays very slowly, due to the contribution of the destabilizing logarithmic term. So for $\Lambda_0 B_0 > 2/5$, the mode can be interpreted as a double-beam plasma instability with the two
singularities dominating at different times. As \( \omega_L \) approaches \( \omega_G \), the simple pole dominates; however, as \( \omega_L \) further decreases and becomes smaller than \( \omega_G \) by a finite amount, the contribution from the simple pole is negligible, while the logarithmic term still contributes to destabilizing the EGAM.\(^8\) The strong instability at \( \omega_L(t) \simeq \omega_G \) may also provide an explanation for the fast growth of EGAM observed experimentally. We note also that the frequency of the unstable LBB can be significantly larger than \( \omega_G \), as is shown in Fig. 1. This may explain the higher-frequency branch of EGAM observed in LHD.\(^16\) Another evidence that the present theory can be applied to explain the LHD experiment is that, in the discharges with electron cyclotron resonance heating (ECH), the dependence of the observed EGAM frequency on plasma temperature is weaker. This is because the slowing down rate is lower in the ECH discharges due to higher electron temperature, such that the unstable mode is a typical “beam-branch” with the frequency determined by the transit frequency of EPs.\(^8\)

The real frequency and growth rate for \( \Lambda_0B_0 < 2/5 \) are shown in Figs. 3 and 4, respectively. In this case, the logarithmic term is stabilizing;\(^8\) thus, the EGAM discussed here is similar to beam-plasma instability, which, however, has a double pole instead of the simple pole of the present case.

We can see that the behaviors of the three branches are similar to the case with \( \Lambda_0B_0 > 2/5 \). However, when \( \omega_L \) becomes smaller than \( \omega_G \) by a finite amount, the growth rate of LBB decreases to zero as the contribution of the simple pole becomes vanishingly small, similar to that of a beam-plasma instability. But it is clearly shown that, due to the contribution from the positive gradient in the lower-energy end of the distribution function, EGAM can still be driven unstable here, unlike the case of the fully slowed-down EP beam.

In conclusion, we elucidate the mechanism for GAM excitation by a not fully slowed down EP beam, which can be applied to explain the experimental observations in Large Helical Device. There are two kinds of singularities in the dispersion relation: one is the logarithmic singularity typical of a slowing-down distribution function and the other one is the simple pole from the low- and high-energy end of the distribution function. It is found that only the lower beam branch with the frequency determined by the low-energy limit of the distribution function is unstable, consistent with the fact that the upper beam branch frequency is larger than the GAM frequency as in the case of experimental interest. For \( \Lambda_0B_0 > 2/5 \), both singularities are destabilizing, with the simple pole dominating as \( \omega_L \simeq \omega_G \); and the contribution of the logarithmic singularity taking over as \( \omega_L \) becomes much smaller than \( \omega_G \). On the other hand, for \( \Lambda_0B_0 < 2/5 \), only the simple pole is destabilizing, and the behavior of EGAM is very similar to a beam-plasma instability. In both cases, the real frequency of the unstable lower beam branch can be significantly higher than the GAM frequency, providing an explanation for experimental observations. Note that the real frequency of the unstable lower beam branch can be higher than local GAM frequency. This does not conflict with the argument we used to explain why the upper beam branch is always stable, since the simple pole at \( \omega \simeq \omega_0 \) is always stabilizing, so one needs the unstable branch to have a frequency lower than local GAM frequency for the instability of the logarithmic term, while the lower beam branch is unstable at \( \omega_L > \omega_G \) due to the contribution of the simple pole. Another novel result of this work is that the EGAM drive is much stronger than that predicted in earlier
theories,\textsuperscript{7,8} since the destabilizing simple pole singularity obtained in the present case is much stronger than the logarithmic singularity from a fully slowed down EP distribution function.

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