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Abstract
The instability of the geodesic acoustic mode (GAM) in tokamak turbulence is analyzed. It can be induced by dynamic shearing of the ambient turbulence by GAMs combined with the poloidal inhomogeneity of the turbulent flux. The dispersion relation is derived. The competition of the drive mechanism and the damping by turbulence viscosity is discussed. The GAMs are more unstable for high safety factors.

Keywords: zonal flow, geodesic acoustic mode, dynamic shearing of turbulence, nonlinear instability, poloidal asymmetry for turbulence

1. Introduction

In toroidal plasmas, turbulent fluctuations in the range of drift wave frequencies are induced by plasma inhomogeneity. They induce turbulent transport of plasma and energy. At the same time, they cause perturbations of radial electric field which are constant on magnetic surface. Such electric field in semi-micro scale length on the other hand suppresses the turbulent fluctuations. The mutual interaction between the turbulent fluctuations and semi-micro scale electric field is the central issue in the study of high-temperature toroidal plasmas [1,2].

Among the electric field perturbations of semi-micro scale length, the geodesic acoustic mode (GAM) [3] has attracted attentions. This is an oscillating perturbation of radial electric field which is constant on magnetic surface. The associated perturbations in density and in flow velocity along the field line have dependence on the poloidal angle. The $E \times B$ velocity by a radial electric field which is constant on magnetic surface requires the secondary return flow if it is static [1, 4], as has been shown for zonal flows [5]. However, the temporal oscillation can compensate the divergence of poloidally-independent $E \times B$ velocity, and the secondary return flow is much smaller for the GAM. This difference of the magnitude of the secondary flow along the magnetic field, which is subject to the damping by the turbulent shear viscosity, leads to the difference in excitation conditions of zonal flow and GAM. The direct nonlinear simulation (DNS) has verified that GAM is excited by turbulence [6, 7] more dominantly excited by turbulence than zonal flow when the safety factor q becomes higher [8]. The influence of oscillating shear flow (such as GAM) on drift wave turbulence has been studied (e.g., [2]), and the mechanism of turbulent shearing that induces GAM has been found in DNS [8]. The mechanism of combination of geodesic curvature and turbulence shearing may be called Winsor-Hallatschek (W-H) mechanism. Recently, experimental investigation of GAM oscillation has been in progress [9-11], and theoretical understanding of the excitation mechanism attracts attentions.

In this article, an analytic discussion is made for the onset of GAM by W-H mechanism. The dispersion relation is derived. The toroidal effect being combined with turbulence shearing is found to cause the GAM instability. The competition between the driving mechanism and damping by turbulent shear viscosity is also discussed. It is found that the GAMs are more unstable for high safety factors in tokamaks.

2. Models

We consider a tight-aspect-ratio tokamak with circular-cross section. The drive of poloidal flow by poloidally-asymmetric cross field transport has been explored by Stringer [12]. The analysis has been extended by Hassam and others in [13]. Following [13], we employ the continuity equation and equation of motion in the fluid model as

$$\frac{\partial}{\partial t} n + \nabla \cdot n \boldsymbol{V}_{\perp} + \nabla_{//} n V_{//} = S - \nabla \cdot \Gamma \tag{1}$$

$$n m_i \left(\frac{\partial}{\partial t} \boldsymbol{V} + \boldsymbol{V} \cdot \nabla \boldsymbol{V} \right) + \nabla p - \boldsymbol{J} \times \boldsymbol{B} = S m_i \boldsymbol{V} \tag{2}$$

together with the charge neutrality condition

$$\nabla \cdot \boldsymbol{J} = 0 \tag{3}$$

and the Ohm's law $\boldsymbol{E} + \boldsymbol{V} \times \boldsymbol{B} = 0$ as governing equations. Here, n is the density, $\boldsymbol{V}$ is the velocity, $\boldsymbol{J}$ is the current, $p = nT$ is the pressure, and the temperature gradient is neglected for simplicity. The terms S and Γ represent the (equilibrium) particle source and flux. In toroidal plasmas, the source and the divergence of the perpendicular flux can be inhomogeneous in poloidal direction. Here we write $F \equiv \left(S - \nabla \cdot \Gamma \right) / n_0$ symbolically as

$$F = F_0(r) + F_1(r, \theta) + \tilde{F}(r, \theta, t) \tag{4}$$

where n_0 is the average density, $F_0(r) + F_1(r, \theta)$ is a stationary part including the poloidal asymmetry and $\tilde{F}(r, \theta, t)$ is the time-varying component associated with GAM. (r, θ) denote the radial and poloidal coordinates. Owing to the presence of the poloidally asymmetric term $F_1(r, \theta)$, the poloidally asymmetric stationary flow $V_{\|}(r, \theta)$ arises as

$$V_{\|,1} = qR \int^{\theta} F_1(r, \theta) \mathrm{d}\theta \,. \tag{5}$$

Particle conservation requires $\int_0^{2\pi} F_1(r, \theta) \mathrm{d}\theta = 0$.

The linearized equation for the poloidally symmetric $E \times B$ motion $\tilde{V}_{E \times B}$ has been derived in [13] in the leading orders of $\varepsilon = r / R$, the inverse aspect ratio. Keeping the term of $\tilde{F}(r, \theta, t)$ and turbulent shear viscosity effect, Eqs.(1)-(3) are rewritten as

$$\frac{\partial}{\partial t} \tilde{N} - \frac{2}{R} \sin\theta \, \tilde{V}_{E \times B} + \nabla_{\|} \tilde{V}_{\|} = \tilde{F}(r, \theta, t) \,, \tag{6}$$

$$\frac{\partial}{\partial t} \tilde{V}_{\|} + \frac{1}{r} \tilde{V}_{E \times B} \frac{\partial}{\partial \theta} V_{\|1} + c_s^2 \nabla_{\|} \tilde{N} - \mu_{\|} \nabla_{\perp}^2 \tilde{V}_{\|} = 0 \,, \tag{7}$$

$$\frac{\partial}{\partial t} \tilde{V}_{E\times B} + \frac{2}{R} c_s^2 \oint \frac{d\theta}{2\pi} \sin\theta \, \tilde{N} = 0 \, , \tag{8}$$

where $\tilde{N} = \tilde{n} / n_0$ and tilde denotes the perturbation with the time dependence $\exp(-i\omega t)$. The $E \times B$ velocity perturbation $\tilde{V}_{E\times B}$ is constant on a magnetic surface. Other perturbed quantities, $\tilde{N}$, $\tilde{V}_{\parallel}$ and $\tilde{F}$ have poloidal-angle-dependence. This set of equations (6)-(8) describes the GAM (in addition to the Stringer spin-up as has been discussed in [13]). For the GAM, the density perturbation has a $\sin\theta$ dependence, i.e., $\tilde{N}(r, \theta; t) = \tilde{N}(r; t) \sin\theta$. In the absence of $\tilde{F}(r, \theta, t)$, $V_{\parallel, 1}$ and $\mu_{\parallel}$, the dispersion relation of GAM [3] is given from Eqs.(6)-(8) as

$$\omega^2 = \omega_{GAM}^2 \equiv \frac{c_s^2}{R^2}\left(2 + \frac{1}{q^2}\right) . \tag{9}$$

3. Instability of geodesic acoustic mode

There are two ways of driving GAM by toroidicity in F. The first is through a stationary flow $V_{\parallel 1}$ in Eq.(7). Neglecting $\tilde{F}(r, \theta, t)$, one has from Eqs.(6) and (7) as

$$\tilde{N} = \frac{1}{\omega(\omega + i\nu) - c_s^2 k_\parallel^2}\left(\frac{-1}{r}\frac{\partial F_1}{\partial\theta} + \frac{2i(\omega + i\nu)}{R}\sin\theta\right)\tilde{V}_{E\times B} \, , \tag{10}$$

where $k_\parallel = 1/qR$ holds for this perturbation, $\nu = \mu_\parallel K_\perp^2$, $K_\perp^2 = K_r^2 - r^{-2}\partial^2/\partial\theta^2$ and K_r is the radial modenumber of GAM. Substituting it into Eq.(8), the dispersion relation is given as

$$\omega = \frac{2c_s^2}{R^2}\frac{(\omega + i\nu) - i\gamma_{asym}}{\omega(\omega + i\nu) - c_s^2 k_\parallel^2} \tag{11}$$

where the driving term by the asymmetric flux is given as

$$\gamma_{asym} = \frac{R}{r}\int_0^{2\pi}\frac{d\theta}{2\pi}\cos\theta \, F_1(\theta) \, . \tag{12}$$

This relation has been derived in [13] in conjunction with the Stringer spin-up problem. In addition to the Stringer spin-up, Eq.(11) also describes the GAM instability. In the limit of weak growth rate and damping rate, $\omega_{GAM} >> \mathrm{Im}(\omega), \gamma_{asym}, \nu$, the growth rate is given as

$$\mathrm{Im}\left(\omega\right) = \frac{-q^2\gamma_{\mathrm{asym}}}{1+2q^2} - \frac{\mu_{\|}K_{\perp}^2}{2\left(1+2q^2\right)} . \tag{13}$$

This result shows that the GAM can be excited when the divergence of cross-field flux has minimum at the inside of torus, $F_1 \sim -\cos\theta$. Owing to the damping by parallel viscosity $\mu_{\|}$, this instability condition is hardly satisfied.

The second mechanism of the GAM excitation is a dynamic shearing of the flux by waves. This W-H mechanism has been found by numerical simulations in [8]. The analytic modelling is derived as follows. In toroidal plasmas, the phase of microfluctuation is weakly dependent in the direction of majour radius and varies mainly in the vertical direction. (See Fig.1) Under this circumstance, the imposition of sheared poloidal rotation induces the change of radial wavenumber k_r of drift wave fluctuations. We denote the angle between the mid plane and the wave front by θ_1. In the absence of GAM, one has $\theta_1 \simeq 0$ and $k_r \simeq k_{\perp 0}\sin\theta$ holds at the poloidal location θ, where $k_{\perp 0}$ is a characteristic wave number of drift wave fluctuations in the absence of GAM. The phase front starts to rotate by the imposition of the zonal flow, and the change of k_r is given by the relation $\mathrm{d}k_r/\mathrm{d}t = -\left(k_{\perp 0}\cos\theta\right)\mathrm{d}\theta_1/\mathrm{d}t$. That is, the quantity k_r^2 evolves as

$$\mathrm{d}k_r^2/\mathrm{d}t = 2\,k_r\,\mathrm{d}k_r/\mathrm{d}t = -\left(2\,k_{\perp 0}^2\cos\theta\sin\theta\right)\mathrm{d}\theta_1/\mathrm{d}t \tag{14}$$

up to the first order of θ_1. This corresponds to the first order response to GAM perturbation, because the angular frequency is given as $\mathrm{d}\theta_1/\mathrm{d}t = \partial\tilde{V}_{E\times B}/\partial r$. The perturbation of k_r^2, $\delta\left(k_r^2\right)$, is given as

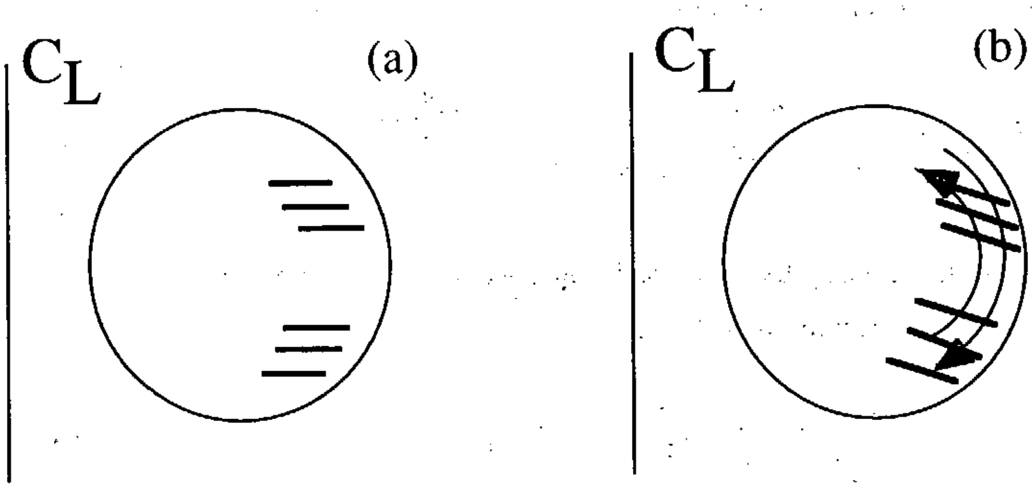

Fig.1 Schematic illustration of dynamic shearing of ambient fluctuations by GAM. Contours of phases of toroidal fluctuations are shown by thick lines in (a). In the presence of poloidally-symmetric shearing motion, the phase contours are modified with up-down asymmetry.

$$\delta\left(k_r^2\right) = -k_{\perp 0}^2 \sin 2\theta \int^t dt \frac{\partial}{\partial r} \tilde{V}_{E\times B} , \qquad (15)$$

As has been discussed in refs.[14, 15], the drift wave action $\left(1 + k_\perp^2 \rho_s^2\right)^2 \left|\phi_{\text{micro}}^2\right|$ is a quasi conserved during the process of modulation of perpendicular wavenumbers. ($\left|\phi_{\text{micro}}^2\right|$ is the fluctuation amplitude of drift waves.) Thus, $\left|\phi_{\text{micro}}^2\right|$ and turbulent flux Γ are modulated in time associated with the oscillating $E \times B$ flow. From the conservation of the drift wave action, the modification of turbulence amplitude, $\delta\left|\phi_{\text{micro}}^2\right|$, is given as

$$\delta\left|\phi_{\text{micro}}^2\right| \simeq -2 \frac{\rho_s^2 \, \delta\left(k_r^2\right)}{\left(1 + k_{\perp 0}^2 \rho_s^2\right)} \left|\phi_{\text{micro}}^2\right| . \qquad (16)$$

When one studies cases of strong turbulence, one has a relation $\Gamma \propto \sqrt{\left|\phi_{\text{micro}}^2\right|}$, i.e., $\delta\Gamma/\Gamma = \left(1/2\right)\delta\left|\phi_{\text{micro}}^2\right| / \left|\phi_{\text{micro}}^2\right|$. That is, the modulation of the flux is given as

$$\delta\Gamma \simeq -\frac{\rho_s^2 \, \delta\left(k_r^2\right)}{\left(1 + k_{\perp 0}^2 \rho_s^2\right)} \Gamma . \qquad (17)$$

(If one studies the quasilinear limit, $\Gamma \propto \left|\phi_{\text{micro}}^2\right|$, a relative modulation of the flux is given as $\delta\Gamma/\Gamma \simeq \delta\left|\phi_{\text{micro}}^2\right| / \left|\phi_{\text{micro}}^2\right|$. An additional factor 2 appears in the right hand side of Eq.(17).) The modification of F is evaluated straightforwardly. Noting that the radial derivative of $\delta\left(k_r^2\right)$ is much sharper than that of Γ/n , modulation of F , $\tilde{F} = -\partial\left(\delta\Gamma / n\right)/\partial r$, is estimated as

$$\tilde{F} \simeq \frac{\rho_s^2}{\left(1 + k_{\perp 0}^2 \rho_s^2\right)} \frac{\Gamma}{n} \frac{\partial \delta\left(k_r^2\right)}{\partial r} = \frac{-\rho_s^2}{\left(1 + k_{\perp 0}^2 \rho_s^2\right)} L F \frac{\partial \delta\left(k_r^2\right)}{\partial r} . \qquad (18)$$

where L is the scale length of large-scale radial variation of F_0 , i.e., $\Gamma/n = -LF$. Substituting Eq.(15) into Eq.(18),

$$\tilde{F} = \frac{\rho_s^2 k_{\perp 0}^2}{\left(1 + k_{\perp 0}^2 \rho_s^2\right)} L F \sin 2\theta \int^t dt \frac{\partial^2}{\partial r^2} \tilde{V}_{E\times B} . \qquad (19)$$

We have the $\sin\theta$ dependent term of $\tilde{F}\left(r, \theta, t\right)$ as

$$\tilde{F}(r,\theta,t) = \langle 2\,F\cos\theta\rangle \sin\theta\, L\left(\frac{k_{\perp 0}^2\rho_s^2}{\left(1+k_{\perp 0}^2\rho_s^2\right)}\int^t \mathrm{d}t\,\frac{\partial^2}{\partial r^2}\tilde{V}_{E\times B}\right). \tag{20}$$

where $\langle F\cos\theta\rangle$ is the poloidal average of $F\cos\theta$ and An evaluation $\langle F\cos\theta\rangle \simeq F_0/2$ if one uses the estimate $F_1 \simeq F_0$. Equations (6), (7) and (20) give

$$\tilde{N} = \frac{1}{\omega(\omega+i\nu) - c_s^2 k_\parallel^2}\left(\frac{-1}{r}\frac{\partial F_1}{\partial\theta} + i(\omega+i\nu)\left(\frac{2}{R} + i\,u_{\mathrm{WH}}\frac{K_r^2}{\omega}\right)\sin\theta\right)\tilde{V}_{E\times B}\,, \tag{21}$$

where

$$u_{\mathrm{WH}} = \frac{k_{\perp 0}^2\rho_s^2}{\left(1+k_{\perp 0}^2\rho_s^2\right)}F_0 L \tag{22}$$

is the coefficient that denotes the magnitude of the W-H effect. (Note that u_{WH} becomes factor 2 larger than Eq.(22) in the case of quasi-linear diffusion.) The dispersion relation is given as

$$\omega = \frac{2c_s^2}{R^2}\,\frac{(\omega+i\nu)\left(1 + i\,\frac{R u_{\mathrm{WH}}}{2}\frac{K_r^2}{\omega}\right) - i\gamma_{\mathrm{asym}}}{\omega(\omega+i\nu) - c_s^2 k_\parallel^2}\,. \tag{23}$$

In the limit of weak growth rate or damping rate,
$\omega_{\mathrm{GAM}} \gg \mathrm{Im}(\omega),\ \gamma_{\mathrm{asym}},\ \nu,\ R u_{\mathrm{WH}} K_r^2 \omega_{\mathrm{GAM}}^{-1}$, the growth rate is given as

$$\mathrm{Im}(\omega) = \frac{q^2}{1+2q^2}\left(\frac{R u_{\mathrm{WH}} K_r^2}{2} - \gamma_{\mathrm{asym}}\right) - \frac{\mu_\parallel K_r^2}{2\left(1+2q^2\right)}\,. \tag{24}$$

The order of magnitude estimate for the value of $R\,u_{\mathrm{WH}}$ gives a relation with γ_{asym} as $R u_{\mathrm{WH}} L^{-2} \sim \gamma_{\mathrm{asym}}$. Compared to the static driving term γ_{asym} , the W-H effect has dependence on K_r^2 so that this destabilizing term can be larger for short-wave length modes, $K_r^2 L^2 \gg 1$. The turbulent shear viscosity term, the last term in RHS of Eq.(24), has the same dependence on K_r^2 and on turbulence amplitude as the $R u_{\mathrm{WH}} K_r^2$ term. If the condition

$$q^2 R u_{\mathrm{WH}} > \mu_{\|} \tag{25}$$

holds, the drive by the dynamic shearing exceeds the stabilization by the shear viscosity, and the excitation of GAM is expected to occur. An order estimate of the divergence of the particle flux may be employed as $F_0 \simeq DL^{-2}$ where D is a turbulent transport coefficient. Substituting this estimate in to Eq.(21), one has

$$u_{\mathrm{WH}} = k_{\perp 0}^2 \rho_s^2 \, D \, L^{-1} \tag{26}$$

for $k_{\perp 0}^2 \rho_s^2 < 1$. Combining Eqs.(25) and (26), one has the condition for GAM instability by W-H mechanism as

$$\frac{q^2 R}{L} k_{\perp 0}^2 \rho_s^2 > \frac{\mu_{\|}}{D} \, . \tag{27}$$

Noting that $\mu_{\|}/D$ is close to unity, the GAM is induced by turbulence for the higher q cases that satisfy $q^2 > q_c^2 \sim LR^{-1} k_{\perp 0}^{-2} \rho_s^{-2}$. If one considers the application to the ion-temperature-gradient (ITG) driven turbulence, the microfluctuations are excited in the range of $k_{\perp 0}^2 \rho_s^2 \sim 0.1$. The GAM excitation is expected to occur when

$$q^2 > q_c^2 \sim 10 \, \frac{L}{R} \tag{28}$$

holds.

Summary

We have studied the excitation of GAM by micro turbulence in toroidal plasmas. The GAM can be excited by the toroidicity induced turbulence through dynamic shearing by GAM. The competition between the damping via turbulent shear viscosity is examined. It is found that the GAM is easily excited in high-q cases that satisfy $q^2 > q_c^2 \sim LR^{-1} k_{\perp 0}^{-2} \rho_s^{-2}$. For the case of the ITG turbulence, the instability condition is estimated as Eq.(28). This condition is more easily satisfied near the edge of higher-q discharges. It should be noticed that the condition of GAM excitation by W-H mechanism is derived here by using simple fluid equations with isothermal dynamics. The qualitative feature of destabilization by dynamic shearing works even if the model of temperature perturbation is improved, but the condition like Eq.(28) will be quantitatively modified. The rate of parallel flow damping $\mu_{\|} K_r^2$ is affected by the model of parallel ion heat response [4].

It might be also worthwhile to note that the coupling with sound wave is unimportant for the GAM instability by the W-H mechanism. If the coupling with sound wave is switched-off, the $k_{\|}^2 c_s^2$ term is neglected in Eq.(23). (This approximation, i.e., $\omega^2 >> k_{\|}^2 c_s^2$ holds in the denominator of Eq.(23), requires $q^2 >> 1$ to hold.) The term $\gamma_{\rm asym}$ is less important compared to the W-H mechanism for standard parameters and can be neglected in Eq.(23). Under such simplifications, Eq.(23) reduces to $\omega^2 = \omega_{\rm GAM}^2 \left(1 + i\, R u_{\rm WH} K_r^2 / 2\omega\right)$, yielding $\mathrm{Im}\left(\omega\right) = R u_{\rm WH} K_r^2/4$. This is essentially the hgih-q limit of Eq.(24). It seems consistent with observation in direct simulations, i.e., the runs for the edge parameters show that the sound wave is completely unimportant for $q > 3$, and switching it off does not change the result of GAM excitation.

It has been shown that the zonal flow is damped by the shear viscosity in the case of high q plasmas [4]. The result in this article shows that the change from zonal flow to GAM is expected to occur as q increases. The effects of GAM for suppression of turbulence are weaker than that of zonal flows [2, 6]. The change from zonal flow to GAM near edge is attributed to one of the reasons of strong turbulence near plasma edge.

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