Dynamics of low frequency zonal flow driven by geodesic acoustic modes

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Abstract. A new method of controlling turbulent transport is proposed. The method is focused on low frequency zonal flows (ZFs) driven by geodesic acoustic modes (GAMs). In order to describe their dynamics with experimental parameters, the damping rates of GAMs and ZFs are introduced into fluid models. The coupling equations among turbulence, GAMs and ZFs are derived, and the parameter region where ZFs is effectively driven is revealed. It is also shown that it is possible to control ZFs by neutral beam injection (NBI), since the growth rate of GAMs are affected by energetic particles. Furthermore the impacts of GAMs on the partition of heating power is addressed.

1. Introduction

In order to realize the nuclear fusion, plasmas with high ion temperature have to be maintained. There are two ways to increase ion temperature: one is to suppress the transport, and another is to heat bulk ions.

The turbulent transport is strongly affected by low frequency zonal flows (ZFs) and geodesic acoustic modes (GAMs) which are destabilized by micro-scale turbulent fluctuations. In order to understand the turbulent transport, the energy partition among turbulence, ZFs, and GAMs is an essential issue, because the effects of ZFs and GAMs on transport differ. The coupling between turbulence and ZFs, and the one between turbulence and GAMs are investigated in experiments, simulations and theories [1]. The competition among turbulence, ZF and GAM has been numerically studied [2]. Recently, GAMs are shown to be composed of radially-propagating eigenmodes [3]. As a consequence, ZF and GAM interact not only indirectly through turbulence, but also directly [4]. This direct coupling between GAM and ZF affects the partition energy among turbulence, ZF, and GAM. The understanding of dynamics of ZF which is driven by GAMs is required, because the transport can be controlled by changing the amplitude of ZF.

Recently, another novel role of GAMs has been pointed out, namely that ions can be heated by GAMs [5]. GAMs can also be excited by energetic particles [6]. The energy exchange rate from energetic particles to bulk ion via GAMs can contribute to substantial ion heating. A large amplitude of GAMs has been observed in Large Helical Device (LHD) plasma which is sustained by Neutral Beam Injection (NBI) [7] and the impacts of these GAM are discussed here.

We investigate the nonlinear dynamics of ZFs driven by GAMs. The outline of this paper is as follows. In Sec. 2, model equations for analyses are described. In Sec. 3, the nonlinear ZFs dynamics is described. In Sec. 4, the discussion on the role of GAMs on heating partition is described. Summary is given in Sec. 5.

2. Models

The fluid equations are employed, and the damping rates of GAMs and ZFs are introduced as parameters, in which kinetic processes are taken into account. A focus is placed upon the response of Reynolds stress, which is evaluated from the action conservation equation. Then, the quadripartite coupling equations for turbulence, radially inward and outward propagating GAMs and ZFs driven by GAMs are derived.

2.1. Basic equations

We consider a high aspect ratio, circular cross section toroidal plasma. The magnetic field is given as

$$\mathbf{B} = \frac{B_0}{1+\varepsilon\cos\theta}\left(\mathbf{e}_\xi + \frac{\varepsilon}{q}\mathbf{e}_\theta\right), \qquad (1)$$

where B_0 is the strength of magnetic field, $\mathbf{e}_\xi, \mathbf{e}_\theta$ express the unit vectors in toroidal and poloidal directions, q is the safety factor, and ε is the inverse aspect ratio ($\varepsilon << 1$). The fluid equations in this magnetic field configuration are written as

$$\frac{\partial U}{\partial t} - \frac{2c_s^2}{R}\nabla_r \oint \frac{d\theta}{2\pi} n \sin\theta + \gamma(\omega)U = -\nabla_r(\mathbf{v}\cdot\nabla)\mathrm{v}_\theta \qquad (2)$$

$$\frac{\partial \mathrm{v}_\parallel}{\partial t} - \mu_\parallel \nabla_\perp^2 \mathrm{v}_\parallel + c_s^2 \nabla_\parallel n = -\mathrm{v}_\theta \nabla_\theta \mathrm{v}_\parallel - \mathrm{v}_\parallel \nabla_\parallel \mathrm{v}_\parallel \qquad (3)$$

$$\frac{\partial n}{\partial t} - \frac{2}{R}\nabla_r^{-1} U \sin\theta + \nabla_\parallel \mathrm{v}_\parallel = -\mathrm{v}_\theta \nabla_\theta n + \frac{2}{R}\mathrm{v}_\theta n \sin\theta - \mathrm{v}_\parallel \nabla_\parallel n \qquad (4)$$

Here, U, $\mathbf{v}, n$ are the toroidal component of vorticity (magnetic surface averaged), velocity field, and the perturbed density normalized by the equilibrium density, respectively. The plasma major radius is denoted by R, the sound velocity is c_s, and the parallel turbulent viscosity is expressed by $\mu_\parallel$, which is treated as a parameter. The damping rate is introduced as $\gamma(\omega)$ in the vorticity equation. As for GAMs, the Landau damping [8] and the effect of energetic particles are considered. As for ZFs, the ion-ion collisional damping rate is introduced. $\gamma(\omega)$ is introduced here as

$$\gamma(\omega_G) = -\delta_h + \frac{\mathrm{v}_T\sqrt{\pi}}{2R} q \left\{ \left(\Omega_G^4 + (1+2\tau_e)\Omega_G^2\right) e^{-\Omega_G^2} + \frac{q_r^2 \rho_T^2 q^2}{4} e^{-\Omega_G^2/4} \left(\frac{\Omega_G^6}{128} + \frac{1+\tau_e}{16}\Omega_G^4 \right) \right\} \qquad (5)$$

$$\gamma(\omega = 0) = \frac{n_i e^4 \ln\Lambda}{12\pi^{3/2}\varepsilon_0^2 m_i^{1/2} T_i^{3/2}} \qquad (6)$$

In Eq. (5), δ_h is the driving term from energetic particles [5], v_T is the ion thermal velocity, ρ_T is the ion gyro-radius calculated by the ion thermal velocity, q is the safety factor, τ_e is the electron temperature normalized by ion temperature, Ω_G is defined as $q\sqrt{7/4+\tau_e}$, and q_r is the radial wavenumber of GAMs.

For the transparency of argument, we focus on Reynolds stress among other nonlinear terms, which are less effective on generating quasi-modes in the present circumstance [9]. As for parallel velocity and density, only their linear responses to the vorticity are introduced to the dynamics of GAMs and ZFs. Eliminating n and $\mathrm{v}_\parallel$ from Eqs. (2)-(4), the equation is obtained as

$$\frac{\partial U}{\partial t} - \frac{2c_s^2}{R^2}\left(-i\omega + \frac{c_s^2 \nabla_\parallel^2}{i\omega + \mu_\parallel \nabla_\perp^2} \right) U + \gamma(\omega)U = -\nabla_r(\mathbf{v}\cdot\nabla)v_\theta . \qquad (7)$$

GAMs are assumed to have counter-propagating components with monochromatic wavenumber,

$$U = \left(U_{\omega,q_r} e^{-i\omega t} + U_{-\omega,q_r} e^{i\omega t} \right) e^{iq_r r} + c.c. \qquad (8)$$

Here, we choose q_r and ω to be positive. We consider spatial separation between the scales of equilibrium density or temperature L_T and GAM wavelength scale, $q_r L_T >> 1$, and the choice of sign of q_r does not change the physics of the solution.

2.2. Coupling equations among turbulence, GAMs and ZFs

The Reynolds stress is evaluated from the action conservation equation. The model for the Reynolds stress is given as [1]

$$\left\langle \nabla_r (\mathbf{v}\cdot\nabla)\mathrm{v}_\theta \right\rangle = \frac{q_r^2}{B^2}\int \frac{k_r k_\theta}{(1+k_\perp^2\rho_s^2)^2} N_k d^2k. \tag{9}$$

Here, the brackets denote time averaging over a time scale much longer than the characteristic time scale of turbulence. The turbulence wavenumber is denoted as k_r, k_θ and $k_\perp$. The action of turbulence is denoted by $N_k = (1+k_\perp^2\rho_s^2)^2|\phi_k|^2$. The action which is modulated by GAMs has been obtained in [10]. The quasi-linear response to GAMs is

$$N^{(1)}_{\omega,q_r} = k_\theta R(\omega,q_r) U_{\omega,q_r}\frac{\partial N^{(0)}}{\partial k_r}, \tag{10}$$

where the response function is defined as $R(\omega,q_r) = i\left(\omega - q_r v_{gr} + i\Delta\omega_k\right)^{-1}$. Here, $\mathbf{v}_g$ is the group velocity and $\Delta\omega_k$ is the nonlinear damping rate of turbulence, respectively. The higher order nonlinearity of the action can be expressed as

$$N^{(2m+1)}_{\omega,q_r} = k_\theta R(\omega,q_r)\left(U^*_{\omega,q_r}\frac{\partial N^{(2m)}_{2\omega,2q_r}}{\partial k_r} + U^*_{-\omega,q_r}\frac{\partial N^{(2m)}_{0,2q_r}}{\partial k_r}\right). \tag{11}$$

Here, U^* denotes the complex conjugate of U. The GAM component appears in the odd-order (2m+1), and the higher harmonics and ZF components appear in the even-order (2m), where m is an integer. ZF component is produced by the nonlinear coupling of radially-counter-propagating GAMs. We neglect the components with high wavenumber such as $(3\omega,3q_r),(0,4q_r),\cdots$, since the higher wavenumber components are strongly damped. Truncating the nonlinear terms up to the third order, Reynolds stress is calculated by using Eq. (9).

Here, we assume the strong turbulence regime. The nonlinear decorrelation rate of turbulence can be written as

$$\Delta\omega = \omega_* \hat{\phi}, \tag{12}$$

where $\omega_* \equiv k_\theta T/eBL_n$ is the drift wave frequency, and $\hat{\phi} \equiv k_\perp^2 L_n e\phi/k_\theta T$ is the normalized drift wave potential. Using Eq. (12), the coupling equations among turbulence, GAMs and ZFs are obtained as

$$\frac{\partial\hat{\phi}^2}{\partial\hat{t}} = \hat{\gamma}_L\hat{\phi}^2 - \hat{\omega}_*\hat{\phi}^3 - 2\eta(\xi^2+\hat{\phi}^2)^{-1}\hat{\phi}^4\left(\left|\hat{U}_P\right|^2 + \left|\hat{U}_M\right|^2\right) - 2\eta\hat{\phi}^2\left|\hat{U}_Z\right|^2 \tag{13a}$$

$$\frac{\partial\hat{U}_P}{\partial\hat{t}} = -i\left(1+\frac{1}{2q^2}\right)\hat{U}_P + \eta(\hat{\phi}+i\xi)^{-1}\hat{\phi}^3\hat{U}_P - \frac{\eta\mu}{2(1+2q^2)}\hat{\phi}^2\hat{U}_P - \hat{\gamma}(\omega_G)\hat{U}_P$$
$$- \eta s(\hat{\phi}+i\xi)^{-2}\left\{(\hat{\phi}+2i\xi)^{-1}\hat{\phi}^3\left|\hat{U}_P\right|^2 + 2\hat{\phi}^2\left|\hat{U}_M\right|^2\right\}\hat{U}_P \tag{13b}$$

$$\frac{\partial\hat{U}_M}{\partial\hat{t}} = i\left(1+\frac{1}{2q^2}\right)\hat{U}_M + \eta(\hat{\phi}-i\xi)^{-1}\hat{\phi}^3\hat{U}_M - \frac{\eta\mu}{2(1+2q^2)}\hat{\phi}^2\hat{U}_M - \hat{\gamma}(\omega_G)\hat{U}_M$$
$$- \eta s(\hat{\phi}-i\xi)^{-2}\left\{(\hat{\phi}-2i\xi)^{-1}\hat{\phi}^3\left|\hat{U}_M\right|^2 + 2\hat{\phi}^2\left|\hat{U}_P\right|^2\right\}\hat{U}_M \tag{13c}$$

$$\frac{\partial\hat{U}_Z}{\partial\hat{t}} = 4\eta\hat{\phi}^2\hat{U}_Z - 4\eta\mu(1+2q^2)\hat{\phi}^2\hat{U}_Z - \hat{\gamma}(\omega=0)\hat{U}_Z - \alpha\eta\hat{\phi}^2(\hat{\phi}+i\xi)^{-1}\hat{U}_P\hat{U}_M, \tag{13d}$$

where γ_L is the linear growth rate of turbulence. The time $\hat{t}$ is normalized by GAM frequency as $\hat{t} = \omega_G t$, where ω_G is the GAM frequency defined as $\omega_G \equiv \sqrt{2}c_s/R$. The vorticity is normalized by drift wave frequency as $\hat{U} = U/\omega_*$. The subscripts P, M and Z indicates $(\omega_G, q_r), (-\omega_G, q_r)$ and $(0, 2q_r)$, which are radially outward and inward propagating GAMs and ZF driven by GAMs, respectively. The coefficients $\eta, \xi, \mu, s, \alpha$ are defined as

$$\eta = \frac{q_r^2}{B^2}\frac{R}{\sqrt{2}c_s}\int \frac{k_\theta^2}{(1+k_\perp^2\rho_s^2)^2}\frac{N^{(0)}}{\Delta\omega}d^2k, \quad \xi = \omega_G/\omega_*, \quad \mu = \mu_\parallel q_r^2 \frac{R}{\eta\hat{\phi}^2\sqrt{2}c_s}$$

$$s = \frac{4k_\theta^2\rho_s^2}{1+k_\perp^2\rho_s^2}, \quad \alpha = 8k_\theta/k_r.$$

The coupling equations Eqs. (13a)-(13d) are extended from those in [4] to include the effects of phase information and of temporal evolution of turbulence energy.

3. Nonlinear dynamics of GAMs and ZFs

We consider the situation where GAMs are driven by turbulence, but ZFs are not excited by quasilinear process of turbulence. The stationary state is described in Sec. 3.1. The nonlinear response of ZFs to a pulse perturbation of GAMs without considering the effect of energetic particles is shown in Sec. 3.2. The controllability of ZFs by NBI modulation is given in Sec. 3.3.

3.1. Stationary state

The stationary state is determined by replacing the time derivative by zero in Eqs. (13a)-(13d). It is worth noting that radially inward and outward propagating GAMs can not exist together in the stationary state. ZFs driven by GAMs can not exist in the stationary state because the driving term becomes zero in Eq. (13d).

3.2. Dynamic response of ZFs without energetic particles effect

In the first step of analysis, the nonlinear response of ZFs to a pulse perturbation of GAMs is studied. This analysis is a prototypical model for the experimental observations where GAMs are excited in an intermittent manner [11]. The time evolutions of the response of fields are shown in Fig. 1. In this study, the pulse of inward-propagating GAM ($\hat{U}_M$) is imposed on a stationary state of outward-propagating GAM ($\hat{U}_P$). After the pulse is applied, $\hat{U}_M$ monotonically decreases, and the fields are relaxed to the initial stationary state. During the time when $\hat{U}_M$ have finite amplitude, $\hat{U}_M$ and $\hat{U}_P$ coexist, and the ZF is excited. ZF affects the turbulence, and the amplitude of $\hat{U}_P$ decreases. In order to understand the behavior of ZF analytically, the linearization of deviation from the stationary state is carried out in the coupling equations Eqs. (13a)-(13d). The deviation from the stationary state is written as

$$\hat{U}_P = \hat{U}_{P,s} + \Delta\hat{U}_P, \hat{U}_M = \Delta\hat{U}_M, \hat{U}_Z = \Delta\hat{U}_Z, \hat{\phi} = \hat{\phi}_s + \Delta\hat{\phi}, \qquad (14)$$

where $U_{P,s}$ denotes the saturated U_P. The linearized coupling equations are obtained as

$$\frac{\partial\Delta\hat{\phi}}{\partial\hat{t}} = -\hat{\gamma}_\phi\Delta\hat{\phi} + \hat{C}_{\phi P}(\hat{U}_{P,s}^*\Delta\hat{U}_P + \hat{U}_{P,s}\Delta\hat{U}_P^*) \qquad (15a)$$

$$\frac{\partial\Delta\hat{U}_P}{\partial\hat{t}} = -i\hat{\omega}_P\Delta\hat{U}_P + \hat{C}_{P\phi}\Delta\phi \qquad (15b)$$

$$\frac{\partial\Delta\hat{U}_M}{\partial\hat{t}} = (-i\hat{\omega}_M - \hat{\gamma}_M)\Delta\hat{U}_M \qquad (15c)$$

$$\frac{\partial\Delta\hat{U}_Z}{\partial\hat{t}} = -\hat{\gamma}_Z\Delta\hat{U}_Z + \hat{C}_Z\hat{U}_{P,s}\Delta\hat{U}_M. \qquad (15d)$$

Here, the coupling coefficients are

$$\hat{\gamma}_{\phi} = -\hat{\gamma}_L + \frac{3}{2}\hat{\omega}_* \hat{\phi}_s^2 + 2\eta(\hat{\phi}_s^2 + \xi^2)^{-1}\left\{2 - \hat{\phi}_s^2(\hat{\phi}_s^2 + \xi^2)^{-1}\right\}\hat{\phi}_s^2 \left|\hat{U}_{P,s}\right|^2$$

$$\hat{C}_{\phi P} = \eta(\hat{\phi}_s^2 + \xi^2)^{-1}\hat{\phi}_s^3$$

$$\hat{\omega}_P = \left(1 + \frac{1}{2q^2}\right) + \eta\xi(\hat{\phi}_s^2 + \xi^2)^{-1}\hat{\phi}_s^3 - 2s\eta\xi\hat{\phi}_s^4(\hat{\phi}_s^2 + \xi^2)^{-2}(\hat{\phi}_s^2 + 4\xi^2)^{-1}(2\hat{\phi}_s^2 - \xi^2)\left|\hat{U}_{P,s}\right|^2$$

$$\hat{C}_{P\phi} = \eta(\hat{\phi}_s + i\xi)^{-1}\hat{\phi}_s^2\left\{3 - \hat{\phi}_s(\hat{\phi}_s + i\xi)^{-1}\right\}\hat{U}_{P,s}$$

$$- s\eta(\hat{\phi}_s + i\xi)^{-2}(\hat{\phi}_s + 2i\xi)^{-1}\hat{\phi}_s^3\left\{3 - 2\hat{\phi}_s(\hat{\phi}_s + i\xi)^{-1} - \hat{\phi}_s(\hat{\phi}_s + 2i\xi)^{-1}\right\}\left|\hat{U}_{P,s}\right|^2\hat{U}_{P,s}$$

$$\hat{\omega}_M = -\left(1 + \frac{1}{2q^2}\right) - \eta\xi(\hat{\phi}_s^2 + \xi^2)^{-1}\hat{\phi}_s^3 - 4s\eta\xi\hat{\phi}_s^3(\hat{\phi}_s^2 + \xi^2)^{-2}\left|\hat{U}_{P,s}\right|^2$$

$$\hat{\gamma}_M = s\eta(\hat{\phi}_s^2 + \xi^2)^{-2}\hat{\phi}_s^2\left\{2(\hat{\phi}_s^2 - \xi^2) - (\hat{\phi}_s^2 + 4\xi^2)^{-1}(5\hat{\phi}_s^2 - \xi^2)\hat{\phi}_s^2\right\}\left|\hat{U}_{P,s}\right|^2$$

$$\hat{\gamma}_Z = \hat{\gamma}(\omega = 0) - 4\eta\left\{1 - (1 + 2q^2)\mu\right\}\hat{\phi}_s^2$$

$$\hat{C}_Z = -\alpha\eta\hat{\phi}_s^2(\hat{\phi}_s + i\xi)^{-1}.$$

The analytical solution of the responses of ΔU_M and ΔU_Z can be obtained as

$$\Delta\hat{U}_M = \Delta\hat{U}_{M,0}e^{-i\hat{\omega}_M\hat{t} - \hat{\gamma}_M\hat{t}} \qquad (16a)$$

$$\Delta\hat{U}_Z = \frac{\hat{C}_Z\left|\hat{U}_{P,s}\right|\Delta\hat{U}_{M,0}}{-i(\hat{\omega}_P + \hat{\omega}_M) - \hat{\gamma}_M + \hat{\gamma}_Z}e^{-\hat{\gamma}_Z\hat{t}}\left(e^{-i(\hat{\omega}_P + \hat{\omega}_M)\hat{t} - \hat{\gamma}_M\hat{t} + \hat{\gamma}_Z\hat{t}} - 1\right). \qquad (16b)$$

The characteristic time of this relaxation process is dominated by γ_M and γ_Z. The life time of the excited ZF $\hat{\tau}_{ZF}$ can be estimated as

$$\hat{\tau}_{ZF} \sim \hat{\gamma}_M^{-1} + \hat{\gamma}_Z^{-1}. \qquad (17)$$

The dependence of the maximum amplitude of ZFs on the pulse height is shown in Fig. 2. The driven ZFs increases almost linearly for a small pulse. When the pulse height becomes large, the large ZFs are excited. ZF with large amplitude strongly affects turbulence, and then the numerical and analytical solutions deviate. The safety factor dependence of ZFs is shown in Fig. 3. The amplitude of ZFs decreases with the safety factor, since the damping rate of ZF is an increasing function of the safety factor. Substantial ZFs can be induced by radially-counter-propagating GAMs.

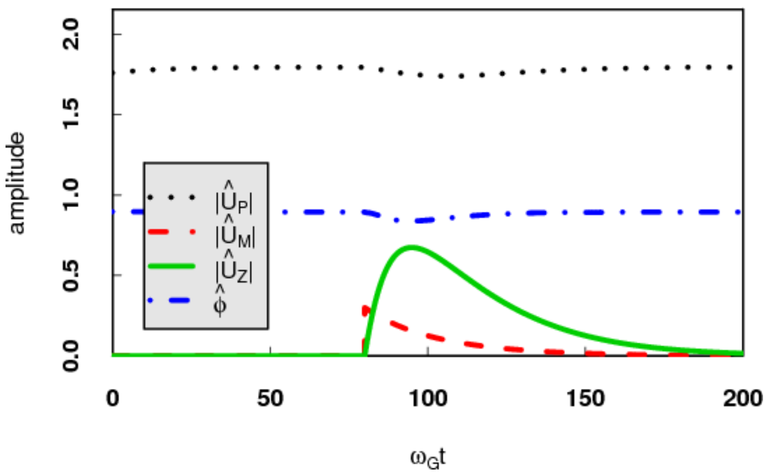

Fig. 1: Response of turbulence, GAMs and ZF to pulse perturbation of $\hat{U}_M$
The parameters used here are $q = 3, \tau_e = 1, q_r\rho_T = 0.1, \eta = 0.1, \mu = 0.1,$

$\xi = 0.25, s = 0.24$.

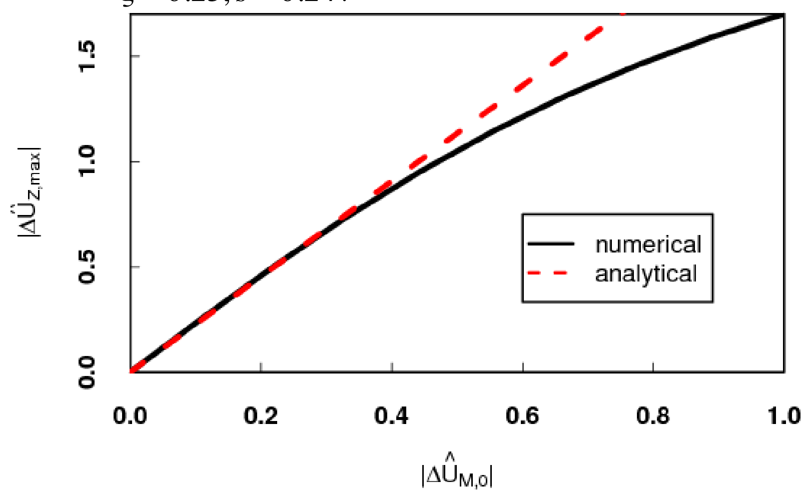

Fig. 2: Initial pulse dependence of driven ZFs
Black line shows the result of numerical calculation of Eqs. (13a)-(13d). Red line shows the result from analytical expression Eq. (16b).

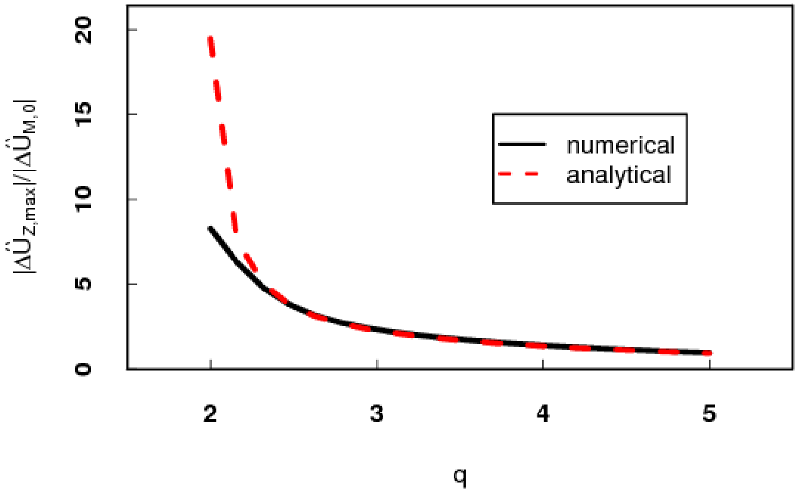

Fig. 3: Safety factor dependence of driven ZFs
Black line shows the result of numerical calculation of Eqs. (13a)-(13d). Red line shows the result from analytical expression Eq. (16b).

3.3. Controlling ZFs by NIB modulation

The controllability of ZFs by NBI is discussed by introducing effects of energetic particles on the linear growth rate of GAMs. Here, we consider the situation when the power of NBI is modulated temporally, so that the energetic particle effect δ_h can be written as

$$\hat{\delta}_h = \delta_{ex} \sin^2(\Omega_{ex}\hat{t} + \Theta). \qquad (18)$$

Here, δ_{ex} is the growth rate of GAMs by energetic particles, which is related to the injection power of NBI, Ω_{ex} is the modulation frequency of NBI, and Θ is the phase delay of pulse injection to modulation period. The analytical expression for the linear response ΔU_Z is obtained as

$$\Delta\hat{U}_Z = \hat{C}_Z \left|\hat{U}_{P,s}\right| \Delta\hat{U}_{M,0} e^{-\hat{\gamma}_Z \hat{t}} \int \exp\left[\left(-i\hat{\omega}_M - \hat{\gamma}_M + \hat{\gamma}_Z + \frac{\delta_{ex}}{2}\right)\hat{t} - \frac{\delta_{ex}}{4\Omega_{ex}}\left\{\sin(2\Omega_{ex}\hat{t} + \Theta) - \sin\Theta\right\}\right] d\hat{t}. \quad (19)$$

The amplitude of ZFs can be controlled by NBI modulation. When the relation between the onset of pulse of ΔU_M and the phase of the modulation satisfies the condition $\Theta \sim \pi/2$,

ΔU_Z becomes larger by several percent, (in the case that $\delta_{ex} = 0.03, \Omega_{ex} = 0.1, \Theta = 0.6\pi$), which is shown in Fig. 4.

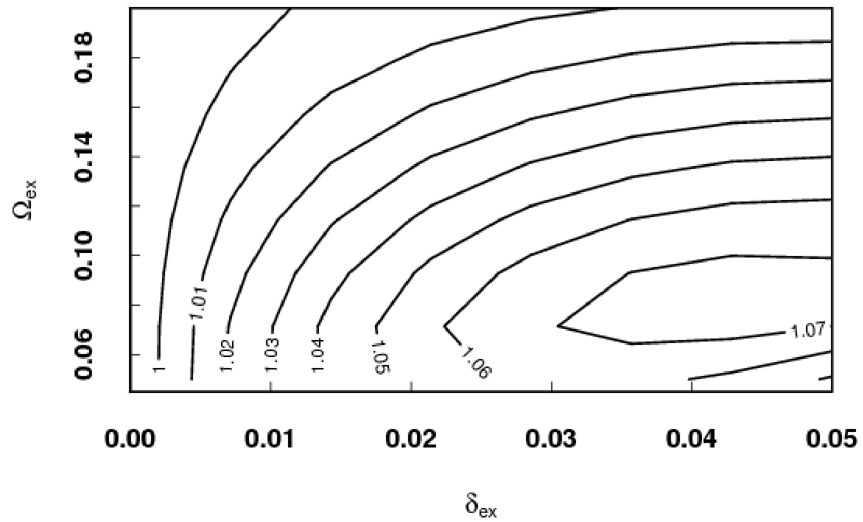

Fig. 4: Control parameters dependence of maximum amplitude of ZF normalized by $\Delta\hat{U}_{Z,\mathrm{ref}}$. Here $\Delta\hat{U}_{Z,\mathrm{ref}}$ is the result of $\delta_{ex} \to 0$.

4. Discussion on the role of GAMs on heating partition

GAMs have another important characteristic for the dynamics of main plasmas. GAMs transfer the energy to bulk ions through Landau damping. Here, we show the characteristics of ion heating by GAMs that are excited by energetic ions. The heating rate of main ions by GAMs $P_{G\to i}$ is given within the framework of quasi-linear theory as [5]

$$P_{G\to i} = \frac{n_i\sqrt{m_i T_i}q_r^2|\phi_G|^2}{RB^2}\hat{\gamma}_G. \qquad (20)$$

Here, ϕ_G is the electrostatic potential of GAMs, and $\hat{\gamma}_G$ is the damping rate of GAM by the Landau damping which is normalized by v_T/R. The heating efficiency of GAMs is proportional to the inverse square of the strength of magnetic field and the inverse of the plasma major radius, and it strongly depends on the safety factor and the electron temperature normalized to the ion temperature. It is proportional to the square of the amplitude and the radial wavenumber of GAMs, respectively. Since the heating efficiency of GAMs is determined by the square of GAMs amplitude, the nonlinear dynamics of GAMs is important. The order of magnitude estimate for the impacts of GAM channeling is discussed here, taking an example of LHD experiments. The beam driven GAMs are observed in LHD plasma which is sustained by NBI heating of $P_{\mathrm{NBI}} = 400$ [kW/m^{-3}]. The observed GAMs have a voltage around $\phi_G = 1$ [kV] [6]. The plasma parameters are R=3.9 [m], a=0.65 [m], $T_e = T_i = 3$[keV], $n_i = 1.5\times 10^{19}$ [m^{-3}], B=3[T]. The radial electric field of GAMs is estimated as $E_r = |q_r||\phi_G| \sim 10^3|q_r|$ [V/m]. The ratio between ion heating effect by GAMs and electron heating effect by NBI can be evaluated as $P_{G\to i}/P_{\mathrm{NBI}}$~0.1, where we assume $q_r \sim 10$ [m^{-1}]. The bulk ion heating by GAMs has a considerable effect on ion energy in real plasmas.

5. Summary

We investigate the nonlinear dynamics of ZFs driven by GAMs, where GAMs are driven by turbulence. Based on the fluid model with the introduction of Landau damping and energetic particle effect for GAMs, and of ion-ion collisional damping for ZFs, the coupling equations among turbulence, GAMs and ZFs are derived. ZFs driven by GAMs are generated effectively at lower safety factor. Substantial ZFs can be induced by radially-counter-propagating GAMs. The pulse response of ZFs with NBI modulation is investigated. When the onset of pulse of GAM and the phase of the modulation corresponds, the amplitude of ZF driven by GAMs becomes larger. This study contributes to the search for the improved confinement of plasmas. Another impact of GAMs on confined plasmas is also discussed. GAMs can affect the partition of injection power between main ions and electrons. These analyses illuminate the importance of nonlinear dynamics of ZFs, GAMs and turbulence.

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Appendix

[1] P. H. Diamond, S.-I. Itoh, K. Itoh, T. S. Hahm, Plasma Phys. Control. Fusion **47**, R35 (2005).
[2] N. Miyato, Y. Kishimoto, Phys. of Plasmas **11**, 5557 (2004). K. Miki, Y. Kishimoto, N. Miyato, J. Q. Li, Phy. Rev. Lett. **99**, 145003 (2007).
[3] K. Itoh, et. al., PFR 1, 037 (2006).
[4] M. Sasaki, K. Itoh, A. Ejiri, Y. Takase, Plasma Phys. Control. Fusion **51**, 085002 (2009).
[5] M. Sasaki, K. Itoh, S.-I. Itoh, to be submitted to Plasma Phys. Control. Fusion.
[6] H. L. Berk, T. Zhou, Nucl. Fusion **50**, 035007 (2010).
[7] T. Ido, et. al., European Physical Society 37th, I4,112 (2010).
[8] H. Sugama, T. H. Watanabe, Phys. of Plasmas **13**, 012501 (2006).
[9] M. Sasaki, K. Itoh, A. Ejiri, Y. Takase, Phys. of Plasmas **16**, 022306 (2009).
[10] K. Itoh, et. al., Phys. of Plasmas **12**, 062303 (2005).
[11] A. Fujisawa, et. al., Nucl. Fusion, **47**, S718 (2007).