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Coherent structure of zonal flow and nonlinear saturation

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Nonlinear structure of the zonal flow in toroidal plasma is investigated. It is found that the turbulent drive of zonal flow starts to decrease at high velocity shear of zonal flow. By this mechanism, the zonal flow evolves into a stable stationary structure in turbulent plasmas. Flow velocity and radial scale length are obtained. Analytic model explains simulational observations.

keywords: zonal flow, plasma turbulence, collisionless saturation, coherent structure, wave kinetic equation

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Turbulence of the magnetized and high-temperature plasmas is subject of intense studies. This is a typical example of structure formation in systems far from equilibrium. It has been pointed out that micro-fluctuations induces semi-micro perturbation of flow [1]. In particular, zonal flows, which are constant on a toroidal magnetic surface but change rapidly in the radial direction [2,3], are known to be induced by micro turbulence in tokamaks and play important role in determining turbulence and turbulent transport.

One of the key issues is the mechanism that regulates the structure of the induced zonal flows. For instance, it remains unanswered whether induced small-scale zonal flows evolve into a large-scale flows. The saturation mechanisms of zonal flow velocity in the collisionless plasma are discussed: e.g., binary collisions play a role [3], and possibility of secondary instabilities has been pursued [4-6]. These are important in the wider field of physics. For instance, they are related to the problem whether the small-scale dynamo magnetic field can evolve into large scale magnetic field or not [7]. In this direction of research, ref.[8] studied theoretically the possibility that the zonal flows evolve into a kink-soliton like structure, and the condensation of micro modes into global modes has been studied by direct nonlinear simulations (DNS) [9]. However, the accessibility to the particular solution in [8] is not clear. It is known that the toroidal geometry is crucial in determining the structure of turbulence and flows [10,11]. In many DNS [12], it has been observed that, in collisionless limit, the turbulence is often completely quenched by the induced zonal flow. In toroidal plasmas, the return flow along the magnetic field line exists, and a stationary state of turbulence and zonal flow has been found in DNS [11].

In this article, we analyze the nonlinear state of zonal flow in toroidal plasmas. It is shown that, in turbulent toroidal plasmas, the perpendicular transport of perpendicular momentum shows nonlinear saturation on the zonal flow shear, while the perpendicular transport of parallel momentum does not. That is, the drive of zonal flow starts to decrease at high velocity but the damping by turbulent viscosity of parallel flow does not. By this, the zonal flow evolves into a nonlinear stationary state. The scale length and saturation level of the zonal flow are analyzed. We first present a theoretical model and then show the verification by comparing with results in DNS.

The growth of the zonal flow in the presence of the drift-wave turbulence has been discussed by use of the time scale separation. The autocorrelation times of the drift wave fluctuations are assumed to be much faster than the evolution time of the zonal flow. In the slow time scale, the evolution of the zonal flow is governed by the equation [8]

$$\frac{\partial}{\partial t}\left(\frac{\partial}{\partial r} V_Z\right) = \frac{\partial^2}{\partial r^2}\frac{c^2}{B^2}\int d^2k \, \frac{-k_\theta k_r}{\left(1 + k_\perp^2 \rho_s^2\right)^2}\bar{N}_k - \gamma_{damp}\frac{\partial}{\partial r} V_Z \qquad (1)$$

where V_Z is the zonal flow velocity, $\bar{N}_k$ is a slow modulation of drift-wave action N_k, which is induced by V_Z, and γ_{damp} denotes the damping rate of zonal flow by other processes. In the slow time scale, N_k satisfies the eikonal equation

$$\frac{\partial}{\partial t} N_k + \frac{\partial \omega_k}{\partial k} \cdot \frac{\partial N_k}{\partial x} - \frac{\partial \omega_k}{\partial x} \cdot \frac{\partial N_k}{\partial k} = 0 \tag{2}$$

A linear response has been obtained from Eq.(2) as $\bar{N}_k^{(1)} = \frac{\partial}{\partial r}\left(k_\theta V_Z\right) R(K, \Omega) \frac{\partial N_k}{\partial k_r}$. Here $R(K, \Omega) = i/\left(\Omega - K\, \partial\omega/\partial k_r + i\Delta\omega_k\right)$ is the response function, $\Delta\omega_k$ is the nonlinear broadening and the zonal flow has a slow dependence as $\exp\left(iKr - i\Omega t\right)$. The first order term gives the diffusion-like form $\gamma_Z = D_{rr}K^2 = -D_{rr}\,\partial^2/\partial r^2$ in Eq.(1) with

$$D_{rr} = \frac{c^2}{B^2} \int \mathrm{d}^2k\, \frac{R(K, \Omega) k_\theta^2 k_r}{\left(1 + k_\perp^2 \rho_s^2\right)^2} \frac{\partial N_k}{\partial k_r}, \tag{3}$$

i.e., the zonal flow growth [8]. The higher order responses with respect to V_Z, $\bar{N}_k^{(2)}$ and $\bar{N}_k^{(3)}$ can be calculated from the relation $\bar{N}_k^{(n)} = k_\theta \left(\partial V_Z/\partial r\right) R(K, \Omega)\, \partial \bar{N}_k^{(n-1)}/\partial k_r$. For a wide spectrum of fluctuations, one has $R(K, \Omega) \to 1/\Delta\omega_k$ and obtains the leading diffusion term of Eq.(3). In such a case, $R(K, \Omega)$ has an approximate symmetry with respect to k_r. The contribution from the second order term is small from the consideration of symmetry, and the next order one comes from the third order term as

$$\bar{N}_k^{(3)} \simeq \left(\frac{\partial}{\partial r} V_Z\right)^3 R(K, \Omega)^3 k_\theta^3 \frac{\partial^3 N_k}{\partial k_r^3}. \tag{4}$$

Substituting Eq.(4) into Eq.(1) with Eq.(3), one obtains the equation for the vorticity $U = \partial V_Z/\partial r$ as

$$\frac{\partial}{\partial t} U = -D_{rr} \frac{\partial^2}{\partial r^2} U + D_3 \frac{\partial^2}{\partial r^2} U^3 - \gamma_{damp} U, \tag{5}$$

$$D_3 = -\frac{c^2}{B^2} \int \mathrm{d}^2k\, \frac{R(K, \Omega)^3 k_\theta^4 k_r}{\left(1 + k_\perp^2 \rho_s^2\right)^2} \frac{\partial^3 N_k}{\partial k_r^3}. \tag{6}$$

Noticing the fact that the spectral function is peaked near $k_r \simeq 0$, the sign in the definition of D_3 is chosen such that D_3 is positive when D_{rr} is positive.

Next, the most unstable wavenumber of zonal flow is considered. The zonal flow growth rate γ_Z does not continue to grow at larger K when the dispersion effect of the drift waves on the zonal flow is introduced. Let us consider the zonal flow excitation in toroidal plasmas. It has been pointed out in ref. [13] that γ_Z is written as

$$\gamma_Z = D_{rr} K^2\left(1 - K^2/K_0^2\right) \tag{7}$$

where K_0^2 represents characteristic scale where the dispersion suppresses the zonal flow instability. An explicit form of K_0^2 is given in [13]. An additional damping mechanism is explained as follows; The $E \times B$ flow in toroidal plasma is associated with the secondary flow. As is shown in [10, 11], the viscous damping of the secondary flow due to toroidicity works as the damping rate of the zonal flow, in addition to the conventional collisional damping. This damping rate is rewritten, when collisional damping is negligibly small, as [10]

$$\gamma_{\text{damp}} = \mu_{\|}\left(1 + 2q^2\right)K^2 \tag{8}$$

where $\mu_{\|}$ is the turbulent shear viscosity for the flow along the field line and q is the safety factor. (The coefficient $1 + 2q^2$ is replaced by $1 + 1.6q^2/\sqrt{\varepsilon}$ in the collisionless limit. This dependence on the collisionality is put aside for a transparency of the argument.) Taking into account Eqs.(7) and (8), Eq.(1) is written in an explicit form as

$$\frac{\partial}{\partial t} U + D_{rr}\left(\frac{\partial^2}{\partial r^2} U + K_0^{-2}\frac{\partial^4}{\partial r^4} U\right) - D_3 \frac{\partial^2}{\partial r^2} U^3 - \mu_{\|}\left(1 + 2q^2\right)\frac{\partial^2}{\partial r^2} U = 0 \ . \tag{9}$$

The zonal flow is excited if the condition $D_{rr} > \mu_{\|}\left(1 + 2q^2\right)$ is satisfied. Both the zonal-flow driving coefficient D_{rr} and the shear viscosity $\mu_{\|}$ are given by the drift wave spectrum N_k . The ratio of $\mu \equiv \mu_{\|}\left(1 + 2q^2\right)D_{rr}^{-1}$ is a function of the spectral shape of drift wave turbulence and geometrical factor such as q . We use normalized variables $x = r/L$, $\tau = t/t_Z$ and $u = U/U_0$, where $L^{-2} = K_0^2\left(1 - \mu\right)$, $t_Z = D_{rr}^{-1} K_0^{-2}\left(1 - \mu\right)^{-2}$ and $U_0^2 = D_{rr} D_3^{-1}\left(1 - \mu\right)$. Equation (9) is rewritten as

$$\frac{\partial}{\partial \tau} u + \frac{\partial^2}{\partial x^2} u - \frac{\partial^2}{\partial x^2} u^3 + \frac{\partial^4}{\partial x^4} u = 0 \ . \tag{10}$$

The short wave length components with $K^2 L^2 > 1$ are stabilized by the higher-order derivative term. The flow is generated in the long wave-length region of $K^2 L^2 < 1$, and the zonal flow energy is saturated by nonlinearity and by higher-order dissipation.

We investigate the case that the flow is generated form the state with small noise level where no net flow exists, $\int dx\, u = 0$. Conservation of total momentum holds for the periodic boundary condition and the flow evolves satisfying the condition $\int dx\, u = 0$. Stationary solution of Eq.(10) in the domain $0 < x < d$, for the periodic boundary condition, is given by an elliptic integral as

$$\int \left(1 - 2u^2 + u^4 - \kappa^2\right)^{-1/2} du = \pm \frac{x}{\sqrt{2}} \tag{11}$$

where κ is an integral constants satisfying $0 \le \kappa < 1$. The integral constant κ is given by the periodicity $\int_{-u_c}^{u_c} \left(1 - 2u^2 + u^4 - \kappa^2\right)^{-1/2} du = d/2\sqrt{2}n$ ($u_c = \sqrt{1 - \kappa}$, and $n = 1, 2, 3, \cdots$). The temporal evolution of Eq. (10) is solved numerically. It is shown that the growth is dominated by the component which has the largest linear growth rate. That is, the integer n is given by the one which is closest to $d/n = 4\sqrt{2}\pi$. Figure 1 illustrates the stable stationary state. The peak value of $u(x)$, is given as $u_c \simeq 0.95$. Compared to a simple sinusoidal function (eigen function of linear operator), the result in Fig.1 has much weaker curvature at the peak and is closer to a piecewise constant function.

The normalized function $u(x)$ is of the order of unity, so that the characteristic values of vorticity and scale length l are given as

$$U_0 = D_{rr}^{1/2} D_3^{-1/2} \left(1 - \mu\right)^{1/2} , \tag{12a}$$

$$l = K_0^{-1} \left(1 - \mu\right)^{-1/2} . \tag{12b}$$

The amplitude of zonal flow is given as $V_0 = l U_0$. From Eqs.(3) and (6), one has an

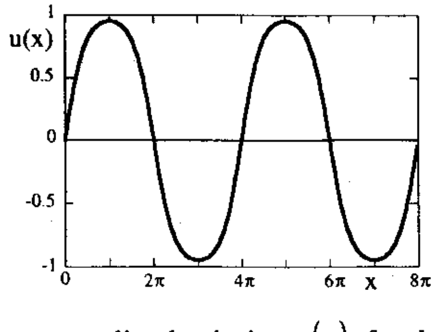

Fig.1 Stationary state of the normalized solution $u(x)$ for the case of $d = 8\pi$. Radial length x and vorticity u are normalized values.

order of magnitude estimate of D_3/D_{rr}. The ratio D_3/D_{rr} is characterized by the weighting function $D_3/D_{rr} \simeq R(K, \Omega)^2 k_\theta^2 (\partial^2/\partial k_r^2)$, which is evaluated as $\left| R(K, \Omega)^2 k_\theta^2 \, \partial^2/\partial k_r^2 \right| \sim \Delta\omega_k^{-2} k^2 \ell^2$. This consideration gives an estimate $D_3 \sim k_\theta^2 \ell^2 \Delta\omega_k^{-2} D_{rr}$. The absolute value of the ratio D_3/D_{rr} depends on the anisotropy of the spectral function N_k, but such a quantitative argument is not taken into account here. Substituting this ratio D_3/D_{rr} into Eq.(12a), one has an estimate $U_0 = v_z \ell^{-1} (1 - \mu)^{1/2}$, where $v_z = \Delta\omega_k k_\theta^{-1}$. The flow velocity $V_0 = U_0 \ell$ is given as

$$V_0 = v_z (1 - \mu)^{1/2} . \tag{13}$$

The fluctuation level of the drift-wave turbulence is determined in the fast time scale, and is expressed by use of the inhomogeneity of the zonal flow as $I = I_0 \left(1 + \tau_{ac}^2 (\mathrm{d}V/\mathrm{d}r)^2 \right)^{-1}$ where I is the magnitude of the drift wave fluctuations (i.e., Fourier space integral of N_k), I_0 is that in the absence of the zonal flow, and τ_{ac} is the autocorrelation time of the fluctuations of drift frequency range [1]. If U_0 approaches τ_{ac}, suppression of turbulence becomes effective. From Eqs.(12b) and (13), one has

$$I = \frac{I_0}{1 + \tau_{ac}^2 v_z^2 K_0^2 (1 - \mu)^2} . \tag{14}$$

From these results we have the following finding. First, the nonlinear mechanism for saturation of zonal flow is effective, and the flow velocity is sustained in the level of v_z. The saturation level of the zonal flow does not explicitly depend on the magnitude of the instability drive of the drift-wave turbulence which is represented by I_0. The drift-wave fluctuation level I increases with I_0. The amplitude of zonal flow also becomes larger by the increment of I_0, through the enhancement of the nonlinear broadening $\Delta\omega_k$ under stronger fluctuations. The radial shape of $V_Z(r)$ deviates from a sinusoidal function and is given by an elliptic integral function. The turbulent viscosity of the parallel flow has an important role for the evolution of the zonal flow. The dependencies of the velocity and scale length of zonal flow and of the fluctuation level on $\mu_\parallel (1 + 2q^2)$ are given by Eqs.(12b), (13), and (14). The flow is annihilated if $\mu_\parallel (1 + 2q^2) > D_{rr}$ holds. These are in contrast to the case of the collisional damping. In the case of collisional damping, the fluctuation level I is independent (only weakly dependent) of I_0 when the zonal flow is generated [3].

It is worth comparing the solution in this article with those in Ref.[8]. A kink-soliton like solution was found in [8] when the solution is searched for in an infinite domain. Equation (11) includes a particular solution of a kink-like soliton in an infinite domain with $\kappa = 0$. The solution in this article is stable against a perturbation. The one in ref.[8] can be unstable against a long-wave length perturbations.

This result can be tested by comparing with the result of direct numerical simulation. Analytic formula gives relation $U_0 \lambda^2 \cong 16\pi^2 v_z K_0^{-1}$ and the wave length $\lambda = 4\pi K_0^{-1} \left(1 - \mu\right)^{-1/2}$. With the estimate of strong turbulence $v_z = \Delta\omega_k k_\theta^{-1} \cong v_{di}$ and $K_0 \simeq s\rho_i^{-1}$, we have $\lambda \sim 4\pi s^{-1} \left(1 - \mu\right)^{-1/2} \rho_i$ and $U_0 \cong s\left(1 - \mu\right) v_{di}\, \rho_i^{-1}$. (v_{di} : the diamagnetic drift velocity and s : magnetic shear). This result is tested by the result of a three-dimensional nonlinear simulation of the ion-temperature-gradient (ITG) mode turbulence based on two fluid models [11]. In this simulation, the dynamics of the electrostatic potential, ion temperature and ion parallel velocity are followed in toroidal geometry with adiabatic response for electrons. Radial width of simulation domain is $120\rho_i$ and a realistic ITG dynamics was obtained in this case by switching off the unrealistically high parallel fluid heat conduction.. Parameters are $\varepsilon_n \equiv 2L_n/R = 0.9$, $L_n/L_{Ti} = 3.1$, $q = 1.4$ ($q = 0.7$ for zonal flow component in order to reduce the damping of zonal flow), and $s = 0.8$. (L_n and L_{Ti} are gradient scale lengths of density and temperature, respectively.) Details are explained in [11]. Figure 2 illustrates the radial distribution of the vorticity associated with the zonal flow, $\mathrm{d}\langle v_y \rangle/\mathrm{d}r$, where $\langle \cdots \rangle$ denotes the average over the magnetic surface and $r-$ and $y-$ coordinates are taken in the radial and poloidal directions, respectively. The simulation result confirms this theoretical modelling. Firstly, the radial distribution of the vorticity shows the flattened quasi-periodic form. This is an extreme of relatively low dispersion and high linear drive of the analytic result. Second, the periodic length is about $\sim 20\, \rho_i$ and is in the rage of the theoretical prediction. Third, the magnitude of the vorticity is $\mathrm{d}\langle v_y \rangle/\mathrm{d}r \sim 0.6 v_{di}\rho_i^{-1}$. This value is also in the rage of theoretical prediction, $U_0 \sim 0.8\left(1 - \mu\right) v_{di}\rho_i^{-1}$. Note that μ is kept smaller than unity in the simulation by lowering q -value for the zonal flow component.

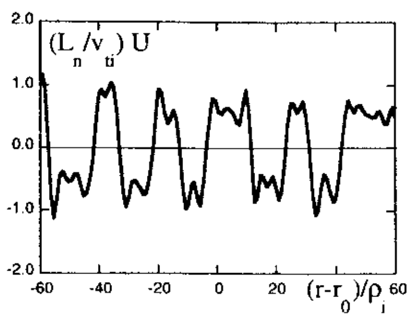

Fig.2 Radial distribution of vorticity of zonal flow U in the DNS. Snap shot in the stationary state shown. Origin of radius r_0 is chosen at the center of simulation box.

In summary, we studied the nonlinear structure of the zonal flow in toroidal plasma which is in the turbulent state. It is found that the drive of zonal flow starts to decrease at high level of velocity shear of zonal flow, but the viscosity damping of parallel flow does not. By this nonlinear mechanism, the zonal flows evolve into a stable stationary structure. The level of flow velocity and scale length are obtained. This result of the coherent structure of the zonal flow in turbulent plasma is confirmed by DNS. Thus we have clarified the nonlinear structure of zonal flow driven by turbulence in the range of drift waves.

The structure of the zonal flow which is obtained here critically depends on the two facts: (i) saturation (reduction) of zonal flow drive at the high level of zonal flow velocity (ii) the turbulence viscosity of parallel flow does not show this saturation at high zonal flow velocity. This theoretical modelling needs verification by DNS. For this purpose, DNS has been performed and the Reynolds stress has been measured. We have made a detailed measurement of the Reynolds stress in DNS, and obtained direct confirmations of (i) and (ii) including the dependencies of Eq.(9). Details will be reported in forthcoming article [14].

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