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Dynamics of Zonal Flows in Helical Systems

H. Sugama and T.-H. Watanabe

National Institute for Fusion Science, Graduate University for Advanced Studies, Toki 509-5292, Japan
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A theory for describing collisionless long-time behavior of zonal flows in helical systems is presented and its validity is verified by gyrokinetic-Vlasov simulation. It is shown that, under the influence of particles trapped in helical ripples, the response of zonal flows to a given source becomes weaker for lower radial wave numbers and deeper helical ripples while a high-level zonal-flow response, which is not affected by helical-ripple-trapped particles, can be maintained for a longer time by reducing their bounce-averaged radial drift velocity. This implies a possibility that helical configurations optimized for reducing neoclassical ripple transport can simultaneously enhance zonal flows which lower anomalous transport.

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Zonal flows are observed in numerous natural systems such as atmospheric currents, while in fusion science they are intensively investigated as an attractive mechanism for realizing a good plasma confinement [1]. Rosenbluth and Hinton [2] showed that initial $\mathbf{E} \times \mathbf{B}$ rotation in tokamaks is not fully damped by collisionless processes, but it approaches a finite value. Collisional decay of zonal flows occurs in the long course of time [3] although the residual zonal flows in a collisionless time scale still influence the turbulent transport. Since zonal flows are a key issue for improved confinement in helical systems as well [4,5], it is necessary to examine how helical geometries affect zonal-flow damping. In the present work, collisionless zonal-flow dynamics in helical systems is investigated. In the same manner as in Rosenbluth and Hinton [2], we here treat the ion-temperature-gradient (ITG) turbulence [6] as a known source and analytically derive the response kernel which relates the zonal-flow potential to the source and also represents dependence on an initially given zonal flow. We also verify the validity of the derived response kernel by a recently developed gyrokinetic-Vlasov-simulation code [7].

In helical configurations, the radial drift motion of particles trapped in helical ripples yields neoclassical ripple transport in the weak collisionality regime [8,9]. We show that this radial drift also causes a significant difference between long-time zonal-flow behavior in helical systems and that in tokamaks. It is observed in the large helical device (LHD) [10] that not only neoclassical but also anomalous transport is reduced by the inward shift of the magnetic axis which decreases the radial drift of helical-ripple-trapped particles but increases the unfavorable magnetic curvature to destabilize pressure-gradient-driven instabilities such as the ITG mode [11–13]. Our study suggests that helical configurations optimized for reduction of the neoclassical ripple transport may simultaneously lower the anomalous transport through enhancing the zonal-flow level.

We use the toroidal coordinates (r, θ, ζ) , where r , θ , and ζ denote the flux-surface label, the poloidal angle, and the

toroidal angle, respectively. The magnetic field is written as $\mathbf{B} = \nabla\psi(r) \times \nabla(\theta - \zeta/q(r))$, where $2\pi\psi(r)$ is equal to the toroidal flux within the flux surface labeled r and $q(r)$ represents the safety factor. Following Shaing and Hokin [9], we here consider helical systems with the magnetic field strength written by a function of poloidal and toroidal angles (its r dependence is not shown here for simplicity) as $B = B_0[1 - \epsilon_{10} \cos\theta - \epsilon_{L0} \cos(L\theta) - \sum_{n=0, \pm 1, \dots} \epsilon_h^{(n)} \times \cos\{(L+n)\theta - M\zeta\}] = B_0[1 - \epsilon_T(\theta) - \epsilon_H(\theta) \cos\{L\theta - M\zeta + \chi_H(\theta)\}]$, where $\epsilon_T(\theta) = \epsilon_{10} \cos\theta + \epsilon_{L0} \cos(L\theta)$, $\epsilon_H(\theta) = \sqrt{C^2(\theta) + D^2(\theta)}$, $\chi_H(\theta) = \arctan[D(\theta)/C(\theta)]$, $C(\theta) = \sum_{n=0, \pm 1, \dots} \epsilon_h^{(n)} \cos(n\theta)$, $D(\theta) = \sum_{n=\pm 1, \dots} \epsilon_h^{(n)} \times \sin(n\theta)$, and M (L) is the toroidal (main poloidal) period number of the helical field. For the LHD, $L = 2$ and $M = 10$. Here, we assume that $l/(qM) \ll 1$. Multiple-helicity effects can be included in the function $\epsilon_H(\theta)$.

The gyrokinetic equation [14] for the zonal-flow component with the perpendicular wave number vector $\mathbf{k}_\perp = k_r \nabla r$ is given by

$$\left(\frac{\partial}{\partial t} + v_\parallel \mathbf{b} \cdot \nabla + i\omega_D \right) g_{\mathbf{k}_\perp} = \frac{e}{T} F_0 J_0(k_\perp \rho) \frac{\partial \phi_{\mathbf{k}_\perp}}{\partial t} + S_{\mathbf{k}_\perp} F_0, \quad (1)$$

where $J_0(k_\perp \rho)$ is the zeroth-order Bessel function, $\rho = v_\perp/\Omega$ is the gyroradius, and $\Omega = eB/(mc)$ is the gyrofrequency. Here, subscripts to represent particle species are dropped for simplicity. The equilibrium distribution function F_0 is assumed to be given by the local Maxwellian and the perturbed particle distribution function $\delta f \equiv f - F_0$ is written in terms of the electrostatic potential ϕ and the solution g of Eq. (1) as $\delta f = -(e\phi/T)F_0 + g \exp(-i\mathbf{k}_\perp \cdot \boldsymbol{\rho})$, where $\boldsymbol{\rho} = \mathbf{b} \times \mathbf{v}/\Omega$. The drift frequency ω_D is defined by $\omega_D \equiv \mathbf{k}_\perp \cdot \mathbf{v}_d \equiv k_r v_{dr}$, where $v_{dr} = \mathbf{v}_d \cdot \nabla r$ is the radial component of the guiding-center drift velocity. The source term $S_{\mathbf{k}_\perp} F_0$ on the right-hand side of Eq. (1) represents the $\mathbf{E} \times \mathbf{B}$ nonlinearity and is written as $S_{\mathbf{k}_\perp} F_0 = (c/B) \sum_{\mathbf{k}'_\perp + \mathbf{k}''_\perp = \mathbf{k}_\perp} [\mathbf{b} \cdot (\mathbf{k}'_\perp \times \mathbf{k}''_\perp)] J_0(k'_\perp \rho) \phi_{\mathbf{k}'_\perp} g_{\mathbf{k}''_\perp}$.

The trapping parameter κ is defined by $\kappa^2 = [1 - \lambda B_0 \{1 - \epsilon_T(\theta) - \epsilon_H(\theta)\}] / [2\lambda B_0 \epsilon_H(\theta)]$ with $\lambda \equiv \mu/w$, where $w \equiv \frac{1}{2}mv^2$ and $\mu \equiv mv_\perp^2/(2B)$ represent the kinetic energy and the magnetic moment, respectively. Then, particles trapped in helical ripples are characterized by $\kappa^2 < 1$. Using $l/(qM) \ll 1$, we approximate the field line element dl by $R_0 d\zeta$, where R_0 denotes the major radius of the toroid. Then, the orbital average within a helical ripple is defined by $\bar{A} = \int (R_0 d\zeta / |v_\parallel|) A / \tau_h$, where $\tau_h = \int (R_0 d\zeta / |v_\parallel|)$; for $\kappa^2 < 1$, the integral $\int d\zeta$ goes over a closed orbit while, for $\kappa^2 > 1$, it goes a

whole helical ripple. Using the longitudinal adiabatic invariant J [9] given by $J = 16(R_0/M)(\mu B_0 \epsilon_H/m)^{1/2} \times [E(\kappa) - (1 - \kappa^2)K(\kappa)]$ for $\kappa^2 < 1$ and $J = 8(R_0/M) \times (\mu B_0 \epsilon_H/m)^{1/2} \kappa E(\kappa^{-1})$ for $\kappa^2 > 1$ with the complete elliptic integrals $K(\kappa)$ and $E(\kappa)$, the orbital average of the radial drift velocity within a helical ripple is given by $\bar{v}_{dr} = (mc/e\psi'\tau_h)(\partial J/\partial\theta)$, where $\psi' = d\psi/dr$ and $\tau_h = m(\partial J/\partial w)$. The drift frequency ω_D is expressed as $\omega_D \equiv k_r(\bar{v}_{dr} + v_\parallel \mathbf{b} \cdot \nabla \delta_r)$, where $\delta_r = \int^l (dl/v_\parallel)(v_{dr} - \bar{v}_{dr})$ represents the radial displacement of the guiding center from the helical-ripple-averaged radial position. Then, Eq. (1) is rewritten as

$$\left(\frac{\partial}{\partial t} + v_\parallel \mathbf{b} \cdot \nabla + ik_r \bar{v}_{dr}\right)(g_{\mathbf{k}_\perp} e^{ik_r \delta_r}) = \frac{e}{T} F_0 e^{ik_r \delta_r} J_0 \frac{\partial \phi_{\mathbf{k}_\perp}}{\partial t} + e^{ik_r \delta_r} S_{\mathbf{k}_\perp} F_0. \quad (2)$$

We here consider the long-time behavior of zonal flows. Then, in Eq. (2), the time-derivative terms, the radial guiding-center drift term, and the source term are smaller than the parallel streaming term such that they are regarded as of the higher order. The parallel derivative is rewritten as $\mathbf{b} \cdot \nabla \simeq R_0^{-1}(\partial/\partial\zeta + q^{-1}\partial/\partial\theta)$. Here, we treat the poloidal field as a higher-order quantity than the toroidal field. Based on these orderings, we expand $g_{\mathbf{k}_\perp} e^{ik_r \delta_r}$ as $g_{\mathbf{k}_\perp} e^{ik_r \delta_r} = h_0 + h_1 + \dots$ and obtain the lowest-order equation $(v_\parallel/R_0)(\partial h_0/\partial\zeta) = 0$ from Eq. (2). Thus, we can write $h_0 = h_0(t, r, \theta, w, \mu, \sigma)$, where the dependence on $\sigma = v_\parallel/|v_\parallel|$ disappears for $\kappa^2 < 1$. The first-order equation is written as

$$\frac{v_\parallel}{R_0} \frac{\partial h_1}{\partial \zeta} = -\left(\frac{\partial}{\partial t} + \frac{v_\parallel}{R_0 q} \frac{\partial}{\partial \theta} + ik_r \bar{v}_{dr}\right)h_0 + \frac{e}{T} F_0 e^{ik_r \delta_r} J_0 \frac{\partial \phi_{\mathbf{k}_\perp}}{\partial t} + e^{ik_r \delta_r} S_{\mathbf{k}_\perp} F_0. \quad (3)$$

For particles trapped in a helical ripple ($\kappa^2 < 1$), the orbital average of Eq. (3) and its time integral yield

$$h_0(t) = h_0(0) e^{-ik_r \bar{v}_{dr} t} + \int_0^t dt' e^{-ik_r \bar{v}_{dr}(t-t')} F_0 \left[\frac{e}{T} \left(e^{ik_r \delta_r} J_0 \frac{\partial \phi_{\mathbf{k}_\perp}(t')}{\partial t'} \right) + \overline{(e^{ik_r \delta_r} S_{\mathbf{k}_\perp}(t'))} \right]. \quad (4)$$

When $\kappa^2 > 1$, using the periodic condition $h_1(\zeta + 2\pi/M) = h_1(\zeta)$ and taking the orbital average of Eq. (3) within a helical ripple give

$$\left(\frac{\partial}{\partial t} + \omega_\theta \frac{\partial}{\partial \theta}\right)(e^{ik_r \Delta_r} h_0) = e^{ik_r \Delta_r} F_0 \left[\frac{e}{T} \left(e^{ik_r \delta_r} J_0 \frac{\partial \phi_{\mathbf{k}_\perp}}{\partial t} \right) + \overline{(e^{ik_r \delta_r} S_{\mathbf{k}_\perp})} \right], \quad (5)$$

where $\omega_\theta = 2\pi\sigma/(qM\tau_h)$ is the helical-ripple-averaged poloidal angular velocity and $\Delta_r = \sigma(qM/2\pi)(mc/e\psi')(J - J_t)$ with J_t defined later represents the radial displacement of the helical-ripple-averaged guiding-center position. For $\kappa^2 > 1$, particles are classified into two types, particles trapped by the toroidicity and passing particles. For these particles, we regard $\omega_\theta \partial(e^{ik_r \Delta_r} h_0)/\partial\theta$ as a dominant term in Eq. (5) based on the long-time ordering and expand $e^{ik_r \Delta_r} h_0$ as $e^{ik_r \Delta_r} h_0 = \eta_0 + \eta_1 + \dots$, where η_0 is independent of θ because it satisfies the lowest-order equation $\omega_\theta \partial \eta_0/\partial\theta = 0$. The solubility condition for η_1 is derived from Eq. (5) and integrated in time to give

$$\eta_0(t) = \eta_0(0) - \frac{e}{T} F_0 \langle e^{ik_r \Delta_r} \overline{(e^{ik_r \delta_r} J_0 \phi_{\mathbf{k}_\perp}(0))} \rangle_{\text{po}} + F_0 \left\langle e^{ik_r \Delta_r} \left[\overline{e^{ik_r \delta_r} \left\{ J_0 \frac{e \phi_{\mathbf{k}_\perp}(t)}{T} + R_{\mathbf{k}_\perp}(t) \right\}} \right] \right\rangle_{\text{po}}, \quad (6)$$

where $R_{\mathbf{k}_\perp}(t) \equiv \int_0^t dt' S_{\mathbf{k}_\perp}(t')$ and the poloidal-orbit average $\langle A \rangle_{\text{po}}$ is defined by $\langle A \rangle_{\text{po}} = \frac{1}{2} \sum_{\sigma=\pm 1} \times \int_{-\theta_t}^{\theta_t} (d\theta/|\omega_\theta|) A / \int_{-\theta_t}^{\theta_t} (d\theta/|\omega_\theta|)$ for toroidally trapped particles and $\langle A \rangle_{\text{po}} = \int_0^{2\pi} (d\theta/|\omega_\theta|) A / \int_0^{2\pi} (d\theta/|\omega_\theta|)$ for passing particles with θ_t given by the condition $\kappa(\theta = \theta_t) = 1$ which is equivalent to $\omega_\theta(\theta = \theta_t) = 0$. Now, J_t is defined by $J_t = J(\theta = \theta_t)$ for toroidally trapped particles and by $J_t = J(\theta = \pi)$ for passing particles.

The electrostatic potential $\phi_{\mathbf{k}_\perp}$ is determined by the quasineutrality condition, $-n_0 e \phi_{\mathbf{k}_\perp} / T_i + \int d^3 v J_0 g_{\mathbf{k}_\perp} = n_0 e \phi_{\mathbf{k}_\perp} / T_e + \int d^3 v g_{\mathbf{k}_\perp}$, where the small electron gyro-radius limit $k_\perp \rho_e \rightarrow 0$ is considered. In the lowest or-

der of the long-time ordering, we substitute Eq. (4) into $g_{\mathbf{k}_\perp} = e^{-ik_r \delta_r} h_0$ for $\kappa^2 < 1$ and Eq. (6) into $g_{\mathbf{k}_\perp} = e^{-ik_r \delta_r} e^{-ik_r \Delta_r} \eta_0$ for $\kappa^2 > 1$ in order to evaluate the non-adiabatic parts of the density perturbations. We find from Eq. (4) that effects of $\bar{v}_{dr}$ on the density of helical-ripple-trapped particles strongly depend on time t . Let us define a characteristic transition time τ_c by $\tau_c \sim 1/|k_r \bar{v}_{dr}|$, where $\bar{v}_{dr}$ is evaluated by considering helical-ripple-trapped thermal particles with $\mu B_0 \sim T$, $\kappa \sim 1$, and $\theta \sim \pi/2$.

When $t \ll \tau_c$, effects of $\bar{v}_{dr}$ are weak and the quasineutrality condition is written as

$$\begin{aligned}
& n_0 e \left(\frac{1}{T_i} + \frac{1}{T_e} \right) \phi_{\mathbf{k}_\perp}(t) \\
& - \frac{e}{T_i} \int_{\kappa^2 < 1} d^3 \mathbf{v} F_{i0} J_0 e^{-ik_r \delta_r} \overline{(e^{ik_r \delta_r} J_0 \phi_{\mathbf{k}_\perp}(t))} \\
& - \frac{e}{T_i} \int_{\kappa^2 > 1} d^3 \mathbf{v} F_{i0} e^{-ik_r \Delta_r} e^{-ik_r \delta_r} J_0 \langle e^{ik_r \Delta_r} \overline{(e^{ik_r \delta_r} J_0 \phi_{\mathbf{k}_\perp}(t))} \rangle_{\text{po}} \\
& - \frac{e}{T_e} \int_{\kappa^2 < 1} d^3 \mathbf{v} F_{e0} \overline{\phi_{\mathbf{k}_\perp}(t)} \\
& - \frac{e}{T_e} \int_{\kappa^2 > 1} d^3 \mathbf{v} F_{e0} \langle \phi_{\mathbf{k}_\perp}(t) \rangle_{\text{po}} = \sigma_<, \quad (7)
\end{aligned}$$

where $\sigma_<$ is given by the initial values $\phi_{\mathbf{k}_\perp}(0)$, $h_0(0)$, and $\eta_0(0)$ as well as the time integral of the $\mathbf{E} \times \mathbf{B}$ nonlinear source terms $R_{\mathbf{k}_\perp}(t) = \int_0^t dt' S_{\mathbf{k}_\perp}(t')$. Here, the radial displacement of the electron guiding center is neglected because of the small electron mass. Representing Eq. (7) by $\mathcal{L}\phi_{\mathbf{k}_\perp} = \sigma_<$ and defining the Hermitian inner product by $(u, v) \equiv \langle u^* v \rangle$, where $\langle \cdot \rangle$ denotes the flux-surface average, we find that the operator $\mathcal{L}$ is self-adjoint, $(u, \mathcal{L}v) = (\mathcal{L}u, v)$, and that $(u, \mathcal{L}u) \geq 0$. Then, the variational principle for $\mathcal{L}\phi_{\mathbf{k}_\perp} = \sigma_<$ is given by $\delta V = 0$, where $V \equiv (\phi_{\mathbf{k}_\perp}, \mathcal{L}\phi_{\mathbf{k}_\perp}) / |(\phi_{\mathbf{k}_\perp}, \sigma_< \phi_{\mathbf{k}_\perp})|^2$.

Now, we assume $k_\perp \rho$ and $k_r \Delta_r$ to be small and use them as expansion parameters. We neglect $k_r \delta_r$ because generally δ_r is much smaller than ρ . The source $\sigma_<$ is considered to be of order $k_\perp^2 \rho^2$. Then, from the lowest-order equation $(\phi_{\mathbf{k}_\perp}, \mathcal{L}_0 \phi_{\mathbf{k}_\perp}) = 0$, we can show that $\phi_{\mathbf{k}_\perp}$ is a

flux-surface function, $\partial \phi_{\mathbf{k}_\perp} / \partial \zeta = \partial \phi_{\mathbf{k}_\perp} / \partial \theta = 0$. The next-order equation $(\phi_{\mathbf{k}_\perp}, \mathcal{L}_1 \phi_{\mathbf{k}_\perp}) = (\phi_{\mathbf{k}_\perp}, \sigma_<)$ gives $e \phi_{\mathbf{k}_\perp} / T_i = \langle \sigma_< \rangle / \mathcal{D}_<$, where $\mathcal{D}_< = \langle \int d^3 \mathbf{v} F_{i0} [\frac{1}{2} k_\perp^2 \rho^2 + k_r^2 \langle \Delta_r^2 \rangle_{\text{po}} - \langle \Delta_r \rangle_{\text{po}}^2 H(\kappa^2 - 1)] \rangle$ and $H(x) = 1$ (for $x > 0$), 0 (for $x < 0$). Here, the second group of terms in the integrand represent the neoclassical polarization effect due to toroidally trapped particles with $\kappa^2 > 1$.

To the lowest order in $k_\perp^2 \rho^2$, electron contributions to $\sigma_<$ are neglected. The initial values $h_{i0}(0)$ and $\eta_{i0}(0)$ in Eq. (7) are given by $h_{i0}(0) = e^{ik_r \delta_r} g_{i\mathbf{k}_\perp}(0)$ and $\eta_{i0}(0) = \langle e^{ik_r \Delta_r} \overline{(e^{ik_r \delta_r} g_{i\mathbf{k}_\perp}(0))} \rangle_{\text{po}}$. We assume the initial perturbed ion gyrocenter distribution function to take the Maxwellian form, $\delta f_{i\mathbf{k}_\perp}^{(\text{gyro})}(0) \equiv -J_0(e \phi_{\mathbf{k}_\perp}(0) / T_i) F_{i0} + g_{i\mathbf{k}_\perp}(0) = (\delta n_{i\mathbf{k}_\perp}^{(\text{gyro})}(0) / n_0) F_{i0}$. The quasineutrality condition gives $\delta n_{i\mathbf{k}_\perp}^{(\text{gyro})}(0) = n_0(k_\perp^2 a_i^2)(e \phi_{\mathbf{k}_\perp}(0) / T_i)$ with $a_i^2 = T_i / (m_i \Omega_i^2)$. Then, we obtain $\langle \sigma_< \rangle = n_0 \langle k_\perp^2 a_i^2 \rangle e \phi_{\mathbf{k}_\perp}(0) / T_i + \langle \int d^3 \mathbf{v} F_{i0} R_{i\mathbf{k}_\perp}(t) \rangle$ and the long-time behavior of the zonal-flow potential for $t \ll \tau_c$,

$$\frac{e \phi_{\mathbf{k}_\perp}(t)}{T_i} = \mathcal{K}_< \left[\frac{e \phi_{\mathbf{k}_\perp}(0)}{T_i} + \frac{\int_0^t dt' \langle \int d^3 \mathbf{v} F_{i0} S_{i\mathbf{k}_\perp}(t') \rangle}{n_0 \langle k_\perp^2 a_i^2 \rangle} \right], \quad (8)$$

where the response kernel for $t \ll \tau_c$ is represented by

$$\mathcal{K}_< = 1 / (1 + G) \quad (9)$$

and

$$\begin{aligned}
G = & \frac{12}{\pi^3} B_0 R_0^2 q^2 \left\langle \frac{B^2}{|\nabla \psi|^2} \right\rangle \left[\int_0^{1/B_M} d\lambda \oint \frac{d\theta}{2\pi} (2\lambda B_0 \epsilon_H)^{-1/2} \kappa^{-1} K(\kappa^{-1}) \right] \left\{ (2\lambda B_0 \epsilon_H)^{1/2} \kappa E(\kappa^{-1}) \right. \\
& \left. - \frac{\oint \frac{d\theta}{2\pi} K(\kappa^{-1}) E(\kappa^{-1})}{\oint \frac{d\theta}{2\pi} (2\lambda B_0 \epsilon_H)^{-1/2} \kappa^{-1} K(\kappa^{-1})} \right\}^2 + \int_{1/B_M}^{1/B'_M} d\lambda \int_{\kappa^2(\theta) > 1} \frac{d\theta}{2\pi} (2\lambda B_0 \epsilon_H)^{1/2} \kappa K(\kappa^{-1}) \left\{ E(\kappa^{-1}) - \frac{1}{\kappa} \left(\frac{\epsilon_H(\theta)}{\epsilon_H} \right)^{1/2} \right\}^2 \right]. \quad (10)
\end{aligned}$$

The geometrical factor G measures the ratio of the neo-classical polarization due to toroidally trapped particles to the classical polarization. Here B_M denotes the maximum field strength over the flux surface and B'_M represents the minimum value of local maximum field strengths within each helical ripple.

Next, when $t \gg \tau_c$, the density of nonadiabatic helical-ripple-trapped particles is strongly damped because of phase mixing caused by the bounce-averaged radial drift motion [see Eq. (4)]. Then, the quasineutrality condition is given by Eq. (7) with the velocity-space integrals over the $\kappa^2 < 1$ region dropped. Employing the same procedures used in deriving Eqs. (8)–(10), $\phi_{\mathbf{k}_\perp}$ is shown to be again a flux-surface function to the lowest order in $k_\perp^2 \rho^2$ and $\epsilon_H^{1/2}$, and we obtain $e \phi_{\mathbf{k}_\perp} / T_i = \langle \sigma_> \rangle / \mathcal{D}_>$, where $\mathcal{D}_> \equiv \mathcal{D}_< + (2/\pi)(1 - \langle k_\perp^2 a_i^2 \rangle + T_i/T_e) \langle (2\epsilon_H)^{1/2} \rangle$ and $\langle \sigma_> \rangle \equiv \langle \sigma_< \rangle - (2/\pi) \langle (2\epsilon_H)^{1/2} \rangle \langle k_\perp^2 a_i^2 \rangle [n_0 e \phi_{\mathbf{k}_\perp}(0) / T_i]$. Finally, the long-time behavior of the zonal-flow potential for $t \gg \tau_c$ is given by

$$\begin{aligned}
\frac{e \phi_{\mathbf{k}_\perp}(t)}{T_i} = & \mathcal{K}_> \left[\frac{e \phi_{\mathbf{k}_\perp}(0)}{T_i} \right. \\
& \left. + \frac{\int_0^t dt' \langle \int_{\kappa^2 > 1} d^3 \mathbf{v} F_{i0} S_{i\mathbf{k}_\perp}(t') \rangle}{n_0 \langle k_\perp^2 a_i^2 \rangle \{1 - (2/\pi) \langle (2\epsilon_H)^{1/2} \rangle\}} \right] \quad (11)
\end{aligned}$$

and

$$\begin{aligned}
\mathcal{K}_> = & \langle k_\perp^2 a_i^2 \rangle [1 - (2/\pi) \langle (2\epsilon_H)^{1/2} \rangle] \\
& \times \{ \langle k_\perp^2 a_i^2 \rangle [1 - (2/\pi) \langle (2\epsilon_H)^{1/2} \rangle + G] \\
& + (2/\pi)(1 + T_i/T_e) \langle (2\epsilon_H)^{1/2} \rangle \}^{-1}. \quad (12)
\end{aligned}$$

Here, terms proportional to $\langle (2\epsilon_H)^{1/2} \rangle$ are derived from suppressing the density perturbations of the nonadiabatic helical-ripple-trapped particles. A term with T_i/T_e appears in the response kernel $\mathcal{K}_>$ for $t \gg \tau_c$ because not only ions but also electrons influence the quasineutrality condition through their helical-ripple-bounce-averaged radial drift motion. The dependence on electrons and on the radial wave number shown in Eq. (12) is not seen in the

FIG. 1. Time evolution of the zonal-flow potential obtained by the gyrokinetic-Vlasov simulation for a helical system with $L = 2$, $M = 10$, $q = 1.5$, $\epsilon_r = \epsilon_h = 0.1$, and $k_r a_i = 0.131$. A dashed horizontal line corresponds to $\mathcal{K}_{>}$ given by Eq. (12) for $t > \tau_c$.

tokamak case. In the axisymmetric limit $\epsilon_H \rightarrow +0$ with $\epsilon_T = \epsilon_i \cos\theta$, we obtain $G \rightarrow 1.6q^2/\epsilon_i^{1/2}$, which reduces both Eqs. (9) and (12) to the Rosenbluth-Hinton formula $\mathcal{K}_{R-H} = 1/(1 + 1.6q^2/\epsilon_i^{1/2})$ [2].

In order to examine the analytical results shown above, a linearized ion gyrokinetic equation combined with the quasineutrality condition is numerically solved by a toroidal flux-tube gyrokinetic-Vlasov code [7]. Since the perturbed electron density is simply calculated by using $n_{e\mathbf{k}_\perp} = (n_0 e / T_e)(\phi_{\mathbf{k}_\perp} - \langle \phi_{\mathbf{k}_\perp} \rangle)$ in our simulations, the term proportional to T_i/T_e in Eq. (12) should be dropped when comparing that formula with the simulation results. Here, we consider the $L = 2/M = 10$ single-helicity case, in which $\epsilon_h^{(n \neq 0)} = 0$ and therefore $\epsilon_H = \epsilon_h^{(0)} \equiv \epsilon_h$ is independent of θ . We also put $\epsilon_t \equiv \epsilon_{i0}$ and $\epsilon_{L0} = 0$ so that $\epsilon_T = \epsilon_i \cos\theta$. The initial perturbed ion gyrocenter distribution function is given by the Maxwellian form. We define the radial coordinate r by $\psi = B_0 r^2/2$ and use $\bar{v}_{dr} = -(c\mu/eB_0)\sin\theta$, $\langle k_\perp^2 a_i^2 \rangle \simeq k_r^2 a_i^2$, and $\tau_c \simeq (k_r c T_i / e B_0 R_0)^{-1} = (R_0 / v_{ti}) / (k_r a_i)$, where $a_i \simeq v_{ti} / \Omega_{i0}$, $\Omega_{i0} = e B_0 / (m_i c)$, and $v_{ti} = (T_i / m_i)^{1/2}$.

Time evolution of the zonal-flow potential obtained by the simulation is plotted by a solid curve in Fig. 1, where $\epsilon_t = 0.1$, $\epsilon_h = 0.1$, $q = 1.5$, and $k_r a_i = 0.131$ are used. Here, a dashed horizontal line represents the response kernel $\mathcal{K}_{>}$ given by Eq. (12) for $t > \tau_c (= 7.6 R_0 v_{ti})$. It is seen that, after oscillations of the geodesic acoustic mode (GAM) [15] are damped, the zonal-flow amplitude approaches the predicted value $\mathcal{K}_{>} = 0.038$, which is smaller than $\mathcal{K}_{<} = 0.39$ and $\mathcal{K}_{R-H} = 0.081$ for the used parameters. Under the conditions used in our simulation, the GAM oscillations dominate the zonal-flow evolution for $t < \tau_c$ so that we cannot identify $\mathcal{K}_{<}$ given by Eq. (8) which describes the long-time behavior for $t \ll \tau_c$ with rapid phenomena such as the GAM neglected. It is confirmed from other simulations for $k_r a_i = 0.0654, 0.131, 0.196$ and $\epsilon_h = 0.05, 0.1, 0.2$ that Eq. (12) agrees with the

long-time limit of $\langle \phi_{\mathbf{k}_\perp}(t) \rangle / \langle \phi_{\mathbf{k}_\perp}(0) \rangle$ obtained by the simulations within an error of about 15% at most. A better agreement between the simulation and theoretical results is verified for lower $k_r a_i$ and smaller ϵ_h because these parameters are assumed to be small in deriving the analytical results.

In conclusion, we have shown how the collisionless long-time behavior of zonal flows in helical systems is influenced by the bounce-averaged radial drift motion of helical-ripple-trapped particles. It is predicted that, under the influence of helical-ripple-trapped particles, for the lower radial wave numbers, the long-time limit of the zonal-flow potential amplitude becomes smaller although simultaneously the characteristic transition time $\tau_c (\sim 1/k_r |\bar{v}_{dr}|)$ becomes longer. In some optimized helical configurations such as quasipoloidally symmetric systems [16,17], which significantly reduce neoclassical transport by suppressing both $|\bar{v}_{dr}|$ and G , we expect the response kernels $\mathcal{K}_{>}$, $\mathcal{K}_{<}$, and τ_c to increase such that large zonal flows can be maintained for a long-time period, which contribute to a reduction of anomalous transport as well.

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- [1] P.H. Diamond *et al.*, in Proceedings of the 20th IAEA Fusion Energy Conference, Vilamoura, Portugal, 2004, OV/2-1 (to be published).
 - [2] M.N. Rosenbluth and F.L. Hinton, Phys. Rev. Lett. **80**, 724 (1998).
 - [3] F.L. Hinton and M.N. Rosenbluth, Plasma Phys. Controlled Fusion **41**, A653 (1999).
 - [4] A. Fujisawa *et al.*, Phys. Rev. Lett. **93**, 165002 (2004).
 - [5] A. Hasegawa and M. Wakatani, Phys. Rev. Lett. **59**, 1581 (1987); H. Sugama, M. Wakatani, and A. Hasegawa, Phys. Fluids **31**, 1601 (1988).
 - [6] W. Horton, Rev. Mod. Phys. **71**, 735 (1999).
 - [7] T.-H. Watanabe and H. Sugama, in 20th IAEA Fusion Energy Conference, Vilamoura, Portugal, 2004, TH/8-3Rb (Ref. [1]).
 - [8] M. Wakatani, *Stellarator and Heliotron Devices* (Oxford University Press, Oxford, 1998), Chap. 7.
 - [9] K.C. Shaing and S.A. Hokin, Phys. Fluids **26**, 2136 (1983).
 - [10] O. Motojima *et al.*, Nucl. Fusion **43**, 1674 (2003).
 - [11] H. Yamada *et al.*, Plasma Phys. Controlled Fusion **43**, A55 (2001).
 - [12] T. Kuroda and H. Sugama, J. Phys. Soc. Jpn. **70**, 2235 (2001).
 - [13] G. Rewoldt *et al.*, Nucl. Fusion **42**, 1047 (2002).
 - [14] E. A. Frieman and L. Chen, Phys. Fluids **25**, 502 (1982).
 - [15] N. Winsor, J.L. Johnson, and J.J. Dawson, Phys. Fluids **11**, 2448 (1968).
 - [16] D.A. Spong *et al.*, Nucl. Fusion **41**, 711 (2001).
 - [17] M. Yokoyama, J. Plasma Fusion Res. **78**, 205 (2002).