Dynamics of Zonal Flows in Helical Systems

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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 $B = B_0[1 - \epsilon_{10}\cos\theta - \epsilon_{L0}\cos(L\theta) - \sum_{n=0,\pm1,\ldots}\epsilon_h^{(n)}\times\cos\{(L+n)\theta - M\zeta\}] = B_0[1 - \epsilon_T(\theta) - \epsilon_H(\theta)\cos\{L\theta - M\zeta + \chi_H(\theta)\}]$, where $\epsilon_T(\theta) = \epsilon_{10}\cos\theta + \epsilon_{L0}\cos(L\theta)$, $\epsilon_H(\theta) = \sqrt{C^2(\theta) + D^2(\theta)}$, $\chi_H(\theta) = \arctan[D(\theta)/C(\theta)]$, $C(\theta) = \sum_{n=0,\pm1,\ldots}\epsilon_h^{(n)}\cos(n\theta)$, $D(\theta) = \sum_{n=\pm1,\ldots}\epsilon_h^{(n)}\times\sin(n\theta)$, 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

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

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 $\overline{A} = \int (R_0 d\zeta/|v_\parallel|)A/\tau_h$, 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 $\overline{v}_{dr} = (mc/e\psi'\tau_h)(\partial J/\partial \theta)$, where $\psi' = d\psi/dr$ and $\tau_h = m(\partial J/\partial w)$. The drift frequency ω_D is expressed as $\omega_D \equiv k_r(\overline{v}_{dr} + v_\parallel \mathbf{b}\cdot\nabla\delta_r)$, where $\delta_r = \int^l (dl/v_\parallel)(v_{dr} - \overline{v}_{dr})$ represents the radial displacement of the guiding center from the helical-ripple-averaged radial position. Then, Eq. (1) is rewritten as

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

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 $g_{\mathbf{k}_\perp} e^{ik_r\delta_r} = h_0 + h_1 + \cdots$ 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

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

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

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

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

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

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 $e^{ik_r\Delta_r}h_0 = \eta_0 + \eta_1 + \cdots$, 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

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

where $R_{\mathbf{k}_\perp}(t) \equiv \int_0^t dt' S_{\mathbf{k}_\perp}(t')$ and the poloidal-orbit average $\langle A\rangle_{\rm po}$ is defined by $\langle A\rangle_{\rm po} = \frac{1}{2}\sum_{\sigma=\pm1} \times \int_{-\theta_t}^{\theta_t}(d\theta/|\omega_\theta|)A/\int_{-\theta_t}^{\theta_t}(d\theta/|\omega_\theta|)$ for toroidally trapped particles and $\langle A\rangle_{\rm po} = \int_0^{2\pi}(d\theta/|\omega_\theta|)A/\int_0^{2\pi}(d\theta/|\omega_\theta|)$ 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, $-n_0 e\phi_{\mathbf{k}_\perp}/T_i + \int d^3v J_0 g_{i\mathbf{k}_\perp} = n_0 e\phi_{\mathbf{k}_\perp}/T_e + \int d^3v g_{e\mathbf{k}_\perp}$, where the small electron gyroradius limit $k_\perp\rho_e \to 0$ is considered. In the lowest order 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 nonadiabatic parts of the density perturbations. We find from Eq. (4) that effects of $\overline{v}_{dr}$ on the density of helical-ripple-trapped particles strongly depend on time t. Let us define a characteristic transition time τ_c by $\tau_c \sim 1/|k_r\overline{v}_{dr}|$, where $\overline{v}_{dr}$ 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 $\overline{v}_{dr}$ are weak and the quasineutrality condition is written as

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

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 $\mathcal{D}_<=\langle\int d^3vF_{i0}[\frac{1}{2}k_\perp^2\rho^2+k_r^2\{\langle\Delta_r^2\rangle_{\rm po}-\langle\Delta_r\rangle_{\rm po}^2\}H(\kappa^2-1)]\rangle$ 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 η_{i0} in Eq. (7) are given by $h_{i0}(0)=\overline{e^{ik_r\delta_r}g_{i\mathbf{k}_\perp}(0)}$ and $\eta_{i0}(0)=\langle e^{ik_r\Delta_r}\overline{(e^{ik_r\delta_r}g_{i\mathbf{k}_\perp}(0))}\rangle_{\rm po}$. We assume the initial perturbed ion gyrocenter distribution function to take the Maxwellian form, $\delta f_{i\mathbf{k}_\perp}^{(\rm gyro)}(0)\equiv -J_0(e\phi_{\mathbf{k}_\perp}(0)/T_i)F_{i0}+g_{i\mathbf{k}_\perp}(0)=(\delta n_{i\mathbf{k}_\perp}^{(\rm gyro)}(0)/n_0)F_{i0}$. The quasineutrality condition gives $\delta n_{i\mathbf{k}_\perp}^{(\rm gyro)}(0)=n_0(k_\perp^2a_i^2)(e\phi_{\mathbf{k}_\perp}(0)/T_i)$ with $a_i^2=T_i/(m_i\Omega_i^2)$. Then, we obtain $\langle\sigma_<\rangle=n_0\langle k_\perp^2a_i^2\rangle e\phi_{\mathbf{k}_\perp}(0)/T_i+\langle\int d^3vF_{i0}R_{i\mathbf{k}_\perp}(t)\rangle$ and the long-time behavior of the zonal-flow potential for $t\ll\tau_c$,

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

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

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

and

$$G=\frac{12}{\pi^3}B_0R_0^2q^2\left\langle\frac{B^2}{|\nabla\psi|^2}\right\rangle\left[\int_0^{1/B_M}d\lambda\oint\frac{d\theta}{2\pi}(2\lambda B_0\epsilon_H)^{-1/2}\kappa^{-1}K(\kappa^{-1})\left\{(2\lambda B_0\epsilon_H)^{1/2}\kappa E(\kappa^{-1})-\frac{\oint\frac{d\theta}{2\pi}K(\kappa^{-1})E(\kappa^{-1})}{\oint\frac{d\theta}{2\pi}(2\lambda B_0\epsilon_H)^{-1/2}\kappa^{-1}K(\kappa^{-1})}\right\}^2+\int_{1/B_M}^{1/B'_m}d\lambda\int_{\kappa^2(\theta)>1}\frac{d\theta}{2\pi}(2\lambda B_0\epsilon_H)^{1/2}\kappa K(\kappa^{-1})\left\{E(\kappa^{-1})-\frac{1}{\kappa}\left(\frac{\epsilon_H(\theta_t)}{\epsilon_H}\right)^{1/2}\right\}^2\right]. \quad (10)$$

The geometrical factor G measures the ratio of the neoclassical 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 $\mathcal{D}_>\equiv\mathcal{D}_<+(2/\pi)(1-\langle k_\perp^2a_i^2\rangle+T_i/T_e)\langle(2\epsilon_H)^{1/2}\rangle$ and $\langle\sigma_>\rangle\equiv\langle\sigma_<\rangle-(2/\pi)\langle(2\epsilon_H)^{1/2}\rangle\langle k_\perp^2a_i^2\rangle[n_0e\phi_{\mathbf{k}_\perp}(0)/T_i]$. Finally, the long-time behavior of the zonal-flow potential for $t\gg\tau_c$ is given by

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

and

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

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

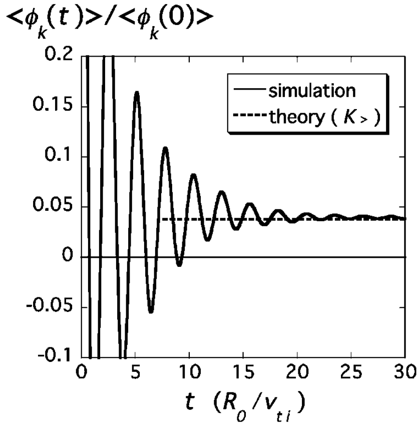

FIG. 1. Time evolution of the zonal-flow potential obtained by the gyrokinetic-Vlasov simulation for a helical system with $L = 2$, $M = 10$, $q = 1.5$, $\epsilon_t = \epsilon_h = 0.1$, and $k_r a_i = 0.131$. A dashed horizontal line corresponds to $\mathcal{K}_>$ given by Eq. (12) for $t > \tau_c$.

tokamak case. In the axisymmetric limit $\epsilon_H \rightarrow +0$ with $\epsilon_T = \epsilon_t \cos\theta$, we obtain $G \rightarrow 1.6q^2/\epsilon_t^{1/2}$, which reduces both Eqs. (9) and (12) to the Rosenbluth-Hinton formula $\mathcal{K}_{\text{R-H}} = 1/(1 + 1.6q^2/\epsilon_t^{1/2})$ [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 $\epsilon_t \equiv \epsilon_{10}$ and $\epsilon_{L0} = 0$ so that $\epsilon_T = \epsilon_t \cos\theta$. 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 $\overline{v}_{dr} = -(c\mu/eR_0)\sin\theta$, $\langle k_\perp^2 a_i^2\rangle \simeq k_r^2 a_i^2$, and $\tau_c \simeq (k_r c T_i/eB_0 R_0)^{-1} = (R_0/v_{ti})/(k_r a_i)$, where $a_i \simeq v_{ti}/\Omega_{i0}$, $\Omega_{i0} = eB_0/(m_i c)$, 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 $\mathcal{K}_{\text{R-H}} = 0.081$ 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 $\tau_c(\sim 1/k_r|\overline{v}_{dr}|)$ becomes longer. In some optimized helical configurations such as quasipoloidally symmetric systems [16,17],which significantly reduce neoclassical transport by suppressing both $|\overline{v}_{dr}|$ 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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