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The magnetic Rayleigh–Taylor instability and flute waves at the ion Larmor radius scales

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The theory of flute waves (with arbitrary spatial scales compared to the ion Larmor radius) driven by the Rayleigh–Taylor instability (RTI) is developed. Both the kinetic and hydrodynamic models are considered. In this way we have extended the previous analysis of RTI carried out in the long wavelength limit. It is found that complete finite ion Larmor radius stabilization is absent when the ion diamagnetic velocity attains the ion gravitation drift velocity. The hydrodynamic approach allowed us to deduce a new set of nonlinear equations for flute waves with arbitrary spatial scales. It is shown that the previously deduced equations are inadequate when the wavelength becomes of the order of the ion Larmor radius. In the linear limit a Fourier transform of these equations yields the dispersion relation which in the so-called Padé approximation corresponds to the results of the fully kinetic treatment. The development of such a theory gives us enough grounds for an adequate description of the RTI stabilization by the finite ion Larmor radius effect. © *2011 American Institute of Physics*. [doi:10.1063/1.3554661]

I. INTRODUCTION

The flute waves driven by the magnetic Rayleigh–Taylor instability (RTI) are believed to be plausibly responsible for the anomalous transport in magnetic confinement plasmas. This RTI is the counterpart of the classic RTI of an inhomogeneous fluid in a gravitational field. The gravity force may alternatively be replaced by a force due to the centrifugal acceleration of the particles (ions and electrons) moving along the curved magnetic field lines. In this case the instability is termed magnetic-curvature-driven RTI. It has been recently recognized that coherent large-scale structures such as zonal flows and streamers, arising in small scale flute microturbulence, can play an essential role in plasma transport.[1,2] The finite ion Larmor radius (FLR) effects substantially influence the evolution of the flute wave dynamics since they lead to the stabilization of the magnetic RTI.[3–5]

The previous theories of RTI and nonlinear structures of the flute waves[1–13] were restricted to the long wavelength limit where the wave spatial scale is much larger than the ion Larmor radius. In what follows, we shall use a local Cartesian coordinate system where the z-axis is along the external magnetic field $\mathbf{B}_0$, the x-axis is along the direction of the plasma inhomogeneity and the direction of the gravitational acceleration $\mathbf{g}$, and where the y-axis completes the triad. In low beta plasmas the dispersion relation and the growth rate of the flute (interchange) waves are[3–5]

$$\omega(\omega - \omega_{iD} - \omega_g) + \frac{k_y^2}{k_\perp^2}\kappa_N g = 0, \tag{1}$$

$$\frac{\gamma}{\Gamma} = \left[1 - \frac{z}{4\sigma}(1+\eta+\sigma)^2\right]^{1/2}, \tag{2}$$

where ω is the wave frequency, $k_\perp^2 = k_x^2 + k_y^2$, k_x and k_y are the components of the wave vector in the x and y directions, $\omega_{iD} = k_y v_{iD}$, $\omega_g = k_y v_g$, $\mathbf{v}_{iD} = -\rho_i^2 \omega_{ci}(\kappa_N + \kappa_T)\hat{\mathbf{y}}$ is the equilibrium ion diamagnetic drift velocity, $\kappa_N = -d \ln n_{0i}/dx > 0$, $\sigma = g/\kappa_N v_{Ti}^2 > 0$, $\kappa_T = -d \ln T_{0i}/dx$, n_{0i} and T_{0i} are the equilibrium ion number density and temperature, respectively, $\rho_i = (T_{0i}/m_i\omega_{ci}^2)^{1/2}$ is the ion Larmor radius, $\omega_{ci} = ZeB_0/m_i$ is the ion cyclotron frequency, m_i is the ion mass, Z is the ion charge number, e is the magnitude of the elementary charge, $\hat{\mathbf{y}}$ is the unit vector along the y axis, $\mathbf{v}_g = -(g/\omega_{ci})\hat{\mathbf{y}}$ is the ion gravitational drift velocity, $\eta = \kappa_T/\kappa_N$, $\Gamma = |k_y/k_\perp|(g\kappa_N)^{1/2}$ is the maximum value of the RTI growth rate, $v_T = (T_{0i}/m_i)^{1/2}$ is the ion thermal velocity, and $z = k_\perp^2 \rho_i^2$.

From Eq. (2) it follows that instability stabilization occurs for finite z in a plasma with finite $v_N/v_g = \kappa_N v_T^2/g$, where $v_N = -\rho_i^2 \omega_{ci}\kappa_N$ is the ion diamagnetic drift velocity due to the ion density inhomogeneity. Since Eqs. (1) and (2) were obtained in the long wavelength approximation, $z \ll 1$, the question of how the results will change qualitatively and quantitatively at $z \simeq 1$ still remains open.

The development of kinetic and hydrodynamic theories of flute waves with the spatial scales of the order of the ion Larmor radius propagating in a low-β plasmas is the main purpose of the present study. Furthermore, the effects of finite ratio v_N/v_g are also taken into account.

The manuscript is organized as follows. Section II describes the results of the kinetic approach in the linear approximation. A closed set of equations describing the nonlinear dynamics of the flute waves with arbitrary spatial scales is deduced in Sec. III. Our discussion and conclusions are found in Sec. IV.

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II. KINETIC DESCRIPTION

Using the fully kinetic description one can obtain the linear response of the normalized ion density perturbation to the variation of the electrostatic field as[5,14–16]

$$\delta n = -\frac{\omega_N}{\omega'}\Phi - \left[\alpha_{11}\left(1-\frac{\omega_N}{\omega'}\right) - \alpha_{12}\frac{\omega_T}{\omega'}\right]\Phi, \tag{3}$$

where $\delta n = \tilde{n}_i/n_{0i}$, $\tilde{n}_i = n_i - n_{0i}$ and n_{0i} are the perturbed and unperturbed ion number densities, $\omega_{N,T} = -k_y\rho_i^2\omega_{ci}\kappa_{N,T}$ is the ion drift frequency due to inhomogeneity of the ion density or temperature, $\omega' = \omega - k_y v_g$, $\Phi = Ze\varphi/T_{0i}$ is the normalized electrostatic potential, and $\mathbf{E}_\perp = -\nabla_\perp\varphi$. Furthermore, the subscript $\perp$ denotes the vector component perpendicular to the ambient magnetic field, $\alpha_{11} = 1-\Gamma_0(z)$, $\alpha_{12} = z\Gamma_0(z)[1-I_1(z)/I_0(z)]$, $\Gamma_0(z) \equiv \exp(-z)I_0(z)$, where $I_{0,1}$ is the modified Bessel functions of the first kind. In contrast to Refs. 5 and 14–16 we have taken into account that due to the ion gravitational drift the wave frequency is Doppler-shifted by the value $k_y v_g$.

The variation of the normalized electron number density is

$$\delta n_e = -\frac{\omega_N}{\omega}\Phi, \tag{4}$$

where $\delta n_e = \tilde{n}_e/n_{0e}$, $\tilde{n}_e = n_e - n_{0e}$ and n_{0e} are the perturbed and unperturbed electron number densities, respectively.

Invoking the quasineutrality, $\delta n_e = \delta n$, we obtain the linear dispersion relation for the flute waves,

$$\omega\left(\omega - \omega_N - \frac{\alpha_{12}}{\alpha_{11}}\omega_T - \omega_g\right) + \frac{z}{\alpha_{11}}\frac{k_y^2}{k_\perp^2}\kappa_N g = 0, \tag{5}$$

and the expression for the normalized growth rate,

$$\frac{\gamma}{\Gamma} = \left[\frac{z}{\alpha_{11}} - \frac{z}{4\sigma}\left(1+\frac{\alpha_{12}}{\alpha_{11}}\eta + \sigma\right)^2\right]^{1/2}. \tag{6}$$

In the long wavelength approximation, $z \ll 1$, i.e., when $\alpha_{12} \simeq \alpha_{11} \simeq z$, the dispersion relation (5) and the growth rate (6) coincide with those given by Eqs. (1) and (2). As it was mentioned before our consideration is valid only for low beta plasmas. The effects of finite beta were recently discussed in Refs. 12 and 13. It should be noted that derivation of Eqs. (1) and (2) requires the use of the ion gyroviscosity or the generalized Ohm's law in an extended magnetohydrodynamic (MHD) model.[4,7]

From Eq. (6) one easily finds that the cut-off wave number (at which $\gamma = 0$) is defined from the condition

$$\frac{4\sigma}{1-e^{-z}I_0(z)} = \left(1+\sigma+\frac{ze^{-z}I_0(z)(1-I_1(z)/I_2(z))}{1-e^{-z}I_0(z)}\eta\right)^2. \tag{7}$$

In a plasma with uniform temperature ($\eta = 0$) it reduces to a very simple form,

$$e^{-z}I_0(z) = \frac{(1-\sigma)^2}{(1+\sigma)^2}. \tag{8}$$

Equation (8) possesses the finite solution for $z = z_{cr}$ except for the case when $\sigma = 1$. The latter assumes that the critical wave number for complete FLR stabilization tends toward infinity and thus such stabilization in this case is absent. The incorporation of the finite values of η also does not lead to disappearance of singularity and plasma remains unstable when $\sigma = 1$. A similar conclusion about the absence of complete FLR stabilization of RTI has been made also in Ref. 13. However, such a statement has been based in this paper on the extended MHD equations which are valid only in the long wavelength limit.

III. HYDRODYNAMIC DESCRIPTION

On one hand, the description of waves with spatial scales of the order of the ion Larmor radius requires kinetic treatment. On the other hand, the hydrodynamic approximation can help us in clarifying the physics of the involved phenomena. In order to formulate the relevant hydrodynamic equations we decompose the ion velocity as

$$\mathbf{v}_i \simeq \mathbf{v}_E + \mathbf{v}_{iD} + \mathbf{v}_g + \mathbf{v}_E^P + \mathbf{v}_{iD}^P + \mathbf{v}_g^P, \tag{9}$$

where the low frequency limit is assumed. Furthermore, $\mathbf{v}_E$ is the $\mathbf{E}\times\mathbf{B}$ drift velocity, $\mathbf{v}_{iD} = (n_{0i}m_i\omega_{ci})^{-1}(\hat{\mathbf{z}}\times\nabla p)$ is the ion diamagnetic velocity, $\hat{\mathbf{z}}$ is the unit vector along the external magnetic field $\mathbf{B}_0$, p is the ion pressure, and $\mathbf{v}_E^P$, $\mathbf{v}_{iD}^P$, and $\mathbf{v}_g^P$ stand for the polarization parts of the ion velocity related to the drift velocities $\mathbf{v}_E$, $\mathbf{v}_{iD}$, and $\mathbf{v}_g$ through the relations

$$\mathbf{v}_E^P = \frac{1}{\omega_{ci}}\left(\hat{\mathbf{z}}\times\frac{d\mathbf{v}_E}{dt}\right), \quad \mathbf{v}_{iD}^P = \frac{1}{\omega_{ci}}\left(\hat{\mathbf{z}}\times\frac{d\mathbf{v}_{iD}}{dt}\right), \quad \text{and} \quad \mathbf{v}_g^P = \frac{1}{\omega_{ci}}\left(\hat{\mathbf{z}}\times\frac{d\mathbf{v}_g}{dt}\right). \tag{10}$$

Following Refs. 14, 15, and 17 we note that the convective time derivative d/dt in Eq. (10) is $d/dt \equiv \partial/\partial t + (\mathbf{v}_E + \mathbf{v}_g)\cdot\nabla_\perp$. The absence of the contribution from $\mathbf{v}_{iD}\cdot\nabla_\perp$ term in the round brackets reflects the "gyroviscous cancellation" effect according to which the correction to the ion diamagnetic velocity due to the magnetic viscosity cancels the contribution from advective diamagnetic part of the ion velocity $\mathbf{v}_{iD}\cdot\nabla_\perp$. After substitution of Eqs. (9) and (10) into the ion continuity equation in the dimensionless form, one obtains

$$\frac{d}{d\tau}(1-\nabla_\perp^2)\delta n - \frac{d}{d\tau}\nabla_\perp^2\delta T + \hat{\kappa}_N\frac{\partial\Phi}{\partial Y} - \frac{d}{d\tau}\nabla_\perp^2\Phi + \delta n\left(\hat{\kappa}_T\frac{\partial\delta n}{\partial Y} - \hat{\kappa}_N\frac{\partial\delta T}{\partial Y}\right) - \{\nabla_\perp\Phi, \nabla_\perp\delta p\} = -\hat{\mu}\nabla_\perp^4\Phi, \tag{11}$$

where the spatiotemporal variables are normalized to ρ_i and ω_{ci}^{-1}, i.e., $(X,Y) = (x,y)\rho_i^{-1}$, $\tau = t\omega_{ci}$, $\hat{\kappa}_N = \kappa_N\rho_i$, $\hat{g} = g/(\omega_{ci}v_T)$, $d/d\tau = \partial/\partial\tau + \{\Phi,\ldots\} - \hat{g}\,\partial/\partial Y$, $\delta T = (T-T_0)/T_0$ is the variation of the normalized ion temperature, $\hat{\mu} = (3/10)\nu/\omega_{ci}$ is the dimensionless dynamic collision viscosity, and ν is the ion-ion collision frequency. Furthermore, $\{f,g\} = (\partial f/\partial X)\partial g/\partial Y - (\partial f/\partial Y)\partial g/\partial X$ denotes the Poisson bracket. Here the collisional term is retained keeping in mind the possible use of these equations for numerical simulations in which the so-called numerical viscosity can be introduced.

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Equation (11) should be supplemented by the ion thermal balance condition

$$n_i\left(\frac{d}{dt}+\mathbf{v}_{iD}\cdot\nabla_\perp\right)T+n_iT\nabla\cdot\mathbf{v}_i+\nabla\cdot\mathbf{q}_\perp=0, \tag{12}$$

where the ion thermal flux $\mathbf{q}_\perp$ is given by

$$\mathbf{q}_\perp=\frac{2p_i}{m_i\omega_{ci}}(\hat{\mathbf{z}}\times\nabla_\perp T)-\frac{2}{m_i\omega_{ci}^2}\frac{d}{dt}(p\nabla_\perp T)-\frac{2\nu p}{m_i\omega_{ci}^2}\nabla_\perp T. \tag{13}$$

Equations (12) and (13) have been written in a form that takes into account the two-dimensional character of the ion motion in the wave field with the effective ratio of specific heats $\gamma=2$, see Refs. 14 and 15. The first term on the right of Eq. (13) represents the standard thermal flux in Braginskii hydrodynamics, the second one is an addition corresponding to the so-called polarization part of the ion thermal flux in the Grad hydrodynamics. Finally, the third term describes the collisional thermal flux arising due to the ion-ion collisions.

In dimensionless form the equation for the variation of the ion temperature is

$$\begin{aligned}&\frac{d}{d\tau}(1-3\nabla_\perp^2)\delta T-\frac{d}{d\tau}\nabla_\perp^2\delta n-\frac{d}{d\tau}\nabla_\perp^2\Phi+\hat{\kappa}_T\frac{\partial\Phi}{\partial Y}\\&\quad+\delta p\left(\hat{\kappa}_T\frac{\partial\delta n}{\partial Y}-\hat{\kappa}_N\frac{\partial\delta T}{\partial Y}\right)+\{\nabla_\perp\Phi,\nabla_\perp(\delta p+2\delta T)\}\\&\quad=\hat{\nu}\nabla_\perp^2\delta T,\end{aligned} \tag{14}$$

where $\hat{\kappa}_T=\kappa_T\rho_i$, $\hat{\nu}=2\nu_{ii}/\omega_{ci}$ is the normalized ion-ion frequency, and $\delta p=\tilde{p}/p_{i0}$ is the normalized ion pressure perturbation.

Equations (11) and (14) should be supplemented by the electron continuity equation. For that purpose we decompose the electron velocity as $\mathbf{v}_e=\mathbf{v}_E+\mathbf{v}_{eD}$. Here $\mathbf{v}_{eD}=Z\tau_i\rho_i^2\omega_{ci}\kappa_N\hat{\mathbf{y}}$ is the electron diamagnetic drift velocity, and $\tau_i=T_e/T_{i0}$ is the ratio of the electron to ion temperature. In the dimensionless form the electron continuity equation is

$$\frac{\partial\delta n}{\partial\tau}+\hat{\kappa}_N\frac{\partial\Phi}{\partial Y}=\{\delta n,\Phi\}+\hat{D}\nabla_\perp^2\delta n, \tag{15}$$

where $\hat{D}\equiv(m_e/m_i)Z\tau_i\hat{\nu}_{ee}$ is the dimensionless diffusion coefficient, $\hat{\nu}_{ee}=\nu_{ee}/\omega_{ci}$ is the normalized electron-electron collision frequency, and m_e is the electron mass.

Equations (11), (14), and (15) constitute a closed set of coupled equations for δn, Φ, and δT describing nonlinear flute waves with arbitrary spatial scales. These equations can be used for the study of nonlinear dynamics of magnetic RTI. In the general case they are quite cumbersome and can be solved only numerically. We note that a similar set but containing only two equations for δn and Φ (the ion and electron continuity equations) has been analyzed in Refs. 4 and 6–11. They can be obtained from our equations by setting $\delta T\to 0$ and rearranging Eq. (11) with the help of Eq. (15). This results in the following equation:

$$\begin{aligned}&\frac{\partial\nabla_\perp^2\Phi}{\partial\tau}-(\hat{\kappa}_N+\hat{g})\frac{\partial\nabla_\perp^2\Phi}{\partial Y}-\hat{g}\frac{\partial(1-\nabla_\perp^2)\delta n}{\partial Y}\\&\quad+\nabla_\perp\{\delta n,\nabla_\perp\Phi\}+\{\Phi,\nabla_\perp^2\Phi\}-\hat{\kappa}_T\delta n\frac{\partial\delta n}{\partial Y}\\&\quad=\hat{\mu}\nabla_\perp^4\Phi+\hat{D}\nabla_\perp^2(1-\nabla_\perp^2)\delta n.\end{aligned} \tag{16}$$

Furthermore, if one neglects $\hat{g}$ relative to $\hat{\kappa}_N$ (this corresponds to $\sigma\ll 1$) and scalar nonlinearity (the last term on the left) in the large-scale approximation ($\nabla_\perp^2\ll 1$), Eq. (16) reduces to the corresponding equation of Refs. 4 and 6–11. This assumes that the equation for the thermal balance equation (14) remains outside the current consideration. At this point a few comments are in order. In Ref. 14 it has noted that such procedure is inadequate since Eqs. (14) and (16) become inconsistent and the ion temperature variations cannot be neglected. Thus, the rigorous consideration of the problem at hand requires consideration of the full system containing three equations for three variables. Summarizing, we note that the detailed picture of the nonlinear dynamics of the magnetic RTI demands quite complicated numerical modeling based on our newly deduced equations (11), (14), and (15).

In the case of collisionless plasmas and in the linear approximation using a Fourier transform of these equations, one can obtain the linear response of the ion number density as

$$\delta n=-\frac{\omega_N}{\omega'}\Phi-\left[\alpha_{11}^H\left(1-\frac{\omega_N}{\omega'}\right)-\alpha_{12}^H\frac{\omega_T}{\omega'}\right]\Phi, \tag{17}$$

where $\alpha_{11}^H=z(1+2z)(1+4z+2z^2)^{-1}$, $\alpha_{12}^H=z(1+4z+2z^2)^{-1}$, and the superscript H denotes the quantities obtained in the hydrodynamic approximation. In the Padé approximation the expression for δn coincides with that obtained in the framework of the fully kinetic treatment equation (3) when $\alpha_{11}\simeq\alpha_{11}^H$ and $\alpha_{12}\simeq\alpha_{12}^H$. Therefore, the dispersion relation of the flute waves and the expression for the RTI normalized growth rate are

$$\omega^2-\omega\left(\omega_N+\frac{\alpha_{12}^H}{\alpha_{11}^H}\omega_T+\omega_g\right)+\frac{k_y^2}{k_\perp^2}\frac{z}{\alpha_{11}^H}\kappa_N g=0 \tag{18}$$

and

$$\frac{\gamma}{\Gamma}=\left[\frac{z}{\alpha_{11}^H}-\frac{z}{4\sigma}\left(1+\frac{\alpha_{12}^H}{\alpha_{11}^H}\eta+\sigma\right)^2\right]^{1/2}. \tag{19}$$

Figures 1–4 show the dependence of the normalized growth rates γ/Γ versus $k_\perp\rho_i$. The parameter $\sigma=\omega_g/\omega_N=g/\kappa_N v_T^2$ takes the values 1/5 and 1/2. The solid, dashed, and dashed-dotted lines correspond to classic (large scale), kinetic, and hydrodynamic (Padé approximation) approaches, respectively. Figures 1 and 3 correspond to plasmas with uniform ion temperature, $\eta=\kappa_T/\kappa_N=0$, and Figs. 2 and 4 to plasmas with $\eta=0.5$. It is seen that the normalized growth rates obtained in the kinetic and model hydrodynamic approaches have qualitatively similar dependencies on $k_\perp\rho_i$ as compared to the classical results obtained in the long wavelength approximation.[3–5] They differ only quantitatively. The insta-

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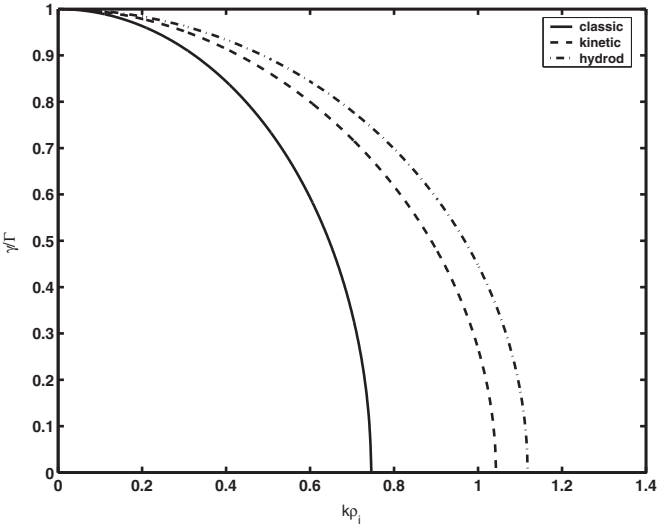

FIG. 1. The normalized RTI growth rate as a function of $k_\perp \rho_i$ in plasmas with homogeneous ion temperature ($\eta=0$) at $\sigma=g/\kappa_N v_{Ti}^2$ equal to 1/5. The solid, dashed, and dashed-dotted lines correspond to the classic (long-scale approximation), kinetic, and model hydrodynamic (Padé approximation) approximations, respectively.

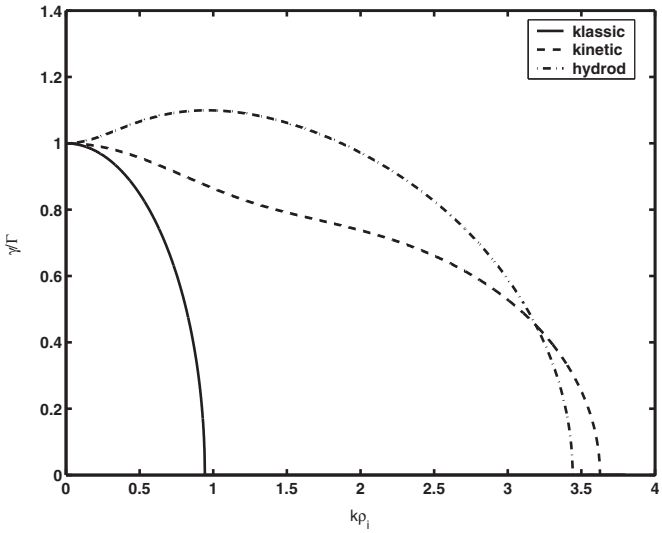

FIG. 3. The normalized RTI growth rate as a function of $k_\perp \rho_i$ in plasmas with homogeneous ion temperature ($\eta=0$) at $\sigma=1/2$. The solid, dashed, and dashed-dotted lines describe the results obtained in the classic, kinetic, and model hydrodynamic approximations, respectively.

bility region in the **k** space defined in kinetics or model hydrodynamics is wider than that obtained in Refs. 3–5. One can see that the results of the model hydrodynamic description are in reasonable agreement with the kinetics in a plasma with $\sigma=1/5-1/2$. Comparing Figs. 1 and 3 with Figs. 2 and 4 one sees that the instability regions become smaller with an increase in η.

IV. SUMMARY

The present analysis is an extension of the previous study of nonlinear flute waves in a low-β plasma with nonzero ion temperature gradient that was limited to consideration of waves with spatial scales larger than the ion Larmor radius ρ_i. The results of a kinetic description of the flute waves and RTI have been generalized to the case of arbitrary wave scales. We have extended the previous analysis of nonlinear flute waves[6–13] using two-fluid hydrodynamics that describes waves with arbitrary spatial scales in the so-called Padé approximation. A new closed set of Eqs. (11), (14), and (15) describing the nonlinear dynamics of flute waves with arbitrary spatial scales and incorporating the ion temperature perturbations has been derived. Particular attention has been paid to waves with spatial scales of the order of the ion Larmor radius, $k_\perp \rho_i \simeq 1$. Furthermore, it was shown that magnetic RTI is not stabilized in a low β plasma when $\sigma=1$.

Our two-fluid magnetohydrodynamics which adequately describes the ion perturbations in low beta plasmas with arbitrary ratio of the characteristic spatial scales to the ion Larmor radius can give us a possibility for numerical simulation of nonlinear stage of the magnetic RTI.

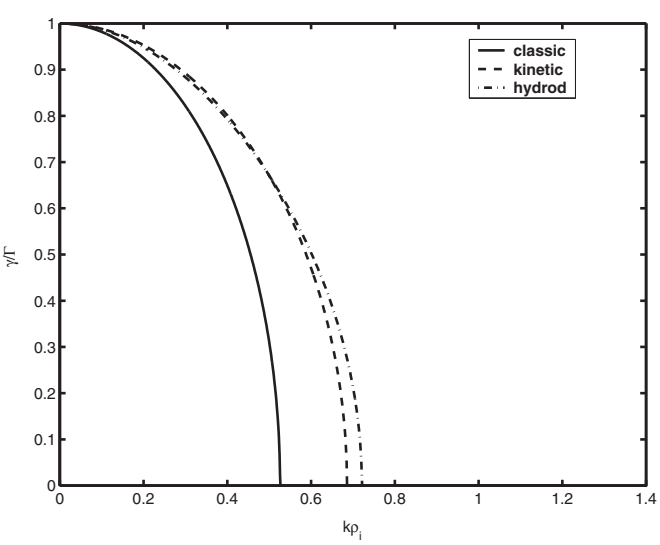

FIG. 2. The same as in Fig. 1 but for a plasma with inhomogeneous equilibrium ion temperature when $\eta=0.5$.

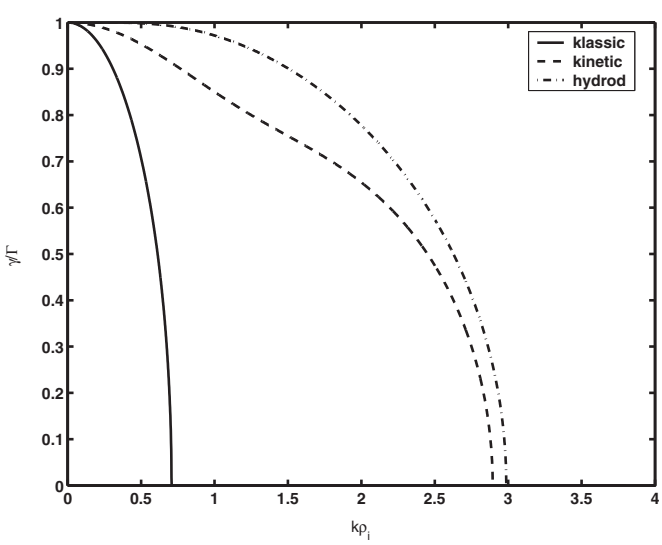

FIG. 4. The same as in Fig. 3 but in a plasma with inhomogeneous equilibrium ion temperature when $\eta=0.5$.

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