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On the Stability of Interchange Mode Turbulence with respect to Excitation of Large-Scale Flows

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Abstract. The coupled dynamics of small- and large-scale flute type perturbations in a magnetized plasma is considered. It is shown that the small-scale wave packets are unstable with respect to the low frequency and long wavelength perturbations. These perturbations are accompanied by the excitation of the long wavelength modes of the density and velocity. One can distinguish two regimes of this parametric instability similar to the case of beam-plasma instability. The instability is of the resonant type (*kinetic* regime), when the instability growth rate is smaller than the spectral width of the small-scale flute fluctuations. In the opposite case, the instability becomes the coherent hydrodynamic type (*hydrodynamic* regime), so that all harmonics grow coherently. It is suggested this generic instability is responsible for the generation of mean flows, so called zonal flows and streamers, in magnetized plasma.

I. Introduction.

It is already well accepted that nonlinear interaction of short-scale fluctuations in magnetized plasma can generate low-frequency, large-scale nonlinear structures (so called zonal flows and streamers) that play an important role in controlling plasma transport properties in magnetic confinement systems.¹⁻³ The transfer of wave energy towards the long wave length region and the formation of large-scale structures are results of the well-known inverse cascade guaranteed in two-dimensional (and quasi-two-dimensional) fluids by the conservation of energy and enstrophy.³ Zonal flows⁴ are defined here as symmetric structures with a finite scale in the direction of plasma inhomogeneity, significantly larger than the scale of the underlying small-scale turbulence. Streamers⁵ are convective cells which complement zonal flows, in that they have short extent in the direction of translation symmetry, i.e. perpendicular to plasma gradients, and are elongated structures in the direction of plasma inhomogeneity. Detailed investigation of the origin and dynamics of these flows is currently of great importance and presents a major challenge to the theory of plasma turbulence.

The mechanism for flow acceleration is commonly attributed to the effect of Reynolds forces generated by small-scale fluctuations.¹ It was shown that the flow acceleration will occur if the underlying small-scale turbulence supports waves propagating in the direction of plasma inhomogeneity, and gradients in the turbulent Reynolds stress have to exist. When the large-scale structures are excited, they form an environment for the parent waves. This leads us to a scenario of self-regulated plasma turbulence in which long wavelength structures “live” in a background of short wavelength turbulence. On the other hand, these large-scale structures can modulate and regulate the turbulence dynamics itself via nonlinear coupling to the small-scale fluctuations. So, the parent waves and secondary flows form a self-regulating system and cannot be addressed in isolation.⁶

Here, we investigate the *coupled* dynamics of large-and small-scale flute perturbations in magnetized plasma. For this turbulence the contribution of density fluctuations and finite

ion Larmor radius effects are significant and must be taken into account.⁷ The evolution equations for mean flow generation are obtained by averaging the model equations for flute modes over fast small-scales. These flows are, effectively, sheared flow layers which strain and distort the microscopic fluctuations they coexist with.⁸ Therefore the principal self-regulatory mechanism of wave turbulence is shearing by self-generated flows. Then, the propagation of flute modes in weakly inhomogeneous media with slowly varying parameters due to flow induced velocity and density perturbations can conveniently be described with the help of a wave kinetic equation.

The flow generation and shearing occur via triad interaction, i.e. those with $\mathbf{k} + \mathbf{k}' + \mathbf{q} \sim 0$. We show, having this fact in the mind, that the spontaneous excitation of large-scale flows by small-scale flute mode turbulence can be described as instability, similar to a parametric instability of a pump wave.⁹ Namely, the small-scale wave packets are unstable with respect to the long wavelength perturbations. These perturbations are accompanied by the excitation of the long wavelength modes of the density and velocity, i.e. zonal flows and streamers. One can distinguish two regimes of the instability similar to the case of plasma-beam instabilities.^{9,10} The instability is of the resonant type (*kinetic regime*), when the instability growth rate is smaller than the spectral width of the small-scale flute fluctuations. In the opposite case, the instability becomes the coherent hydrodynamic type (*hydrodynamic regime*), so that all harmonics grow coherently.

The rest of the paper is organized as follows: In Sec.II, the model equations for flute oscillations are presented and linear properties of flute modes are outlined. In Section III, the nonlinear dynamics of coupled system is investigated. It includes formulation of evolution equations for zonal flows and streamers generation, and the wave kinetic equation describing the interaction of the small-scale flute turbulence and large-scale shear flows. The quasi-linear and stability analysis is performed in Sec.IV. Finally, the short discussion of the results obtained is given in Section V.

II. Basic equations

Flute modes are low frequency oscillations ($\omega/\Omega_i \ll 1$ with $\Omega_i = eB_0/m_i c$ and ω being the frequency of the flute mode) of a weakly inhomogeneous magnetized plasma that are uniform in the direction of the external magnetic field ($k_{\parallel} = 0$). To describe the flute modes in such a plasma we use the two-fluid macroscopic equations as basic equations. The analysis will be performed for the most straightforward case of the slab geometry $(x, y) \equiv \mathbf{r}$. Suppose that such a quite dense ($4\pi n_0 m_i c^2/B_0^2 > 1$) plasma is immersed in a constant magnetic field $\mathbf{B}_0 = (0, 0, B_0)$. The curvature of magnetic field lines which always exists in magnetic confinement devices is imitated by the fictional gravity $\mathbf{g} = (g(x), 0, 0)$. Because of treating a low- β plasma, we assume the electric field of the perturbation to be potential, $\mathbf{E} = -\nabla\Phi$. The plasma density is expressed as $N(\mathbf{r}, t) = n_0(x) + \delta n(\mathbf{r}, t)$, where the equilibrium plasma density $n_0(x)$ is assumed to be slowly varying with x over a characteristic length-scale. For the low-frequency perturbations we can neglect the charge separation and, instead of the Poisson equation, use the condition of quasineutrality.⁷

It is well known that the finite ion temperature effect is important for the theory of flute modes. Therefore the ion fluid is considered as having a finite temperature, for simplicity $T_i = T_e = T$ here. Then, to the first order in the parameter ω/Ω_i we reduce our basic equations to a pair of coupled nonlinear equations for the dimensionless density $n(\mathbf{r}, t) = \delta n/n_0$ and the potential $\Phi(\mathbf{r}, t)$,¹¹

$$\frac{\partial n}{\partial t} + \frac{\kappa c}{B_0} \frac{\partial \Phi}{\partial y} = -\frac{c}{B_0} \{\Phi, n\}, \quad (1)$$

$$\frac{c}{B_0} \left[\frac{\partial}{\partial t} \nabla_{\perp}^2 \Phi - v_* \frac{\partial}{\partial y} \nabla_{\perp}^2 \Phi \right] + g \frac{\partial n}{\partial y} = -\frac{c^2}{B_0^2} \left[\{\Phi, \nabla_{\perp}^2 \Phi\} + \frac{T}{e} \text{div}\{n, \nabla_{\perp} \Phi\} \right]. \quad (2)$$

Here we have neglected the inertia of electrons, the effect of a temperature gradient and the collisional viscosity of the stress tensor; $v_* = \frac{\kappa c T}{e B_0}$ the ion diamagnetic drift velocity,

$\kappa \equiv -\frac{d \ln n_0}{dx} > 0$ and $g = \frac{T}{m_i R}$, R is the characteristic-scale length of magnetic field

inhomogeneity. The $\{\}$ brackets denote the Poisson bracket, $\{a, b\} = \mathbf{z} \cdot [\nabla a \times \nabla b]$, and $\nabla \equiv \nabla_{\perp}$. The finite Larmor radius (FLR) effect is incorporated to the lowest order by the term proportional to T .

Analysis of Eqs.(1)-(2) shows that there are important differences in our model in contrast to the well-known Hasegawa-Mima equation describing the electrostatic drift wave turbulence. First, basic equations consist of two fields associated with the perturbed potential and the perturbed density. Second, for drift wave turbulence the nonlinearity appears from the ion polarization drift. For the flute mode turbulence the contribution of density fluctuations to the ion polarization drift is significant and is taken into account by the second term on RHS of Eq.(2). The polarization drift nonlinearity is well studied in the drift model and is known to cascade power towards long-scales. However, the diamagnetic component of polarization nonlinearity is known to sufficiently modify cascade properties. Moreover, the right-hand side (RHS) of Eq.(1) contains also so-called convective $\mathbf{E} \times \mathbf{B}$ nonlinearity that behaves in exactly the opposite fashion cascading power into the shorter scale length.^{12,13}

Since the linear physics of the flute modes is well documented,⁷ we provide only a brief overview here. Linearizing Eqs.(1)-(2) for small perturbations $(n, \Phi) \sim \exp(-i\omega t + i\mathbf{k} \cdot \mathbf{r})$, we can obtain the dispersion relation for the flute modes with the wave (linear) eigenfrequency being

$$\omega_k = \omega_k^{\text{Re}} + i\gamma_k, \quad \omega_k^{\text{Re}} = -\frac{k_y v_*}{2}, \quad \gamma_k = \frac{k_y v_*}{2} \gamma, \quad \gamma = \left(\frac{4L_n}{R} \frac{1}{k^2 \rho_i^2} - 1 \right)^{1/2}, \quad (3)$$

where $k^2 = k_x^2 + k_y^2$, $k \gg \kappa$, and $\rho_i = V_{Ti} / \Omega_i$ is the ion Larmor radius. If $T = 0$, then Eq.(3)

describes pure unstable flute modes with a growth rate $\gamma_0 = \frac{k_y}{k} \sqrt{\kappa g}$. It is clear that the linear growth rate vanishes for modes with $k_y = 0$. Thus, convective cells that have short poloidal extend can be generated by the linear instability mechanism, whereas poloidally symmetric flows can be generated only through some nonlinear interaction mechanism. Accounting for the FLR effect leads to the stabilization of flute instabilities for $k^2 \rho_i^2 \geq 4L_n/R$, where $L_n \equiv \kappa^{-1}$ and R are the characteristic-scale lengths of plasma and magnetic field inhomogeneities, respectively.

III. Coupled dynamics of small-and large-scale flute modes

To understand nonlinear dynamics of the system under consideration we now carry out an analytical analysis of Eqs.(1)-(2). To begin with we employ the ansatz of scale separation

to distinguish between the small-scale fluctuations of flute modes and the large-scale shear flow pattern generated in the final state. To describe the evolution of the coupled system (wave turbulence + large-scale plasma flows) we represent the electrostatic potential Φ as a sum of a large-scale flow $\bar{\Phi}$ quantity and a small-scale turbulent part $\tilde{\Phi}$. A similar representation has been chosen for the plasma density:

$$\Phi = \bar{\Phi} + \tilde{\Phi}, \quad n = \bar{n} + \tilde{n}. \quad (4)$$

As we are interested in mode coupling, it is useful to expand $\tilde{\Phi}$ and \tilde{n} into a Fourier series:

$$\tilde{\Phi}(\mathbf{r}, t) = \sum_k \Phi_k(t) \exp(i\mathbf{k} \cdot \mathbf{r}), \quad \tilde{n}(\mathbf{r}, t) = \sum_k n_k(t) \exp(i\mathbf{k} \cdot \mathbf{r}). \quad (5a)$$

The requirement that Φ and n must be real imposes the following constraints on the complex Fourier amplitudes: $\Phi_{-k} = \Phi_k^*$, $n_{-k} = n_k^*$. The large-scale flows can likewise be represented as

$$\bar{\Phi}(\mathbf{r}, t) = \Phi_q(t) \exp(i\mathbf{q} \cdot \mathbf{r}) + c.c., \quad \bar{n}(\mathbf{r}, t) = n_q(t) \exp(i\mathbf{q} \cdot \mathbf{r}) + c.c. \quad (5b)$$

where the condition $|q| \ll |k|$ holds. Zonal flow structures are characterized by $q_x \neq 0$ and streamer-like structures have $q_y \neq 0$. In order to describe the formation of these large-scale flows by small-scale flute modes, we perform the average over the fast-scales the basic Eqs. (1),(2). Then for zonal flows, the evolution equations (1) and (2) become decoupled and are reduced to

$$\frac{\partial n_q}{\partial t} = 0, \quad (6)$$

$$\frac{\partial \Phi_q}{\partial t} = \frac{c}{B_0} \sum_k k_x k_y \left(1 - \frac{R}{2L_n} k^2 \rho_i^2 \right) |\Phi_k|^2. \quad (7)$$

It is explicitly seen from Eq.(6) that the mean plasma density associated with zonal flows does not evolve with time, i.e. can be considered as a constant.

Corresponding evolution equations for streamers after averaging over the fast-scales have the form

$$\frac{\partial n_q}{\partial t} + i q_y \frac{\kappa c}{B_0} \Phi_q = 0, \quad (8)$$

$$\left(\frac{\partial}{\partial t} - i q_y v_* \right) \Phi_q - i \frac{g B_0}{q_y c} n_q = - \frac{c}{B_0} \sum_k k_x k_y \left(1 - \frac{R}{2L_n} k^2 \rho_i^2 \right) |\Phi_k|^2. \quad (9)$$

It is clear from Eqs. (8)-(9) that density and potential fluctuations associated with streamers remain coupled, and dynamics of streamers is covered by more complicated equations in comparison with the case of zonal flows. Eq.(8) shows that, in contrast to the case of zero-frequency zonal flows, streamers always have a real frequency, sufficiently lower however than that of parent waves.

The flute mode turbulent medium is now a weakly inhomogeneous medium with slowly varying parameters. The sources of this slow spatial and temporal variations are flow induced velocity and density perturbations. The propagation of flute modes in such weakly inhomogeneous medium can be conveniently described with the help of a wave kinetic equation for the wave action density $N_k(\mathbf{r}, t)$ in \mathbf{r} - \mathbf{k} space¹⁴

$$\frac{\partial N_k}{\partial t} + \frac{\partial \omega_{k,r}^{NL}}{\partial \mathbf{k}} \frac{\partial N_k}{\partial \mathbf{r}} - \frac{\partial \omega_{k,r}^{NL}}{\partial \mathbf{r}} \frac{\partial N_k}{\partial \mathbf{k}} = 2\gamma_k N_k - St\{N_k\}. \quad (10)$$

where

$$N_k = k^2 \rho_i^2 \frac{4R}{L_n} \left(1 - k^2 \rho_i^2 \frac{R}{4L_n} \right) \left| \frac{e\Phi_k}{T} \right|^2$$

The linear frequency of flute modes entering this equation is modified in the presence of the flows because of the Doppler shift induced by the flow velocity,

$$\omega_{k,r}^{NL} = \omega_k^{\text{Re}} + \mathbf{k}(\mathbf{V}_0 + \mathbf{V}_1), \quad \mathbf{V}_0 = -\frac{c}{2B_0} [\nabla \bar{\Phi} \times \mathbf{z}], \quad \mathbf{V}_1 = -\frac{cT}{4eB_0} [\nabla \bar{n} \times \mathbf{z}]. \quad (11)$$

Here \mathbf{V}_0 is associated with the Doppler effect introduced by $\mathbf{E} \times \mathbf{B}$ flows ($\bar{\Phi}$), and \mathbf{V}_1 corresponds to the lowest-order FLR corrections due to finite ion temperature. Thus, the small-scale turbulence will be sheared by both the flow shear and the diamagnetic effect.

IX. Quasilinear and stability analysis

a) Quasilinear theory

We assume that the wave spectrum consists of an equilibrium part $N_k^{(0)}$ averaged over the short spatial scales, but evolving in time and \mathbf{k} , and a perturbed part \tilde{N}_k , i.e. $N_k = N_k^{(0)} + \tilde{N}_k$. For simplicity, we take $St(N_k)$ in the simplest, quadratic form, $St(N_k) = \Delta\omega_k N_k^2$, with $\Delta\omega_k$ being a nonlinear damping rate, accounting for local nonlinear coupling. Thus, the stationary spectrum can be established in the equilibrium, $N_k^{(0)} = 2\gamma_k / \Delta\omega_k$, and the RHS of Eq.(10) is close to zero.¹⁰

The quantity \tilde{N}_k may be straightforwardly calculated via linearization of the wave kinetic equation (10):

$$\frac{\partial \tilde{N}_k}{\partial t} + \mathbf{v}_g \frac{\partial \tilde{N}_k}{\partial \mathbf{r}} - \frac{\partial}{\partial \mathbf{r}} (\mathbf{k}\mathbf{v}_f) \cdot \frac{\partial N_k^{(0)}}{\partial \mathbf{k}} = -2\gamma_k \tilde{N}_k \quad (12)$$

whereas quasi-linear equation for $N_k^{(0)}$ has a form:

$$\frac{\partial N_k^{(0)}}{\partial t} - \frac{\partial}{\partial \mathbf{r}} (\mathbf{k}\mathbf{v}_f) \cdot \frac{\partial \tilde{N}_k}{\partial \mathbf{k}} = 0 \quad (13)$$

with $(\overline{ab..})$ in the second term meaning again the averaging over the short spatial scales. Then, assuming the temporal ($\Omega \ll \omega_k$) and spatial ($\mathbf{q} \ll \mathbf{k}$) variations for small perturbations $(\tilde{N}_k, \bar{\Phi}, \bar{n}) \sim \exp(-i\Omega T + i\mathbf{q} \cdot \mathbf{r})$, modulation of \tilde{N}_k is calculated from Eq.(12), to yield

$$\tilde{N}_k = -\frac{c}{2B_0} k_y q_x^2 \Phi_q \frac{\partial N_k^{(0)}}{\partial k_x} R(\Omega, q_x), \text{ for zonal flows} \quad (14a)$$

$$\tilde{N}_k = \frac{c}{2B_0} k_x q_y^2 \Phi_q \frac{\partial N_k^{(0)}}{\partial k_x} R(\Omega, q_y), \text{ for streamers} \quad (14b)$$

Here the quantity $R(\Omega, q)$ is the response function,

$$R(\Omega, q) = \frac{i}{\Omega - \mathbf{q}\mathbf{v}_g + i\delta\omega_k}, \quad (15)$$

with \mathbf{v}_g being the flute mode group velocity, and $\delta\omega_k$ is the total decorrelation frequency which may also include the linear growth rate and a nonlinear frequency shift. In the weakly nonlinear regime

$$R(\Omega, q) \rightarrow \pi\delta(\Omega - \mathbf{q}\mathbf{v}_g).$$

For the wide spectrum of fluctuations, one obtains

$$R(\Omega, q) \rightarrow 1/\delta\omega_k.$$

The broad spectrum of large-scale structures regulates turbulence by the process of random shearing, which may be represented, in the spirit of quasilinear theory, by diffusion in \mathbf{k} space. Such shearing is now understood to be the key mechanism that governs the self-regulation and saturation of wave turbulence.

Substituting the expression for \tilde{N}_k into the Eq.(13) and using the standard quasi-linear theory, we can rewrite the equation for $N_k^{(0)}$ as diffusion-like equation:

$$\frac{\partial N_k^{(0)}}{\partial t} - \frac{\partial}{\partial \mathbf{k}} D_{\mathbf{k}} \frac{\partial N_k^{(0)}}{\partial \mathbf{k}} = 0, \quad (16)$$

where $D_{\mathbf{k}}$ notes the diffusivity of waves by large-scale flows. Indeed, the \mathbf{k} -space diffusion coefficient for the zonal flows ($q_x \neq 0$) after simple manipulations can be written in the form:

$$D_{k_x}^{ZF} = \frac{c^2 T^2}{4e^2 B_0^2} \sum_q k_y^2 q_x^4 \left(\left| \frac{e\Phi_q}{T} \right|^2 + \frac{1}{4} |n_q|^2 \right) R(\Omega, q). \quad (17)$$

Thus, the zonal flows constitute a random strain field which randomly refracts flute modes, causing a diffusive increase in k_x which in turn enhances their coupling to small-scale (high- k_x) dissipation.

In contrast to zonal flows, shearing by streamer structures ($q_y \neq 0$) refers to poloidal shearing of radial streamer flows, rather than the usual process of radial shearing of poloidal flows. In this case, the stochastic refraction of eddies by poloidally sheared streamer flows is described as

$$D_{k_y}^{St} = \frac{c^2 T^2}{4e^2 B_0^2} \sum_q k_x^2 q_y^4 \left(\left| \frac{e\Phi_q}{T} \right|^2 + \frac{1}{4} |n_q|^2 \right) R(\Omega, q). \quad (18)$$

The second term in parentheses of Eqs.(17),(18) comes from density modulation (diamagnetic effect).

b) Stability analysis

The spontaneous excitation of large-scale flows by small-scale flute mode turbulence can be described as *instability*, similar to a parametric instability of a pump wave: the small scale wave packets are unstable with increment γ_q with respect to the long wavelength perturbation. We distinguish two regimes of this instability. The instability is of resonant type (*kinetic* regime) if the growth rate $\gamma_q < \delta\omega_k$, the spectral width of the flute mode turbulence. In this case, coupled Eqs. (6) – (9) and (12) can be solved to show that modulations of the wave packets and flows are unstable. Indeed, using (14a) and (14b) in Eqs.(6) – (9) we obtain the following growth rate

$$\gamma_q^{ZF} = -\left(\frac{cT}{eB_0}\right)^2 \frac{L_N}{R} q_x^2 \int d\mathbf{k} \frac{k_y^2}{8k^2 \rho_i^2} k_x \frac{\partial N_k^{(0)}}{\partial k_x} R(\Omega, q), \quad (19a)$$

for zonal flows, and

$$\gamma_q^{Sr} = -\frac{\Omega^{Sr}}{2\Omega^{Sr} + q_y v^*} \left(\frac{cT}{eB_0}\right)^2 \frac{L_N}{R} q_y^2 \int d\mathbf{k} \frac{k_x^2}{8k^2 \rho_i^2} k_y \frac{\partial N_k^{(0)}}{\partial k_y} R(\Omega, q) \quad (19b)$$

for streamers. Here Ω^{Sr} is the real frequency of streamers defined by

$$\Omega^{Sr} = -\frac{q_y v^*}{2} \left(1 \mp \sqrt{1 - \frac{4L_n}{R} \frac{1}{q_y^2 v_*^2}} \right)$$

The resonant type of instability is obtained from (19a) and (19b) by using the resonant response function $R(\Omega, q) = i/(\Omega - \mathbf{q}\mathbf{v}_g + i\delta\omega_k) \rightarrow \pi\delta(\Omega - \mathbf{q}\mathbf{v}_g)$. Then, the instability may be interpreted as a result of the resonant interaction of the flute wave packet with slow modulations of the mean flow. The similarity with Landau resonance in the kinetic plasma theory is evident. In both cases, the conditions for growth, $\gamma_q > 0$, are equivalent to

inequalities $k_x \frac{\partial N_k^{(0)}}{\partial k_x} < 0$ and $k_y \frac{\partial N_k^{(0)}}{\partial k_y} < 0$ for zonal flows and streamers correspondingly.

These conditions are typically satisfied in flute mode turbulence and therefore no population inversion is required.

When the growth rate γ_q becomes large compared to the characteristic frequency spread $\delta\omega_k$ for the background flute mode fluctuations, individual N_k components contribute to the instability coherently, so called *hydrodynamic* regime. In order to elucidate the physics of this instability, we consider a simple case of a monochromatic wave packet with $N_k^{(0)} = N_0 \delta(\mathbf{k} - \mathbf{k}_0)$. Then, the dispersion relation for the instability (we consider the zonal flow excitation only) takes the form

$$(\Omega - q_x v_g)^2 = \frac{L_n}{16R} q_x^2 k_y^2 v_T^2 \frac{N_0}{k_y v_*} \frac{\partial v_g}{\partial k_x} \quad (20)$$

It is seen from Eq.(20) that the criterion for the instability is

$$\frac{N_0}{k_y} \frac{\partial v_g}{\partial k_x} < 0 \quad (21)$$

and thus reduces to a general Lighthill criterion $\frac{\partial^2 \omega_k}{\partial k_x^2} < 0$ for modulation instability. It can

readily be seen that the growth rate of the coherent instability, $\gamma_q^{coh} \propto q\sqrt{N_0}$, is large

compared to that of resonant instability, $\gamma_q^{res} \propto q^2 k \frac{\partial N_k^{(0)}}{\partial k}$.

Y. Conclusions

We have shown by performing the analysis similar to one used in the theory of parametric instabilities that small-scale flute mode packets are unstable with respect to the long wavelength perturbations. These perturbations are accompanied by the excitation of the long wavelength modes of velocity and density, i.e. zonal flows and streamers.

Depending on the relation between the nonlinear growth rate and the spectral width of the background flute mode turbulence, the instability is of the resonant type (kinetic regime) or coherent type (hydrodynamic regime). These mechanisms constitute two regimes of the

same process (generation of large-scale flows), similar to the case of beam-plasma instabilities.

It follows from the analysis, that the kinetic regime of the generation of large-scale flows is characterized by the growth rate smaller than that of hydrodynamic regime. Moreover, the criterion for the coherent instability is similar to the general Lighthill criterion for modulation instability.

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