Large Scale Instabilities in the Electromagnetic Drift Wave Turbulence and Transport

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Large scale structures play an important role in self-organization of drift wave turbulence. Large scale perturbations of plasma flow and magnetic field are spontaneously generated in generic electromagnetic drift wave turbulence via the action of Reynolds stress and electromotive force. Initial large scale perturbations are amplifed by positive feedback due to the modulations of wave packets by the shearing effect of the large scale flow and/or by the perturbed large scale magnetic field. As a result, the propagation of small scale wave packets is accompanied by the instability of a low frequency, long wavelength components. Anomalous transport due to drift wave turbulence may also be unstable with respect to the large scale perturbations of plasma profile. In this case the instability occurs as a result of a positive feedback response of the anomalous flux to the large scale variations of plasma temperature.

INTRODUCTION

It has recently been realized that generation of the large scale shear flow (zonal flow) plays an important role in self-regulation of the drift wave turbulence [1-3]. Spontaneously excited large scale flows occur as a result of the intrinsic non-ambipolarity of the radial plasma flow, or, in other words, due to the radial momentum flux [4,5]. The transfer of wave energy towards the long wavelength region and the formation of large scale structures (zonal flows and convective cells) may be viewed as a result of the inverse cascade in two-dimensional and quasi two-dimensional geostrophic fluids [6]. The strongly sheared flow associated with localized structures leads to turbulence suppression and enhancement of confinement in a tokamak. Zonal flows are defined here as poloidal and toroidally symmetric ($q_z = q_\theta = 0$) perturbations with a finite radial scale q_r^{-1} larger than the scale of the underlying small scale turbulence, $q_r \ll k_r$, $\mathbf{q}$ is the wave vector for large scale motions, $\mathbf{k}$ is the wave-vector of small scale turbulence, and r, θ, and z are axis of a straight cylindrical tokamak.

Opposite to zonal flows is another class of coherent large scale structures, streamers which are radially elongated formation with short poloidal wavelength, $q_\theta >> q_r$. The latter structures may be an underlying mechanism for nonlocal, radially extended transport events or avalanches that have been recently identified in turbulent transport [7-10]. Both types of structures are nonlinearly generated in drift wave turbulence.

In the electromagnetic case, generation of the large scale magnetic field is also possible [11-13] and can be considered as a fast dynamo process [14]. Magnetic structures (magnetic islands and magnetic streamers) have been studied as a possible mechanism of regulation and enhancement of the electron transport [15,16].

We consider the dynamics of small scale drift wave turbulence coupled with slow, large scale perturbations of the electrostatic potential, magnetic field, and plasma temperature. We demonstrate how the large scale coherent structures in plasma flow, magnetic field and plasma temperature may be spontaneously generated. For simplicity we consider them separately, though in general case, the large scale magnetic field, shear flows and transport events may be coupled [17].

SHEAR FLOW INSTABILITY

Instabilities of the shear flow can be considered by using a simple two-dimensional model for electron drift waves

$$\left(\frac{\partial}{\partial t} + \mathbf{V}_0 \cdot \nabla\right) \frac{e\widetilde{\phi}}{T_e} + \mathbf{V}_* \cdot \nabla \frac{e\widetilde{\phi}}{T_e}$$
$$-\rho_s^2 \left(\frac{\partial}{\partial t} + \mathbf{V}_0 \cdot \nabla + \widetilde{\mathbf{V}}_E \cdot \nabla\right) \nabla_\perp^2 \frac{e\phi}{T_e} = 0. \quad (1)$$

Here, ρ_s^2 is the ion-sound Larmor radius and $\mathbf{V}_* = \widehat{\theta} V_*$ is the electron diamagnetic drift velocity. Nonlinear equation (1) is similar to the Hasegawa-Mima model except the term $\mathbf{V}_0 \cdot \nabla(e\widetilde{\phi}/T_e)$ which is retained because the plasma density do not follow Boltzman distribution for large scale modes (similar situation occurs in the sheared magnetic field for modes with $k_\| \to 0$ at the rational surface). Coupled dynamics of large scale flow and small scale turbulence are considered, so that the electrostatic potential ϕ is a sum of fluctuating $\widetilde{\phi}$ and mean $\overline{\phi}$ quantities, $\mathbf{V}_0 = c\mathbf{b} \times \nabla\overline{\phi}/B_0$, $\widetilde{\mathbf{V}}_E = c\mathbf{b} \times \nabla\widetilde{\phi}/B_0$. The mean potential is the average of the total potential over fast, small scale variables and depends only on slow variables $\mathbf{X}$ and T, $\overline{\phi} = \overline{\phi}(\mathbf{X}, T)$. The evolution equation for the mean flow is

$$\frac{\partial}{\partial T}\ \nabla_\perp^2 \overline{\phi} = -\frac{c}{B_0}\ \overline{\left\{\widetilde{\phi}, \nabla_\perp^2 \widetilde{\phi}\right\}}\ , \quad (2)$$

which shows that the large scale flow is driven by small scale fluctuations via Reynolds stress forces. Here, $\{a, b\} = \partial_r a \partial_\theta b - \partial_\theta a \partial_r b$ is the Poisson bracket.

Coupling of small scale fluctuations to the mean flow is described by the kinetic equation for wave packets

$$\frac{\partial N_k}{\partial T} + \frac{\partial \omega_k}{\partial \mathbf{k}} \cdot \frac{\partial N_k}{\partial \mathbf{X}} - \frac{\partial \omega_k}{\partial \mathbf{X}} \cdot \frac{\partial N_k}{\partial \mathbf{k}} = S, \qquad (3)$$

where $N_k = N_k(\mathbf{X}, T)$ is the adiabatic action invariant, and the exact form of N_k is model dependent. For the model given by Eq. (1), the wave frequency is $\omega_k = k_\theta V_0 + \omega_k^l$, where $\omega_k^l = k_\theta V_*/(1 + k_\perp^2 \rho^2)$ is the local wave frequency, and the mean flow $\mathbf{V}_0$ enters the total frequency ω_k as a simple Doppler shift.

The source term in (3) describes the wave growth and damping due to linear and nonlinear mechanisms. We assume that small scale turbulence is close to a stationary state, so that $S \to 0$.

Coupled equations (2,3) can be solved to show that the modulations of the wave packets and zonal flow $\mathbf{V}_0$ are unstable [2]. We consider equations (2),(3) linearized for small perturbations $(\widetilde{N}_k, \overline{\phi}) \sim \exp(-i\Omega T + iqr)$,where $q \equiv q_r = -i\partial/\partial r$ is the radial wave vector of the large scale perturbation. Then, Eq. (2) takes the form

$$-i\Omega\overline{\phi} = \frac{c}{B_0} \int k_r k_\theta \left|\phi_k\right|^2 d^2k. \qquad (4)$$

The modulation of $\widetilde{N}_k$ is calculated from (3)

$$\widetilde{N}_k = -\frac{c}{B_0} q^2 \overline{\phi} k_\theta \frac{\partial N_k^0}{\partial k_r} \frac{i}{\Omega - qV_g}, \qquad (5)$$

where $V_g = \partial\omega/\partial k_r$. Using (5) in (4), we obtain the following equation [2]

$$-i\Omega = -q^2 c_s^2 \int d^2k \, \frac{k_\theta^2 \rho_s^2}{(1 + k_\perp^2 \rho_s^2)^2} k_r \frac{\partial N_k^0}{\partial k_r} \frac{i}{\Omega - qV_{gr}}. \qquad (6)$$

The resonant type instability is obtained from (5) by using the resonant function $R = i/(\Omega - qV_g) \to \pi\delta(\Omega - qV_g)$ or its broadened counterpart $i/(\Omega - qV_g + \Delta\omega_k)$ (for a white noise source, this can be taken as $1/\Delta\omega_k$, where $\Delta\omega_k$ is nonlinear broadening due to the wave-wave interaction). For the case of the narrow resonant function approximated by a delta-function, the growth rate of the resonant instability is

$$\gamma_q = -q^2 c_s^2 \int d^2k \, \frac{k_\theta^2 \rho_s^2}{(1 + k_\perp^2 \rho_s^2)^2} k_r \frac{\partial N_k^0}{\partial k_r} \pi\delta(\Omega - qV_g). \qquad (7)$$

The condition $\partial N_k^0/\partial k_r < 0$ is required for instability.

When the growth rate of the instability becomes large compared to the characteristic frequency spread for the background fluctuations, individual N_k components contribute to the instability coherently. Insight into this mechanism can be provided by a simple case of a monochromatic wave packet with $N_k^0 = N_0\delta(\mathbf{k} - \mathbf{k}_0)$, with $\mathbf{k}_0 = (k_{r0}, k_{\theta 0})$. In this case after some transformations we obtain [18]

$$(\Omega - qV_{gr})^2 = q^2 c_s^2 \, k_\theta^2 \frac{N_k^0}{2k_\theta V_*} \frac{\partial V_g}{\partial k_r}, \qquad (8)$$

Note that the criterion for the instability is thus $N_k^0 (2k_\theta V_*)^{-1} \partial V_g/\partial k_r < 0$. It can readily be seen that the coherent ("hydrodynamic") instability has a larger growth rate compared to that of the resonant instability (7). Resonant growth of large scale perturbations can be described as negative viscosity instability [19]. In the "hydrodynamic" regime the instability if of the reactive type due to the interaction of negative and positive energy modes.

ROBUST FAST DYNAMO IN ALFVEN TURBULENCE

Modulational instability of small scale electromagnetic fluctuations may also lead to the generation of large scale magnetic structures in a turbulent magnetized plasma [11-17]. The large scale magnetic field is driven by the mean electromotive force term in Ohm's law, $\widetilde{\mathbf{v}} \times \widetilde{\mathbf{B}}$.

As an example we consider a collisionless Alfven wave turbulence in the presence of an ambient magnetic field $\mathbf{B}_0 = B_0\widehat{\mathbf{z}}$. A spontaneous excitation of large scale magnetic fields $\overline{\mathbf{B}} = \nabla\overline{\psi} \times \widehat{\mathbf{z}}$, where $\overline{\psi} = \overline{\psi}(\mathbf{X})$, is a result of coupling of small scale turbulence and the initial perturbation of the mean field. Large scale random magnetic field refracts wave packets of the Alfven waves and, thus, modulates spectrum of the turbulence. Modulated spectrum reacts back on the generated field via correlation between the perturbed small scale components of electrostatic and magnetic potentials. The latter provides electromotive force in the mean Ohm's law.

To consider the evolution of a large scale magnetic field we consider a simple case of the mean Ohm's law with the main contribution $\widetilde{\mathbf{v}} \times \widetilde{\mathbf{B}}$ and neglecting other terms which can be important, such as contrubutions due to the electron pressure

$$\frac{\partial}{\partial t}\overline{\psi} = -\frac{c}{B_0}\widehat{\mathbf{z}} \cdot \overline{\nabla\widehat{\phi} \times \nabla\widehat{\psi}}$$
$$= \frac{c}{B_0}\frac{\partial}{\partial X} R_y - \frac{c}{B_0}\frac{\partial}{\partial Y} R_x \,, \qquad (9)$$

where $R_y = \left(\overline{\sum \widehat{\phi}_k i k_y \widehat{\psi}_{-k}}\right)$, and $R_x = \left(\overline{\sum \widehat{\phi}_k i k_x \widehat{\psi}_{-k}}\right)$. It is obvious that a finite phase between ϕ and ψ is required for the generation of the large scale magnetic field. At a linear stage such a phase shift occurs due to mode grow/damping associated with dissipation effects. Our initial equations do not contain such dissipation/wave particle interaction effects. In two-fluid model dissipative effects will appear in the momentum balance equation as electron viscosity and in the energy balance equation as a heat flux. However, quasineutrality equation (or density conservation equation) in general form remains correct even in the presence of the dissipation. For illustration purposes we consider Alfven waves neglecting all drift effects. Therefore from the quasinetrality equations we have

$$\phi_k = \frac{k_\parallel c\,(\omega_k - i\eta_k)}{{\omega_k}^2 + \eta_k^2} \frac{v_A^2}{c^2}\psi_k, \qquad (10)$$

where ω_k is a real part of the wave frequency, and η_k is the imaginary part of the mode frequency. In the linear (transient) regime η_k is due to the mode growth rate/damping while in the nonlinear steady-state regime η_k is due to the nonlinear mode interaction and negative. As an estimate, eddy turn-over time can be used for this parameter. More detailed estimates can be done using various strong turbulence closure schemes [20].

Then we have

$$R_y = \overline{\sum \widehat{\phi}_k i k_y \widehat{\psi}_{-k}} = \sum \frac{k_{\|} \eta_k k_y}{{\omega_k}^2 + \eta_k^2} \left|\psi_k\right|^2, \qquad (11)$$

and similarly R_x can be calculated. The response of $\left|\psi_k\right|^2$ to the variations of the mean vector potential $\overline{\psi}$ can be found by using the wave kinetic equation in the form

$$\frac{\partial}{\partial t} N_k + \frac{\partial}{\partial \mathbf{k}} \left(\omega_k + \delta k_{\|} \frac{\delta \omega}{\delta k_{\|}}\right) \cdot \frac{\partial}{\partial \mathbf{x}} N_k - \frac{\partial}{\partial \mathbf{x}} \left(\omega_k + \delta k_{\|} \frac{\delta \omega}{\delta k_{\|}}\right) \cdot \frac{\partial}{\partial \mathbf{k}} N_k = 0. \qquad (12)$$

General solution of this equation can be wrritten as an integral along the characteristics $N_k = N_k(\mathbf{x}, \mathbf{k}, t) = N_k^0(\mathbf{k}_0, \mathbf{x}_0)$ where $\mathbf{x}_0 = \mathbf{x}_0(\mathbf{k}, \mathbf{x}, \mathbf{t})$ and $\mathbf{k}_0 = \mathbf{k}_0(\mathbf{k}, \mathbf{x}, \mathbf{t})$ are inverse of the characteristics $\mathbf{x} = \mathbf{x}(\mathbf{x}_0, \mathbf{k}_0, t)$ and $\mathbf{k} = \mathbf{k}(\mathbf{x}_0, \mathbf{k}_0, t)$ defined by the equations

$$\frac{d\mathbf{x}}{dt} = \frac{\partial}{\partial \mathbf{k}} \left(\omega_k + \delta k_{\|} \frac{\delta \omega}{\delta k_{\|}}\right), \qquad (13)$$

$$\frac{d\mathbf{k}}{dt} = -\frac{\partial}{\partial \mathbf{x}} \left(\omega_k + \delta k_{\|} \frac{\delta \omega}{\delta k_{\|}}\right). \qquad (14)$$

In the linear approximation, $\mathbf{k} = \mathbf{k}_0 + \delta \mathbf{k}$,

$$\frac{d\delta \mathbf{k}}{dt} = -\frac{\partial}{\partial \mathbf{x}} \left(\delta k_{\|} \frac{\delta \omega}{\delta k_{\|}}\right), \qquad (15)$$

where $\delta k_{\|} = -B_0^{-1} \widehat{\mathbf{z}} \cdot \nabla \overline{\psi} \times \mathbf{k}$, and $\delta \omega / \delta k_{\|} = k_{\|} v_A^2 / \omega$. Then in the linear approximation

$$\frac{\partial^2}{\partial t^2} \overline{\psi} = \frac{c}{B_0} \left(\frac{\partial}{\partial X} A_y - \frac{\partial}{\partial Y} \sum A_x\right), \qquad (16)$$

where

$$A_y = \sum \frac{k_{\|} k_y \eta_k}{{\omega_k}^2 + \eta_k^2} \frac{\partial}{\partial t} \delta N_k,$$

$$A_x = \frac{k_{\|} k_x \eta_k}{{\omega_k}^2 + \eta_k^2} \frac{\partial}{\partial t} \delta N_k.$$

Perturbations of the wave action density can be found

$$\delta N_k = N_k^0(\mathbf{k} - \delta \mathbf{k}) - N_k^0(\mathbf{k}) \simeq -\frac{\partial N_k^0}{\partial \mathbf{k}} \cdot \delta \mathbf{k}$$

Then for $\partial / \partial X = 0$ we have

$$\frac{\partial^2}{\partial t^2} \overline{\psi} = -\frac{\partial^3}{\partial Y^3} \overline{\psi} \sum \frac{\eta_k}{\omega_k^2 + \eta_k^2} \frac{k_x^2 k_z^2}{\omega} \frac{\partial N_k^0}{\partial k_y}. \qquad (17)$$

Remarkable is that the instability occurs for arbitrary signs of $\partial N_k^0 / \partial k_y$ and $\mathbf{q}$ that can be interpreted as a robust alpha-effect. Further investigation is warranted here to explore whether this type of dynamo may by a subject of quenching due to the backreaction of large scale magnetic field [21].

In the above analysis we have implicitly assumed that the wave action invariant can be written $N_k \sim \left|\psi_k\right|^2$. In general, more exact expression is required that will change the final expression for the growth rate, though it does not affect the basic mechanism of the instability.

LARGE SCALE TRANSPORT EVENT INSTABILITY

Recently, detailed studies of the distrubution of transport events in turbulent plasmas have shown that relatively infrequent but large transport events may provide a dominant contrubution to the overall plasma transport [7,8]. Dynamical model of anomalous transport due to weakly excited long wavelength modes, "avalanches" or streamers, has been suggested in Ref. 22. Nonlinear instability leading to the formation of strongly anisotropic, poloidally localized streamers has been investigated in Refs. 7 and 9. Here we study the large scale instability of plasma temperature profile as a result of a positive feedback from background small scale fluctuations. As an example of the generation of large scale instabilities in the anomalous transport we consider a model transport equation (temperature) with a local diffusive operator

$$\frac{\partial}{\partial t} T + \frac{\partial}{\partial x} D \frac{\partial}{\partial x} T = P. \qquad (18)$$

We assume that a local diffusion coefficient can be written in the form $D = D_0 \sum N_k$, where N_k is suitably normalized dimensionless action invariant $N_k \sim e \left|\phi_k\right|^2 / T$. We consider a stability of the diffusion equation (18) with respect to global perturbations on a scale length comparable to the device scale and much larger than the turbulence correlation length. Under this assumption the evolution of the background turbulent field can be modeled by the wave kinetic equation (3). For simplicity we consider a simple one dimensional case. Then general solution of (3) can be written in the form of the integral over trajectories $N_k(k, x, t) = N_k^0(k_0, x_0) = N_k^0\left(k_0(k, x, t), x_0(k, x, t), t\right)$, where $k_0(k, x, t)$ and $x_0(k, x, t)$ are inverse functions to $k = k(k_0, x_0, t)$ and $x = x(k_0, x_0, t)$ which are solutions of the ray equations (subscript x is omitted below)

To illustrate the instability we consider perturbations δT and δN_k. Then we have

$$\frac{\partial}{\partial t} \delta T + \frac{\partial}{\partial x} D_0 \sum \delta N_k \frac{\partial T_0}{\partial x} = 0. \qquad (19)$$

One can show that the effect of the variation of the temperature under constant D can be neglected for typical parameters.

The perturbation of the wave action can be found as $\delta N_k = N_k^0(k - \delta k) - N_k^0(k) = -\delta k \partial N_k^0/\delta k$ where is found from equations

$$\frac{d}{dt}\delta k = -\frac{\partial}{\partial x}\delta\omega = -\frac{\partial}{\partial x}\frac{\partial\omega}{\partial T}\delta T, \qquad (20)$$

$$\frac{d}{dt} = \frac{\partial}{\partial t} + \frac{\partial\omega}{\partial k}\frac{\partial}{\partial x}.$$

Then in the Fourier form we have

$$\delta k = \frac{q}{\Omega - qV_g}\frac{\partial\omega}{\partial T}\delta T. \qquad (21)$$

We obtain for the contribution of the individual wave vector to the growth rate

$$\Omega(\Omega - qV_g) = -q^2 D_0 \frac{\partial\omega}{\partial T} T_0^{'} \frac{\partial N_k^0}{\partial k_x}. \qquad (22)$$

Instability occurs for $\partial\omega/\partial T \; T_0^{'}\partial N_k^0/\partial k_x > 0$.

One can show that further evolution of the instability within Eqs (18) and (3) leads to the formation of singularities [22].

SUMMARY

It is shown here that large scale structures such as strongly sheared flow, magnetic islands and streamers, as well as localized transport events may develop as a result of secondary instabilities of the saturated drift wave turbulence. The instabilities develop due to a positive feedback on large scale formations from modulations of wave packets. We examined instability for the basic drift waves as well as for the electromagnetic fluctuations where generation of large scale magnetic field is possible. We have also shown that anomalous transport may be unstable leading to development of global plasma profile perturbations that may be relevant to large scale transport events, or avalanches.

Long wavelength structures, such as decsribed above, appear to be a crucial element of the self-organized drift wave turbulence due to a dominant role of nonlocal interactions [23]. The large scale modes are driven by the energy cascade from small scales, while the spectrum of small scale drift wave fluctuations is modified by the back reaction of large scales (such as random shearing [2]). It results in the coupled nonlinear system for large and small scale components. Fixed points of this system define saturated states of the drift wave turbulence and respective transport coefficients. Transitions between different fixed states may be related to sudden changes in transport (bursting phenomena) frequently observed in gyrokinetic simulations of tokamak transport [7].

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