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Zonal Flow and Stremaer Generation by Small – Scale Drift-Alfven Waves in Ionosphere Plasma

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Abstract

In the present work, the generation of large-scale zonal flows and streamers by modulationally unstable short-scale drift-Alfven waves in the ionosphere is investigated. Positive feedback in the system is achieved via modulation of the skin size drift-Alfven waves by the large-scale zonal flow. The conditions for the instability development and possibility of the generation of large-scale structures are determined. The instability pumps the energy of primarily small-scale Alfven waves into that of the large-scale zonal structures which is typical for an inverse turbulent cascade. Energy pumping into the large-scale region noticeably depends also on the width of the pumping wave spectrum and with an increase of the width of the initial wave spectrum the instability can be suppressed. It is assumed that the investigated mechanism can refer directly to the generation of mean flow in the atmosphere of the rotating planets and the magnetized plasma.

***Key words*:** Skin-size perturbations; Small and large scale turbulence; inverse cascade; Zonal flow; large scale magnetic field; pumping of energy with respect to scales. Reynolds and Maxwell stresses.

Introduction

Large scale coherent structures such as zonal flows and streamers can play important roles in transport behavior of magnetized plasmas. They present the integral parts of the collective activity of the majority of the planetary atmospheres and are manifested in the form of the large-scale low-frequency modes, propagating along the parallels (Busse, 1994; Aubert, et al., 2002). In toroidally enhanced plasmas, radially localized and poloidally elongated zonal flows can suppress (Lin et al, 1998; Li and Kashimoto, 2004), but poloidally localized and radially elongated streamers can enhance (Diamond et al, 2001; Yamada et al, 2010), the radial particle transport. There is evidence that these coherent structures are excited by drift and/or Alfv´en type of waves and the corresponding turbulence at time scales below the ion cyclotron frequency (Smolyakov et al, 2000; Kaladze et al, 2005). Earlier studies have focused mainly on waves with spatial scales larger than the ion gyroradius.

On the other hand, oscillations with spatial scales smaller than the ion gyroradius, such as the small-scale drift-Alfv´en waves (SSDAWs), can also efficiently drive coherent structures (Kaladze et al, 2007). In particular, it is found that pump SSDAWs propagating in the poloidal direction can excite zonal flows most efficiently (Kaladze et al, 2007), and the SSDAWs can be in either the electron or ion diamagnetic drift directions. However, up to now there is little investigation on the generation of streamers by SSDAWs.

The previous authors made the trials of investigations of the special features of the zonal flow generation by means of drift-Alfven type fluctuation on the basis of three sufficiently simplified models, describing nonlinear interaction between these modes: the first, a class of the models in which the effect of the ion temperature is negligible and only the effect of the so-called finite Larmor radius of ions according to the electron temperature (Guzdar, et al., 2001; Lakhin, 2003) is taken into account; the second model, where both disturbances, the primary small-scale as well as the large-scale zonal disturbances, have characteristic scale less than a Larmor radius of ions ρ_i (Smolyakov, et al., 2002); and the third class of the models, where finite Larmour radius of ions are considered neglecting the skin size inertial effects (Lakhin, 2004; Mikhailovskii, et al, 2006 b; Shukla, 2005; Kaladze et al, 2013). Although in the work (Pokhotelov, et al., 2003), generation of the zonal flow was studied by inertial Alfven fluctuations. But, it was made in uniform plasma neglecting finiteness of a Larmor radius of electrons, ions (T_e, $T_i \to 0$). One of the important wave modes in non-uniform magnetized space (Stasiewicz, et al., 2000; Sahraoui, et al., 2006; Narita, et al., 2007) as well as in laboratory (Gekelman, 1999) plasma media are electromagnetic small-scale drift-Alfven (SSDA) modes with the transverse wavelengths, small in comparison with a Larmor radius of ions (Aburjania et al, 2009) These small scale fluctuations can generate large-scale zonal modes and streamers in the space and as well as in the laboratory plasma. Moreover, the contemporary theory of anomalous transfers (Kadomtsev, Pogutse, 1984; Aburjania, 2006; Aburjania, 1990) predicts, that the anomalous thermal conductivity and diffusion in the plasma medium may be stipulated, in essence, by the processes with the characteristic wavelength $\lambda_\perp$ of the order of collision-less skin length λ_s, $\lambda_\perp = 2\pi / k_\perp \sim \lambda_s = c/\omega_{Pe}$, where $k_\perp$ is transversal (according to external equilibrium magnetic field) wave number of perturbations, $\omega_{Pe} = \left(4\pi e^2 n_0 / m_e\right)^{1/2}$ is a plasma frequency. In this connection, description of the nonlinear wave processes on the scales $\lambda_s \sim \lambda_\perp < \rho_i$ appears necessary.

In this paper we consider excitation of streamers from modulational instability of SSDAWs. A nonlinear equation describing the coupling of coherent structures and SSDAWs propagating in arbitrary directions is obtained.

1. Initial dynamic equations

The nonlinear equations describing the coupling of the SSDAWs with coherent structures in terms of the scalar potential ϕ and the parallel component A of the vector potential are (Aburjania et al, 2009)

$$\frac{\partial A}{\partial t} + V_{*e}\frac{\partial A}{\partial y} + c(1+\tau)\nabla_{||}\phi - \lambda_s^2 \frac{d}{dt}\Delta_\perp A = 0, \tag{1}$$

$$\frac{d}{dt}\phi + V_{*i}\frac{\partial \phi}{\partial y} - \frac{V_{Te}^2}{c\tau}\lambda_s^2 \nabla_{||}\Delta_\perp A = 0. \tag{2}$$

Here $V_{*e,i} = \mp c T_{e,i}\kappa_n/(eB_0)$ are electron and ion drift velocities respectively, where subscript "e" means electrons and "i" means ions; $V_{Te} = (T_e/m_e)^{1/2}$ – electrons' thermal velocity; $\tau = T_e/T_i$, $\nabla_{||} = \partial/\partial z - B_0^{-1}(\nabla A \times \nabla)_z$. Getting (1), (2) ion longitudinal motion is neglected and it is supposed that longitudinal current $J_{||}$ are caused by plasma electrons, $J_{||} = -c\Delta_\perp A / 4\pi$.

For threee-wave interaction involving a pump SSDAW and their satellites (secondary SS modes) we can write:

$$X = \tilde{X} + \hat{X} + \overline{X}, \tag{3}$$

where

$$\tilde{X} = \sum_k \left[\tilde{X}_+(k)\exp(i\mathbf{k}\cdot\mathbf{r} - i\omega_{\mathbf{k}}t) + \tilde{X}_-(k)\exp(-i\mathbf{k}\cdot\mathbf{r} + i\omega_{\mathbf{k}}t)\right], \tag{4}$$

describes a spectrum of SSDA initial pumping modes, $\mathbf{k} = (k_x, k_y, k_z)$, ω-wave vector and frequency of the initial modes, amplitude satisfies the condition $\tilde{X}_- = \tilde{X}_+^*$, where asterisk indicates a complex conjugation,

$$\hat{X} = \sum_k \left[\hat{X}_+(k)\exp(i\mathbf{k}_+\cdot\mathbf{r} - i\omega_+ t) + \hat{X}_-(k)\exp(i\mathbf{k}_-\cdot\mathbf{r} + i\omega_- t) + \text{c.c.}\right], \tag{5}$$

describes the small scale satellite (secondary) modes and

$$\overline{X} = \overline{X}_0\exp(-i\Omega t + i\mathbf{q}\cdot\mathbf{r}) + \text{c.c.}, \tag{6}$$

describes zonal flows. Laws of energy and impulse conservation is written in the next form: $\omega_\pm = \Omega \pm \omega_k$ and $\mathbf{k}_\pm = \boldsymbol{q}\cdot\mathbf{r} \pm \mathbf{k}$, respectively. Thus, the pairs $(\omega_{\mathbf{k}}, \mathbf{k})$ and $(\Omega, \mathbf{q}\cdot\mathbf{r})$ represent frequency and wave vector of SSDA pumping modes and zonal flows, respectively. Amplitude of the zonal modes $\overline{X}_0 \equiv (\overline{A}_0, \overline{\phi}_0)$ is considered to be constant in local approach. Since the coherent structures are perpendicular to the external magnetic field we have $q = q_\perp$. Further analyze will be carried out in the frames of the standard approximation $q_\perp / k_\perp << 1, \Omega/\omega << 1$.

1.1 Generation of the zonal flows by SSDAWs

If we consider a case when $q_y = 0$, e.i. $q_\perp = q_x$, following the standard quasi nonlinear procedure, substituting the expressions (3)-(6) in the equations (1), (2) and neglect small nonlinear terms with high frequency modes we obtain the coherent structures – the zonal flows:

$$(\omega_k - \omega_{*i})\tilde{\phi}_\pm(\mathbf{k}) - \frac{cT_i k_z k_\perp^2}{4\pi e^2 n_0}\tilde{A}_\pm(\mathbf{k}) = 0, \tag{7}$$

$$k_z c(1+\tau)\tilde{\phi}_\pm(\mathbf{k}) - \left[\omega_{\mathbf{k}}(1 + \mathbf{k}_\perp^2\lambda_{\mathbf{s}}^2)\right]\tilde{\mathbf{A}}_\pm(\mathbf{k}) = 0. \tag{8}$$

$$(\omega_\pm \mp \omega_{*i})\hat{\phi}_\pm \mp \frac{cT_i k_z k_{\perp\pm}^2}{4\pi e^2 n_0}\hat{A}_\pm = \mp\alpha_2^\pm \frac{cT_i k_z (k_\perp^2 - q_x^2)}{4\pi e^2 n_0[(1 + k_\perp^2\lambda_s^2)\omega_\pm - \omega_{*e}]}, \tag{9}$$

$$\mp k_z c(1+\tau)\hat{\phi}_\pm + [(1 + k_\perp^2\lambda_s^2)\omega_\pm \mp \omega_{*e}]\hat{A}_\pm = \mp\alpha_1^\pm\left[1 - \alpha_0\frac{\overline{\phi}_0}{\overline{A}_0}\frac{ck_z(1+\tau)}{(1 + k_\perp^2\lambda_s^2)\omega_k - \omega_{*e}}\right]. \tag{10}$$

$$i\Omega\overline{\phi}_0 = \frac{cT_i q_x^2}{4\pi e^2 n_0 B_0}\sum_{\mathbf{k}} k_y R_1(\mathbf{k}), \tag{11}$$

$$-i\Omega\left(1 + q_x^2\lambda_s^2\right)\overline{A}_0 = \frac{cq_x(1+\tau)}{B_0}\sum_{\mathbf{k}} k_y R_2(\mathbf{k}) + \frac{c}{B_0}q_x\lambda_s^2\sum_k k_y R_3(\mathbf{k}). \tag{12}$$

Here

$$\alpha_0 = \frac{1+\tau+k_\perp^2\lambda_s^2}{1+\tau+q_x^2\lambda_s^2}\ ,\ \ \alpha_2^{\pm} = \frac{ic}{B_0}k_y q_x(1+\tau)\overline{A}_0\tilde{\phi}_{\pm},\ \ \alpha_1^{\pm} = \frac{ic}{B_0}k_y q_x\left(1+\tau+q_x^2\lambda_s^2\right)\overline{A}_0\tilde{\phi}_{\pm}\ .$$

$$R_1(\mathbf{k}) = q_x\left(\tilde{A}_-\hat{A}_+ - \tilde{A}_+\hat{A}_-\right) + 2k_x\left(\tilde{A}_-\hat{A}_+ + \tilde{A}_+\hat{A}_-\right), R_2 = \tilde{\varphi}_-\hat{\delta}_+ - \tilde{\varphi}_+\hat{\delta}_-,$$

$$\hat{\delta}_{\pm} = \hat{A}_{\pm} - \frac{ck_z(1+\tau)}{\left(1+k_\perp^2\lambda_s^2\right)\omega_k - \omega_{*e}}\hat{\varphi}_{\pm},$$

$$R_3 = k_\perp^2\left(\tilde{A}_+\hat{\phi}_- + \tilde{\phi}_-\hat{A}_+ - \tilde{A}_-\hat{\phi}_+ - \tilde{\phi}_+\hat{A}_-\right) + q_x^2\left(\tilde{\phi}_-\hat{A}_+ - \tilde{\phi}_+\hat{A}_-\right) + 2q_x k_x\left(\tilde{\phi}_-\hat{A}_+ + \tilde{\phi}_+\hat{A}_-\right).$$

1.2 Dispersion equation of the large scale zonal flows and magnetic fields

Using the above expression general dispersion relation can be obtained:

$$\overline{\phi}_0 = L_1^{\phi}\overline{\phi}_0 + L_1^{A}\overline{A}_0\ , \tag{13}$$

$$\overline{A}_0 = L_2^{\phi}\overline{\phi}_0 + L_2^{A}\overline{A}_0\ . \tag{14}$$

where

$$\left(L_1^{\phi},\ \ L_1^{a},\ L_2^{\phi},\ L_2^{A}\right) = \sum_k \frac{\left(l_1^{\phi},\ l_1^{A},\ l_2^{\phi},\ l_2^{A}\right)}{\left(\Omega - q_x V_g\right)^2}\ . \tag{15}$$

Here $\mathbf{V_g} = \mathbf{V_g}(\mathbf{k})$ - zonal group velocity, defined by equality bellow:

$$V_g = \frac{2k_x}{k_\perp^2}\frac{(\omega_k - \omega_{*i})\left[\left(1+k_\perp^2\lambda_s^2\right)\omega_k - \omega_{*e}\right]}{2\left(1+k_\perp^2\lambda_s^2\right)\omega_k - \left(1+k_\perp^2\lambda_s^2\right)\omega_{*i} - \omega_{*e}}, \tag{16}$$

And the functions $\left(l_1^{\phi},\ l_1^{A},\ l_2^{\phi},\ l_2^{A}\right)$ denote

$$l_1^{\phi} = (1+\tau)\frac{q_x k_x}{k_\perp^2}\frac{(\omega_k - \omega_{*i})\Gamma_0^2}{\left[\left(1+k_\perp^2\lambda_s^2\right)\omega_k - \omega_{*e}\right]\cdot\left[2\left(1+k_\perp^2\lambda_s^2\right)\omega_k - \left(1+k_\perp^2\lambda_s^2\right)\omega_{*i} - \omega_{*e}\right]^2}R_1^{\phi}, \tag{17}$$

$$l_1^{A} = \frac{1}{c}\frac{q_x k_x}{k_z k_\perp^2}\frac{(\omega_k - \omega_{*i})\Gamma_0^2}{\left[2\left(1+k_\perp^2\lambda_s^2\right)\omega_k - \left(1+k_\perp^2\lambda_s^2\right)\omega_{*i} - \omega_{*e}\right]^2}R_1^{A} \tag{18}$$

$$l_2^{\phi} = c(1+\tau)\frac{k_z\Gamma_0^2\left[(1+\tau)R_2^{\phi} + \frac{q_x\lambda_s^2}{\Omega}\left(1+\tau+k_\perp^2\lambda_s^2\right)R_3^{\phi}\right]}{\left[\left(1+k_\perp^2\lambda_s^2\right)\omega_k - \omega_{*e}\right]\cdot\left[2\left(1+k_\perp^2\lambda_s^2\right)\omega_k - \left(1+k_\perp^2\lambda_s^2\right)\omega_{*i} - \omega_{*e}\right]^2}, \tag{19}$$

$$l_2^{A} = -(1+\tau)\frac{\Gamma_0^2\left(R_2^{A} + \frac{q_x\lambda_s^2}{\Omega}R_3^{A}\right)}{\left[2\left(1+k_\perp^2\lambda_s^2\right)\omega_k - \left(1+k_\perp^2\lambda_s^2\right)\omega_{*i} - \omega_{*e}\right]^2}, \tag{20}$$

where

$$\Gamma_0^2 = \frac{c^2 q_x^2 k_y^2}{B_0^2}I_k\ . \tag{21}$$

From the closed system of equations (13) and (14), we simply get the dispersion equation for large scale zonal flows and the magnetic fields:

$$1 - \left(L_1^{\phi} + L_2^{A}\right) + L_1^{\phi}L_2^{A} - L_2^{\phi}L_1^{A} = 0, \tag{22}$$

The dispersion relation of the zonal modes (22) allows an investigation of their generation via continuous spectrum of the initial modes with skin scale, which is the main subject of the traditional theory of such generation, which uses a kinetic equation for the waves, summarized in (Diamond, et al., 2005). Thus, the approach developed in section 3.4, based on dynamic equations of magnetic hydrodynamics of the ionosphere, is an alternative to the approach in (Diamond, et al., 2005) and in our opinion, is more convenient in to realize, also in the interpretation of results obtained based on them. It's obvious that the dispersion relation (22) represents bisquared equation according to $\Omega - q_x V_g$. However, as it will be shown bellow, this equation can be reduced to a squared one for a very interesting range of frequencies Ω of the zonal perturbation.

1.3 Generation of the streamers by SSDAWs

Let us consider the case of coherent structures $q_\perp \neq 0$. We define an angle between the direction of pump wave SSDAW $k_\perp$ and **x** radial direction θ_p, angle between $k_\perp$ and direction of the coherent structure $q_\perp$ - θ_c. Therefore, zonal flows and streamers correspond to $\theta_c = 0^\circ$ and $\theta_c = 90^\circ$, respectively. For $\theta_c \neq 0^\circ$ and $\theta_c \neq 90^\circ$, it corresponds to oblique coherent structure.

$$\left(\Omega - \omega_{*i}^{q}\right)\overline{\Phi}_0 = -\frac{iV_A^2\rho_i^2}{cB_0}[\left(k_x k_{+y} - k_{+x}k_y\right)\left(k_\perp^2 - k_{+\perp}^2\right)A_k^* A_{k+} - \left(k_x k_{-y} - k_{-x}k_y\right)\left(k_\perp^2 - k_{-\perp}^2\right)A_k^* A_{k-}]; \quad (23)$$

$$\left(\Omega - \omega_{*i}^{q}\right)\overline{A}_0 = -\frac{ic}{B_0}\left(1+\frac{1}{\tau}\right)[\left(k_x k_{+y} - k_{+x}k_y\right)\left(A_k^*\Phi_{k+} - A_{k+}\Phi_k^*\right) - \left(k_x k_{-y} - k_{-x}k_y\right)\left(A_k^*\Phi_{k-} - A_{k-}\Phi_k^*\right)]; \quad (24)$$

Where

$$\omega_{*i,e} = V_{*i,e}k_y, \quad \omega_{*i,e}^{q} = V_{*i,e}q_y, \quad \omega_{*i,e}^{\pm} = V_{*i,e}k_{\pm y},$$

$$\alpha = -\frac{ic}{B_0}\left(1+\frac{1}{\tau}\right)\left(k_y q_x - k_x q_y\right)\overline{\Phi}_k,$$

$$D_\pm = \left(\omega_\pm - \omega_{*i}^{\pm}\right)\left(\omega_\pm - \omega_{*e}^{\pm}\right) - \frac{k_{\pm\perp}^2}{k_\perp^2}\left(\omega_k - \omega_{*i}\right)\left(\omega_k - \omega_{*e}\right).$$

We can obtain the coupling equations for the coherent structures:

$$\begin{pmatrix} \frac{1}{\omega_{ci}^2}\left(\Omega - \omega_{*i}\right)D_+D_- - a_1 I_{11} & a_2 I_{12} \\ b_1 I_{21} & \frac{1}{\omega_{ci}^2}\left(\Omega - \omega_{*e}\right)D_+D_- - b_2 I_{22} \end{pmatrix}\begin{pmatrix} \overline{\Phi}_0 \\ \overline{A}_0 \end{pmatrix} = 0, \quad (25)$$

Eq. (25) gives the general dispersion relation of the coherent structures

$$\left[\frac{1}{\omega_{ci}^2}\left(\Omega - \omega_{*i}\right)D_+D_- - a_1 I_{11}\right]\left[\frac{1}{\omega_{ci}^2}\left(\Omega - \omega_{*e}\right)D_+D_- - b_2 I_{22}\right] + a_2 b_1 I_{12} I_{21} = 0. \quad (26)$$

If the nonlinear term $|B_\perp/B_0|^2$ is small and can be neglected, Eq. (26) is reduced to

$$D_+^2 D_-^2 \left(\Omega - \omega_{*i}^0\right)\left(\Omega - \omega_{*e}^0\right) = 0, \tag{27}$$

whose solutions are

$$\begin{aligned}
\Omega_{1\pm} &= \frac{1}{2}\left(-b_+ \pm \sqrt{b_+^2 - 4c_+}\right), \\
\Omega_{2\pm} &= \frac{1}{2}\left(-b_- \pm \sqrt{b_-^2 - 4c_-}\right), \\
\Omega_3 &= \omega_{*i}^0, \\
\Omega_4 &= \omega_{*e}^0,
\end{aligned} \tag{28}$$

Where

$$\begin{aligned}
b_\pm &= \pm\left(2\omega_k - \omega_{*i} - \omega_{*e}\right) - \omega_{*i}^0 - \omega_{*e}^0, \\
c_\pm &= \pm\omega_{*i}^0\left(\omega_k - \omega_{*e}\right) \pm \omega_{*e}^0\left(\omega_k - \omega_{*e}\right) + \omega_{*i}^0\omega_{*e}^0 \\
&\quad - \frac{1}{k_\perp^2}\left(q_\perp^2 \pm 2k_\perp \cdot q_\perp\right)\left(\omega_k - \omega_{*i}\right)\left(\omega_k - \omega_{*e}\right),
\end{aligned}$$

where $\Omega_{1\pm}$ and $\Omega_{2\pm}$ are the solutions of $D_+ = 0$ and $D_- = 0$, respectively.

Discussion of the results and conclusion

In this work, the features of the generation of large scale zonal flows and the streamers due to small scale drift Alfven (SSDA) turbulence in the ionospheric plasma medium are investigated. Two self-consistent interconnected nonlinear equations (5) and (6), determining electrostatic and vector potentials describing the dynamics of the wave structures with finite ion Larmour radius. They are valid for the structures till the skin scale ($k_\perp^2 c^2 / \omega_{pe}^2 \sim 1$). Analysis of these structures was carried out in the frames of nonlinear parametric formalism, analogous to the theory of convective cells generation due to monochromatic pumping waves, but generalized in the sense that instead of separate monochromatic packet we investigated the initial waves arbitrary according to spectrum width. Such modification of the parametric approach reveals a new feature of the interaction of the small scale and the large scale modes . As in the majority of preceding works, we suppose the existence of some linear or nonlinear mechanisms of initial mode excitation responsible for growth of their amplitudes above the fluctuation level. Due to competition between linear and nonlinear effects, some stationary state of the initial waves will form. Herewith, we consider that zonal mode generation treated by us takes place at the end of the corresponding initial modes' stationary level formation. These initial modes have real frequencies.

It must be mentioned that zonal mode generation by SSDA pumping waves is possible (see eq. (26), only at drift effects ($\omega_{*e,i} \neq 0$), but the initial nonlinear equations of these modes (5) and (6) are valid also neglecting these effects, i.e., for uniform plasma.

It is established, that SSDA of the skin size as well as comparably long wavelength one, effectively generate the large scale zonal flows and the streamers. These modes will be excited due to joint nonlinear action of the Reynolds and Maxwell stresses. Physical reason of the zonal flow generation due to electron drift modes represents Reynolds stress. As it comes to quasi electromagnetic ion-drift waves, they as well as kinetic Alfven pumping in uniform plasma, can not excite the large scale zonal modes as a result of whole compensation of moderate Reynolds stresses with Maxwell's ones.

We suppose, the mechanism discussed in this work is applicable for theoretical foundation experimentally observed mean flow generation in the atmospheres of the rotating planets and in magnetized plasma. On the one side, it can play a definite role in boundary layer Alfven turbulence formation (Pokhotelov, et al., 2003). On the other side, parametric instability can cause shear flow generation in laboratory plasma where it can sufficiently impact the drift plasma turbulence and suppress the transfer processes (Smolyakov, et al., 2000; Manfredi, et al., 2001). Thus, this instability can be one of the main nonlinear saturation mechanisms of amplitudes of the wave perturbations in space and laboratory plasma.

This study developed a general nonlinear dispersion equation to describe the generation of coherent structures by SSDA modes. For the meso-scale streamer structure (θ_{coh} = 90°), it is found that its excitation comes from either the nonlinear coupling between two branches $\mathbf{\Omega}_{1+}$ and $\mathbf{\Omega}_{2-}$ for the ion diamagnetic-drift pump wave, or two branches $\mathbf{\Omega}_{1-}$ and $\mathbf{\Omega}_{2+}$ for the electron diamagnetic-drift pump wave. The excitation is most effective when the pump SSDA mode propagates at $\theta_{pump} \sim 30$

The study compared the streamer excitation to the zonal flow excitation. It is shown that like the streamer case both ion and electron diamagnetic-drift pump modes can drive zonal flows. However, the strongest excitation of zonal flows happens at the angle different from that for streamer excitation. Therefore, it proposed that the generation of coherent structures depends strongly on θ_{pump}.

The theory which is developed in this work is applicable for nonlinear dynamics of the ionospheric Alfven resonator (IAR). Ground-based observations of mid-latitudes (Belyaev, et al., 1990; Bosinger, et al., 2002), high-latitudes (Belyaev et al., 1999; Demekhov, et al., 2000) and satellite observations (Grzesiak, 2000; Chaston, et al., 2002) convincingly verify the existence of IAR in the upper ionosphere. Results of ionospheric measurements indicate also (Stasiewicz, et al., 2000), that Alfven waves excited in the IAR actually are not seen as small amplitude linear waves. They always have comparably large amplitudes and represent eigen modes of IAR in a strongly nonlinear state. Let’s estimate effectiveness of the mechanism considered for IAR regions (Pokhotelov, et al., 2003). For characteristic parameters of IAR (Pokhotelov, et al., 2003): $E \sim 10^{-2}$v/m, $B_0 \sim 3\times10^{-5}$T, $\lambda_s \sim 100$ m, $k_\perp\lambda_s \sim 1$, $q_x \sim 0.01k_x$, we get $\gamma_{opt} \sim 10^5\ s^{-1}$. Thus, the considered parametric instability pumps energy of SSDA pumping waves into the energies of the large scale zonal flow or streamers in $\sim 10^{-5}$ seconds. A more detailed qualitative comparison of the theoretical results obtained through the observed and experimental data is outside the scope of this work and is a subject for a separate publication.

Thus, our analysis shows that the parametric instability in the ionosphere is developed simply, and can become a sufficient nonlinear mechanism of energy pumping from small-scale drift Alfven turbulence into large-scale (or meso-scale) zonal flows and the streamers. Such energy distribution leads to a small-scale turbulence level decreasing and to noticeable weakening of anomalous transfer processes in the medium. So, parametrically excited meso-scale flows can present sufficient parts of elements of the structural plasma turbulence in the ionosphere and lower magnetosphere of the Earth.

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ზონალური დინებებისა და სტრიმერების გენერაცია მცირე-მასშტაბიანი დრეიფული ალფენის ტალღებით იონოსფერულ პლაზმაში

ხ. ჩარგაზია

რეზიუმე

ნაშრომში შესწავლილია დიდმასშტაბიანი ზონალური დინებებისა და სტრიმერების გენერაცია მოდულაციურად არამდგრადი მცირე მასშტაბიანი დრეიფული ალფენიოს ტალღებით იონოსფეროში.გამოვლენილია დადებითი უკუკავშირი სკინ-სისქის დრეიფულ ალფენის ტალღასა და დიდმასშტაბიან ზონალურ დინებებს შორის. განსაზღვრულია არამდგრადობის განვითარებისა და დიდმაშტაბიანი სტრუქტურების გენერაციის პირობა. არამდგრადობა განაპირობებს ენერგიის გადაქაჩვას მნიშვნელოვნად მცირე მასშტაბიანი ალფენის ტალღიდან დიდმასშტაბიან ზონალურ დინებებში, რაც დამახასიათებელია ტურბულენტური უკუკასკადისათვის. ენერგიის გადატანა შედარებით დიდმასშტაბიან არეში დამოკიდებულია საწყისი ტალღის სპექტრის სიგანეზე, რომლის ზრდასთან ერთად, არამდგრადობა ითრგუნება. აღნიშნული ენერგიის გადატანის მექანიზმი შეიძლება გამოყენებულ იქნას ფონური დინების გენერაციის შემთხვევისათვის მბრუნავი პლანეტების ატმოსფეროებში და ასევე მაგნიტური ველებისათვის.

Генерация зональных течений и стримеров мелко масштабными дрейфого Альвеновскими волнами в ионосферной плазме

Х. Чаргазия

Резюме

В работе изучена генерация крупномасштабных зональных течений и стримеров с модуляционной неустойчивостью мелкомасштабных дрейфовых Альвеоновских волн в ионосферной плазме. Выявлена позитивная обратная связь между дрейфовыми Альвеоновскими волнами скиновой толщины и зональными течениями. Определены условия развития неустойчивости и генерации крупномасштабных структур. Неустойчивость определяет перекачку энергии с мелкомасштабных дрейфовых Альвеоновских волн в зональных течений, что свойственно турбулентному обратному каскаду. Перекачка энергии в относительно крупномасштабной области зависит от ширины спектра начальных волн, с ростом каторого неустойчивость убывает. Рассмотренный механизм перекачки энергии может быть использован для случая генерации фонового течения в атмосферах вращающихся планет, а также для магнитных полей.