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Dynamics of large-scale vortical structures in electron-positron-ion plasmas

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To describe the nonlinear propagation of electrostatic drift waves the generalized Hasegawa–Mima equation containing both vector (Jacobian) and scalar (Korteweg–de Vries-type) nonlinearities is obtained for electron-positron-ion plasmas. The drift waves are supposed to have arbitrary wavelengths (as compared with the Larmour radius of plasma ions at the plasma electron temperature). Temperature inhomogeneity of electrons and positrons is taken into account. Spatial increase in the linear plasma-potential perturbations in the direction of density and temperature inhomogeneities is shown. Self-organization mechanism of large-scale drift solitary vortices is considered. It is shown that the existence of positrons in plasma enriches the class of solutions of the generalized Hasegawa–Mima equation. © *2009 American Institute of Physics*.
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Due to the long lifetime of positrons, most astrophysical[1–3] and laboratory plasmas[4–6] become an admixture of electrons, positrons, and ions. Three-component electron-positron-ion (EPI) plasma can also be produced by electrons accelerated to relativistic speeds either in the fields of high-power laser beams[7] or in the large amplitude wake fields that are generated by ultrashort laser pulses in plasmas.[8,9] The three-component EPI plasmas have also been created in laboratories[10,11] by injecting positrons in EI systems. Therefore, the study of EPI plasmas is important to understand the behavior of both astrophysical and laboratory plasmas. To grasp the basic physics of the three-component EPI plasmas several theoretical investigations have been made[12–21] within the framework of multifluid theory which are applied to both astrophysical and laboratory plasmas.

The dynamics of low-frequency waves studied in usual EI plasmas is generally modified in EPI plasmas. These modifications depend on the concentration ratios of different species as well as on the magnitudes of their temperatures. The study of low-frequency long wavelength drift waves is of great interest because of its applications to many laboratory, space, and astrophysical systems. Such interest in drift waves is also connected with the consequence that their existence explains anomalous transport of a plasma transverse to a magnetic field. Zonal flows, generated by drift modes, also play a major role in controlling the level of anomalous transport in magnetic confinement systems. Different aspects of electrostatic drift waves dynamics have been considered in EPI plasma,[22–26] where electron and positron temperatures have been assumed to be equal and constant in space. Besides, the modified Hasegawa–Mima (HM) equation obtained is not correct. In Refs. 22 and 23 solitary vortical structures are also found as a solution of the corresponding nonlinear HM equation. In these papers it is assumed that the characteristic size of the considered nonlinear structures is less or of the order of the ion Larmour radius at the plasma electron temperature. Such structures may be described in the framework of classical HM equation containing only vectorial (Jacobian) nonlinearity. In other words classical HM equation only describes small-scale structures. Vortices in a strongly magnetized EPI plasma in the electrostatic limit were investigated by Shukla *et al.*[27] via the HM equation with a vector nonlinearity and vortex solutions of two different classes, viz., vortex chain and a double vortex were shown to exist. Extension to the large-scale electrostatic drift nonlinear structures (having dimensions larger than the characteristic Larmour radius of plasma ions) was discussed by Nezlin and Chernikov.[28] It was shown that such structures are described in the framework of the generalized HM equation containing in addition to the vectorial nonlinearity a new scalar nonlinearity of the Korteweg–de Vries (KdV) type. According to the new self-organization mechanism, solitary structures are formed by mutual compensation of wave dispersion by both scalar and vector nonlinearities. As a result, in the general case, a solitary structure becomes essentially anisotropic and is a superposition of monopolar and dipolar vortices.

In this report the generalized HM equation valid for arbitrary sizes of vortical structures is obtained and the nonlinear dynamics of large-scale solitary vortical structures on the electrostatic drift waves propagating in EPI plasma is discussed. In addition temperatures of electrons and positrons are assumed to be arbitrary.

Let us consider, in the electrostatic approximation, the quasi-two-dimensional motion of a quasineutral EPI plasma. We consider a local perturbation (with respect to the unperturbed plasma environment) of the plasma potential $\varphi(t,x,y)$ and assume that the external magnetic field $\vec{B}_0$ is taken in the $\hat{z}$ direction.

The unperturbed plasma densities of electrons and positrons $n_{eo}(x)$, $n_{po}(x)$ and corresponding temperatures $T_{eo}(x)$, $T_{po}(x)$ are inhomogeneous and assumed to decrease monotonously along the x axis. The ions are considered "cold" and the quasineutrality condition in equilibrium

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$$n_{eo}(x) = Zn_{io}(x) + n_{po}(x) \tag{1}$$

is fulfilled, where Z is the charge number of positive ions.

Let us assume that in this system the plasma density perturbation arises (corresponding to the plasma-potential perturbation φ), which excites a drift wave. Assume that the plasma motion in the (x,y) plane is sufficiently slow, so that electrons and positrons (fast moving along the magnetic field) follow the Boltzmann equilibrium. Then from the plasma quasineutrality condition (1), the ion density is defined by the relationship

$$Zn_i = n_{eo}(x)e^{e\varphi/T_e(x)} - n_{po}(x)e^{-e\varphi/T_p(x)}. \tag{2}$$

The equation of motion for the plasma-ion component under the action of the crossed electric (caused by the perturbed potential φ) and magnetic, $\vec{B}_0$, fields has the form

$$\vec{v}_t + (\vec{v}\cdot\nabla)\vec{v} = \vec{v}\times\vec{\omega}_{Bi} - \frac{Ze}{M}\nabla\varphi, \tag{3}$$

where $\vec{\omega}_{Bi}=Ze\vec{B}_o/M$ is the cyclotron frequency of ions, Ze and M are the ion charge and mass, respectively, and $\vec{v}(u,v,0)$ is the velocity vector which has the velocity components u and v (along the x and y axes, respectively). Note that the $\vec{\omega}_{Bi}$ vector is directed along the magnetic field, that is, in the direction opposite to that of the angular velocity vector of the ion Larmour rotation.

We consider a quasi-two-dimensional motion, and therefore the velocity perturbation along the z axis is absent. Using the continuity equation for ions with Eq. (3), we may obtain the equation of conservation of the so-called potential vorticity, which is defined as

$$\frac{d}{dt}\left(\frac{\varsigma + \omega_{Bi}}{n}\right) = 0, \tag{4}$$

where $d/dt=\partial_t+u\partial_x+v\partial_y$ and the vorticity $(\vec{\varsigma}=\nabla\times\vec{v})_z=v_x - u_y$. Here the lower indices refer to differentiation with respect to x, y, and t, the cyclotron frequency $\omega_{Bi}=\vec{e}_z\cdot\vec{\omega}_{Bi}$ and $\vec{e}_z$ is the unit-vector in the z direction.

The drift wave regime in plasma takes place when the inertia terms in the Euler equation (4) are small compared to the Lorentz force, due to which the equations of motion, in the first approximation, are the equations of equilibrium between the magnetic force and the force caused by the potential gradient; such an approximation implies the existence of a small parameter,

$$\frac{\omega}{\omega_{Bi}} \ll 1, \tag{5}$$

where ω is the characteristic frequency of the perturbation. We also take into account the polarization drift (which is a higher order term) and in accordance with Eq. (5) we get the following expression for the ion velocity:

$$\vec{v}_\perp = -\frac{Ze}{M\omega_{Bi}}\nabla\varphi\times\vec{e}_Z - \frac{Ze}{M\omega_{Bi}^2}\left(\frac{\partial}{\partial t} + \frac{Ze}{M\omega_{Bi}}\vec{e}_Z\times\nabla_\perp\varphi\cdot\nabla_\perp\right)\nabla_\perp\varphi. \tag{6}$$

Here the subscript $\perp$ denotes the plane perpendicular to the external magnetic field.

Substituting Eqs. (2) and (6) into Eq. (4) we obtain the following generalized (containing both vector and scalar nonlinearities) HM equation for an EPI plasma:

$$\begin{aligned}
&-r_L^2\left(1-\frac{n_{po}}{n_{eo}}\right)\frac{\partial\Delta_\perp\varphi}{\partial t} + \frac{1}{Z}\left(1+\frac{n_{po}}{n_{eo}}\frac{T_e}{T_p}\right)\frac{\partial\varphi}{\partial t} - r_L^2\omega_{Bi}\frac{n'_{eo}-n'_{po}}{n_{eo}}\frac{\partial\varphi}{\partial y} - r_L^2\frac{n'_{eo}-n'_{po}}{n_{eo}}\frac{\partial^2\varphi}{\partial t\,\partial x} - r_L^2\frac{e}{T_e}\left(1+\frac{n_{po}}{n_{eo}}\frac{T_e}{T_p}\right)\varphi\frac{\partial\Delta_\perp\varphi}{\partial t} - r_L^4\frac{Ze\omega_{Bi}}{T_e}\\
&\times\left(1-\frac{n_{po}}{n_{eo}}\right)J(\varphi,\Delta_\perp\varphi) + r_L^2\frac{e}{T_e}\left(1+\frac{n_{po}}{n_{eo}}\frac{T_e}{T_p}\right)\Delta_\perp\varphi\frac{\partial\varphi}{\partial t} + \frac{e}{2Z}\left(\frac{1}{T_e}-\frac{n_{po}}{n_{eo}}\frac{T_e}{T_p^2}\right)\frac{\partial\varphi^2}{\partial t} - r_L^4\frac{Ze\omega_{Bi}}{T_e}\frac{n'_{eo}-n'_{po}}{n_{eo}}\Delta_\perp\varphi\frac{\partial\varphi}{\partial y} - r_L^2\frac{e}{2T_e}\\
&\times\left(1+\frac{n_{po}}{n_{eo}}\frac{T_e}{T_p}\right)\frac{\partial(\nabla_\perp\varphi)^2}{\partial t} - r_L^4\frac{Ze\omega_{Bi}}{T_e}\frac{n'_{eo}-n'_{po}}{n_{eo}}J\left(\varphi,\frac{\partial\varphi}{\partial x}\right) - \frac{e}{2M\omega_{Bi}}\left(\frac{n'_{eo}}{n_{eo}}+\frac{n'_{po}}{n_{eo}}\frac{T_e}{T_p}-\frac{T'_e}{T_e}-\frac{n_{po}}{n_{eo}}\frac{T_eT'_p}{T_p^2}\right)\frac{\partial\varphi^2}{\partial y} = 0.
\end{aligned} \tag{7}$$

In obtaining the above equation the inequality (5) and condition $e\varphi/T_e \ll 1$ were used to retain the dominant nonlinear terms. Thus Eq. (7) is valid for arbitrary sizes of vortical structures. In Eq. (7) $J(a,b)=\partial a/\partial x\,\partial b/\partial y - \partial a/\partial y\,\partial b/\partial x$ is the Jacobian (vector nonlinearity), the operator $\Delta_\perp=\nabla_\perp^2=\partial^2/\partial_x^2+\partial^2/\partial_y^2$, $r_L=(T_e/M\omega_{Bi}^2)^{1/2}$ is the ion Larmour radius at electron temperature T_e, and the prime denotes derivative with respect to the x variable.

In the linear regime Eq. (7) reduces to

$$\frac{Z^2}{M\omega_{Bi}}\frac{n_{io}}{n_{eo}}\frac{\partial\Delta_\perp\varphi}{\partial t} - \frac{\omega_{Bi}}{T_e}\left(1+\frac{n_{po}}{n_{eo}}\frac{T_e}{T_p}\right)\frac{\partial\varphi}{\partial t} + \frac{Z^2}{M}\frac{n'_{io}}{n_{eo}}\frac{\partial\varphi}{\partial y} + \frac{Z^2}{M\omega_{Bi}}\frac{n'_{io}}{n_{eo}}\frac{\partial^2\varphi}{\partial t\,\partial x} = 0. \tag{8}$$

The last term (which is of the order of ω/ω_{Bi}) is kept in order to obtain the spatial structure of the electrostatic drift waves. Indeed, if we introduce the inhomogeneity length $1/L = |n'_{io}/n_{io}|$ and look for the propagation of drift plane waves of the form

$$\varphi \sim e^{\mp x/2L} e^{ik_x x + ik_y y - i\omega_o t}, \tag{9}$$

we get for the electrostatic drift frequency

$$\omega_o = \frac{\mp \dfrac{Z^2 T_e}{M\omega_{Bi}} \dfrac{k_y}{n_{eo}} \dfrac{1}{L}}{\dfrac{1}{n_{io}}\left(1 + \dfrac{n_{po}}{n_{eo}} \dfrac{T_e}{T_p}\right) + Z^2 \dfrac{r_L^2}{n_{eo}}\left(k_\perp^2 + \dfrac{1}{4L^2}\right)}. \tag{10}$$

Here $k_\perp^2 = k_x^2 + k_y^2$ and the $\mp$ sign in Eqs. (9) and (10) corresponds to positive and negative signs of n'_{io}/n_{io}. Thus the drift wave potential is spatially unstable in the direction of the inhomogeneity. The drift wave structure given by Eqs. (9) and (10) is similar to that of acoustic-gravity waves propagating in an inhomogeneous atmosphere embedded in a gravitational field.[29]

It is seen from Eq. (7) that the classical HM equation (containing only vector nonlinearity) can be isolated, which has the form

$$\begin{aligned} &-\frac{Ze}{M\omega_{Bi}}\left(1 - \frac{n_{po}}{n_{eo}}\right)\frac{\partial \Delta_\perp \varphi}{\partial t} + \frac{e\omega_{Bi}}{T_e}\left(1 + \frac{n_{po}}{n_{eo}}\frac{T_e}{T_p}\right)\frac{\partial \varphi}{\partial t} \\ &\quad - \frac{Ze}{M}\frac{n'_{eo} - n'_{po}}{n_{eo}}\frac{\partial \varphi}{\partial y} - \frac{Ze}{M\omega_{Bi}}\frac{n'_{eo} - n'_{po}}{n_{eo}}\frac{\partial^2 \varphi}{\partial t\, \partial x} \\ &\quad - \frac{Z^2 e^2}{M^2 \omega_{Bi}^2}\left(1 - \frac{n_{po}}{n_{eo}}\right) J(\varphi, \Delta_\perp \varphi) = 0. \end{aligned} \tag{11}$$

The HM equation obtained above is valid if the following inequalities are fulfilled:

$$\frac{a^2}{r_L^2} \ll \frac{\omega_{Bi}}{\omega}, \quad \frac{L}{a}; \quad \frac{a^4}{r_L^4} \ll \frac{\omega_{Bi}}{\omega}; \quad \frac{L}{a} \gg 1. \tag{12}$$

Here a is the perpendicular size of the structure and L is the characteristic scale of the inhomogeneity. Thus, the classical HM equation is valid only when the ratio of the characteristic size of the considered nonlinear structures and the ion Larmour radius (at the plasma electron temperature) satisfies the inequalities (12) and the smallness of this ratio as given in Ref. 28 need not to be fulfilled. Thus the classical HM equation, describes only "small-scale" structures. The classical HM equation predicts the existence of a solitary structure in the form of a dipolar vortex that is a cyclonic-anticyclonic pair. Solitary monopolar vortices (cyclones, anticyclones) are absent in the framework of the classical HM equation.

We now look for large-scale vortical structures,

$$\frac{a}{r_L} \gg 1. \tag{13}$$

In this limit Eq. (7) reduces to the following generalized HM equation for the potential perturbation φ:

$$\begin{aligned} &\left(\frac{n_{eo}}{n_{io}} + \frac{n_{po}}{n_{io}}\frac{T_e}{T_p}\right)\frac{\partial \varphi}{\partial t} - Z^2 r_L^2 \frac{\partial \nabla_\perp^2 \varphi}{\partial t} - Z^2 r_L^2 \omega_{Bi} \frac{n'_{io}}{n_{io}}\frac{\partial \varphi}{\partial y} \\ &\quad - Z^2 r_L^2 \frac{n'_{io}}{n_{io}}\frac{\partial^2 \varphi}{\partial t\, \partial x} + \frac{e}{2n_{io}}\left(\frac{n_{eo}}{T_e} - \frac{n_{po} T_e}{T_p^2}\right)\frac{\partial \varphi^2}{\partial t} \\ &\quad - \frac{Z^3 e \omega_{Bi} r_L^4}{T_e} J(\varphi, \Delta_\perp \varphi) - \frac{Ze\omega_{Bi} r_L^2}{2LT_e}\frac{\partial \varphi^2}{\partial y} = 0. \end{aligned} \tag{14}$$

Here we have introduced the inhomogeneity length L defined as

$$\begin{aligned} \frac{1}{L} &= \frac{1}{n_{io}}\left(n'_{eo} + n'_{po}\frac{T_e}{T_p} - n_{eo}\frac{T'_e}{T_e} - n_{po}\frac{T_e T'_p}{T_p^2}\right) \\ &= \frac{T_e}{n_{io}}\left(\frac{n_{eo}}{T_e} + \frac{n_{po}}{T_p}\right)'. \end{aligned} \tag{15}$$

In the generalized HM equation (14) we have kept the (so-called scalar) nonlinear term $\propto \partial \varphi^2/\partial y$ which is analogous to the KdV nonlinearity. The scalar and vector nonlinearities are of the same order when

$$\frac{a^2}{r_L^2} \sim \frac{L}{a} \gg 1. \tag{16}$$

Analogous to Eq. (14) a more simplified equation obtained in Ref. 22 with the expression for $1/L$ analogous to Eq. (15) for equal electron and positron temperatures is not correct.

Thus in the process of self-organization of large-scale vortical solitary structures, both nonlinearities, scalar and vector, taking place are important. Large-scale structures having dimensions larger than the characteristic Larmour radius of plasma ions are influenced by the scalar nonlinearity. It is very essential to elucidate the mechanism of formation of solitary drift vortical structures due to the competition of dispersion and nonlinearity.[28] The problem is as follows: In the solitary solitons of the KdV type, the dispersive spreading of a wave packet is balanced by its nonlinear steepening. This equilibrium, in particular, causes quite a definite relationship between the packet width and amplitude: The larger the amplitude (stronger nonlinearity), the smaller the width (stronger dispersion). However, the situation is more complicated when it was discovered in the model experiments[30] that the characteristic size of the structures under consideration was not dependent on the amplitude. This evidenced a more complicated dynamics of the structure formation compared to the simple compensation of dispersion by nonlinearity. When a nonlinear dynamic equation [e.g., obtained here in Eq. (14)] involves two types of nonlinearities (scalar and vectorial) the new type of self-organization of solitary structures was revealed.[28] In particular, if at a given structure size the structure amplitude is too large, i.e., the scalar nonlinearity exceeds dispersion, then the additional dispersive compensation of the scalar nonlinearity is provided by the vector nonlinearity, which prevents the structure from undergoing unlimited steepening. In this case, the vector nonlinearity "works" against the scalar one. In the other case, when at a given structure size the structure amplitude turns out to be too small, so that the wave dispersion exceeds the scalar nonlinearity, the vector scalar nonlinearities work against dispersion together. Thus, the mechanism for self-organization of solitary structures is associated with the compensation of wave dispersion by both the scalar and vector nonlinearities. As a result, a solitary structure is in general intrinsically anisotropic and contains a circular (monopolar) vortex superimposed on a dipole perturbation. The degree of anisotropy increases sharply as the size of the vortex approaches the intermediate size (16). When the scalar nonlinearity prevails over the vector one [in case of sufficiently large sizes, see Eq. (16)] only monopolar structures exist. Such solitary structures of monopole type were first found in laboratory modeling of solitary Rossby vortices.[30] Numerical calculations of the generalized HM equation containing both scalar and vector nonlinearities were performed in Ref. 31. It was shown here that a large-scale dipole vortex splits into two monopoles (a cyclone and an anticyclone), where a vortex of one polarity is long lived whereas the vortex of the opposite polarity disperses. In case of drift waves, only the anticyclones that propagate faster than the maximum velocity of the corresponding linear waves survive.

In the case when

$$\frac{\omega}{\omega_{Bi}} \gg \frac{r_L^4}{a^4}, \quad \frac{r_L^2}{a^2}\frac{a}{L} \ll 1, \tag{17}$$

we may be convinced that the first nonlinear term $\propto \partial\varphi^2/\partial t$ in Eq. (14) is more important than others and we get the following equation with only the scalar nonlinearity:

$$\left(\frac{n_{eo}}{n_{io}} + \frac{n_{po}}{n_{io}}\frac{T_e}{T_p}\right)\frac{\partial\varphi}{\partial t} - Z^2 r_L^2 \frac{\partial \nabla_\perp^2 \varphi}{\partial t} - Z^2 r_L^2 \omega_{Bi}\frac{n'_{io}}{n_{io}}\frac{\partial\varphi}{\partial y} + \frac{e}{2n_{io}}\left(\frac{n_{eo}}{T_e} - \frac{n_{po}T_e}{T_p^2}\right)\frac{\partial\varphi^2}{\partial t} = 0. \tag{18}$$

If the structure drifts with constant velocity v_D then considering the stationary case of propagation, one may obtain $\partial\varphi/\partial t \to -v_D \partial\varphi/\partial y$. In this way we may easily obtain the following equation:

$$\nabla_\perp^2 \varphi = \left[\frac{1}{Z^2 r_L^2}\left(\frac{n_{eo}}{n_{io}} + \frac{n_{po}}{n_{io}}\frac{T_e}{T_p}\right) + \omega_{Bi}\frac{n'_{io}}{n_{io}}\frac{1}{v_D}\right]\varphi + \frac{e}{2n_{io}Z^2 r_L^2}\left(\frac{n_{eo}}{T_e} - \frac{n_{po}T_e}{T_p^2}\right)\varphi^2. \tag{19}$$

According to the above mentioned discussion, this equation with only scalar nonlinearity describes solitary monopole vortical structures. Approximate analytical solution and numerical simulation of Eq. (19) confirm the solution in the form of monopole vortical structures as it is has the same form as Eq. (30) of Ref. 32.

In conclusion, in the present note we have discussed a new self-organization mechanism of formation of large-scale electrostatic drift vortical structures in EPI plasmas based on the competition between scalar and vector nonlinearities. As a result the solitary structure thus formed is intrinsically anisotropic and contains monopole vortex superimposed on a dipole perturbation. We have shown that the dynamics of low-frequency waves studied in usual EI plasmas is generally modified in EPI plasmas. These modifications depend on the concentration ratios of different species as well as on the magnitudes of their temperatures. Temperature inhomogeneity of electrons and positrons is taken into account. A new class of corresponding differential equations (14) and (18) with appropriate validity conditions (13), (16), and (17) is obtained. The generalized HM equation [Eq. (7)] valid for arbitrary sizes of structures is obtained. We have shown that due to the existence of positrons in the plasma, the sign of the derivative $Zn'_{io}(x) = n'_{eo}(x) - n'_{po}(x)$ may change which in turn enriches the class of solutions of the generalized HM equation. The new spatial structure given by Eq. (9) for the drift waves with dispersion relation given by Eq. (10) is obtained. Finally we note that the range of validity of the classical HM equation is revised and is given by Eq. (12).

Our results should be useful for understanding the properties of three-component EPI plasmas in laboratory experiments,[5,10,11] where positrons are used as probes and these types of electrostatic fluctuations are utilized for the diagnostic of EPI plasmas. Also our results are related to the localized nonlinear electrostatic structures, which may be connected to the observed large-scale density inhomogeneities of the universe.[1] After the identification of the drift modes and corresponding large-scale structures described in the given paper, their eventual observation should be used in the diagnostic of EPI plasma.

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