Dynamical behaviour of transverse plasmons in pair plasmas*

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This paper analytically investigates the nonlinear behaviour of transverse plasmons in pair plasmas on the basis of the nonlinear governing equations obtained from Vlasov–Maxwell equations. It shows that high frequency transverse plasmons are modulationally unstable with respect to the uniform state of the pair plasma. Such an instability would cause wave field collapse into a localized region. During the collapse process, ponderomotive expulsion is greatly enhanced for the increase of wave field strength, leading to the formation of localized density cavitons which are significant for the future experimental research in the interaction between high frequency electromagnetic waves and pair plasmas.

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1. Introduction

In recent years, collective wave modes and nonlinear phenomena in pair plasmas have aroused a great deal of interest for great progress in pair plasma creation. Pair plasmas are composed of negatively and positively charged particles with equal mass and absolute charge. Electron–positron plasmas as the most prominent member of pair plasma family are not easy to be generated and maintained because the annihilation time is short in comparison with the plasma period. Fortunately, a dense pair-ion plasma, consisting of equal-mass, positive and negative fullerene ions is produced in laboratory.[1−3] The pair-ion plasma, unlike the electron–positron plasma, has a long enough lifetime which makes experimental studies of the collective and nonlinear process in pair plasma under controlled conditions possible.

Pair plasmas are quite different from conventional electron–ion plasmas. Due to the absence of mass difference, the pair particles respond on the same time scale, i.e., the natural separation of time scale associated with high and low frequency motions no longer exists. Therefore the properties of waves in pair plasmas are distinct from those in conventional plasmas. For example, there is no Faraday rotation in pair plasmas,[4] with the immediate consequence that parallel propagating linear electromagnetic waves are not circularly but linearly polarised.[5] In addition, the well studied ion-acoustic waves have no counterpart in unmagnetized pair plasmas.

The unmagnetized pair plasmas only support two eigen wave modes:[6] the transverse electromagnetic mode and the electrostatic Langmuir mode. The dispersion relation for transverse mode with frequency very closed to $\omega_{\rm p}$ is[7]

$$\omega^2 = k^2c^2 + \omega_{\rm p}^2 \approx \omega_{\rm p}^2, \quad \omega >> kc, \tag{1}$$

where $\omega_{\rm p}^2 = \omega_{\rm pe}^2 + \omega_{\rm pp}^2$ is the combined plasma frequency; $\omega_{\rm pe}$ and $\omega_{\rm pp}$ are plasma frequency of negatively and positively charged particle, respectively. This wave is undamped in a collisionless plasma because its phase velocity always exceeds the speed of light c. The group velocity of the mode is very small, it is extremely difficult for this oscillation to escape the plasma. As a result, there will be strong nonlinear wave–wave wave–particle interactions. It is convenient to call the electromagnetic mode of Eq. (1) transverse plasmon.[8] In this paper, we will investigate theoretically the dynamical behaviour of transverse plasmons with dispersion relation Eq. (1).

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Although in the present situation it is extremely difficult in laboratory to excite electromagnetic modes of pair plasmas on account of the large ion mass, the relatively low density and temperature. Yet, an attempt is made to generate hydrogen pair-ion plasmas consisting of the light H^+–H^- ions, which are suitable for studying novel collective phenomena in the range of high frequencies.[3] The nonlinear interactions of transverse plasmons with plasmas assume an important role in laboratory, space and astrophysical applications, especially in laser plasma interaction. There are many papers respectively into the interactions of electromagnetic waves with pair plasmas, most of them are based on the fluid model.[9−18] However, when the characteristic scale of interest in the problem is smaller in comparison with the mean free path of charged particles, the kinetic effects start to dominate and fluid theory is not relevant. In this case, Vlasov–Maxwell equations must be employed to describe the wave–wave and wave–particle interactions. Actually, pair plasma produced in laboratory lies in collisionless regimes.[2,19]

Recently, considering the nonlinear wave–wave and wave–particle interactions, a set of nonlinear governing equations which can describe the interaction of transverse plasmons with pair plasmas have been obtained on the basis of Vlasov–Maxwell equations in Ref. [20]. In which, the numerical results show that as a result of modulational instability, the wave field will collapse and form a localized structure where strong transverse plasmons are trapped. However, in the numerical calculation a special initial condition is used. In order to obtain a general knowledge about the nonlinear behaviour of transverse plasmons in pair plasmas, in this paper we are devoted to the study of the modulational instability and the collapse behaviour of transverse plasmons from the governing equation in Ref. [20].

2. Dynamical behaviour of transverse plasmons

2.1. Nonlinear governing equation

On the basis of Vlasov–Maxwell equations, in consideration of the nonlinear wave–wave, wave–particle interactions, we have come to a set of nonlinear equations which can describe the nonlinear behaviour of high frequency transverse plasmons in unmagnetized isotropic pair plasmas,[20]

$$\mathrm{i}\frac{\partial}{\partial\tau}\tilde{\boldsymbol{E}}(\tilde{\boldsymbol{r}},\tau)-\nabla\times\nabla\times\tilde{\boldsymbol{E}}(\tilde{\boldsymbol{r}},\tau)-\tilde{n}(\tilde{\boldsymbol{r}},\tau)\tilde{\boldsymbol{E}}(\tilde{\boldsymbol{r}},\tau)=0, \tag{2}$$

$$\tilde{n}(\tilde{\boldsymbol{r}},\tau)=-|\tilde{\boldsymbol{E}}(\tilde{\boldsymbol{r}},\tau)|^2, \tag{3}$$

where $\tilde{\boldsymbol{r}}$, τ, $\tilde{n}$ and $\tilde{\boldsymbol{E}}$ are dimensionless position vector, time, low frequency density perturbation and slow varying envelope for transverse plasmons, respectively, their relation to dimensional ones are as follows:

$$\tilde{\boldsymbol{r}}=\frac{\omega_{\mathrm{p}}}{c}\boldsymbol{r},\quad \tau=\frac{\omega_{\mathrm{p}}}{2}t,\quad \tilde{n}(\tilde{\boldsymbol{r}},\tau)=\frac{n(\boldsymbol{r},t)}{2n_0},$$
$$\tilde{\boldsymbol{E}}(\tilde{\boldsymbol{r}},\tau)=\frac{\boldsymbol{E}(\boldsymbol{r},t)}{\sqrt{8\pi n_0 T_0}}. \tag{4}$$

In deriving equations (2) and (3) we have assumed that both of pair particles have the same density n_0 and temperature T_0 in equilibrium state. Equations (2) and (3) characterize the nonlinear coupling of transverse plasmons with quasistationary density perturbation driven by the ponderomotive force. It is obvious, as seen from Eq. (3), that the high-frequency plasmons are trapped inside local zone, where the ponderomotive force acts on the plasma so as to create a density depression. The last term in the left-hand side of Eq. (2) represents the nonlinear frequency shift caused by ion-density fluctuations, i.e., the usual modulation perturbation resulting from density rarefaction by the ponderomotive force. Substituting Eq. (3) into Eq. (2) one can obtain

$$\mathrm{i}\frac{\partial}{\partial\tau}\tilde{\boldsymbol{E}}(\tilde{\boldsymbol{r}},\tau)-\nabla\times\nabla\times\tilde{\boldsymbol{E}}(\tilde{\boldsymbol{r}},\tau)+|\tilde{\boldsymbol{E}}(\tilde{\boldsymbol{r}},\tau)|^2\tilde{\boldsymbol{E}}(\tilde{\boldsymbol{r}},\tau)=0. \tag{5}$$

It may be seen that the wave field of transverse plasmons is completely determined by Eq. (5). The numerical results of the nonlinear equation (5) show that the field will collapse into localized structure due to modulational instabilities.[20] In the following we will analytically prove that it is the case.

2.2. Modulational instability

Now, we study the stability of a finite amplitude monochromatic transverse waves in the Lyapunov sense with respect to Eq. (5). The following form of plane wave,

$$\tilde{\boldsymbol{E}}_{\mathrm{I}}=\tilde{\boldsymbol{E}}_0\exp[\mathrm{i}(\tilde{\boldsymbol{k}}_0\cdot\tilde{\boldsymbol{r}}-\tilde{\omega}_0\tau)], \tag{6}$$

is an exact solution of Eq. (5), provided that

$$\tilde{\omega}_0=\tilde{k}_0^2-\tilde{E}_0^2,\quad \tilde{\boldsymbol{k}}_0\cdot\tilde{\boldsymbol{E}}_0\propto\tilde{\boldsymbol{k}}_0\cdot\boldsymbol{e}_0=0, \tag{7}$$

where $\tilde{\boldsymbol{k}}_0$ is dimensionless wave vector, $\tilde{\boldsymbol{E}}_0 = \tilde{E}_0\boldsymbol{e}_0$, $\tilde{E}_0$ and $\boldsymbol{e}_0$ are real amplitude of transverse plasmons and real unit vector respectively,

$$|\tilde{\boldsymbol{E}}_0|^2 = \frac{|\boldsymbol{E}_0|^2}{8\pi n_0 T_0} << 1.$$

In order to investigate the stability of the above solution, we assume the presence of small perturbation ($\delta\tilde{\boldsymbol{E}}$) to the solution, if the disturbance is amplified, then the solution is unstable in the Lyapunov sense. Linearizing Eq. (5) in terms of the perturbation, we obtain

$$\begin{aligned}&\mathrm{i}(\delta\tilde{\boldsymbol{E}})_\tau - \nabla\times\nabla\times\delta\tilde{\boldsymbol{E}} + |\tilde{\boldsymbol{E}}_\mathrm{I}|^2\delta\tilde{\boldsymbol{E}}\\&+\tilde{\boldsymbol{E}}_\mathrm{I}(\tilde{\boldsymbol{E}}_\mathrm{I}\cdot\delta\tilde{\boldsymbol{E}}^* + \tilde{\boldsymbol{E}}_\mathrm{I}^*\cdot\delta\tilde{E}) = 0.\end{aligned} \tag{8}$$

We study the perturbation with the form

$$\begin{aligned}\delta\tilde{\boldsymbol{E}} = \frac{1}{2}\{&\tilde{\boldsymbol{E}}\exp[\mathrm{i}(\tilde{\boldsymbol{k}}_+\cdot\xi - \tilde{\omega}_+\tau)]\\&+\tilde{\boldsymbol{E}}^+\exp[-\mathrm{i}(\tilde{\boldsymbol{k}}_-\cdot\tilde{\boldsymbol{r}} - \tilde{\omega}_-\tau)]\},\end{aligned} \tag{9}$$

in which $\tilde{\boldsymbol{E}}$, $\tilde{\boldsymbol{E}}^+$ are transverse perturbations,

$$\tilde{\boldsymbol{E}} = \tilde{E}\boldsymbol{e}_1, \quad \tilde{\boldsymbol{E}}^+ = \tilde{E}^*\boldsymbol{e}_2^+, \quad \boldsymbol{e}_1\perp\tilde{\boldsymbol{k}}_+, \quad \boldsymbol{e}_2^+\perp\tilde{\boldsymbol{k}}_-, \tag{10}$$

where $\boldsymbol{e}_1$ and $\boldsymbol{e}_2^+$ are real unit vectors, and

$$\tilde{\boldsymbol{k}}_\pm = \tilde{\boldsymbol{k}} \pm \tilde{\boldsymbol{k}}_0, \quad \tilde{\omega}_\pm = \tilde{\omega} \pm \tilde{\omega}_0. \tag{11}$$

Substituting Eqs. (9) and (6) into Eq. (8), one can obtain

$$(\tilde{\omega}_- + \tilde{k}_-^2 - 2|\tilde{\boldsymbol{E}}_0|^2)\boldsymbol{e}_2^+\cdot\boldsymbol{e}_0 = |\tilde{\boldsymbol{E}}_0|^2\frac{\tilde{E}}{\tilde{E}^*}\boldsymbol{e}_0\cdot\boldsymbol{e}_1, \tag{12}$$

$$(\tilde{\omega}_+ - \tilde{k}_+^2 + 2|\tilde{\boldsymbol{E}}_0|^2)\boldsymbol{e}_0\cdot\boldsymbol{e}_1 = -|\tilde{\boldsymbol{E}}_0|^2\frac{\tilde{E}^*}{\tilde{E}}\boldsymbol{e}_0\cdot\boldsymbol{e}_2^+. \tag{13}$$

Combining Eqs. (12) and (13), we obtain the dispersion relation,

$$\begin{aligned}1 = |\tilde{\boldsymbol{E}}_0|^2\Big[&\frac{1}{\tilde{\omega} - \tilde{k}_0^2 + (\tilde{\boldsymbol{k}} - \tilde{\boldsymbol{k}}_0)^2}\\&-\frac{1}{\tilde{\omega} + \tilde{k}_0^2 - (\tilde{\boldsymbol{k}} + \tilde{\boldsymbol{k}}_0)^2}\Big].\end{aligned} \tag{14}$$

The dispersion equation (14) can be reduced to

$$(\tilde{\omega} - 2\tilde{k}\tilde{k}_0\cos\theta)^2 = \tilde{k}^4 - 2\tilde{k}^2|\tilde{\boldsymbol{E}}_0|^2, \tag{15}$$

where θ is the angle between $\tilde{\boldsymbol{k}}$ and $\tilde{\boldsymbol{k}}_0$. From above equation one can easily find when $\tilde{k}^2 < 2|\tilde{\boldsymbol{E}}_0|^2$ an instability occurs. In this case, the growth rate $\tilde{\gamma}$ is

$$\tilde{\gamma} = \tilde{k}\sqrt{2|\tilde{\boldsymbol{E}}_0|^2 - \tilde{k}^2}. \tag{16}$$

The maximum growth rate $\tilde{\gamma}_\mathrm{max}$ and the corresponding wave number $\tilde{k}_\mathrm{max}$ can be derived from Eq. (16) as follows:

$$\tilde{\gamma}_\mathrm{max} = |\tilde{\boldsymbol{E}}_0|^2, \quad \tilde{k}_\mathrm{max} = |\tilde{\boldsymbol{E}}_0|. \tag{17}$$

In terms of dimensional quantities, the maximum growth rate γ_max and characteristic scale l_c of the field are

$$\gamma_\mathrm{max} = \frac{1}{2}|\tilde{\boldsymbol{E}}_0|^2\omega_\mathrm{p}, \quad l_\mathrm{c} = \frac{2\pi}{\tilde{k}_\mathrm{max}} = \frac{2\pi}{|\tilde{\boldsymbol{E}}_0|^2}\frac{c}{\omega_\mathrm{p}}. \tag{18}$$

For modulation instability $\tilde{k} >> \tilde{k}_0$, the following condition should be satisfied:

$$\frac{\tilde{k}_0^2}{2} << |\tilde{\boldsymbol{E}}_0|^2. \tag{19}$$

From Eq. (18) one can easily obtain the result that stronger pump wave intensity will result in greater growth rate and smaller characteristic spatial scale. Equation (19) shows that for the longer wavelength of monochromatic pump wave there is a lower threshold for modulation instability to take place. This means the existence of a long wave instability in the interaction of transverse plasmons with pair plasmas.

2.3. Collapse of transverse plasmons

In the above we have proved that transverse plasmons in pair plasmas are modulationally unstable, now we will identify that the modulational instability will cause the collapse of wave field. If $\psi = \tilde{\boldsymbol{E}}$, $\psi_\tau = \tilde{\boldsymbol{E}}_\tau$, $\psi_{\tilde{r}_f} = \partial\tilde{\boldsymbol{E}}/\partial\tilde{r}_f$ ($\tilde{r}_f$ denotes the three component of vector $\tilde{\boldsymbol{r}}$, $f = 1, 2, 3$) and their conjugates are independent field functions, one can construct the following Lagrangian density:

$$\begin{aligned}L = \frac{\mathrm{i}}{2}&(\tilde{\boldsymbol{E}}^*\cdot\tilde{\boldsymbol{E}}_\tau - \tilde{\boldsymbol{E}}_\tau^*\cdot\tilde{\boldsymbol{E}})\\&-(\nabla\times\tilde{\boldsymbol{E}})\cdot(\nabla\times\tilde{\boldsymbol{E}}^*) + \frac{1}{2}|\tilde{\boldsymbol{E}}|^4,\end{aligned} \tag{20}$$

from which the Euler equation

$$\frac{\partial}{\partial\tilde{r}_f}\frac{\partial L}{\partial\psi_{\tilde{r}_f}} + \frac{\partial}{\partial t}\frac{\partial L}{\partial\psi_\tau} - \frac{\partial L}{\partial\psi} = 0 \tag{21}$$

gives Eq. (5).

The Lagrangian density is invariant for the phase transformations $\tilde{\boldsymbol{E}} \to \mathrm{e}^{\mathrm{i}\,\mu}\tilde{\boldsymbol{E}}$ and $\tilde{\boldsymbol{E}}^* \to \mathrm{e}^{-\mathrm{i}\,\mu}\tilde{\boldsymbol{E}}^*$. Thus, there is the following conservation equation:

$$\frac{\partial|\tilde{\boldsymbol{E}}|^2}{\partial t} + \nabla\cdot\boldsymbol{S} = 0, \tag{22}$$

in which the current density

$$\boldsymbol{S} = \mathrm{i}[\tilde{E}_j^* \nabla \tilde{E}_j - \tilde{E}_j \nabla \tilde{E}_j^* - \nabla \times (\tilde{\boldsymbol{E}} \times \tilde{\boldsymbol{E}}^*)], \quad (23)$$

where $j = 1, 2, 3$, any repeated index in a product of vectors, matrices and tensors implies a summation over the range of index values.

Equation (5) is also invariant with respect to space–time translations, so it leads to the momentum conservation

$$\frac{\partial \boldsymbol{P}}{\partial t} + \nabla \cdot \boldsymbol{G} = 0, \quad (24)$$

in which the momentum density $\boldsymbol{P}$ and the stress tensor $\boldsymbol{G}$ are given by

$$\boldsymbol{P} = \frac{\mathrm{i}}{2}(\tilde{E}_j \nabla \tilde{E}_j^* - \tilde{E}_j^* \nabla \tilde{E}_j), \quad (25)$$

$$\begin{aligned} G_{ij} = & -L\delta_{ij} + \left(\tilde{E}_{m,x_j} \frac{\partial \tilde{E}_m^*}{\partial x_i} - \tilde{E}_{m,x_j} \frac{\partial \tilde{E}_i^*}{\partial x_m}\right) \\ & + \left(\tilde{E}_{m,x_j}^* \frac{\partial \tilde{E}_m}{\partial x_i} - \tilde{E}_{m,x_j}^* \frac{\partial \tilde{E}_i}{\partial x_m}\right), \end{aligned} \quad (26)$$

and the energy conservation

$$\frac{\partial \varepsilon}{\partial t} + \nabla \cdot \boldsymbol{h} = 0, \quad (27)$$

where the energy density ε and the energy current density $\boldsymbol{h}$ are respectively

$$\varepsilon = (\nabla \times \tilde{\boldsymbol{E}}) \cdot (\nabla \times \tilde{\boldsymbol{E}}^*) - \frac{1}{2}|\tilde{\boldsymbol{E}}|^4, \quad (28)$$

$$\boldsymbol{h} = \tilde{\boldsymbol{E}}_\tau \times (\nabla \times \tilde{\boldsymbol{E}}^*) + \tilde{\boldsymbol{E}}_\tau^* \times (\nabla \times \tilde{\boldsymbol{E}}). \quad (29)$$

For localized field, Eqs. (22), (24) and (27) may be integrated to give the following conserved quantities for the plasmon number, the momentum and the energy.

$$N = \int |\tilde{\boldsymbol{E}}|^2 \mathrm{d}\tilde{\boldsymbol{r}}, \quad (30)$$

$$\boldsymbol{P}_0 = \frac{\mathrm{i}}{2} \int (\tilde{E}_j^* \nabla \tilde{E}_j - \tilde{E}_j \nabla \tilde{E}_j^*) \mathrm{d}\tilde{\boldsymbol{r}}, \quad (31)$$

$$\varepsilon_0 = \int (\nabla \times \tilde{\boldsymbol{E}}) \cdot (\nabla \times \tilde{\boldsymbol{E}}^*) - \frac{1}{2}|\tilde{\boldsymbol{E}}|^4 \mathrm{d}\tilde{\boldsymbol{r}}. \quad (32)$$

Following the method in Refs. [21] and [22], one can introduce the spatial average of any function $f(\boldsymbol{r})$ with respect to the plasmon number

$$\langle f(\tilde{\boldsymbol{r}}) \rangle = \int f(\tilde{\boldsymbol{r}}) \frac{|\tilde{\boldsymbol{E}}|^2}{N} \mathrm{d}\tilde{\boldsymbol{r}}. \quad (33)$$

Combining Eqs. (22), (23) and (31), we obtain

$$\langle \tilde{\boldsymbol{r}} \rangle_t = \int \tilde{\boldsymbol{r}} \frac{|\tilde{\boldsymbol{E}}|_t^2}{N} \mathrm{d}^3\tilde{\boldsymbol{r}} = \frac{2\boldsymbol{p}_0}{N}. \quad (34)$$

In the derivation of Eq. (34), we have used the approximation that all the field quantities are zero at infinity. Similarly, one can obtain

$$\begin{aligned} \frac{\partial^2}{\partial t^2} \langle \delta \tilde{\boldsymbol{r}}^2 \rangle = & \frac{\partial^2}{\partial t^2} \int (\tilde{\boldsymbol{r}} - \langle \tilde{\boldsymbol{r}} \rangle)^2 \frac{|\tilde{\boldsymbol{E}}|^2}{N} \mathrm{d}^3\tilde{\boldsymbol{r}} \\ = & 12\frac{\varepsilon_0}{N} - 8\left(\frac{\boldsymbol{p}_0}{N}\right)^2 \\ & - \frac{5}{N} \int |\nabla \times \tilde{\boldsymbol{E}}|^2 \mathrm{d}^3\tilde{\boldsymbol{r}}. \end{aligned} \quad (35)$$

Integrating Eq. (35) twice over time, one can obtain the spatial width of a wave packet

$$\langle \delta \tilde{\boldsymbol{r}}^2 \rangle = A\tau^2 + B\tau + C - \frac{5}{N} \int_0^\tau \mathrm{d}\tau_1 \int_0^{\tau_1} I(\tau) \mathrm{d}\tau, \quad (36)$$

in which, $A = \frac{4}{N}(3\varepsilon_0 - 2\frac{|\boldsymbol{p}_0|^2}{N})$, $I(\tau) = \int |\nabla \times \tilde{\boldsymbol{E}}|^2 \mathrm{d}^3\tilde{\boldsymbol{r}}$, B and C are integration constants.

Since $I(\tau)$ are positive definite, it follows that when $A < 0$, $\langle \delta \tilde{\boldsymbol{r}}^2 \rangle$ will collapse to such a value in a finite time until the field amplitude becomes too large for Eq. (5) to be true. During the collapsing, the field strength will increase, this fact can be inferred from the conservation of plasmon number N (see Eq. (30)).

3. Conclusions

In this paper we have analytically examined the dynamical behaviour of transverse plasmons in pair plasmas from Eq. (5). It turns out that the high frequency transverse plasmons are modulationally unstable. For a uniform pump field of the plasmon, the instability would lead the wave field to collapse into a narrow region. During the collapse process, the strength of localized wave field would increase. From Eq. (3), one can easily come to that with the increasing of field amplitude, enhanced ponderomotive expulsion will give rise to a localized density hole in which wave fields are trapped. These results are in agreement with those in numerical calculations.[20]

The maximum growth rate and the characteristic spatial scale for wave field and density caviton are given in Eq. (18). If pair particle density n_0 varies from 10^{10} to 10^{20} cm^{-3} for $\mathrm{H}^+ - \mathrm{H}^-$ plasmas, under the limiting condition that the energy density of transverse plasmons tends to approximately equal the thermal one of pair plasma, i.e., $|\tilde{\boldsymbol{E}}_0| \sim 1$, the characteristic scale of density caviton lies in the range from 1.8×10^3 to 1.8×10^{-2} cm. It means that this localized density hole may be observed in laboratory if the intensity of high frequency electromagnetic waves is strong enough and the density of pair plasmas is high enough. The obtained results are useful for the near future experimental research on the nonlinear behaviour of electromagnetic wave in pair plasmas.

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