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Drift eigenmodes in plasmas with negative ions

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Abstract. The instabilities associated with drift waves can be suppressed (or even removed) by introducing negative ions into the usual electron ion plasma. The presence of negative ions causes a decrease in the amplitude of the perturbed electrostatic fields which are mainly produced by the difference of electron and ion masses. It is shown that the presence of negative ions is a stabilizing effect for drift eigenmode in a slab geometry and the frequency of this electrostatic mode is reduced. In a pure pair-ion plasma the convective cell mode can become unstable because of shear flow under certain conditions.

1. Introduction

The pair-ion (PI) plasmas are not completely free of electrons in general. Saleem (2007) has presented a criterion to define a pure PI plasma. It puts a limit on the ratio of equilibrium density of electrons (n_{e0}) to the density of positive ions (n_{+0}) to ignore the electron dynamics in the plasma system, viz $\frac{n_{e0}}{n_{+0}} \ll \frac{m_e}{m_i}$, where m_e and m_i are the electron and ion masses, respectively.

A group of scientists (Oohara and Hatakeyama 2003, 2007; Oohara et al. 2005, 2007) has been trying to produce pure PI fullerene $C_{60}^{\pm}$ and Hydrogen ($H^{\pm}$) plasmas in laboratories. The magnetized fullerene PI plasma ($C_{60}^{\pm}$) has been produced in a cylinder with density $n \approx 10^8$ cm^{-3} and temperature $T_+ = T_- = T = 0.5$ eV (Oohara et al. 2005). Recently the low temperature hydrogen plasma ($H^+ + H^-$) has also been produced and confined magnetically with density $n \approx 10^{10}$ cm^{-3} (Oohara and Hatakeyama 2007). In the experiments with $C_{60}^{\pm}$ and $H^{\pm}$ plasmas, the acoustic waves and plasma waves have been investigated. Some strange behavior of excited electrostatic waves has also been quoted (Oohara et al. 2005). Saleem et al. (2006) showed that the experimental observations do not confirm that the plasma is a pure PI system. These authors used the quasi-neutrality to investigate ion acoustic wave (IAW) in a pair-ion–electron (PIE) plasma. Later, Saleem (2006) pointed out that the quasi-neutrality is not a good approximation for PIE plasmas. He also showed that the ion Landau damping of IAW reduces due to the presence of negative ions, and hence these waves can be exited easily in such systems. The frequency of the IAW can increase as the electron density reduces corresponding to certain range of plasma parameters. Therefore, larger frequency of the low frequency electrostatic mode observed in the experiment is itself an indication of the presence of electrons in the fullerene plasma. Moreover, the finite size of the system

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can also affect the behavior of externally excited perturbations. The wavelengths of the perturbations are comparable to the size of the system, therefore linear theory does not seem to be applicable to the experiment. A large number of papers has appeared on linear and nonlinear dynamics of (PI) and PIE plasmas (Schamel & Luque 2005; Shukla & Khan 2005; Vranjes and Poedts 2005; Kourakis et al. 2006; Verheest 2006; Mahajan and Shatashvili 2008). Shukla and Stenflo (2005) have considered finite amplitude waves in PI plasmas. The electrostatic drift eigenmodes have not been investigated theoretically in PIE plasmas. The presence of negative ions in high density low temperature plasma has been investigated in a hydrogen plasma in the context of nuclear fusion research by Shibata and Kasaki (2008). The density of negative ions (H^-) gives an insight into the importance of recombination processes activated by negative ions in the diverter plasma. The mechanism of 'plasma detachment' is proposed as a possible solution to reduce the excess heat to the diverter plate by Mathews (1995) and Takamura et al. (2002). Mohach et al. (2002) applied the negative ions (H^-), extracted from plasma sources, to generate charged and neutral particle beams from accelerator systems. Weiland (2000) discussed the reactive and dissipative instabilities associated with drift waves in fusion plasmas.

Levinton et al. (1995) and McCarthy and Maurer (1998) experimentally observed in electron–ion plasma that the ions flow along the external magnetic field $\mathbf{B_0}$ lines in tokamaks. The reason for ion shear flow along field lines has not been known. Assuming this fact to be given, the drift wave behavior has been investigated in a bounded slab geometry by Sen et al. (2000) for cold ions $T_i \ll T_e$, where T_e (T_i) are electron (ion) temperatures, respectively.

In the present work, the drift eigenmodes are investigated in a bounded (PIE) plasma. It is shown that the presence of negative ions suppresses the drift eigenmode in a slab geometry. Here we discuss the effects of negative ions on drift waves and investigate the local perturbation of pure PI plasmas in a simpler way. The electron temperature is also not assumed to be very high otherwise negative ions will be converted into positive ions. The collisions are neglected for simplicity. Such a study can be useful for long time confinement of cold plasmas for academic studies. Since it is not easy to get rid of electrons in the plasma, therefore we investigate the dynamics of both PIE and pure PI plasmas in a slab geometry on the lines of Sen et al. (2000). Furthermore we shall investigate the local perturbation of pure PI plasma with shear flow.

It is expected that the presence of negative ions can suppress the drift instabilities. If the cause of ions shear flow along $\mathbf{B}_0$ is due to some kind of zero order electric field $\mathbf{E}_0$ produced due to electrons interactions with the walls of the container, then negative and positive ions can have opposite shear flows. This is a speculation and therefore, we just consider both the possibilities; that is in Case I both polarity ions flow in the same direction along $\mathbf{B}_0$ and in Case II the positive ions flow along $\mathbf{B}_0$ and negative ions in opposite direction. In addition to eigenvalue problem in a slab geometry for cold positive and negative ions, we also investigate the local shear flow instability in pure PI plasma similar to electron–ion plasma case discussed first by D'Angelo (1965) and later with some modifications by Saleem et al. (2007). Both these research works are based on plasma fluid equations. In the drift ordering ($\partial_t < \Omega_i$) (where $\Omega_i = \frac{q_i B_0}{m_i c}$ is the ion gyro frequency) the electrons are assumed to follow Boltzmann density distribution.

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2. PIE plasma

Using the two fluid theory in slab geometry, the magnetic field is assumed to be sheared as $\mathbf{B}_0 = \mathbf{B}_0(\hat{\mathbf{z}} + \frac{x}{L_z}\hat{\mathbf{y}}) = \mathbf{B}_0(\hat{\mathbf{e}_\parallel})$ where we define $\kappa_B = |\frac{1}{\mathbf{B}_0}\frac{d\mathbf{B}_0}{dx}|$ and $L_s = \kappa_B^{-1}$ is the scale length of magnetic shear. The analogy of x, y and z directions in slab geometry with tokamak configuration is established with the radial, poloidal and toroidal directions, respectively. The background plasma has inhomogeneities along x-axis and the perturbations has the form $\varphi_{(x,t)} = \varphi_{(x)}e^{i(k_y y + k_z z - \omega t)}$. The magnetic shear can be produced by a zero-order current along z-axis because the shear flow velocities of ions are much smaller than the electrons, i.e., $v_{\pm 0(x)} \ll v_{eo(x)}$ and subscripts $(\pm)$ denote the positive and negative ions; respectively. Then

$$\mathbf{J}_0(x) = e(n_{+0}\mathbf{v}_{+0} - n_{-0}\mathbf{v}_{-0} - n_{e0}\mathbf{v}_{e0}) = e(n_{+0}v_{+0} - n_{-0}v_{-0} - n_{e0}v_{e0})\hat{\mathbf{z}}, \quad (2.1)$$

whereas we may have $\mathbf{v}_{-0} = \mathbf{v}_{+0}$ or $\mathbf{v}_{-0} = -\mathbf{v}_{+0}$. Furthermore, we assume $\nabla n_{j0} = -\hat{\mathbf{x}}\frac{dn_{j0}}{dx}$. The momentum equations for ions can be written as,

$$(\partial_t + \mathbf{v}_\pm.\nabla)\mathbf{v}_\pm = \frac{q_\pm}{m_\pm}\left(\mathbf{E} + \frac{1}{c}\mathbf{v}_\pm \times B\right) - \frac{1}{m_\pm n_\pm}(\nabla p_\pm + \nabla.\Pi_\pm), \quad (2.2)$$

where q and m are charge and mass of particles, respectively, and c denotes the speed of light. Assuming $\mathbf{E} = -\nabla\varphi$, we take the cross product of (2.2) with $\hat{\mathbf{e}_\parallel}$ to have perpendicular perturbed velocities of cold ions as,

$$\mathbf{v}_{\pm\perp} = \frac{c}{\mathbf{B}_0}\mathbf{E}_\perp \times \hat{\mathbf{e}_\parallel} \mp \frac{1}{\Omega_i}(\partial_t + \mathbf{v}_\pm.\nabla)\mathbf{v}_{\pm\perp} \times \hat{\mathbf{e}_\parallel} = \mathbf{v}_E + \mathbf{v}_{\pm p}, \quad (2.3)$$

where $\omega_i = \frac{|q_\pm|\mathbf{B}_0}{m_i c}$ and we assume $m_+ = m_- = m_i$. In drift ordering we have $|\mathbf{v}_{\pm p}| \approx \epsilon|\mathbf{v}_E|$ where $\epsilon \ll 1$. Therefore,

$$\nabla.(n_\pm \mathbf{v}_{\pm p}) \approx n_{\pm 0}\nabla.\mathbf{v}_{\pm p} \approx \mp\frac{c}{\mathbf{B}_0}\frac{1}{\Omega_i}(\partial_t \pm \mathbf{v}_{\pm\parallel 0}.\nabla)\nabla_\perp^2\varphi. \quad (2.4)$$

First we consider the case I with $\mathbf{v}_{-0} = \mathbf{v}_{+0}$ and hence the parallel components of (2.2) become for +ve and –ve cold ions, respectively, as

$$(\omega - \omega_\pm)v_{\pm\parallel} \approx (\pm c_s^2 k_\parallel)\Phi - (D_e k_y A_\pm)\Phi, \quad (2.5)$$

where $A_+ = \frac{1}{\Omega_i}\frac{dv_{+\parallel 0}}{dx}, A_- = \frac{1}{\Omega_i}\frac{dv_{-\parallel 0}}{dx}, c_s^2 = \frac{T_e}{m_i}, D_e = \frac{cT_e}{e\mathbf{B}_0}, \omega_\pm = v_{\pm 0}k_\parallel$ and $\Phi = \frac{e\varphi}{T_e}$. The continuity equations can be written for $\pm$ ions as,

$$(\partial_t + v_{\pm\parallel 0}\partial_\parallel)n_\pm + \nabla n_{\pm 0}.\mathbf{v}_E + n_{\pm 0}\nabla_\perp.\mathbf{v}_{\pm p} + n_{\pm 0}\partial_\parallel v_{\pm\parallel} = 0. \quad (2.6)$$

These equations yield for +ve and −ve ions, respectively,

$$-\rho_s^2 n_{+0}\Omega_+^2\nabla_\perp^2\Phi + \Omega_+^2 n_+ - D_e k_y\frac{dn_{+0}}{dx}\Omega_+\Phi - n_{+0}\left\{-D_e k_y k_\parallel\frac{dv_{+\parallel 0}}{dx} + c_s^2 k_\parallel\right\}\Phi = 0 \quad (2.7)$$

and

$$\rho_s^2 n_{-0}\Omega_-^2\nabla_\perp^2\Phi + \Omega_-^2 n_- - D_e k_y\frac{dn_{-0}}{dx}\Omega_-\Phi - n_{-0}\left\{-D_e k_y k_\parallel\frac{dv_{-\parallel 0}}{dx} - c_s^2 k_\parallel\right\}\Phi = 0. \quad (2.8)$$

Let us define $\kappa_{\pm n} = |\frac{1}{n_{\pm 0}}\frac{dn_{\pm 0}}{dx}|, L_{\pm n} = \kappa_{\pm n}^{-1}$, and $\omega_\pm^*(x) = -D_e k_y \kappa_{\pm n}$. We further define $k_\parallel' = \frac{k_y}{L_s} = k_y\frac{x}{L_s}$ where $k_\parallel$ increases with x (the distance from the mode rational surface defined by $x = 0$ where $\mathbf{k}.\mathbf{B} = 0$). Then (2.7) and (2.8) can be

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written, respectively, as

$$(\rho_s^2\nabla_\perp^2\varphi)n_{+0} - n_+ + n_{+0}\frac{\omega_+^*}{\Omega_+}\Phi - n_{+0}\frac{\rho_s k_y}{\Omega_+}\frac{dv_{+\|0}}{dx}\frac{x}{x_{+s}}\Phi + n_{+0}\frac{x^2}{x_{+s}^2}\Phi = 0 \qquad (2.9)$$

and

$$(\rho_s^2\nabla_\perp^2\varphi)n_{-0} - n_- - n_{-0}\frac{\omega_-^*}{\Omega_-}\Phi - n_{-0}\frac{\rho_s k_y}{\Omega_-}\frac{dv_{-\|0}}{dx}\frac{x}{x_{-s}}\Phi + n_{-0}\frac{x^2}{x_{-s}^2}\Phi = 0, \qquad (2.10)$$

where $\omega_\pm^* = \pm k_y\rho_s c_s\frac{1}{L_{\pm n}}$ and $\rho_s = \frac{c_s}{\Omega_i}$. These two equations can be combined to have,

$$\rho_s^2\nabla_\perp^2\Phi(n_{+0}+n_{-0}) - (n_+ - n_-) + \frac{\omega_e^*}{\Omega_+}n_{+0}\Phi - \frac{\omega_e^*}{\Omega_-}n_{-0}\Phi - n_{+0}\frac{\rho_s k_y}{\Omega_+}\frac{dv_{+\|0}}{dx}\frac{x}{x_{+s}}$$

$$\Phi - n_{-0}\frac{\rho_s k_y}{\Omega_+}\frac{dv_{-\|0}}{dx}\frac{x}{x_{-s}}\Phi + \left(n_{+0}\frac{x^2}{x_{-s}^2} + n_{-0}\frac{x^2}{x_s^2}\right)\Phi = 0, \qquad (2.11)$$

where $\omega_+^* = \omega_-^* = \omega_e^*$. Now we consider the case I with $\mathbf{v}_{+\|0} = \mathbf{v}_{-\|0} = \mathbf{v}_0$. Then we have, $\Omega_0 = \Omega_+ = \Omega_-$, and $x_{+s} = x_{-s} = x_s$. Hence (2.11) becomes,

$$\rho_s^2\nabla_\perp^2\Phi(n_{+0}+n_{-0}) - (n_+ - n_-) + \frac{\omega_e^*}{\Omega_0}(n_{+0}-n_{-0})\Phi - (n_{+0}+n_{-0})\frac{\rho_s k_y}{\Omega_0}\frac{dv_0}{dx}\frac{x}{x_s}\Phi$$

$$+ (n_{+0}+n_{-0})\frac{x^2}{x_s^2}\Phi = 0. \qquad (2.12)$$

In the slow time scale, we assume electrons to follow the Boltzmann distribution,

$$n_e \approx n_{e0}(x)e^{\frac{e\varphi}{T_e}}. \qquad (2.13)$$

Poisson equation can be written as,

$$(1 - \lambda_{De}^2\nabla^2)\Phi = \frac{1}{n_{e0}}(n_+ - n_-), \qquad (2.14)$$

where $\lambda_{De}^2 = \frac{T_e}{4\pi n_{e0}e^2}$. In deriving (2.14) we have assumed $\frac{e\varphi}{T_e} < 1$ which is quite possible even for $T_e \approx 1ev$. This follows from the assumption that the perturbed electrostatic field energy is lesser than the equilibrium electron thermal energy. The plasma is assumed to be low-density and hence collisions are neglected for simplicity. But in low temperature denser plasmas, collisions can be very important. Since $|\partial_z| \ll |\nabla_\perp|$, therefore (2.12–2.14) yield,

$$L_0^2(d_x^2 - k_y^2)\Phi - \Phi + \frac{\omega_e^*}{\Omega_0}\Phi - \frac{\rho_s k_y}{\Omega_0}\frac{dv_0}{dx}\frac{x}{x_s}\Phi + N_0\frac{x^2}{x_s^2}\Phi = 0, \qquad (2.15)$$

where $L_0^2 = (N_0\rho_s^2 + \lambda_{De}^2)$ and $N_0 = \frac{(n_{+0}-n_{-0})}{n_{e0}}$. Equation (2.15) can be rewritten as,

$$L_0^2\frac{d^2\Phi}{dx^2} + (\Lambda + Px^2 - Qx)\Phi = 0, \qquad (2.16)$$

where

$$P = N_0\frac{L_n^2}{\rho_s^2 L_s^2} - 2\frac{v_{1100}}{c_s}\frac{L_n^2}{L_{v2}^2}\frac{1}{\rho_s L_s} - \frac{1}{L_n^2},$$

$$Q = \frac{v_{1100}}{c_s}\frac{L_n^2}{L_{v1}}\frac{1}{\rho_s L_s},$$

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and

$$\Lambda = \frac{\omega_0^*}{\Omega_0} - k_y^2 L_0^2 - 1.$$

The above relations have been obtained by assuming the following quadratic shear flow profile in slab geometry used also by Sen et al. (2000),

$$v_{110}(x) = v_{1100} + v_{110}^{'} x + \frac{1}{2} v_{110}^{''} x^2,$$

where $v_{110}^{'} = \frac{v_{1100}}{Lv_1}$ and $\frac{1}{2} v_{110}^{''} = \frac{v_{1100}}{Lv_2}$. It may be mentioned here that in a cylindrical geometry the shear flows are commonly assumed to have Gaussian profiles. Note that we have assumed $\omega_e^*(x) = \omega_0^*(1 - \frac{x^2}{L_n})$ similar to Sen et al. (2000) where $\omega_0^* = v_0^* k_y$ and $v_0^* = |\frac{\rho_s c_s}{L_n}|$. If $N_0 = 1$, i.e., $n_{-0} = 0$, then (2.16) is the same as (2) of Sen et al. 2000. But it is not a matter of N_0 only which seems to be a very minor modification to the electron–ion plasma case. Physically, the large N_0 indicates that quasi-neutrality does not remain valid because λ_{De} becomes large. Therefore we have to use Poisson equation as has been pointed out by Saleem (2006).

As N_0 becomes very large, the (2.16) changes and in case of pure PI plasma it gives a local shear flow instability. In usual EI plasma $\lambda_{De} < \rho_s$ holds because $v_A < c$ (where v_A is Alfven speed). Within the fluid framework the limit $\rho_s k \leqslant 1$ should hold, in general. Therefore, we have $\lambda_{De}^2 k^2 \ll 1$ and quasi-neutrality is used. But as n_{e0} reduces and N_0 increases, the terms Q and Λ become very small. This case will be discussed later. The solution of (2.16) for $P < 0$ is given as,

$$\Phi = \Phi_0 exp\left\{ -\frac{|P|^{\frac{1}{2}}}{2L_0}(x - x_0)^2 \right\}, \tag{2.17}$$

where $x_0 = \frac{Q}{2|P|}$. This solution satisfies the physical boundary condition $\Phi \to 0$ as $x \to \pm\infty$. The perturbation does not propagate in this case. For $0 < P$, the solution becomes,

$$\Phi = \Phi_0 exp\left\{ -i\frac{|P|^{\frac{1}{2}}}{2L_0}(x + x_0)^2 \right\}. \tag{2.18}$$

This is a non-localized mode propagating outward. It is expected that the drift wave is generated near $k_\parallel = 0$ where the electron shielding is inefficient and then propagate outwards with increasing x. Then the $k_\parallel$ increases such that $k_\parallel v_{ti} \approx \omega$ (where v_{ti} is ion thermal velocity), the Landeu damping sets in and wave is absorbed. The wave does not reach the boundary and hence it is generally called the outgoing boundary condition. The perturbation will decay in the absence of any energy source. For this mode we find,

$$\Omega_0 = \left(\frac{\omega_0^*}{H_0} - i\frac{L_0 |P|^{\frac{1}{2}}}{H_0} \Omega_0 \right), \tag{2.19}$$

where $H_0 = (1 + L_0^2 k_y^2 + \frac{Q}{4|P'|})$.

Since Ω_0 is doppler-shifted frequency, and if we define $\omega = \omega_r + i\gamma$, then we get

$$\omega_r = \left(\omega_0 + \frac{\omega_0^*}{H_0} \right) \tag{2.20}$$

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and

$$\gamma \approx -\frac{L_0|P|^{\frac{1}{2}}}{H_0}\Omega_0 \approx -\frac{L_0|P|^{\frac{1}{2}}\omega_0^*}{H_0}, \tag{2.21}$$

where $\Omega_0 \approx \omega_0^*$ has been approximated in the imaginary term of (2.19).

The damping rate is proportional to ω_0^* and

$$\omega_0^* \propto \frac{d}{dx}(n_{+0} - n_{-0}). \tag{2.22}$$

If drift wave is derived by some mechanism and the growth rate is γ_0, then the perturbation will be stable if

$$\gamma_0 < \gamma. \tag{2.23}$$

The real frequency ω_0^* of drift waves decreases due to the presence of negative ions, i.e., $n_{-0} \neq 0$. Therefore very low frequency drift waves may be produced in such systems. The growth rate γ_0 of drift wave may also become smaller due to $n_{-0} \neq 0$, therefore a smaller damping rate γ in (2.21) does not necessarily mean that the wave will be destabilized due to negative ions. Rather the drift wave may disappear when $n_{+0} \approx n_{-0}$. Hence the addition of negative ions in electron–ion plasma is a stabilizing effect. In pair-ion plasmas, the drift wave will indicate the presence of electrons in the system.

For $n_{-0} \neq 0$ and $v_0(x) = 0$, (2.19) becomes

$$\omega = \frac{\omega_0^*}{1 + \rho_s^2 k_y^2}\left(1 - i\frac{L_n}{L_s}\right), \tag{2.24}$$

which is the same as (5.38) of Weiland (2000) for electron–ion plasma. In the quasi-neutrality limit $\lambda_{De}^2 k^2$ has been neglected while in the case of PIE plasma we are not allowed to neglect this term, in general.

In case II, where we have $v_{+110} = -v_{-110}$, the (2.15) has a modification in terms containing shear flow effect and becomes,

$$L_0^2(d_x^2 - k_y^2)\Phi - \Phi + \frac{\omega_e^*}{\Omega_0}\Phi - N_0\frac{\rho_s k_y}{\Omega_0}\frac{dv_0}{dx}\frac{x}{x_s}\Phi + N_0\frac{x^2}{x_s^2}\Phi = 0. \tag{2.25}$$

The solutions will be similar to the above case.

3. Pure PI plasma

If the plasma fulfills the criterion of a pure PI plasma presented by Saleem (2007), then electron temperature T_e becomes irrelevant and we write (2.7) and (2.8) in unnormalized form, respectively, as

$$-\frac{1}{m_i}\frac{1}{\Omega_i^2}n_{+0}\Omega_0^2\nabla_\perp^2(e\varphi) + \Omega_0^2 n_+ - \frac{c}{e\mathbf{B}_0}k_y\frac{dn_{+0}}{dx}\Omega_0(e\varphi)$$
$$-n_{+0}\left\{-\frac{c}{e\mathbf{B}_0}k_y k_\parallel\frac{dv_0}{dx} + \frac{1}{m_i}k_\parallel^2\right\}(e\varphi) = 0 \tag{3.1}$$

and

$$\frac{1}{m_i}\frac{1}{\Omega_i^2}n_{-0}\Omega_0^2\nabla_\perp^2(e\varphi) + \Omega_0^2 n_- - \frac{c}{e\mathbf{B}_0}k_y\frac{dn_0}{dx}\Omega_0(e\varphi)$$
$$-n_{-0}\left\{-\frac{c}{e\mathbf{B}_0}k_y k_\parallel\frac{dv_0}{dx} - \frac{1}{m_i}k_\parallel^2\right\}(e\varphi) = 0. \tag{3.2}$$

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where we have assumed $v_{+110} = v_{-110}$. In this case $\mathbf{B}_0 = \mathbf{B}_0\hat{\mathbf{z}}$ is constant and the effects of density gradient and shear flow disappear. The local analysis yields the pair ion convective cell (PPCC) mode (Saleem et al. 2006),

$$\Omega_0^2 = \frac{k_\parallel^2}{k_\perp^2}\Omega_i^2. \tag{3.3}$$

If the shear flow of positive and negative ions is in opposite directions, then the PPCC mode becomes unstable and the local linear dispersion relation turns out to be,

$$\Omega_0^2 = \left(1 - \frac{k_y}{k_\parallel}A_0\right)\Omega_i^2\frac{k_\parallel^2}{k^2}, \tag{3.4}$$

where $A_0 = \frac{1}{\Omega_i}\frac{dv_0}{dx}$. The instability condition turns out to be the same as that of usual shear flow instability found long ago (D' Angelo 1965),

$$\frac{k_\parallel}{k_y} < A_0 \tag{3.5}$$

However, for $A_0 < \frac{k_\parallel}{k_y}$, we have a stable PPCC mode modified due to shear flow.

4. Conclusions

It has been shown that the presence of negative ions in plasmas reduces the electrostatic field and hence suppresses the associated instabilities as well as damping rates. The instabilities associated with plasma inhomogenities are the most important hurdles in plasma confinement. These instabilities arise mainly due to the difference in dynamics of ions and electrons because of their different masses. The drift wave type instabilities disappear in a pure pair ion (PI) plasma. The presence of negative ions reduces the magnitude of electrostatic field and it helps in suppressing the drift wave instabilities.

In a pair-ion-electron (PIE) plasma the drift wave amplitude driven by shear flow is reduced compared to usual (EI) plasma because ω_0^* reduces as can be seen in (2.21). The frequency of the local drift wave also decreases as the density of electrons decreases (Saleem 2007). Therefore the instabilities associated with drift waves are automatically suppressed. Since negative ions can be converted into positive ions in PIE plasma if the electron temperature is high, therefore the present investigation is applicable to low temperature plasmas. The long time cold plasma confinement in toroidal devices can be useful for academic studies. However the pair ion convective cell mode can become unstable locally due to shear flow if positive and negative ions move in opposite directions. If they flow in the same direction, then the instability does not appear.

The effects of shear flow also disappear in PI plasmas like the effects of density inhomogeneity. Therefore pure PI plasma of deuterium or tritium can be a very good source of fusion. However, it does not seem easy to produce PI plasma at high temperatures needed for fusion.

In PI plasma experiments, the efforts should also be made to detect drift waves. The present investigation can be useful in this regard. The existence of drift waves will confirm or otherwise the presence of electrons in PI plasmas similar to the case of IAWs as has been pointed out by Saleem (2006).

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