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Effects of the particle spin polarisation on the unstable modes in the anisotropic dense system

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Abstract. Polarisation of the particle spin can be an important problem for different plasmas. In this article, the contribution of the electron spin on the growth rate of the temperature anisotropy of electromagnetic instabilities has been investigated. Results show that polarisation of the electron spin will restrict the instability growth rate while instability can survive due to the spin-depolarised electrons even when the requested temperature anisotropy is vanished. Instability can reach the damping state exponentially due to the spin-polarised electrons while it can grow linearly due to the spin-depolarised (the semi-classical) electrons.

Keywords. Temperature anisotropy; electromagnetic instability; spin-polarised properties; unstable modes.

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1. Introduction

The electromagnetic instabilities in the velocity space have always been an interesting issue for researchers because of the possibility to generate strong magnetic fields in the inertial confinement fusion (ICF) plasmas [1–5]. Controlling the instabilities has been one of the important problems in the transport of the energetic particle to the compressed fuel core [6,7]. Recent studies show the importance of polarisation of the particle spin to obtain high energy gain and fraction of the burned fuel can be provided in a low ignition threshold and requested driver energy by increasing the thermonuclear cross-section for the spin-polarised fuel initially [8,9]. On the other hand, the fuel can be depolarised when the spin–orbit interactions of electron–triton and electron–^{3}He are present where this can be small because of the low cross-section compared to the fusion interaction cross-section [10]. The impact of the particle spin polarisation on the electromagnetic instabilities present in the velocity space depends on the electromagnetic instabilities. The basic theory will be introduced initially according to two different kinetic models for spin-polarised and depolarised systems in §2. The semiclassical properties of the particle spin will be used for depolarised cases while full quantum properties will be used in the spin-polarised cases. Finally, results and conclusion will be presented in §3.

2. Basic model

There have been many studies in fluid models [11–13] where the magnetic dipole force and the magnetisation current associated with the electron spin have been included while, there have been more accurate models applying kinetic theory [13]. Here, the theoretical model is based on the kinetic theory where, two different models have been applied for investigating the electron spin properties. Studies show that a Wigner transform in phase space and a Q-transform in spin space can produce a suitable physical evolution equation for fully quantum plasmas when the spin degrees of freedom are included [14]. The main purpose of this study is to compare the polarisation effects of the particle spin. It is assumed that the de Broglie wavelength is smaller than the plasma long scale so that the spin effects can have more important contributions than other quantum effects (such as the Fermi pressure and the particle dispersive effects). Therefore, the evolution of the particle distribution function, $f(r, v, \hat{s}, t)$, can be obtained by [15]

$$\frac{\partial f}{\partial t} + \mathbf{V}\cdot\nabla_{\mathbf{x}} f + \left[\frac{q_e}{m_e}(\mathbf{E}+\mathbf{V}\times\mathbf{B}) + \frac{\mu_e}{m_e}\nabla_{\mathbf{x}}(\hat{\mathbf{S}}\cdot\mathbf{B}+\mathbf{B}\cdot\nabla_{\hat{\mathbf{s}}})\right]\cdot\nabla_{\mathbf{v}} f + \frac{2\mu_e}{\hbar}(\hat{\mathbf{S}}\times\mathbf{B})\cdot\nabla_{\hat{\mathbf{s}}} f = 0, \quad (1)$$

where $q_e = -e$, $\mu_e = -g e\hbar/4m_e$, $g = 2.002319$, m_e, $\hat{\mathbf{S}}$ and $\hbar$ are the charge of an electron, the electron magnetic moment, the electron spin g factor, the electron mass and the unit spin vector and the Planck's constant divided by 2π, respectively. The contributions of the plasma ions have been ignored because of their larger mass and smaller magnetic moment than the electron. The spin terms proportional to $\hat{\mathbf{S}} \cdot \mathbf{B}$ and $\hat{\mathbf{S}} \times \mathbf{B}$ are due to the magnetic dipole force and the spin precession, respectively. Notice that, spin is an independent quantity and the term proportional to $\mathbf{B} \cdot \nabla_{\hat{\mathbf{s}}}$ will be a modification of the magnetic dipole force which is due to the spread out nature of the spin probability distribution. The quantities $\mathbf{E}$ and $\mathbf{B}$ are the electric and magnetic fields, respectively. The proposed unstable waves are the transverse electromagnetic instabilities that propagate parallel to the homogeneous background magnetic field by the wave number $\mathbf{K} = k_z\hat{z}$.

In the opposite case (for investigating depolarisation of the spin effect), it will be assumed that the spin effects are due to the electron spin that may survive even when macroscopic variations occur. As the plasma system can behave classically, in this case, interesting effects will be due to the semiclassical properties of the electron spin. Therefore, suitable physical evolution of the particle distribution can be obtained as the Vlasov model where, delocalisation of the particle wave function cannot be too large. Thus it can be introduced as [16]

$$\frac{\partial f}{\partial t} + \mathbf{V} \cdot \nabla_{\mathbf{x}} f + \left[\frac{q_e}{m_e} (\mathbf{E} + \mathbf{V} \times \mathbf{B}) + \frac{2\mu_e}{m_e \hbar} \nabla_{\mathbf{x}} (\mathbf{S} \cdot \mathbf{B}) \right] \cdot \nabla_{\mathbf{v}} f + \frac{2\mu_e}{\hbar} (\mathbf{S} \times \mathbf{B}) \cdot \nabla_{\mathbf{s}} f = 0, \tag{2}$$

where the quantity $\mathbf{S}$ is the spin vector with a fixed length $|S| = \hbar/2$. The function f is the sum of the perturbed and unperturbed distribution functions of the particle. The perturbed quantities are approximated by a standard ansatz of quasimonochromatic harmonic variation, for example according to $f_1 = \tilde{f}_1 \exp[i(\mathbf{K} \cdot \mathbf{X} - \omega t)]$. The perturbed distribution function, $\tilde{f}_1$, can be solved in terms of the orthogonal eigenfunctions to the right-hand side operator so that

$$\tilde{f}_1 = \sum_{a=-\infty}^{+\infty} \sum_{b=-\infty}^{+\infty} \frac{1}{2\pi} g_{ab}(v_\perp, v_z, \theta_s) \times \exp(-ia\varphi_v) \exp(-ib\varphi_s). \tag{3}$$

Different quantum effects such as the Fermi–Dirac static, Landau quantisation and spin-splitting can be contained in the unperturbed distribution function for quantum plasmas. The first two cases can be affected by the plasma long scale and the de Broglie wavelength. Here, the chemical potential is large and difference between the nearby Landau levels is smaller than thermal energy. These effects can be suppressed where the velocity distribution function of the electrons approaches the classical Maxwellian. In addition, the probability distribution of the spin-up and spin-down states of the electron populations due to the spin splitting effect can be still important. Thus, the distribution can be approximated for the spin-polarised plasmas as [17]

$$f_0 = n_0 \left[\frac{m_e^{3/2}}{4\pi (2\pi K_B)^{3/2} T_\perp T_z^{1/2}} \right] \times \exp\left(-m \left(\frac{v_\perp^2}{2T_\perp K_B} + \frac{v_Z^2}{2T_Z K_B} \right) \right) \times \left[\exp\left(\frac{\mu_e B_0}{K_B T_{sp}} \right) + \exp\left(-\frac{\mu_e B_0}{K_B T_{sp}} \right) \right]^{-1} \times \sum_{\nu=+,-} \exp\left(\frac{\nu \mu_e B_0}{K_B T_{sp}} \right) (1 + \nu \cos\theta_s), \tag{4}$$

where the signs $+$ and $-$ indicate the spin-up and spin-down states respectively and the requested temperature anisotropy by the instability is included. For the spin-depolarised cases, the probability distribution of the spin states is equal to unity as the unperturbed distribution function is of the form:

$$f_0(v_\perp, v_z, \theta_s) = \frac{n_0}{4\pi T_\perp T_z^{1/2}} \left(\frac{m_e}{2\pi K_B} \right)^{3/2} \frac{\alpha}{\sinh\alpha} \times \exp\left(-\frac{m v_\perp^2}{2K_B T_\perp} - \frac{m v_z^2}{2K_B T_z} \right) \times \exp(-\alpha \cos\theta_s), \tag{5}$$

where $\alpha = \mu_e B_0 / K_B T_z$. The quantity T_{sp} defines spin temperature which can be introduced when the coupling between the spin freedom is strong and the spin can behave as an independent variable. For semiclassical cases, this will be replaced by a suitable thermal temperature.

Studying the instability requires evaluating a suitable dispersion relation. In this order, derivation of the general dispersion function is based on the relation $\det(D_{ij}) = 0$ where

$$D_{ij} = \delta_{ij} \left(1 - \frac{c^2 k^2}{\omega^2} \right) - \frac{k_i k_j c^2}{\omega^2} \delta_{ij} + \frac{i}{\varepsilon_0 \omega} \sigma_{ij}.$$

The quantities δ_{ij} and ε_0 are the Kronecker delta and the permittivity constant, respectively. The quantity σ_{ij} is the conductivity tensor, which is proportional to the current density as $J_i = \sum_j \sigma_{ij} E_j$. The current density can be approximated as

$$\mathbf{J} = \int q_e \mathbf{V} f_1 \mathrm{d}^2 s \mathrm{d}^3 v + \nabla \times \mathbf{M}. \tag{6}$$

The first and second terms define the free and magnetisation current due to the electron population where the contributions of the plasma ions have been ignored because of their larger mass and smaller magnetic moment than the electron. Suitable definition of the magnetisation current will be important. For the spin-polarised plasmas, the magnetisation current can be different from the semi-classical cases by a factor 3 due to the available quantum states. The magnetisation current can be approximated in the spin-polarised and semi-classical cases (the spin-depolarised cases) respectively as follows:

$$\mathbf{J}_{\text{M-pol}} = \sum_{\nu=+,-} \nabla \times \left(3\mu_e \int \hat{\mathbf{S}} f_{1\nu} \mathrm{d}^2 s \mathrm{d}^3 v \right) \tag{7}$$

and

$$\mathbf{J}_{\text{M-depol}} = \mu_e \nabla \times \int \hat{\mathbf{S}} f_1 \mathrm{d}^2 s \mathrm{d}^3 v. \tag{8}$$

It is well known that for a parallel propagation to an external magnetic field, in addition to a longitudinal electrostatic mode, there are two other modes, the right and circular polarisation, so that a linear polarised mode approximately can be suitable for introducing the dispersion relations of the unstable modes for wave frequency much higher than the electron cyclotron frequency, or for the wave frequency much lower than the ion cyclotron frequency. Here, effects of the ions are ignored because of the above-mentioned reasons. Therefore, only the cyclotron effects of the electrons are considered and it will be assumed that, the wave frequency is much higher than the electron cyclotron frequency.

To be specific, we consider waves with a polarisation $\mathbf{E} = E_1 \hat{y}$ so that $\mathbf{B} = B_0 \hat{z} - \frac{k_z}{\omega} E_1 \hat{i}$. This polarisation can be justified only when $\sigma_{yy} \gg \sigma_{xy}, \sigma_{zy}$, so that dispersion relation can be presented (by the Ampere law) as follows:

$$\omega^2 - c^2 k_z^2 + \frac{i\omega}{\varepsilon_0} \sigma_{yy} = 0. \tag{9}$$

Finally, the dispersion equations of the unstable modes have been derived in the spin-polarised [17] and spin-depolarised cases respectively, as follows:

$$\begin{aligned}
&\omega^2 - c^2 k_z^2 + \sum_{\nu=+,-} \frac{q_e^2}{4 m_e \epsilon_0} \\
&\quad \times \int \frac{(\omega - v_z k_z)}{(\omega - v_z k_z - \omega_{ce})} v_\perp \frac{\partial F_{0\nu}}{\partial v_\perp} \mathrm{d}^3 v \\
&\quad + \sum_{\nu=+,-} \frac{q_e^2}{4 m_e \epsilon_0} \int \frac{(\omega - v_z k_z)}{(\omega - v_z k_z + \omega_{ce})} v_\perp \frac{\partial F_{0\nu}}{\partial v_\perp} \mathrm{d}^3 v \\
&\quad + \sum_{\nu=+,-} \frac{q_e^2}{4 m_e \epsilon_0} \int \frac{v_\perp^2 k_z}{(\omega - v_z k_z - \omega_{ce})} \frac{\partial F_{0\nu}}{\partial v_z} \mathrm{d}^3 v \\
&\quad + \sum_{\nu=+,-} \frac{q_e^2}{4 m_e \epsilon_0} \int \frac{v_\perp^2 k_z}{(\omega - v_z k_z + \omega_{ce})} \frac{\partial F_{0\nu}}{\partial v_z} \mathrm{d}^3 v \\
&\quad + \sum_{\nu=+,-} \frac{\mu_e^2}{2 m_e \epsilon_0} \int \frac{k_z^3}{(\omega - v_z k_z - \omega_{cg})} \frac{\partial F_{0\nu}}{\partial v_z} \mathrm{d}^3 v \\
&\quad + \sum_{\nu=+,-} \frac{\mu_e^2}{2 m_e \epsilon_0} \int \frac{k_z^3}{(\omega - v_z k_z + \omega_{cg})} \frac{\partial F_{0\nu}}{\partial v_z} \mathrm{d}^3 v \\
&\quad - \sum_{\nu=+,-} \frac{\mu_e^2}{\hbar \epsilon_0} \int \frac{k_z^2}{(\omega - v_z k_z - \omega_{cg})} \nu F_{0\nu} \mathrm{d}^3 v \\
&\quad + \sum_{\nu=+,-} \frac{\mu_e^2}{\hbar \epsilon_0} \int \frac{k_z^2}{(\omega - v_z k_z + \omega_{cg})} \nu F_{0\nu} \mathrm{d}^3 v = 0
\end{aligned} \tag{10}$$

and

$$\begin{aligned}
&\omega^2 - c^2 k_z^2 + \pi \frac{q_e^2}{m_e \varepsilon_0} \frac{\sinh\alpha}{\alpha} \\
&\quad \times \int \frac{v_\perp (\omega - v_z k_z)}{(\omega - v_z k_z - \omega_{ce})} \frac{\partial f_0}{\partial v_\perp} \mathrm{d}^3 v \\
&\quad - \pi \frac{q_e^2}{m_e \varepsilon_0} \frac{\sinh\alpha}{\alpha} \int \frac{v_\perp v_z k_z}{(\omega - v_z k_z + \omega_{ce})} \frac{\partial f_0}{\partial v_\perp} \mathrm{d}^3 v \\
&\quad + \pi \frac{q_e^2}{m_e \varepsilon_0} \frac{\sinh\alpha}{\alpha} \int \frac{v_\perp^2 k_z}{(\omega - v_z k_z - \omega_{ce})} \frac{\partial f_0}{\partial v_z} \mathrm{d}^3 v \\
&\quad + \pi \frac{q_e^2}{m_e \varepsilon_0} \frac{\sinh\alpha}{\alpha} \int \frac{v_\perp^2 k_z}{(\omega - v_z k_z + \omega_{ce})} \frac{\partial f_0}{\partial v_z} \mathrm{d}^3 v \\
&\quad + \pi \frac{q_e^2 \omega}{m_e \varepsilon_0} \frac{\sinh\alpha}{\alpha} \int \frac{v_\perp}{(\omega - v_z k_z + \omega_{ce})} \frac{\partial f_0}{\partial v_\perp} \mathrm{d}^3 v \\
&\quad - 2\pi \frac{\mu_e^2}{m_e \alpha^2 \varepsilon_0} \left(\frac{\sinh\alpha}{\alpha} - \cosh\alpha \right) \\
&\quad \times \int \frac{k_z^3}{(\omega - v_z k_z - \omega_{cg})} \frac{\partial f_0}{\partial v_z} \mathrm{d}^3 v \\
&\quad - 2\pi \frac{\mu_e^2}{m_e \alpha^2 \varepsilon_0} \left(\frac{\sinh\alpha}{\alpha} - \cosh\alpha \right) \\
&\quad \times \int \frac{k_z^3}{(\omega - v_z k_z + \omega_{cg})} \frac{\partial f_0}{\partial v_z} \mathrm{d}^3 v \\
&\quad - 4\pi \frac{\mu_e^2}{\hbar \alpha \varepsilon_0} \left(\frac{\sinh\alpha}{\alpha} - \cosh\alpha \right) \\
&\quad \times \int \frac{k_z^2}{(\omega - v_z k_z - \omega_{cg})} f_0 \mathrm{d}^3 v
\end{aligned}$$

$$+ 4\pi \frac{\mu_e^2}{\hbar\alpha\varepsilon_0}\left(\frac{\sinh\alpha}{\alpha} - \cosh\alpha\right) \times \int \frac{k_z^2}{(\omega - v_z k_z + \omega_{cg})} f_0 \mathrm{d}^3 v = 0. \tag{11}$$

The integration $\mathrm{d}\Omega = \mathrm{d}^2 s \mathrm{d}^3 v$ includes ordinary spherical coordinates in the spin space by two degrees of freedom, and the cylindrical coordinates in the velocity space by three degrees of freedom.

The dispersion equation of the unstable modes can be rewritten in the spin-polarised [17] and spin-depolarised cases respectively as

$$\begin{aligned}
&\omega^2 - c^2 k_z^2 - \omega_{pe}^2 \\
&+ \frac{g^2\hbar^2 k_z^2}{32 m_e K_B T_z}\omega_{pe}^2 \left[2 + \zeta_1 Z(\zeta_1) + \zeta_2 Z(\zeta_2)\right] \\
&+ \frac{g^2\hbar k_z}{16\sqrt{2 m_e K_B T_z}}\omega_{pe}^2 \tanh\left(\frac{\mu_e B_0}{K_B T_{sp}}\right) \\
&\times \left[Z(\zeta_1) - Z(\zeta_2)\right] \\
&+ \frac{\sqrt{m_e}}{2\sqrt{2K_B T_z}}\frac{\omega_{ce}}{k_z}\omega_{pe}^2 \left[Z(\zeta_3) - Z(\zeta_4)\right] \\
&+ \frac{T_\perp}{2T_z}\omega_{pe}^2 \left[2 + \zeta_3 Z(\zeta_3) + \zeta_4 Z(\zeta_4)\right] = 0
\end{aligned} \tag{12}$$

and

$$\begin{aligned}
&\omega^2 - c^2 k_z^2 - \omega_{pe}^2 + \frac{T_\perp}{T_z}\omega_{pe}^2 \\
&+ \frac{g^2}{16}\frac{\hbar^2 k_z^2}{m_e K_B T_z}\omega_{pe}^2 \left(\frac{\coth\alpha}{\alpha} - \frac{1}{\alpha^2}\right) \\
&- \frac{g^2}{32}\frac{\hbar^2 k_z^2}{m_e \alpha^2 K_B T_z}\omega_{pe}^2 \left[\zeta_1 Z(\zeta_1) + \zeta_2 Z(\zeta_2)\right] \\
&+ \frac{g^2}{32}\frac{\hbar^2 k_z^2}{m_e K_B T_z}\frac{\coth\alpha}{\alpha}\omega_{pe}^2 \left[\zeta_1 Z(\zeta_1) + \zeta_2 Z(\zeta_2)\right] \\
&+ \frac{g^2}{16}\frac{\hbar k_z}{\sqrt{2 m_e K_B T_z}}\omega_{pe}^2 \coth\alpha \left[Z(\zeta_2) - Z(\zeta_1)\right] \\
&+ \frac{g^2}{16}\frac{\hbar k_z}{\alpha\sqrt{2 m_e K_B T_z}}\omega_{pe}^2 \left[Z(\zeta_1) - Z(\zeta_2)\right] \\
&+ \frac{\sqrt{m_e}\omega_{ce}}{2\sqrt{2K_B T_z} k_z}\omega_{pe}^2 \left[Z(\zeta_3) - Z(\zeta_4)\right] \\
&+ \frac{T_\perp}{T_z}\frac{\omega_{pe}^2}{2}\left[\zeta_3 Z(\zeta_3) + \zeta_4 Z(\zeta_4)\right] = 0.
\end{aligned} \tag{13}$$

Here, the function $Z\,(M = \zeta_1, \zeta_2, \zeta_3, \zeta_4)$ is the well-known dispersion function of the plasma with arguments

$$\zeta_1 = \frac{(\omega - \omega_{cg})}{k_z}\sqrt{\frac{m_e}{2K_B T_z}},$$

$$\zeta_2 = \frac{(\omega + \omega_{cg})}{k_z}\sqrt{\frac{m_e}{2K_B T_z}},$$

$$\zeta_3 = \frac{(\omega - \omega_{ce})}{k_z}\sqrt{\frac{m_e}{2K_B T_z}}$$

and

$$\zeta_4 = \frac{(\omega + \omega_{ce})}{k_z}\sqrt{\frac{m_e}{2K_B T_z}}.$$

Let us restrict the article to subliminal non-resonant instabilities because it is necessary for the arguments of the plasma dispersion function to be much smaller than unity so that the thermal energy can be larger than the magnetic field energy. In this situation, the plasma dispersion function can be approximated as $Z(M) = -2M + \cdots + i\pi$. In addition, the wave frequency, ω, is the sum of the real parts defining the wave fluctuations and the imaginary parts introducing the growth rate of the instability in the complex plane. Therefore, the instability growth rate and the fluctuation frequency have been derived by separating the wave frequency and applying the approximated dispersion function of the plasma and ignoring the second and higher order of the plasma dispersion arguments;

$$\begin{aligned}
\omega_{\text{i-pol}} = {} & \frac{v_{th,z} k_z}{\sqrt{\pi}}\left[1 - \frac{(1 + (c^2 k_z^2/\omega_{pe}^2))}{T_\perp/T_z}\right] \\
& \times \left[1 + \frac{g^2\hbar^2 k_z^2}{16 m_e K_B T_\perp}\right]^{-1} \\
& + \frac{v_{th,z} k_z}{\sqrt{\pi}}\frac{g^2\hbar^2 k_z^2}{16 m_e K_B T_\perp} \\
& \times \left[1 + \frac{2m_e}{\hbar k_z^2}\omega_{cg}\tanh\left(\frac{\mu_e B_0}{K_B T_{sp}}\right)\right] \\
& \times \left[1 + \frac{g^2\hbar^2 k_z^2}{16 m_e K_B T_\perp}\right]^{-1}
\end{aligned} \tag{14}$$

and

$$\begin{aligned}
\omega_{\text{i-depol}} = {} & \sqrt{\frac{2K_B T_z}{\pi m_e}} k_z \left[1 - \frac{(1 + (c^2 k_z^2/\omega_{pe}^2))}{T_\perp/T_z}\right. \\
& + \frac{g^2}{16}\frac{\hbar^2 k_z^2}{m_e K_B T_\perp}\left(\frac{\coth\alpha}{\alpha} - \frac{1}{\alpha^2}\right) \\
& \left. - \frac{g^2}{8}\frac{\hbar}{K_B T_\perp}\alpha\omega_{cg}\left(\frac{\coth\alpha}{\alpha} - \frac{1}{\alpha^2}\right)\right] \\
& \times \left[1 + \frac{g^2}{16}\frac{\hbar^2 k_z^2}{m_e K_B T_\perp}\left(\frac{\coth\alpha}{\alpha} - \frac{1}{\alpha^2}\right)\right]^{-1}.
\end{aligned} \tag{15}$$

The real (fluctuation) frequency is equal to zero for two cases. For non-resonant instability, it is necessary for the instability growth rate to be larger than the fluctuation.

Therefore, the instability will be restricted to wave number satisfying the following relations (the spin-polarised and spin-depolarised system, respectively):

$$\frac{T_\perp}{T_z}\left(1+\frac{g^2\hbar}{8K_BT_\perp}\omega_{cg}\tanh\left(\frac{\mu_eB_0}{K_BT_{sp}}\right)\right) \gg \left[1+\frac{c^2k_z^2}{\omega_{pe}^2}-\frac{g^2\hbar^2k_z^2}{16m_eK_BT_z}\right] \tag{16}$$

and

$$\frac{T_\perp}{T_z}\left[1-\frac{g^2}{8}\frac{\hbar}{K_BT_\perp}\alpha\omega_{cg}\left(\frac{\coth\alpha}{\alpha}-\frac{1}{\alpha^2}\right)\right] \gg \left[1+\frac{c^2k_z^2}{\omega_{pe}^2}-\frac{g^2}{16}\frac{h^2k_z^2}{m_eK_BT_z}\left(\frac{\coth\alpha}{\alpha}-\frac{1}{\alpha^2}\right)\right]. \tag{17}$$

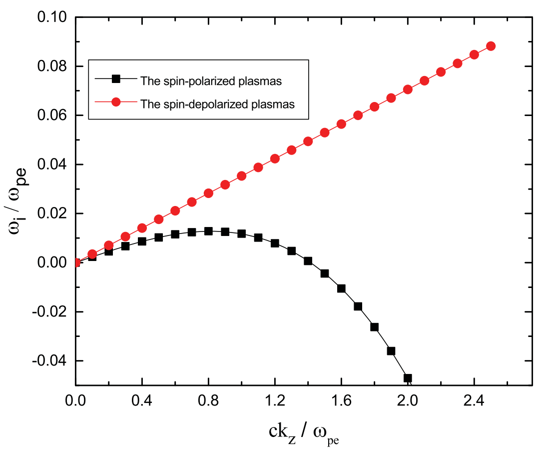

Figure 1. The normalised growth rate of the instability, ω_i/ω_{pe} as a function of wave number, ck_z/ω_{pe}, for the spin-polarised and spin-depolarised plasmas in the fixed $T_z = 1000$ eV, $n_0 = 10^{30}\,\mathrm{m}^{-3}$, $T_\perp/T_z = 3$, $T_{sp}/T_z = 2$ and $B_0 = 1$ T.

Both parts of eqs (14) and (15) include the classical and quantum parts so that the known classical results can be obtained in the absence of the spin effects.

3. Discussion and conclusion

The polarisation effects of the particle spin will be an important and interesting issue for inertial confinement fusion where the spin–orbit interactions can lead to depolarisation of the spin-polarised fuel where its cross-section can be small compared to the cross-section of the nuclear fusion interactions. In this study, the obtained growth rate of the instability are simulated for the ICF plasma conditions in two different cases of the spin-polarised and spin-depolarised plasmas. Results for the spin-polarised cases show that, the growth rate of the instability (normalised to the plasma frequency) first increases exponentially at the light frequency smaller than the plasma frequency up to when the light frequency reaches the plasma frequency. After that, it decreases exponentially for the light frequency higher than the plasma frequency and goes to negative value (reaching damping state). This situation is different from the spin-depolarised cases. For the spin-depolarised cases, it is observed that the normalised growth rate of the instability has a linear relation with the normalised light frequency. In fact, in this case, the instability grows continuously without any decrease in the values of the normalised growth rate while there is no restriction for the light frequency normalised to the plasma frequency (see figure 1). Indeed, spin-depolarised electrons can provide the free energy required by the instability so that, the instability can be present and grow even when the initial temperature anisotropy is stopped. Figure 2 illustrates conditions governing the growth of the instability for the spin-polarised and spin-depolarised electrons.

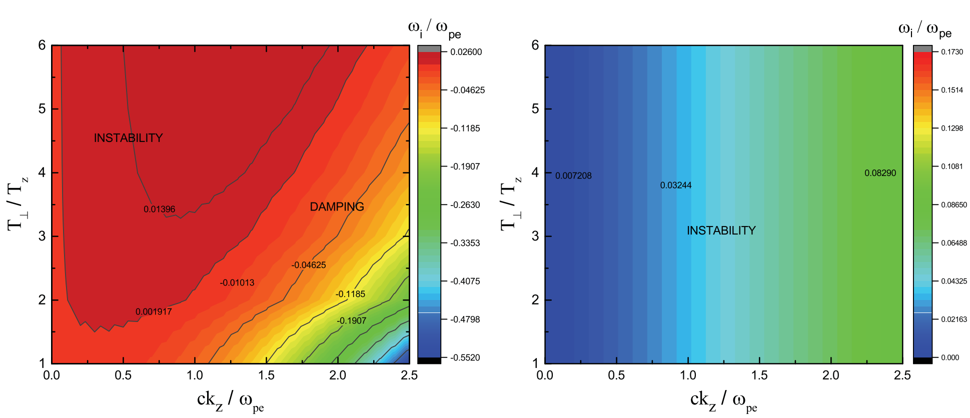

Figure 2. The domain of temperature anisotropy instabilities for different values of $T_\perp/T_z$, in the fixed $B_0 = 1$ T, $T_z = 1000$ eV, $n_0 = 10^{30}\,\mathrm{m}^{-3}$ and $T_{sp}/T_z = 2$ in the spin-polarised (left-hand side) and spin-depolarised (right-hand side) plasmas.

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