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Firehose instability in space plasmas with bi-kappa distributions

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ABSTRACT

Context. The existence of suprathermal charged particle populations in space plasma is frequently confirmed by interplanetary missions. In general, the velocity distribution functions are anisotropic, field aligned (gyrotropic) with two temperatures, parallel ($T_\parallel$) and perpendicular ($T_\perp$) to the ambient magnetic field $\boldsymbol{B}_0$.

Aims. Here, the dispersion properties of the firehose instability, which relaxes an anisotropic electron distribution function ($T_\parallel > T_\perp$) of bi-*kappa* type, are investigated for the first time.

Methods. The Solar wind is generally accepted to be a collisionless plasma and, therefore, the dispersion formalism is constructed on the basis of the kinetic Vlasov-Maxwell equations. The general dispersion relations are derived in terms of the modified plasma dispersion function.

Results. Simple analytical forms are obtained for the dispersion relation of the firehose instability and the instability criterion is derived. The exact numerical evaluation shows a significant departure of the dispersion curves from those obtained for a bi-Maxwellian plasma.

Conclusions. While the maximum growth rate is slightly diminished, the instability extends to large wave-numbers in the presence of suprathermal particles. Thus, this instability is more likely to be found in space plasmas with an anisotropic distribution of bi-*kappa* type. If all other parameters are known, measuring the instability growth time enables the determination of the spectral index κ.

Key words. plasmas – instabilities – Sun: coronal mass ejections (CMEs) – Sun: flares – Sun: solar wind

1. Introduction

There is consistent observational evidence that nonthermal particle populations are ubiquitous in the near Earth space plasma (for an overview, see in Zouganelis et al. 2005, and the references therein). Suprathermal populations are described by the so-called κ distribution function (introduced by Vasyliunas 1968), i.e. a power law in particle speed, with high-energy tails, and which degenerates into a Maxwellian distribution function as the spectral index becomes very large ($\kappa \to \infty$). Moreover, such deviations from the Maxwellian distribution are expected to exist in any low-density plasma in our Universe, where binary collisions are sufficiently rare. Instead, the wave-particle interactions can be responsible for the energization of plasma particles and the occurrence of κ distributions. In an ambient magnetic field, plasma particles gain energy through the cyclotron resonance and transit time damping (Landau resonance) of the linear waves (Fisk 1976). Thus, the small (linear) amplitude MHD modes, Alfvén and fast magnetosonic waves are able to accelerate both the electrons and the protons in the solar flares (Miller 1991, 1997), and in the inner magnetosphere (Summers & Ma 2000), while the whistler mode can enhance the energy of the electrons in the Earth's foreshock (Ma & Summers 1998). In the presence of large amplitude waves and plasma turbulence the nonlinear mechanisms of dissipation (e.g., the nonlinear Landau damping) also can be responsible for the energization of plasma particles (Miller 1991, Leubner 2000, Shizgal 2007). However, the MHD turbulent cascades channel the energy to small scales where the quasi-linear interactions with particles are again possible (Miller 1991, 1997). The non-thermal features of the particle distribution functions can also result from heat flows and temperature anisotropies (Leubner & Viñas 1986). The use of the generalized *kappa*-type distribution functions enables a more realistic description of the turbulent fields and the acceleration of plasma particles in such environments (Summers & Thorne 1991; Maksimovic et al. 1997; Mace & Hellberg 2003; Lazar et al. 2008b).

In magnetized interplanetary plasmas, the particle distribution functions are in general anisotropic, i.e. depend on the direction in velocity space (Pilipp et al. 1990; Salem et al. 2003), and, therefore, unstable to the excitation of plasma instabilities (Paesold & Benz 1999; Fahr & Siewert 2007, for a review of the fundamental properties of anisotropic plasma instabilities and their applications in space plasma physics, we recommend the textbook of Gary 1993). Although considerable progress has been made, questions still exist about the origin of the particle velocity anisotropies and the relaxation mechanisms in space plasmas (Stverak et al. 2008). Recently it has been shown how the bulk energy of the solar wind ions is converted by developing pronounced pitchangle anisotropies when the ions pass over the solar wind MHD termination shock (Fahr & Siewert 2007). Such anisotropic distribution functions are mirror-mode unstable and up to 60 percent of the thermal ion energy can be converted by this instability into magnetoacoustic turbulence energy. However, at large solar distances mechanisms of adiabatic cooling can be identified (e.g., the conservation of the two magnetic CGL invariants), which prevent the solar wind from developing strong pressure anisotropies (Fahr & Siewert 2008).

The preference for the acceleration of plasma particles along the background magnetic field is a common feature of the different acceleration models in solar flares and in the solar wind (see the Introduction of Paesold & Benz 1999, for a brief description of these mechanisms of acceleration). Thus, the velocity

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distribution function becomes anisotropic with a substantial amount of free energy residing in the direction of the magnetic field that may give rise to the excitation of the firehose instability.

The electron firehose mode propagates along the ambient magnetic field and the instability typically arises when the parallel temperature is much larger than the transverse temperature ($T_\parallel > T_\perp$). In this case, the approximate instability criterion of Hollweg & Völk (1970) is fulfilled:

$$1 - A_e\,\beta_{e,\parallel} < 0, \tag{1}$$

where $A_e = 1 - (T_{e,\perp}/T_{e,\parallel})^2$ is the electron temperature anisotropy, and $\beta_{e,\parallel} = 8\pi n k_B T_{e,\parallel}/B_0^2$ is the ratio of the thermal to the magnetic pressure. According to condition (1), an instability can occur only for positive but sub-unitary anisotropies, i.e. $0 < A_e < 1$, and consequently, for a parallel beta larger than unity.

Since the frequency of the firehose mode is in the range of the ion gyrofrequency, the electrons are non-resonant and the instability is non-resonantly driven by the electron temperature anisotropy (Hollweg & Völk 1970). The ions, however, are resonant and they can gain energy making the firehose instability responsible for the transfer of the electron energy to the ions. The first simulations of the electron firehose instability have also shown that the generated magnetic field fluctuations scatter the electrons, reducing their anisotropy (Gary & Nishimura 2003; Paesold & Benz 2003).

Presently, it is believed that the electron firehose instability presents two branches (Li & Habbal 2000): one with a finite real frequency, $\omega_r \neq 0$ and quasi-parallel propagation ($\theta = \widehat{\boldsymbol{k}, \boldsymbol{B}_0} < 30°$), and another one that is non-resonant, with $\omega_r = 0$ and quasi-perpendicular propagation ($\theta > 30°$). Both of these branches have recently been confirmed by two-dimensional Particle-In-Cell (PIC) simulations with a physical mass-ratio (Camporeale & Burges 2008). Here, we should remark that the second branch describing a purely-growing and quasi-perpendicular instability is a Weibel-like instability (Weibel 1959), as long as it is non-resonantly excited by a temperature anisotropy and has the wave-vector normal to the high temperature axis.

Potential applications of the kinetic electron firehose instability to solar flares or to the solar wind are discussed in Paesold & Benz (1999) and Li & Habbal (2000). The main idea is that this instability can constrain the increase of the electron temperature anisotropy, thus explaining the observations. Indeed, the anisotropic velocity distribution can be subjected to a pitchangle isotropisation by quasilinear interactions of electrons with plasma waves and other dissipative processes. The electron firehose instability itself can be the most efficient mechanism of temperature isotropisation in solar flares limiting electron anisotropy and thus providing the necessary condition for further acceleration (Messmer 2002).

In the present paper, we briefly describe the dispersion properties of the electron firehose instability driven by an anisotropic electron distribution function ($T_\parallel > T_\perp$) of the bi-*kappa* type, and propagating along the ambient magnetic field ($\theta = 0$). The plasma is assumed to be quiescent, without large amplitude oscillations and turbulence, so that we restrict this discussion to a quasilinear description of this instability.

2. Dispersion theory

2.1. Parallel transverse modes

For the transverse waves propagating *along* the ambient magnetic field, i.e., $\boldsymbol{k} \parallel \boldsymbol{B}_0$, (along the y-axis in the present model), we take the general dispersion relation from Tautz & Schlickeiser (2005)

$$1 - \frac{k^2c^2}{\omega^2} + \frac{\pi}{\omega^2}\sum_a \omega_{p,a}^2 \int_{-\infty}^{\infty} \frac{dv_\parallel}{\omega - kv_\parallel \pm \Omega_a} \times \int_0^\infty dv_\perp v_\perp^2 \left[(\omega - kv_\parallel)\frac{\partial F_{a,\kappa}}{\partial v_\perp} + kv_\perp \frac{\partial F_{a,\kappa}}{\partial v_\parallel}\right] = 0, \tag{2}$$

where ω and k denote the frequency and the wave number of the plasma modes, respectively, and c is the speed of light in a vacuum, $\Omega_a = q_a B_0/(m_a c)$ is the (non-relativistic) gyrofrequency, $\omega_{p,a} = (4\pi n_a e^2/m_a)^{1/2}$ is the plasma frequency for the particles of species a, and the $\pm$ are used to identify the circularly polarized electromagnetic waves with right-hand and left-hand polarizations, respectively.

2.2. Anisotropic bi-kappa distribution function

We consider a bi-kappa velocity distribution function of plasma particles

$$F_\kappa = \frac{1}{\pi^{3/2}\theta_\perp^2\theta_\parallel}\frac{\Gamma(\kappa+1)}{\kappa^{3/2}\Gamma(\kappa-1/2)}\left[1 + \frac{v_\parallel^2}{\kappa\theta_\parallel^2} + \frac{v_\perp^2}{\kappa\theta_\perp^2}\right]^{-\kappa-1} \tag{3}$$

that describes the initially unperturbed plasma system, and which is normalized by $\int d^3v\, F_\kappa = 1$. We use polar coordinates in the particle velocity space

$$(v_x, v_y, v_z) = (v_\perp \cos\phi, v_\perp \sin\phi, v_\parallel), \tag{4}$$

and the corresponding thermal velocities take different values as follows

$$\theta_\perp = \left(1 - \frac{3}{2\kappa}\right)^{1/2} v_{T_\perp}; \quad \theta_\parallel = \left(1 - \frac{3}{2\kappa}\right)^{1/2} v_{T_\parallel}, \tag{5}$$

where

$$v_{T_\perp}^2 = \frac{2k_BT_\perp}{m}; \quad v_{T_\parallel}^2 = \frac{2k_BT_\parallel}{m} \tag{6}$$

denote the perpendicular and parallel thermal velocities, respectively. Such an anisotropic distribution function, known as a bi-kappa distribution function, has been introduced by Summers and Thorne (1991), and it approaches the bi-Maxwellian distribution function as the spectral index reaches the limit $\kappa \to \infty$.

2.3. Firehose instability

The electron firehose instability arises when the plasma is hotter along the direction of the stationary magnetic field, i.e. when $\theta_\parallel > \theta_\perp$. This instability is left-handed (LH) circularly polarized at $\boldsymbol{k} \times \boldsymbol{B}_0 = 0$, and the electrons are non-resonant, $g_e^- \gg 1$. Only the ions are cyclotron resonant, but their contribution is minimized here assuming that, initially, the ion distribution function is isotropic. The ion firehose instability develops for similar conditions of the ion temperature anisotropy, $\theta_\parallel > \theta_\perp$, and saturates at maximum growth rates comparable to the electron firehose instability (viz. of the order of the ion gyrofrequency). However, because the ion response could be either inhibited or enhanced by the resonant dissipation of the electron firehose instability, the interplay of these two instabilities is more complex and it will be investigated in detail separately.

We now insert the distribution function (3) in Eq. (2) and derive the dispersion relation of the LH transverse modes

$$1 - \frac{k^2c^2}{\omega^2} + \sum_a \frac{\omega_{p,a}^2}{\omega^2}\left\{\frac{\omega}{k\theta_{a,\|}} Z_\kappa^0(g_a) - A_a\left[1 + g_a Z_\kappa^0(g_a)\right]\right\} = 0, \quad (7)$$

with

$$g_a = \frac{\omega - \Omega_a}{k\theta_{a,\|}}, \quad (8)$$

and where the temperature anisotropy is defined as

$$A = 1 - \frac{\theta_\perp^2}{\theta_\|^2} = 1 - \frac{T_\perp}{T_\|}\cdot \quad (9)$$

The new modified dispersion function

$$Z_\kappa^0(g) = \frac{1}{\pi^{1/2}\kappa^{1/2}} \frac{\Gamma(\kappa)}{\Gamma\left(\kappa - \frac{1}{2}\right)} \int_{-\infty}^{\infty} \mathrm{d}x \, \frac{(1 + x^2/\kappa)^{-\kappa}}{x - g}, \quad (10)$$

for $\Im(f) > 0$, usually describes the dispersion properties of the transverse waves and instabilities (Lazar et al. 2008a), and is related to the modified dispersion function Z_κ of Summers & Thorne (1991) by (Lazar et al. 2008a)

$$Z_\kappa^0(g) = \left(1 + \frac{g^2}{\kappa}\right) Z_\kappa(g) + \frac{g}{\kappa}\left(1 - \frac{1}{2\kappa}\right)\cdot \quad (11)$$

Both of these modified dispersion functions approach the standard plasma dispersion function Z of Fried & Conte (1961) in the limit of $\kappa \to \infty$, and Eq. (7) reduces exactly to the dispersion relation of parallel transverse waves in a plasma with bi-Maxwellian anisotropic distribution functions, see Eq. (7.1.6) from Gary (1993).

Consider a plasma with anisotropic electrons and isotropic ions (protons), $T_{e,\perp} = T_{i,\perp} = T_{i,\|}$. The dispersion relation (7) then simplifies to

$$1 - \frac{k^2c^2}{\omega^2} + \frac{\omega_{p,e}^2}{\omega^2}\left\{\frac{\omega}{k\theta_{e,\|}} Z_\kappa^0(g_e) - A_e\left[1 + g_e Z_\kappa^0(g_e)\right]\right\}$$
$$+ \frac{\omega_{p,i}^2}{\omega^2}\frac{\omega}{k\theta_{i,\|}} Z_\kappa^0(g_i) = 0. \quad (12)$$

For Maxwellian distributions, with the same assumption of anisotropic electrons and isotropic ions, the dispersion relation (12) transforms, in the limit of $\kappa \to \infty$, to

$$1 - \frac{k^2c^2}{\omega^2} + \frac{\omega_{p,e}^2}{\omega^2}\left\{\frac{\omega}{kv_{T_{e,\|}}} Z(f_e) - A_e\left[1 + f_e Z(f_e)\right]\right\}$$
$$+ \frac{\omega_{p,i}^2}{\omega^2}\frac{\omega}{kv_{T_{i,\|}}} Z(f_i) = 0, \quad (13)$$

where $Z(f)$ is the standard plasma dispersion function Z of Fried & Conte (1961) with the argument

$$f_a = \frac{\omega - \Omega_a}{kv_{T_{a,\|}}}\cdot \quad (14)$$

The electrons are non-resonant and, therefore, the arguments (8) or (14) can be considered sufficiently large, $g_e \gg 1$. This allows us to use the power series of the modified plasma dispersion function from Summers & Thorne (1991), and to find, for example, for two values of $\kappa = 2, 4$,

$$Z_2(g \gg 1) \simeq -\frac{3}{4g}\left(1 + \frac{2}{3g^2}\right),$$

$$Z_4(g \gg 1) \simeq -\frac{7}{8g}\left(1 + \frac{4}{7g^2}\right)\cdot \quad (15)$$

Combining the last Eqs. (15) with (11), we obtain for both values of κ the same expression of

$$Z_{2,4}^0(g \gg 1) \simeq -\frac{1}{g}\left(1 + \frac{1}{2g^2}\right)\cdot \quad (16)$$

In the limit of large arguments, $g \gg 1$, the second-order approximation of the modified dispersion function Z_κ^0 does not depend explicitly on the spectral index κ, but only through the factor

$$\alpha_\kappa = \left(1 - \frac{3}{2\kappa}\right)^{1/2}, \quad \lim_{\kappa\to\infty} \alpha_\kappa = 1, \quad (17)$$

which relates θ to the thermal velocity, v_T, in Eq. (5). Thus, we can use Eq. (16) to simplify Eq. (12) and obtain

$$1 - \frac{k^2c^2}{\omega^2} - \frac{\omega_{p,e}^2}{\omega^2}\left[\frac{\omega}{\omega - \Omega_e}\left(1 + \frac{1}{2(g_e)^2}\right) - \frac{A_e}{2(g_e)^2}\right]$$
$$+ \frac{\omega_{p,i}^2}{\omega^2}\frac{\omega}{k\theta_i} Z_\kappa^0(g_i) = 0. \quad (18)$$

This is the dispersion relation of the firehose mode, which is unstable under certain conditions that we derive below.

2.4. Instability criterion

Here we follow the same formalism as Hollweg & Völk (1970) and generalize the firehose instability criterion for anisotropic particle velocity distributions of the bi-*kappa* type. We use some reasonable simplifications based on what we are looking for: the firehose instability has a frequency ($\omega = \omega_r + \imath\omega_i$) in the vicinity of the ion cyclotron frequency, and therefore $\omega_i < \omega_r \ll |\Omega_e|$, and $\omega^2 \ll k^2c^2$. Moreover, for simplicity, the protons can be considered Maxwellian (as sustained by many observations) and isotropic in the velocity space without any contribution to the instability. Besides that, the protons are resonant and we can assume a small argument $f_i < 1$ of the dispersion function (Maxwellian) that in this case approaches $Z(f < 1) \simeq \imath\pi^{1/2}\exp(-f^2) \simeq \imath\pi^{1/2}$. With these presumptions, the dispersion relation (18) becomes

$$\frac{k^2c^2}{\omega_{p,i}^2} + \frac{\omega_{p,e}^2}{\omega_{p,i}^2}\left[\frac{\omega}{|\Omega_e|} - \frac{A_e k^2\theta_{e,\|}^2}{2\Omega_e^2}\right] - \imath\pi^{1/2}\frac{\omega}{kv_{T_i}} = 0 \quad (19)$$

and for the real and imaginary part of the frequency scaled to the ion cyclotron frequency, we derive

$$\frac{\omega_r}{\Omega_i} = -\frac{k^2c^2}{\omega_{p,i}^2}\left(1 - \frac{A_e}{2}\alpha_\kappa\beta_{e,\|}\right)\left(1 + \pi\frac{\Omega_i^2}{k^2v_{T_i}^2}\right)^{-1} \quad (20)$$

and

$$\frac{\omega_i}{\Omega_i} = -\pi^{1/2}\frac{k^2c^2}{\omega_{p,i}^2}\left(1-\frac{A_e}{2}\alpha_\kappa\beta_{e,\parallel}\right)\left(\frac{kv_{T_i}}{\Omega_i}+\pi\frac{\Omega_i}{kv_{T_i}}\right)^{-1}, \quad (21)$$

respectively. The firehose instability will occur for

$$1-\frac{A_e}{2}\alpha_\kappa\beta_{e,\parallel}<0, \quad (22)$$

and for a given anisotropy, condition (22) is satisfied only for spectral indexes larger than the threshold

$$\kappa>\kappa_t=\frac{3}{2}\left(1-\frac{2}{A_e\beta_{e,\parallel}}\right)^{-1}. \quad (23)$$

If the anisotropy chosen here (in accordance with the earlier estimates of Paesold & Benz 1999) is $T_{e,\parallel}/T_{e,\perp}=20$, we find $\kappa>\kappa_t\simeq 3.93$, and indeed, in Fig. 1b, the instability exists only for $\kappa\geqslant 4$. For high values of the parallel beta, the spectral index is sufficiently large $\kappa>\kappa_t>3/2$, so that we keep away from the critical value $\kappa_c=3/2$, where the distribution function (3) collapses and the temperature (5) becomes undefined.

While Eq. (22) gives the instability criterion in a plasma with a bi-*kappa* temperature anisotropy, for an infinitely large $\kappa\to\infty$ we find the *exact* criterion of firehose instability for a bi-Maxwellian distribution

$$1-\frac{A_e}{2}\beta_{e,\parallel}<0, \quad (24)$$

which upgrades the approximative condition (1) of Hollweg & Völk (1970).

Keeping an ion distribution function of the kappa-type (as suggested by the recent works of Fisk & Gloeckler 2006, 2007), only the numerical factor $\pi^{1/2}\simeq 1.77$ in the last term of the dispersion relation (19) will increase slightly, for example, to 2.02 and 2.82 for $\kappa=4$ and $\kappa=2$, respectively. But this term has no influence on the instability criterion in Eq. (22). This is confirmed in the next section by comparing the exact numerical solutions.

3. Numerical solutions

In the previous section, we derived the analytical (approximative) forms of the frequency and the growth rate. Here, however, to be rigorous, we present the exact numerical solutions of the firehose instabilities. If the protons are Maxwellian, Eq. (18) transforms in terms of the standard dispersion function (Fried & Conte 1961) into

$$1-\frac{k^2c^2}{\omega^2}-\frac{\omega_{p,e}^2}{\omega^2}\left[\frac{\omega}{\omega-\Omega_e}\left(1+\frac{1}{2(g_e)^2}\right)-\frac{A_e}{2(g_e)^2}\right] + \frac{\omega_{p,i}^2}{\omega^2}\frac{\omega}{kv_{T,i}}Z(f_i)=0. \quad (25)$$

The real and imaginary solutions of Eq. (25) are shown with solid lines in Fig. 1. We consider three different cases of bi-kappa distributions of electrons, corresponding to $\kappa=4,5,6$. For comparison, we indicated with dashed lines the reproduced numerical solutions from Paesold & Benz (1999) obtained for the same values of plasma parameters (typical for solar flares), but for a bi-Maxwellian distribution function ($\kappa\to\infty$) of the electrons.

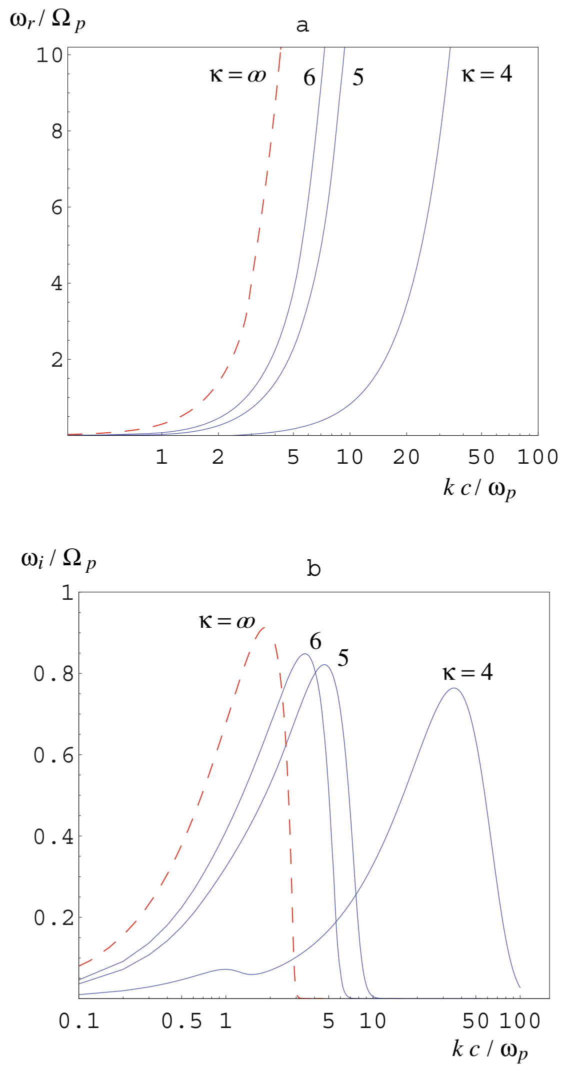

Fig. 1. The real and imaginary solutions of Eq. (25) are plotted with solid lines for three values of $\kappa=4,5,6$, and with dashed lines for an infinitely large $\kappa\to\infty$. The electron-proton plasma parameters are $n_e=5\times10^{10}$ cm^{-3}, $T_{e,\perp}=T_{p,\perp}=T_{p,\parallel}=10^7$ K, $T_{e,\parallel}=T_{e,\perp}=20$, $B_0=100$ G.

In this case, the maximum growth rates cannot be obtained by simply imposing $d\omega_i/dk=0$ in Eq. (21), or Eq. (25), because the instability is saturated by two mechanisms at the same time, either by scattering electrons or by heating ions (Messmer 2002; Paesold & Benz 2003). However, the growth rates shown in Fig. 1 are only slightly diminished by the non-Maxwellian distributions of the electrons. Contrary to this, the range of instability is extended to larger wave-numbers. Again, the so-called cutoff wave number, k_c, which limits the existence of the firehose instability to wave numbers $k<k_c$, cannot be found by simply imposing condition $\omega_i=0$ in Eq. (25), because the growth rate decreases drastically but it does not vanish completely at $k=k_c$. Therefore, here, we only suppose that k_c can be derived from the resonance condition for the ions, $f_i\sim 1$, in the dispersion Eq. (25), and we leave the exact calculation of the maximum growth rates and the cutoff wave numbers for a detailed analysis elsewhere.

In Fig. 2, with dotted and solid lines, the dispersion curves (a) and the growth rates (b) given by Eqs. (18) and (25), respectively, are compared. We have assumed that the electrons are described by a bi-$\kappa=4$ anisotropic distribution function, and

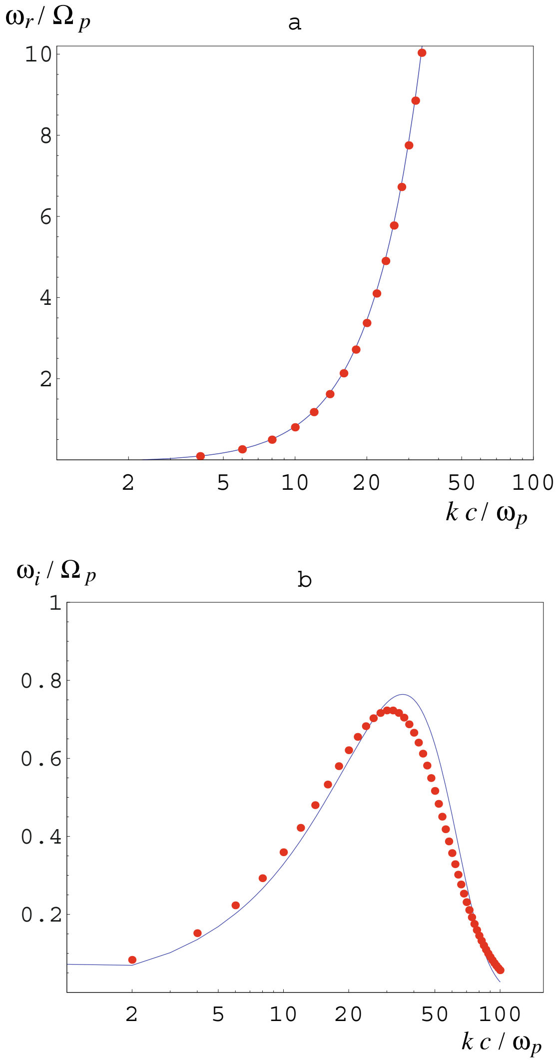

Fig. 2. Comparison of numerical solutions given by the Eqs. (18) (dotted lines) and (25) (solid lines) for an index $\kappa = 4$. The plasma parameters are the same as in Fig. 1.

in Eq. (18), the ions are described by an isotropic $\kappa = 4$ distribution function. A small deviation is observed only for the growth rates, which are slightly decreased by the non-Maxwellian ions.

The dispersion relations provided above, for example Eqs. (18) and (25), allow for a more extensive analysis of the dispersion properties and the stability of different plasma systems if we combine electron and ion distribution functions with diverse spectral indexes as given by the observations.

4. Conclusions

1. We have investigated the dispersion properties of the firehose instability driven by an anisotropic bi-*kappa* type electron distribution function ($T_\parallel > T_\perp$).
2. The general integral dispersion relations have been derived in terms of the modified plasma dispersion function.
3. Simple analytical forms of the dispersion relations are obtained and the firehose instability criterion is provided.
4. The exact numerical evaluation of the dispersion curves and the growth rates have proved a strong dependence on the spectral index κ, and a significant departure from the unstable solutions of a bi-Maxwellian plasma.
5. While the maximum growth rate is slightly diminished, the instability extends to large wave-numbers. In the presence of suprathermal particles, the wave-number corresponding to the maximum growth rate as well as the cutoff wave-number increase considerably.
6. This instability is thus more likely to be found in space plasmas with an anisotropic distribution of the bi-*kappa* type.
7. If all other parameters are known, measuring the instability growth time makes it possible to determine the spectral index κ.

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