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This paper is the first approach for analyzing the influence of κ-distributed particles on the modified two-stream instability (MTSI). It is assumed that the plasma consists of a magnetized Maxwellian electron contribution and unmagnetized κ-distributed ions drifting across the electrons. Within an electrostatic approximation, the influence of the κ parameter on the maximum growth rate of the MTSI is evaluated for the special case of parallel drift velocity and wave propagation. © *2005 American Institute of Physics*. [DOI: 10.1063/1.2065370]

I. INTRODUCTION

There are many occasions in space plasmas where ion and electron contributions have a relative velocity to each other. Such a situation in a plasma leads to the excitation of various instabilities. A complete and systematic picture based on an asymptotic study of the general dispersion relation is presented in Ref. 1. Among the many instabilities that can be driven by a current flowing in a plasma, this paper is devoted to the so-called modified two-stream instability (MTSI). The stability of a plasma with respect to the MTSI is determined by various parameters, including the amount and direction of the relative velocity, the ratio of the temperatures of the different plasma populations, the temperature anisotropy of the respective particle species, as well as the plasma β, i.e., the ratio of magnetic to thermal pressures in the ambient plasma.

Considerable theoretical effort has been given to analyzing the particular features of the MTSI. A complete treatise of the electrostatic approximation in the linear regime can be found in Ref. 2, where all main features appearing due to a relative motion between electrons and ions are given in a very comprehensive way. The first full electromagnetic approach is presented in Refs. 3 and 4, where the relevant dispersion relation has been derived and solved numerically. The authors show that, in general, the electromagnetic effects reduce the growth rate of the MTSI. Nonlinear investigations revealed that the MTSI is responsible for considerable plasma heating.[5–7] Recently, performing shock simulations with a one-dimensional full particle code, the appearance of a MTSI at the foot of a quasiperpendicular supercritical shock has been shown in Refs. 8 and 9. An important limiting case associated with the MTSI is the so-called ion-Weibel instability (see Refs. 10–12). This instability belongs to the same dispersion curve as the MTSI and occurs due to the electromagnetic response of the ion contribution.

We note that the studies mentioned above mainly focus on the MTSI as an application to the terrestrial bow shock and/or the current sheet in the geomagnetic tail. However, there are more interesting application fields of the MTSI, e.g., the flow of the solar wind around unmagnetized bodies. Examples of such studies are given in Refs. 13 and 14 where the MTSI in application to Mars and Venus has been considered, providing theoretical explanation for the observed electromagnetic activity in the vicinity of an unmagnetized planet.

All of the studies mentioned above considered the case when the particle contributions are distributed according to a Maxwellian distribution function. In general, however, measured distribution functions may deviate considerably from a Maxwellian one. First examples of such distribution functions were given in Ref. 15, where a family of power law distribution was introduced in order to model the suprathermal particles as observed by the Orbiting Geophysical Observatory 1 and 2. From then on, numerous *in situ* observations revealed the presence of nonthermal plasma contributions in a variety of astrophysical plasma environments (see Ref. 16 and references therein). Motivated by the experimental suggestion of the validity of the family of κ distributions, considerable theoretical effort has been made to adapt the existing kinetic theory of plasma waves to this particular distribution function. In analogy to the well-known plasma dispersion function for a Maxwellian plasma, which proved to be a useful tool in the analysis of wave propagation and microinstabilities, Summers and Thorne[17] were the first to introduce the so-called modified plasma dispersion function (MPDF) based on a κ-distribution function for integer κ. This result was generalized to a real κ in Ref. 18, where the proportionality of the MPDF to Gauss' hypergeometric function was also revealed. The dispersion equation for electrostatic waves involving the MPDF has been successfully solved in a magnetized plasma with a κ distribution in Ref. 19 for calculating the electrostatic Bernstein modes. The κ distribution has also been applied to the theoretical analysis of the mirror instability,[20–22] where it has been shown that such a distribution effectively lowers the instability threshold. In addition to the adoption of the κ distribution to the existing theory of plasma waves and instabilities, there exist several approaches providing a theoretical explanation of the occurrence of power law distributions itself (see Refs. 23–26).

In our study, we apply a classical approach for the description of the MTSI within the framework of linear theory. Within this approach, a coordinate system is employed in

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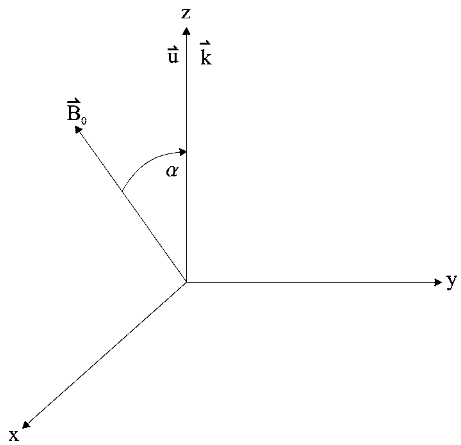

FIG. 1. Sketch of the applied coordinate system.

which the electrons are assumed to be at rest and the ions have a relative motion $\mathbf{u}$. Since the unstable waves considered have frequencies well above the ion gyrofrequency, the ions are assumed to be unmagnetized, in contrast to the electrons that are strongly magnetized. In order to keep the analysis as transparent as possible, we impose three additional assumptions. We restrict ourselves to the consideration of the electrostatic problem. This assumption simplifies the analysis considerably as in the case of a potential electric field it is sufficient to consider the dielectric response function instead of the whole susceptibility tensor. In addition, we assume the electrons to be Maxwellian, whereas the ions are taken to obey a κ-distribution function. Finally, we assume the wave vector $\mathbf{k}$ to be parallel to the flow velocity of the ions.

The organization of the paper is as follows. In Sec. II we focus on the applied coordinate system and examine the explicit expression for the susceptibility of the electron contribution. Section III discusses a short derivation of the susceptibility for the unmagnetized κ-distributed ions. In Sec. IV we present the numerical solution of the dispersion relation. The paper ends with a discussion of the obtained results as well as the outlook for future studies.

II. DIELECTRIC RESPONSE FUNCTION OF THE ELECTRONS

A coordinate system as sketched in Fig. 1 is adopted in the current study. We assume the wave vector $\mathbf{k}$ as well as the flow direction of the ions to lie along the z axis. The magnetic field, $\mathbf{B_0}$, lies in the xz plane, inclined towards the z axis under the angle α. Thus, we have

$$\mathbf{k}=\begin{pmatrix}0\\0\\k_z,\end{pmatrix},\quad \mathbf{u}=\begin{pmatrix}0\\0\\u_z,\end{pmatrix},\quad \mathbf{B}_0=\begin{pmatrix}B_0\sin\alpha\\0\\B_0\cos\alpha\end{pmatrix}. \tag{2.1}$$

The choice of this particular coordinate system is due to the fact that it keeps the analysis for the ion dielectric response function as simple as possible since there is just one component of the wave vector present. This simplification goes along with a more complex form of the electron susceptibility.

As a next step, we turn our attention towards the derivation of the susceptibility for the electron contribution. For convenience, we neglect subscript e in this section as all physical quantities refer to the electrons only. We follow the analysis performed in Refs. 27 and 28 and adopted to our particularly chosen coordinate system. As we consider electrostatic perturbations only, the following solution for the disturbed distribution function, f_1, can be related to the zero-order distribution function, f_0, via the Vlasov equation:[27]

$$f_1(\mathbf{r},\mathbf{v},t)=\frac{q}{m}\int_{-\infty}^{t}\mathbf{E}_1(\mathbf{r}',t')\cdot\nabla_{v'}f_0(\mathbf{v}')dt', \tag{2.2}$$

where q and m refer to charge and mass of the particle, respectively, and $\mathbf{E}$ refers to the electric field. The integration is meant to be performed along the particle trajectories. The wave electric field is assumed to be

$$\mathbf{E}_1=\mathbf{E}e^{i(\mathbf{k}\cdot\mathbf{r}'-\omega t)}, \tag{2.3}$$

where ω is the frequency. The trajectory that reaches $\mathbf{v}'=\mathbf{v}$ at $t'=t$ is governed by the equation

$$\frac{d\mathbf{v}'}{dt}=\frac{q}{m}\mathbf{v}'\times\mathbf{B}_0. \tag{2.4}$$

Solving this equation allows us to get for the particle velocities

$$v_x'=v_x(\cos^2\alpha\cos\Omega\tau+\sin^2\alpha)-v_y\cos\alpha\sin\Omega\tau + v_z\cos\alpha\sin\alpha(1-\cos\Omega\tau), \tag{2.5}$$

$$v_y'=v_y\cos\Omega\tau+(v_x\cos\alpha-v_z\sin\alpha)\sin\Omega\tau, \tag{2.6}$$

$$v_z'=v_1x\cos\alpha\sin\alpha(1-\cos\Omega\tau)+v_y\sin\alpha\sin\Omega\tau + v_z(\cos^2\alpha+\sin^2\alpha\cos\Omega\tau), \tag{2.7}$$

with Ω being the electron cyclotron frequency and $\tau=t-t'$. For the spatial coordinates we get

$$x'=x+\frac{v_y}{\Omega}\cos\alpha(1-\cos\Omega\tau)-\tau\sin\alpha(v_z\cos\alpha + v_x\sin\alpha)-\frac{\sin\Omega\tau}{\Omega}\cos\alpha(v_x\cos\alpha-v_z\sin\alpha), \tag{2.8}$$

$$y'=y-\frac{v_y}{\Omega}\sin\Omega\tau+\frac{v_x\cos\alpha-v_z\sin\alpha}{\Omega}(\cos\Omega\tau-1), \tag{2.9}$$

$$z'=z-\frac{v_y}{\Omega}\sin\alpha(1-\cos\Omega\tau)-\frac{1}{2}\tau(v_z+v_z\cos 2\alpha + v_x\sin 2\alpha)+\frac{\sin\alpha}{\Omega}\sin\Omega\tau(v_x\cos\alpha+v_z\sin\alpha). \tag{2.10}$$

Applying the traditional analysis for obtaining the susceptibility, as is described in many textbooks (e.g., Refs. 27 and

28), we arrive at the following expression for the electron susceptibility

$$\chi = \frac{2\omega_p^2}{k_z^2 v_t^2} \sum_{n=-\infty}^{\infty} e^{-\lambda} I_n(\lambda) \left((1 + \eta Z[\eta]) - \frac{n}{\cos\alpha} \frac{\Omega}{k_z v_t} Z[\eta] \right), \tag{2.11}$$

where I_n denotes the Bessel function of the first kind, and v_t and ω_p refer to thermal velocity and plasma frequency, respectively. Quantities λ and η are given as

$$\lambda = \frac{k_z^2 v_t^2}{4\Omega^2}(1 - \cos 2\alpha), \tag{2.12}$$

$$\eta = (\omega + n\Omega)\left(\frac{k_z^2 v_t^2}{2}(1 + \cos 2\alpha) \right)^{-1/2}. \tag{2.13}$$

In the cold limit, i.e., $v_t \to 0$, the susceptibility leads to the following asymptotic expression:

$$\chi = -\frac{\omega_p^2}{\omega^2 - \Omega^2} + \frac{\omega_p^2}{\omega^2} \frac{\Omega^2}{\omega^2 - \Omega^2} \cos^2\alpha. \tag{2.14}$$

III. DIELECTRIC RESPONSE FUNCTION OF THE IONS

As in the previous section, for convenience, here we drop subscript i, which denotes the particle species. For the detailed derivation of the dielectric response function, we refer to Refs. 17–19. Here, we briefly outline the main formalism of how the ion susceptibility is obtained. In general, the susceptibility in an unmagnetized plasma is given as[2]

$$\chi = \frac{4\pi e^2}{mk^2} \int \frac{1}{\omega - \mathbf{k} \cdot \mathbf{v}} \mathbf{k} \cdot \frac{\partial f_0}{\partial \mathbf{v}} d\mathbf{v}. \tag{3.1}$$

For our study, the zero-order distribution function is assumed to be of the following form:[19]

$$f_0(\mathbf{v}) = A_\kappa \left(1 + \frac{1}{\kappa v_\kappa^2}[v_x^2 + v_y^2 + (v_z - u_z)^2] \right)^{-\kappa-1}, \tag{3.2}$$

where κ is the spectral index, and

$$A_\kappa = \left(\frac{1}{\pi \kappa v_\kappa^2} \right)^{3/2} \frac{\Gamma[\kappa + 1]}{\Gamma[\kappa - 1/2]}, \quad \kappa > \frac{3}{2}, \tag{3.3}$$

$$v_\kappa^2 = 2 \frac{\kappa - \frac{3}{2}}{\kappa} \frac{T}{m}. \tag{3.4}$$

Here, Γ denotes the gamma function. Important features associated with the κ distribution are that it tends to a Maxwellian distribution for the limit $\kappa \to \infty$. A second important property of the distribution function as defined in Eq. (3.2) is that it reflects an inverse power for high velocities and thus allows one to model superthermal particle contributions within the plasma.

The main reason why we apply the distribution function in its current form is that it is consistent with those of Refs. 17–19, and thus allows a simple comparison of the resulting dielectric response function. However, there is a problem with this particular form of distribution function, which will be outlined in the following. Since $2T/m$ denotes the thermal spread of the Maxwellian, the expression v_κ^2 denotes the corresponding modified quantity for the κ distribution. For $\kappa < \infty$ the thermal spread of the κ function appears to be less than the one of the corresponding Maxwellian [see Eq. (3.4)] and thus is inconsistent in view of the suprathermal wings of the distribution generated for finite values of κ. In other words, the modified (effective) temperature decreases for decreasing κ values as compared to the Maxwellian temperature. In principle, this should be just the opposite. A distribution function resembling the described feature as well as a theoretical explanation for its occurrence is given in Refs. 16, 26, and 29. The application of such a form of distribution function is beyond the scope of the current paper, however, it will be addressed in a forthcoming work.

Substitution of Eq. (3.2) into (3.1) leads to the following expression determining the susceptibility for the considered ion contribution:

$$\chi = -\frac{4\pi e^2}{mk_z} \int \frac{1}{\omega - k_z v_z} \frac{2\mathcal{A}_\kappa}{v_\kappa^2} \frac{\kappa + 1}{\kappa}(v_z - u_z) \times \left(1 + \frac{1}{\kappa v_\kappa^2}[v_x^2 + v_y^2 + (v_z - u_z)^2] \right)^{-\kappa-2} d^3v. \tag{3.5}$$

Thus, there are three integrals that need to be performed. The integrals over v_x and v_y can be solved in analogy to[30]

$$\int_{-\infty}^{\infty} \frac{1}{(ax^2 + 2bx + c)^n} = \frac{\Gamma[n - 1/2]}{\Gamma[1/2]} \frac{\pi a^{n-1}}{(n-1)!(ac - b^2)^{n-1/2}}, \quad a > 0, \quad ac > b^2. \tag{3.6}$$

Adapting the above expression to our requirements leaves the following integral:

$$\chi = -\frac{4\pi e^2}{mk_z} 2\mathcal{A}_\kappa \pi \frac{\kappa + 1}{(\kappa v_\kappa^2)^{-\kappa-1}} \frac{\Gamma[\kappa + 1]}{\Gamma[\kappa + 2]} \int_{-\infty}^{\infty} \frac{1}{\omega - k_z v_z} \times (v_z - u_z)[\kappa v_\kappa^2 + (v_z - u_z)^2]^{-\kappa-1} dv_z. \tag{3.7}$$

The above integral can be solved via application of[2]

$$(\omega - k_z v_z)^{-1} = -i \int_0^\infty \exp[i\tau(\omega - k_z v_z)] d\tau, \tag{3.8}$$

and[30]

$$\int_0^\infty x(x^2 + \beta^2)^{\mu - 1/2} \sin(ax) dx = \sqrt{\pi}\beta \left(\frac{2\beta}{a} \right)^\mu \frac{1}{\Gamma[1/2 - \mu]} K_{\mu+1}(a\beta), \tag{3.9}$$

where

$$a > 0, \quad \mathrm{Re}[\beta] > 0, \quad \mathrm{Re}[\mu] < 0, \tag{3.10}$$

and K denotes the modified Bessel function according to

$$K_\mu = \frac{\pi}{2}\frac{I_{-\mu} - I_\mu}{\sin(\mu\pi)}. \tag{3.11}$$

With the help of the previously introduced relation, one can finally obtain the following expression for the susceptibility:

$$\chi = 2\frac{\omega_p^2}{k_z^2 v_\kappa^2}\left(\frac{\kappa - 1/2}{\kappa} + w' Z_\kappa[w']\right). \tag{3.12}$$

Here

$$w' = \frac{\omega - k_z u_z}{k_z v_\kappa}, \tag{3.13}$$

and Z_κ denotes the plasma dispersion for a Lorentzian distributed plasma, given as

$$Z_\kappa(w') = \frac{i}{2^{\kappa-1/2}\kappa^{3/2}\Gamma[\kappa - 1/2]} \times \int_0^\infty \exp\left[i\frac{w's}{\kappa^{1/2}}\right] s^{\kappa+1/2} K_{\kappa+1/2}[s] ds, \tag{3.14}$$

which has been obtained by Refs. 17–19, except that here the bulk flow of the ions is included. One convenient property of the above representation is, as already pointed out in Ref. 19, that it allows a simple comparison with the well-known results obtained by applying a Maxwell distribution.

As in the case of the electron susceptibility, our next step consists of considering the limiting case of vanishing temperature. For large arguments, the following Taylor series is valid for the MPDF (Ref. 18):

$$Z_\kappa(w') = -i\frac{\sqrt{\pi}}{\kappa^{3/2}}\frac{\Gamma(\kappa+1)}{\cos(\pi\kappa)\Gamma(\kappa - 1/2)}\left(\frac{w'}{i\sqrt{\kappa}} - 1\right)^{-(\kappa+1)} \times \left(\frac{w'}{i\sqrt{\kappa}} + 1\right)^{-(\kappa+1)} - \frac{2\kappa - 1}{2\kappa}\frac{1}{w'} \times \sum_{n=0}^{\infty} \kappa^n \frac{\sqrt{\pi}}{\Gamma(-n+1/2)}\frac{\Gamma(-\kappa+1/2)}{\Gamma(-\kappa+1/2+n)}\xi^{-2n}. \tag{3.15}$$

Applying the above expansion and taking the limit for $v_\kappa \to 0$ allows us to obtain the following familiar expression:

$$\chi = -\frac{\omega_p^2}{(\omega - k_z u_z)^2}. \tag{3.16}$$

IV. NUMERICAL SOLUTION

In the following, we assume the parameter λ to be much smaller than unity, which allows us to truncate the infinite sum occurring in the electron response function. In other words, we consider the effects due to the electron cyclotron harmonics $n=0$ and $n=\pm 1$. The contribution from the higher harmonics are neglected, allowing us to find for the dispersion equation

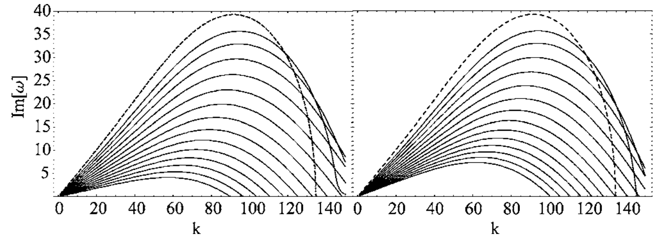

FIG. 2. Growth rate as a function of k for Maxwellian (left) and κ-distributed (right) ions in normalized units. The plasma β varies from 0.02 to 0.15 with a step size of 0.01. The dotted profiles correspond to the case of vanishing temperatures.

$$1 + \frac{2\omega_{pe}^2}{k_z^2 v_t^2}\sum_{n=-1}^{1} e^{-\lambda} I_n(\lambda)\left((1 + \eta Z[\eta]) - \frac{n}{\cos\alpha}\frac{\Omega_e}{k_z v_t} Z[\eta]\right) + 2\frac{\omega_{pi}^2}{k_z^2 v_\kappa^2}\left(\frac{\kappa - 1/2}{\kappa} + w' Z_\kappa[w']\right) = 0. \tag{4.1}$$

The occurring modified Bessel functions are approximated by using the following Taylor series:

$$e^{-\lambda} I_{-1}[\lambda] \approx \frac{\lambda}{2}, \quad e^{-\lambda} I_0[\lambda] \approx 1 - \lambda, \quad e^{-\lambda} I_1[\lambda] \approx \frac{\lambda}{2}. \tag{4.2}$$

For computional convenience, we introduce dimensionless quantities according to

$$\omega = \Omega_i \tilde{\omega}, \quad k = \frac{\Omega_i}{v_A}\tilde{k}, \quad \tau = \frac{\omega_{pe}^2}{\Omega_e^2}, \quad u = v_A \tilde{u}, \quad \mu = \frac{m_i}{m_e}. \tag{4.3}$$

Applying the above normalization leaves the dispersion equation as a function of several input parameters for which we take values[9]

$$\tilde{u} = 2, \quad \mu = 1836, \quad \tau = 2 \times 10^4, \quad \alpha = 85^\circ. \tag{4.4}$$

These values correspond to the terrestrial bow shock and remain fixed throughout the whole numerical analysis. In addition, the dispersion equation is a function of the electron plasma beta and the spectral index, which are the two key parameters of the present study.

As a first step in presenting numerical solutions to the dispersion equation, we refer to Fig. 2, where the growth rate for the MTSI is shown for a Maxwellian (left panel) and a κ distribution (right panel). Here, the growth rate is shown as a function of k and for various values of the plasma β, which is assumed to be equal for electrons and ions, and κ remains fixed at $\kappa=3$. The dotted line in this figure corresponds to the cold case, where the usual applied asymptotic expansion for small arguments has been used. In this figure, one can clearly see how the thermal effects within the plasma effectively reduce the growth rate of the MTSI, as has already been obtained by others (e.g., Ref. 3), for both applied distribution functions. However, from this figure we also find that the larger the plasma β, the more important the impact of the κ parameter. The latter becomes more transparent with regard to Fig. 3. Here the maximum growth rate, $\mathrm{Im}[\omega_M]$, is shown

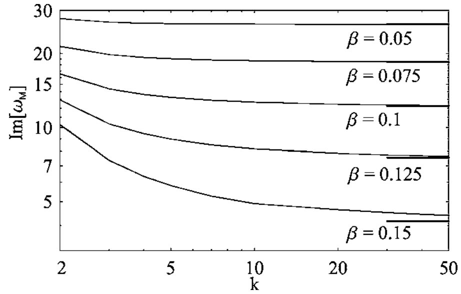

FIG. 3. Maximum growth rate as a function of κ for various values of the plasma β.

as a function of κ for various values of the plasma β. As a side note, we want to explain that what we call maximum growth rate refers to our particular assumed one-dimensional problem. In general, the maximum growth rate of the MTSI occurs at oblique angles and does not have to coincide with the direction of the bulk flow of the ions. The thick black horizontal lines in Fig. 3 indicate the solutions obtained by applying a Maxwell distribution. As a test of the numerical calculations we see that the solutions, obtained by applying the κ distribution function, tend toward the solutions obtained via a Maxwellian one for large values of κ. In addition, we note that for a rather cold plasma, even for smaller values of κ, there is virtually no difference in the maximum growth obtained from the two different distribution functions. However, the larger the value of the plasma β, the more significant becomes the influence of the assumption of κ-distributed ions. For instance, for β=0.15 and κ=2, the obtained maximum growth rate is more than twice as large as in the case of a Maxwellian distribution function. This feature can also be seen in Fig. 4, where the growth rate is shown as a function of k and κ for β=0.15. The thick black profile corresponding to the Maxwellian solution is put into this figure at κ=50 for comparison.

V. DISCUSSION AND OUTLOOK

In this paper we have presented our first approach in identifying the effects related to κ-distributed ions on the growth rate of the MTSI. For this, we basically applied a traditional approach, assuming the electrons were strongly magnetized whereas the ions were taken to be unmagnetized. A frame of reference was adopted in which the electrons were assumed at rest and the ions moved across them. A particular feature of the applied coordinate system was that the wave vector was assumed to be aligned along a coordinate axis whereas the zero-order magnetic field is inclined toward it. This assumption keeps the analysis for the ion dielectric response function as simple as possible as there is only one component of the wave vector present. After deriving the dispersion relation for the case of a potential electric field and parallel bulk flow and wave propagation, we present numerical solutions for the growth rate for various input parameters.

We find that the effects associated with κ distributed ions are directly related to the plasma beta. For a rather cold plasma, the influence of the spectral index on the growth rate is weak, and the obtained maximum growth rates are comparable over a relatively wide range of κ values. However, for a larger value of the plasma β, i.e., $\beta > 0.1$, the impact of κ distributed ions can be quite significant. We find that the maximum growth rate increases for decreasing κ. The solution obtained via a Maxwellian distribution function acts as a lower limit. It is shown that the maximum growth rate can even be doubled for sufficiently large β and small κ.

As this work presents the first approach in analyzing the effects of the κ distribution function on the MTSI, several simplifying assumptions are made. These assumptions include the consideration of the electrostatic problem, the electrons are still taken to be Maxwellian, and the wave vector is assumed to be parallel to the bulk flow. Future studies are devoted to the consideration of the more general problems by dropping the respective assumptions step by step.

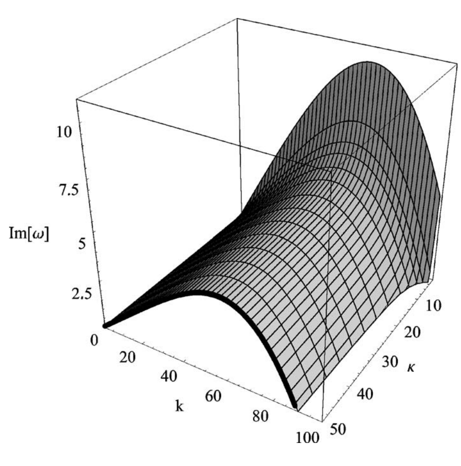

FIG. 4. Growth rate as a function of k and κ in normalized units for a fixed value of the plasma β, i.e., β=0.15. The thick black profile corresponds to the solution for Maxwellian ions.

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