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Effect of q-non-extensive distribution of electrons on the plasma sheath floating potential

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In this paper, a collisionless unmagnetized plasma sheath consisting of electrons following non-extensive q-distribution, and cold mobile inertial ions is studied in the stationary state. In this type of plasma with non-Maxwellian electron distribution (Tsallis statistical mechanics), the effective electron temperature ($T_{e,\mathrm{eff}}$) and electron screening temperature ($T_{e,*}$) are evaluated. The other plasma sheath phenomena such as the Bohm sheath criterion, Debye shielding, floating potential, and sheath length are investigated in the presence of q-non-extensive velocity-distributed electrons. It is observed that above-mentioned phenomena depend significantly on the non-extensive parameter q.

1. Introduction

Plasma sheath is a non-neutral region that normally appears at plasma boundaries to balance the loss of the electrons and ions in the plasma. It is one of the well-known aspects of the confined plasma. In isotropic plasma with equal numbers of positive ions and electrons, the electrons are far more mobile than the ions. The plasma will therefore be charged positively with respect to the grounded walls and a so-called sheath is formed (Hershkowitz 2005). The knowledge of plasma sheath properties is important for understanding the plasma surface interactions (Ghomi et al. 2006, 2007; Sharifian and Shokri 2007, 2008, 2010), Langmuir probe characteristics, spacecraft charging (Parker 1978; Conrad 1987; Cooke and Katz 1988; Mandell et al. 2006), etc. Many papers have been written about the formation of sheath in plasmas confined to a finite volume or in dusty plasmas around the surface of a small astronomical body (Baishya and Das 2003). However, the investigation of the sheath formation in the astrophysical environment apparently requires consideration of the interaction of exposed surfaces with solar wind plasma in which the plasma streaming velocity is very high (Goldstein et al. 2005).

Pines et al. (2010) have formulated a novel kinetic theory of sheath formation around the absorbing planar surface immersed in the solar wind plasma in the case in which the photoelectric effect has been neglected. They have also applied this theory to the cases of planar, cylindrical, and spherical electrodes immersed into collisionless plasma. Furthermore, they have proposed a full kinetic formulation of the

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sheath formation in solar wind plasma around an arbitrary-oriented planar absorbing surface with floating potential or any induced potential (Pines et al. 2010).

Baishya et al. have investigated the dynamics of the negatively charged dust particles in a plasma sheath above a surface in space (Baishya and Das 2003). Dusty plasma as a unique multi-component plasma is composed of dispersed macroscopic dust-charged grains in any given parent plasma background thus forming a colloidal-type suspension (Baishya et al. 1999). The dust grains with finite size make a qualitative modification of the plasma particles' conventional charging mechanism. This is a common aspect in astrophysical problems (Mendis and Rosenberg 1994), as well as in laboratory plasmas (Selwyn et al. 1990; Sheehan et al. 1990; Boufendi et al. 1993), causing new findings in the dynamics of the plasma due to contamination by the charged dust grains (Baishya et al. 1999).

Most of the papers, which investigated the plasma sheath, assumed that the plasma electron distribution is Maxwellian (Emmert and Henry 1992; Edelberg and Aydil 1999; Wang et al. 2005; Ghomi et al., 2006; Sharifian and Shokri 2007, 2008, 2010). However, it is well known that the Maxwell distribution is valid only for the macroscopic ergodic equilibrium state. Moreover, the Maxwell distribution is insufficient to explain the long-range interactions in collisionless unmagnetized plasma in which a non-equilibrium stationary state exists. Space plasma observations and gas discharge plasmas at low gas pressures clearly indicate that the electrons are far away from their thermodynamic equilibrium (Shukla et al. 1986; Ghosh and Bharuthram 2008; Pakzad 2009; Pakzad and Javidan 2011). In the space plasmas, with the use of probes or spacecraft itself as a large probe, it is shown that one deals with a non-equilibrium situation. The electrons and ions could be accelerated toward the plasma boundary by weak electric fields which are produced in the plasma. Therefore, the electron distribution may well differ from a Maxwellian distribution, because there is a flow of electrons toward the wall, probe or spacecraft. In gas discharge plasma at low gas pressure, the electrons and ions are lost continuously to the walls. Also, a similar situation exists in laboratory dusty plasmas. There is a continuous flow of electrons and ions to the dust grains where they recombine. Therefore, in these cases, the electron and ion distributions are not exactly Maxwellian.

In recent years, a great deal of attention has been paid to them based on extensive statistical mechanics in accordance with deviations of the Boltzmann–Gibbs–Shannon (B–G–S) entropic measure. An appropriate non-extensive generalization of the B–G–S entropy for statistical equilibrium has first been recognized by Rényi (1955) and subsequently proposed by Tsallis (1988), suitably extending the standard additivity of the entropies to the nonlinear non-extensive case in which the entropic index q determines the non-extensivity degree of the considered system ($q = 1$ corresponds to the standard extensive B–G–S statistics). Non-extensive statistics were successfully exerted on a number of astrophysical plasma (Amour and Tribeche 2010; Tribeche et al. 2010; Bains et al. 2011a, b; Parvin et al. 2011; Eslami et al. 2011a; Eslami et al. 2011b; Liu et al. 2011; Pakzad 2011a; Pakzad and Tribeche 2011; Reza Pakzad 2011; Ghosh et al. 2012; Roy et al. 2012a, b). In statistical mechanics and thermodynamics, non-extensive systems are defined as systems for which the sum of the entropies of the respective parts is different from the entropy of the whole. In other words, the generalized entropy of the whole is greater than the sum of the entropies of the parts if ($q < 1$) (super-extensivity), whereas the generalized entropy of the system is smaller than the sum of the entropies of the parts if $q > 1$ (subextensivity). In acco- rdance with the attestations (Lima et al. 2000; Abe et al. 2001; Kaniadakis 2001; Wada 2002; Leubner 2008; Amour and Tribeche 2010; Tribeche et al. 2010;

Bains et al., 2011a; Eslami et al., 2011a; Eslami et al. 2011c; Pakzad 2011b; Sahu 2011; Tribeche and Merriche 2011; Ghosh et al., 2012), the q-entropy might provide a convenient frame for the analysis of many astrophysical scenarios, such as solar neutrino or wind problem, stellar polytropes, and peculiar velocity distribution of galaxy cluster. It might be noted that the q-distribution is non-normalizable for $q < -1$. In the extensive limiting case ($q \to 1$), the q-distribution reduces to the well-known Maxwell–Boltzmann velocity distribution (Ghosh et al. 2012).

In the conventional hydrodynamic treatment of plasma sheath formation, the well-known Bohm sheath criterion is formulated as the necessary existence condition for the sheath. From this criterion, ions have to enter the sheath region with a supersonic velocity, assuming that they are being accelerated by electric fields in a 'pre-sheath' region, with a typical size of sheath region of the order of a few electron Debye lengths (Riemann 1991, 2009; Pines et al., 2010).

Aside from its applicability to the sheath formation, this research might be particularly connected to the moon exploration at the dark side and even more important at the quite large shadowed craters and terminator region just beyond the optical terminator line. This particular region is of intense interest for upcoming activities of exploration. Accurate analysis of the extension and strength of electric fields in space charge plasma sheath is important for understanding the dust transport and electrostatic levitation of charged dust grains (Stubbs et al. 2006) as well as for the safety of electronic and other equipment (Pines et al., 2010).

To the best of our knowledge, sheath in plasma with non-extensive electrons has never been addressed in the plasma literature. For plasmas in which the electron distribution is not exactly Maxwellian, the constant T_e is not the 'electron temperature' but rather a formal global parameter having the dimension of temperature. For the sake of completeness, effective electron temperature ($T_{e,\mathrm{eff}}$) and electron screening temperature ($T_{e,*}$) definitions have been evaluated and resulting expressions are presented. Also, we have investigated the effect of q-non-extensive distributed electrons on the other characteristics of the plasma sheath in the time-independent state such as the Bohm sheath criterion, Debye shielding, floating potential, and sheath length. This paper is organized as follows: in the next section, we present the basic equations of the hydrodynamics model of the plasma sheath by considering the q-distribution function. Our results and conclusions are presented in Sec. 3.

2. Basic equations

It has been considered that the charge-neutral collisionless unmagnetized plasma in contact with a planar wall consists of ions and electrons. The time-independent system of equations for the dynamics of a one-dimensional plasma sheath is governed by

$$\frac{\partial}{\partial x}(n_i u_i) = 0, \tag{1}$$

$$u_i \frac{\partial u_i}{\partial x} = -\frac{e}{M}\frac{\partial \phi}{\partial x}, \tag{2}$$

$$\frac{\partial^2 \phi}{\partial x^2} = -\frac{e}{\varepsilon_0}[n_i - n_e], \tag{3}$$

where n_i and n_e are, respectively, the ion and electron density, u_i is the ion velocity, M is the mass of ion, e is the magnitude of the electron charge, and ϕ is the electrostatic potential.

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To model the effects of electron non-extensivity, it has been referred to the following q-distribution function given by (Silva et al. 1998; Lima et al., 2000; Amour and Tribeche 2010)

$$f_e(q, v_e, T_e, \phi) = C_q \left[1 - (q-1)\left(\frac{mv_e^2}{2T_e} - \frac{e\phi}{T_e}\right)\right]^{1/(q-1)}, \tag{4}$$

where the constant of normalization is

$$C_q = \begin{cases} n_{e0}\dfrac{\Gamma\left(\frac{1}{1-q}\right)}{\Gamma\left(\frac{1}{1-q}-\frac{1}{2}\right)}\left(\dfrac{m(1-q)}{2\pi T_e}\right)^{1/2}, \text{ for } -1 < q < 1 \\ n_{e0}\left(\frac{1+q}{2}\right)\dfrac{\Gamma\left(\frac{1}{q-1}+\frac{1}{2}\right)}{\Gamma\left(\frac{1}{q-1}\right)}\left(\dfrac{m(q-1)}{2\pi T_e}\right)^{1/2}, \text{ for } q > 1 \end{cases},$$

and where q stands for the strength of non-extensivity, Γ for the standard gamma function, and T_e is measured in energetic units. It may be useful to note that for $q < -1$, the q-distribution (4) is non-normalizable (Lima et al. 2000). In the extensive limiting case ($q = 1$), (4) reduces to the well-known Maxwell–Boltzmann velocity distribution.

Integrating $f_e(v_e, T_e, \phi)$ over the velocity space and noting that for $q > 1$, (4) exhibits a thermal cutoff on the maximum value allowed for the velocity of the particles, given by

$$v_{\max} = \left(\frac{2T_e}{m}\left(\frac{1}{q-1} - \frac{e\phi}{T_e}\right)\right)^{1/2}, \tag{5}$$

which can be used to obtain the electron density

$$\begin{aligned} n_e(q, T_e, \phi) &= \begin{cases} \displaystyle\int_{-\infty}^{+\infty} f_e(q, v_e, T_e, \phi) dv_e, \text{ for } -1 < q < 1 \\ \displaystyle\int_{-v_{\max}}^{+v_{\max}} f_e(q, v_e, T_e, \phi) dv_e, \text{ for } q > 1 \end{cases} \\ &= n_e\left[1 + (q-1)\frac{e\phi}{T_e}\right]^{(1/q-1)+(1/2)}, \end{aligned} \tag{6}$$

and the electron flow velocity will be

$$u_e(q, T_e, \phi) = \frac{1}{n_e(q, T_e, \phi)} \begin{cases} \displaystyle\int_{-\infty}^{+\infty} v_e f_e(q, v_e, T_e, \phi) dv_e, \text{ for } -1 < q < 1 \\ \displaystyle\int_{-v_{\max}}^{+v_{\max}} v_e f_e(q, v_e, T_e, \phi)\, dv_e, \text{ for } q > 1. \end{cases}$$

We re-emphasize that the constant T_e appearing throughout this work can only be considered as a formal parameter with the dimension of energy, whereas the (effective) electron temperature $T_{e,\text{eff}}$ (7) is a measure of local electron internal energy. Hence, the distinction between the formal parameter T_e and the really meaningful physical field quantity $T_{e,\text{eff}}$ is by no means just a formal issue but rather a substantial one and absolutely indispensable for a complete presentation and understanding of the problem at hand. The effective temperature is defined for both Maxwellian and non-Maxwellian systems and for the former one yields the 'usual' thermodynamic temperature. In conclusion, the real ('effective') electron temperature is defined as

$$T_{e,\text{eff}}(q,T_e,\phi)=\frac{m_e}{n_e(q,T_e,\phi)}\begin{cases}\displaystyle\int_{-\infty}^{+\infty}[v_e-u_e(q,T_e,\phi)]^2 f_e(q,v_e,T_e,\phi)\,dv_e, & \text{for } 1/3<q<1\\ \displaystyle\int_{-v_{\max}}^{+v_{\max}}[v_e-u_e(q,T_e,\phi)]^2 f_e(q,v_e,T_e,\phi)\,dv_e, & \text{for } q>1\end{cases}$$
$$=-\frac{\Gamma\left(-\frac{3}{2}+\frac{1}{1-q}\right)}{\Gamma\left(\frac{1+q}{2-2q}\right)}\left(\frac{T_e}{(q-1)}+\mathrm{e}\phi\right)$$

or

$$\frac{T_{e,\text{eff}}(q,T_e,\phi)}{T_e}=-\frac{\Gamma\left(-\frac{3}{2}+\frac{1}{1-q}\right)}{\Gamma\left(\frac{1+q}{2-2q}\right)}\frac{1}{(q-1)}[1+(q-1)\,e\phi/T_e]. \tag{7}$$

Moreover, in the present problem, there is yet another relevant temperature appearing in the Bohm criterion, namely the 'electron screening temperature' (Riemann 2000),

$$T_{e,*}(q,T_e,\phi)=\frac{en_e(q,T_e,\phi)}{dn_e(q,T_e,\phi)/d\phi}=\frac{1}{(1/q-1)+(1/2)}\left(\frac{T_e}{q-1}+e\phi\right)$$

or

$$\frac{T_{e,*}(q,T_e,\phi)}{T_e}=\frac{2}{(q+1)}[1+(q-1)e\phi/T_e]. \tag{8}$$

The characteristic length scale in plasma is the electron Debye length, λ_{De}, which is the distance scale over which significant charge densities can spontaneously exist. Applying non-extensive electron density (6) at temperature T_e and the method which has been used in chapter 2 of Lieberman and Lichtenberg (1994), the electron Debye length in plasma with non-extensive electrons will be found as follows:

$$\lambda_{De}(q,T_e,n)=(2\varepsilon_0 T_e/(q+1)e^2 n)^{1/2}. \tag{9}$$

Assuming (a) non-extensive electron distribution (6) at temperature T_e, (b) cold ions ($T_i=0$), and (c) quasi-neutrality approximation, setting $n_{e0}(x=0)=n_{i0}\ (x=0)=n_0$ at the plasma–sheath interface (interface between essentially neutral and non-neutral region), defining the zero of the potential ϕ at $x=0$, and taking the ions to have a velocity u_0 there, ion energy conservation (assuming no collision) then gives

$$\frac{1}{2}Mu^2(x)=\frac{1}{2}Mu_0^2-e\phi(x). \tag{10}$$

The continuity of ion flux (1) without ionization in the sheath is

$$n_i(x)u(x)=n_0u_0, \tag{11}$$

where n_0 is the ion density at the sheath edge. Solving for u from (10) and substituting in (11), we obtain

$$n_i(x)=n_0\left(1-\frac{2e\phi(x)}{Mu_0^2}\right)^{-1/2}. \tag{12}$$

Substituting n_i and n_e into Poisson's equation (3), we obtain

$$\frac{d^2\phi(x)}{dx^2}=\frac{en_0}{\varepsilon_0}\left\{\left[1+(q-1)\frac{e\phi(x)}{T_e}\right]^{(1/q-1)+(1/2)}-\left[1-\frac{e\phi(x)}{E_0}\right]^{-1/2}\right\}, \tag{13}$$

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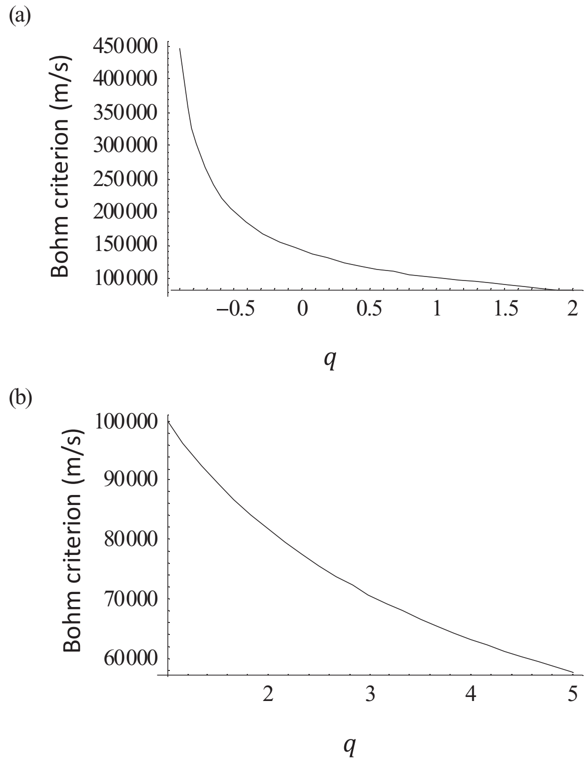

FIGURE 1. Variation of the Bohm criterion with the electron entropic index q.

where $E_0 = \frac{1}{2}Mu_0^2$ is the initial ion energy. The generalized Bohm criterion states that the ions always enter the sheath region with a velocity equal to or larger than the ion-acoustic speed (Riemann 2000),

$$u_0 \geqslant u_B = \left[\frac{T_{e,*} + \gamma_i T_i}{M}\right]^{1/2}\Bigg|_{\phi=0}. \tag{14}$$

In plasma with cold ions ($T_i = 0$), the ion-acoustic speed reduces to

$$u_B = \left[\frac{T_{e,*}}{M}\right]^{1/2}\Bigg|_{\phi=0}, \tag{15}$$

and this leads to the form of the Bohm criterion,

$$u_0 \geqslant u_B = \left[\frac{T_e}{M(q+1)}\right]^{1/2}. \tag{16}$$

Using the quasi-neutrality approximation at the plasma–sheath interface and q-non-extensive distribution of electrons, the electron flux to the wall will be

$$j_e(q, T_e, \phi_f) = j_{e0}\left[1 + (q-1)\frac{e\phi_f}{T_e}\right]^{(1/q-1)+(1/2)}, \tag{17}$$

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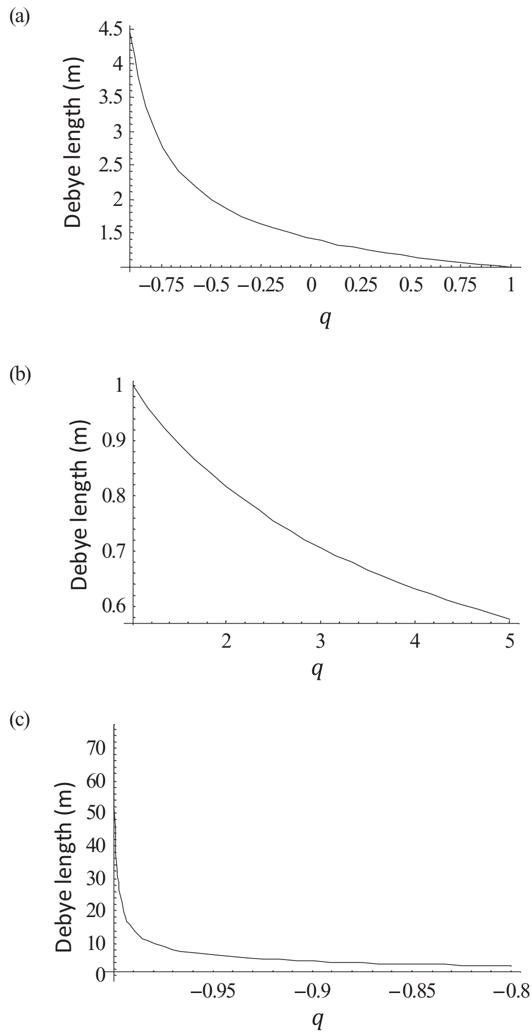

FIGURE 2. Variation of the Debye length (λ_D) with q in different intervals of q. (a) $1 < q < 5$, (b) $-0.8 < q < 1$, and (c) $-1 < q < -0.8$.

where the floating potential (ϕ_f) is the wall potential at which the total current vanishes. Here, $j_{e0} = n_0(T_e/2\pi m)^{1/2}$ is the chaotic flux at the sheath edge, where m is the electron mass.

Assuming the continuity of the ion flux (1) at the stationary state and no ionization in the sheath region, the positive ion flux to the surface in the sheath will be

$$j_i = n_0 u_0 = n_0 \left[\frac{2T_e}{M(q+1)}\right]^{1/2}. \tag{18}$$

From the model, the floating potential is defined as the electric potential reached by an isolated wall inside the plasma. First, expressions for the current densities collected at the wall due to each kind of particles (ions and electrons) are obtained using the plasma fluid theory. Then, we equate the negative (17) and positive (18) charge current densities ($j_i = j_e$) at the wall. Finally, the floating potential and its

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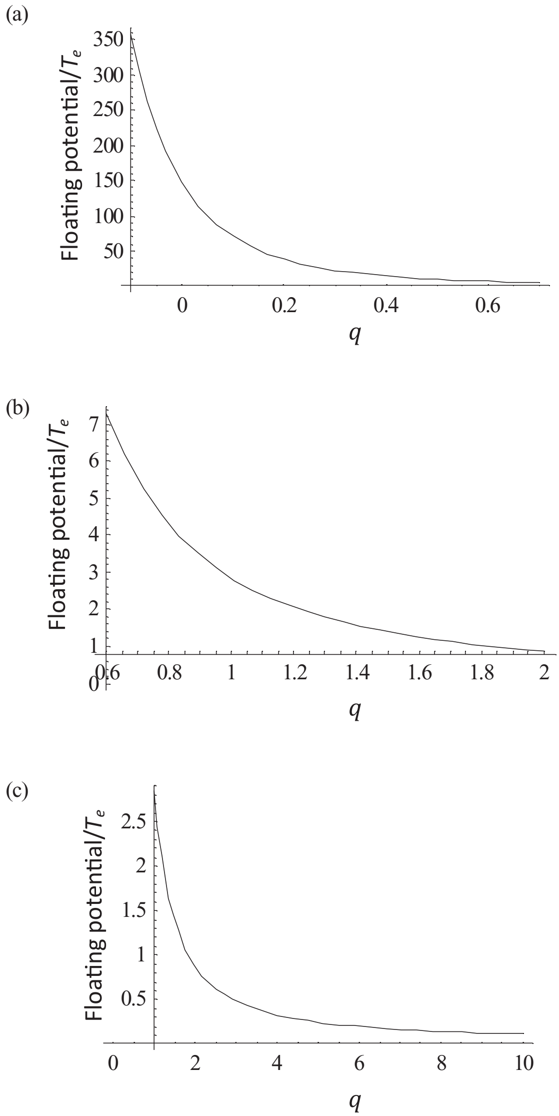

FIGURE 3. Floating potential of the plasma sheath as a function of q.

dependency on the non-extensivity of electrons q can be obtained as

$$\frac{e\phi_f}{T_e} = (1/q - 1)\{[4\pi m/(q+1)M]^{q-1/q+1} - 1\}. \tag{19}$$

The sheath thickness has been defined as a function of floating potential as shown below (Franklin 2003; Akihiro 2004; Demidov et al. 2005):

$$L_{Sh} = \left[(\sqrt{2}/3)\left[\left(1 + 2\left|\frac{e\phi_f}{T_e}\right|\right)^{\frac{1}{2}} - 2\right]^{\frac{3}{2}} - (2\sqrt{2})\left[\left(1 + 2\left|\frac{e\phi_f}{T_e}\right|\right)^{\frac{1}{2}} - 2\right]^{\frac{1}{2}}\right]\lambda_{De}. \tag{20}$$

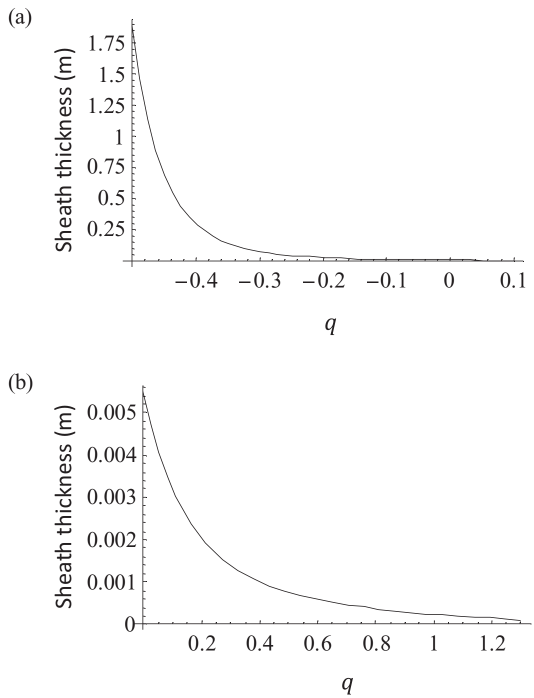

FIGURE 4. Thickness of the near-wall sheath as a function of the values of the non-extensive q-parameter.

3. Results, discussions and conclusions

We now proceed with the presentation of our results. We have considered a two-component collisionless unmagnetized plasma system consisting of electrons following non-extensive q-distribution, and cold mobile inertial ions.

In Fig. 1, the variation of the Bohm criterion with the electron entropic index q is shown. It can be seen that an increase of q leads to a decrease of the Bohm criterion value. Comparing Fig. 1(a) with Fig. 1(b) shows that the electron entropic index q effect on the value of the Bohm criterion for $q < 1$ is stronger than for $q < 1$.

Figures 2(a)–(c) display the variation of the Debye length (λ_D) with $q(-1 < q < 5)$ for different intervals of q. Apparently, the Debye length is found to increase as $q \to -1$. To see the effect of the electron non-extensivity, we have considered here the three possible ranges of the non-extensive parameter (q), viz., $1 < q < 5$, $-0.8 < q < 1$, and $-1 < q < -0.8$. For $q < -0.9$, the enhancement of λ_D becomes more rapid as q decreases (see Fig. 2(c)).

The effect of q on the floating potential of the plasma sheath is presented in Figures 3(a)–(c). It can be seen that as q increases, i.e., the non-extensive character of the plasma becomes important, the floating potential decreases. For $q > 1$, this falloff of the floating potential becomes very slow as q increases (see Fig. 3(c)).

For the sake of complete study of the plasma sheath, we have plotted the thickness of the near-wall sheath as a function of the values of the non-extensive q-parameter in Fig. 4. Its behavior is consistent with the results displayed in Fig. 3, where the effective floating potential is found to increase as $q \to -1$. It can be seen that as q increases, the value of the near-wall sheath thickness decreases.

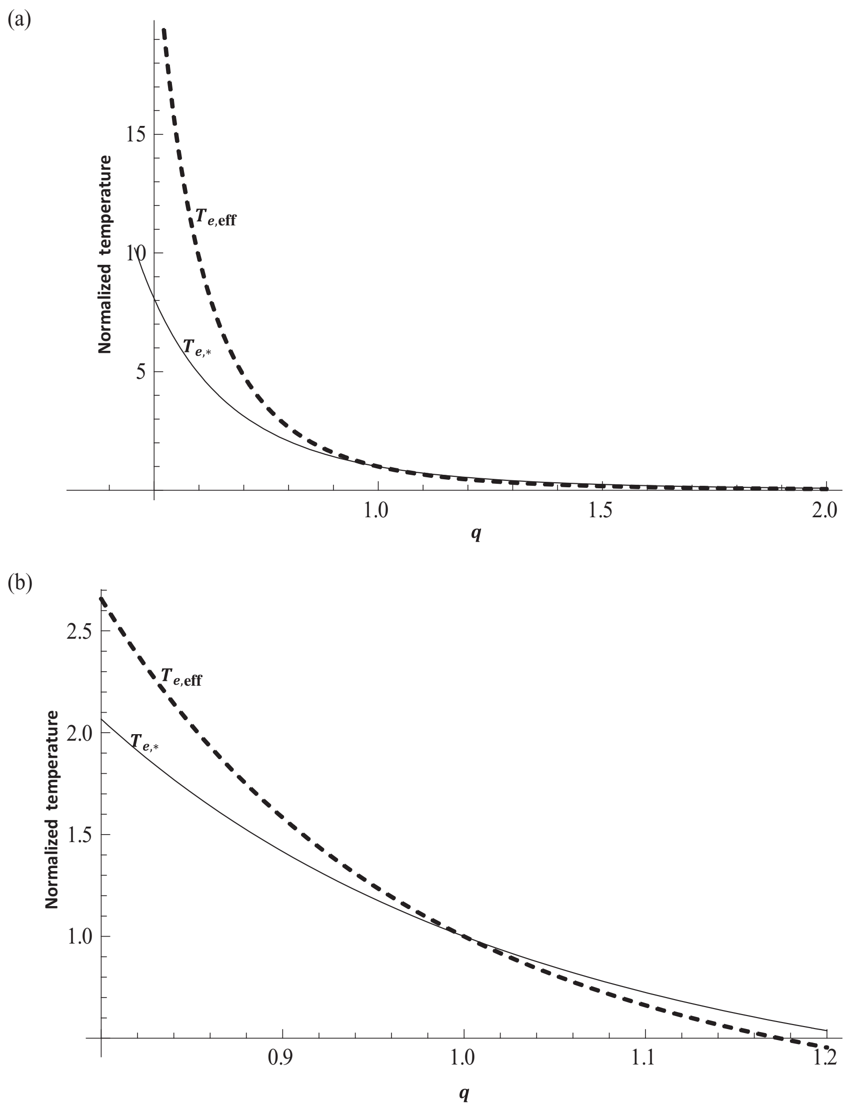

FIGURE 5. The effective electron and electron screening temperature near the wall as a function of electron entropic index q.

The real ('effective') electron ($T_{e,\mathrm{eff}}$) and electron screening ($T_{e,*}$) temperatures near the wall as a function of q have been illustrated in Fig. 5. It can be seen that both normalized temperatures are equal to 1 at $q = 1$. This result has been predictable because in the extensive limiting case ($q \to 1$) the q-distribution reduces to the well-known Maxwell–Boltzmann velocity distribution in which the effective electron and electron screening temperatures are identical. Figure 5(a) illustrates that as the non-extensive character,q, grows to values more than unity ($q > 1$), the effective electron and electron screening temperatures tend toward zero. This behavior is related to the fact that the density of high-energy electrons which are in the tail of the electron distribution function decreases with increase of the non-extensive character, q. From Figs. 5(a) and (b), it could be seen that the effective electron and electron screening temperatures are very different when $q < 1$. For example, at $q = 1/3$, the normalized $T_{e,*}$ becomes about 8, but normalized $T_{e,\mathrm{eff}}$ tends to infinity.

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