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Electrostatic effects during neutral beam propagation through plasmas

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Abstract. In this paper several aspects of the interaction of 1-D neutral beams with plasmas are considered. In order to clearly understand the dual rôles of a background plasma in collisionally decelerating the beam and its response to the consequently generated E-field, we examine two cases: an artificial unresponsive background plasma, that corresponds closely to the case of an unionized gas; and a realistic background plasma, whose charges respond to the presence of any E-field. In the former case, the electric field results solely in extremely small scale oscillation of the beam electrons about the protons as both decelerate collisionally and, although electron runaway is possible, the numbers involved are found to be negligibly small. However, collisional separation of the beam electrons and protons does occur in the case of a realistic background plasma, since E easily drives a neutralising dense plasma electron current. Despite the charge separation, runaway is found to be impossible in this case because of the masking effect that the background plasma electrons have on the separated charges unless the plasma is extremely tenuous compared to the beam, where the plasma behaviour is close to that of a near-neutral gas. These effects are shown by approximate analytic mean particle solutions and confirmed by simulations from an electrostatic particle code, which can deal with collective effects. Additional numerical simulations are also performed to investigate the generation of Langmuir waves by a neutral beam. The consequences of these results are discussed for flare neutral beams and corresponding radiation signatures. As regards hard X-ray production by runaways, although some very high energy electrons may result, the number is far too small to be relevant to solar flare HXR burst production. The beam and plasma oscillations may however be relevant to radio bursts.

Key words: acceleration of particles – Sun: flares – Sun: X-rays, gamma rays – Sun: radio radiation – plasmas

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1. Introduction

Fast particle beams play a very important role in many physical processes observed in space, e.g. in solar flares (Kundu et al., 1989) and cosmic jets (Benford 1985, Blandford, Begelman and Rees 1982). Such beams can be registered directly by particle detectors in interplanetary space (Lin et al. 1986), or indirectly through their electromagnetic emissions. Beams can be charged (electron, proton and ion beams) or ionised but neutral, having both electron and proton components of the same density and velocity. An important aspect of such a neutral beam is that it does not constitute an electric current, with possible implications for the global electrodynamics of the acceleration process (e.g. Martens 1988). Partly for this reason neutral beams have been invoked as playing a central role in some outstanding astrophysical problems. In particular Simnett and Haines (1990) have claimed that injection of a neutral beam whose energy flux is initially in the ions may lead to hard X-ray bremsstrahlung production by acceleration of runaway electrons in an electric field generated by collisional separation of charges in the beam. Secondly Lesch et al (1989) have proposed that similar charge separation effects, resulting from injection of a fast neutral flow into a cool gas ring around a galactic centre, constitute a current which creates a magnetic field and that this may answer the fundamental cosmic problem of production of a seed magnetic field for dynamo action (cf. also Biermann 1950). Other astrophysical applications may include the interaction of galactic (AGN) and stellar (T-Tauri jets) with their ISM/IGM environments. Though we do not address the issue here, a key question is how such jets manage to remain collimated and rectilinear so far from their sources. While these ideas are interesting and stimulating, in none of these analyses is the electrodynamics treated self consistently. This is in part because the complexity of real cosmic situations demands approximations and estimates. It is clear, however, that improvement in our understanding of neutral beam propagation is essential to these important astrophysical problems.

In this paper we shed light on these problems by examining the behaviour of neutral beams in some simplified situations to enable progressive insight into real situations. In particular we obtain solutions for a series of idealised situations both by analytic and particle simulation treatments (using an electrostatic particle code) and compare the results of these two approaches with those claimed earlier by others. One element which we take care to include is the displacement current term which is frequently ignored, since it can play a central role in the transient behaviour which can be crucial to the nature of any final quasi-steady state reached as was noted by Brown and Bingham (1984) in relation to return current neutralisation of charged beams (cf. Miller 1982, Oss and van den Oord 1995).

Due to more effective collisional deceleration of the electron component of the neutral beam, compared to the protons, electric fields are generated which act against the collisional electron-deceleration force. Intuitively, the appearance of this electric field E can usually be viewed as being generated by a non-zero charge density or, equivalently, by the displacement current. However, in an infinite homogeneous beam, for which the charge density is everywhere zero, there exists an E field due to the commonly neglected displacement current $\partial E/\partial t = -j$, where the current density j is due to the differing beam electron and proton velocities due to collisions. In this case, the charge density picture is intuitively confusing, but in some sense it can be understood that the E-field is generated by charges at infinity. No magnetic fields are involved, as shall be proved below.

In our treatment we start by examining the behaviour of electrons, ions and electric fields generated by the interaction of a neutral 'beam' with a dense background plasma whose particles are 'immobile' so that it provides collisional drag on the beam particles but does not produce a background current in response to the electric field. This case corresponds very closely to that of an unionised gas, the only difference being in the constants appearing in the collisional term. We shall refer to this artificial background as an *unresponsive plasma*. First we consider the homogeneous situation in which the 'beam' exists everywhere, electrons and ions being 'launched' initially with the same speed. Then we consider a true 'injection' problem where the beam is semi-infinite and beam head effects appear explicitly. Next, we consider a realistic background plasma which responds to the electric field by producing a plasma current which can significantly modify the system's behaviour as noted by Simnett and Haines (1990). In our analytic treatment of these problems the beam is treated as having unique particle speeds, so a test particle approach is required to treat electron runaway. However, our particle simulations have some randomisation of speeds and therefore should exhibit the runaway effects proposed by Simnett and Haines (1990).

2. Electric field generation by a neutral beam in an unresponsive plasma

2.1. Infinite homogeneous beam case

In this case the background plasma only exerts a collisional drag on the beam particles so that electrodynamic effects are confined to the beam. Since the system is infinite and homogeneous and the beam is purely $1-D$, any $\mathbf{E}$ arising must be uniform in space so that $\nabla \times \mathbf{E} = \partial \mathbf{B}/\partial \mathbf{t} = \mathbf{0}$ so that no magnetic field $\mathbf{B}$ exists if $\mathbf{B} = \mathbf{0}$ at $t = 0$. Even then, this purely temporal problem is complicated by the large range of timescales involved - from the electron plasma period up to the proton collision time. Since the proton collision time is long compared to the electron collision time and extremely long compared to the electron plasma period, we first simplify the problem by considering the proton speed to remain constant and the electron speed to vary by only a small fraction of the initial beam (and proton) speed so that the electron collision term remains constant.

That is, we start from the equations (in a mean particle description of plasma processes):

$$m_e \frac{dv_e}{dt} = -eE - C_e(v_e), \tag{1}$$

$$m_p \frac{dv_p}{dt} = eE - C_p(v_p), \tag{2}$$

$$\frac{dE}{dt} = 4\pi ne(v_e - v_p), \tag{3}$$

where n is the beam density, v_e and v_p are electron and proton velocities, and C_e and C_p are collisional drag forces on electrons and protons due to the surrounding dense gas. Throughout this paper, primed quantities refer to the background plasma particles whereas unprimed quantities refer to the beam particles. The subscripts e and p are exclusively used to denote electron and proton respectively.

Based on our previous remarks, these equations simplify to:

$$m_e \frac{dv_e}{dt} = -eE - C_e, \tag{4}$$

$$\frac{dE}{dt} = 4\pi ne(v_e - v_0), \tag{5}$$

where v_0 is the constant proton speed.

These equations can be rewritten into one for the (oscillatory) motion of the relative proton-electron velocity $u = v_e - v_p$, viz.

$$m_e \frac{d^2u}{dt^2} + 4\pi e^2 nu = 0, \tag{6}$$

the solution of which is

$$u = -\frac{C_e}{\omega_e m_e} \sin \omega_e t, \tag{7}$$

where ω_e is the angular beam electron plasma frequency with $u(0) = 0$ and $\frac{du}{dt}(0) = -\frac{C_e}{m_e}$ since $E(0) = 0$. The corresponding electric field can be expressed as

$$E = \frac{C_e}{e}(\cos \omega_e t - 1). \tag{8}$$

It follows that the electrons here are fully dragged by the protons due to the mean electric field E= -eC_e as stated by Simnett and Haines (1990), though oscillating about them at angular frequency ω_p with very tiny space and velocity amplitudes which can be written

$$\frac{\Delta v_e}{v_0} = \frac{\Delta x}{l} = 9 \times 10^{-10} \frac{n'/(10^{10}\mathrm{cm}^{-3})}{(n/(10^8\mathrm{cm}^{-3}))^{1/2}} \frac{1}{(v_0/(10^{10}\mathrm{cm\ s}^{-1}))^3} \quad (9)$$

where l is the nonthermal beam Debye length v_0/ω_e and we have used

$$C_e = \frac{12\pi e^4 \Lambda n'}{m_e v_e^2} \quad (10)$$

Here Λ is the Coulomb logarithm, n' is the density of the background gas, e is the charge, m_e is the electron mass and v_e the beam electron velocity. The factor 12 in this expression accounts for both the effect of the electron-electron collisions (8 parts) and for the electron-proton collisions (4 parts).

A more complete numerical solution of the full analytic mean particle equations, including both C_p and C_e and their dependences on velocity, shows exactly the same behaviour with the electrons oscillating about the protons as they decelerate together due to the combined dragged force C_e+C_p, the electric field increasing as the velocities fall and the Cs increase.

2.2. Semi-infinite beam case

We now consider the more realistic case of injection of a semi-infinite neutral beam rather than an infinite one ‘launched’ homogeneously. When such a neutral beam is injected and penetrates into a dense gas then, due to collisions, its electron component is again decelerated first. This deceleration force is equivalent to that of an electric field of magnitude E_c given by

$$E_c = -C_e/e. \quad (11)$$

The beam electron velocity initially (before E becomes significant) declines with time due to collisions according to $v_e(t) = (v_0^3 - 3Kt)^{1/3}$, where $v_e(t=0) = v_0$ is the initial electron velocity, and $K = 12\pi e^4 \Lambda n'/m_e^2 = m_e C_e {v_{be}}^2$ so that E increases as the electrons slow down. Due to the relative deceleration of electrons, a positively charged layer forms at the head of the finite neutral beam and at the tail, a negatively charged layer forms so that an electric field is generated along the whole beam. In terms of the initially growing charges in the widening end layers this field can be expressed as

$$E_d(t) = 4\pi n e d(t), \quad (12)$$

where $n = n_e = n_p$ is the neutral beam density initially equal to electron and proton beam densities everywhere, and the thickness of the charged layers is $d = v_0 t + \frac{1}{4K}(v_0^3 - 3Kt)^{4/3} - \frac{1}{4K}v_0^4$. This electric field continues to grow, with the charge layer thickness d, until E reaches the value E_c, when d takes the tiny value Δx given in Eq. 9 above, and is strong enough to balance the differential collisional drag force. Thereafter the beam electrons are dragged along on average at the proton velocity, the two species slowing down together as proton collisional losses progress, but with electron oscillations superposed on this motion. Note that in this analysis, the beam width is taken to be large compared to its length (cf Oss and van der Oord 1995 and Miller 1982).

2.3. Acceleration of electrons

Though the charge separation induced by differential collisional drag is very small, the resulting E field along the beam is sufficient to offset the collisional drag on beam electrons and will also act to accelerate some of them which deviate randomly from the exact equilibrium speed given by $C_e(v_e) = eE$. But, as can be seen from a comparison of the electric field E and the collisional force which depends on electron velocities, the acceleration is possible only if an electron has a velocity greater than that of the monoenergetic beam average. If a test electron has a velocity exactly equal to the beam velocity then forces are in equilibrium, and the test electron is moving with the velocity of beam. For initial velocities smaller than the beam velocity, electrons are decelerated.

In Simnett and Haines (1990), a calculation was performed to find the maximum possible energy of an accelerated runaway electron. The first step was to find Φ, the potential due to the assumed double layer. This step is questionable given their subsequent analysis since, as we shall show, this potential cannot remain unaffected if significant runaway occurs. A second less fundamental objection is that the m_i in the denominator of the collision term in their Eq. (8) should actually be m_e. The consequence of this last error is to reduce the maximum possible electron energy to 75% of the proton energy rather than 100%, as they stated.

Runaway electrons with high energies are in principle possible in an unresponsive plasma, if a particle has a speed greater than the average beam speed. The key issue now is how many electrons can be involved in this. Till now we considered only one test electron, or more generally a number sufficiently small not to affect the electrodynamic equations we used to find the neutral beam generated E field. But clearly, with an increasing number of runaway electrons their electrodynamic effect must be taken into account. Here we make a rough estimate of the upper limit to the number of runaway electrons which can be produced before their effect on E is significant. According to the previous scenario, protons of the neutral beam effectively pull electrons by a “charge separation” electric field

$$E_c = -\frac{12\pi e^3 \Lambda n'}{m_e v_0^2}. \quad (13)$$

Any runaway electrons accelerated in this field would effectively result in a double layer of charges with associated electric field

$$E_d = 4\pi n_r e x_r, \quad (14)$$

where n_r is the density of runaway electrons, and x_r is the length of the runaway electron pulse. In order for the runaways not to modify the E field accelerating them we require that $E_d \ll E_c$. This gives the following constraint on the fraction of electrons that can runaway from a beam of length L:

$$\frac{n_r x_r}{nL} \ll \frac{3e^2 \Lambda n'}{nLm_e v_0^2} = \frac{\Lambda}{N}\frac{l}{L}. \quad (15)$$

where $N = (4\pi/3)nl^3$ is the beam 'plasma parameter' (number of beam electrons in a Debye sphere). Since l is a small length, of the order of cm in solar conditions, and since N is always very large (typically of the order of 10^{10}) compared to Λ ($\sim$ 20) the runaway fraction is constrained to be extremely small.

We conclude, therefore, that runaway acceleration of electrons, whatever their origin, in the E field of a neutral beam, for the case of an unresponsive background plasma, has no relevance to the generation of flare hard X-ray emission. Such emission would require (from the basic bremsstrahlung equation - Brown 1972) very large electron fluxes, carrying a total electron beam power of near the total impulsive impulsive flare power (Brown 1971, Hoyng et al 1976 and others).

3. Electric field generation by a neutral beam in a realistic plasma

3.1. Background plasma response - analytic description

Assuming again that protons are infinitely heavy (so $C_p = 0$), then the motions of the beam and background electrons in the mean particle description can be expressed as:

$$m_e \frac{dv_e}{dt} = -eE - C(v_e), \quad (16)$$

$$m_e \frac{dv_e'}{dt} = -eE - C(v_e'), \quad (17)$$

$$\frac{dE}{dt} = 4\pi ne(v_e - v_0) + 4\pi en'v_e', \quad (18)$$

where n' is the plasma density, while v_e and v_e' are beam and plasma electron velocities. (Note that here this set of mean particle equations is used solely for the study of collisional effects and cannot describe strong collective effects such as the two-stream instability).

First, let us consider times sufficiently short so that $v_e - v_0 \ll v_0$ and so $C(v_e) \approx C_0$, and $C(v_e') \approx 0$.

Then the previous equations can be rewritten:

$$m_e \frac{dv_e}{dt} = -eE - C_0, \quad (19)$$

$$m_e \frac{dv_e'}{dt} = -eE, \quad (20)$$

$$\frac{dE}{dt} = 4\pi ne(v_e - v_0) + 4\pi en'v_e'. \quad (21)$$

Elimination of dE/dt from Eqs. (16-18) gives:

$$\frac{d^2 v_e}{dt^2} = -\frac{4\pi e^2}{m_e}(n'v_e' + n(v_e - v_0)) = \frac{d^2 v_e'}{dt^2}, \quad (22)$$

the solution of which is

$$v_e = v_e' - \frac{C_0 t}{m_e} + v_0. \quad (23)$$

Now, using Eq. (23), the equation for the second time derivative of the relative beam electron velocity $u = v_e - v_0$ is

$$\frac{d^2 u}{dt^2} = -\omega^2 u + \frac{4\pi n' e^2 C_0 t}{m_e^2}, \quad (24)$$

where $\omega^2 = 4\pi(n' + n)e^2/m_e$ is the combined plasma frequency of beam + plasma electrons.

The solution of Eq. (24) is

$$u(t) = -\frac{C_0}{m_e}\frac{n'}{n' + n}(t + \frac{n}{n'}\frac{\sin \omega t}{\omega}). \quad (25)$$

Now using Eq. (23) again and the definition of u we can express the velocity of the background plasma as

$$v_e' = \frac{C_0}{m_e}\frac{n}{n' + n}(t - \frac{\sin \omega t}{\omega}). \quad (26)$$

From these equations it can be seen that, as the beam electrons decelerate, the background plasma electrons react to the electric field $E(t)$ (generated by the beam separation), increasing their velocities in the beam direction. Oscillatory terms appear in both equations (for v_e and v_e') and have the same amplitude whereas the secular terms differ in size by a factor n'/n, as should be the case for a return current (see below).

The above describes the early evolution on plasma timescales, driven by the ('infinitely' massive) protons. The electric field created again oscillates, as for an unresponsive plasma background, but it is reduced compared to (8) by the factor $n/(n + n')$ since the free plasma electrons now provide part of the total zero current condition ($j_b + j_p + dE/dt = 0$). The beam electrons are thus enabled to separate from the protons with a differential speed growing with time at a rate depending on n'/n. In the limit $n' \to 0$ the previous neutral gas solution is recovered.

One must next ask how the system evolves on longer timescales and, in particular, whether it approaches a steady state after, and averaged over, times long compared to ω. If we set all time derivatives in Eqs. (20) and (21) to zero we get

$$-eE = C(v_e) = C(v_e') \quad (27)$$

and

$$nv_e + n'v_e' = nv_0 \quad (28)$$

Since $C(v)$ is monotonic (above thermal speeds) $C(v_e) = C(v_e')$ implies $v_e = v_e' = nv_0/(n + n')$ so that the beam and plasma electrons become indistinguishable, the motion of both being

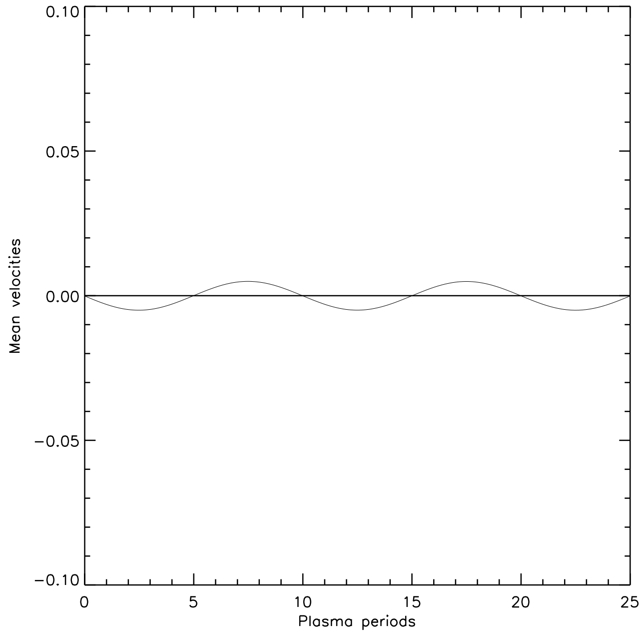

Fig. 1. Mean velocities of proton (full thick line) and electron (full thin line) components of neutral beam for the case with the unresponsive background plasma.

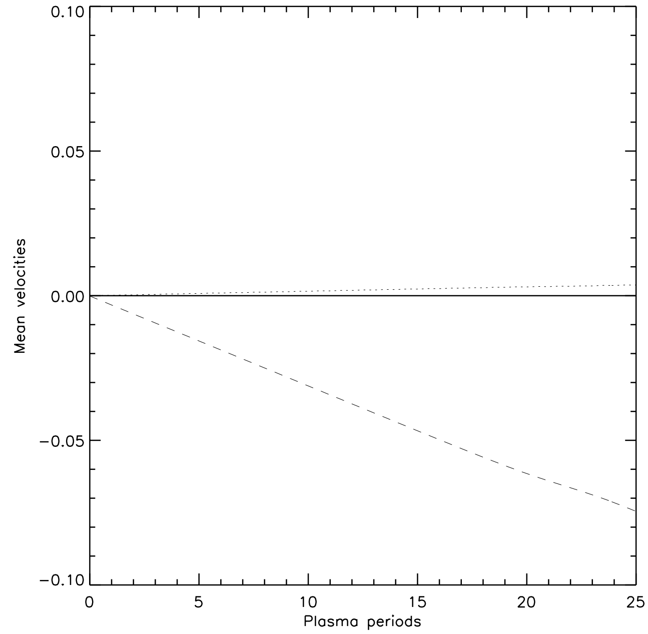

Fig. 2. Mean velocities of proton (full line) and electron (dashed line) components of neutral beam, and that of background plasma (multiplied 10 times - dotted line) for the case with the background plasma.

described by Ohm's Law $-eE = C(v_e)$ and the total current being zero i.e. $j_{pe} + j_{be} = -j_{bp}$ (the ambient protons remaining stationary). Whether this steady state is approached in practice depends on whether it is established on a timescale short compared to the beam lifetime as set by the actual stopping time of finite mass protons t_{sp} (during which the ambient protons will also react to E and make a small correction to j). By Eq. (24) this will occur on a time t^* over which $\Delta u \approx -v_0$ or

$$t^* \approx \frac{m_e v_0}{C_0}\frac{n' + n}{n'} = t_{se}\frac{n_p + n}{n'} = t_{sp}\frac{m_e}{m_p}\frac{n' + n}{n'} \qquad (29)$$

Thus provided $n/n' \ll m_p/m_e$ (ie the beam is not much denser than the plasma) we expect the beam electrons to be essentially decoupled from the beam protons by the action of the plasma electron drift current to reduce E, with negligible runaway. In solar flare beam and plasma conditions this condition is very amply met.

3.2. Numerical simulation

To check these conclusions we simulated the above processes by a numerical electrostatic particle code (see Birdsall and Langdon, 1985; Peratt, 1992). Firstly, we consider the case of a neutral beam in an unresponsive plasma and then proceed to consider a realistic background plasma. A homogeneous infinite neutral beam, as studied in Sect. 2.1 cannot be simulated in a finite space code, even with periodic boundary conditions. The reason is that in the code there can be no charges at infinity and it is such charges which can be thought of as the source of E in this problem. To get around this, we made a simulation using a code with periodic boundary conditions, but with the neutral beam shorter than the length of the system so that beam-end effects are present. For the length of the system we took L=20π (N_G = 256 grid points) with the beam initially between $L/4$ and $3L/4$, i.e. occupying half of the system length. We assume the protons to be infinitely heavy. As convenient numerical values, the beam and background densities were taken to be 1 and 100 respectively, the ratio of the total numbers of plasma and beam electrons in the system then being 200, because the beam occupies only one half of the system.

For the case of a neutral beam in an unresponsive plasma we considered the problem in the beam frame and took the initial velocities of both electron and proton components to be zero. Then at every time step Δt=0.2 (31.4 steps = plasma period of background plasma = $2\pi/\omega_e'$), to simulate the effect of collisions, we reduce the velocity of beam electrons by 10^{-4} (where the velocity is expressed in the code velocity unit ($L/(N_G \Delta t)$). The result is shown in Fig. 1, where the oscillatory motion of the electron component of the neutral beam with the beam plasma period can be seen. This result confirms our previous conclusions of Sect. 2.1.

Now we investigate the case where the background plasma can respond. In this case, however, as shown in the following, there is a strong collisionless effect - the two-stream instability - not reproducible in a mean particle approach. In reality, though, a number of factors serve to reduce the effect of the two stream instability, eg a distribution of beam particle velocities. However, to examine the collisional effect numerically, we need to perform our simulation in an artificial reference frame where the neutral beam, which has zero velocity relative to the background plasma, is nevertheless forced to experience a collisional drag. If the angular plasma frequency of the background plasma is 1

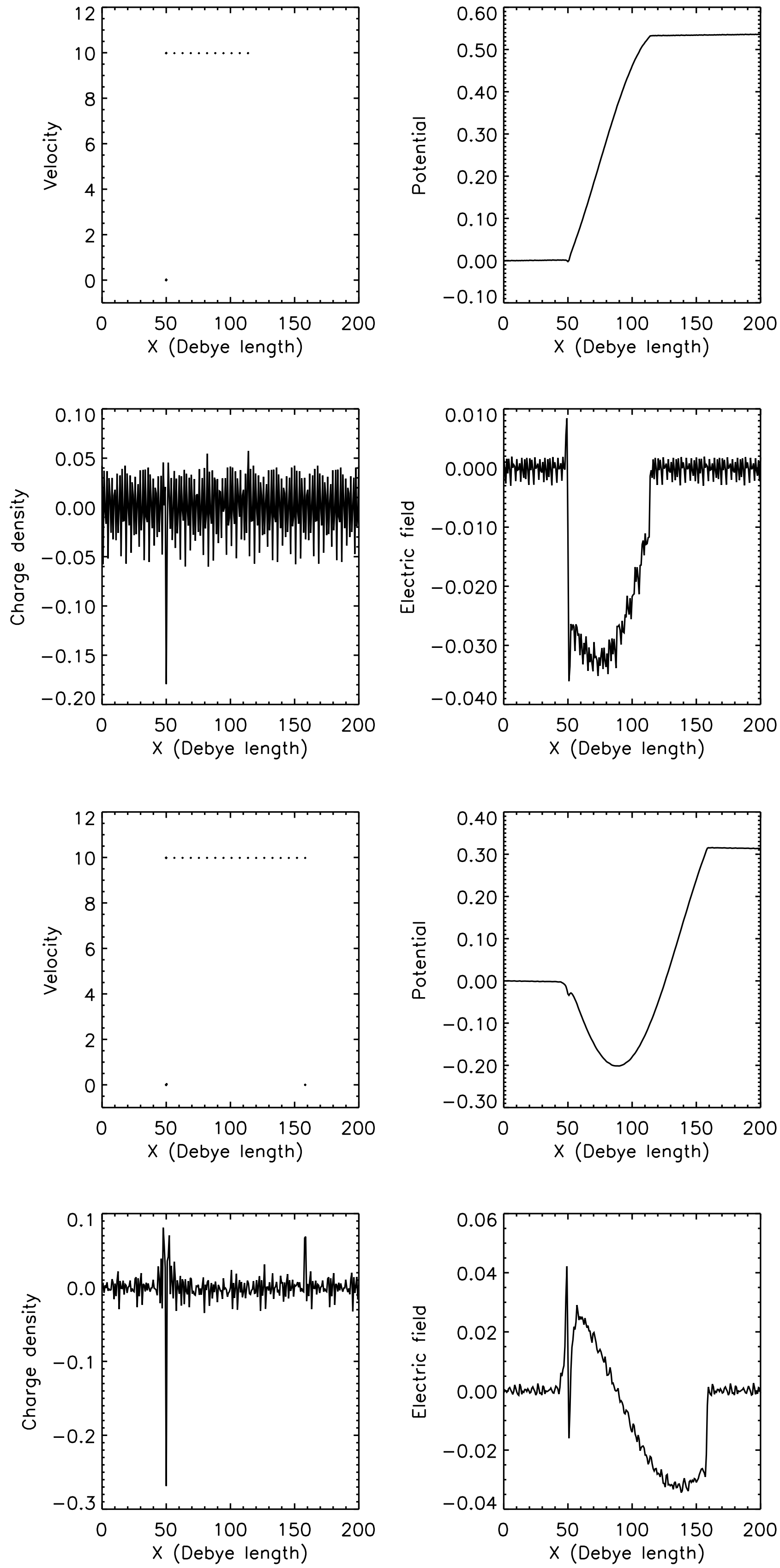

Fig. 3a and b. The velocity and positions of protons and electrons, the potential, charge density and electric field during the neutralization process of the neutral beam by the background plasma at 0.318 T (**a**), and 0.636 T (**b**), where T is the plasma period.

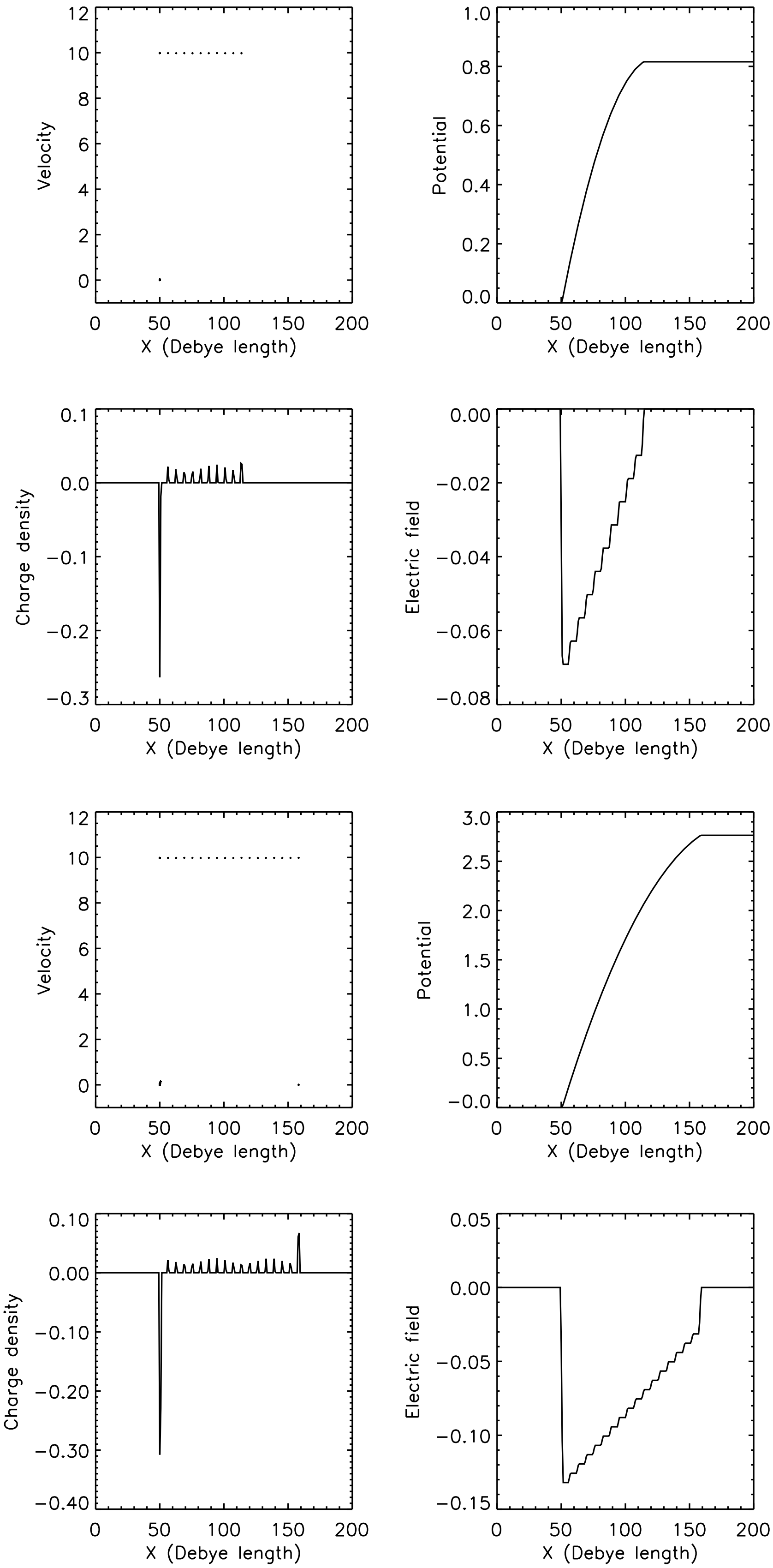

Fig. 4a and b. The velocity and positions of protons and electrons, the potential, charge density and electric field during the evolution of the neutral beam in dense plasma, but without the neutralization process by the background plasma at 0.318 T (**a**), and 0.636 T (**b**), where T is plasma period.

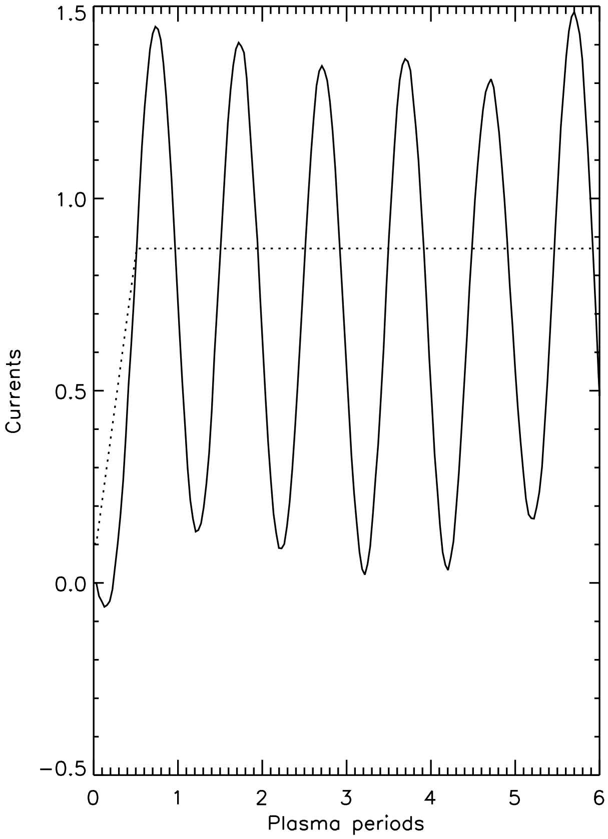

Fig. 5. The time evolution of plasma (full line) and beam (dotted line) currents. The plasma current has opposite sign.

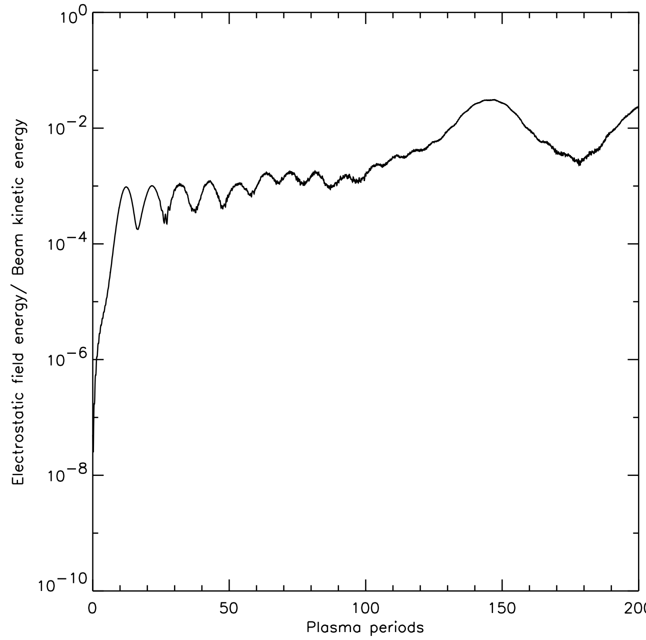

Fig. 6. The time evolution of electrostatic field energy in the simulated neutral beam-plasma system.

and that of the beam is 0.1 then the beam/plasma density ratio is the same as the beam/gas density ratio in the previous case. Because all parameters of this simulation are the same as in the unresponsive plasma case, these situations can be simply compared. Results are shown in Fig. 2, where the mean velocities of the proton beam component (full line), electron beam component (dashed line), and background plasma electrons (dotted line - values are multiplied by 10 for better visualization) are depicted. Weak plasma oscillations of the background plasma were recognized in the evolution of the electrostatic field energy.

This simulation, despite being artificial so as to suppress collective effects, confirms clearly the very important effect of the previous analysis. That is, in the presence of a background plasma, electron and proton components of a neutral beam become spatially separated due to collisions. This is because the growth of E arising from the beam electron-proton separation current is neutralized by a background plasma electron current - i.e. $dE/dt = -(j' + j)$ instead of just $dE/dt = -j$.

3.3. The electric field at the beam head

It was shown in Sect. 3.1 that collisional separation of electron and proton components of a neutral beam is possible only in the presence of a sufficiently dense background plasma ($n'/n \gg m_e/m_p$). Here we follow in detail the process of charge separation at the electron and proton head fronts. We used the following numerical simulation parameters: numerical system length $L = 20\pi = 200\lambda_D$, where $\lambda_D = v_e^T/\omega_e'$ is the Debye length, v_e^T is the thermal plasma electron velocity; 10000 and 10000 numerical electrons and protons representing background plasma; time step $\Delta t = 0.2\omega_e'^{-1}$; proton-electron mass ratio is 100.

At position 50 λ_D along the beam we mimic/exaggerate the collisional separation process in the neutral beam as follows: one electron is stopped there (to confine the process under study within our limited numerical space we adopt an abrupt stop to zero velocity); a corresponding proton is allowed to propagate with the velocity $v = 10$ in code unit $(L/(N_G\Delta t))$ to the right up to the position 156 λ_D, where it is suddenly stopped. The evolution of the system is shown in Figs. 3a and 3b, for the velocities and positions of electrons and protons of the neutral beam, the potential, together with the charge density and electric field, for two times very early in the process. To see the neutralizing effect of the background plasma we add Figs. 4a and 4b where the same variables at the same times are shown for the beam evolution, but without the background plasma effects. We see that a double layer is formed with an electric field, where the electrons stop, of $E_e = 4\pi eFt$ and where the protons stop of $E_p = 4\pi e(Ft - n_{bi}d)$ where $F = n_b v_b$ is the particle flux.

Comparing Figs. 3 and 4 we see the plasma neutralizing effect. Near the position x=50 λ_D the electric field starts to be screened – see sharp peaks of the electric field at this position (Figs. 3a and 3b).

Simultaneously, the plasma between the two stopping places starts to oscillate with the plasma frequency. This oscillation process can be seen also in Fig. 5, where the plasma and beam currents are depicted. The beam current then increases up to a

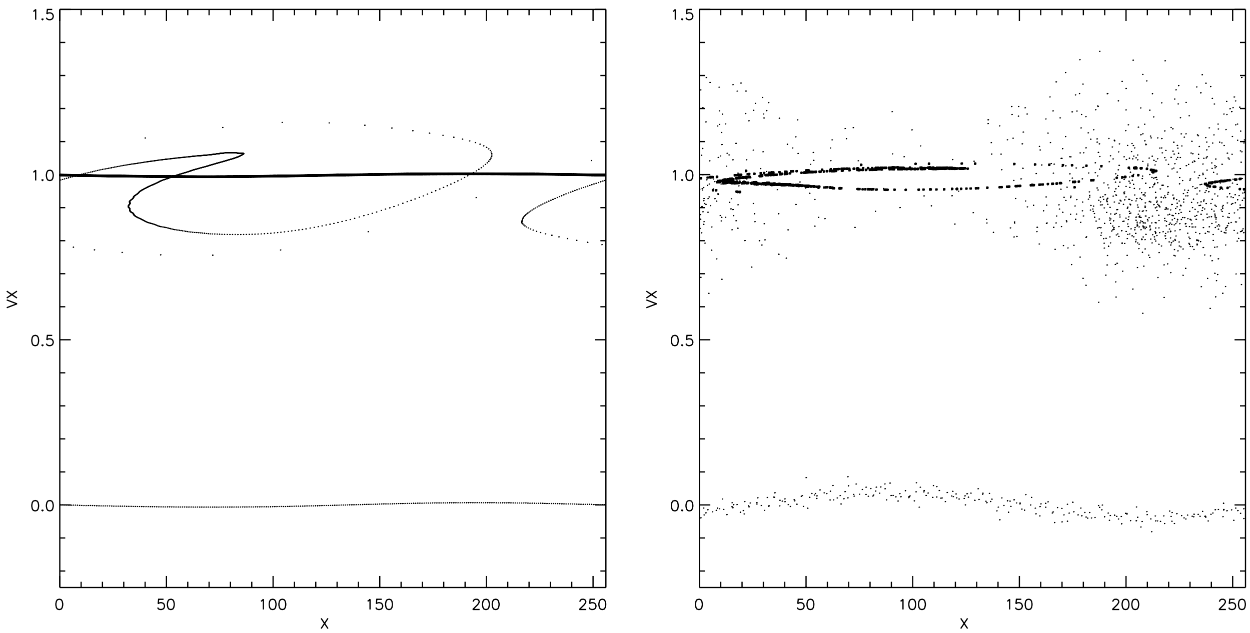

Fig. 7a and b. The phase space at times 15.9 T (**a**), and 159.2 T (**b**) for the neutral beam-plasma system. The electrons and protons are represented by the smaller and bigger dots respectively.

constant value over the finite propagation time of protons from where the electrons are stopped to where the protons are stopped. On the other hand, the plasma (return) current oscillates from zero to a value of twice the beam current so that the total mean charge current is zero and the instantaneous total current including the displacement current is zero - i.e. $j' + j + dE/dt = 0$, as discussed previously.

4. Generation of Langmuir waves

Here we consider the full behaviour of collective effects by using a 1-D electrostatic particle code similar to that in the paper by Messerotti and Karlický (1991). The case of an infinite neutral beam will be compared with the pure electron and proton beam cases. The length of the system, which was chosen to be equal to the wavelength of the most unstable wave, was set to 2π. The proton-electron mass ratio was chosen as 100; the electron and proton beam frequencies are $\omega_e^2/\omega_e'^2 = 10^{-3}$ and $\omega_p^2/\omega_e'^2 = 10^{-5}$. Cold monoenergetic beams were considered. The time evolution of the electrostatic field energy is shown in Fig. 6. From these results it is evident that Langmuir waves are generated in two steps. Firstly, the neutral beam behaves like a pure electron beam and later on like a proton beam. This means that at the beginning the growth rate of the instability is as high as that for an electron beam. Furthermore, the first saturation level of Langmuir waves is given by electron trapping. On the other hand, the high level of Langmuir waves at the second saturation level is determined by the beam protons. The phase space diagrams (Figs. 7a and 7b) show that firstly electrons are trapped, and after a mixing phase (t= 70-110 T, where T is the plasma period) the proton-beam instability starts and protons are trapped in the intense Langmuir wave. These numerical results fully agree with the theoretical estimates given by formulae for the growth rates, saturation levels and trapping periods, as follows (Mikhailovskii, 1974; Drummond et al., 1970; LeQueau and Roux, 1987):

$$\frac{\gamma}{\omega_e'} = \frac{3^{1/2}}{2^{4/3}} \left(\frac{n}{n_p}\right)^{1/3} \left(\frac{m_e}{m_b}\right)^{1/3}, \tag{30}$$

$$W_s = \frac{1}{2} n' m_b v_b^2 \left(\frac{R}{2}\right)^{1/3}, \tag{31}$$

where $R = \omega_{pb}^2/\omega_e'^2$,

$$P = \frac{2\pi}{(keE/m_b)^{1/2}}, \tag{32}$$

where γ is the growth rate of the two-stream instability, W_s is the wave energy density at the saturation level, P is the trapping period, m_b is the mass of beam particles, v_b is the beam velocity, while k and E are the k-vector and the electric field amplitude of Langmuir waves, respectively.

These considerations are purely intended to indicate some of the complication likely to arise when high densities of coherent plasma waves occur in beam/plasma/gas interaction. We have not considered in any detail how wave scattering may allow stable propagation of beams but we observe that such interaction may determine the ultimate fate of the beam and may be crucial to an understanding of coherent radiation diagnostics such as Type III bursts.

5. Conclusions

We have considered the propagation of a monoenergetic neutral beam travelling in a dense background plasma, comparing our results with those of Simnett and Haines (1990). In the first instance, we treated the problem without including the response of the background plasma, as was done in Simnett and Haines (1990) (they only considered it *a posteriori*, concluding that it only reduced the size and steepness of the potential drop). Although electrons do suffer a greater collisional deceleration, careful consideration of the generated electric field shows that the protons will always "drag" the electrons along with them. In this case therefore, no significant separation of the beam electrons and protons can occur, the charge separation being oscillatory in nature and on a scale extremely small compared to the beam non-thermal Debye length. A particular consequence of this result is that the formation of a spatially extended double layer, as assumed by Simnett and Haines (1990), is not possible under solar chromospheric conditions, invalidating their subsequently proposed mechanism for electron runaway. In terms of the unresponsive plasma description presented in this paper, runaway is still possible in principle, with some electrons gaining as much as 75% of the beam proton energy (the 100% quoted by Simnett and Haines (1990) arose from an erroneous factor of m_e/m_p appearing in their proton equation of motion). However, in order to avoid "destroying" the self-consistent electric field that accelerates them, only a very small number of such electrons can actually runaway. We conclude, therefore, that electron runaway from neutral beams is negligible in this scenario.

When the response of the background plasma is considered, the end conclusion is the same, but for different reasons. As the beam electrons are decelerated, the generated electric field serves to accelerate the background plasma electrons, which in turn serves to neutralise the current. During this initial stage, both beam and plasma electrons undergo oscillations of the same amplitude and frequency. Ultimately, the beam electrons will be decelerated until they have the same mean speed as the plasma electrons. The two populations of electrons will then become indistinguishable, with the beam protons proceeding alone until stopped by collisions. In this more realistic scenario, the response of the background plasma masks the charge separation, and in doing so makes runaway again negligible.

All of the results described above were obtained using analytic mean particle methods. To add weight to our conclusions, we compared our results to those from an electrostatic particle code, confirming all the above results. In addition to this confirmation, further numerical simulations were carried out to investigate the collective effects. It was shown that a monoenergetic neutral beam generates Langmuir waves in two steps, firstly, like an electron beam and then like a proton beam. This means that at the very beginning the instability growth rate is high, corresponding to an electron beam instability. Then, after a mixing phase, a proton beam instability starts with a much lower growth rate. The final high saturation level of plasma waves corresponds to that of a proton beam.

These effects of neutral beam propagation have interesting consequences for interpretation of radio emission from chromospheric layers during the flare impulsive phase as we discuss in a subsequent paper. According to Simnett and Haines (1990) they also provide a new mechanism for production of flare hard X-ray bursts, though they do not make it clear how. Emission of deka-keV HXRs needs deka-keV electrons with, in a non-thermal beam model, a very large particle flux (Brown 1971). If these electrons were part of a neutral beam the accompanying protons would carry an energy flux m_p/m_e times larger and far in excess of the total flare power available. If, as Simnett and Haines suggest, the HXR electrons are runaways accelerated by the E fields created by the neutral beam, then even if the E field were capable of such acceleration, the runaway would have to be extremely energy efficient, transferring most of the flare power from the protons to the runaways and a large electron beam current. Such a mechanism does not eliminate the need for intense electron beams, though it produces them via neutral and proton beam processes rather than directly as a first generation acceleration product.

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