PARTICLE-IN-CELL SIMULATION OF A STRONG DOUBLE LAYER IN A NONRELATIVISTIC PLASMA FLOW: ELECTRON ACCELERATION TO ULTARELATIVISTIC SPEEDS

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ABSTRACT

Two charge- and current-neutral plasma beams are modeled with a one-dimensional particle-in-cell simulation. The beams are uniform and unbounded. The relative speed between both beams is 0.4c. One beam is composed of electrons and protons, and the other of protons and negatively charged oxygen (dust). All species have the temperature 9.1 keV. A Buneman instability develops between the electrons of the first beam and the protons of the second beam. The wave traps the electrons, which form plasmons. The plasmons couple energy into the ion acoustic waves, which trap the protons of the second beam. A structure similar to a proton phase-space hole develops, which grows through its interaction with the oxygen and the heated electrons into a rarefaction pulse. This pulse drives a double layer, which accelerates a beam of electrons to about 50 MeV, which is comparable to the proton kinetic energy. The proton distribution eventually evolves into an electrostatic shock. Beams of charged particles moving at such speeds may occur in the foreshock of supernova remnant (SNR) shocks. This double layer is thus potentially relevant for the electron acceleration (injection) into the diffusive shock acceleration by SNR shocks.

Key words: acceleration of particles – methods: numerical – plasmas

Online-only material: color figures, mpeg animations

1. INTRODUCTION

The blast shell of a supernova remnant (SNR) can accelerate charged particles to cosmic-ray energies through diffusive shock acceleration by the SNR shocks (Uchihama et al. 2007; Reynolds 2008; Hillas 2005). Particles are scattered upstream and downstream of a shock and some particles cross the shock repeatedly, gaining energy each time (Bell 1978; Krymsky 1977; Blandford & Ostriker 1978; Malkov & Drury 2001). Diffusive shock acceleration requires a seed population of particles with energies well in excess of the thermal ones, since only such particles can cross the shock repeatedly and be scattered efficiently by the magnetohydrodynamic waves on either side of the shock. The presence of shock-accelerated electrons is evidenced by the radio emissions of SNRs. It is not clear yet how the electrons can reach the energy threshold for diffusive shock acceleration; that requires their kinetic energies are comparable to those of the ions. Some electrons must be accelerated locally, that is close to the SNR shock or close to one of its precursors (Reynolds 2008; Hillas 2005).

This injection of electrons is thought to be accomplished by the interplay of energetic ions with electrons and dust (Ellison et al. 1997) by means of plasma instabilities (Cargill & Papadopoulos 1988). The instabilities occur in the upstream plasma, just ahead of the SNR shock. This region is called the foreshock. One source of the ion beams is the downstream plasma. Ions can leak through the shock boundary and outrun it (Malkov 1998). The shock can also reflect a substantial fraction of the upstream ions. If the reflection is specular, then the ions can cross the foreshock at twice the shock speed. Reflected and leaked ion beams can be observed in situ at the Earth bow shock (Eastwood et al. 2005). Particle-in-cell (PIC) simulations of fast plasma shocks suggest that they may also occur at SNR shocks when the upstream magnetic field is orthogonal to the shock normal (Hoshino & Shimada 2002; Lembge et al. 2004; Chapman et al. 2005; Amano & Hoshino 2007). Precursor shocks move ahead of the main SNR shock. They can reach mildly relativistic speeds for the most extreme supernova explosions (Kulkarni et al. 1998). Such shocks can also reflect ions in the presence of a quasi-parallel magnetic field (Dieckmann et al. 2008c).

We focus in this work on ion beam instabilities in the foreshock, but we do not model the shock itself. We rely on a foreshock model that has been widely used, e.g., in the works by Shimada & Hoshino (2003, 2004) and Dieckmann et al. (2007, 2008a). This model considers an ion beam, which moves through the upstream plasma, consisting of cool electrons and protons. The system is spatially uniform. Often a second, counterpropagating ion beam is introduced, which cancels the current of the first one. This second beam can be due to ions that have been reflected by a perpendicular shock and which return to the shock, after they have been rotated by the upstream magnetic field (McClements et al. 1997). Here, a beam of comoving negatively charged oxygen ions cancels the current of a proton beam.

A foreshock model that assumes the spatial uniformity of the plasma with respect to the mean speed, the plasma flow direction and temperature can be used only if all processes develop on scales well below an ion Larmor radius and an inverse ion cyclotron frequency in the case of a perpendicular shock. Otherwise, the rotation of the beams by the magnetic field and the shock restructuring (Hoshino & Shimada 2002) has to be taken into account. The foreshock model may be representative over larger scales if the shock is not a perpendicular one because ions can expand much farther if they move along the magnetic field, as it is observed at the Earth’s bow shock. This foreshock model is, however, not self-consistent because it decouples the foreshock physics from that of the shock. It is nevertheless useful because we can examine with it the acceleration of the...
upstream electrons as a function of the bulk parameters, such as the beam temperatures and mean speeds and the background magnetic field. We can also separate individual processes from the global shock dynamics that involve coupling over many scales (Hoshino 2003). This simplifies their identification and interpretation. The periodic boundary conditions also allow us to reduce the box size, making PIC simulations computationally efficient, even for realistic ion-to-electron mass ratios. The insight gained from such simulations can then guide us in the selection of plasma parameters for the more realistic and demanding PIC simulations of full shocks.

Simulation studies employing this foreshock model have revealed that proton beam-driven instabilities are not efficient electron accelerators. An obstacle is the mass difference between the electrons and the protons. Most plasma instabilities saturate by accelerating the electrons. The fields are too weak to accelerate the ions and the heated electrons quench most of the known plasma instabilities, e.g., through Landau damping. Often only a few percent of the ion energy is transferred to the electrons (Shimada & Hoshino 2003, 2004; Dieckmann et al. 2007, 2008a). Initially, this is also the case for ion beams that expand into the upstream plasma (Dieckmann et al. 2008b) prior to the formation of a shock. Certain mechanisms (Kuramitsu & Krasnoselskikh 2005; Katsouleas & Dawson 1983) can, in principle, accelerate the electrons to relativistic speeds, but this remains yet to be demonstrated with self-consistent PIC simulations. However, the electrons, that have been accelerated by an instability like the Buneman instability (BI; Buneman 1958), are often spatially bunched, for example in the form of phase-space holes (Roberts & Berk 1967; Korn & Schamel 1996a, 1996b; Dieckmann et al. 2007). Electron bunches or plasmons have a much higher inertia than electrons and they modulate the ions to give ion acoustic waves. The ion acoustic waves can trigger further processes, one of which we consider in this work.

We investigate a plasma consisting of two spatially uniform beams. One beam is composed of protons and electrons and the second one of protons and oxygen carrying a single negative charge. Each of the beams is charge and current neutral. Both beams move at a relative speed of 0.4c. All species have the temperature 9.1 keV. These initial conditions are the simplest ones that take into account three particle species and they permit a first assessment of the impact of heavier ions. Section 2 discusses in more detail these initial conditions and their potential relevance for a foreshock plasma. Section 3 describes the simulation results, which can be summarized as follows.

A BI develops between the electrons and the protons of the second beam. The BI saturates by forming electron phase-space holes and these plasmons feed wave energy into the ion acoustic wave (Mendonca & Bingham 2002; Mendonca et al. 2005). Proton phase-space holes develop once the ion acoustic waves are strong enough to trap the beam protons. Proton phase-space holes are, strictly speaking, electrostatic potentials in which the trapped protons gyrate on closed orbits. Here we use this term also for structures in the proton phase-space distribution that will eventually develop into phase-space holes. Such structures are characterized by filaments that have been accelerated out of the proton bulk distribution. The electrons eventually thermalize and reach a mildly relativistic thermal speed, which is stabilizing the system for some time. The BI plays no further role after the electron thermalization. We examine in detail a proton phase-space hole that grows into being a proton density pulse with a negative density (Infeld et al. 1989; Infeld & Rowlands 1990). The growth of this proton phase-space hole may be caused by its interaction with the negative oxygen (Eliasson & Shukla 2004) or by the disruption of the electron current (Smith 1982). The pulse triggers a double layer and evolves subsequently into an electrostatic shock.

The electron phase-space structure and evolution of this double layer resembles that found by a recent simulation study by Newman et al. (2001). That study addressed the electron acceleration in the Earth’s auroral ionosphere. However, our simulation yields the much stronger electron acceleration that is necessary for injecting the electrons into the diffusive shock acceleration at SNR shocks. The proton phase-space distribution also differs from the standard picture of double layers outlined, for example, in the review papers by Smith (1982) and Raadu & Rasmussen (1988). Our simulation shows no protons that are trapped by the double-layer potential. The drift speed of the free-streaming protons with respect to the double layer also exceeds the proton thermal speed by a factor of 100, which is rather extreme. The substantial free energy stored in the difference of the mean velocities of both proton beams causes the extreme electron acceleration up to 50 MeV. The double layer converts directed proton energy into directed electron energy, without heating up the plasma. Electrons can thus be accelerated by the double layer to much higher speeds than by a shock because the latter also transfers flow energy into heat.

We discuss our simulation results in Section 4, relate it to some previous investigations of double layers, and discuss the necessary future research.

2. THE INITIAL CONDITIONS AND THE SIMULATION METHOD

2.1. Motivation of Initial Conditions

The bulk plasma of a supernova blast shell can expand into the ambient medium at a speed of up to 0.1c–0.2c, at least during its early phase (Vink et al. 2006; Fransson et al. 2002). A fraction of the blast shell plasma may expand even faster. A plasma flow speed of 0.9c has, for example, been observed for a subshell ejected by a particularly violent supernova explosion (Kulkarni et al. 1998). Shocks form, where the expanding plasma impacts on the ambient plasma, the interstellar medium. The typical shock speeds should be comparable to the bulk plasma speed, but some precursor shocks may be considerably faster. In particular, the fastest SNR shocks moving with 0.9c are filamentary because the energy dissipation is accomplished by the nonplanar filamentation modes and the mixed modes (Lee & Lampe 1973). However, a flow-aligned magnetic field results in planar shocks even for such high shock speeds (Dieckmann et al. 2008c). Lower flow speeds decrease the growth rates of the mixed and filamentation modes and the shocks should resemble the nonrelativistic ones found in the solar system plasma.

We assume that at least the fast precursor shocks of SNR blast shells can move at a speed of 0.2c–0.4c in the upstream reference frame. The relativistic momentum of the particles with such a speed is equal to twice the equivalent for the fastest observed bulk flows and only a fifth of the extreme one observed by Kulkarni et al. (1998). It may thus not be unrealistic for the precursors of typical SNR shocks. We, furthermore, assume here that the foreshock region is resembling that of the Earth bow shock (Eastwood et al. 2005) and that beams of ions are present with a speed exceeding that of the shock. We consider here a beam speed of 0.4c. The shock would move at a slightly lower speed into the ambient plasma, if ion leaking through the shock
were the beam source, and the shock speed would be 0.2c, if the beam originates from the specular reflection of the ambient plasma.

We examine the interaction between two spatially uniform and unbounded plasma beams, consisting of one negatively and one positively charged species each. One beam consists of electrons and protons and the second one of protons and oxygen carrying a single negative charge. The electrons and protons of the first beam represent the ambient plasma, into which the SNR shock is expanding. The protons and the oxygen from the second beam originate from the downstream plasma or from a reflection of ambient plasma by the shock. We expect to find the two proton components from the ambient plasma and from the leaked or shock-reflected ions in the foreshock, which is evidenced for example by the PIC simulation by Dieckmann et al. (2008c).

We initially assume that the electron thermal spread is much less than the beam speed in the foreshock, which contradicts previous PIC simulation studies of full shocks. However, the nonlinear processes we consider here develop after the BI has saturated. The BI establishes a hot electron distribution in the simulation and modulates the ion beams in a realistic and self-consistent manner. Our simulation setup contrasts the one by Newman et al. (2001), where the beams have been marginally stable with respect to the BI.

Charged dust and heavy ions exist in the supernova ejecta and in the interstellar medium (Ellison et al. 1997; Meikle et al. 2006). Their ionization state and their relative densities and distributions are, however, unclear. It is typically only the dust, which has a negative time-dependent charge (Momemii et al. 2007; Verheest & Pillay 2008), while the oxygen is positively charged (Katsuda et al. 2008). The negatively charged molecules, which occur in laboratory (electronegative) plasma experiments (Charles 2007), are not likely to be present close to the SNR shocks with its intense electromagnetic radiation and with its high plasma temperature. The oxygen with its single negative charge may thus not be realistic.

We introduce the oxygen here for the following reasons. It introduces an asymmetry in the plasma constituents between the ambient plasma and the leaked or shock-reflected plasma. The plasma, that has leaked into the upstream from the supernova blast shell (downstream) plasma, is likely to have a composition that differs from that of the ambient plasma. Even the shock-reflected ambient plasma is likely to differ from the ambient plasma. The ion reflection by the shock should first depend on the ion charge-to-mass ratio and this process is thus selective. Second, the ions in the shock-reflected ion beam have been picked up at an earlier time. The compositions of the shock-reflected plasma and of the ambient plasma will thus differ if the ambient plasma is not spatially uniform. The negative oxygen also removes the need for a counterpropagating beam of positive ions to cancel the beam current. The latter is only present at perpendicular shocks. Our simulation setup is more suitable for (quasi-)parallel shocks because of the long simulation time and because a flow-aligned magnetic field favors the (flow-aligned) electrostatic processes we focus on here over the multidimensional electromagnetic instabilities (Dieckmann et al. 2008c). Replacing the electrons of one beam with the oxygen also alters the phase speed of the BI such that a stronger interaction between the wave and the proton beams is possible. It results in proton phase-space holes. This increases the probability with which secondary processes triggered by these proton phase-space holes develop. This is necessary because rare processes, which can develop in the vast foreshock regions of astrophysical shocks, may not do so in the small simulation boxes. Furthermore, the lower charge-to-mass ratio of the oxygen ions introduces the potentially important extra timescale (Eliasson & Shukla 2004), which is realistic because SNR outflows do not consist solely of electrons and protons. The oxygen may finally be a placeholder for the negatively charged dust.

We have selected a density of the oxygen beam that equals the density of one proton beam and we have omitted any further ion species to keep the system simple, while demonstrating the potential importance of heavier ions in PIC simulations of the foreshock. No magnetic field is introduced for the same reason. We will expand our simulation studies to include more realistic beam systems and magnetic fields in future work.

2.2. Physical Parameters

We use a spatially uniform foreshock model, which has been employed in previous works (Shimada & Hoshino 2003, 2004; Dieckmann et al. 2007, 2008a), and with the plasma parameters that we have motivated above. The simulation box is placed well ahead of the shock in the upstream (ambient) plasma, which forms the first beam. The shock will not enter the simulation box at any time. The presence of the shock does, however, give rise to the second plasma beam that consists of leaked or shock-reflected ions. We consider the reference frame, in which both plasma beams move in opposite x directions and with equal speed moduli $v_b$. This reference frame would be the shock frame, if the shock would be time-stationary in its rest frame, and if the ions would have been specularly reflected.

Beam 1 is the ambient plasma into which the shock expands, and consists of electrons with the plasma frequency $\Omega_e = (e^2 n_e / m_e e_0)^{1/2}$, where $e$, $m_e$, and $n_e$ are the elementary charge, the electron rest mass, and the electron number density, respectively, and $e_0$ is the dielectric constant. The equally dense protons with the mass $m_p = 1836 m_e$ have the plasma frequency $\Omega_p = R_p^{1/2} \Omega_e$, where $R_p = m_e / m_p$. Both species move at the mean speed $\langle v \rangle = -v_b$ with $v_b = c / 5$. Beam 2 corresponds to the leaked or shock-reflected ions and is composed of protons with $\Omega_p$ and oxygen, which carries a single negative charge. The oxygen plasma frequency is $\Omega_o = \Omega_p / 4$. The mean speed of both species of beam 2 is $\langle v \rangle = v_o$. All species have the temperature $T = 9.1$ keV, which gives the electron thermal speed $v_e = (k T / m_e)^{1/2} = v_b / 1.5$.

The electron thermal speed is not high enough to suppress the instabilities.

We calculate numerically the growth rate map in the wavenumber space defined by the flow-aligned $k_x$ and one perpendicular component $k_z$. The system is symmetric around $v_b$ and $k_z$ can be neglected. We normalize the wavenumbers to $Z_e = k_x v_o / \Omega_e$ and $Z_p = k_x v_o / \Omega_p$, where the relative speed difference of both beams is $\approx 2v_b$. Figure 1 shows the growth rate $\Omega_g / \Omega_e$ for our plasma parameters that has been computed with a warm fluid model ( Bret et al. 2008). The growth rate maximum is found on the branch of unstable waves that goes over into the BI for $Z_e \to 0$. The maximum growth rate is $\Omega_g / \Omega_e = 0.049$ at $Z_e = 0$ and $Z_p = 0.59$. The growth rate decreases for increasing $Z_e$. Figure 1 reveals the presence of a second wave branch at $Z_e \ll 1$ that expands to high $Z_e$. These are the filamentation modes (Lee & Lampe 1973) that are important, if the beam speeds approach the speed of light. Their growth rate is here an order of magnitude below that of the BI.

In what follows, we concentrate on the $x$ direction, along which the most unstable wave is found. By neglecting the $T \neq 0$ and the heavy species of each beam, the dispersion relation for
In the box frame. The estimates for the beam-aligned electrostatic modes can be cast into the simple linear growth rate equation:

\[ 1 - \frac{\Omega_f^2}{(\omega + k v_b)^2} - \frac{\Omega_p^2}{(\omega - k v_b)^2} = 0. \] (1)

The wavenumber and frequency of the most unstable wave are determined as \( k_u = \Omega_f/2v_b \) and \( \omega_u = \Omega_f \) in the reference frame of beam 1. The Doppler shift will reduce \( \omega_u \) to \( \Omega_f/2 \) in the reference frame we consider. The wavelength \( \lambda_u = 2\pi/k_u \). The phase speed \( v_{ph} = \omega_u/k_u \approx 2v_b \) in the reference frame of beam 1 or \( v_b \) in the box frame. The estimates for \( \omega_u \) and \( k_u \), as well as the growth rate \( \Omega_{BC}/\Omega_f \approx (3^{1/2}/2^{1/4})(\Omega_f/\Omega_e)^{2^{1/4}} \approx 0.056 \) of the cold BI (Buneman 1958) are close to those computed with the warm fluid model. The growth rate in the warm plasma model is lower only by about 10% and the corresponding wavenumber is increased by 20%.

Replacing the oxygen with electrons gives symmetric beams, which lets the phase speed of the waves vanish in the box frame. A resonant interaction between the electrostatic wave and the protons is more difficult to achieve because the electrostatic wave would have a relatively high frequency in the rest frame of each proton beam. Our simulation will show that the electrostatic waves resulting from the BI, which are comoving with beam 2, are nevertheless strong enough to modulate the protons of beam 1. The modulation of the protons of beam 2 is, however, stronger and nonlinear processes develop on this beam.

In what follows, we consider only the \( x \)-direction due to limitations of the computer time. While this direction can resolve the fastest-growing mode, the growth rate of the oblique modes in Figure 1 is not negligible. The one-dimensional PIC simulation we perform below can still give an important first qualitative insight into the developing nonlinear processes. A comparison of the results presented here with future PIC simulations, which resolve a second spatial dimension, can furthermore provide information about how the instabilities and nonlinear processes develop on the dimensionality of the system. This would not be possible with two-dimensional PIC simulations alone.

2.3. The Simulation Code

The PIC simulation method reviewed by Dawson (1983) solves the Vlasov–Maxwell set of equations with the method of characteristics. It approximates the phase-space density \( f(x, v, t) \) of a kinetic and collisionless plasma with an ensemble of computational particles (CPs). The trajectories of the CPs are prescribed by the Lorentz equation of motion. The CPs correspond to volume elements of the space phase and are not physical particles. The charge \( q_{cp} \) and mass \( m_{cp} \) of a CP can thus differ from those of the particles it represents. The charge-to-mass ratio \( q_{cp}/m_{cp} \) must, however, equal that of the physical particles. The ensemble properties of the CPs approximate well those of the physical plasma.

The Maxwell–Lorentz set of equations can be normalized, which provides us with results that are independent of the plasma density. This is convenient with respect to astrophysical plasma, for which accurate density information is often unavailable. Let \( E_p, B_p \) be the electric and the magnetic fields and let \( p_p, J_p \) be the charge and the macroscopic current. The subscript \( p \) says that the quantities have physical units. The charge and current are related to the phase-space density of the species \( i \) as \( \rho_{p,i} = q_i \int f_i(x_p, v_p) dv_p \) and \( J_{p,i} = \int v_p f_i(x_p, v_p) dv_p \). The total charge and current are \( \rho_p = \sum \rho_{p,i} \) and \( J_p = \sum J_{p,i} \).

We normalize \( E = eE_p/(\rho_{pe} \Omega_e), B = eB_p/(\rho_{pe} \Omega_e), \rho = \rho_p/(\rho_{pe}), \) and \( J = J_p/(\rho_{pe} \Omega_e) \). The differential operators become \( \nabla = \lambda_e \nabla \) and \( (d/dt) = \Omega_e^{-1}(d/dt) \), where the electron skin depth \( \lambda_e = c/\Omega_e \). The normalized position \( x = x_p/\lambda_e \) and time \( t = t_p \Omega_e \). The normalized velocity and momentum of the \( i \)-th CP are \( v_i = v_{p,i}/c \) and \( p_i = m_{ei}/m_e = \tilde{m}_{cp,i} \Gamma(v_i) v_i \). The mass and charge of a CP are \( \tilde{q}_{cp} = q_{cp}/e \) and \( \tilde{m}_{cp} = m_{cp}/m_e \). Then,

\[ \nabla \times B = J + \frac{\partial E}{\partial t} \] (2)

\[ \nabla \times E = - \frac{\partial B}{\partial t} \] (3)

\[ \nabla \cdot B = 0, \nabla \cdot E = \rho \] (4)

\[ \frac{d p_i}{dt} = \tilde{q}_{cp,i} (\tilde{E}[x_i] + v_i \times \tilde{B}[x_i]), \frac{d x_i}{dt} = v_i. \] (5)

In what follows, we will express all quantities in normalized units, except in the figure and movie captions and labels. There we state explicitly the normalization.

The PIC codes represent \( E, B, J \), and \( \rho \) on a grid. The CPs move freely. The electric and magnetic fields have to be updated on each grid cell. The electric and magnetic fields have to be evaluated at the position of the CP to update its momentum. The charge and current of each species \( i \) are prescribed by the Lorentz equation, which updates the three-dimensional vector of the relativistic momentum for an ensemble of CPs.

The simulation box with the length \( L = 2930 \) in units of \( \lambda_e \) can resolve 1170 wavelengths \( \lambda_u \) and it is aligned with the \( x \)-direction. The box is resolved by \( N_p = 7 \times 10^5 \) grid cells, each with the length \( \Delta x \). The electron Debye length in physical units is \( \lambda_D = v_e/\Omega_e \). The \( \Delta x \) corresponds to 0.31 electron Debye lengths for our initial conditions. Each of the four species is represented by 250 CPs per grid cell, giving a total of \( N_p = 1.75 \times 10^7 \) per species. The system is followed in time for \( t_s = 2600 \), which...
is subdivided into $10^5$ time steps $\Delta_t$. The beams propagate the distance $t_S v_b < L/5$ and even a light pulse can cross $L$ just once. Effects due to the periodic boundary conditions of the simulation box cannot occur.

### 3. SIMULATION RESULTS

The electrostatic BI develops between the electrons of beam 1 and the protons of beam 2. However, the secondary instabilities give rise to electric fields, which will let all the particle species interact. The energy densities normalized to $n_em_ec^2$ of the electrostatic field $E_F(t) = N_g^{-1} \sum_{j=1}^{N_g} E_j(t, t)$ and of the particle species $i$ are $E_i(t) = N_e^{-1} \Delta_t^{-3} \sum_{j=1}^{N_g} \tilde{m}_i[\Gamma(v_j) - 1]$, where the summation is over all particles $j$ of the species $i$ with the normalized mass $\tilde{m}_i$ and a normalized speed $v_j(t)$ giving $\Gamma(v_j) = (1 - v_j^2)^{-1/2}$. The $E_F, E_1$ are normalized to $E_1(t = 0)$, the energy densities of the ions are normalized to their respective initial values and their time evolution is displayed by Figure 2.

The $E_F \propto E_1^2$ grows in the interval $t < 100$ at an exponential rate $2\Omega_e \approx 0.06 \Omega_e$, which is below the expected $2\Omega_B = 0.098 \Omega_e$. The growth rate $\Omega_B$ is calculated for a sine wave. The $L \gg \lambda_u$ and the broad $k_u$ spectrum of unstable waves in Figure 1 implies, however, that we integrate over a mixture of waves, all of which have a growth rate $\approx \Omega_B$ and the initial growth of $E_F$ cannot be compared directly to $\Omega_B$.

The electron trapping lets $E_F(t)$ saturate at $t \approx 100$ and causes the growth of $E_1(t)$. Thereafter, $E_F(t)$ decreases and reaches a meta-equilibrium at $t \approx 10^3$. The $E_1(t)$ is almost constant until $t \approx 1500$, when $E_1(t)$ and $E_F(t)$ start to grow again. All ion beams have lost less than $1\%$ of their energy prior to $t \approx 1500$, confirming that the BI is not an efficient electron accelerator. Then, $E_2(t)$ and $E_3(t)$ start to change and even $E_4(t)$ has decreased by $1\%$ at $t \approx 2600$.

Figure 3 shows the phase-space distributions of the trapped electrons and of the protons of beam 2 at $t = 145$ in a subinterval of the simulation box. The trapped electrons gyrate around a potential that is moving with the mean speed $v_b$ of beam 2, which is also the phase speed of the Buneman wave. Only a fraction of the electrons is trapped. The electrons of the untrapped bulk population move on oscillatory phase-space paths. The electron phase-space hole in the interval $0 < x/\lambda_u < 1$ has the expected size, while the electron phase-space hole in the interval $2 < x/\lambda_u < 3.5$ is larger. Electron phase-space holes

![Figure 3](image-url)
increase their size by coalescence and their mutual interaction will deform them (Roberts & Berk 1967). The mean momentum of the proton beam is weakly modulated. The initially periodic train of electron phase-space holes will coalesce, until only a dissipative equilibrium and solitary electron phase-space holes remain (Korn & Schamel 1996a, 1996b).

Figure 4 confirms this. A simulation box interval is selected at the time $t = 10^3$, in which an electron depletion at $x/\lambda_u \approx 9$ separates the turbulent region $x/\lambda_u < 8$ from a smooth plateau at larger $x$. The latter constitutes a dissipative equilibrium. The supplementary Movie 1 corresponds at its starting time to Figure 4(a). It follows the time-evolution of the electron phase-space distribution at the position of the electron depletion. The color scale is the 10-logarithm of the number of CPs. Note that the absolute position in the box is given in Movie 1, while Figure 4 uses a relative position. Movie 1 demonstrates that it is this density depletion which is responsible for the turbulent phase-space region. A spatial correlation is visible in Figure 4 between the electron depletion and a phase-space hole in the protons of beam 2. This explains why the electron depletion is stationary in the reference frame of Movie 1, which moves with $v_b$ in the box frame. The unperturbed electron distribution at $x/\lambda_u > 10$ in Figure 4(a) results in a spatially almost uniform electron flow or current in the reference frame moving with $v_b$. This disruption of the flow or current by the proton phase-space hole can destabilize the latter (Dupree 1986).

The electron temperature is now relativistically high along the simulation direction, by which the linear damping of the ion acoustic waves is reduced and these modes can accumulate wave energy. The electron temperature orthogonal to the simulation direction is unchanged. This thermal anisotropy of the electrons causes the growth of waves by the Weibel instability with a wavevector that is orthogonal to the direction, along which the electrons are hot (Weibel 1959; Morse & Nielson 1971). It is suppressed here by the one-dimensional geometry that excludes such wavevectors. Note that the thermal anisotropy is due to the saturation of the BI and the Weibel instability is thus a secondary instability. It differs from the beam-driven filamentation instability, which destabilizes the waves with $Z_i \ll 1$ in Figure 1 already at $t = 0$.

Figure 5 plots the $E_x$, the density of the proton beam 2, and that of the electrons at the same location and time as the phase-space distributions have been sampled in Figure 4. The electric field at $x/\lambda_u \approx 9$ is attractive for the protons and repels the electrons, which is in agreement with the observed phase-space

Figure 4. Phase-space distribution of the electrons (a) and of the protons of beam 2 (b) at $t \Omega_e = 10^3$; the color scale shows the 10-logarithm of the number of CPs. The electrons display a plateau distribution for $x/\lambda_u > 10$ and vortices for $x/\lambda_u < 8$. The plateau is split at $x/\lambda_u \approx 9$. The proton beam shows tenuous filaments, which are protons trapped in the ion acoustic waves. A phase-space hole is forming at $x/\lambda_u \approx 9$.

(A color version of this figure is available in the online journal. Mpeg animations of this figure are available in the online journal.)

Figure 5. Normalized electrostatic field $eE_x/(\Omega_e c m_e)$ (a), the proton density of beam 2 (b), and the electron density (c). The densities are normalized to their mean. The electric field shows a bipolar oscillation at $x/\lambda_u \approx 9$, which modulates the densities of the protons and electrons. The time is $t \Omega_e = 10^3$. 

Figure 6. Phase-space distributions at $t \Omega_e = 2000$ of the electrons (a), the protons of beam 2 (b), and of beam 1 (c). The color scale is the 10-logarithm of the number of CPs. The electron distribution and the protons of beam 1 evidence a double layer in the interval $15 < x / \lambda_u < 20$. The protons of beam 1 correspond to the untrapped ions of the double layer. The protons of beam 2 would be the trapped ions, but here they form a pulse.

(A color version of this figure is available in the online journal.)

structure in Figure 4. The electric field amplitude at this position is not unusually large, while the plasma density depletion is. The oxygen density is unaffected at this time, as we show below.

The component energy densities in Figure 2 demonstrate that the equilibrium between $E_F(t)$ and $E_1(t)$ breaks at $t \approx 1500$ and that a secondary instability sets in after that time. An examination of the full phase-space data shows that this secondary instability is spatially localized. Two similar but spatially well separated structures develop in the phase-space distribution of the protons of beam 2 in the box. That two structures rather than one grow suggests that their growth is not accidental. Figure 6 displays the larger of these two structures at $t = 2000$. A different subinterval of the box than in Figure 4 is investigated, since the proton phase-space structures convect with $v_b$. We refer to this proton structure as a pulse. This pulse constitutes a density depletion (Infeld et al. 1989; Infeld & Rowlands 1990). The supplementary Movie 2 animates in time the development of the pulse out of the proton phase-space hole. The movie starts from the proton distribution shown in Figure 4(b). The color scale shows the 10-logarithm of the number of CPs. Movie 2 displays absolute box coordinates and its reference frame moves with $v_b$ through the box. It demonstrates that the proton phase-space hole is stationary for a long time, before it develops into a rapidly growing pulse. It is this pulse that is responsible for the energy changes in Figure 2. We have to emphasize here that the probability for the development of a pulse out of the ensemble of proton phase-space holes is not high. Our box length holds $\approx 10^2$–$10^3$ filamentary structures to which we refer as proton phase-space holes. This is because their size is comparable to or somewhat larger than that of the initial $\approx 10^3$ electron phase-space holes, as we can see from Figure 4. Two pulses grow out of this ensemble of proton phase-space holes. We may not observe such pulses in shorter simulation boxes.

The electron phase-space distribution reveals in Figure 6 a double layer. This double layer accelerates on $5 \lambda_u$ or $12.5 \lambda_e$ a beam of electrons out of the bulk to about 7.5 MeV. This is the characteristic spatial scale of a double layer (Smith 1982). Movie 1 demonstrates that the mean energy of the electron beam grows in time. It is thus not a stationary double layer but one with a floating potential. The beam temperature remains unchanged. The energy the accelerating electrons gain far exceeds the thermal energy of the electrons, even after the BI has heated the electrons. The beam of accelerated electrons then interacts with the bulk electrons through a two-stream instability, which is responsible for the beam momentum oscillations in the interval $x / \lambda_u < 15$ in Figure 6. Such a secondary two-stream instability has also been found by Newman et al. (2001), although the electron beam was much slower there. The electric field of the double layer is strong enough to decrease the mean speed mod ulus of the proton beam 1, which supplies the energy for the electron acceleration.

The accelerating electrons and the protons of beam 1 correspond to the freely streaming particle species in the double-layer picture. The trapped electron species are shown in Figure 6(a) at $x / \lambda_u < 15$ and at low $|p_x|$. The protons of beam 2 should represent the trapped species, according to the usual phase-space structure of a double layer (Smith 1982; Raadu & Rasmussen 1988). Here they are not trapped and they form a pulse instead. The difference in the velocities is also much larger for the proton beams in the double layer in Figure 6 than the typical speed gap between trapped and untrapped protons. The Bohm criterion discussed, for example, by Smith (1982) compares the thermal speeds of the species to their relative drift speed and a strong double layer requires a drift that is somewhat faster than the thermal speed. The double layer in our simulation moves with beam 2. The drift speed of the protons of beam 1 with respect to the double layer exceeds their thermal speed by the extreme
factor $>10^2$. The huge directed proton flow energy enables the immense electron acceleration. Movie 2 is stopped when the protons of beams 1 and 2 start to mix by the formation of an electrostatic shock. Movie 1 is stopped when the electron beam energy exceeds significantly the one resolved by the movie window.

The electron distribution at $t = 2600$ is illustrated in Figure 7. The double layer has not expanded in space, but the electrons are now accelerated to 50 MeV in the box frame. This energy is exceeding even the 20 MeV energy of the protons in the same reference frame and it is comparable to the kinetic energy stored in the drift motion of both proton beams. The total density of the ultrarelativistic beam of untrapped electrons in Figure 7(b) is about 40% of the density of the bulk (trapped) electrons in Figure 7(c). The momentum distribution of the bulk electrons is practically identical to that in Figure 4. The most energetic electrons in the trapped electron population reach an energy of about 2 MeV and the bulk has less than 1 MeV. The kinetic energy the beam electrons have gained through the double-layer potential exceeds this thermal energy by almost 2 orders of magnitude. The ultrarelativistic electron beam would drive the unresolved and faster growing oblique mode instability (Bret 2006), rather than the two-stream instability that is evidenced by Movie 1. The type of instability should, however, not be important for the evolution of the electron double layer because it forms behind it.

Figure 8 displays the phase-space distributions of both proton beams in the same spatial interval as Figure 7. The proton distributions of both beams have merged and a large electrostatic shock forms at $x/\lambda_u \approx 55$. A second electrostatic shock involving only the protons of beam 2 occurs at $x/\lambda_u \approx 92$. Movie 2 illustrates that the second shock is driven by the rapid expansion of the proton density pulse. It is remarkable that some protons of beam 2 have increased their $p_z$ momentum by a factor of 2.6 through the electrostatic field, which is also shown in Figure 8. The electron double layer is thus also a proton accelerator. The electrostatic field of the double layer does not change its sign, which contrasts the bipolar fields of phase-space holes, like in Figure 5.

Figure 9 displays the density distribution of the oxygen ions at three times, $t = 1000$, $1600$, and $2000$. The frame is shifted with beam 2 and the proton phase-space hole is located at $(x - v_b t)/\lambda_u \approx 20$ in the chosen frame. The oxygen depletion at $t = 10^3$ is comparable to the statistical fluctuations of the number of CPs. The integration of the density over three grid cells reduces these statistical fluctuations by a factor of $\sqrt{3}$, but keeps the density depletion value unchanged. The proton phase-space hole grows into the pulse at $t \approx 1500$. Figure 9(b) reveals that the oxygen density modulation is significant at around this time. The sudden growth of the proton phase-space hole may be a consequence of this oxygen depletion or the oxygen depletion...
could be caused by the large electric fields of the pulse. The proton pulse and the oxygen depletion are intertwined. As the proton pulse grows, the structure and amplitude of the oxygen depletion remain qualitatively unchanged, but the depletion expands in space. This may reflect the observation from Movie 2 that the shape of the proton pulse is practically unchanged until \( t \approx 2000 \).

4. DISCUSSION

Our motivation has been to identify a process that could accelerate electrons in the foreshock of an SNR shock to energies, with which they can undergo the diffusive shock acceleration to ultrarelativistic energies. The radio synchrotron emissions of SNR shocks evidence the presence of such electrons (Uchiyama et al. 2007) but their origin remains unclear. The electrons are injected into the diffusive shock acceleration, if their kinetic energies are comparable to the typical proton kinetic energies (Reynolds 2008; Hillas 2005). It is thought that the electrons are accelerated by instabilities, which are driven by ion beams in the foreshock region of SNR shocks (Cargill & Papadopoulos 1988). The source mechanism of the ion beams is either the specular reflection of upstream ions by the shock or the leaking of downstream ions into the upstream plasma. The BI (Buneman 1958) invoked by Cargill & Papadopoulos (1988) is, however, usually too weak to inject the electrons. It can only transfer a few percent of the ion energy to the electrons.

The work presented here examines the secondary instabilities triggered by the BI with a PIC code (Dawson 1983; Eastwood 1991). The initial conditions of the simulation have been based on the foreshock model employed in several related papers (Shimada & Hoshino 2003, 2004; Dieckmann et al. 2007, 2008a). We have considered two unmagnetized, interpenetrating, uniform, and unbounded beams that have been composed of three particle species. Beam 1 has contained protons and electrons. Beam 1 represents in our model the ambient plasma ahead of the expanding SNR shock. Beam 2 has been formed by protons and oxygen with a single negative charge. It is composed in our model of the downstream ions that have leaked through the shock or of the ions that have been reflected by the shock. The oxygen ions we have modeled can be found in the SNR ejecta, although not in the high density and in the ionization number we have assumed here (Katsuda et al. 2008). Our negative charge is more representative of dust, which we also find in supernova blast shells (Meikle et al. 2006). The simulation box has been placed in the foreshock region in which we expect to find such beams. We have, however, not resolved the shock and our model is thus not self-consistent. Each beam has been charge and current neutral. The relative speed between both beams has been set to 0.4c. The main blast shell of an SNR shock moves with speeds <0.2c (Vink et al. 2006; Fransson et al. 2002) and the beam speed would be too high if ion leaking or specular reflection would be the source of the ions. Precursor shocks may, however, be faster and could give rise to such a beam speed. Subshells of the SNR ejecta with speeds up to 0.9c have been observed at particularly violent supernovae (Kulkarni et al. 1998). The initial temperature of all species has been set to 9.1 keV.

One spatial x-direction has been resolved by the PIC simulation, which implies that the wave spectrum driven by the BI is monodirectional. The solution of the linear dispersion relation has shown that the exponential growth rate of the waves is largest along the resolved beam flow direction. The filamentation modes (Lee & Lampe 1973), that are important for relativistic plasma collisions, have been negligible. However, the growth rate of the modes with wavevectors that are oblique to the beam velocity vector grow almost as fast as the BI and they are important. Even if the BI would be dominant, it would couple to oblique wave modes by means of weak turbulence (Ziebell et al. 2008). The restriction to one dimension thus limits the realism of the simulation, but it first allows us to resolve all relevant spatiotemporal scales and, second, it facilitates the identification of important physical processes, which can then be examined in more realistic settings by future simulation studies. The large box has ensured that the periodic boundary conditions have not affected the results and that the secondary processes could develop with a reasonable probability.

The BI has grown in the simulation box and has saturated by the trapping of electrons (Roberts & Berk 1967), which groups the electrons into phase-space holes that give rise to charge density modulations. The latter can behave as quasi-particles and are thus referred to as plasmons. We could match the mean (group) speed of these plasmons with that of the protons of beam 2 by selecting negatively charged oxygen rather than electrons. Such a resonance facilitates the coupling of the plasmon energy into the ion acoustic wave mode (Mendonca & Bingham 2002; Mondaca et al. 2005). Eventually the electron distribution thermalized and a dissipative equilibrium developed between the electric fields and the electrons (Korn & Schamel 1996a, 1996b). The ion beams had lost only a small fraction of their energy, as in previous works (Shimada & Hoshino 2003, 2004; Dieckmann et al. 2007, 2008a); the electrons moved consequently only at the mildly relativistic speeds that are not sufficient to inject them into the diffusive shock acceleration.

The BI has accelerated the electrons along the beam flow direction but it has left unchanged the temperature orthogonal to it. The resulting thermal anisotropy of the electron distribution would trigger an instability similar to that proposed by Weibel (1959) and modeled numerically by Morse & Nielson (1971), resulting in the growth of magnetic fields and multidimensional field structures that cannot be represented by a one-dimensional PIC simulation. The Weibel instability is a secondary instability that develops after the BI has saturated. It has available the energy stored in the thermal anisotropy of the electron distribution, which is small compared to the ion beam energy. This secondary instability is thus probably not a strong electron accelerator. We leave its investigation to future work.

The electronic plasmons have pumped the ion acoustic mode until its fields became strong enough to interact nonlinearly with the protons of beam 2. This interaction has resulted in the formation of density filaments in phase space that have been accelerated out of the proton beam. We refer to these filamentary structures as phase-space holes, although they have not yet formed in full. As the time progressed, two large density rarefaction pulses (Infeld et al. 1989; Infeld & Rowlands 1990) have grown out of the proton phase-space holes, which are moving with beam 2. We have examined in more detail the larger pulse. We have identified two potential causes for its emergence. The animation of the electron phase-space distribution close to this proton structure has revealed a vortex field. The electron flow becomes turbulent as it passes the proton phase-space hole and the electron current is disrupted, which can lead to the growth of phase-space holes (Dupree 1986) and to the development of double layers (Smith 1982). The long time between the formation of the proton phase-space hole and of the pulse has suggested also a second possible cause, namely the
oxygen response. We have found that the growth of the proton density pulse coincided with the displacement of a substantial fraction of the oxygen ions. It is unclear though if the oxygen displacement is the cause or the consequence of the emergence of the pulse. A potentially related effect has been observed by Eliasson & Shukla (2004). There it was shown that an electron phase-space hole was ejected by a density depletion in the background ions, which its electric fields excavated.

The growth of the proton density pulse has resulted in a double layer that accelerated the electrons to ultrarelativistic speeds. Double layers are particularly efficient accelerators because they can convert directed flow energy from the protons to the electrons, without heating up the plasma. The electrons have reached a peak energy of 50 MeV, which is comparable to the energy stored in the relative drift motion of both proton beams. The accelerated electrons have interacted with the bulk of the electrons and nonrelativistic for the protons and it is thus a semirelativistic one. Apparently, the different scaling properties of relativistic and nonrelativistic double layers (Carlqvist 1982) do not exclude their simultaneous presence. The potential is a floating one and the system is initially current neutral.

It is clear that our restriction to one spatial dimension, the beam speed that is relatively high, and the usage of a dense, negatively charged ion species limits the realism of the simulation and the direct applicability of our results to SNR shock acceleration. Future simulation work must thus address several open issues. First, what is the minimum speed for the beams? This has to be done in the form of parametric simulation studies because it is not clear for which combinations of the beam speeds and temperatures the double layers can form. Existing analytic studies (Verheest & Pillay 2008) do not cover the case we consider here. The work presented here indicates that the double layers may form with a relatively small probability out of the proton structures, which has to be taken into account for further studies. The effects of the obliquely propagating modes driven by the initial beam instability have to be examined, as well as those due to the Weibel instability, which we would get after the initial BI has saturated. This requires at least two-dimensional PIC simulations. Even multidimensional PIC simulations would not correctly reproduce the plasma dynamics. The massive electron acceleration will result in Weibel and filamentation instabilities and, thus, in strong electromagnetic fields. The electrons would emit synchrotron radiation and electrostatic bremsstrahlung (Schlickeiser 2003; Fleishman & Toptygin 2007). The balance of these processes (Schlickeiser & Lerche 2007) requires a method to deal with the radiation. Finally, other ion species and plasma compositions have to be examined to assess how easily double layers can form and to identify whether it is the electron current disruption or the ion displacement that destabilizes the proton phase-space holes.

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