Formation of a localized acceleration potential during magnetic reconnection with a guide field

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Magnetic reconnection is a basic plasma physics process that allows magnetic field lines to change topology in the presence of a plasma. Although it is localized it controls large scale explosive plasma phenomena such as coronal mass ejections, magnetic storms in the Earth’s magnetotail and internal disruptions in magnetic fusion devices.1,2

During reconnection on the sun, the foot points where reconnecting magnetic loops intersect the photosphere light up in x rays. In the Earth’s magnetotail, energetic electrons up to 300 keV have recently been observed by the Wind spacecraft.3 This is conclusive evidence that reconnection is associated with a significant production of energetic electrons. It has been proposed that the electrons are energized by direct acceleration in the reconnection electric field along the magnetic X-lines.4,5

However, numerical simulations of reconnection provide significant information about the physics and structure of the reconnection layer, which is inconsistent with the direct acceleration models. In particular particle in cell (PIC) simulation is a useful tool for such investigations because it includes the kinetic effects of both the electrons and the ions. Here we introduce the acceleration potential, \( \Phi_0 \), as a measure of the work done by parallel electric fields on electrons streaming along magnetic field lines. This provides a new and more complete framework for understanding the energization of electrons by direct acceleration in guide field reconnection. We define the reconnection region as the area where two fluid effects are important, extending about 50 \( d_e \), upstream of the X-line and about 200 \( d_e \), downstream of the X-line. In this region a localized structure in \( \Phi_0 \) is documented, caused by a sign reversal in \( E_i \). The analysis shows how this localized structure limits direct acceleration and causes extensive trapping of thermal electrons.

The electron dynamics and energization are studied based on profiles of an open boundary PIC simulation including 2 \times 10^9 particles in a domain \( xy=3072 \times 3072 \) cells=569 \times 569 c/\( \omega_{pe} \). The simulation is translationally symmetric in the z-direction and is characterized by the following parameters: \( m_i/m_e=360 \), \( T_i/T_e=2 \), \( B_{\text{guide}}=0.5B_0 \), \( \omega_{pe}/\omega_{ce}=2.0 \), background density=0.30\( n_0 \) (peak Harris density), \( v_{th,e}/c=0.20 \), and an initial current sheet width of about 0.4 \( d_e \).

We consider the fields observed at the single time slice of Fig. 1, during a time interval without significant evolution in the profiles. As seen in Fig. 1 the asymmetric current profile is not aligned with the superimposed contours of magnetic flux. This type of current profile is typical of kinetic simulations of reconnection with a guide magnetic field. The enhancement in the plasma density is also asymmetric, exhibiting a quadrupolar structure similar to that discussed in Refs. 7 and 8. In the inflow region the electrons are isotropic with \( T_i=T_e \). However, as the electrons enter the reconnection region \( T_i \) is increased (by a factor of about 4), while \( T_e \) is reduced proportionately to the decrease in the magnetic field, \( B \). In turn, the profile of \( B \) is strongly influenced by the strong quadrupolar structure in \( B_z \). The inductive electric field, \( E_z \), is nearly uniform throughout the ion diffusion region and the in-plane potential includes the familiar quadrupolar structure required for having \( E \cdot B \sim 0 \) outside the reconnection region.

In Ref. 9 the details of the electron distributions observed by the Wind spacecraft are accounted for by hypothesizing the trapping of electrons by an in-plane electrostatic potential. The simulation presented here (with \( B_{\text{guide}}/B_0 \) and \( T_i/T_e \) relevant to the Wind event) confirms that trapping is important. However, the trapping is not caused only by an electrostatic potential, but also by the inductive electric field, \( E_z \), in combination with the details of the magnetic field geometry.

To explore the role of \( E_i \) we introduce the parallel acceleration potential \( \Phi_0 \) as a “pseudopotential” that measures the work done by \( E_i \) on a passing electron as it escapes the
simulation box in a straight shot along a magnetic field line. Therefore, mathematically $\Phi_1$ is defined as

$$\Phi_1(x) = \int_{x}^{\infty} \mathbf{E} \cdot d\mathbf{l}, \tag{1}$$

where the integration is carried out from the point $x$ along the magnetic field to the boundary of the simulation box. Note that in contrast with the electrostatic potential $\Phi$ in Fig. 1(h), $\Phi_1$, in addition to the in-plane electric fields, includes important contributions from the inductive electric field, $E_z$.

We find that $\Phi_1$ is nearly independent of whether $d\mathbf{l}$ is taken parallel or antiparallel to $\mathbf{B}$. This is because the integral of $\mathbf{E} \cdot d\mathbf{l}$ over its full length in the simulation domain is small, $\int_{x}^{\infty} \mathbf{E} \cdot d\mathbf{l} < T_e/e$. This result can be understood on the basis of the adiabatic fluid model of Ref. 10, which shows that $\Phi_1 = \Phi_1(n, B)$ is a function of only $n$ and $B$ and increases monotonically with $n$ and $1/B$. As a static magnetic field line is followed through the reconnection region, its two ends outside the region will have similar values of $n$ and $B$. The values of $\Phi_1$ outside the region will therefore also be similar. The localized maximum of $\Phi_1$ occurs because of the localized maxima of both $n$ and $1/B$ in the region. This, in turn, dictates the sign reversal in $E_z$, which develops to boost the electron density in the regions of high ion density and low magnetic field strength.

In Fig. 2(a) contours of constant $e\Phi_1$ are shown. Outside the reconnection region the contributions to $\Phi_1$ from $E_z$ and the in-plane electric fields cancel reflecting that $E_z \sim 0$. However, as seen in Fig. 3(a), inside the region $E_z$ is finite and changes sign where $n$ and $1/B$ peak. Thus, consistent with the discussion above the largest values of $\Phi_1$ coincide with the regions of density enhancement and small $B$; here $e\Phi_1$ reaches an amplitude of about $5T_e$. Considering a number of different simulations (not included), we find similar localized structures for $\Phi_1$.

In the reconnection region where the density is enhanced nearly all electrons bounce multiple times due to the localized structure and large amplitude ($> T_e/e$) of $\Phi_1$. The
FIG. 3. (Color online) (a) Contours of constant values of $E_\parallel$, (b) zoomed-in view of $\Phi_\parallel$ for the area indicated in Fig. 2(a). (c) Number of electrons with $\varepsilon > 0.5 m_\text{e} c^2$ evaluated on a grid with $d_x$ resolution. (d) Energy distributions for the location marked by “O” in (c). The dashed lines represent $f(\varepsilon) = f_d(\varepsilon)$ and $f(\varepsilon) = f_d(\varepsilon - 5 T_e)$, respectively. Note that while the simulation is fully relativistic, it applies that $f_d > \exp(-\varepsilon/\gamma_{\text{th}}^2)$, where $\gamma$ is the particle velocity and $v_{\text{th}}$ is the thermal velocity.

kinetic properties of these electrons are analyzed by considering the action integral

$$J = \oint v_x dl = \oint \sqrt{\frac{2 E_\parallel}{m}} dl.$$  \hspace{1cm} (2)

This integral is over one particle bounce motion, where $l$ represents the length measured along the orbit and $E_\parallel$ is the parallel kinetic energy. $J$ is an adiabatic invariant and is conserved if the bounce motion is sufficiently rapid. Following the notation of Ref. 11, $J$ is expanded in a series of increasing powers of $\varepsilon = m_e c$ giving $J = J_0 + \varepsilon J_1 + \ldots$. In Ref. 11 it is shown that $J_0$ oscillates and vanishes at the orbit bounce points. Furthermore, $J_0$ is obtained simply by integration along field lines (not the actual electron orbits). Let $x_0$ be a bounce point ($E_\parallel = 0$) of an electron. The parallel kinetic energy at a point $x$ on the same field line can then be calculated as

$$E_\parallel(x, x_0) = e[\Phi_\parallel(x) - \Phi_\parallel(x_0)] + \mu [B(x_0) - B(x)].$$  \hspace{1cm} (3)

It is here assumed that $\mu$ is conserved, which will be justified below. The action integral can then be expressed as

$$J_0(x_0) = \oint \sqrt{\frac{2 E_\parallel(x', x_0)}{m}} dl'.$$  \hspace{1cm} (4)

For a given bounce point, $x_0$, the points of integration, $x'$, are those on the field line where $E_\parallel(x', x_0) > 0$.

In Fig. 2(b) contours of constant $J_0(x_0)$ are shown for $x_0$ represented by the points in the $xy$-plane. The contours coincide with loci of the electron bounce points and therefore provide a direct visualization of how the trapped trajectories are channeled by the ion diffusion region. The asymmetry in $\Phi_\parallel$ causes the electrons entering the upper (lower) inflow region to exit through the left (right) outflow region generally without making direct contact with the magnetic X-line. Contours for one particular value of $J_0$ are highlighted by the white lines. An electron orbit corresponding to this value of $J_0$ (and $\mu = T_e / B_{\text{line}}$) is shown in yellow. Inside the reconnection region the bounce points coincide with the white contours of $J_0$ and we conclude that $J_0$ is conserved in this region. In turn, given $J_0$ is obtained based on $E_\parallel$ in Eq. (3), this emphasizes the importance of $\Phi_\parallel$ to the motion of the trapped electrons.

In Fig. 2(b) $J_0$ is normalized such that in the areas along the separator where $J_0 > 1$ the level of trapping is insignificant and free streaming electrons can carry parallel currents.1,12 Meanwhile, in the extended areas where $J_0 < 1$ the dynamics of the trapped electrons dominate. To evaluate the heating of the trapped electrons we consider the adiabatic invariants, $\mu$ and $J_0$. From conservation of $\mu = E_\parallel / B$ it is clear that the perpendicular particle energy $E_\perp$ must be reduced in the inner reconnection region where the lowest values of $B$ are observed. A comparison of Figs. 1(d) and 1(e) clearly shows that $T_e$ is proportional to $B$ (consistent with $\mu$-conservation for most electrons). The evolution of the parallel kinetic energy is controlled by $J_0$. As is evident in Fig. 2(d), when the trapped electrons enter the reconnection region their bounce lengths (measured along field lines) are reduced. As a consequence [see Fig. 2(e)] the parallel velocity is increased in order for $J_0$ to be conserved. Consistently, in Fig. 1(c) we observed that $T_e$ increases by about a factor of 4 inside the current channel.

The acceleration potential is also important for the energy balance of the passing electrons. In the limit where the electrons strictly follow the field lines, $\Phi_\parallel(x)$ is a measure of the gain in energy of the passing electrons. Because $B_z = B_x = 0$ at the $x$-line, it follows that $|v_x|/|v_y| \rightarrow B_\perp / \sqrt{B_t^2 + B_\perp^2} \rightarrow \infty$ at this point. Thus, $\Phi_\parallel$ as defined in Eq. (1) has a logarithmic singularity which can be observed in the zoomed-in view of $\Phi_\parallel$ given in Fig. 3(b). In a narrow region around the singularity less than $d_x / 4$ wide, we have $e \Phi_\parallel > 6 T_e$ along two of the four separators, and electrons can therefore acquire large energies by direct acceleration.13 However, because the singularity is only logarithmic and very narrow the number of electrons accelerated here is small.

Excluding the singular region we find that $\int_{\infty}^\infty E_x \cdot dl < T_e / e$, which suggests that electrons will gain very little energy in a single pass through the reconnection region. However, because the reconnecting field lines are “moving,” the passing electrons do not strictly follow the static field lines. This allows a significant group of passing electrons to be accelerated as they enter the reconnection region, while avoiding the decelerating electric field as they leave the region. The electron trajectories in Fig. 3(a) are representative for this heating mechanism. The locations in the simulation domain where the heating mechanism is effective are shown in Fig. 3(c), which contains a color map of the number of electrons recorded in $1 d_x \times 1 d_y$ cells with energies $E_\parallel > 0.5 m_e c^2$. The combination of the sign reversal of $E_\parallel$ and the drift of the field lines causes energization of the electrons along all four separators (rather than just the two separators of the logarithmic singularity).

The passing electrons that enter the simulation domain at optimal locations will acquire an energy gain on the order of...
$e\Phi_{\text{max}} \sim 5T_e$, where $\Phi_{\text{max}}$ is the maximal value of $\Phi_t$ in the bulk of the ion diffusion region. Thus, for any point in the simulation domain, the kinetic energy of a passing electron is characterized by

$$\mathcal{E}(x) \leq \mathcal{E}_0 + e\Phi_{\text{max}},$$

where $\mathcal{E}_0$ is the kinetic energy the electron had when it entered the simulation at the domain boundary. This can be tested directly against the distribution of electrons in the simulation. According to Liouville’s theorem, $df/dt = 0$, the phase space density, $f$ along a particle orbit is conserved such that $f(\mathcal{E}(x)) = f_0(\mathcal{E}_0)$, where $f_0$ is the distribution function applied at the simulation boundary. Thus, since $\mathcal{E}_0 \sim E(x) - e\Phi_{\text{max}}$ and $df_0/d\mathcal{E} < 0$ we find that

$$f(\mathcal{E}) \leq f_0(\mathcal{E} - e\Phi_{\text{max}}).$$

In Fig. 3(d) the full line represents $f(\mathcal{E})$ obtained directly from the particle data within the small region $(3d_i \times 3d_i)$ marked by the circle in Fig. 3(c). The two dashed theoretical curves are $f_0(\mathcal{E})$ and $f_0(\mathcal{E} - 5T_e)$, respectively. Consistent with Eq. (6), we verified that throughout the computational domain these two curves bracket the simulated $f(\mathcal{E})$, as expected from Eq. (6).

Although a significant relativistic population is present in Fig. 3(d) this cannot account for observations of fast electrons in nature. Rather, the relativistic electrons in the simulation are just a result of the large value of $v_{\parallel}/c \approx 0.2$ (typical for PIC-simulations) combined with the energy boost of $e\Phi_t \sim 5T_e$. For example, a boost of $5T_e$ for the electrons observed in the Wind event, would only increase the temperature to 2 keV and is not consistent with the observation of electron acceleration up to 300 keV. Other possible mechanisms, such as Fermi acceleration or perhaps three-dimensional effects in reconnection, must therefore be explored to account for electron heating. Interestingly, observations consistent with Fermi acceleration have recently been reported for an event where the Cluster spacecraft encountered a series of magnetic islands in the Earth’s magnetosphere.

The formation of a localized acceleration potential and the described physics of the trapped electrons are not dependent on the choice of boundary conditions. This is evident from Fig. 4, which illustrates that $\Phi_t$ in a 3D code simulation with periodic boundaries is similar to $\Phi_t$ in Fig. 3(a).

In summary, the documented sign reversal in $E_v$ leads to the formation of a localized acceleration potential, $\Phi_t$, which controls the energization of both trapped and passing electrons. In contrast with models where the in-plane electric fields are neglected, we find that single pass energization is limited and is bounded by $e\Phi_{\text{max}} \sim 5T_e$ (per electron). However, given the sign reversal in $E_v$ at this level of energization is possible along all four separators.

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