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Particle simulation study of driven magnetic reconnection in a collisionless plasma

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Driven magnetic reconnection in a collisionless plasma, "collisionless driven reconnection," is investigated by means of two-and-one-half-dimensional particle simulation. Magnetic reconnection develops in two steps, i.e., slow reconnection, which takes place in the early stage of the compression when the current layer is compressed as thin as the orbit amplitude of an ion meandering motion (ion current layer), and subsequent fast reconnection, which takes place in the late stage when the electron current is concentrated into the narrow region with a spatial scale comparable to the orbit amplitude of an electron meandering motion (electron current layer). The global dynamic evolution of magnetic reconnection is controlled by the physics of the ion current layer. The maximum reconnection rate is roughly in proportion to the driving electric field. It is also found that both ion heating and electron heating take place in accordance with the formation of two current layers and the ion temperature becomes two or more times as high as the electron temperature.

I. INTRODUCTION

Magnetohydrodynamic (MHD) studies[1,2] have disclosed that driven magnetic reconnection plays an essential role on the energy relaxation and the self-organization of a magnetically confined plasma. Under the influence of a driving source and a small amount of electrical resistivity, magnetic reconnection takes place in a time scale much shorter than the resistive time scale, and the reconnection rate is primarily determined by the driving electric field.[3] This process can lead to fast energy conversion from the field energy to the particle energy, as well as a topological change of magnetic field.[1–3] On the other hand, energetically active phenomena[4] triggered by magnetic reconnection are often observed in a high-temperature, rarefied plasma, in which binary collisions between particles are negligible, namely, in a collisionless plasma. It is not so easy to explain how an electric field along the equilibrium current is generated in the neutral sheet of collisionless plasma. The concept of an anomalous resistivity,[5] which originates from the wave–particle interaction or the stochasticity of particle orbit,[6] has been introduced to explain collisionless reconnection. It is widely considered that tearing instability[7–17] is an effective mechanism for collisionless magnetic reconnection. Particle simulation[13–17] has been applied to investigate the collisionless process of magnetic reconnection for a variety of magnetic configurations. Leboeuf *et al.*[14] discussed the collisionless magnetic reconnection by means of two-dimensional particle simulation when a small amount of resistivity due to particle collisions was introduced in a periodic simulation domain. They pointed out the importance of the finite Larmor radius effect in current penetration. Hewett *et al.*[15] also examined collisionless magnetic reconnection by using an implicit particle simulation code for the case where there existed no external driving force, and the initial equilibrium had a current layer the width of which was as thin as the electron skin depth. They found that the time evolution was strongly dependent on the ratio of ion to electron mass.

When a system is subjected to an external driving source, or it is macroscopically unstable, it evolves dynamically with time,[3] and thus, the physical parameters, which characterize the state of a system, become a function of time. In this paper we report the simulation of collisionless magnetic reconnection for the case where a plasma is initially in an MHD equilibrium, and a driving plasma flow is supplied into a system through the boundary of a simulation domain. Then, the current layer is compressed by the convergent plasma flow, and its spatial scale length changes with time. After some period the system dynamically evolves into a kinetic regime in which the finite Larmor radius effect becomes active. In considering such a dynamical evolution, it is important to solve self-consistently the nonlinear phenomena, including electron dynamics, ion dynamics, and the dynamical evolution of the electromagnetic field. For this purpose, we have developed an electromagnetic particle simulation code,[18–20] which can describe not only a finite ion Larmor radius effect, but also a finite electron Larmor radius effect in an open system. The purpose of this paper is then to demonstrate the dynamical evolution of driven magnetic reconnection in a collisionless plasma and to clarify the fine mechanism leading to collisionless driven magnetic reconnection. The initial condition and the simulation model are described in Sec. II. We show that two different types of magnetic reconnection take place as a system evolves dynamically. In order to investigate the physical mechanism, the dependence of the simulation results on the ion mass, the electron mass, and the driving electric field is examined by carrying out several simulation runs. The detailed simulation results are discussed in Sec. III. The detailed analysis of the simulation results reveals that the temporal evolution of magnetic reconnection scales as $E_0^{-1/2}M_i^{1/4}$, and the maximum reconnection rate is roughly in proportion to E_0, where E_0 is the driving electric field at the boundary and M_i is the ion mass. Finally, we give a summary of this paper and a physical meaning of the scaling law in collisionless driven magnetic reconnection, in connection with the dynamics of

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the current layer in Sec. IV. Comparisons of the simulation results with theoretical studies of the tearing instability are also given in Sec. IV.

II. SIMULATION MODEL

Let us study externally driven magnetic reconnection of a collisionless plasma by means of two-and-one-half-dimensional particle simulation. The semi-implicit method[18–20] is used for time advancing. Because both electrons and ions are treated as particles in this method, a finite electron Larmor radius effect can be described, as well as a finite ion Larmor radius effect. The physical quantities are assumed to be translationally symmetric in the z direction ($\partial/\partial z=0$). The basic equations to be solved are the equations of motion,

$$\frac{d(\gamma_j \mathbf{v}_j)}{dt}=\frac{q_j}{m_j}\left(\mathbf{E}+\frac{\mathbf{v}_j}{c}\times\mathbf{B}\right), \tag{1}$$

$$\frac{d\mathbf{x}_j}{dt}=\mathbf{v}_j, \tag{2}$$

and the Maxwell equations,

$$\frac{1}{c}\frac{\partial \mathbf{B}}{\partial t}=-\nabla\times\mathbf{E}, \tag{3}$$

$$\frac{1}{c}\frac{\partial \mathbf{E}}{\partial t}=\nabla\times\mathbf{B}-4\pi\mathbf{j}, \tag{4}$$

$$\nabla\cdot\mathbf{B}=0, \tag{5}$$

$$\nabla\cdot\mathbf{E}=4\pi\rho, \tag{6}$$

where $\mathbf{x}_j(t)$, $\mathbf{v}_j(t)$, m_j, and q_j are the position, the velocity, the rest mass, and the charge of the jth particle, and the relativistic γ factor of the jth particle is defined by

$$\gamma_j=1/\sqrt{1-(\mathbf{v}_j\cdot\mathbf{v}_j)/c^2}. \tag{7}$$

The current density $\mathbf{j}(\mathbf{x},\mathbf{t})$ and the charge density $\rho(\mathbf{x},t)$ are obtained by summing over all the particles, namely,

$$\mathbf{j}(\mathbf{x},\mathbf{t})=\sum_{j=1}^{N}\frac{q_j\mathbf{v}_j(t)}{c}S[\mathbf{x}-\mathbf{x}_j(t)], \tag{8}$$

$$\rho(\mathbf{x},t)=\sum_{j=1}^{N}q_jS[\mathbf{x}-\mathbf{x}_j(t)], \tag{9}$$

where N is the total number of particles, and $S(\mathbf{x})$ is the form function of particles that is expressed by a triangle with the base length equal to 2.0 times the grid separation.

As an initial condition, we adopt a one-dimensional equilibrium with the Harris-type antiparallel magnetic configuration as

$$B_x(y)=B_0\tanh(y/L), \tag{10}$$

$$P(y)=B_0^2/8\pi\,\mathrm{sech}^2(y/L), \tag{11}$$

where L is the scale height along the y axis. There is a magnetically neutral sheet along the mid-horizontal line ($y=0$) in the initial equilibrium. Let us assume that the initial particle distribution is given by a shifted Maxwellian with a spatially constant temperature, and the average particle velocity is equal to the diamagnetic drift velocity. Then, the particle position and the particle velocity are determined from the pressure profile $P(y)$ and the current density $\mathbf{j}(\mathbf{x},0)$ ($=[0,0,-B_0/(4\pi L)\mathrm{sech}^2(y/L)]$). Because both an ion and an electron are loaded at the same spatial position, there is no electric field in the initial profile.

In order to drive magnetic reconnection at the center of the simulation domain, we adopt an input boundary condition[19] at the boundary of the y axis ($y=\pm y_b$) and a periodic boundary condition at the boundary of the x axis ($x=\pm x_b$). At the input boundary, the y derivatives of the field quantities are equal to zero, and the plasma is smoothly supplied with the $\mathbf{E}\times\mathbf{B}$ drift velocity into the simulation domain. The driving electric field $\mathbf{E}_d(x,t)=[0,0,E_{dz}(x,t)]$ at the input boundaries is taken to be zero at $t=0$, and gradually increases for $0<t<t_A$. After this period ($t>t_A$), $E_{dz}(x,t)$ is described by a constant profile as

$$E_{dz}(x,t)=E_0(\epsilon_f+(1-\epsilon_f)\{\cos[\phi(x)]+1\}/2), \tag{12}$$

where

$$\phi(x)=\begin{cases}-\pi, & \text{if } x<-x_{b1},\\ \pi x/x_{b1}, & \text{if } x_{b1}>x>-x_{b1},\\ \pi, & \text{if } x>x_{b1};\end{cases}$$

$\epsilon_f=0.12$, $x_{b1}=x_b$, $t_A=2y_b/v_{A0}$, and v_{A0} is the initial average Alfvén velocity. As a condition at the boundary of the x axis, we have examined two kinds of boundary conditions, i.e., (1) a free boundary condition,[19] where the x derivatives of the field quantities are equal to zero and a particle can pass freely through the boundary; and (2) a periodic boundary condition. It is found *a posteriori* that we have almost the same results for both cases if the simulation box is long enough along the x axis, say, $x_b/y_b\geqslant 2.5$. In this paper, therefore, the ratio of the side lengths of the simulation box x_b/y_b is fixed to 3, and the periodic condition is adopted at the boundary of the x axis.

It is important to examine the numerical accuracy of the simulation model before stepping in the analysis of simulation results. Let us show one example of the numerical tests to find appropriate simulation parameters. Figure 1 shows the temporal evolution of the total magnetic energy for the cases, where the total numbers of particles are equal to 120 000, 180 000, 240 000, and 360 000, respectively. Since these simulation results are found to little depend on the particle number, it is appropriate to fix the total number of particles to 240 000 in order to obtain the simulation result independent of the particle number. The simulation domain is implemented on a (66×129) point grid. The time step width Δt is determined, so as to satisfy the Courant–Friedrichs–Lewy condition[21] for the electromagnetic wave in vacuum; the relation $\Delta t=0.5\ \Delta y/c$ holds in our simulation, where Δy is the grid separation along the y axis.

III. SIMULATION RESULTS

We carry out several simulation runs with different physical parameters. The relation $L>\rho_i>c/\omega_{pe}>\Delta y$ holds for all cases, where ρ_i is the ion Larmor radius and ω_{pe} is the

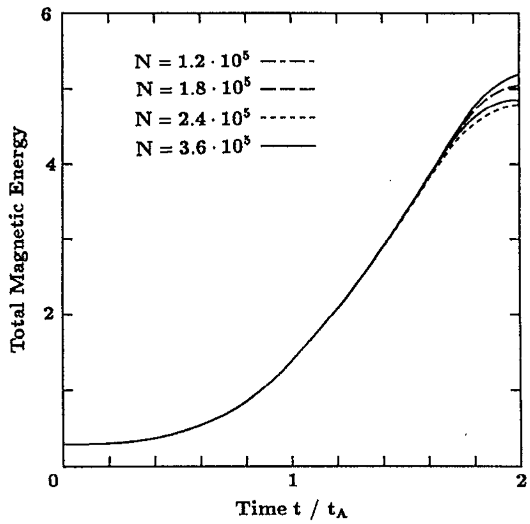

FIG. 1. Temporal evolution of the total magnetic energy for four test runs, where the total numbers of particles are equal to 120 000 (dot–dashed line), 180 000 (dashed line), 240 000 (dotted line), and 360 000 (solid line), respectively.

electron plasma frequency. Let us first discuss the overview of collisionless driven reconnection by using the simulation result for case A, where the ratio of ion to electron mass M_i/M_e is 50, the half-width d_{h0} of the initial mass density profile is $21.9\ \Delta y$, the initial ion Larmor radius ρ_{i0} is $10.4\ \Delta y$, the driving electric field E_0 is $-0.04B_0$, $\omega_{pe}/\omega_{ce}=5.0$, $c/(\omega_{pe}\ \Delta y)=2.0$, $\omega_{ce}\ \Delta t=0.05$, and $\omega_{ce}(=eB_0/M_ec)$ is the electron cyclotron frequency. In this case, the maximum value of the input flow velocity reaches 0.4 of the local Alfvén velocity. Figure 2 shows four snapshots of magnetic flux contours (left) and vector plots of the average ion velocity (right) in the (x,y) plane at $t=0$, $t=1.22t_A$, $t=1.78t_A$, and $t=2.12t_A$, where the magnetic flux contours less than the initial value at the input boundary are plotted with the dotted lines. A magnetically neutral sheet exists along the mid-horizontal line ($y=0$), and no bulk ion flow exists in the (x,y) plane at $t=0$ (top panel). An x-shaped structure of magnetic separatrix becomes visible as a result of magnetic reconnection after the period of $t=1.22t_A$. The region occupied by the reconnected flux spreads over the whole simulation domain at this time. A fast directed flow arises from the X point after magnetic reconnection sets in, and carries the reconnected flux toward the boundaries of the x axis ($x=\pm x_b$). The reconnection point does not always stay at the midpoint of the simulation domain, but it shifts slightly along the magnetically neutral sheet with time. One can find in the bottom panel that a shock structure appears in the ion flow pattern.

A. Two-step evolution of collisionless reconnection

Generation of the electric field along the equilibrium current is needed for driving magnetic reconnection. Figures 3(a) and 3(b) show (a) the temporal evolution of the z component of the electric field, which is the reconnection electric field, and (b) that of the mass density ρ_m (solid line) and the z component of the current density $-j_z$ (dot–dashed line) at the X point for case A. The spatial profile of the electric field

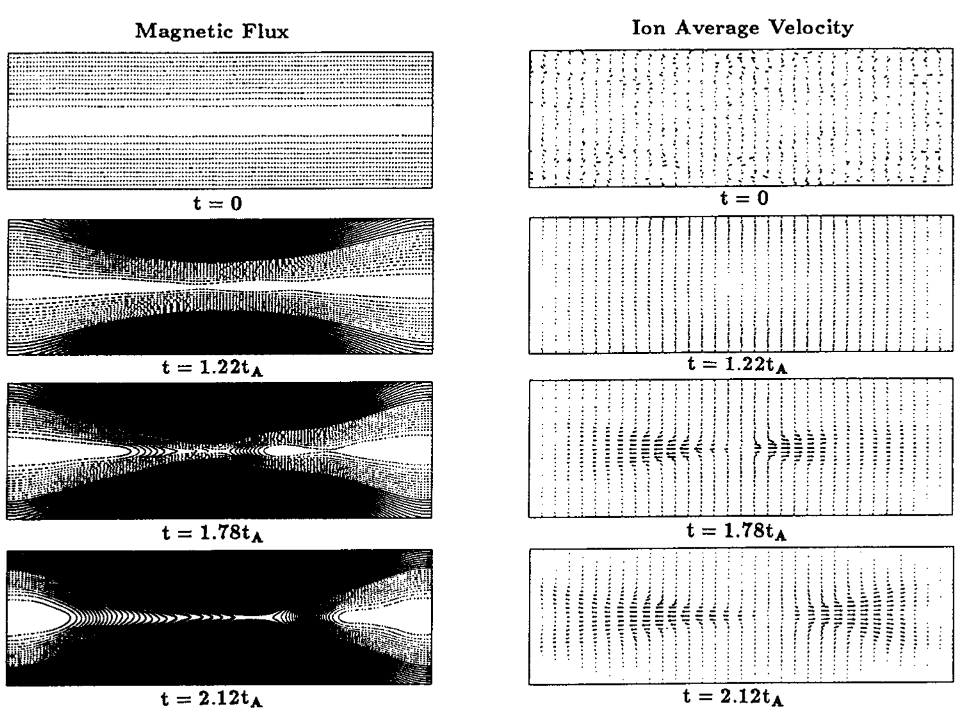

FIG. 2. Time sequences of magnetic flux contours (left) and vector plots of the average ion velocity (right) in the (x,y) plane at $t=0$, $t=1.22t_A$, $t=1.78t_A$, and $t=2.12t_A$ for case A, where the magnetic flux contours less than the initial value at the input boundary ($y=\pm y_b$) are plotted with the dotted lines.

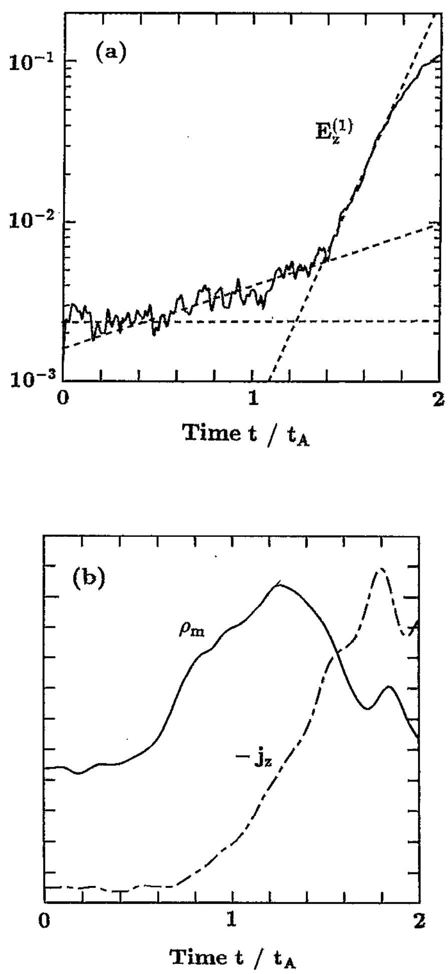

FIG. 3. Temporal evolutions of (a) the reconnection electric field, and (b) the mass density ρ_m (solid line) and the z component of the current density $-j_z$ (dot–dashed line) at the X point for case A. Note that the amplitude $E_z^{(1)}$ of the $n=1$ mode along the neutral sheet is plotted on a logarithmic scale in (a), while the mass density and the current density are plotted on a linear scale in (b). Three fitting functions of $E_z^{(1)}(t)=c_0(0<t<0.6t_A)$, $E_z^{(1)}(t)\propto\exp(0.92t/t_A)$ $(0.6t_A<t<1.3t_A)$, and $E_z^{(1)}(t)\propto\exp(6.0t/t_A)$ $(1.4t_A<t<1.7t_A)$ are also plotted in (a) for comparison.

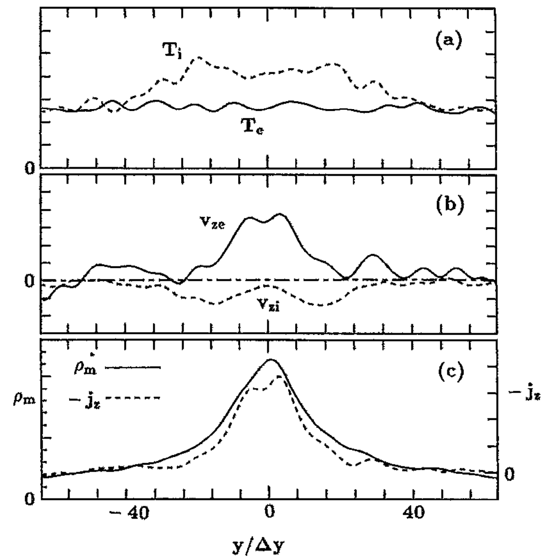

FIG. 4. Spatial profiles of (a) the electron temperature T_e (solid line) and the ion temperature T_i (dotted line), (b) the z component of the average electron velocity (solid line) and the z component of the average ion velocity (dotted line), and (c) the mass density (solid line) and the z component of the current density (dotted line) at $t=1.0t_A$ for case A, where the profiles along the vertical line passing on the reconnection point are plotted.

$E_z(x,0)$ along the neutral sheet is expanded into the Fourier series, and the amplitude $E_z^{(1)}$ of the $n=1$ mode is plotted on a logarithmic scale in Fig. 3(a), while the mass density and the current density are plotted on a linear scale in Fig. 3(b). There are three temporal phases in the evolution of the reconnection electric field $E_z^{(1)}$, i.e., the initial phase $(0<t<0.6t_A)$, the slow reconnection phase $(0.6t_A<t<1.3t_A)$, and the fast reconnection phase $(1.3t_A<t<1.8t_A)$. The reconnection electric field $E_z^{(1)}$ remains at the noise level in the initial phase. The electric field begins to grow slowly when both the mass density and the current density start increasing in the current layer as a result of the compression by the convergent plasma flow in the slow reconnection phase. The growth rate γ_g of the electric field is estimated to be $0.92/t_A$ or $0.103\omega_{ci}$, where $\omega_{ci}(=eB_0/M_ic)$ is the ion cyclotron frequency. The inclination of the growth curve steepens suddenly after the mass density reaches the maximum value at $t\approx1.3t_A$. The mass density tends to decrease slowly in the fast reconnection phase, while the current density continues to increase. The growth rate of the electric field in the fast reconnection phase is estimated to be $6.0/t_A$ or $0.013\omega_{ce}$. This value is much larger than the growth rate in the slow reconnection phase. These results suggest that magnetic reconnection in the fast reconnection phase is controlled by a different mechanism from that in the slow reconnection phase.

Let us examine the behaviors of the physical quantities in the slow reconnection phase. Figures 4(a)–4(c) show (a) the spatial profiles of the electron temperature T_e (solid line) and the ion temperature T_i (dotted line); (b) those of the z component of the average electron velocity (solid line) and the z component of the average ion velocity (dotted line); and (c) those of the mass density (solid line) and the z component of the current density (dotted line) at $t=1.0t_A$, where the profiles are plotted along the vertical line passing on the reconnection point $[(x,y)=(-2.5\ \Delta x,0)]$, and the particle temperature is obtained by assuming that the distribution function is approximated by the shifted Maxwellian. The equilibrium current is initially dominated by the diamagnetic component. In other words, the electron current is equal to the ion current because $T_e=T_i$ at $t=0$. The electron current becomes dominant over the ion current because of its small inertia as time goes on. It is worth noting that the mass density profile has the same width as the current density profile. In the vicinity of the magnetically neutral sheet, an ion executes a meandering motion[20] due to the finite ion Larmor radius effect. This motion creates the negative ion current at the edge region of the current layer. The ion motion is rather unconstrained by the magnetic field in the inner

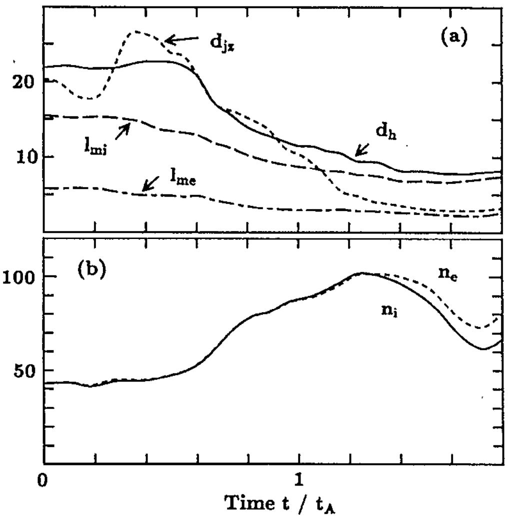

FIG. 5. Temporal evolutions of (a) the half-width d_h of the mass density profile (solid line), the half-width d_{jz} of the current density profile (dotted line) along the y axis, the average orbit amplitude l_{mi} of meandering ions (dashed line) and the average orbit amplitude l_{me} of meandering electrons (dot–dashed line); and (b) the electron number density (dotted line) and the ion number density (solid line) at the X point for case A, where the spatial scales d_h, d_{jz}, l_{mi}, and l_{me} are normalized by the grid separation Δy along the y axis.

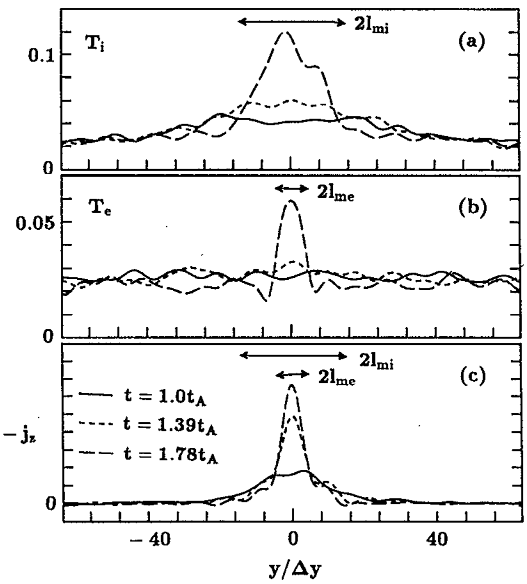

FIG. 6. Spatial profiles of (a) the ion temperature, (b) the electron temperature, and (c) the z component of the current density at $t=1.0t_A$ (solid line), $t=1.39t_A$ (dotted line), and $t=1.78t_A$ (dashed line) for case A, where the profiles along the vertical line passing on the reconnection point are plotted. The average orbit amplitude l_{mi} of meandering ions and the average orbit amplitude l_{me} of meandering electrons that are estimated at $t=1.78t_A$ are also drawn with arrows for comparison.

region of the current layer while it is strongly restricted by the motion of magnetic field outside the current layer. When the current layer is compressed by increasing magnetic pressure, ions inside the current layer act as if they are compressed by two walls approaching each other. The interaction between ions and the moving magnetic walls results in the ion heating through the Fermi mechanism, as was seen in Fig. 4(a). On the contrary, an electron moves together with the magnetic field without interacting with the magnetic wall, because the electron Larmor radius ρ_e is much smaller than the width of the current layer. Thus, the electron temperature remains almost constant in the slow reconnection phase.

Let us examine what happens in the transition period between the slow reconnection phase and the fast reconnection phase. Figures 5(a) and 5(b) show (a) the temporal evolutions of the half-width d_h of the mass density profile (solid line), the half-width d_{jz} of the current density profile (dotted line) along the y axis, the average orbit amplitude l_{mi} of meandering ions (dashed line), and the average orbit amplitude l_{me} of meandering electrons (dot–dashed line); and (b) those of the electron number density (dotted line) and the ion number density (solid line) at the X point, where $l_{mi} = \sqrt{d_h \rho_i}$ and $l_{me} = \sqrt{d_h \rho_e}$. The spatial scales d_h, d_{jz}, l_{mi}, and l_{me} in Fig. 5(a) are normalized by Δy, and $d_{jz} \approx d_h$ $(=0.863L)$ at $t=0$. Both the mass density profile and the current density profile have an initial half-width of about $2\rho_{i0}$, where the initial ion Larmor radius ρ_{i0} is equal to $10.4\ \Delta y$. Both d_h and d_{jz} decrease with the same rate in the compression time scale in the first half of the slow reconnection phase. The change of d_h slows down, and it approaches l_{mi} in the latter half of the slow reconnection phase. The half-width d_{jz} continues to decrease during the slow reconnection phase, and it approaches l_{me} at the end of the slow reconnection phase. Both d_i and d_{jz} remain almost constant in the fast reconnection phase, where $d_{jz} \approx 3\ \Delta y$ and $d_h \approx 8\ \Delta y$. Thus, the electric current is localized in a narrower region than the mass density in the fast reconnection phase.

The number density at the X point begins to increase at $t \approx 0.6t_A$, as is seen in Fig. 5(b). Both the electron density and the ion density increase with the same growth rate during the slow reconnection phase. The growth of the number density slows down and the density becomes maximum at the end of the slow reconnection phase. In the fast reconnection phase, the electron density becomes dominant over the ion density. Comparing Fig. 5(a) with Fig. 5(b), one can find that the charge separation becomes distinct after the scale height of the mass density profile becomes nearly equal to the ion Larmor radius, i.e., $d_h \approx l_{mi} = \sqrt{d_h \rho_i}$. This phenomenon can be easily understood by the finite ion Larmor radius effect in the vicinity of the neutral sheet.[14] That is, most of the ions in the current layer become unmagnetized when $d_h \approx \rho_i$, while the electrons remain magnetized. Therefore, the input flow (Poynting flux) no longer works on thinning the ion current layer, but acts to increase the ion temperature in it. For the electrons, however, compression keeps working. Consequently, the charge neutrality condition is violated in the central region of the current layer, when the fast reconnection phase sets in.

Let us examine the behaviors of the physical quantities in the fast reconnection phase in detail. Figures 6(a)–6(c) show (a) the ion temperature profiles; (b) the electron temperature profiles; and (c) the spatial profiles of the z component of the current density at three different periods, where the solid, dotted, and dashed lines correspond to the spatial profiles along the vertical line (y axis) passing on the reconnection point at $t=1.0t_A$, $t=1.39t_A$, and $t=1.78t_A$, respectively. The ion temperature takes a broad profile with a spatial scale comparable to l_{mi} and its peak value increases monotonously during both the slow reconnection phase and the fast reconnection phase. On the contrary, the electron temperature keeps a fairly flat profile during the slow reconnection phase, but increases suddenly in the narrow region comparable to the orbit amplitude l_{me} of an electron meandering motion as soon as the fast reconnection phase starts. The ion temperature is about two times as high as the electron temperature in the central region of the current layer in the fast reconnection phase.

It is interesting to notice in Fig. 6(c) that there are two spatial scales in the current density profile, i.e., the orbit amplitude l_{mi} of an ion meandering motion and the orbit amplitude l_{me} of an electron meandering motion. The current density profile with the large spatial scale of l_{mi}, "ion current layer," develops in the slow reconnection phase and is maintained in the fast reconnection phase. The peaked profile with the small spatial scale of l_{me}, "electron current layer," emerges inside the large-scale profile and grows rapidly in the fast reconnection phase. The separation of the current density into two, broad and narrow, spatial profiles is due to the dominance of either the finite ion Larmor radius effect or the finite electron Larmor radius effect. The ion current is dominant in the outer region of the current layer, while the electron current is dominant in the inner region in the fast reconnection phase.

Comparing Fig. 6(b) with Fig. 6(c), one can find that abrupt electron heating takes place in the electron current layer in the final phase of fast reconnection, i.e., $t=1.78t_A$. This phenomenon indicates that as the current layer is compressed as thin as l_{me}, the driving flow acts no longer to thin the current layer, but acts to increase the temperature of meandering electrons in it. A similar observation is made for the ions when the current layer becomes as thin as the ion Larmor radius.

Figure 7 shows contour maps of the mass density (top), the z component of the current density (second), the electron temperature (third), and the ion temperature (bottom) in the (x,y) plane at $t=1.78t_A$, where contours with a smaller value than the average are plotted with the dotted lines. Most of the plasma that exists inside the current layer in the slow reconnection phase is carried away by the divergent plasma flow toward the boundaries of the x axis (Fig. 2), and thus the mass density tends to take a fairly flat profile in the current layer in the fast reconnection phase. The electric current is enhanced in the long narrow region between two peaks of the mass density. The half-width of the current layer is comparable to the orbit amplitude of the electron meandering motion, as was seen in Figs. 5(a) and 6(c). The electron temperature profile is very similar to the current profile, since the electron heating takes place exclusively in the electron meandering layer, which is equivalent to the peaked current layer. In contrast, the ion temperature takes a broad profile with the half-width nearly equal to l_{mi}, which is caused by the compressional heating of meandering ions by the convergent plasma flow.

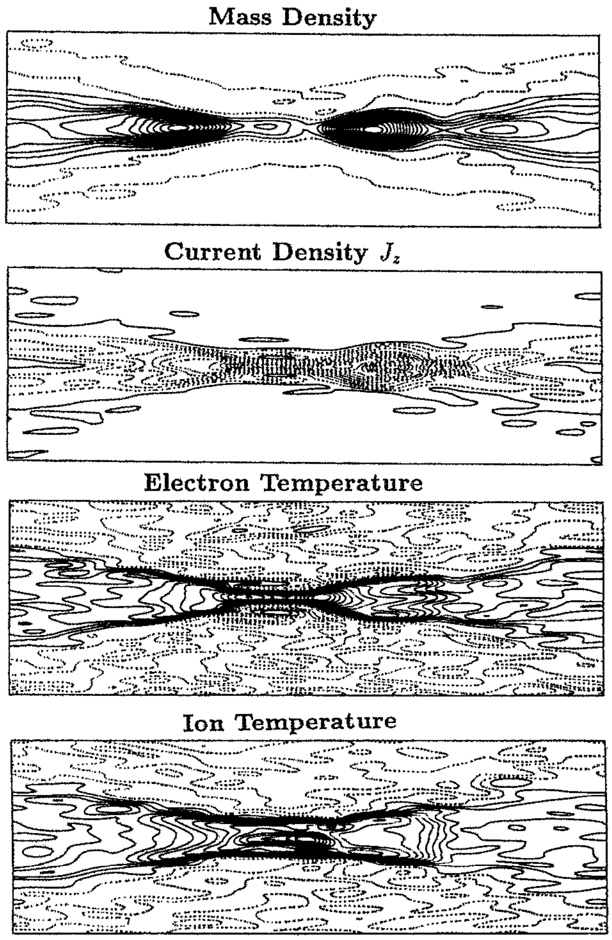

FIG. 7. Contour maps of the mass density, the z component of the current density, the electron temperature, and the ion temperature in the (x,y) plane at $t=1.78t_A$ for case A, where a contour with a smaller value than the average is plotted with the dotted line.

B. Dependence on ion mass

It is important to examine the parameter dependence of the simulation results in considering the physical mechanism of collisionless driven reconnection. We carry out several simulation runs with various sets of physical parameters. The simulation parameters are listed in Table I. The dependence on the ion mass is examined by comparing the simulation results for case A, case B, and case C, where the simulation parameters are chosen so that the parameters associated with the ion mass, say, the ion Larmor radius, the Alfvén velocity, and so on, are different, while keeping the others fixed. One aspect of the results is shown in Fig. 8, where the temporal evolutions of the reconnection electric field $E_z^{(1)}$ for case A (dotted line), case B (solid line), and case C (dashed line) are plotted on a logarithmic scale. Note in Fig. 8 that the time is normalized by using the electron cyclotron frequency ω_{ce},

TABLE I. Simulation parameters.

Case	M_i/M_e	$\rho_{i0}/\Delta y$	$d_{h0}/\Delta y$	E_0/B_0	ω_{pe}/ω_{ce}
A	50	10.4	21.9	−0.04	5.0
B	25	7.4	21.9	−0.04	5.0
C	100	15.1	21.9	−0.04	5.0
D	25	10.4	21.9	−0.04	3.5
E	100	10.4	21.9	−0.04	7.07
F	50	10.4	21.9	−0.01	2.0
G	50	10.4	21.9	−0.02	2.0

which is independent of the ion mass. The growth of the electric field becomes slower as the ion mass or ion Larmor radius increases. This result corresponds to the fact that the compression speed of the current layer decreases as the ion mass increases. On the other hand, the saturated reconnection rate or the peak value of the electric field is almost independent of the ion mass. By changing the normalization of time, it is found that three evolution curves agree fairly well with one another if a quantity proportional to $M_i^{1/4}$ is introduced as a normalization of time, as is shown in Fig. 9. That is, the whole temporal evolution of collisionless driven reconnection is characterized by a time scale in proportion to $M_i^{1/4}$, irrespective of whether the evolution stage is the slow or fast reconnection stage. Figures 10(a) and 10(b) show (a) the temporal evolutions of the ion temperature, and (b) those of the electron temperature at the reconnection point for the same cases as Fig. 8, where a closed circle, an open circle, and an open square correspond to the simulation results for case A, case B, and case C, respectively. The ion temperature increases in proportion to $M_i^{1/2}$ in the fast reconnection phase, while the electron heating takes place independent of the ion mass. These results suggest that the dynamical evolution of collisionless driven reconnection and the ion heating are controlled by a physical mechanism associated with

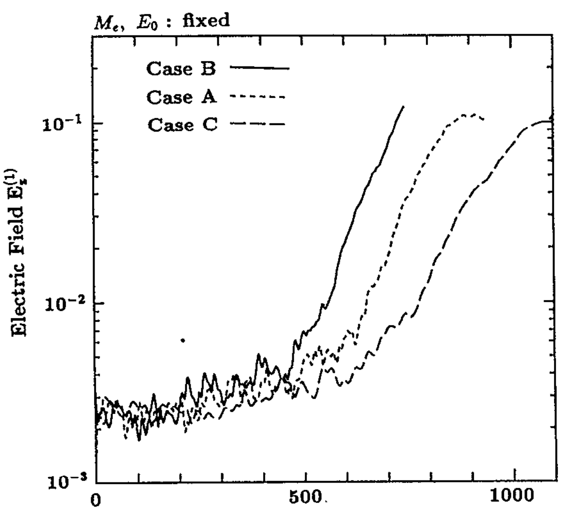

FIG. 8. Temporal evolutions of the reconnection electric field for three different ion masses with a fixed electron mass and a fixed driving electric field, i.e., $\rho_{i0}/\Delta y$=7.4 (case B), 10.4 (case A), and 15.1 (case C).

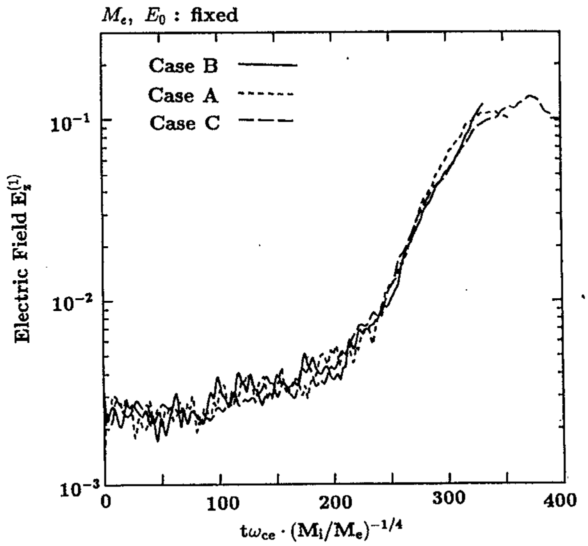

FIG. 9. The same figure as Fig. 8, but the abscissa is replaced by $t\omega_{ce}(M_i/M_e)^{-1/4}$. Note that the electron mass is fixed for three cases.

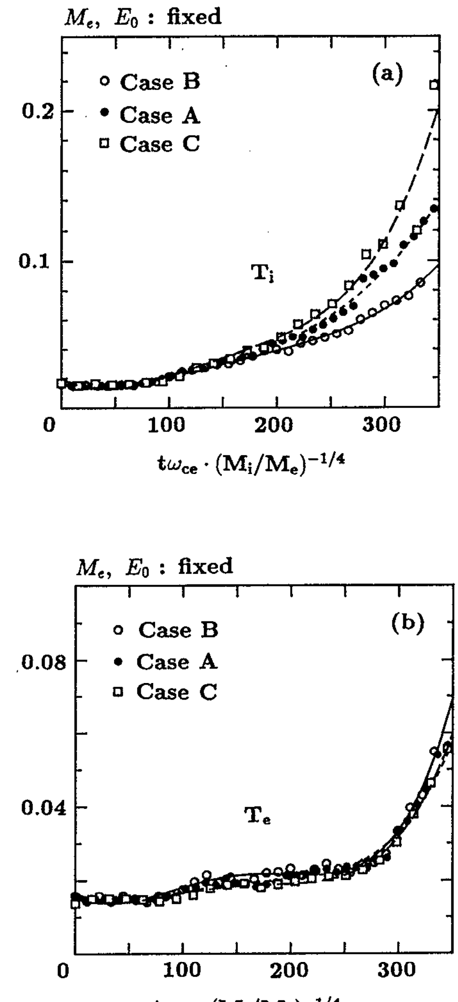

FIG. 10. Temporal evolutions of (a) the ion temperature and (b) the electron temperature at the reconnection point for the same cases as Fig. 8, where an open circle, a closed circle, and an open square correspond to the simulation results for case B, case A, and case C, respectively.

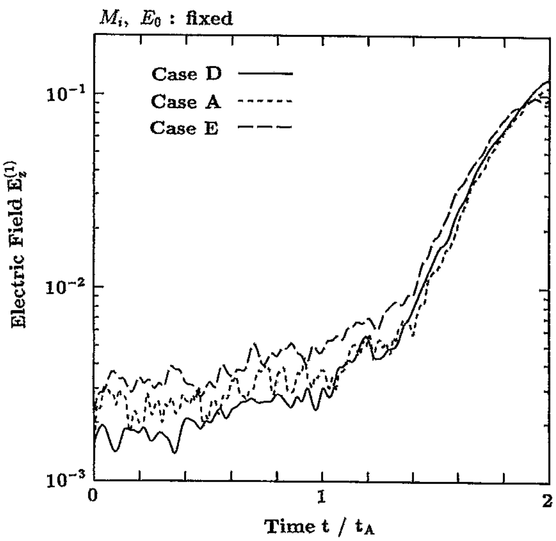

FIG. 11. Temporal evolutions of the reconnection electric field for three different electron masses with a fixed ion mass and a fixed driving electric field, i.e., $\omega_{pe}/\omega_{ce}=3.5$ (case D), 5.0 (case A), and 7.07 (case E).

the ion dynamics, but that the electron heating is caused by a mechanism independent of the ion mass. We will give a phenomenological explanation of these results in connection with the dynamics of the ion current layer in the next section.

C. Dependence on electron mass

The dependence on the electron mass is examined by comparing the simulation results for case A, case D, and case E, where the ion mass is fixed. Figure 11 shows the temporal evolutions of the reconnection electric field $E_z^{(1)}$ for case A (dotted line), case D (solid line), and case E (dashed line), where the time is normalized by the Alfvén transit time. In contrast to the cases shown in Fig. 8, there is no significant difference between the three cases, except for the fact that the noise level in the initial phase depends weakly on the electron mass. In other words, the development of the reconnection field depends little on the electron mass, although most of the electric current in the vicinity of the reconnection point is carried by electrons and the width of the current layer is given by the small spatial scale associated with the electron mass in the fast reconnection phase. Figures 12(a) and 12(b) show (a) the temporal evolutions of the ion temperature; and (b) those of the electron temperature at the reconnection point for the same cases as Fig. 11, where an open circle, a closed circle, and an open square correspond to the results for case D, case A, and case E, respectively. The electron temperature increases slowly as soon as the slow reconnection phase starts, and a small but significant difference appears between the three cases. The electron temperature remains nearly constant for a while, and then increases rapidly in the fast reconnection phase. The electron heating takes place more efficiently as the electron mass becomes larger. It is found that the electron temperature scales as $M_e^{1/4}$ in the fast reconnection phase. In contrast to the electron temperature, the ion temperature increases independent of M_e. These results indicate that while the ion heating is caused by the compressional heating of meandering ions, the electron heating is a process of pure electron dynamics. The proportionality of the electron temperature to an electron mass, $M_e^{1/4}$, suggests that the electron heating must result from energization of meandering electrons in the narrow current layer, subject to the incoming Poynting flux from the external region.

D. Dependence on driving electric field

According to MHD studies,[3] the reconnection rate depends strongly on the driving electric field E_0. Let us examine the dependence of the reconnection rate in a collisionless plasma on the driving electric field. Figure 13 shows the temporal evolutions of the reconnection electric field $E_z^{(1)}$ for case A (dotted line), case F (solid line), and case G (dashed line), where the time is normalized by the Alfvén transit time t_A. Magnetic reconnection starts earlier and the electric field grows faster as the driving electric field becomes larger. The

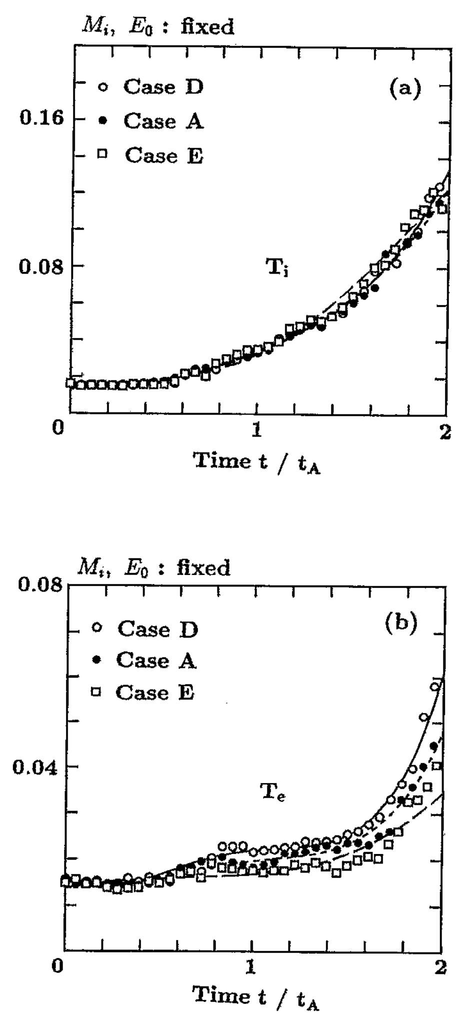

FIG. 12. Temporal evolutions of (a) the ion temperature and (b) the electron temperature at the reconnection point for the same cases as Fig. 11, where an open circle, a closed circle, and an open square correspond to the simulation results for case D, case A, and case E, respectively.

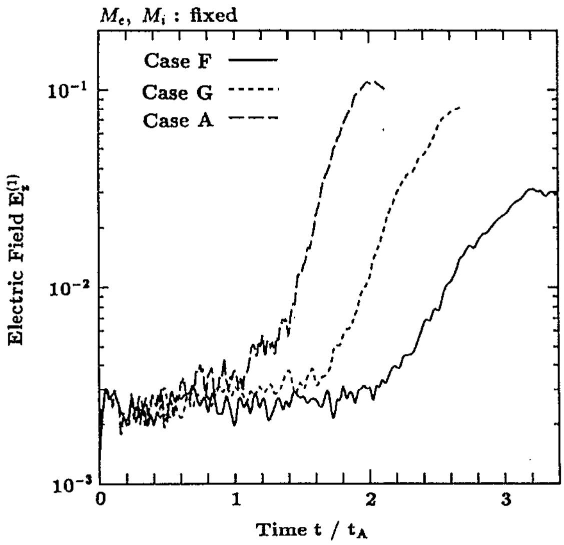

FIG. 13. Temporal evolutions of the reconnection electric field for three different driving fields, with a fixed ion mass and a fixed electron mass, i.e., $E_0/B_0=-0.01$ (case F), -0.02 (case G), and -0.04 (case A).

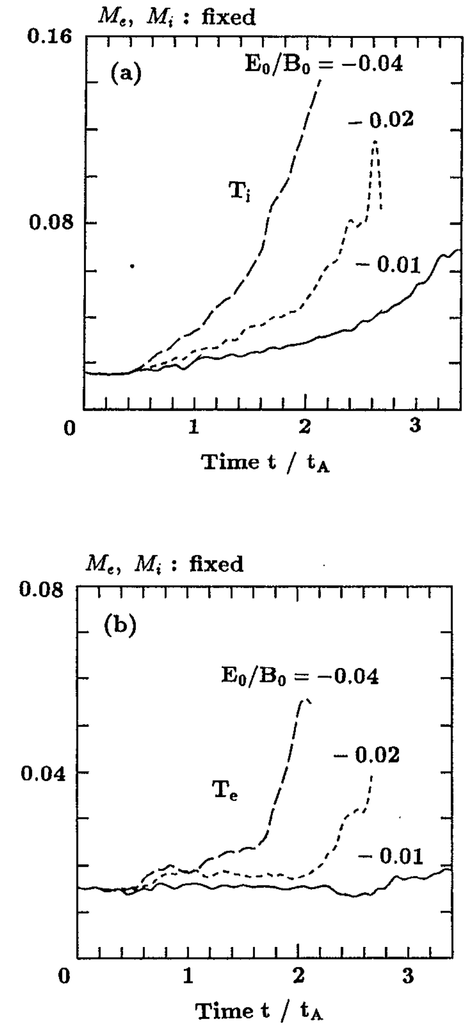

FIG. 14. Temporal evolutions of (a) the ion temperature and (b) the electron temperature at the reconnection point for the same cases as Fig. 13.

detailed examination reveals that the evolution time of magnetic reconnection scales as $E_0^{-1/2}$. This reflects the fact that the compression speed of the current layer depends strongly on the input rate of the magnetic flux through the boundary, i.e., the driving electric field E_0. The saturated reconnection rate or the peak value of electric field is roughly in proportion to E_0. This result is in good agreement with that of the MHD studies,[3] where the reconnection rate is approximately given by the external driving electric field and depends little on the value of the electric resistivity.

Figures 14(a) and 14(b) show (a) the temporal evolutions of the ion temperature; and (b) those of the electron temperature at the reconnection point for the same cases as Fig. 13, where the temperature is plotted in the same normalization unit. Both the electron heating and the ion heating take place more efficiently as the driving electric field becomes larger. The ion temperature is always higher than the electron temperature, irrespective of the driving electric field, except in the initial phase. The ion temperature is two or more times as high as the electron temperature in the fast reconnection phase. These results indicate that the plasma compression becomes stronger as the driving electric field increases, and thus, reconnection and collisionless plasma heating take place more efficiently. We will briefly touch on a physical meaning of the scaling law with E_0 in the next section.

IV. SUMMARY AND DISCUSSIONS

We have investigated driven magnetic reconnection in a collisionless plasma by means of two-and-one-half-dimensional particle simulation for the case where the current layer is compressed by the convergent plasma flow, which is supplied through the boundary by the external electric field. The main results are summarized as follows.

(1) Driven magnetic reconnection consists of three temporal phases, i.e., the initial phase in which the electric field remains at the noise level ($t<0.6t_A$), the slow reconnection phase ($0.6t_A<t<1.3t_A$), and the fast reconnection phase ($1.3t_A<t<1.8t_A$).

(2) The reconnection electric field grows at the reconnection point with a slow growth rate of the order of roughly $0.1\omega_{ci}$, while both the current layer and the mass density are equally being compressed.

(3) When the width of the current layer is compressed as thin as the orbit amplitude of the ion meandering motion, the ions in the current layer become unmagnetized while the electrons remain magnetized. Only magnetized drifting electrons can carry the Poynting (magnetic) flux inside the current layer, thereby enhancing the electron current in a narrow layer with the spatial scale of the orbit amplitude of meandering electrons. Thus, the current layer is split into ion and electron current layers in the fast reconnection phase. Magnetic reconnection takes place with a fast growth rate of $0.01\omega_{ce}$, in good accordance with this enhanced narrow electron current layer.

(4) The evolution time of magnetic reconnection scales as $E_0^{-1/2}M_i^{1/4}$, and depends little on the electron mass. That is, magnetic reconnection develops faster as the ion mass decreases and/or the driving electric field increases. The

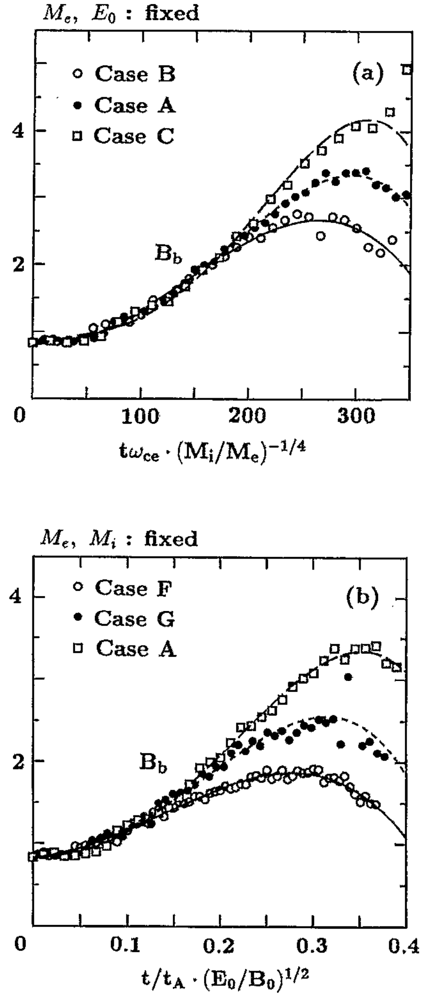

FIG. 15. Temporal evolutions of the magnetic field at the current boundary (a) for the same cases as Fig. 8, and (b) for the same cases as Fig. 13. Here the time is normalized by $(M_i/M_e)^{1/4}/\omega_{ce}$ in (a), and by $(E_0/B_0)^{-1/2}t_A$ in (b).

maximum reconnection rate is roughly in proportion to the driving electric field, but independent of both the ion mass and the electron mass.

(5) Only the ion heating takes place inside the current layer in the slow reconnection phase. In the fast reconnection phase, the ion heating takes place inside the broad ion current layer, while the electron heating takes place only inside the narrow electron current layer. The ion temperature is two or more times as high as the electron temperature during the reconnection process.

Here, let us give a phenomenological explanation of the scaling law of $E_0^{-1/2}M_i^{1/4}$ by using a simple model of the current layer. Suppose that both an ion and an electron are magnetized outside the ion current layer, while an ion is unmagnetized and an electron is magnetized inside the ion current layer. A convergent plasma flow carries the magnetic field from the boundary of the simulation box toward the ion current layer and compresses it. Since magnetized drifting electrons can penetrate through the ion current layer, magnetic flux can be carried inside.[14,22] Thus, the ion current layer dynamically evolves while satisfying the balance between the plasma compression by the convergent plasma flow and the penetration of the magnetic flux into the ion current layer.

TABLE II. Magnetic field at the ion current boundary.

Case	A	B	C	F	G
M_i/M_e	50	25	100	50	50
E_0/B_0	−0.04	−0.04	−0.04	−0.01	−0.02
B_b/B_0	3.4	2.7	4.1	1.9	2.5

Figures 15(a) and 15(b) show (a) the temporal evolutions of the magnetic field B_b at the current boundary for case A (closed circle), case B (open circle), and case C (open square); and (b) those for case A (open square), case F (open circle), and case G (closed circle). The time is normalized by $(M_i/M_e)^{1/4}/\omega_{ce}$ in Fig. 15(a), and by $(E_0/B_0)^{-1/2}t_A$ in Fig. 15(b). It is found in Figs. 15(a) and 15(b) that the maximum value of magnetic field B_b at the current boundary, which is listed in Table II, changes, according to the scaling law of $B_b \propto E_0^{1/2}M_i^{1/4}$. Though not shown here, the electric field E_b at the current boundary is observed to be proportional to the driving electric filed E_0. Since the penetration speed v_{pn} of magnetic field into the current layer is represented by the $\mathbf{E}\times\mathbf{B}$ drift velocity,[14] we have the relation as $v_{pn} \propto E_0^{1/2}M_i^{-1/4}$. This scaling law is in good agreement with the inverse of the scaling law of the dynamical evolution time of magnetic reconnection. We conclude that the penetration speed at the ion current boundary determines the whole dynamical evolution of collisionless driven reconnection.

It has been clarified by our simulation study that there exist two types of physical mechanism, leading to magnetic reconnection in a collisionless plasma. One is responsible for the slow reconnection process associated with the ion meandering motion, and the other is responsible for the fast reconnection process associated with the electron meandering motion. When the width of the current layer becomes comparable to the orbit amplitude of the ion meandering motion, the rate of magnetic reconnection due to the ion meandering motion is dominant over that due to the electron meandering motion. The penetration of the magnetic field into the current layer splits the current layer into the ion one and the electron one by compressing only electrons within the ion current layer. The Poynting flux into the ion current layer results in the increase of the ion temperature and the enhancement of the electric current. Since the ion temperature is always higher than the electron temperature, the magnetic pressure is roughly balanced with the ion thermal pressure. The pressure balance reduces to $T_i \propto E_0^{1/2}M_i^{1/2}$ where we use the observed scaling law of the number density n_{in} at the center of the current layer, i.e., $n_{in} \propto E_0^{1/2}$. This scaling law of the ion temperature is consistent with that of the simulation results shown in Figs. 10(a) and 14(a). When the electron current layer is compressed as thin as the orbit amplitude of the electron meandering motion, the thinning of the electron current layer is stopped. After this period, the rate of magnetic reconnection due to the electron meandering motion is dominant over that due to the ion meandering motion. The electric field E_z at the reconnection point increases with the growth rate of fast reconnection until it reaches the input rate of magnetic field at the boundary of the simulation domain, i.e., the driving electric field E_0. Thus, the maximum reconnection rate is roughly in proportion to the driving electric field.

We have discussed the simulation results without specifying the physical mechanism responsible for driven magnetic reconnection in a collisionless plasma. It is widely believed that a collisionless tearing instability[7–17] is an effective mechanism for collisionless magnetic reconnection. Let us first consider the triggering mechanism of the collisionless reconnection in the slow phase. The linear growth rate of the ion tearing instability[17] is estimated to be $\gamma_i \approx 0.08\omega_{ci}$ for a thin current layer ($d_{jz} \approx 2\rho_{i0}$). This value is in good agreement with the simulation results of $\gamma_g \approx 0.103\omega_{ci}$ in the slow reconnection phase. This agreement suggests that the slow reconnection can be triggered by the ion tearing instability associated with the ion meandering motion.

One can immediately come up with the electron tearing instability[7,8,10–12] for the fast reconnection process. The linear growth rate γ_e of the electron tearing instability[7] is estimated to be $\gamma_e \approx 0.02\omega_{ce}$ if the relations $d_{jz} \approx l_{me}$ and $\gamma_e \approx (\rho_e/d_{jz})^{5/2}\omega_{ce}$ are used. This value appears consistent with the simulation results of $\gamma_g \approx 0.013\omega_{ce}$ in the fast reconnection phase, and the half-width of the electron current layer is in good agreement with the theoretical value of $l_{me} = \sqrt{d_h \rho_e}$ inside which the resonant wave–particle interaction takes place.[11] Thus, fast reconnection may be attributable to the electron tearing instability. If so, however, we have a conflict with the obtained scaling law as $E_0^{-1/2} M_i^{1/4}$, which is independent of the electron dynamics. Remember, however, that all the electron parameters, such as the electron cyclotron frequency and the electron Larmor radius are not time independent, but strongly time variable. It is quite likely that the local electron parameters are so adjusted as to match the external macroscopic dynamics, and the apparent dynamical behavior appears to be controlled by the ion dynamics, although, in fact, the evolution is governed by the local electron dynamics, which is highly susceptible to the surrounding macroscopic condition.

Before concluding this paper, let us point out two important problems to be studied in the near future. First, an electron flow with a large bulk velocity along the z axis was observed in the fast reconnection phase in our simulation. This fact suggests a possibility that the electron flow would excite kinetic waves, such as the lower-hybrid-drift mode,[5] thereby anomalous resistivity being generated to cause collisionless magnetic reconnection. Investigation of this effect requires a really three-dimensional treatment. Second, in the fast reconnection phase, the electron current layer is found to become as thin as the orbit amplitude of the electron meandering orbit. However, its width is three times the grid size along the y axis, i.e., $d_{jz} \approx 3\ \Delta y$. Furthermore, the electron skin depth c/ω_{pe} is typically equal to $2\ \Delta y$. These values are close to the lower limit of the spatial resolution in the numerical calculation, i.e., the grid separation Δy. Thus, we are now carrying out simulation runs with a higher spatial resolution than that in the present study in order to make the simulation results more physically comprehensive and reliable.

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