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Experimental study of the dynamics of a thin current sheet

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Abstract

Many plasmas in natural settings or in laboratory experiments carry currents. In magnetized plasmas the currents can be narrow field-aligned filaments as small as the electron inertial length in the transverse dimension or fill the entire plasma column. Currents can take the form of sheets, again with the transverse dimension the narrow one. Are laminar sheets of electric current in a magnetized plasma stable? This became an important issue in the 1960s when current-carrying plasmas became key in the quest for thermonuclear fusion. The subject is still under study today. The conditions necessary for the onset for tearing are known, the key issue is that of the final state. Is there a final state? One possibility is a collection of stable tubes of current. On the other hand, is the interaction between the current filaments which are the byproduct endless, or does it go on to become chaotic? The subject of three-dimensional current systems is intriguing, rich in a variety of phenomena on multiple scale sizes and frequencies, and relevant to fusion studies, solar physics, space plasmas and astrophysical phenomena. In this study a long (δz = 11 m) and narrow (δx = 1 cm, δy = 20 cm) current sheet is generated in a background magnetoplasma capable of supporting Alfvén waves. The current is observed to rapidly tear into a series of magnetic islands when viewed in a cross-sectional plane, but they are in essence three-dimensional flux ropes. At the onset of the current, magnetic field line reconnection is observed between the flux ropes. The sheet on the whole is kink-unstable, and after kinking exhibits large-scale, low-frequency rotation about the background field with an amplitude that grows with distance from the source of the current. Three-dimensional data of the magnetic and electric fields is acquired throughout the duration of the experiment and the parallel resistivity is derived from it. The parallel resistivity, for the most part, is not largest in the reconnection regions, but peaks in the neighborhood of large current gradients. At early times a quasi-separatrix layer (QSL) is observed where the current sheet tears, but later on a QSL of larger value, not obviously associated with reconnection, is measured at the edge of the current sheet. This QSL enhancement is connected with the rapidly spatially diverging magnetic fields in the moving sheet (ropes).

Keywords: flux ropes, 3D reconnection, 3D current systems, tearing mode, basic plasma experiment

(Some figures may appear in colour only in the online journal)
1. Introduction

Magnetic field line reconnection has been considered an important topic in plasma physics for decades. Reconnection is, simply stated, the forcing together of oppositely directed magnetic fields by a variety of means such as plasma flows or the temporal evolution of faraway currents. In three dimensions this may involve only a small portion of the total field. The topic of field line reconnection has also been studied extensively in theory [1], and computer simulations [2], of various complexities as well as in experiments [3]. The magnetic field in two-dimensional reconnection (or a component of it in 3D) becomes zero locally and reappears as energy in waves, flows, jetting particles, or heat. The source of all of this is an electric field, which develops when magnetic flux is annihilated. In fact, the reconnection rate is defined as the electric field generated by the process aligned along the magnetic field lines in the system. In this experiment it was directly measured. The early theories of Parker [4], Sweet [5], and Petschek [6], and computer simulations of reconnection were two-dimensional. Sweet-Parker (SP) type models became a standard of comparison, and whenever the reconnection rate could be estimated it was compared to the SP rate. When experiments were limited and computer technology prohibited anything else it made sense to study 2D current sheets and gain a foothold on the problem. This is no longer the case.

While historically much of the research in reconnection physics has focused on two-dimensional models, this is rapidly changing. Solar physicists have realized that structures in the corona, and magnetic reconnection associated with them, are fully three-dimensional and have tailored their simulations to deal with this. The concept of quasi-separatrix layers (QSLs) [7], which is used in this work, was developed to address reconnection in the most general case where there is no magnetic null point in the system. Finally one must appreciate that reconnection does not happen in isolation. The closest system to the 2D case is thought to be in the Earth’s magnetotail. It is often the case that there is a weak guide field (of the order of the transverse fields) in the tail, but often the field is even smaller. Even in this case computer simulations have revealed a variety of three-dimensional effects. The magnetic fields are part of dynamic current systems; the currents and how they close cannot be ignored (though they usually are). The currents can heat electrons on their own, and cause instabilities making it difficult to separate what is due to reconnection and what is not. Often experiments start with magnetic fields forced to merge by external circuits or plasmas forced together, and simulations often involve oppositely directed magnetic fields that are pushed into one another by artificial means. Reconnection can occur in a spontaneous and straightforward way when a sheet of current tears.

Current sheets and tearing modes in current sheets have been studied for quite a long time both theoretically [8], and experimentally [9]. Tearing involves the filamentation of a sheet of current into two or more smaller current channels. If these channels subsequently bump into one another they can magnetically reconnect at the impact sites. These could be at points or along a line or surface. There has been a parallel track of research on magnetic flux ropes. Magnetic ropes consist of field lines with pitch varying with radius across the rope diameter, a result of currents flowing along the field. They are seldom, if ever, motionless. It is a common occurrence for flux ropes to have companions. In the solar corona flux ropes frequently occur in pairs and move toward or away from one another as their footprints, which are frozen onto the solar surface, move about. X-ray and XUV images of the Sun in which multiple rope-like structures appear are not uncommon. When the footprints of flux ropes on the solar surface move they can cause the ropes to collide, and partially reconnect in the corona [10]. If a current sheet is entrained in a background magnetic field, the growth of the tearing instability will give rise to a series of flux ropes, which may strongly interact as they grow to sufficient amplitude [11]. A key point that must be taken into account in simulations of the tearing mode is that non-periodic systems are very different from periodic ones [12]. It is obvious that the subjects of current sheet tearing, magnetic flux ropes and reconnection are not separate. In this work, these subjects are explored experimentally for a thin current sheet that is unstable to the tearing instability, leading to the formation of 3D flux ropes. The magnetic field and plasma currents as well as $\vec{E} = -\nabla \phi - \frac{\partial \vec{A}}{\partial t}$ are measured throughout the plasma volume. Reconnection occurs but it is not the star player on the court. Instead, the dynamics is dominated by the large-scale motion of the kink-unstable flux ropes after the current sheet tears. The total electric field and parallel plasma resistivity are measured and are largest in the region of current gradients. The parallel resistivity is not large compared to the classical resistivity in regions where reconnection occurs.

2. Overview of tearing mode theory

One of the most general and far-reaching ways in which magnetic reconnection can be triggered is through the tearing instability, which is driven by gradients in the current density and spontaneously gives rise to topological changes in the magnetic field structure. Similar to nonlinear reconnection processes, the growth of tearing instabilities depends upon the cross-scale coupling between the outer equilibrium scales, which are normally well-described by ideal MHD, and the inner resonance layer where $\vec{k} \cdot \vec{B} = 0$ and resistivity or kinetic effects are important. Within the original derivation of the tearing instability [8], both the outer region and inner resonance layer were treated with MHD. This classic paper was followed by many others, which extended the domain of tearing instabilities into both two-fluid and collisionless kinetic regimes (see, e.g., [13], and references therein). While these theories assume periodic boundary conditions appropriate to toroidal fusion machines, variations on the basic periodic theory have been applied to space plasmas, including the Earth’s magnetotail [14], as well as the dayside magnetopause [15]. However, since the equilibrium structures of interest in space and astrophysical plasmas are never periodic,
the influence of boundary conditions on the development of tearing modes has also been explored. In particular, the influence of line-tied boundary conditions has been considered within both the MHD [12, 16], and kinetic [17], descriptions of tearing. To extend the theory to nonperiodic systems, the basic approach involves superimposing multiple eigenmodes from the periodic theory in order to satisfy the line-tied boundary condition in the nonperiodic case. In sufficiently large systems, one can often superimpose just two tearing eigenmodes from the periodic theory, which is sufficient to approximately satisfy the boundary conditions. The results from these papers suggest that in large space and astrophysical plasmas, the stability constraints arising from boundary conditions are not overly restrictive, and thus the basic scaling predictions for the tearing mode growth rate will survive in non-periodic systems.

3. Experimental setup

The experiment was done in the upgraded Large Plasma Device (LaPD) at UCLA [18]. The machine is pictured in figure 1.

The background plasma is helium (L = 17 m, dia = 60 cm, n = 1.0 × 10^{19} cm^{-3}, Te = 4 eV, T_i = 0.5 eV) and is pulsed at a 1 Hz repetition rate (t_{on} = 10 ms). The plasma is produced with a DC discharge (V_D = 70 V, I_D = 4 kA) with a cathode-anode separation of 50 cm. The background plasma is
current-free and quiescent ($\delta n/n = 5\%$), and flux ropes and current sheets do not occur naturally within it. To generate the current sheet a second plasma source is required. The source is based on the use of a lanthanum hexaboride cathode ($\text{LaB}_6$) [19], which has a large emissivity but must be nearly twice as hot ($T_{\text{LaB}_6} \approx 1750 \, \text{°C}$) as barium oxide cathodes. A schematic diagram of the experimental setup is shown in figure 2. A current sheet is formed by placing a carbon slot ($\delta x = 1 \, \text{cm}$, $\delta y = 20 \, \text{cm}$) in a port 260 cm in front of the cathode and 11.1 meters from a molybdenum anode ($\delta x = \delta y = 25 \, \text{cm}$) and then applying a pulsed voltage ($V_{\text{slot}} = 260 \, \text{V}$, $I_{\text{slot}} = 390 \, \text{A}$, $\delta t_{\text{slot}} = 3.5 \, \text{ms}$) between them. The carbon slot is electrically ‘floating.’ The slot current is on during the background plasma discharge. The plasma is diagnosed with swept Langmuir probes ($n$, $T_e$), magnetic probes ($\frac{\partial B}{\partial t}$) and emissive probes ($V_{\text{pot}}$) and a vacuum ultraviolet spectrometer (fast electrons). The repeatability of the experiment, coupled with the relatively high repetition rate, enables the collection of time-dependent, volumetric data sets.

4. Results

4.1. Magnetic fields and currents

The magnetic field is measured using a three-axis differentially wound magnetic pickup probe [20]. The probe was moved to 3025 locations on a plane transverse to the magnetic field ($\delta x = \delta y = 0.5 \, \text{cm}$) and on 11 such planes parallel to the background magnetic field ($\delta z = 64 \, \text{cm}$). Data was acquired every 0.32μs for 12,000 time steps and at 10 experimental shots for each position. An emissive probe designed to measure the plasma potential at the relatively high densities in this experiment [21], acquired data at the same spatial locations as magnetic fields. A second three-axis magnetic probe was fixed in space ($\delta x = 6$, $\delta y = 0$, $\delta z = 870 \, \text{cm}$) for the entire data run. This is in the middle of the machine and slightly off center. The fixed probe acquired data every shot while the other probes were moved throughout the machine. A correlation technique was used to determine a new $\delta t = 0$, corresponding to the time the current sheet undergoes coherent oscillations rather than the time the rope currents are switched on. This is in essence a conditional trigger. Figure 3 shows ten temporal sequences of the magnetic field, acquired at a given location.

The magnetic field data is not divergence-free because the probe’s axes are not perfectly aligned with the machine axis. Once the data is aligned and averaged it is ‘divergence cleaned’ [22]. This matters in the precise calculation of field lines. The vector potential $\mathbf{A}$ is calculated by integrating in the Fourier domain using the Coulomb gauge. After Fourier inversion the vector potential is splined with $C^2$ continuous splines. Using the splined vector potential the divergence cleaned magnetic field is calculated using

$$\mathbf{B} = \nabla \times \mathbf{A},$$  

Figure 3. Ten experimental shots at a spatial location in the current sheet as a function of time. The upper figure shows the ten data shots and the lower figure shows them aligned using a correlation with the fixed probe. The dark line is the ten-shot average in each case. The conditional trigger allows proper averaging.
which now identically ensures that the divergence of $B$ is zero.

Figure 4 shows one component of the magnetic field, $B_x$, near the center of the data plane at two different $z$ locations. There are small oscillations in the first 1.5 ms, which are barely visible, but the 16-bit digitizer captures them as well as the large oscillations at later times ($t > 1.5$ ms). The oscillations grow spatially and are associated with the motion of the entire current sheet. The RMS value is quite large, on the order of 10 gauss or five percent of the background magnetic field.

The conditional averaging technique works well after the monochromatic oscillations appear ($t > 1.5$ ms, which we will refer to as the second epoch). At earlier times both the conditional trigger and a straight average of the ten data shots were compared. It turns out that the magnetic islands which occur in the first epoch were most clearly visible using simple averaging. This makes sense as the coherent mode does not exist in the first epoch, and the time scale of island development was much faster than that of the low-frequency oscillations.
4.2. Current sheet tearing

The current sheet tears very quickly, in approximately an Alfvén transit time along the length of the sheet. Figure 5 shows the transverse vector magnetic field $B_{xy}$, $(\psi_0)$ field in a plane 8.63 meters from the source of the sheet just after the current sheet is turned on. As the current builds up, the axial magnetic field generated by the current can rise to five percent of background field and the transverse field as large as ten percent. In reconnection terminology, this implies the guide field is approximately ten times larger than the reconnecting field. When the current sheet breaks up, the individual flux rope currents are inclined with respect to $B_z$, and 'X' and 'O' points at a given distance from $z = 0$ show up when the viewing plane is tilted. This will be illustrated; however, one can get a sense of the sheet breakup by examining the field in transverse planes due to the axial current, $J_z$. The law of Biot and Savart states

$$\mathbf{B}_o(\mathbf{r}) = \frac{\mu_0}{4\pi} \int \frac{\mathbf{J} \times \mathbf{r'}}{|\mathbf{r} - \mathbf{r'}|^2} \; dr',$$

in which $\mathbf{r}$ is a vector in the transverse plane of interest. Fifteen microseconds after the current is switched on the current sheet begins to tear. This is roughly the Alfvén transit time between the cathode and anode.

The tearing of the current sheet is apparent in figure 6. Here the transverse magnetic vectors are shown in two planes 2.4 meters apart and at $\tau = 762.1 \mu s$ after the sheet is switched on (the current is switched on at $\tau = 0$). The transparent yellow surface is an isosurface of current density $J_z = -0.3 \text{ A/cm}^2$. Two large islands have formed at the ends of the sheet and several small islands in the center.

It is instructive to examine the components of the current along the background magnetic field. Figures 3 and 4 suggest that the temporal regimes ($0 \leq \tau \leq 1.5 \text{ ms}$) and $\tau > 1.5 \text{ ms}$ contain different phenomena. The large, nearly monochromatic oscillation in the magnetic field occurs in the second epoch. The current density along the background magnetic field, $J_z$, is shown in figure 7. Data is displayed on three transverse planes, two at the ends and one in the center of the volume data was acquired within. At the earlier time (left-hand column of figure 7) the current sheet is no longer a slab (figure 5) but has developed several local maxima, which are at different locations as one moves away from the source. A small tilt of the current sheet develops with distance. At a later time (right-hand column of figure 7) the current sheet has rotated by 90 degrees 3.84 meters from its origin, and at 7 meters has multiple islands. Return currents are visible at every time. The complex morphology of the current and evolves in space and time.

The current ropes are fully three-dimensional, and the proper way to search for magnetic islands is within the volume and not on transverse planes. Once multiple islands exist the experiment becomes one on the interaction of flux ropes [23]. Divergence cleaning of the magnetic field data
allows calculation of the magnetic fields, currents and other quantities anywhere within the volume in which we acquired data. To view the true axial current one must look in a plane which is slightly tilted with respect to the machine axis because the currents are tilted. In figure 8 the data plane has been tilted by three degrees about the horizontal axis. This small tilt makes the two current channels that the sheet has broken into stand out sharply. If one were to plot \( \rho B \) in this plane it would surround the currents just as in figure 6. Traditionally one associates the out-of-plane vector potential with X and O points, but when it is plotted with the current in figure 8 only an O point is visible. Changing the tilt angle over a wide range never results in the appearance of X points, which speaks against the presence of reconnection. A low level of reconnection is present, however. The perpendicular magnetic field, consistent with the currents in figure 6, has the requisite island geometry. The correspondence of \( A_z \) to magnetic field lines is only rigorous in 2D and there is no sign of magnetic islands in the \( A_z \) data here. In a fully three-dimensional system it is not appropriate to associate contours of \( A_z \) with reconnecting field lines.

The transverse magnetic fields due to the currents are on the order of 10 gauss or five percent of the background field. The transverse currents, however, have a large contribution due to diamagnetic currents, as found in earlier experiments [19]. For the present conditions the diamagnetic current \( \vec{J} = \nabla \rho \times \vec{B} / B^2 \) is estimated to be on the order of 1–2 A-cm\(^{-2}\) in our conditions. The electrons emitted from the current source will follow the field lines, but the total current (including the diamagnetic currents on the order of ten percent of the parallel

Figure 7. Axial current density displayed on three planes and at two different times in the experiment. \( J_z \) at \( \delta t = 108 \mu s \) is shown in (a)–(c), and (d)–(f) display current density acquired at \( \delta t = 3.34 \) ms. The current source is at \( \delta z = 0 \) and the anode is at \( \delta z = 10.86 \) m. The Alfvén transit time along the length of the volume is 155 \( \mu s \) and an ion gyro period is 13 \( \mu s \).
Also shown are contours of the magnetic vector potential \( A_z \) (perpendicular to the tilted plane). The contour values are indicated on the diagram and the color bars. The data is shown at \( \delta z = 2.56 \text{ m} \) and \( \tau = 390 \mu \text{s} \) during the reconnection phase. The data plane has been tilted by 3 degrees about the \( x \)-axis and the current density and vector potential normal to this plane were recalculated using the 3D data set. Two islands in the current are clearly visible but this is not reflected in \( A_z \).

Figure 9 shows a single frame of a movie taken with a fast framing camera. The bright lines in the images are due to light from fast electrons striking ions and neutrals in the background. The currents resemble those of \( \delta z \) near its edge. The camera was outside the vacuum system. In this limit, the tearing mode growth rate derived from kinetic theory [17] is of the form [24],

\[
\frac{\gamma}{\Omega_c} \approx \frac{B_z}{B_0} k \lambda \sqrt{\beta_e} \frac{m_i}{m_e} \left( \frac{\Delta \lambda}{\lambda} \right)^3 \left( \frac{d_s}{\lambda} \right)^3
\]

where \( \Omega_c \) is the ion cyclotron frequency computed with reconnecting component of the field \( B_x \approx 20 \text{ G} \), \( \lambda \approx 3d_e \) is the half-thickness of the layer, \( B_x \approx 200 \) is the guide field, \( \beta_e \approx 0.1 \) (assuming \( T_e = 10^4 \text{ eV} \) and \( n \approx 10^{19} \text{ cm}^{-3} \)) and for helium plasmas \( m_i/m_e = 7344 \), and \( \Delta' \) is the tearing mode index which characterizes the drive from the equilibrium [8] and has been computed for force-free current sheets [17, 24]. Typically, the fastest-growing modes within these thin sheets [24] occur for the wave vector \( k \lambda \approx 0.5 \) corresponding to \( \Delta' \lambda \approx 3 \), resulting in \( \gamma \approx 3.6 \Omega_{ci} \approx 1.7 \times 10^2/\text{sec} \) for expected growth rate. This corresponds to growth time \( 1/\gamma \approx 5.8 \mu s \). This time scale is roughly consistent with experiment, since visible tearing in the current sheet is first apparent at \( \approx 15 \mu s \). Furthermore, the number of flux ropes observed in the experiment corresponds to a wave vector \( k \lambda \approx 0.47 \), which is close to the expected value.

There is a great deal of dynamics associated with the current sheet, but one of the first things that come to mind is the role of magnetic field line reconnection. This process can cause ion jetting, plasma flows, and electron heating among other things when magnetic field energy is converted into other forms [25]. The electrons and ions in the current sheet do get hotter and plasma flows occur, but is this the result of reconnection alone? We first explore reconnection in this experiment and later compare it with other forms of energization.

Figure 8. Shaded surface showing current density in amps-cm\(^{-2}\). Also shown are contours of the magnetic vector potential \( A_z \) (perpendicular to the tilted plane). The contour values are indicated on the diagram and the color bars. The data is shown at \( \delta z = 2.56 \text{ m} \) and \( \tau = 390 \mu \text{s} \) during the reconnection phase. The data plane has been tilted by 3 degrees about the \( x \)-axis and the current density and vector potential normal to this plane were recalculated using the 3D data set. Two islands in the current are clearly visible but this is not reflected in \( A_z \).

Figure 9. False-color image of the current sheet taken with a fast framing camera. The blue region is brightest, corresponding to flux ropes that the sheet has broken up into. The exposure time is \( 1 \mu s \) and the picture was taken at \( \tau = 2.28 \mu s \). The object at the rear is the mesh anode (figure 2) located at \( \delta z = 11 \text{ m} \). The anode is square, 25 cm on a side.
4.3. Magnetic field line reconnection

When reconnection occurs several things happen. There is a change in the local magnetic topology, induced electric fields caused the changing magnetic field appear, and energy is released. The reconnection process is accompanied by the generation of a QSL [26]. As we will see, QSLs can form without reconnection as well. In this experiment as well as in the case of flux ropes near the Sun, the reconnection process is fully three-dimensional [27], and simplified pictures of X points no longer work. To illustrate that something is happening, figure 11 presents the transverse magnetic field on a plane \( z = 4.48 \) meters at several times. (Also see figure 4 in reference [6].) Figures such as this can be somewhat deceiving, however, because as shown above the magnetic islands can be strongest in planes that are tilted with respect to the guide field and reconnection geometries should be searched for within them.

The reconnection event in the transverse plane of figure 11 is not representative of the reconnection geometry. Reconnection is fully three-dimensional and occurs between field lines stretched along the magnetic field. Figure 12 shows field lines at the same location (\( z = 4.48 \) m) but with the third component of the magnetic field of the flux ropes as well as the background magnetic field included. This component is larger than the transverse fields, and when it is included the reconnection geometry is not evident. Now the X and O points are completely invisible.

The tearing is fully three-dimensional, as seen in figure 12 (which is a 3D picture).

One method used to identify reconnection regions is the identification of QSLs (sometimes called squashing factors). QSLs [28], are flux tubes of field lines that start off very close to one another on a plane transverse to the background magnetic field (plane A) and wind up very far apart on a plane further away (plane B). The value of the QSL is indicative of the strength of the wandering. For example, if \( Q = 100 \) then field lines at a remote location can be 50 times farther apart than at the start. Field lines have separated if \( Q > 2 \), but only larger values are of great interest. It is believed that a large \( Q \) will occur if there is a region of magnetic field line reconnection somewhere between A and B. In this case the field (or transverse components of it) can rapidly change from field line to field line. \( Q \) can also be large if there are large gradients in the current but no reconnection. In figure 11 there is certainly a reconnection, at \( t_0 \) and at \( z = 4.48 \) m. What is the QSL in this case? In this experiment, when magnetic islands such as those in figure 11 are clearly visible, the value of \( Q \) is very small; in this case, \( Q \sim 6 \) and peaks in the center of the current sheet. Figure 13(a) displays contours of \( Q \) at two
times. At the earlier time, which corresponds to that of figure 11, the QSL is found inside the current sheet, possibly because of internal magnetic islands. The value of $Q$ is small, and can be associated with a low level of reconnection. In simulations of solar loops $Q$ (also with values on the order of 5) were observed between the currents associated with two flux ropes; [29], they were also observed in laboratory experiments [30]. Also shown in figure 13 is the reconnection rate or the integral of the inductive electric field along field lines, $\varepsilon = -\int \frac{\partial A}{\partial r} \cdot dl$, where $dl$ is the incremental distance along the local magnetic field.

In figure 13(a) the largest value of $\varepsilon$ (1 V) occurs between two locations at which $Q$ is largest. The X point in figure 11 is close to this location. At later times the value of $Q$ becomes much larger. In figure 13(b), acquired at $\tau = 860 \mu s$, $Q$ has grown to 20 but it is now largest at the edge of the current sheet, while the somewhat smaller reconnection electric field remains largest in the center. For $\tau > 1.5 \text{ ms}$, $40 < Q < 120$, and $Q$ remains largest at the narrow edges of the current sheet. There does not seem to be evidence for reconnection at these later times, and the enhanced value of $Q$ is associated with large gradients in the current at these locations.

4.4. Electric field

When magnetic flux changes in a reconnection event it is by definition accompanied by a change in the vector potential $\mathbf{A}$. Using equation (1) the vector potential throughout the data volume was calculated from the 3D current. The inductive electric field, $\mathbf{E} = -\frac{\partial \mathbf{A}}{\partial t}$, is then evaluated. This must be compared to the electric field from gradients in the plasma potential. The plasma potential, $V_p$, was measured with an emissive probe, and the electrostatic component of $E$ derived from $E_e = -\nabla V_p$. In some 2D models of reconnection this is ignored and $E$ is purely inductive at the x-line. The emissive probe measured the plasma potential at the same locations and times that the magnetic field was collected. Figure 14 is an example of $E_e$ close to the source of the sheet.

The two components of the electric field are measured in different ways. The electrostatic component is measured...
using an emissive probe, and the electromagnetic component is derived from the measured three-dimensional magnetic field. The temporal dependence of both components is shown in figure 15 at the same transverse position \((x, y) = (2.0, 3.5 \text{ cm})\) where \((0, 0)\) is the center of the current sheet. In figure 15 the field is shown on two planes two meters apart. The large fluctuations are associated with gross motion of the sheet.

The total electric field on two planes as a function of time is shown in figure 16. The electric field can be positive (pointing towards the source of the current sheet) and the sign of \(\mathbf{E} \cdot \mathbf{J}\) can change.

### 4.5. Plasma resistivity

On the plane \(\delta z = 4.48 \text{ m}\), the largest contribution to the axial electric field is from the plasma potential, not the induced electric field. Later in time \((t > 1 \text{ ms})\) the current sheet becomes unstable to a low-frequency mode, and this is reflected in oscillations of the electric field. When the oscillation is fully developed the induced electric field is comparable to the electrostatic field. This late time mode has little to do with reconnection. A characteristic of this mode is rotation of the entire sheet in the \(x\)-\(y\) plane at a frequency of 2.7 kHz. Ohm’s law for plasmas can be written as:

\[
\frac{m_e}{ne^2} \frac{\partial \mathbf{J}}{\partial t} = \mathbf{E} - \frac{1}{ne} \nabla P_e + \mathbf{u} \times \mathbf{B} - \frac{1}{ne} \mathbf{J} \times \mathbf{B} - \eta_1 \mathbf{J}_1 - \eta_2 \mathbf{J}_2.
\]

Here \(\mathbf{J}\) is the plasma current, \(P_e\) the electron pressure, \(\mathbf{u}\) the plasma drift velocity, and \(\eta\) is the resistivity. The plasma

**Figure 12.** Magnetic field lines at \(\tau = 236 \mu\text{s}\). This image was taken from the same vantage point as figure 11. There is no indication of an X point. The marker arrow on the upper right is located at \(\delta z = 4.48 \text{ m}\). This is an anaglyph best viewed with red/blue glasses.

**Figure 13.** Shaded surface of \(Q\), showing the squashing factor at two times and contour lines (in volts) of the induced electric field integrated along 3D magnetic field lines.
density and electron temperature were not measured throughout the volume, but measurements at fixed points on several planes indicate that the axial pressure gradient for this experiment was small and we neglect the pressure gradient term. This is not the case for intense flux rope experiments where the ropes heated the electrons and completely ionized the background plasma. The sheet current was one-tenth that in past flux rope experiments \[31\], and the density and temperature in the sheet were much greater than the background plasma, resulting in very small pressure gradients. The plasma flow perpendicular to the background field is at most 0.6 \(c_s\), as seen in a previous experiment \[31\], and the equivalent perpendicular electric field from the \(\vec{u}_\perp \times \vec{B} \approx \vec{u}_\perp \times \hat{B}\) term is on the order of 1.8 V cm\(^{-1}\). The Hall term \(\frac{1}{\mu} \vec{J} \times \hat{B}\), is at most 0.03 V cm\(^{-1}\) perpendicular to the background field. This is three percent of the potential gradient in figure 14. The flow term in the equation for \(\eta_{\parallel}\) cannot be ignored. In this data run we did not acquire volumetric flows, and the study of perpendicular resistivity will be done in future work. In fact one must be careful in evaluating the perpendicular resistivity without including flows and the large perpendicular electric field from space charge gradients. Experiments that leave these out could be in error \[32\]. One must also be careful in assuming that one can calculate \(\nabla V_p\) using differences in the floating potential and then deduce electric fields, because in the presence of currents or non-Maxwellian electron distributions they are not the same \[11\].

There is no Hall term for parallel flows, and one can evaluate the parallel component of equation \(4\). In this case the parallel Ohm’s law can be reduced to:

\[
\eta_{\parallel} \vec{J} = -\nabla \phi - \frac{\partial \vec{A}_{\parallel}}{\partial t},
\]

Figure 14. The electrostatic component of the field in a data plane 128 cm from the slot. The potential in the center of the slot current (which has tilted) is negative with respect to the background plasma. The largest transverse field is of order 1 V cm\(^{-1}\), as shown by the marker arrow in the figure. The electric field parallel to \(B_0\) is two orders of magnitude smaller. Further away the current sheet becomes structured, as does the electric field.

As there are volumetric measurements of the terms on the right of equation \(4\), the plasma resistivity can be evaluated and compared to the classical Spitzer resistivity,

\[
\eta_{\text{Spitzer}} = \frac{e^2 \ln \Lambda \left( \frac{K T_e}{m_e} \right)^{\frac{1}{2}}}{3\varepsilon_0 (2\pi)^{3} m_e},
\]

which is due to Coulomb collisions.

In the center of the current sheet \(\eta_{\text{Spitzer}} = 1.8 \text{ m}\Omega \text{ cm}\). Figure 17 shows the plasma resistivity at time \(\tau = 940 \mu s\) calculated from the plasma current and electric field data using

\[
\eta_{\parallel} = \frac{\vec{E}_{\parallel} \cdot \vec{J}_{\parallel}}{J_{||}^2},
\]

where the electric field is the right-hand side of equation \(5\).

Before calculating the resistivity it was necessary to carefully align the data sets from different axial (\(z\)) planes. The initial alignment was done by a transit looking down the machine axis. It was discovered that this was good only to 2–3 mm in the transverse direction \(\delta \xi\). The misalignment caused \((\vec{E} \cdot \vec{J})_0\) to be negative in several regions. The resistivity was calculated using the electric field and current parallel to the local value of the magnetic field. The background magnetic field accounts for most of this but there is a contribution of several percent due to the currents in and around the sheet. The alignment did away with most of this contribution. The plasma potential at early times exhibited a narrow negative peak in the center of the current channel, which was used to determine the final alignment. The potentials were also aligned to the center of the narrow plasma current from the slot. In addition, if the power was smaller than 10 mW cm\(^{-2}\) (the power in the center of the current sheet is 20–100 times this) it was considered to be in
the noise level and the resistivity was set to the Spitzer value. This prevented division by zero when evaluating equation (7).

Figure 17 displays the resistivity as a shaded surface along with contours of current parallel to the background field \( \mathbf{j} = \mathbf{j}_z \). The current is negative, denoting that it is carried by electrons. The uniform background magnetic field points out of the page. The resistivity is largest on the edge of the current sheet, not where the current itself is largest.

What is the cause of the anomalous resistivity in the current sheet? It is not reconnection; there is not enough annihilated flux. The obvious question is, what causes the anomalous resistivity shown in figure 17? The resistivity persists throughout as long as the current is present and reconnection is sporadic. Thus, it seems that magnetic field reconnection is probably not responsible for the enhanced value of \( \eta \). Figure 18 shows the parallel resistivity at a given spatial location on two different planes. The resistivity varies greatly from point to point (figure 17) but in the current channel it is significantly larger than the classical value. The oscillations due to the flux ropes are obvious on further planes.

The regions of larger resistivity are localized to the current sheet. In fact, the average resistivity throughout the volume is only slightly above classical, as shown in figure 19.

The power injected into the current sheet is over 45 kW (190 A, 240 V). The sheet is pulsed on for 3.4 ms and the total energy supplied is 155 J. Only a fraction of this is absorbed in the plasma. Figure 20 shows the power within the plasma from the evaluation of \( \int E \cdot j \, dV \) integrated over the volume the data was acquired in. Here \( \overline{E} \) consists of the sum of electrostatic and electromagnetic components. This is roughly one percent of the injected power, which is no surprise since most of the voltage drops on sheaths near the cathode and anode. The power supply voltage does not reflect the voltage
in the system. The total energy in the current system is 3.91 joules over its lifetime. The efficiency (Power \( E_j \) \( J \) /Power discharge) is 2.5%. The particle energy in the current sheet, \( \int nK(T_e + T_i) \) dV, is approximately 1 J or 25% of the energy in the current. The energy also takes other forms such as a considerable amount of ultraviolet light generated, plasma flows, and energy in the fields associated with the current sheet. Most of the energy flows through the system, as the data indicates that the energy confinement time is \( \tau_p = 1 \) J/25 kJ s\(^{-1}\) = 40 \( \mu \)s. The Coulomb ion collision time \([33]\), is on the order of 5 \( \mu \)s.

Data was acquired over an axial range of 7 meters, but the volume containing the current sheet was 11 meters long, so the total power is 40–50 percent higher.

In Figure 20 it is apparent that the power has a low-frequency oscillation of approximately 3 kHz, which is the same as the flapping motion of the flux ropes. A higher-frequency oscillation that persists throughout the time the current is on is barely visible. Both the high- and low-frequency oscillations are seen in the parallel resistivity, as well as in the electric and magnetic fields. The input power is much larger than this.

---

**Figure 16.** The z component of the total electric field \(- \nabla \phi_p - \frac{\partial \mathbf{A}}{\partial \mathbf{r}}\) at a given spatial position on two planes, and as a function of time. Close to the source at 64 cm the electric field and current (not shown) are parallel. The electric field oscillates when the current sheet starts to flap. The error bar is based on the RMS value of the plasma potential over 10 shots at each location.

**Figure 17.** Plasma resistivity shown as a shaded surface on a transverse plane (\( \delta z = 1.28 \) m) at \( \tau = 940 \) \( \mu \)s. Contours of current density are displayed to show its relationship to the resistivity.
The low-frequency mode, which reflects the current sheet motion, is visible in both the electric and magnetic field measurements. These were done using very different detectors. The frequency spectra of these quantities have a sharp peak at 3 kHz. Figure 21 shows the spatial location of the 3 kHz component of \( E_y = -\nabla V_p - \frac{\partial A}{\partial x} \). The other components of the electric field have the same spatial distribution.

The resistivity, QSL, and field lines are summarized in figure 22. The magenta surfaces are \( Q = 40 \) at \( \tau = 1.42 \) ms after the sheet begins. Magnetic field lines within the QSL are drawn in black. It is clear that they diverge as one moves away from \( \delta z = 0 \) (the origin of the sheet).

Some current lines (striped filaments) are included. The current sheet becomes wider at larger distances. The QSL is largest at the top and bottom of the current sheet. At this later time there is no appreciable magnetic field line reconnection.

### 4.6. Motion of the current sheet

About 1.2 to 1.5 ms after the current sheet forms it begins to rotate violently. This is reflected in the fields shown in figures 15 and 16, the power dissipation, and the global resistivity (figure 19). This motion is also illustrated in a hodogram (figure 23) which depicts the position of two magnetic field lines which start on a plane 7.04 m away at \( x \) and \( y \) locations (1.3, 3.0) and (3.0, 1.75) respectively. The
endpoints of the field lines are followed for 2.7 ms. The field lines start at $\delta z = .64$ m and are displayed on a plane at $\delta z = 7.68$ m. They rotate counterclockwise, which is in the electron diamagnetic drift direction.

The motion of the current sheet, illustrated in figure 24, is similar to previous experiments with two or three circular cross-section flux ropes [19]. The plane on the upper right is a surface plot of the current density. The field lines near the start of the current sheet are in the same position because the current sheet is frozen into the slot at $\delta z = 0$. If one examines the field lines on the lower left one sees the red and blue lines have moved upward (and rotated) at the later time. They have also spread out in space. They are tracked by the lower, magenta field lines. This motion is in accordance with the individual field line motion shown in figure 23. The QSL moves with the field lines but it is always much larger on the narrow edge (top and bottom) of the current (figure 22). The QSL is large because the field lines diverge spatially at the top and bottom, not because there is magnetic field line reconnection. The contours of either component of the electric field at these times and far from the source, or anywhere else in the volume, show that they do not peak where the QSL is largest.

This is an important observation, because the QSL is often identified with magnetic field line reconnection in the solar physics community. Figure 25 depicts the QSL at four different times after the current sheet begins to move as a whole. The QSL is rendered as four surfaces, which are colored differently at each time step. In all cases $Q = 25$. The perpendicular magnetic field is shown on two planes (at the first time step). The QSL moves sideways. The current sheet twists about itself and the QSL follows.

5. Summary

A thin current sheet ($\psi = 20$) was created in a background-magnetized plasma and observed to tear into filaments. The tearing happens rapidly and some corrugation of the current can be observed within several Alfvén transit times, $\tau_A$ (50 ns in the background plasma), along the system length ($L = 11$ m). Reconnection happened throughout the plasma volume but did not lead to strong QSLs or anomalously high resistivity. After about 1.5 ms the distorted current sheet exhibits large-amplitude, low-frequency (3 kHz) oscillations as it rotates about the background field. At this time two large QSLs ($20 < Q < 120$) form at the narrow edges of the current sheet. They are not connected to a reconnection event in any obvious way; rather, they result from a large divergence of field lines in the current sheet as it moves and spreads across the background field.

The parallel resistivity was measured using emissive and magnetic probes. It was localized in the gradients of the current and was one to two orders of magnitude above the classical value. The resistivity averaged over the entire plasma, however, was less than twice classical. The enhanced resistivity is most likely due to current-driven instabilities ($f_{\text{lower Hybrid}} = 6.5$ MHz, $f_{\text{pe}} = 9$ GHz), but the diagnostics we employed did not have the bandwidth to detect them. One such instability previously studied in a current sheet is ion sound. The ion acoustic instability occurs when the electrons drift with respect to ions [34],

$$\frac{v_{\text{drift}}}{c_s} > 1 + \frac{m_i}{m_e}\left(\frac{T_i}{T_e}\right)^{3/2}e^{-\frac{T_e}{2T_i}}.$$  (8)

In this experiment the ion temperature was not measured, but it has been measured using laser-induced fluorescence in a previous flux rope experiment [30]. The ratio of $T_e/T_i$ is in the range of 4–10 in this experiment. The electron drift speed determined from the measured current density in the sheet is such that $\frac{v_{\text{drift}}}{c_s} \approx 6–10$, which is in the range to make the instability switch on. Investigation of high-frequency instabilities will be part of future work in this area.

Figure 20. The Ohmic dissipation term derived from the measured current and electric field. This is integrated over the plasma volume. Note that the contribution from the parallel current is twice that of the perpendicular currents.
Figure 21. Spatial location of the 3 kHz oscillation of one component, $E_y$, of the electric field on a plane 4.48 m from the source of the neutral sheet. Close to the origin of the current the oscillation is smaller in amplitude and located within the current sheet. At larger distances the oscillation spreads over the plane.

Figure 22. The QSL ($Q = 40$), is shown along with $\vec{B}$ on a plane 4 m from the origin of the sheet. An isosurface of enhanced resistivity is shown. At this time the resistivity is large in the first two meters and largest at the edge of the current channel (see figure 17).
In a previous reconnection experiment involving a current sheet the ion sound instability was identified [35]. All the terms in the force density equation were measured, and it was discovered that an additional drag term of the form $\nabla \cdot \mathbf{V}$ had to be included to balance the equation. The scattering was attributed to short-wavelength $\lambda \approx \lambda_0$ ion sound waves. The frequency spectra of electrostatic fluctuation extended up to the ion plasma frequency. A two-probe cross-correlation measurement with sub-millimeter resolution established that the waves obeyed the ion acoustic dispersion relation. That experiment, now 32 years old, was done at much lower density and background magnetic field. The ions were not magnetized in the current sheet, but it is entirely possible the same physics is going on in this experiment.

The current necessary for a cylindrical plasma current to kink was first derived by Kruskal [36], and Shafranov [37], to be

$$I_{KS} = 10 \frac{\pi \alpha^2 B_z}{L},$$  \hfill (9)

where $I_{KS}$ is in A, $B_z$ in G and length in cm.

This predicts that a flux rope with an area of the sheet in this experiment (20 cm$^2$) and a background magnetic field of 200 G will be unstable for currents above 37 A. The current in this experiment is ten times this. The geometry is not that of a cylinder, but there is no question that this is well above the necessary area to kink. Once the plasma does kink it is difficult to predict what will happen next. In the analysis of the motion and kinking of a slender plasma column, Ruytov et al
[38], predict that a column with one fixed boundary and a second free boundary will oscillate with a frequency of

\[ \omega = \pm \frac{\pi V_{A2}}{2L} \left( 1 - \frac{2k_0L}{\pi^2} \right), \]

where \( V_{A2} = \sqrt{2} V_A \) at the kink mode threshold. Under our conditions the real part of the frequency is 2.7 kHz, which is close to what is observed. The current sheet is not a slender tube, and direct comparison with equation (10) can only yield approximate solutions. Ruytov et al [38] also show that in the linear state a current tube is wound around a cylinder whose radius increases axially from the fixed boundary condition. This is also similar to what is observed. After kinking the motion of the sheet becomes quite large far away from its source, and one expects that it could only be modeled by an appropriate computer simulation.

6. Conclusions

There is an MHD (plasma fluid) equation often used to describe plasmas. This equation is similar to the Navier–Stokes equation used to describe moving liquids (fluids) and the atmosphere. The plasma equation contains the magnetic field, and Navier–Stokes does not, but both are highly nonlinear and not amenable to analytic solutions. The kinetic theory of plasmas, which is rigorously correct, is even worse. Any one of these theories can lead us down the road to turbulence, or worse still chaos [31]. How do we cope? One tactic is to ignore the details and instead calculate the spectrum of modes, often looking for power laws. Different theories predict different power laws. In situations where detailed measurements are not possible, such as in fusion plasmas, or in space where there may be only one to five satellites, power laws have achieved the height of popularity. Often a frequency spectrum (of density, magnetic field, etc) will have several regions, each with a different power law. What this means is hotly debated. The other side of the coin has been a movement to construct simple models, which can be drawn as a cartoon with only several lines. This was a good start 30 or more years ago, but not much has changed in that department. One of the first reactions a plasma physicist may have when told that you are studying the evolution of a current sheet is that it is a reconnection experiment.

This experiment illustrates that the phrase ‘magnetic reconnection’ by itself cannot properly describe the rich variety of processes that occur when a current sheet tears into flux ropes. One must be aware of how the currents close in three dimensions and whether many phenomena are due to current-driven instabilities and not reconnection alone. They happen in tandem. In this experiment, higher-frequency phenomena such as Whistler and Langmuir waves were not measured. They are most likely present. Langmuir waves are often present when electron distribution functions are non-Maxwellian or when beamlets are present. They are often accompanied by very small structures such as electron phase space holes [39]. Phase space holes have localized electric fields, which can scatter electrons. Whether these high-frequency phenomena such as Whistler and Langmuir waves are often present when electron distribution functions are non-Maxwellian or when beamlets are present. They are often accompanied by very small structures such as electron phase space holes [39]. Phase space holes have localized electric fields, which can scatter electrons. Whether these high-frequency phenomena affect lower frequency phenomena like those in this study is an open question.

Three-dimensional magnetic and electric fields are complicated and change rapidly. If one were to examine the time history of the magnetic field components in this experiment at three or four locations, they would not have a chance of reconstructing the field topology from the data. However, this is routinely done in space plasmas where the interpretation of data is highly model-dependent. Data not in accordance with a model is rejected as uninteresting, or dismissed because...
maybe something else is happening. This is dangerous in the
sense that progress is impeded, sometimes for decades. Instead of truth, one settles for ‘truthiness.’ Hopefully in the future nano-satellites will be deployed by the thousands, and phenomena in space will be revisited in a new light. However, well-diagnosed, reproducible laboratory experiments are here now. They offer not only the chance to study a sheet of current, but a chance to revisit the field of nonlinear fluids from a fresh perspective.

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