Experimental observations of the tearing of an electron current sheet

W. Gekelman and H. Pfister
Department of Physics, University of California, Los Angeles, California 90024-1547
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A neutral magnetic sheet, in which the current is carried mainly by the electrons, is set up in a laboratory plasma. By forcing the current through a thin slot, the ratio of the length to height \( t \) of the sheet may be varied; the current is observed to tear when \( t \gtrsim 30 \). The structure of the magnetic islands and their associated currents is fully three dimensional, although a linear two-dimensional theory gives a very good estimate of the tearing mode growth time. Tearing is accompanied by the generation of significant Hall currents, and magnetic disturbances are observed to propagate at the whistler wave speed.

I. INTRODUCTION

A problem of fundamental interest in plasma physics is the behavior of a current channel in a magnetoplasma, especially if the self-magnetic field of the current is comparable to the background field. A plasma, unlike a solid conductor, is free to move and change its properties, making current a dynamic entity. Consider a sheet of field aligned current propagating as shown in Fig. 1(a). This is unstable to the tearing mode [Fig. 1(b)] which reorganizes it into a series of current filaments. In the plane transverse to the flow, the magnetic field topology consists of a series of X and O points; this is a lower magnetic energy state than the sheet. Finally, one may consider adjacent pairs of current filaments. If the self-field of each channel is significant they attract one another via the force \( \mathbf{J}_1 \times \mathbf{B}_2 \), where \( \mathbf{B}_2 \) is the magnetic field due to channel 2 at the location of \( \mathbf{J}_1 \). In an idealized situation consisting of an infinite chain of equally spaced filaments, the forces would balance. In practice, some channels carry more current than others and coalescence should occur. Many of the theoretical models and simulations of either the tearing mode or of coalescence have used a magnetohydrodynamic (MHD) picture. In the case of tearing, a number of models with differing boundary conditions, geometries, and plasma parameters have been developed. Since these models predict widely varying growth rates for the instability, it will be instructive to briefly review them and single out what is most applicable to our experimental solution. Finally, none of the work in tearing mode theory of coalescence is fully three dimensional; however, the experiment is. As the data are presented these differences will be highlighted. The first generation of tearing mode theories were for nearly MHD plasmas.\(^1,2\) They considered a planar geometry with antiparallel field lines as in Fig. 1(a). As seen in the figure, initial conditions consist of a thin resistive layer in the center and a periodic electromagnetic perturbation across the current sheet. If the plasma is resistive in the central region where \( \mathbf{B}_1 \rightarrow 0 \), and flux is lost as a result of field line reconnection, the perturbation can grow. This, in turn, results in greater field diffusion into that region of space. The tearing mode growth rate in these theories depends upon the Alfvén velocity, i.e., how rapidly magnetic fields can move into the resistive region. The ions in these models are magnetized every-

where outside of the sheet, and the tearing mode frequency is always less than the ion gyrofrequency. In general, if there are no large electrostatic fields within the tearing layer, the problem is determined by boundary conditions at the layer. In fact, the current within this layer is carried exclusively by electrons and the problem may be attacked by a single fluid MHD theory,\(^1-3\) or with a kinetic approach.\(^4-7\) In either case the growth rate is proportional to the product of the electron thermal velocity and the square of the electron skin depth. Characteristic ion scale lengths do not enter. This is true even with magnetic shear, i.e., the inclusion of an axial magnetic field (\( B_z \)). A considerable body of literature exists on effects such as the inclusion of a normal (\( B_z \)) component of the magnetic field,\(^3\) high beta effects,\(^8,9\) kinetic theories,\(^10\) and anisotropic particle distributions.\(^7\) It is best to defer the discussion of this literature until the experimental data are presented. However, we note that much of the work centers around finding conditions in which growth is rapid enough to explain spacecraft or fusion data. Over the past few years a quiescent high beta laboratory plasma has been used to study magnetic field line reconnection in great detail.\(^8-14\) A striking result of this work is that what appears to be a geometrically stable neutral sheet (over many Alfvén transit times) is

![FIG. 1. Three possible magnetic field topologies (solid lines) resulting from (a) current that flows in a neutral sheet and (b) a tearing mode that reorders the neutral sheet current into filaments, each of which produces a magnetic \( \times \) point. The \( \times \) points separate neighboring current channels (c), a circular cross section of current that could be the result of a merging of pairs of filaments.](image-url)
FIG. 2. Schematic of the experiment. The dc discharge plasma is produced when the cathode (at \( y = 0 \)) is biased negatively. A plasma current, which flows from cathode to a grounded anode, is induced by external currents flowing above and below the column in the reconnection plates. The slotted grid at \( y = 100 \) cm is used to constrict the plasma current. The magnetic probe may be moved anywhere in the volume bounded by the plates, the anode, and the grid.

seething with pressure anisotropies, unstable waves, anomalous resistivity, magnetic turbulence, and runaway electrons. On the other hand, the spontaneous tearing of this current layer into multiple X and O points was never observed. A clue came from a series of computer simulations by Leboeuf et al.\(^{15} \) They found that the stability of the topology depended on the crossing angle of the transverse magnetic field lines at the X point which initially was at the center of their simulation. Experimentally this implies that a critical ratio of length to height, \( t = L_x/L_y \), of a neutral sheet must be exceeded for tearing to occur. (It turns out that this condition is embedded in the theory as well.) When the experimental apparatus was properly configured (Sec. II) to allow a variation of \( t \), fully three-dimensional tearing was observed. Section III details the experimental observations with emphasis on the three-dimensional currents and magnetic fields. A comparison of the measured growth rate with theory is the subject of Sec. IV. The summary will attempt to place these results within the framework of current theory.

**TABLE I. Basic Parameters.**

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
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<tbody>
<tr>
<td>( n = 10^{12} \text{ cm}^{-3} )</td>
<td>( \text{argon} \times 10^{-4} \text{ Torr} )</td>
</tr>
<tr>
<td>( \delta = \epsilon/\omega_{pe} )</td>
<td>0.5 cm</td>
</tr>
<tr>
<td>( f_{pe} )</td>
<td>10 GHz</td>
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<td>( v_{pe} )</td>
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<tr>
<td>( kT_e = 7-11 \text{ eV} )</td>
<td>( kT_i \approx 2 \text{ eV} )</td>
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<td>( f_{pe} )</td>
<td>30 MHz</td>
</tr>
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<td>( f_{pe} )</td>
<td>56 MHz</td>
</tr>
<tr>
<td>( \Lambda_{nmf} \approx 2 \text{ m} )</td>
<td></td>
</tr>
<tr>
<td>( \Delta \approx 30 \mu \text{m} )</td>
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**II. EXPERIMENTAL ARRANGEMENT**

Figure 2 is a schematic of the experimental configuration. The 2 m long argon discharge plasma is generated between a negatively biased cathode (1 m diam, \( V_{bias} = -40 \text{ V} \)) and a grounded anode 1 cm in front of it. It is uniform, quiescent (\( \delta n/n = 1\% \)), essentially collisionless (\( \Lambda_{nmf} \approx 2 \text{ m} \)), and highly reproducible in pulses of duration \( t_p \) = 5 msec, repeated every \( t_s = 1.5 \text{ sec} \). Plasma parameters are given in Table I. The plasma is immersed in a uniform background magnetic field of \( B_y = 20 \text{ G} \). This models the component of the interplanetary magnetic field that is at a right angle to the neutral sheet. A time varying \( (t_{rise} = 100 \mu \text{sec}) \) transverse magnetic field \((0 < B < 20 \text{ G})\) is applied by passing a current \( I_{vac} \) through two insulated plates on the plasma boundary. The induced electric field \( E_y \) causes an electron drift from the cathode resulting in a plasma current \((I_x \approx 1 \text{ kA})\) and neutral sheet formation. Current is collected on the unbiased end anode. The thickness of the neutral sheet is controlled by biasing a slotted grid in the center of the plasma column to a large negative potential \((-V_y = 200 \text{ V} > kT_e/e)\) after formation of the sheet. It is observed that a significant amount (70\%) of the current is funneled through the slot and is thus constricted to the slot width.

Magnetic field measurements are obtained with three mutually orthogonal magnetic loop probes mounted on a coaxial telescoping axial shaft that allows them to be placed anywhere within the plasma volume. In this experiment magnetic fields are obtained in situ at 5440 spatial locations.
and 1024 time steps ($dt = 50$ nsec) in the volume to the left of the grid ($y > 100$ cm). At each location an ensemble average of ten experimental discharges is taken for each component. Data are acquired and stored with a state-of-the-art acquisition system. The analog traces from the probes are digitized with 20 MHz analog to digital converters housed in a CAMAC system that is serviced by an LSI 11/23 computer. The LSI serves as a slave to a VAX 11/750 computer and is linked to it by a specialized high-speed (230 kByte/sec sustained transfer rate) network. The data are placed in large arrays that may be accessed in real time (the computers must do all calculations on a faster time scale than the experimental repetition rate) by an on-line array processor (for digital filtering and averaging) and other concurrent computer processes (i.e., graphics). The VAX is linked via a second network to a graphics laboratory (Fig. 3) that consists of VAXstation II, DMA linked to a Raster Technologies 3-D solid modeling engine (model 1/380). This system is able to render the data into the color images presented in this paper. The VAX is also networked to the MFE and San Diego Supercomputer centers where CRAY computers are used for manipulation of the final data sets.

III. EXPERIMENTAL OBSERVATION

Figure 4 shows the components of $B$ detected by a probe positioned 30 cm from the slot ($y = 130$ cm). At $t = 0$ the current flow in the plates bounding the plasma begins. The constriction grid with a slot width $l_s = 1$ cm is biased negatively at $t = t_g$ ($t_g = 17$ $\mu$sec); a dramatic change in the relative magnitude of the field components is then observed. At this position the magnitude of $B_y$ dominates all other time varying components within several microseconds. From a single vantage point such as a magnetic probe in a large chamber, or on a satellite in space, little can be said about large changes in topology. "Signatures" that a satellite has crossed a neutral sheet (or vice versa) are model dependent.

![FIG. 4. A ten discharge ensemble average of magnetic field data at a fixed position within the device. The reconnection experiment begins at $t = 0$; the slotted grid is biased at $t = t_g$. The three channels are recorded simultaneously at a 50 nsec/sample.](image)

![FIG. 6. Magnetic field lines at $t_g = 4.75$ $\mu$sec in four transverse planes for increasing values of $y$ ($y = 100$ cm is the slotted grid location).](image)

It is also not possible to visualize the three-dimensional field from inspection of several (or many) of the thousands of such traces acquired throughout the volume. For purpose of orientation it is instructive to examine data in several planes before the fully three-dimensional situation is presented. Figure 5 shows the transverse vector field ($B_x, B_z$)
close to the grid (i.e., at \( y = 105 \) cm) just after it has been biased \((t_b = 1.25 \mu \text{sec})\), but prior to tearing. The topology is that of a “classic” neutral sheet with a half-width of order 1 cm. The length to half-width ratio \((l_x/l_y)\) is about 30.

Magnetic field lines are drawn for four transverse planes at \( t_x = 4.75 \mu \text{sec} \) in Fig. 6. The existence of more than one island can clearly be seen at \( y = 110 \) cm and \( y = 130 \) cm [Figs. 6(b) and 6(d), respectively], and only one large island at \( y = 125 \) cm [Fig. 6(c)]. The case at \( y = 105 \) cm [Fig. 6(a)] is intermediate in that there is a suggestion that an island may appear to the left of the island centered at \((x,z) = (21.5)\) cm if one was to look upstream or downstream. These data were then synthesized into the three-dimensional image displayed in Fig. 7. The cyan axis drawn down the center is the \( y \) or device axis. The magenta markers delineate steps of 2 cm. The outer semitransparent blue surface is close to the separatrix between closed field lines that bound the islands within it and open lines that leave the plasma and encircle the current carrying plates. The inner red surface is the boundary of the innermost island(s). This diagram was generated from transverse field lines on 32 parallel planes. Although the diagram is qualitative, it shows the structure of the magnetic field to be complex, fully three dimensional, and exhibiting merging and spatial growth. It will be shown that the largest current flows within the islands, and its three-dimensional structure, reflect that of Fig. 7. Finally, the temporal growth of the islands at one transverse plane \((y = 115 \) cm) is shown in Fig. 8. Here field lines are superimposed on the vector field. The topology changes from an initial sheet-like structure to a two island configuration, and then to a single island. By viewing the intermediate time steps one clearly observes that the island on the left moves toward the island on the right and merges with its companion. One also sees that \( B \), near the islands is small \((\sim 0.5 \mu \text{T})\) in comparison with the guide field. The tearing takes place within a few microseconds.

The current may be obtained from the fields via \( \mathbf{J} = \left(1/\mu_0\right) \nabla \times \mathbf{B} \). Prior to tearing, the axial current \( J_y \) near the slot has a sheet-like appearance [Fig. 9(a)] and the bulk of the current flows within the neutral sheet. Microseconds later [Fig. 9(b)], the tearing current contours in a plane reflect the topology of the magnetic field (Figs. 6 and 7); the largest axial current density appears at the center of the islands.

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**FIG. 7.** Three-dimensional magnetic field topology at \( t_x = 4.25 \mu \text{sec} \). The viewer’s eyes are located on the cathode side of the slotted grid and above it. The data begins at the grid \((y = 0)\) and extends toward the anode.
When the constrictor grid is pulsed negatively the current sheet at all axial locations beyond it does not contract instantaneously. A thin tongue of current is observed to propagate away from the opening at the whistler wave speed (as shown in Fig. 10), which is a temporal series of axial current profiles taken in a vertical longitudinal plane at $x = 16$ cm. One must note that the disturbance is three dimensional, and current from adjacent $y$-$z$ planes moves in and out of the current displayed in Fig. 10 during this temporal interval. Structures within the current are visible, however, and can be tracked. One is delineated by a thickened contour line. Analysis of the data shows the total current is conserved. The advance of the current shown by the dashed line in Fig. 10 does not proceed at the Alfvén speed since the ions are not magnetized in this instance; it is, in fact, the whistler wave speed ($v_w = 3 \times 10^7$ cm/sec), about one-fifth of the electron thermal velocity. Although the current is carried by the electrons, its associated magnetic perturbation cannot move at the electron characteristic velocity. This makes the mechanism responsible for the current propagation a fascinating one for future studies.

The large jump in the axial component $B_y$ of the mag-

![Figure 8: Temporal development of the transverse ($B_x$, $B_y$) field lines at $y = 115$ cm. The transverse magnetic field vectors are superimposed. The magnitudes may be estimated from the 0.7 G marker arrow in Fig. 5. Each arrow shown here corresponds to a measurement position ($\Delta x = 2$ cm, $\Delta z = 0.5$ cm). In Fig. 5 data were spliced along the $x$ axis.]

![Figure 9: Current density in the plane before and after tearing. (a) Axial current, $J_z$, 5 cm from the slot just before the onset of tearing. The largest current flows within the neutral sheet. Contours are in A/cm$^2$. (b) $J_x$ at the same location after tearing. The current is concentrated in two filaments that coincide with the islands in Fig. 6. (c) Hall currents at $y = 105$ cm. The arrow at the lower left indicates magnitude.]

![Figure 10: Propagation of axial current in the $y$-$z$ plane at $x = 16$ cm. The center of the slot is at $z = 3.5$ cm. The step size on the vertical $z$ axis is 0.5 and 1 cm/div, respectively, on the horizontal $y$ axis. The darkened current contour of 2 A/cm$^2$ moves to the right at approximately the whistler wave speed.]


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magnetic field (Fig. 4) over a large portion of the volume implies that substantial Hall currents may be present. This can be seen in Fig. 9(c), which displays the $J_x$-$J_y$ vector field at the same time and plane as the axial current shown in Fig. 9(b).

A comparison of these figures indicates that the largest axial current ($J_x = 3.75$ A/cm$^2$) in the center of the island $(x,y,z) = (11.0,105.0,5.3)$ is approximately three times greater than the largest Hall current ($J_y \approx 1.3$ A/cm$^2$) at the upper edge of an island $(x,y,z) = (27.0,105.0,6.5)$.

These Hall currents surround the islands and are obviously responsible for the production of the time varying axial field which is of the order of 20% of the guide field and reinforces it.

The experimental data were used to generate a three-dimensional picture of the current after tearing (Fig. 11). Axial currents with densities of $J_y = 2.4, 2.0$, and $1.6$ A/cm$^2$ are shown as nested red, green, and blue surfaces, respectively. The outermost highly transparent (blue) surface indicates that a low level current flows everywhere and is structureless. The inner current surfaces exhibit holes, the magnetic islands (Fig. 7) straddle these gaps in current, i.e., the largest current flows through their centers. Hall currents calculated on the periphery of the three-dimensional contour $J_y = 2.0$ A/cm$^2$ are displayed on five transverse planes as purple arrows that encircle the axial plasma current. The currents spiral around the individual filaments; the average net current spirals as well. A red rectangle drawn to scale with, and positioned at the opening of, the constrictor grid is included on the lower right for orientation.

IV. THEORY

Theoretical work on the stability of a current sheet with respect to tearing started about 16 years ago with the aim of interpreting data from pinch experiments.$^{17,18}$ Since this is not a review article we will sketch a typical growth rate calculation that contains many features of the large body of MHD work. A fully kinetic theory developed by Laval et al.$^4$ gives a similar end result.

Outside the tearing layer which is of width $2z_s$, an electron fluid approximation is used. Ampère's law relates the perturbed vector potential $A_i(z)e^{i(kx - \omega t)}$ to the perturbed current in the current layer:

$$\frac{\partial^2 A_i}{\partial z^2} - k^2 A_i = J_{iy}. \quad (1)$$

The oppositely directed zero order magnetic field in Fig.
The axial current density averaged over $x$ [see Fig. 9(a)] as a function of position $z$ across the neutral sheet prior to tearing. Also shown is the hyperbolic secant current model used in the theory. The two match well for a current half-width of 1.3 cm.

$1(a)$ is modeled by $B_0 = B_0 \tanh(z/\lambda)$. The experimental data, as shown in Fig. 12, indicate a good fit to

$$J_y = -J_{y0} \text{sech}^2(z/\lambda),$$  

(2)  

with $\lambda = 1.3$ cm, $J_{y0} = (1/\mu_0)B/\lambda = 3.1$ A/cm$^2$ is in agreement with the current contours shown in Fig. 9(a) and the magnetic vector field in Fig. 5.

The MHD momentum equation without the inertial term

$$(1/\mu_0)J \times B = \nabla p$$  

(3)  

is integrated over the $x$ direction yielding $p_x = J_xA_j/\mu_0$, where $p_x$ is the perturbed pressure. The contribution from Hall currents is ignored. The $z$ component of Eq. (3) is $\partial p_z/\partial z = -J_zB_z$, when combined with Eq. (1) and the model magnetic field one arrives at

$$\frac{\partial^2 A_1}{\partial z^2} - k^2 A_1 = \frac{2A_1 \text{sech}^2(z/\lambda)}{\lambda^2},$$  

(4)  

with a general solution

$$A_1 = \alpha e^{-kz} \left(1 + \frac{\tanh(z/\lambda)}{k\lambda}\right) + \beta e^{kz} \left(1 - \frac{\tanh(z/\lambda)}{k\lambda}\right).$$  

(5)  

The boundary conditions on (1) require that the induced electric field $E_x$ and the perturbed vertical magnetic field component $B_{z0}$ vanish for large values of $z$ [i.e., $A_1(L) = 0$]. This is certainly true at $z = L$ ($L = 16$ cm is the position of the current plates]. The solutions (5) for $z > 0$ and $z < 0$ cannot be joined without a discontinuity $\Delta'$ in $A_1$ across the tearing layer. The growth rate in the case of magnetic shear is then determined to be

$$\gamma = \frac{v_{\text{th}}}{2l_s} k \delta^2 \Delta', \quad \Delta' = \frac{1}{A_1(z = 0)} \left(\frac{\partial A_1}{\partial z}\right)_{z_s}.$$  

(6)  

Here, $\delta = (c/\omega_p) = 0.5$ cm, $v_{\text{th}}$ is the electron thermal speed ($1.8 \times 10^6$ cm/sec, $kT_e/e = 11$ eV), and $l_s$ is the magnetic shear length

$$l_s = \frac{B_{z0}}{(\partial B_x/\partial z)_{z_s}}, \quad B_{z0} = 20$ Gauss guide field.$$  

(7)  

The $\Delta'$ term is the normalized jump in the vector potential which enters because magnetic energy is converted to particle energy within the tearing layer; no such thing happens in the MHD region surrounding it. With Eqs. (4)–(6) one obtains

$$\Delta' = \frac{1}{k\lambda^2} \left[ \frac{e^{-kz}(s^2 - k^2\lambda^2(1 + t/k\lambda))}{e^{-kz}(1 + t/k\lambda) + (\beta/\alpha) e^{kz}(1 - t/k\lambda)} - (\beta/\alpha) e^{kz}(s^2 - k^2\lambda^2(1 - t/k\lambda)) \right],$$  

(8)  

with

$$s = \text{sech}(z_s/\lambda), \quad t = \tanh(z_s/\lambda),$$  

(9a)  

$$\beta/\alpha = -e^{-2kL} [1 + \tanh(L/\lambda)/k\lambda]/[1 - \tanh(L/\lambda)/k\lambda].$$  

(9b)  

Equations (4) and (5) are evaluated at $z = z_s$, the position at which the electrons are no longer Landau resonant with the wave. A physical picture for this is given by Drake and Lee for the case of a magnetic guide field. Electrons traveling directly down the center of the sheet along the guide field can give energy to the wave since they see a constant electric field on a time scale of the wave growth rate $\tau = 1/\gamma$. If the electron is displaced to $z = z_s$, it moves along a sheared field line and sees an effective $k_1 = kz_s/l_s$. For large $z_s$ (and effective $k_1$) the electron will move in an oscillating wave field and give up no net energy. The resonance condition is that $k \parallel v_{\text{th}} < \gamma$ or

$$(z_s/\lambda) < \gamma/kv_{\text{th}} (B_{y0}/B_x)_{z_s} \leq L.$$  

(10)  

The evaluation of a wavelength for the mode is not straightforward. For a periodic chain of islands it would be either the distance from center to center of neighboring islands or from $X$ point to neighboring $X$ point. From structures in either the magnetic field [Figs. 6(b) and 8(b)] or the current [Fig. 9(b)], one sees that these distances are not equal; there are asymmetries in the data. Inspection of these and other structures suggests a suitable average $\lambda_{\text{ave}} = 20$ cm. Since $k = 2\pi/\lambda_{\text{ave}} = 0.31$ cm$^{-1}$, $\beta/\alpha \rightarrow 0$; this indicates that the
physical boundaries do not play a large role. Therefore, the growth rate becomes
\[
\gamma = \frac{v_{\text{th}} \delta^2}{2l_0 \lambda^2} \left[ \frac{\delta^2 - k^2 \lambda^2 (1 + t/k \lambda)}{1 + \tau/k \lambda} \right].
\] (11a)

Since in the guide field case the resonant interaction region is much smaller than the current half-width \( \lambda \), we can approximate \( s = 1, t = 0 \) to get the familiar expression
\[
\gamma = \left( v_{\text{th}} / 2l_0 \right) \left( \delta^2 / \lambda^2 \right) (1 - k^2 \lambda^2).
\] (11b)

This will be used in (6) to evaluate \( z_e \). The shear length is approximated closely by \( l_0 = \lambda B_{\text{eq}} / B_{\text{c}} (z = 5 \text{ cm}) = 4.34 \text{ cm} \), or can be evaluated directly from the data as 4.3 cm which gives \( \gamma' = 2.5 \times 10^6 \text{ sec}^{-1} \). This approximate growth rate is used in (6) to calculate \( (z_e/\lambda) < 0.13 \). (The actual half-width of the sheet containing the resonant electrons which produce the tearing, is less than 2 mm!) In turn, when the above ratio is used in (9a) a better estimate of the growth rate [Eq. 11(a)] yields \( \gamma = 1.7 \times 10^6 \text{ sec}^{-1} \) and a growth time \( \tau = 1/\gamma = 0.6 \mu\text{sec} \).

V. EXPERIMENTAL GROWTH RATES

The data, as illustrated in Figs. 7, 8, and 11, indicate that the tearing takes place in less than 4 \( \mu\text{sec} \) after the current is squeezed into a narrow sheet. However, tracking the change in the topology of the system is not good enough to generate a reliable number for the growth rate. To this end, the spatial data were Fourier analyzed to find the wavenumbers associated with tearing. The growth rate is then calculated from the temporal evolution of one of these modes. A two-dimensional fast Fourier transform of \( B_z \), taken 0.25 \( \mu\text{sec} \) after the tearing, shows a single peak with a wavelength of \( \lambda_x \approx 40 \text{ cm} \); this corresponds to the initial neutral sheet. At this early time the \( k \) spectra is invariant along the \( y \) axis. The wave number spectra of the current sheet in a plane \( (y = 110 \text{ cm}) \) is displayed in Fig. 13(a). To enhance the contribution of the tearing mode at \( t_e = 2.5 \mu\text{sec} \) the spectra at the pretearing time has been subtracted. The zero order peak is still visible \( (\lambda_x = 40, \lambda_y = \infty) \), however, since the magnitude of all the modes is time dependent. After the current constriction a fairly broad band of modes \( \lambda_x \) and \( \lambda_y \) has developed. In contrast to the theoretical picture, the islands that actually grow are neither perfectly symmetric nor does the neutral sheet extend to infinity along the \( x \) axis. It is therefore difficult to give a precise value for the mode wavelength. In the previous section we argued that \( \lambda_{\text{wave}} = 20 \text{ cm} \) was reasonable; we note a significant component in the spectra there. We follow the temporal evolution of the \( B_y \) component at \( \lambda_x = 20 \text{ cm}, \lambda_y = 8 \text{ cm} \) in Fig. 13(b), which is a semilogarithmic display and includes a best fit straight line in the region of growth \( (t_e + 1.5 \mu\text{sec} < t < t_e + 2.6 \mu\text{sec}) \). The curve gives an e-folding time of 0.97 \( \mu\text{sec} \) that compares favorably with the 0.6 \( \mu\text{sec} \) predicted by theory. Waves in the vicinity of the above mode have similar growth rates; but if we stray too far away \( (\lambda_x > 10 \text{ cm}, \lambda_y > 25 \text{ cm}) \) the temporal behavior may become erratic and even oscillate. This of course is not the complete picture, which is again three dimensional and shown in Fig. 14. At early times, i.e., before tearing, the modes are close to the \( k_x \) and \( k_y \) axis and reflect the un

![FIG. 13. Wavelength spectra of data and time history of mode e growth. (a) Wavenumber spectra for \( B_y \) at \( t_e = 2.5 \mu\text{sec} \) in a single x-z plane. The spectra at \( t_e = 0.25 \mu\text{sec} \), which is a single peak in the lower left-hand corner, has been subtracted. The tic marks have been labeled with the appropriate structure wavelength in cm. (b) The temporal evolution of the mode amplitude \( \lambda \) singled out by a black dot in (a). The curve is drawn as a semilog plot \[ \ln(A(t)) \] and a best line fit is superimposed in the region of growth.

![FIG. 14. Fully three-dimensional wavenumber spectra for the mode at (a) \( t_e = 0.25 \mu\text{sec} \) and (b) \( t_e = 2.5 \mu\text{sec} \). The spectra spills over into \( k_z \) space which reflects the three-dimensionality of the situation (Figs. 7 and 11). Here, \( 0.157 \text{ cm}^{-1} < k_x, k_y < 1.57 \text{ cm}^{-1} \), \( 0.628 \text{ cm}^{-1} < k_z < 1.57 \text{ cm}^{-1} \). The zero order peak at the origin has been suppressed for clarity.](image-url)
VI. DISCUSSION AND CONCLUSIONS

We have observed that a sufficiently "thin" current sheet in a high beta plasma with a guide field will tear into a series of "X" and "O" points. The axial current is largest in the center of the magnetic islands, which is what one would expect intuitively (although the opposite has been observed in a computer simulation\textsuperscript{10}). A very good estimate of the growth time has been derived from the simplest of linear tearing mode theories and agrees fairly well with experimental observations. One must add a caveat in that the experimental situation is far more complex than the theoretical model, which is two dimensional, neglects Hall currents, and considers a final state with a single mode. The observations show that in a given plane the topology of the magnetic field lines or axial current is not symmetric, and the $k$ spectrum is fairly broad. The agreement is therefore good in an average sense for a representative mode. That there is any agreement at all, considering the complexity of Figs. 7 and 11, is both remarkable and important for those doing research on the Earth's magnetotail where in situ measurements such as these are not presently possible. The spatial structure in the data indicates that there is either competition between tearing and magnetic merging, or waves locked into a three-dimensional eigenmode. In a separate series of experiments (to be reported later), a series of current channels without the added complexity of a neutral sheet were observed to merge. Based on these results we believe the first conjecture to be correct, although the final state is not as simple as the one depicted in Fig. 1(c). In all probability there is no final state if the currents remain strong. In this experiment the transformer cannot maintain the induced electric field and the axial current grows weak. At late times $t > t_a + 30\mu$sec the neutral sheet disappears, and only an X point, resulting from the last of the current in the plates bounding the column, remains.

1. The present experiment, which uses a fairly collisionless laboratory plasma of mean free path, $l_{dmp} \approx 2$ m ($R_m \approx 1$), does not approach the large Reynolds numbers seen in space or fusion plasmas ($R_m > 1000$). Furthermore, previous experiments have shown that the electron distribution in a neutral sheet is anisotropic; although such measurements were not repeated in this case, it is likely that non-Maxwellian distributions are the rule. The observed growth rate ($\gamma \approx 1 \times 10^6$ sec$^{-1}$) is of the same order as the collision frequency ($\nu_{el} \approx 1 \times 10^6$ sec$^{-1}$). The growth rate is in the collisional regime of Drake and Lee,\textsuperscript{6} which for our parameters is $\gamma_{coll} \approx 1.5 \times 10^6$ sec$^{-1}$. The collisional growth time $\tau_{coll} \approx 0.7\mu$sec is $\sim 20\%$ slower than what is observed.

We are not far enough into the collisional regime for kinetic effects to significantly affect our results. The experimental situation is midway between these regimes, and the mode structure is rich enough to prevent one from saying the kinetic approach is absolutely necessary here.

It is interesting to explore the origin of the observed current. A small Langmuir probe was scanned across the constrictor grid and $L-V$ traces swept before $(t = t_g - 3\mu$sec) and after $(t = t_g + 6\mu$sec) the slabbed grid was pulsed negative. In the first case no variation in the plasma potential across the slot was detected. When the grid is biased ($V_g = -200$ V) the plasma potential $V_p$ in the center of the current channel was observed to be $5.5$ V positive, with respect to that measured near the slot edge, which results in an electric field $E_z \approx 1$ V/cm and drift velocity $v_d \approx (E_z/B_0) \approx 5 \times 10^6$ cm/sec. This leads to currents $j \approx 0.8$ A/cm$^2$, comparable to the average observed Hall currents. These currents therefore originate in electrostatic fields set up across the neutral sheet. In previous experiments on neutral sheet dynamics,\textsuperscript{9} electrostatic fields of similar magnitude have also been reported.

The whistler wave speed is less than the electron thermal velocity and, as shown, the currents and fields change on the whistler time scale. It is interesting to speculate on the mechanism responsible for current propagation. Electrons drifting at the whistler speed could carry the axial current at these densities, but they would interact strongly with the waves. If electrons initially moved at their thermal speed they would Čerenkov radiate whistlers that, in turn, could trap them. The final current could be composed of electrons involved in a wave–particle interaction. A similar thing has been observed in an electrostatic computer simulation by Decyk et al.\textsuperscript{20} and will be investigated in future experiments.

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