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Magnetic Field Line Reconnection Experiments

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Abstract

A laboratory experiment concerned with the basic physics of magnetic field line reconnection will be discussed. Stimulated by important processes in space plasmas and anomalous transport in fusion plasmas, the work addresses the following topics: Dynamic magnetic fields in a high beta plasma, magnetic turbulence, plasma dynamics and energy transport. First, the formation of magnetic neutral sheets, tearing and island coalescence are shown. Nonstationary magnetic fluctuations are statistically evaluated displaying the correlation tensor $\mathbf{B} \mathbf{B}$ in the $\omega-k$ domain for mode identification. Then, the plasma properties are analyzed with particular emphasis on transport processes. Although the classical fluid flow across the separatrix can be observed, the fluctuation processes strongly modify the plasma dynamics. Direct measurements of the fluid force density and ion acceleration indicate the presence of an anomalous scattering process characterized by an effective scattering tensor $\mathbf{V} \mathbf{r}$. Turbulence also enhances the plasma resistivity $\eta$, by one to two orders of magnitude. Measurements of the three-dimensional electron distribution function $f_e(x, y, z, t)$ using a novel energy analyzer exhibit the formation of runaway electrons in the current sheet. Associated microinstabilities are observed. Finally, a macroscopic disruptive instability of the current sheet is observed. Excess magnetic field energy is converted to a double layer into particle kinetic energy and randomized through beam-plasma instabilities. These laboratory results will be compared with related observations in space and fusion plasmas.

1. Introduction

The reconnection of magnetic field lines at neutral points in plasmas is a problem of fundamental interest in magnetospheric physics [1-3]. In both solar flares and magnetic substorms reconnection is considered to explain the energization of particles at the expense of stored magnetic field [4, 5]. Many theories [6-15] and computer simulations [16-19] have been performed on this subject. Observations in space have recently produced convincing evidence for reconnection events [20-22]. However, information collected from a single spacecraft does not give a detailed picture of the three-dimensional phenomena. Thus, experiments to model reconnection in controlled laboratory plasmas have been undertaken [23-26].

Tearing instabilities and reconnection also play an important role in fusion research, in particular in tokamaks [27, 28], spheromaks [29], reverse field pinches [30], and mirrors [31]. Although the emphasis is on confinement rather than on particle acceleration, the basic MHD phenomena in space and fusion have much in common. Due to the high temperature of fusion plasmas, the diagnostics of magnetic fields is restricted to remote observations.

We have, therefore, designed a reconnection experiment in a laboratory plasma in a parameter regime which optimizes diagnostic access, repeatability, and parameter control [32-39]. Time varying antiparallel magnetic fields are applied to an initially uniform, almost collisionless, large plasma column. The crucial aspects which are investigated are the self-consistent magnetic field topologies, particle heating and acceleration, anomalous scattering processes associated with microinstabilities, and the energy transfer from magnetic fields to particles.

2. Experimental setup

The experiment is performed in a linear device in which a large (1 m dia., 2 m length) low pressure ($10^{-4}$ torr argon, hydrogen or helium) discharge is produced, with, perhaps, the largest oxide coated cathode developed so far, one meter in diameter (Fig. 1). The pulsed plasma $(\tau_{pe} \approx 10^{-2} \text{cm}^{-3}, kT_e \approx 10 kT_e = 5 \text{ to } 30 \text{ eV}, t_{rep} = 1 \text{ s})$ is immersed in a constant spatially uniform bias magnetic field $(\mathbf{B}_0 = 0 \text{ to } 100 \text{ G})$. The emission properties of the cathode are such that the plasma parameters are essentially constant for 80 cm across and 200 cm along $\mathbf{B}_0$. The temporal reproducibility from pulse to pulse is excellent ($\delta n/n \lesssim 5\%$).

The spatial length is then sufficient to reach steady state plasma conditions at which time the reconnection experiments are initiated. Note that we have adopted the magnetospheric coordinate system where $x, y$ are transverse and $z$ is the axial direction. For the reconnection experiment a time-dependent (damped sinusoidal) transverse magnetic field, $\mathbf{B} = (B_x, B_y)$, is established by pulsed axial currents $(I_p \lesssim 20 \text{ kA}, t_e \approx 80 \mu \text{s rise time})$ through two parallel aluminum plates (75 cm wide, 32 cm spacing, 200 cm length, insulated with glass) and located on top and bottom of the plasma column.

The magnetic field geometry in vacuum, sketched in Fig. 1(a), exhibits an X-type neutral point on the axis of the device. In the plasma, however, axial currents are induced $(I_p > 1000 \text{ A})$ which modify the magnetic field topology until a self-consistent current and field pattern is established. The induced plasma current $I_p$ flows opposite to the applied plate current $I_e$ and is largely carried by electrons. The direction of currents and electric fields $(\mathbf{E} = -\mathbf{A} - \mathbf{V}_p)$ are shown in Fig. 1(b) during the first phase of a the current pulse where $dI_p/dr > 0$. Electrons drift from the cathode toward the end anode where they are collected and the current closes via the metallic chamber wall. The cathode emits as many electrons as are collected at the end plate so as to maintain space charge neutrality. Since at the cathode sheath the electron current is limited to $j_e \leq (m_i/m_e)^{1/2} j_i$ where $j_i \approx (kT_e/m_i)^{1/2}$, the electron drift velocity is in general smaller than the thermal velocity, $v_d < (2kT_e/m_i)^{1/2}$. In response to a large induced electric field $(\mathbf{A} > n m_e v_e / e)$ a space charge electric field $-\mathbf{V}_p$ builds up opposite to $-\mathbf{A}$ so as to limit the net field and electron drift.

3. Measurement techniques

The plasma is diagnosed in situ with various probes offering time and space resolution. The data are processed digitally as shown in Fig. 2. The data acquisition system consists of a set of analog-to-digital converters (8 bit, 20 MHz) within which the waveforms are quantized in 1024 time steps. For example, the magnetic field obtained from small orthogonal magnetic loops is recorded vs. time at 400 spatial locations $(x, z)$ on repetitive pulses $(t_{rep} \leq 2 \text{s})$, averaged over 25-80 pulses per position to
obtain statistical information and the typical 20 million data points per run are analyzed with an on-line array processor. Correlation measurements can take up to 48 continuous hours over which several billion mathematical operations are performed. The reduced data are shipped over a direct link to large offsite computers on which the data is reorganized and displayed as vector maps, three dimensional contour maps and color movies. Basic plasma properties such as density, temperature, currents, and potentials are obtained from rapidly swept (\(t_{\text{sweep}} < 2 \mu s\)) differential Langmuir probes [34]. The probe characteristics are evaluated on-line. Probe measurements are compared with microwave interferometry and test wave diagnostics.

Similar to the magnetic field we measure the particle flows with orthogonal differential velocity analysers. Recording typically \(10^6\) traces of ion currents vs. time we can subsequently plot flow fields, density and temperature profiles. Subsequent data evaluation yields, for example, the directed ion energy density \((1/2m_iv_iz^2)\), the ion force density \((n_iz_in_i v^2/dt)\), and flow properties such as the vorticity \((\omega = \nabla \times v_\text{in})\). These data are taken under identical conditions as the electric and magnetic fields and electron temperature so that quantities such as the force on a fluid element \(\mathbf{F} = \mathbf{J} \times \mathbf{B} - \mathbf{V}(n_kT_e)\), and the resistivity \((\eta = E/J)\) may be evaluated; hence, a very complete picture of field and plasma properties emerges.

With a novel directional energy analyser [40] to be described in Section 6 we have measured the three-dimensional electron distribution function \(f_e(v)\). On the basis of this important data we are in a position to analyze, e.g., energy transport processes and velocity-space instabilities.

4. Average magnetic field topologies

During the rise of the applied magnetic field \((0 < t \leq 80 \mu s)\) plasma currents opposite to the applied currents are induced and slow down the flux transfer across the separatrix. Whereas in vacuum the separatrix intersects at right angles \((\nabla \times \mathbf{B} = 0)\) in the plasma it has two contact points joined by a common line along which \(B_1 \approx 0\).

We observe the self-consistent formation of neutral sheets with width-to-thickness ratios \(\Delta x/\Delta z \approx 20\) where the thickness \(\Delta z\), defined by the half width of the current sheet \((j = \nabla \times B_1/\mu_0)\) can be as narrow as \(\Delta z \approx 3 c/\omega_{pe} \approx 2 \text{cm}\). Such sheets have been observed to be stable for \(t \approx 80 \mu s\) which is compared with the Alfvén time from metal plate to current sheet \((t_A \approx 7 \mu s\) in He, \(20 \mu s\) in Ar).

Such a structure is shown in Fig. 3 as a vector plot, the arrow length is proportional to \(B_1\). Each arrow is centered on an experimental grid point and is averaged over an ensemble of 25 shots. The vector potential \(A = \int B_1 dx - \int B_1 dy\) is then evaluated at 1024 time steps so that a temporal history of the merging of magnetic field lines \((A_\phi = \text{const.})\) may be viewed. Several frames, each one half microsecond apart, are shown in Fig. 4 for argon. The neutral sheet is fully established within one Alfvén transit time. At \(t = 67 \mu s\) flux is still being forced vertically inward into the separatrix; field lines moving in this direction touch, reconnect and then move outward in the \(\pm z\) direction. The field dependence on four parameters, \(B(x, y, z, t)\),

![Fig. 1. Schematic view of the experimental device. (a) Cross-sectional view with transverse magnetic field line pattern in vacuum, (b) Side view with characteristic fields and currents. Note that the axis of the device is in the \(y\) direction while the transverse coordinates are \(x, z\), as is custom in magnetospheric physics.](image)

![Fig. 2. Block diagram of the plasma diagnostics performed with various probes and a digital data acquisition system. Fast mass data handling with computers allows one to perform time and space resolved vector measurements, statistical analysis and particle distribution function measurements.](image)

![Fig. 3. Vector field \(B_1(x, z)\) showing a neutral sheet magnetic field topology.](image)
has been displayed in computer-generated movies which show, for the first time, dynamic field line reconnection, tearing, and coalescence of magnetic islands in a real plasma, not in a simulation.

The stability of a current sheet with respect to tearing is a subject of great interest [41–46]. We have observed various degrees of tearing instabilities. In helium, for example, islands are seen to exist within the newly formed neutral sheet. This is displayed in Fig. 5(a) as a field of unit vectors (i.e., direction only). The structures are remarkable as the null regions in which these islands exist are also subject to large fluctuations in time and shot to shot. These islands coalesce within 10 \( \mu \)s. However, directional plots [Fig. 5(b)] indicate that a large island with two directional plots \([\text{Fig. 5(b)}]\) indicate that a large island with two

![Fig. 4. Magnetic field lines (contours of constant axial vector potential \( A_x \)) showing the process of magnetic field line reconnection. Note that two horizontal field lines near the center move in \( \pm z \) direction toward one another, merge at the full point, separate, and form two new vertically connected field lines which move rapidly apart in \( \pm x \) direction. The upper curve is at \( t = 57 \mu s \) with 0.5 \( \mu s \) interval between succeeding ones. \( x' = x + 28, \ z' = z + 12 \).](image)

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![Fig. 5. Unit vector field \( B_y/B_x \) showing the formation of two magnetic islands \([\text{Fig. 5(a)}]\) \( t = 7.2 \mu s \), which in time merge into a single island \([\text{Fig. 5(b)}]\) \( t = 17.2 \mu s \).](image)

\( X \) points at horizontal separatrix (or second order contact points) lie within it.

Major magnetic island formations are observed under two conditions: (a) For large axial magnetic fields \( (B_y > B_x) \) and strong plasma currents \( (I_p > 1500 \text{ A}) \) a large magnetic island occupying most of the grid space shown in Figs. 3–5 is formed instead of a neutral sheet. This configuration is closer to a Tokamak than the magnetotail. (b) In the second phase of the experiment the pulsed external currents become small \( (I_e \rightarrow 0) \) and the magnetic topology is determined chiefly by the plasma currents.

5. Magnetic fluctuations

Although the initial plasma parameters and applied currents are highly reproducible from shot to shot, we observe magnetic field fluctuations inside the plasma. Depending upon the time and spatial position the fluctuations can be comparable to the average value, hence cannot be neglected. Figure 6 shows a magnetic field hodogram, i.e., the trajectory of the vector tip in time. The probe is positioned within the neutral sheet where the magnetic noise is largest. The fluctuations within one shot become apparent as small loops in the vector plot [Fig. 6(b)]. These can be interpreted as a rapidly rotating magnetic field component associated with a wave. The random nature of these waves is apparent in Fig. 6(c) which shows hodograms from three separate shots.

The spatial properties of the fluctuations within one shot cannot be explored with a single probe. We can only obtain statistical properties such as probability distributions, moments, correlations from a large ensemble of shots taken at many positions. This requires the use of a digital data acquisition system. For example, we have recorded 25 shots of \( B_{90} \) at each of 300 positions and calculated on-line both the ensemble mean values of the field magnitude \( (B) \) and direction \( (\theta) \) and the root-mean-square values \( B_{\text{rms}}, \ \theta_{\text{rms}} \). From this extensive data base one concludes that the fluctuations maximize near the neutral sheet region at times of large plasma currents \( (25 \leq t \leq 50 \mu s) \).
In general, both the magnetic field direction and angle exhibit fluctuations with related space-time variations. Since the magnetic turbulence ultimately affects the particle motion (as will be seen in the following section) and possibly heat transport, detailed studies of it have been undertaken. Assuming the noise is comprised of a large number of randomly generated waves, a correlation analysis has been performed using the following procedure: Two identical probe sets measure the vector components of $B_i$. During each reconnection event the signals are digitized by four 20 MHz analog-to-digital converters then fast Fourier transformed and digitally filtered by the on-line array processor (MAP200). An inverse FFT is done to generate the time history of the filtered signals in the frequency band of interest. Eighty consecutive shots are stored in the fast MAP memory and the following operations are performed on the signals filtered at frequency $\omega$:

$$B_{i,\omega,n}(t) = B^{\text{tot}}_{i,\omega,n}(t) - \langle B_{i,\omega}(t) \rangle$$

(1)

where $i = x, z$ component, $B^{\text{tot}}_{i,\omega,n}$ is the total field measured during the $n$th shot and $\langle B_i \rangle = \sum_{n=1}^{80} B^{\text{tot}}_{i,\omega,n}$ is the average field at time $t$. The ensemble number $N = 80$ is sufficiently large for a convergence of the mean values. Analysis of the amplitude probability distribution function using the chi-square test shows a normal distribution to a high degree of confidence [47].

The cross-spectral function (CSF) may be normalized in the form

$$\rho_\omega(t, \tau, \omega) = \frac{(B_{ix}B_{ix2})(\Delta r, t, \tau, \omega)}{(B_{ix,\omega}(t, \tau))^2(B_{ix2,\omega}(\Delta r, t + \tau))^2}$$

(3)

so that it lies between $\rho = \pm 1$. Since the two terms in the denominator are separately stored it is also possible to evaluate the unnormalized correlation function, eq. (2).

Since there are four vector components the correlation is a tensor

$$\rho = \begin{pmatrix} \rho_{x1,x2} & \rho_{x1,x3} \\ \rho_{x1,x2} & \rho_{x1,x3} \end{pmatrix}$$

(4)

The off-diagonal elements are large if over the displacement $(\Delta r, \tau)$ an $x$ component of the field rotates into the $z$ direction or vice versa. The total number of on-line calculations to get $\rho_\omega$ at 400 spatial grid positions, 64 delay times $\tau$, at each of eight starting times $t$ is staggering ($N > 2 \times 10^8$). Nevertheless, we have measured $\rho_\omega$ at several frequencies and in two orthogonal planes.

Figure 7 shows the cross-spectral intensity $\langle B_{ix}B_{ix2} \rangle$ at $f = 1$ MHz in the horizontal $x$-$y$ plane. The fixed reference probe $2$ is located at $(x', y', z') = (34 \text{ cm, 14 cm, 0})$ noting that the axis of the device goes through the center of the plane of Fig. 7. The cathode in this coordinate system is at $y = 150 \text{ cm}$. The three-dimensional display in the upper part and its two-dimensional projection as a contour map are taken at $t = 36.4 \mu$s after the start of the reconnection pulse with signal $B_{ix2}$ delayed by $\tau = 1.3 \mu$s with respect to $B_{ix1}$. The observations show a wave-like pattern in the $x'$-$z'$ plane. Variations in the delay time indicate no preferred direction of wave propagation but an interference of standing and propagating modes.

From this diagram alone identification of the modes cannot be made. The analysis proceeds with a two-dimensional Fourier
transformation of the spatial distribution so as to obtain a wave number spectrum \( \langle B_{k_1}B_{k_2}\rangle(k_x,k_y,\tau) \). As shown in Fig. 8 the magnetic turbulence exhibits a broad range of \( k \) values even at a single frequency \( \omega \). Note that each box in \( k_x,k_y \) corresponds to wavelengths \( \lambda_x = 50\,\text{cm}, \lambda_y = 22\,\text{cm} \) and, of course, larger \( k \)'s to smaller wavelengths. The probes can resolve structure of order 1 cm.

The data is still too complex; however, a further simplification clarifies it immensely. The cross spectral functions are filtered in \( k \) space and the dominant spatial mode in Fig. 8 is then individually examined after an inverse transformation. This is displayed in Fig. 9 for the cross spectral function \( \langle B_{k_1}B_{k_2}\rangle \). In Fig. 9(a) only one of the dominant modes (\( \lambda_x = 50\,\text{cm}, \lambda_y = 5\,\text{cm} \)) from the spectrum of Fig. 8 is retained; the function is sinusoidal with phase fronts nearly parallel to \( B_{k_0} \). As will be discussed below, the phase fronts may be easily tracked as a function of \( \tau \) and the propagation studied. If a second mode is included [Fig. 9(b)] the interference becomes obvious although the contour plots projected below are symmetric. For four modes [Fig. 9(c)] the pattern begins to approach that of Fig. 7 and tracking the motion of peaks as well as comparison with other tensor components of the cross spectral function becomes nearly impossible.

The wave dispersion may therefore be unravelled by examining the cross spectral function at several frequencies in orthogonal planes one mode at a time. We have initiated this by investigating the orthogonal tensor components \( \langle B_{k_1}B_{k_2}\rangle \) and \( \langle B_{k_3}B_{k_4}\rangle \) (\( k_{x_2} \) is measured by the spatially fixed reference probe).

It is found that the polarization of the \( B_1 \) vector is circular in the direction of the electron cyclotron motion around the dominant axial field component \( B_{k_0} \) in the neutral sheet.

As was shown in Fig. 8, the magnetic turbulence exhibits a broad range of \( k \)-values even at a single frequency \( \omega \). On the basis of the frequency regime (\( \omega_{\text{LH}} < \omega < \omega_{\text{pe}} \), where \( \omega_{\text{LH}}/2\pi \approx 0.1\,\text{MHz}, \omega_{\text{pe}}/2\pi \approx 28\,\text{MHz} \)) the oblique whistler modes have the properties of exhibiting a broad range of wave-
The plasma parameters have been measured with the use of rapidly swept probes inserted into the plasma without causing noticeable perturbations. At seven separate times within the applied magnetic field pulse the plasma potential \( \phi_p \), electron density \( n_e \), and temperature \( T_e \) are determined on-line from computer evaluation of ten shot ensemble averages. This is done at several hundred positions in the transverse \( x-z \) plane typically involving the evaluation of about 20000 Langmuir probe traces in one data run.

While prior to the magnetic field pulse the plasma parameters are highly uniform, during the reconnection event strong gradients are generated. Initially (\( t \leq 25 \mu s \)), the electron density rises rapidly due to ionization of residual neutrals but subsequently it levels off in spite of rising electron temperatures indicating neutral gas burn-out. Thus, during most of the experiment, density gradients are caused by fluid drifts. We observe a density minimum at the neutral sheet and maxima \((n_e \approx 2 \times 10^{13} \text{cm}^{-3})\) near the two contact points of the separatrix. There is also an axial density gradient toward the cathode, i.e., in the direction of the ion drift along the induced electric field.

The electron temperature increases strongly but nonuniformly in the reconnection process, ranging from \( \approx 5 \text{ eV} \) in the center of the neutral sheet to \( \approx 15 \text{ eV} \) in argon (30 eV in helium) near the two contact points of the separatrix. Axial temperature gradients develop in the direction of the electron drift but are smaller \((\Delta T_e < 2 \text{ eV})\) than the transverse ones on the average.

Temperature and density data are appropriately summarized by forming the plasma pressure \( p = n k T \), displayed in Fig. 11(a). The ion contribution is neglected since we observe \( T_i \gg T_e \). Two pressure islands near the edges of the neutral sheet are obvious, and there is a pressure minimum in the center of the neutral sheet. In Fig. 11(b) we also show the magnetic pressure \( B_i^2/2\mu_0 \). The forces on the fluid are not in balance leading to an acceleration of the plasma. We have both measured the flow field and the force field directly in order to check the validity of describing the turbulent plasma by the classical fluid equations. First, we discuss the flow measurements. From three orthogonal differential ion collectors the normalized ion drift velocity vector \( v_i/c_s \) is obtained and recorded digitally vs. time and position analogous to the magnetic field [see Fig. 3]. Subsequently, at a given time, a flow velocity field is drawn or streamlines calculated for the predominantly rotational flow. Since the time variation is quantized in \( \approx 10^3 \) small increments \((\Delta t \approx 200 \text{ ns})\) a nearly continuous movie display of the fluid motion is generated.

Figure 12 shows the ion flow field at two characteristic time

\[ T = 2140 \mu \text{sec} \]
\[ \tau = 130 \mu \text{sec} \]

\( \text{Fig. 10. Two-dimensional interferometer traces of the } B_x \text{-components of a magnetic wave at } f = 1 \text{ MHz. The wave is excited at } x = 24 \text{ cm, } y \approx z \approx 8 \text{ cm from a magnetic loop antenna with } B = (0, 0, B_z). \text{ Propagation in the transverse } x-z \text{ plane; } B_{y0} = 10 \text{ G.} \)

\[ \text{Fig. 11. (a) Plasma pressure } p = n k T. \text{ (b) Magnetic pressure } B_i^2/2\mu_0. \text{ Both total pressures are comparable } (p \approx 1) \text{ but not in equilibrium resulting in an acceleration of the plasma.} \]
steps: (a) At early times when the plasma is highly turbulent, and (b) when the turbulence level has decreased although neutral sheet topology still exists. We find the classical flow due to the $J \times B$ force in a neutral sheet geometry well established in Fig. 12(b), but this occurs only after the large pressure gradients [see Fig. 11(a)] and fluctuations have decreased. The maximum outflow velocity, normalized to the local Alfvén speed, is $v_{{\text{max}}} \approx 0.5 v_A$, the inflow velocity is $v_i/v_A \approx 0.2$. The latter is a measure for the rate of steady-state reconnection [49]. At early times ($t \leq 40 \mu s$) the flow field is turbulent in space and time.

We now investigate whether the observed velocities can be explained on the basis of the momentum equation $\rho \text{v} \cdot \text{d} = J \times B - \nabla P$. Magnetic fields, currents and the plasma pressure have been directly measured and the fluid momentum change is readily obtained from the time derivative of the flow field data and local density values. We find that at early times the force density $\mathbf{f} = J \times B - \nabla P$ is much larger than the fluid momentum change which indicates that an additional drag or scattering process modifies the free fluid acceleration.

Assuming that the additional force, $\Delta \mathbf{f} = f - mn(\text{v}/\partial t)$ can be modeled by a scattering term $\Delta \mathbf{f} = mnv^*v$ we can quantitatively evaluate the scattering rate $v^*$. Since force and velocity are, in general, not in parallel, the rate coefficient is a tensor quantity, $v^*$. Physically, such scattering properties can be expected from wave-particle interactions with anisotropic turbulence. The scattering is, of course, not due to ion-neutral collisions which are essentially negligible ($v_{{\text{in}}}^{-1} > 1 \text{ ms}$). Likewise, classical viscosity $\eta \nabla^2 \text{v}$ cannot account for $\Delta \mathbf{f}$. Although it is possible to model $\Delta \mathbf{f}$ by anomalous viscosity coefficient $\eta^*$ we prefer the simpler scattering model since the details of the anomalous processes are not well known. In our model, the fluctuations have to absorb a large momentum change which appears only possible if they are coupled to the axial bias magnetic field, hence are hydromagnetic fluctuations as described earlier. It is also interesting to note that although the ions are unmagnetized they are subject to magnetic forces via the space charge coupling to the magnetized electrons.

Figure 13 shows the spatial distribution of both diagonal ($v_{{\text{ex}}} = v_{{\text{ex}}}$) and off-diagonal ($v_{{\text{ex}}} = -v_{{\text{ex}}}$) scattering tensor components, normalized to the local ion plasma frequency $f_{{\text{pi}}}$. The space–time variation shows that significant scattering is found in the neutral sheet region at early times which correlates with location and time of large plasma currents and electromagnetic fluctuations. At late times ($t \approx 80 \mu s$) when the classical flow pattern is observed [Fig. 12(b)] the anomalous scattering effects are negligible. These results show that the fluid description does not adequately describe the plasma dynamics in a turbulent fluid unless augmented by anomalous terms.

A clear identification of which modes in the turbulence spectrum is responsible for the scattering process cannot yet be made. We have extended the fluctuation studies in frequency and observed a wide range of electrostatic modes. Using two electrostatic probes we have measured with analog correlation techniques, the frequency and wavenumber spectrum and summarized the results in Fig. 14. In contrast to the magnetic whistler mode turbulence the electrostatic noise consists of ion sound modes.

Due to the short scale lengths of the sound modes ($\lambda \rightarrow \lambda_D$ as $\omega \rightarrow \omega_{pe}$) significant scattering of electrons can be expected which manifests itself in an enhanced resistivity. In spite of experimental difficulties due to error propagation we have attempted to measure the plasma resistivity directly.

In the framework of fluid theory the relation between currents and electric fields for slow processes ($\partial / \partial t \ll v$, $\omega_{pe}$) and scalar resistivity is given by the generalized Ohm's law [50]

$$\eta \mathbf{J} + (\mathbf{J} \times \mathbf{B})/ne = \mathbf{E} + v \times \mathbf{B} + (1/ne)\nabla P$$

(5)

Ohmic and Hall currents are driven by pressure gradients, Hall fields and electric fields which are both rotational (induced)
and divergent (space charge) fields. We have measured under identical conditions all fields, currents and plasma parameters and calculated from eq. (5) the resistivity $\eta$ and normalized it to the classical Spitzer resistivity $\eta_s$ [51]. We note that the current density obtained from $J = (1/\mu_0)\nabla \times B$ has both axial and transverse components like the magnetic field. The axial electric field has both inductive and electrostatic contributions of comparable magnitude but opposite sign which are measured simultaneously with a differential dipole-type Langmuir probe [52]. The transverse electric field is dominantly electrostatic and largest in magnitude while the Hall field $v \times B$ is nearly negligible. The results show that the resistivity is spatially highly nonuniform and, when averaged over the $x-z$ plane is given by $\eta/\eta_s \approx 25$. This significant enhancement may be attributed to the fluctuation spectrum in the neutral sheet.

Having identified the basic properties of the fluctuations and their effects on the macroscopic plasma properties one would now like to identify the cause for these microinstabilities and their direct effects on the particle dynamics. The knowledge of the particle distribution function is essential for the study of wave-particle interactions. We have developed the proper tools for making three-dimensional velocity distribution function measurements [40].

Figure 15 shows the measurement principle. A novel directional electron velocity analyzer [Fig. 15(a)] is used which provides for angular resolution by means of a microchannel plate. This plate contains a multitude of long thin parallel holes in a small metal plate admitting electrons only from a narrow cone in velocity space. The analyzer is mounted on a probe shaft and can be scanned along the two polar angles $\theta, \phi$ so as to cover essentially the entire surface in velocity space [Fig. 15(b)].

The electron distribution function $f_e(v_x, v_y)$ observed in the neutral sheet at early times ($t \approx 20 \mu s$) is shown in Fig. 16. A tail of energetic electrons drifting along the neutral line ($B_{\parallel 0}$) in the direction of the current flow is clearly visible. The distribution is marginally stable. Although the density of runaway electrons is small, their importance increases progressively for the higher moments of the distribution, i.e., they carry a significant fraction of the plasma current, determine greatly the mean energy of the electrons, and dominate the heat transport.

As regards velocity space instabilities, a variety of modes can be resonantly driven unstable. These include electrostatic plasma waves ($\omega \approx \omega_{pe}$), oblique electron cyclotron harmonic modes ($\omega > \omega_{ce}$), oblique whistler modes ($\omega < \omega_{ce}$) and oblique ion sound waves ($\omega < \omega_{pi}$). Slow low-frequency modes can also be destabilized by the drift of the main electron population ($v_{\parallel 0}/v_T \approx 0.1$). Subsequent wave–wave interactions can lead to other modes nonresonant with the beam. For example, electromagnetic radiation near the electron plasma frequency ($\omega \approx \omega_{pe}$ $\approx 2 \pi \times 12$ GHz) is observed during the reconnection process. This mode can be generated by the decay of Langmuir waves into sound waves as demonstrated in a related experiment [53]. Thus, the current sheet is a rich source of plasma turbulence. Further analysis of the importance of the various modes in determining the macroscopic plasma properties is required.

Velocity Distribution $F(v)$

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7. Current sheet disruptions and double layers

In the previous sections we have described magnetic field topologies and plasma properties in terms of their average and random behavior. We now turn to events which explosively change the field and plasma properties. These are analogous to magnetic substorms or solar flares which disrupt a quasi steady-state neutral sheet topology on a short time scale [54]. These macroscopic instabilities require the consideration of the total current system with boundary conditions and not only the local properties of the null region. In addition, plasma phenomena such as potential double layers may arise which are outside the classical MHD theories for magnetic field line reconnection [55]. For example, a magnetic substorm is characterized by the partial disruption of the magnetospheric crosstail current. The current flow is not terminated but rather diverted along the magnetic field lines into the polar regions where auroral potential structures build up [56]. At these localized potential drops particles are energized, i.e., the magnetic field energy converted into kinetic particle energy. The particle beams dissipate their energy in the lower ionosphere.

In our laboratory experiment we observe events closely related to the physics of magnetic substorms even though the boundary conditions are not the same [39]. We are interested in the stability of the current sheet when the current density in the central region is raised above its “normal” value, defined as the value corresponding to the formation of the stable current sheet (Fig. 3). This approach is motivated by the events thought to lead to a substorm, i.e., a continuous build-up of the crosstail current due to steady reconnection at the magnetopause [57]. Experimentally, this is accomplished as shown in Fig. 17. The neutral sheet current is terminated on a grounded metallic end anode (32 cm x 75 cm) whose central region (6 cm x 13 cm) is separated and connected to an external d.c. power supply. When the supply voltage is increased \( \left( V_{\text{d.c.}} > 0 \right) \) the current to the center plate \( I_a \) rises until a critical value is reached at which it disrupts spontaneously. The associated processes inside of the plasma are studied \textit{in situ} with probes and particle detectors.

Figure 18 shows the induced center plate current \( I_a(t) \) at different applied d.c. voltages \( V_{\text{d.c.}} \). At low currents a smooth sinusoidal waveform is observed whose shape is determined by the externally applied primary current. When the current is raised to \( I_a \geq 200 \text{A} \) spontaneous sharp current drops develop which, for \( V_{\text{d.c.}} \approx 15 \text{V} \), can lead to a complete current loss. It is found that the disrupted center plate current flows to the large surrounding grounded end anode. Thus, the total induced plasma current is constant. The disruptive instability is localized and involves a redirection of the current flow. It is an instability of the current sheet. If we start with a different topology, e.g., a magnetic island, produced by collecting the total plasma current at the small end plate \( I_a \approx 600 \text{A} \) no disruptions are observed.

The current path through plasma and conductors may be represented by an electrical circuit with inductance \( L \). Due to the current disruptions an inductive voltage \( L \frac{dI}{dt} \) arises which drops off at the location of the current disruption. When we observe the instantaneous plate voltage to ground we find at every current disruption an inductive voltage spike well in excess of the applied d.c. potential \( V_{\text{d.c.}} \approx 10 \text{V} \), \( L \frac{dI}{dt} = 50 \ldots 100 \text{V} \).

With Langmuir probes the local instantaneous plasma potential has been measured in order to determine how the inductive voltage drops off inside the plasma. As shown in Fig. 19 the high positive plasma potential near the anode \( \Phi_p \approx V_a \) decreases abruptly well inside the plasma, i.e., not at a sheath...
and not uniformly distributed. Thus, a potential double layer [58] has been formed \((\Delta \phi \approx |d| f/d \approx 35 \text{ V}; \ d \approx 5 \text{ mm} \approx 100 \lambda_d)\). Its transverse \((x, z)\) dimensions are approximately those of the small anode. It exists only during current disruptions. On the high potential side \((\Delta y < 6 \text{ cm})\) the electron density is greatly \((\sim 50\%)\) reduced. The electron distribution function measured with the directional velocity analyzer described in Section 6 exhibits a beam of free electrons accelerated at the double layer. Beam-plasma instabilities generate bursts of electron plasma waves with \(\omega \approx \omega_{pe} \approx 2 \pi \times 6 \text{ GHz}\). On the low potential side ion energy analyzers detect energetic \((\sim 40 \text{ eV})\) ions streaming away from the plate region into the ambient plasma. Low frequency electrostatic and magnetic turbulence is generated in the surrounding plasma.

The disruptive instability is believed to arise from a density loss in the center of the current sheet, possibly triggered by local plasma heating. The resultant current decrease causes an inductive emf which raises the plasma potential near the end plate and expels ions from the current channel. The feedback between rising potential and decreasing charged particle density leads to the rapid current drop. Counterstreaming ions and electrons enable the formation of the potential double layer. The double layer plays an important role in the energy transfer mechanism. Stored magnetic field energy is converted at the double layer into particle kinetic energy. Microinstabilities lead to the thermalization of the particle beams, to plasma heating and electromagnetic radiation [53] analogous to type III radio bursts [59].

S. Conclusions

The interaction between plasmas and magnetic fields in a reconnecting geometry is a complicated problem in nonlinear physics. In the present laboratory experiment we are attempting to understand the major processes on the basis of careful observations. These have progressed from the average macroscopic current-voltage relations to space and time resolved measurements of fields and plasma parameters. Now a statistical analysis of the fluctuation phenomena is in progress which will lead to a characterization of the various microinstabilities in terms of eigenmodes. In parallel, the plasma diagnostics has been refined to measure directly the particle distribution function \(f(\nu, r, r)\) from which the most basic plasma properties can be derived. The next step will be to correlate the field and particle studies which form a coupled system.

From the data presented it has become obvious that turbulence and particle studies require the measurement of fluctuations depending on many variables. With conventional analog measurement techniques this task would be beyond reasonable limits of time and effort. But with the evolution of high speed computers, one can now investigate a vast parameter space in a finite time and, what is equally important, analyze and condense the information since our imagination of multivariable functions is limited. While computers have been successfully used in solving large numerical tasks to aid theoretical work [60] we have now adapted them to experiments for mass data analysis and reduction. Our results are not simulations but real experimental observations. This development in experimental physics has vastly expanded the capabilities in basic research. It is essential for unraveling the complicated processes taking place in magnetic field line reconnection.

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