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Spheromaks, solar prominences, and Alfvén instability of current sheets

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Three related efforts underway at Caltech are discussed: experimental studies of spheromak formation, experimental simulation of solar prominences, and Alfvén wave instability of current sheets. Spheromak formation has been studied by using a coaxial magnetized plasma gun to inject helicity-bearing plasma into a very large vacuum chamber. The spheromak is formed without a flux conserver and internal λ profiles have been measured. Spheromak-based technology has been used to make laboratory plasmas having the topology and dynamics of solar prominences. The physics of these structures is closely related to spheromaks (low β, force-free, relaxed state equilibrium) but the boundary conditions and symmetry are different. Like spheromaks, the equilibrium involves a balance between hoop forces, pinch forces, and magnetic tension. It is shown theoretically that if a current sheet becomes sufficiently thin (of the order of the ion skin depth or smaller), it becomes kinetically unstable with respect to the emission of Alfvén waves and it is proposed that this wave emission is an important aspect of the dynamics of collisionless reconnection.

1. Introduction

1.1 Force-free equilibria

Three topics related by common underlying physics will be discussed: experimental studies of spheromak formation, experimental simulations of solar prominences, and theoretical analysis of Alfvén instability of current sheets.

Common to these three topics is the low β nature of the plasma equilibrium. Because of the low β, the MHD equation of motion

$$\rho \frac{d\mathbf{U}}{dt} = \mathbf{J} \times \mathbf{B} - \nabla P \tag{1}$$

in equilibrium reduces to simply $\mathbf{J} \times \mathbf{B} = \mathbf{0}$ which implies that $\mu_0 \mathbf{J} = \lambda \mathbf{B}$ where λ is a so-far undetermined scalar function of position. This gives the fundamental force-free relationship (Lundquist, 1950)

$$\nabla \times \mathbf{B} = \lambda \mathbf{B} \tag{2}$$

a simple equation having an infinity of solutions depending on the boundary conditions. If λ is spatially uniform then the curl of Eq. (2) gives a vector Helmholtz equation,

$$\nabla^2 \mathbf{B} + \lambda^2 \mathbf{B} = 0. \tag{3}$$

1.2 Properties of λ

The divergence of Eq. (2) gives

$$\mathbf{B} \cdot \nabla \lambda = 0 \tag{4}$$

which means that λ is constant on a field line but may vary across field lines (i.e., λ may be thought of as a function of the parameter which labels individual field lines). If there are flux surfaces, then λ is a surface quantity (i.e., λ is constant on a flux surface, but may vary across flux surfaces).

Equation (2) has important limiting cases determined by boundary conditions and also by λ profiles. Boundary conditions may have symmetries or lack thereof and also may have open field lines or not. λ may be spatially uniform (i.e., the Taylor (Taylor, 1974) state) or in the more general case may be non-uniform. Helicity flows from high λ to low λ (Jarboe *et al.*, 1985; Fernandez *et al.*, 1989) and so the uniform λ situation corresponds to having no helicity flow. This is analogous to thermodynamics where temperature plays the role of λ and heat flux plays the role of helicity. The Taylor state (Taylor, 1974) (uniform λ state) corresponds to an isolated system in thermal equilibrium (i.e., a state with uniform temperature).

A spheromak is a magnetohydrodynamic (MHD) plasma equilibrium which, like a tokamak, has nested toroidal flux surfaces formed by helical magnetic fields. However, unlike a tokamak a spheromak has zero toroidal field at its surface. This means that no external coils link the spheromak and so allows the spheromak to be simply connected (spheroid-like). This contrasts with a tokamak which is doubly connected (toroid-like) and which has external coils linking the toroid and producing toroidal fields which are non-zero at the surface of the toroid. Spheromaks are of interest in fusion research because, not having a "hole in the doughnut", spheromaks provide a much simpler method for achieving the nested toroidal flux surfaces required for plasma confinement. Traditional laboratory spheromaks (Furth, 1983; Jarboe, 1994; Bellan, 2000) correspond to axisymmetric solutions of Eq. (2) over a finite volume.

Magnetic helicity is conserved on the relaxation (reconnection) time scale, but eventually becomes dissipated on the much slower resistive time scale. Thus, if a spheromak is to be sustained on the resistive time scale, it is necessary to inject helicity to make up for the slow resistive losses. It is also necessary to inject helicity when creating a spheromak. Helicity injection in spheromaks is typically accomplished using a magnetized plasma gun source designed to impose a finite electrostatic potential difference across the ends of an open flux tube that intercepts the bounding surface.

In general, the bounding surface is a flux conserving wall which imposes the condition that the normal magnetic field vanishes except at the helicity injection source (plasma gun). In contrast, solar prominences (Tandberg-Hanssen, 1995) correspond to solutions of Eq. (2) with non-axisymmetric boundary conditions and, instead of having a finite volume, one boundary is a planar surface and the other boundaries are at infinity. Boundary conditions are imposed on the planar surface. λ is usually non-uniform since it is obviously energetically impossible to fill up a semi-infinite space with uniform λ.

Although MHD analysis shows that helicity flows from regions of large λ to regions of small λ (Jarboe *et al.*, 1985), MHD says very little about the actual mechanism(s) causing the flow. In general large λ gradients tend to drive MHD instabilities which have the effect of reducing the λ gradient. Helicity is an extensive quantity and λ is the associated intensive quantity. Helicity flow is a natural consequence of helicity conservation. A conserved extensive quantity (e.g., helicity, mass, energy) is fungible and may be transported from one place to another. Helicity is conserved on the reconnection time scale (i.e., reconnection does not significantly dissipate helicity) but helicity does decay on the much longer resistive time scale (Bellan, 2000).

λ has several related interpretations:

1. λ is an eigenvalue of Eq. (3) and has the dimensions of an inverse length (Jarboe *et al.*, 1985), i.e., $\lambda \sim L^{-1}$ where L is a characteristic length of the system.

2. λ is related to the twist of the field lines in a flux tube. For the particular example of a long cylindrical flux tube, the solution of Eq. (2) is $B_z(r) = \bar{B}J_0(\lambda r)$, $B_\theta(r) = \bar{B}J_1(\lambda r)$ and field line trajectories are given by the relation $rd\theta/B_\theta = dz/B_z$. Thus, for small r the field line twist is $d\theta/dz = B_\theta/rB_z \simeq \lambda/2$.

3. If one integrates Eq. (2) over the cross-section of a flux tube, it is seen that $\lambda = \left(\int \nabla\times\mathbf{B}\cdot d\mathbf{s}\right)/\left(\int \mathbf{B}\cdot d\mathbf{s}\right)$ or $\lambda = \mu_0 I/\Phi$. Thus λ is just the axial current per axial flux of a flux tube.

4. In an isolated system, Eq. (2) can be integrated to give $\mathbf{B} = \lambda\mathbf{A} + \nabla f$ where f is an arbitrary scalar function. The magnetic energy $W = \int (B^2/2\mu_0)d^3r$ for an isolated system with no open field lines can thus be written as

$$
\begin{aligned}
2\mu_0 W &= \int \mathbf{B}\cdot\nabla\times\mathbf{A}d^3r \\
&= \int \left(\mathbf{A}\cdot\nabla\times\mathbf{B} + \nabla\cdot(\mathbf{A}\times(\lambda\mathbf{A}+\nabla f))\right)d^3r \\
&= \int \left(\lambda\mathbf{A}\cdot\mathbf{B} + \nabla\cdot(f\mathbf{B})\right)d^3r \\
&= \lambda K \qquad (5)
\end{aligned}
$$

where $K = \int \mathbf{A}\cdot\mathbf{B}d^3r$ is the magnetic helicity (Bellan, 2000; Turner *et al.*, 1983) of the system. Energy dissipating processes that conserve helicity (e.g., magnetic reconnection) will cause an isolated system to relax to a minimum energy state with the lowest λ allowed by the boundary conditions (Taylor, 1974). Helicity flows from regions of high λ to low λ because such a flow will reduce the energy of the system since $W = \lambda K/2\mu_0$. Because λ scales as an inverse length, helicity flow will involve geometric expansion. Because λ is proportional to energy per helicity and equivalently to I/Φ (also equivalently, λ is proportional to approximate inverse length and to approximate field line twist), a system can reduce its magnetic energy in a helicity-conserving fashion by having helicity flow from small, highly twisted volumes into accessible larger volumes which have less twisted field lines and smaller I/Φ.

5. When open flux tubes intercept biased electrodes, helicity can be injected into the volume of the open flux tube. This is the basis of the magnetized plasma gun. In particular if a power supply is connected to the electrodes, the power supply can drive currents through the open flux tube and so impose a $\lambda_{gun} = \mu_0 I/\Phi$ where I is the driven current in the open flux tube and Φ is the flux intercepting the electrode. The rate of helicity injection into the open flux tube is $2V\Phi$ where V is the electrostatic potential of the electrode intercepting the flux tube (Bellan, 2000; Jensen and Chu, 1984).

2. Spheromak Formation Experiment

In this experiment (Yee and Bellan, 2000) a coaxial magnetized plasma gun is mounted on a vacuum chamber that is much larger than the gun dimensions. There is no flux conserver so the plasma gun is effectively surrounded by empty space. Four distinct operational regimes were observed (Yee and Bellan, 2000) and these regimes were distinguished from each other by the value of λ_{gun} compared to the nominal linear dimension L of the gun (the radial dimension).

The behavior is determined by a competition between two different kinds of magnetic force: (i) the magnetic pressure force, associated with the gun current and proportional to I^2, which tends to push plasma out of the gun and (ii) the magnetic tension force, associated with the gun flux and proportional to Φ^2, which tries to restrain the plasma from exiting the gun.

Operation at different values of λ_{gun} have revealed four distinct experimental regimes:

1. For $\lambda_{gun}L \gg 1$, the magnetic tension is completely inadequate to restrain the magnetic pressure and so plasma explodes out of the gun without restraint (Regime I).

2. For $\lambda_{gun}L$ slightly larger than unity the magnetic pressure is marginally able to push the plasma out and it is found via magnetic probes that closed flux surfaces form, i.e., a spheromak configuration separated from the gun by an x-point; this is called regime II. A very

surprising feature of Regime II is that spheromaks are formed even though there is no flux-conserving wall; the spheromaks move away from the gun with constant velocity and expand self-similarly. Measurements of internal λ show that λR remains constant where R is the major radius. It is believed that inertial effects play the role of an effective wall.

3. For $\lambda_{gun} L$ slightly smaller than unity, the magnetic pressure is marginally unable to push the plasma out, and a twisted bulbous structure is observed which protrudes out from the gun, but does not quite detach; this is called Regime III.

4. In the case where $\lambda_{gun} L \ll 1$, the magnetic tension is so strong that the plasma cannot leave the gun (Regime IV).

Measurements of internal magnetic fields (Yee and Bellan, 2000) show that there is a significant difference between the internal $\lambda = \mu_0 J_\phi / B_\phi$ profile in Regimes II and III. The internal λ is peaked on the magnetic axis of the spheromak in Regime II, indicating that the spheromak is detached from the gun (not being driven by the gun). On the other hand, in Regime III (where the plasma has not quite detached to form a spheromak) the internal λ decreases monotonically from the gun into the plasma indicating that helicity is flowing from the gun into the plasma. These two contrasting observations of λ profiles are in accord with the concept that helicity flows from high λ to low λ.

3. Solar Prominence Simulation Experiment

In the solar prominence experiment (Bellan and Hansen, 1998; Hansen and Bellan, submitted) the magnetized plasma gun has the shape of a horse-shoe magnet, i.e., has two opposite magnet poles separated by a distance L. The vacuum magnetic field produced by the magnet has an arched shape. Applying a capacitor bank across the magnet poles breaks down neutral gas puffed into the region between the poles, creating a plasma. The capacitor then drives a current through the plasma. Initially this current follows the vacuum magnetic field lines, but as the current becomes stronger its hoop force causes the arched field to bulge out. The hoop force (Arzimovich, 1965; Bateman, 1978; Freidberg, 1987; Miyamoto, 1989; Biskamp, 1993; Chen, 1989; Krall *et al.*, 2000) causes a current following a curved path to increase its radius of curvature and is a direct consequence of the mutual repulsion between anti-parallel currents (just as the pinch force is the consequence of the mutual attraction between parallel currents). A circular loop of current (or current hoop) will have anti-parallel currents on opposite sides of the loop which mutually repel each other and cause the radius of the loop to increase. Equivalently, the hoop force can be thought of as the consequence of there being a higher magnetic pressure on the inside of current loop than on the outside (in the extreme example a long solenoid has magnetic pressure on the inside but not on the outside).

As in the spheromak case, field line tension opposes the hoop force in the prominence experiment and what results is the consequence of the competition between these two opposing magnetic forces. The restraining force is due to the vacuum magnetic field while the hoop force is due to gradients of magnetic pressure associated with the plasma current. The field aligned plasma current also causes the current channel to become twisted and to writhe as it goes from one magnet pole to another. In all prominence simulation experiments so far, the plasma has not detached from the gun and so λ is just the λ_{gun}. These experiments are thus analogous to Regime III of the spheromak experiments.

4. Alfvén Instability of Current Sheets

As mentioned earlier, MHD analysis predicts (Jarboe *et al.*, 1985; Fernandez *et al.*, 1989) that helicity flows from high λ to low λ; since λ is proportional to L^{-1}, the helicity flux has the effect of tending to flatten λ gradients. In resistive MHD, one finds that peaked λ profiles are unstable to resistive tearing modes (Furth *et al.*, 1963; Goldston and Rutherford, 1995) which flatten the current profile and hence the λ profile. However, most plasmas of interest (solar, magnetospheric, and laboratory) have insufficient resistivity for resistive tearing modes to be relevant because their growth rate is much too slow (hence the oft-used invocation of 'anomalous' resistivity).

Observational evidence points to some kind of collisionless reconnection process which similarly feeds on λ gradients and flattens these gradients but at a much faster rate than resistive reconnection. Controversy exists on whether these collisionless processes are continuous or transient; either situation may occur depending on boundary conditions, scale sizes, and symmetries. Observations also indicate that there is often an anomalous ion heating associated with reconnection (Mayo *et al.*, 1991; Ono *et al.*, 1996) and that various waves are excited as part of the reconnection (Gekelman and Stenzel, 1984). It has been postulated (Axford and McKenzie, 1992) that the high speed solar wind is accelerated by Alfvén waves originating from tiny reconnection regions in microflares in coronal holes at the base of the corona.

A simple model has been developed (Bellan, 1999) which shows that force-free equilibria with extremely peaked λ profiles are linearly unstable with respect to emission of inertial or kinetic Alfvén waves (Stasiewicz *et al.*, 2000). This model was motivated by an earlier calculation (Bellan, 1998) showing that the transient localized field-aligned current associated with magnetic reconnection must act as an antenna radiating Alfvén waves. The radiation resistance experienced by this antenna was postulated (Bellan, 1998) as constituting the anomalous resistivity required for rapid reconnection. Also, calculations of the field pattern excited by a transient, localized field-aligned current showed that this field pattern was consistent with observed auroral oscillations if one assumed that the source was far away and the disturbance propagated dispersively from the source to the point of observation (Bellan, 1996); this is the behavior one would expect for the waves emitted by a localized reconnection event.

If the equilibrium is low β and the magnetic field lines are locally straight, then the magnetic field magnitude is invariant. In this case all the field can do is rotate its direction. A region of peaked λ confined to a plane corresponds to a current sheet with associated magnetic field rotational discontinuity across the current sheet. MHD analysis of evolving complex field topology shows that current sheets typically develop when flux tubes become braided (Parker, 1983; Longbottom *et al.*, 1998).

If the local angle of rotation of **B** inside a rotational discontinuity is denoted by $\theta(x)$ and the total angle of rotation across the discontinuity is Δ, then the force-free magnetic field in the vicinity of the rotational discontinuity can be written as

$$\mathbf{B}(\mathbf{x}) = \hat{y}B\sin[\theta(x)] + \hat{z}B\cos[\theta(x)] \tag{6}$$

where

$$\theta(x) = \int_0^x \lambda(x')dx' \tag{7}$$

gives the local angle of rotation. Equation (6) is an exact solution of Eq. (2). Suppose the current sheet has width a and assume that λ is uniform in the current sheet and zero outside the current sheet. In this case integration of Eq. (7) across the current sheet gives the relation

$$\Delta = a\lambda. \tag{8}$$

However, using $|B| = const.$ together with Eq. (2) gives $J(x) = \lambda(x)B/\mu_0$. These relations can be combined to give the electron flow velocity $u_{\|e0}$ in the current sheet,

$$u_{\|e0} = \frac{J}{nq} = \frac{\lambda B}{nq\mu_0} = \frac{\Delta B}{anq\mu_0} = v_A\frac{c\Delta}{\omega_{pi}a}. \tag{9}$$

Thus if $a\omega_{pi}/c < \Delta$, the field-aligned electron flow becomes super-Alfvenic and destabilization of Alfvén waves becomes a possibility. It is interesting that this low β scaling of critical current sheet width in terms of the ion skin depth appears similar to high β collisionless reconnection models which also show that current sheet widths smaller than the ion skin depth are necessary (Drake *et al.*, 1994; Bhattacharjee *et al.*, 1999; Yamada *et al.*, 2000) and that instability of collisionless waves is associated with reconnection. Although the scaling with ion skin depth is similar, the physics is somewhat different in the low v. high β cases, because in the low β case instability results from super-Alfvenic field-aligned electron flow whereas in the high β case the current is not field-aligned and there is a null in the magnetic field at the middle of the current sheet.

Destabilization at a rotational discontinuity in a low β plasma is examined by calculating Alfvén wave behavior using a kinetic description for parallel particle motion and taking into account the presence of the beam of field-aligned electrons with streaming velocity $u_{\|e0} = J/nq$ in the current sheet region $-a/2 < x < a/2$. We assume that perturbed quantities vary as $g(x)\exp(ik_\|(x)s - i\omega t)$ where $\omega \ll \omega_{ci}$, $k_\| = k_z\cos\theta$, and s is the distance along **B**. Using a standard Vlasov analysis, the parallel wave current is given by

$$\tilde{J}_\| = \sum_\sigma q_\sigma \int dv_\| v_\| \tilde{f}_\sigma(v_\|) = \frac{i\omega\tilde{E}_\|}{\mu_0 c^2}\sum_\sigma \frac{1}{2k_\|^2\lambda_{D\sigma}^2} Z'\left(\frac{\omega - k_\| u_{\|\sigma 0}}{k_\| v_{T\sigma}}\right) \tag{10}$$

where Z is the plasma dispersion function and $v_{T\sigma} = \sqrt{2\kappa T_\sigma/m_\sigma}$.

Analysis of perpendicular particle dynamics shows that both electrons and ions have identical $\tilde{\mathbf{E}} \times \mathbf{B}$ drifts which therefore do not result in any perpendicular current. The lowest order perpendicular current thus comes from polarization drift, and since this is proportional to ion mass, the ion polarization drift $\tilde{\mathbf{u}}_{i,pol} = (m_i/q_iB^2)\partial\tilde{\mathbf{E}}/\partial\mathbf{t}$ is the dominant contributor to perpendicular current. Thus, the perpendicular wave current is

$$\mu_0\tilde{\mathbf{J}}_\perp = \frac{1}{v_A^2}\frac{\partial\tilde{\mathbf{E}}}{\partial t}. \tag{11}$$

Using Faraday's and Ampere's laws, all field components can be written in terms of $\tilde{A}_\|$. Combining the parallel component of Ampere's law with Eq. (10) gives the Alfvén wave equation (Bellan, 1999)

$$c^2\nabla_\perp^2\tilde{A}_\| = \left(\omega^2 - k_\|^2 v_A^2\right) \cdot \left[\sum_\sigma \frac{1}{2k_\|^2\lambda_{D\sigma}^2} Z'\left(\frac{\omega - k_\| u_{\|\sigma 0}}{k_\| v_{T\sigma}}\right)\right]\tilde{A}_\|. \tag{12}$$

For a uniform plasma this reduces to the kinetic Alfvén wave dispersion relation $\omega^2 = k_\|^2 v_A^2(1 + k_\perp^2\rho_s^2)$ if $v_{Ti} \ll \omega/k_\| \ll v_{Te}$ and to the inertial Alfvén dispersion relation $\omega^2 = k_\|^2 v_A^2/(1 + k_\perp^2 c^2/\omega_{pe}^2)$ if $v_{Te} \ll \omega/k_\|$. If there is no perpendicular dependence then Eq. (12) reduces to the MHD Alfvén wave dispersion $\omega^2 = k_\|^2 v_A^2$. By retaining displacement current, Eq. (12) can be extended to correspond to the standard low-frequency electrostatic dispersion relation which gives ion acoustic waves (Bellan, in press) in a magnetized plasma.

Thus, when there is a sufficiently strong current sheet, the field-aligned electrons flow super-Alfvenically inside the current sheet and so cause the argument of the electron Z' function to become negative in the current sheet. This causes the imaginary part of Z' to produce destabilization (instead of Landau damping) and so within the current sheet there is a source of instability. On the other hand, the waves destabilized inside the thin current sheet propagate out of the current sheet and this loss of wave energy to the exterior region will constitute a loading on the gain mechanism within the current sheet. The entire system (current sheet plus exterior region) will become unstable when the rate of creation of wave energy within the current sheet exceeds the rate of dissipation of wave energy in the exterior region. Numerical solutions of Eq. (12) show that destabilization occurs when the electrons become slightly super-Alfvenic.

Since the Alfvén wave instability feeds on the peaked λ profile, it is expected that non-linear analysis will show that the instability will tend to diminish this peaking and so spread out λ. Thus, the Alfvén instability tends to flatten λ profiles. It is also found that while the Alfvén waves are in general much faster than the ion thermal velocity (in accordance with the low β assumption) the rotation of the magnetic field at the current sheet causes $\omega/k_\|$ inside the current sheet to be slower than outside and so there can be significant Landau damping of destabilized Alfvén waves on ions within the current sheet. The net effect would be to produce beams of energetic ions in the current sheet as a side effect.

This analysis shows that peaked λ profiles become linearly unstable with respect to Alfvén wave emission when the condition

$$\frac{a}{c/\omega_{pi}} < \Delta \tag{13}$$

is satisfied. This indicates that the ion skin depth c/ω_{pi} is the appropriate 'yardstick' for measuring current sheet widths, and that instability results when the sheet width measured in units of ion skin depth is smaller than the angle of rotation of the magnetic field across the current sheet.

5. Summary and Conclusions

If a plasma is approximately force-free and approximately in equilibrium, then its magnetic field is determined by the simple force-free equation $\nabla \times \mathbf{B} = \lambda \mathbf{B}$. The behavior of the solutions to this equation are determined by boundary conditions and by the internal λ profile. Spheromaks and solar prominences are distinguished by their differing boundary conditions: spheromaks have azimuthal symmetry whereas prominences have mirror symmetry about a plane. In principle, spheromaks also differ in being bounded by a flux-conserving wall, but the experiments reported here show that unbounded spheromaks can be formed transiently.

Instabilities tend to flatten the λ profile and so cause helicity to flow from regions of high to low λ. If λ is sharply peaked, this corresponds to a current-sheet and at sufficiently strong peaking, the electron flow velocity in the current sheet becomes super-Alfvenic resulting in the spontaneous emission of Alfvén waves. This Alfvén instability is a form of beam-plasma instability (kinetic instability) and occurs when the current sheet width measured in units of ion collisionless skin depth becomes somewhat smaller than the rotation angle of the magnetic field across the current sheet.

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