Generation of a Reversed-Field Configuration without an Applied Magnetic Field

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(Received 5 July 1978)

A rotating relativistic electron beam, charge but not current neutralized and exceeding the Alfvén limit, has been propagated down a closed metal tube in the absence of an external guide field. The radial equilibrium is determined by the self-fields of the beam and the induced wall currents. The 100-nsec duration beam, upon leaving the system, induces plasma currents that maintain a field-reversed configuration for 18 μsec.

This Letter describes experiments in which a plasma in a reversed-field configuration, with both axial and azimuthal magnetic field components, has been produced inside a closed metal tube in which there is initially no field. The configuration is generated by a rotating relativistic electron beam injected into neutral hydrogen gas, and maintained by plasma currents induced when the beam leaves the system. Previous studies of similar beam-generated configurations[1-4] have all used an initial, externally applied, magnetic field. Reversal of the applied field by up to 4 times has been observed,[3] with a lifetime determined by the L/R decay of the currents in the fully ionized plasma ($n_e = 5\times10^{15}$ cm^{-3}, $T_e \sim 3-5$ eV).[4]

Radial equilibria are possible for both beams and plasmas inside a flux-conserving cylinder without an applied field. Yoshikawa[5] has described the equilibrium of a rotating beam in its self-induced fields, and has shown that in this configuration the beam current, I, is not subject to the Alfvén limit, $I < I_A = 17\,000\beta\gamma$ amperes. Arbitrarily large currents can then flow in a configuration that becomes increasingly force free as $I \gg I_A$.

To produce a rotating beam, an annular beam is first created by a diode in an axial magnetic field. The field is brought to zero a short distance from the anode by using a suitable arrangement of coils to divert the field lines radially outward. (This is known as a "half-cusp.") The interaction of the axial velocity of the beam with the radial component of the field gives the beam an azimuthal component of velocity[6]; the resulting hollow rotating beam thus generates both axial (B_z) and azimuthal (B_θ) magnetic fields. If the beam is injected into a closed metal cylinder, flux conservation requires that there should be an axial magnetic field, B_{z0}, between the beam and the wall in the opposite direction to the axial field, B_{zi}, inside the beam. The equilibrium radius of the beam is then determined by the balance of the magnetic and centrifugal forces on the electrons, and flux conservation.

The beam is injected into neutral gas, which is ionized by collisions with the beam electrons and the strong electric field induced by the rapidly rising magnetic field at the beam head. The gas pressure may be chosen so that the resulting plasma is sufficiently dense to charge neutralize, the beam. Thus the magnetic field of the beam is carried into the plasma. During the beam pulse the plasma is heated and its conductivity increased, so that when the beam leaves the system, currents are induced in the plasma to conserve the magnetic flux. The field of the beam is thus "frozen into" the plasma, and will remain for a time limited only by resistive dissipation of the plasma currents.

This sequence of events has been observed in the experimental apparatus shown in Fig. 1. An annular relativistic electron beam from the modified Triton accelerator[7] [$V = 900$ kV, $I = 110$ kA, $\tau = 100$ nsec full width at half-maximum (FWHM)] is injected through a half-cusp, located at $z = 0$, into a 14.6-cm-diam stainless-steel tube containing neutral hydrogen gas. The half-cusp is formed by a solenoidal coil around the cathode, which

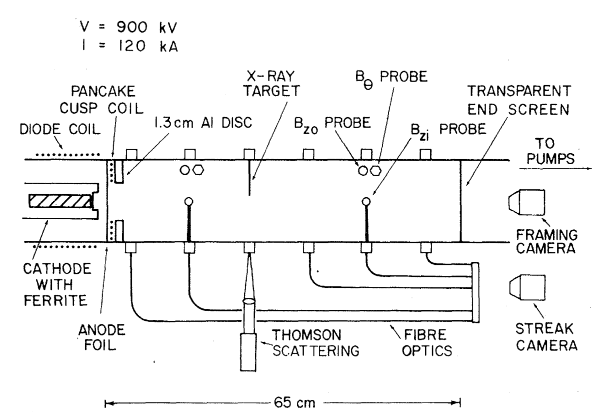

FIG. 1. The experimental facility.

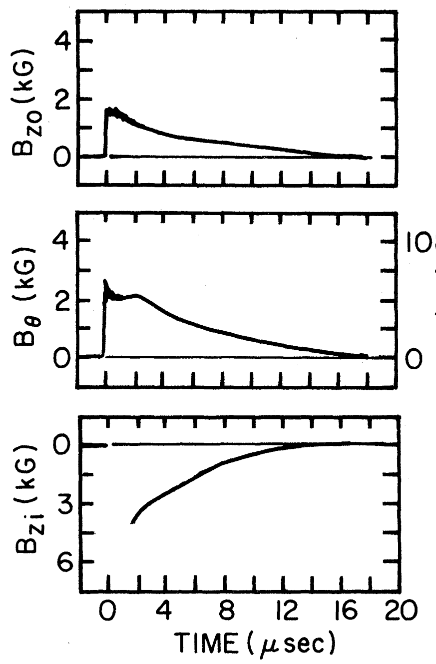

FIG. 2. Output of magnetic probes measuring $B_z(r=6.3\ \text{cm}) \equiv B_{zo}$, $R_\theta(r=6.3$ cm), and $B_z(r=0\ \text{cm}) \equiv B_{zi}$.

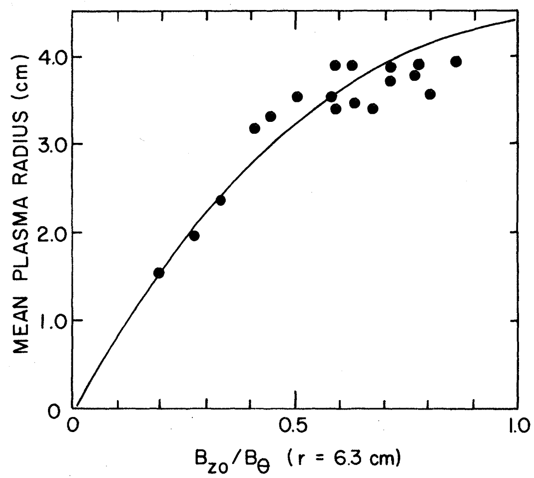

FIG. 3. B_{zo}/B_θ vs mean plasma radius.

contains a 15-cm-long ferrite cylinder, and a flat pancake coil, situated 0.3 cm from the anode foil and 0.2 cm from a 1.3-cm-thick aluminum plate, which excludes magnetic flux during the 400-μsec risetime of the current in the coils. Thus, the field lines emanate from the cathode perpendicular to the emission surface and pass out between the pancake coil and aluminum plate, resulting in a measured B_r axial extent (FWHM) of 1.8 cm. The system is terminated with a transparent brass screen at $z = 65$ cm.

Typical results are shown in Fig. 2. The traces show values of B_θ at $r=6.3$ cm, B_z at $r=6.3$ cm (i.e., B_{zo}), and B_z on axis (i.e., B_{zi}) as measured by three miniature magnetic probes. B_{zo} and B_{zi} are in opposite directions and indicate that a field-reversed configuration persists for 12 μsec. End-on framing photographs show that the plasma has an annular profile (typical mean radius 3.9 cm, annular width 1.5 cm) and is clearly separated from the tube wall (radius 7.3 cm). As the configuration decays, the plasma radius does not change, unlike in the guide-field case.[4] This is to be expected, since without the applied field, all the confining fields decay with the plasma.

The equilibrium position of the plasma differs from that of the beam because of the absence of a centrifugal force term in the radial balance. If the plasma pressure is low, the plasma currents are force free. The equilibrium radius of a thin plasma layer can then be simply found from pressure balance,

$$B_{zi}^2 = B_{zo}^2 + B_\theta^2, \tag{1}$$

combined with flux conservation,

$$B_{zi} r_p^2 + B_{zo}(r_w^2 - r_p^2) = 0, \tag{2}$$

where r_p and r_w are the radii of the plasma and wall, respectively. These equations lead simply to

$$\frac{r_p}{r_w} = \left(\frac{1-\cot^2\alpha}{2}\right)^{1/2}, \quad \frac{B_{\theta w}}{B_{zo}} = \left(\frac{2}{\tan^2\alpha - 1}\right)^{1/2}, \tag{3}$$

where $B_{\theta w}$ is B_θ at the tube wall, and α is the pitch angle of the helical plasma current (note that this model predicts no equilibrium unless $\alpha > 45°$). The pitch angle of the beam may be adjusted by changing the magnetic field in the half-cusp; increasing the field winds the beam into a tighter helix, increasing both B_{zo}/B_θ and the plasma radius. In Fig. 3, B_{zo}/B_θ, measured by magnetic probes at $r = 6.3$ cm, is plotted against the plasma radius, measured from framing photographs. Both quantities are obtained at $t = 2$ μsec. The solid curve is the prediction of the model in Eq. (3), and good agreement with the data is seen. The apparent limitation of the plasma radius at 4 cm was found to be due to the beam hitting the edge of the aluminum plate at the higher half-cusp magnetic fields, resulting also in reduced axial current and a marked decrease in plasma thickness.

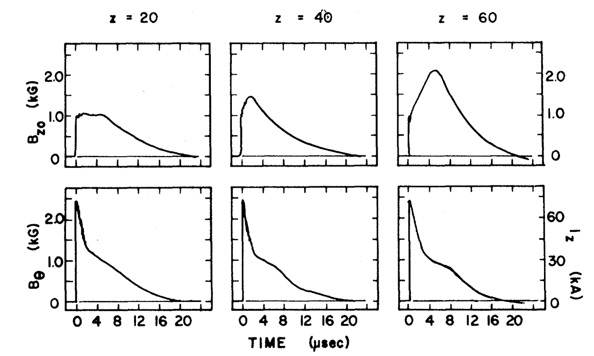

FIG. 4. Magnetic probes measuring B_{zo} (upper) and $B_\theta(r=6.3)$ at three axial positions.

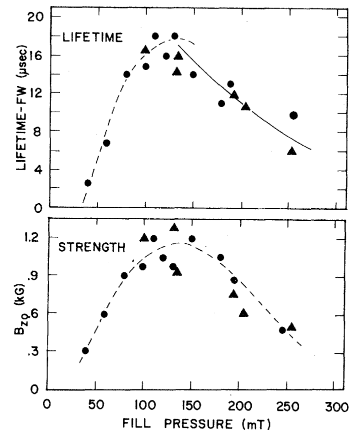

FIG. 5. Layer lifetime and strength vs fill pressure.

The B_{zi} probe, used to verify Eqs. (1) and (2), was found to have a perturbing effect on the plasma and was removed for subsequent measurements, since knowledge of $B_{\theta w}$, B_{zo}, and r_p is adequate to determine the configuration. With the probe removed, the configuration is created uniformly along the full 65-cm length of the tube and persists for approximately 18–20 μsec. This observation is in keeping with side-on streak photography, which shows that the light-emitting region has a similar axial extent. In Fig. 4 signals from identical magnetic probes at $z=20$, 40, and 60 cm are presented. Note that immediately after passage of the beam ($t=0$), the magnetic fields are uniform along the length of the tube. B_θ is shown in units of axial current on the right-hand scale. The current of 75 kA exceeds the Alfvén current ($I_A=43$ kA for 900 kV electrons), thus confirming the prediction of Yoshikawa.[5] This net current is, however, only 68% of the diode current; this loss may be due to some current neutralization of the beam or to some loss in transmission through the half-cusp.

As the configuration decays, B_θ changes uniformly along the tube, suggesting that the configuration is continuous over its length. However, B_{zo} at $z=60$ cm increases by a factor of 2 within the first 4 μsec, indicating that the rotating currents are piling up against the end screen. As B_{zo} at both 20 and 40 cm does not decrease, evidently magnetic energy is being transferred from the aximuthal field to the axial field, and is indicated by the rapid early decrease in B_θ. This observation can be explained by visualizing the plasma currents as a helical coil, which contracts in a manner to minimize its magnetic energy. (The tendency to collect at the end wall is probably due to the asymmetry introduced by a small residual magnetic field that has penetrated the aluminum cusp plate.) The overall lifetime of the configuration is consistent with the classical L/R decay time of the plasma currents, assuming an electron temperature of ~7 eV; it is also comparable to the time for plasma to free stream out the ends of the system.

The configuration strength (in terms of B_{zo}) and lifetime (full width), as determined by a magnetic probe at $z=20$ cm, are plotted as a function of gas pressure in Fig. 5. The data points include measurements taken with (triangles) and without (circles) a 1-mm-diam tungsten wire inserted across a radius of the tube at $z=30$ cm. The results are unaffected by the presence of the wire; since the wire would absorb any trapped beam electrons within 500 nsec, this confirms that the field-reversed configuration is indeed maintained by plasma currents alone.

Both lifetime and strength of the configuration have maxima at hydrogen pressures between 100 and 150 mTorr. Below 100 mTorr insufficient plasma is produced to charge-neutralize the beam, which will not propagate beyond $z=20$ cm. As the pressure is increased above 150 mTorr the beam is probably current neutralized to an increasing extent, while the energy deposited by the beam has to be shared by more particles, resulting in a lower electron temperature, T_e. If limited by classical resistive decay, the lifetime of the configuration would follow $\tau \propto T_e^{3/2}$. Assuming $T_e \propto n^{-1}$ leads to $\tau \propto n^{-3/2}$. The solid line in Fig. 5(a) represents this $n^{-3/2}$ scaling, and is in quite good agreement with the data.

The significance of these observations is that a rotating beam, charge but not current neutralized, with a current $I > I_A$, can (i) propagate with an equilibrium determined by its self-fields, as predicted by Yoshikawa,[5] and (ii) set up a reversed-field plasma configuration by inducing currents in the plasma and wall of a closed, initially field-free, metal tube.

In the present experiments, the configuration resembles a linear reversed-field pinch. It is possible to envisage extensions of this technique to produce plasma configurations with closed field lines. These could be further heated by the injection of intense neutral, electron, or ion beams; or by an imploding liquid metal liner, as in the Naval Research Laboratory LINUS fusion concept.[8]

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Stability of an Anisotropic High-β Tokamak to Ballooning Modes

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(Received 26 May 1978)

We have applied the Kruskal-Oberman energy principle to a simple model of an anisotropic tokamak in which the pressure varies around flux surfaces. We show that the weighting of pressure towards regions of favorable curvature leads to a significant stabilization of the high-n ballooning modes.

Following recent theoretical investigations of the MHD (magnetohydrodynamic) stability of *scalar*-pressure tokamaks,[1-6] it is now generally believed that the upper limit to β is set by the ballooning mode. Apart from its use as an additional heat source, neutral injection has been proposed as a method for "pumping-up" β in the flux-conserving tokamak[7]; it is also fundamental to the counterstreaming ion concept.[8] These applications have led us to consider the MHD stability of an anisotropic model of tokamak to high-n ballooning, n being the toroidal mode number.

Our analysis is based on the Kruskal-Oberman energy principle[9]; using the property of adiabatic invariance, Andreoletti[10] has shown their result to be independent of the form of distribution function. We assume that neutral injection is applied at an angle to the magnetic field such that hot ions are created only in the untrapped region of velocity space, so that the distribution function for the trapped particles is not significantly anisotropic. Then for small inverse aspect ratio, δ, the kinetic term in Kruskal-Oberman is $O(\delta^{7/2})$,[11] whereas the fluid terms are $O(\delta^2)$, when $\beta\sim\delta$. Thus, we drop the kinetic term, anticipating that our general analysis will be applied to a large-aspect-ratio model. Writing the fluid terms in a form as closely analogous to that for scalar pressure[5] as possible, we obtain

$$\delta W = \int d\tau \Big\{ (1-\sigma_-)\vec{Q}_\perp^{\,2} - (1-\sigma_-)\frac{\vec{J}\cdot\vec{B}}{B^2}(\vec{\xi}\times\vec{B}\cdot\vec{Q}) - 2(\vec{\xi}\cdot\vec{\kappa})(\vec{\xi}\cdot\nabla\bar{p}) + B^2(1+\sigma_\perp)\left[\left(1+\frac{1-\sigma_-}{1+\sigma_\perp}\right)\vec{\xi}\cdot\vec{\kappa} + \nabla\cdot\vec{\xi}\right]^2 + B^2\left(\frac{1-\sigma_-}{1+\sigma_\perp}\right)(\sigma_\perp+\sigma_-)(\vec{\xi}\cdot\vec{\kappa})^2 \Big\}, \tag{1}$$

where $\bar{p} = (p_\perp + p_\parallel)/2$, $\sigma_- = (p_\parallel - p_\perp)/B^2$, $\vec{Q} = \mathrm{curl}(\vec{\xi}\times\vec{B})$, $\xi_\parallel \equiv 0$, and $\vec{\kappa}$ denotes the field-line curvature. In order to define $\sigma_\perp$, we introduce the pressurelike moment

$$C = \sum_j m_j \iint \frac{B\,d\mu\,d\epsilon}{v_\parallel}\frac{\partial f_j}{\partial\epsilon}(\mu B)^2, \tag{2}$$