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Fission of a Hot Plasma*

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It has been observed that the plasma formed in a fast B_z compression develops an instability after about a 4- to 5-μsec quiescent period. The consequence of the instability is that the plasma fissions into two well-defined filaments that rotate with an angular velocity $\Omega \sim 10^7$ sec^{-1}. This configuration has been observed for many rotational periods. The plasma appears to remain confined although impurity radiation indicates that some plasma reaches the walls during the fission process. The fission process in a deuterium plasma is accompanied by a rapid decline in neutron production. The origin of the rotation and the development of the instability are discussed.

I. INTRODUCTION

THE fast B_z compression has previously been discussed in detail.[1] In the quiescent state after collapse and after the shock bouncing subsides, the deuterium plasma has the following properties[2]:

Plasma radius	$r_0 \cong 4$ mm,	(streak camera photographs)
Electron temperature	$T_e \cong 1$ kev,	(soft x-ray emission)
Ion temperature	$T_I \cong 1$ kev,	
Electron density	$n \cong 5\times 10^{16}$ cm^{-3}.	(optical continuum)

The external magnetic field reaches a maximum $B_z \sim 80$ kgauss at ~ 5 μsec. The ion cyclotron frequency is estimated to be $\omega_I \cong 2\times 10^8$ sec^{-1} for ions immersed in the plasma at the onset time of the instability $\omega_I \cong 10^8$ sec^{-1}, and the ion Larmor radius is $a_I \cong 1$ min. Initially the plasma has a temperature of 1.2 ev with 50% ionization from a preheater discharge and contains a trapped B_z field in the opposite direction, normally between 3 and 6 kgauss. During the early stages of the compression the internal reversed B_z is observed with probes and is of the same order of magnitude as the external B_z due to flux conservation. The consequences of the dissipation of this trapped field are discussed later. The ion temperature is presumed to be of the same order as the measured electron temperature in *estimating* a_I. The observed neutron yield is also consistent with kilo-electron-volt ions (mean energy), although the ion temperature has not been measured directly.

Streak camera photographs of the plasma were available in 1958.[1] Changes in the apparent radius of the plasma that developed from 2–5 μsec (depending on the coil length) after the compression were thought to be radial oscillations of the plasma cylinder. The present paper is concerned with a new interpretation of these photographs as well as some more recent data.

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† The work of this author was carried out under a joint General Atomic-Texas Atomic Energy Research Foundation program on controlled thermonuclear reactions.
‡ The experimental work reported here was carried out under the sponsorship of the U. S. Atomic Energy Commission and the Office of Naval Research.

[1] A. C. Kolb, *Proceedings of the Second United Nations International Conference on the Peaceful Uses of Atomic Energy, Geneva, 1958* (United Nations, Geneva, 1958), Vol. 31, pp. 328–340; A. C. Kolb, H. R. Griem, and W. R. Faust, *Fourth International Conference on Ionization Phenomena in Gases, Uppsala, Sweden* (North-Holland Publishing Company, Amsterdam, 1960), Vol. II, pp. 1037–1041.

[2] H. R. Griem, A. C. Kolb, W. H. Lupton, and D. T. Phillips, Bull. Am. Phys. Soc. **6**, 205 (1961); *Conference on Plasma Physics and Controlled Nuclear Fusion Research, Salzburg, Austria*, September, 1961, Nuclear Fusion J. (to be published).

II. EXPERIMENTAL DATA

Stereoscopic streak photographs of the plasma are shown in Fig. 1. The plasma luminosity is due to the visible bremsstrahlung continuum, so that it is a measure of electron density. Departures from electrical neutrality may obtain for distances of the order of the Debye length $L_D \sim (kT_e/4\pi ne^2)^{\frac{1}{2}} \sim 10^{-4}$ cm. Since $L_D \ll r_0$ it is reasonable to suppose that a knowledge of the electron density implies a knowledge of the ion density, i.e., that the streak photographs show where the plasma is.

The photographs of Fig. 1 were obtained with two slits placed in the y-z and x-z planes at the same position z. Similar results were obtained at different positions z along the plasma.

It should be noted that the nodes in the side view of Fig. 1 always appear at the same time as the loops in the top view and vice versa. This indicates that the plasma cylinder first flattens and then separates into two well-defined pieces that rotate about the z axis. The probable appearance of the plasma to an observer in a frame of reference rotating with the plasma is illustrated in Fig. 2.

The rate of rotation of the plasma can be obtained directly from the photograph in Fig. 1. The time between nodes is π/Ω so that $\Omega \cong 10^7$ sec^{-1}. It is not possible to obtain the wavelength λ from the present photographs because they are taken at the same z. λ could be obtained by employing slits at different values of z, or with a single slit perpendicular to the x-y plane.

In some of the photographs the two parts of the plasma are consistently distinguishable due, for example, to differences in minor diameter. It is then possible to determine the direction of rotation. The direction of rotation was determined for ten cases and always found to be in the sense that ions would rotate in the external B_z field.

An apparently similar phenomenon has previously been observed for rotating shaped charge jets. When the rotational frequency is large enough to produce fission, flash x-ray photographs[3] of the jet look very similar to Fig. 2. The stability of a rotating liquid mass with a gravitational field was considered by Darwin and Jeans.[4] The formation of binary stars was attributed to the fission of such bodies. Although neither of these cases relate directly to the plasma under consideration, the superficial similarity provided the motivation for further study of the "radial oscillations" and the present interpretation.

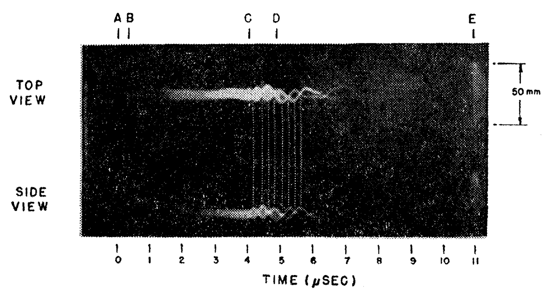

FIG. 1. Stereoscopic streak photographs of the plasma compression and $m=2$ instability mode with -3 kilogauss initial reverse field. (A) Start of discharge; (B) imploding shock reflects off axis; (C) onset of instability; (D) fission into two well-defined filaments; (E) imploding shock wave at beginning of the second half-cycle. The camera slit is imaged to the left of the photograph and defines the relative positions of the two filaments in the side and top views.

III. INTERPRETATION OF THE INSTABILITY

Consider the stability of an infinite cylinder of plasma with a magnetic field $B_z=\alpha_v B_0$ outside the plasma and $B_z=\alpha_p B_0$ inside the plasma. The usual procedure is to Fourier analyze the problem which is equivalent to considering radial displacements of the form

$$\xi=\xi_0 \exp[i(\omega t+m\theta+kz)].$$

In Fig. 3, radial displacements of this type are plotted for constant z and t. $m=0$ and $m=1$ are the usual "sausage" and kink instabilities. Modes for which $m>1$ are usually called flute instabilities. If ξ_0 is large as indicated in Fig. 3 it is apparent that the modes for which $m>1$ can be interpreted as fission modes. Moreover there would be an apparent rotation of the plasma with angular velocity $\dot{\theta}=-\mathrm{Re}\,\omega/m$. Therefore the present observations are consistent with an $m=2$ instability.

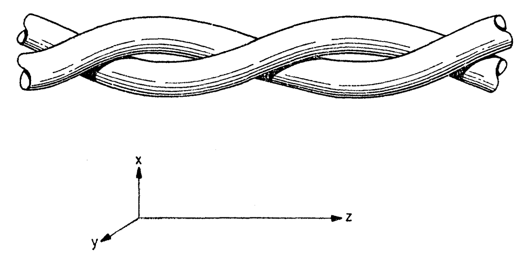

FIG. 2. Plasma configuration in a frame of reference rotating with the plasma after the $m=2$ fission is complete.

The result of a stability analysis with the magnetohydrodynamic approximation[5] is

$$\omega^2=c_A{}^2k^2\left\{\alpha_p{}^2+\alpha_v{}^2\frac{G_{\mu m}(kr_0)H_m(kr_0)-L_m(kr_0)}{H_m(kr_0)[1-G_{\mu m}(kr_0)]}\right\}, \quad (1)$$

where $c_A{}^2=B_0{}^2/4\pi\rho_0$, ρ_0 is the initial plasma density,

$$G_{\mu m}(x)=K_m'(\mu x)I_m'(x)/K_m'(x)I_m'(\mu x),$$

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[3] R. Schall and G. Thomas, *Actes de 2me Congres International de Photographie et Cinematographic Ultra Rapide* (Dunod, Paris, 1956), p. 261.

[4] J. Jeans, Phil. Trans. Roy. Soc. **A67**, (1902), *Astronomy and Cosmogony* (Cambridge University Press, New York, 1929).

[5] M. N. Rosenbluth, Los Alamos Scientific Laboratory Report LA-2030, 1956 (unpublished).

$H_m(x)=I_m(x)/xI_m'(x)$, $L_m(x)=K_m(x)/xK_m(x)$; I_m, K_m are Bessel functions, $I_m'=dI_m/dx$, and $\mu=R/r_0$ where R is the radius of the perfectly conducting boundary. ω^2 is real and positive for all modes and all possible values of α_p, α_v, μ, and B_0. In the limit $\mu\rightarrow\infty$, $kr_0\rightarrow 0$ Eq. (1) has the asymptotic form

$$\begin{aligned}\omega^2&=c_A{}^2k^2[\alpha_p{}^2-\tfrac{1}{2}\alpha_v{}^2(kr_0)^2\ln kr_0-\cdots] \quad \text{for} \quad m=0\\ &=c_A{}^2k^2[\alpha_p{}^2+\alpha_v{}^2] \quad \text{for} \quad m\geq 1.\end{aligned} \tag{2}$$

In particular when $k\rightarrow 0$ the plasma has only marginal stability.

The above result is obtained for an initial state in which there is no macroscopic motion. If we assume an initial rotation such that $V_\varphi=r\Omega$ where Ω is constant, the dispersion relation can be obtained as a special case of the result of Gerjuoy and Rosenbluth.[6] In the limit $\mu\rightarrow\infty$, $kr_0\rightarrow 0$,

$$\begin{aligned}\omega^2&=\alpha_p{}^2c_A{}^2k^2-\tfrac{1}{2}\Omega^2(kr_0)^2 \quad \text{for} \quad m=0\\ \omega+m\Omega&=\Omega[1\pm(1-m)^{\frac{1}{2}}] \quad \text{for} \quad m\geq 1.\end{aligned} \tag{3}$$

For $kr_0\ll 1$ the modes $m=0$ and $m=1$ are stable or their growth rates are proportional to kr_0. For $m=2$, $\omega=-\Omega+i\Omega$ so that the $m=2$ mode grows with an e-folding time of $1/\Omega$ and an apparent rotation of $\dot{\theta}=\Omega/2$. Modes for which $m>2$ are also unstable and their growth rates increase with m.

We would expect nonlinear effects to become important when the amplitude of the perturbation approaches the wavelength. Nonlinear effects are then expected to inhibit further growth of the instability. Direct observations of the growth of instabilities usually show that the long-wavelength perturbations dominate.[7] We therefore expect that perturbations with $k\rightarrow 0$ and the smallest value of m for which the mode is unstable would be observed.

There is an additional reason to expect that modes for which $m>2$ would be suppressed. The radius of the plasma is only about four times an ion Larmor radius. It does not seem possible for fission modes to develop which result in fragments of the plasma that are smaller than an ion Larmor radius; i.e., finite Larmor radius effects should stabilize the modes for which $m\gg 2$.

A quantitative discussion of the finite Larmor radius effects will be given in another paper.[8] For present purposes we shall summarize the results that are relevant to the fast B_z compression. The magnetohydrodynamic approximation is based on an expansion to lowest order in $\lambda=a_i/r_0$, ω/ω_i, ω/ω_p where a_i is the Larmor radius and ω_i the Larmor frequency for particles of species i, and ω_p is the plasma frequency. The result $\omega\sim c_Ak$ of Eq. (2) indicates that for sufficiently small $k\sim a_i/r_0{}^2$, $\omega/\omega_i\sim(a_i/r_0)^2=O(\lambda^2)$ instead of $O(\lambda)$ and we must keep terms like $(a_i/r_0)^2$ to be consistent. A consistent calculation has been carried out for a particular case where there is initially a small radial electric field of order λ of the form $E_x=-(B_z/c)\Omega x$, $E_y=-(B_z/c)\Omega y$. A simple stability criterion obtains in the limit that $k\rightarrow 0$ and $\beta=8\pi P_0/B_z{}^2\rightarrow 0$ where P_0 is the equilibrium plasma pressure. The modes $m=0$ and $m=1$ are marginally stable and modes for which $m>1$ are stable if

$$(1-\sqrt{m})/2<\Omega/\Omega_0<(1+\sqrt{m})/2, \tag{4}$$

where $\Omega_0=2(a_I/r_0)^2\omega_I$, a_I is the Larmor radius, and ω_I is the ion Larmor frequency. For example, if $m=2$ the stability limits are $-0.207<\Omega/\Omega_0<1.207$, and for $m=3$, $-0.366<\Omega/\Omega_0<1.366$. The negative sign corresponds to rotation in the direction of ions in the external field. It is thus apparent that much less rotation can be tolerated in this direction and that the stability limits increase as m increases. The experimental values of Ω and Ω_0 are $\Omega\cong-10^7$ sec^{-1} and $\Omega_0\cong 3\times 10^7$ sec^{-1}, so that $\Omega/\Omega_0\cong-0.3$ which is in satisfactory agreement with the fact that the $m=2$ instability is observed.

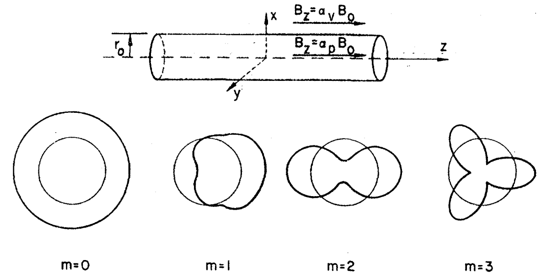

FIG. 3. Radial displacements at a particular z and t for the $m=0$, 1, 2, 3 instability modes.

IV. ORIGIN OF THE ROTATION

The rotational frequency is of the order of Ω_0 which represents an angular momentum of the order of $Mr_0{}^2\Omega_0$ per unit length to be accounted for. M is the mass per unit length of the plasma. For a plasma that is infinite in the z direction, the angular momentum per unit length is

$$L=\sum_i m_i\int\int(xv_y-yv_x)f_i(x,\mathrm{v})d\mathrm{v}dxdy. \tag{5}$$

After the plasma has become detached from the walls, L can only change as a result of leakage. If electrons and ions leak at equal rates the angular momentum carried off by the ions dominates. Each ion removes an angular momentum of about $m_Ia_I{}^2\omega_I=Mr_0{}^2\Omega_0/N$, where N is the number of particles per unit length. It is thus clear that to account for the observed rotational frequency a substantial fraction of the plasma would have to leak out of the ends, and even then the predicted

[6] E. Gerjuoy and M. N. Rosenbluth, Phys. Fluids, 4, 112 (1961). The particular result is that of Eq. (38).

[7] D. J. Albares, N. A. Krall, and C. L. Oxley, Phys. Fluids (to be published).

[8] M. N. Rosenbluth, N. Rostoker, and N. A. Krall, *Conference on Plasma Physics and Controlled Nuclear Fusion Research, Salzburg, Austria, September, 1961*, Nuclear Fusion J. (to be published).

sign of rotation would be wrong. We therefore neglect leakage and assume that the value of L that obtains immediately after detachment from the walls remains constant.

After the fast compression has taken place, the plasma contains a trapped B_z field in a direction opposite to the external field. From the slow compression following the initial implosion one would not expect ion temperatures in excess of about 0.5 kev on the basis of numerical calculations[9,10] (assuming adiabatic compression). The additional heating presumably results from a process such as that discussed by Rosenbluth[11] in which the initially crossed internal and external magnetic fields relax. Such a change in the internal magnetic field could also account for the rotation in the present experiment.

Consider a stationary state of the plasma in which there is a radial electric field $E_r=O(\lambda)$. f_i satisfies the Vlasov equation

$$\mathbf{v}\cdot\nabla f_i+\frac{q_i}{m_i}\left(\mathbf{E}+\frac{1}{c}\mathbf{v}\times\mathbf{B}\right)\cdot\frac{\partial f_i}{\partial\mathbf{v}}=0. \tag{6}$$

A power series expansion in $\lambda=a_i/r_0$ can be obtained for f_i,[12]

$$f_i=f_i^{(0)}(r,v_\perp,v_z^2)+\frac{1}{\omega_i}\left(\frac{\partial f_i^{(0)}}{\partial r}+\frac{q_iE_r}{m_iv_\perp}\frac{\partial f_i^{(0)}}{\partial v_\perp}\right)\frac{xv_y-yv_x}{r}, \tag{7}$$

where $v_\perp^2=v_x^2+v_y^2$. For example, let

$$f_i^{(0)}=n_i(r)(m_i/2\pi\Theta_i)^{\frac{3}{2}}\exp(-m_iv_\perp^2/2\Theta_i). \tag{8}$$

The density $n_i(r)=n(r)$ is independent of species because $\delta n/n\sim(L_D/r_0)^2\ll\lambda$; where L_D is the Debye length. From the appropriate moments of Eq. (7),

$$L=\sum_i\int 2\pi r^2dr\left\{-\frac{cE_r}{B_z}n(r)m_i+\frac{\Theta_i}{\omega_i}\frac{\partial n}{\partial r}\right\}, \tag{9}$$

and

$$J_\theta=-\frac{c}{4\pi}\frac{\partial B_z}{\partial_r}=\frac{c}{B_z}\sum_i\Theta_i\frac{\partial n}{\partial r}. \tag{10}$$

Let

$$Mr_0^2\Omega=-\sum_i\int 2\pi r^2dr\frac{cE_r}{B_z}n(r)m_i,$$

substitute Eq. (10) into Eq. (9), and neglect the electron contribution in the second term. Then

$$L=Mr_0^2\Omega+2\pi\frac{m_I}{e}\frac{\Theta_I}{\Theta_e+\Theta_I}\int J_\theta(r)r^2dr. \tag{11}$$

Immediately after the compression we may assume that Ω is within the stable bounds because the plasma is observed to be stable for about 4 to 5 μsec. The current density J_θ generates the large difference between the external and reversed internal B_z fields. As the internal B_z field disappears, $\int J_\theta(r)r^2dr$ increases (algebraically). There must be a corresponding decrease in Ω to keep L constant. In this way a negative value of Ω can be obtained corresponding to rotation in the direction of ions in the external field. The essential point is that a substantial fraction of the diamagnetic current density is carried by ions. For the particular choice of $f_i^{(0)}$ given by Eq. (8) the fraction of current density carried by ions is $\Theta_I/(\Theta_e+\Theta_I)$. This is a reasonable assumption for $f_i^{(0)}$ since the density is high and the collision times are short during the early stages of the compression. It satisfies all of the relevant equations, but is not unique. The present discussion does not depend on this particular choice, but only on the feature that the ions carry a substantial fraction of the diamagnetic current. This represents a large angular momentum and when the current changes due to the relaxation of the internal field, the corresponding angular momentum appears as a guiding center drift so that the total angular momentum is conserved.

Insufficient information is available on the field-mixing process[1,2,13] to make a prediction of more than the sign and the order of magnitude of Ω. For example, assume that initially $\Omega=0$ and $\partial B_z/\partial r\sim 2B\delta(r-r_1)$, where B is the external field and $r_1<r_0$ is the location of a sheet current that produces the internal reversed B_z. In the final state we simply neglect J_θ because without the internal reversed B_z, J_θ is relatively small. Then, taking $L=$constant and making use of Eq. (11), we have

$$\Omega=-(2\Omega_0/\beta)(r_1/r_0)^2. \tag{12}$$

If we take $\beta=8\pi n\theta_I/B^2$ to be a few tenths (from pressure balance considerations) and assume $r_1=r_0$ this result has the correct sign, but it is about an order of magnitude too large. The order of magnitude would be consistent with the demands of the experimental data if we assume $r_1/r_0<1$ or that the relaxation of the internal fields is incomplete after 5 μsec.

If one associates the time at which the reverse field begins to disappear rapidly with the observed rapid rise in the electron temperature at 2 to 3 μsec, an estimate of r_1/r_0 can be obtained by assuming that the internal trapped field $B_z\cong-50$ kilogauss (estimated from flux conservation) is dissipated and the energy $\pi r_1^2B_z^2/8\pi$ per unit length is converted to thermal energy of $n\pi r_0^2$ particles per unit length.

Calculations[9,10] of the temperature expected from shock heating and adiabatic compression yield $T_e\lesssim 0.5$ kev, while the observed electron temperatures are in the kilovolt range. Consequently, taking $\sim$0.5 kev to

[9] A. C. Kolb, W. R. Faust, and A. D. Anderson, Nuclear Fusion J. (to be published).
[10] K. Hain and A. C. Kolb, *Conference on Plasma Physics and Controlled Nuclear Fusion Research*, Salzburg, Austria, September, 1961, Nuclear Fusion J. (to be published).
[11] M. N. Rosenbluth, *Proceedings of the Second United Nations International Conference on the Peaceful Uses of Atomic Energy, Geneva, 1958* (United Nations, Geneva, 1958), Vol. 31, pp. 85–89.
[12] M. N. Rosenbluth and N. Rostoker, Phys. Fluids **2**, 23 (1959).
[13] A. C. Kolb, C. B. Dobbie, and H. R. Griem, Phys. Rev. Letters **3**, 5 (1959).

be the particle energy due to field mixing

$$(r_1/r_0)^2(B_z^2/8\pi n)\cong 0.5 \text{ kev},$$

one finds $(r_1/r_0)^2$ and $\Omega\cong 2-5\times\Omega_0$ which is of the right order of magnitude.

The comparison between the experimental data on the fission instability and the above calculations support the hypothesis that the dissipation of the reversed trapped field is responsible for both the instability and the observed electron temperatures. Consideration of Eqs. (4) and (12) show that one can expect this instability to appear in situations where $r_1/r_0\sim 1$ if there is some mechanism which causes the reverse field to transfer its energy to the plasma. The time scale for such dissipation will increase with increasing plasma radius if the resistivity is given by the conventional Spitzer formula. However, if the fields interdiffuse rapidly due to some other instability of the reverse current sheet, the fission mode may appear early in time even in experiments where the plasma radius is relatively large.

However, experimental observations[1] of the time at which the fission mode appears in coils of various lengths show that the stable time can be increased by increasing the coil length. With longer coils the confinement time is increased because end losses are decreased[10] and this results in a larger plasma radius at a given time. These observations suggest that the time over which the field mixing occurs increases with r_0 and that perhaps the onset time of the fission instability can be delayed by appropriate scaling.[14]

ACKNOWLEDGMENT

We would like to express our appreciation to Dr. Marshall Rosenbluth for his essential criticism and suggestions.

[14] *Note added in proof.* This instability has not yet been observed in PHAROS, a larger experiment (coil length 180 cm, diameter 10.5 cm, $H_{max}=80$ kgauss in 9 μsec, plasma radius 1.7 cm at peak field) with an initial reverse field of -3.8 kgauss. [A. C. Kolb, H. R. Griem, W. H. Lupton, D. T. Phillips, S. A. Ramsden, E. A. McLean, W. R. Faust, and M. Swartz, *Conference on Plasma Physics and Controlled Nuclear Fusion Research, Salzburg, Austria, September, 1961,* Nuclear Fusion J. (to be published)].

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High-Voltage Glow Discharges in D_2 Gas. I. Diagnostic Measurements*

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Pulsed glow discharges in D_2 gas at currents of the order of 1 amp and voltages in the range 40–80 kv are experimentally examined to determine the mass and energy distribution of ions incident on the cathode, the secondary electron emission due to various particles incident on the cathode, the energy spectrum of the electrons incident on the anode, the potential distribution along the glass walls, and the temperature and density of the electrons in the plasma which develops in the anode end of the discharge. The experimental results are discussed together with available cross-section data from the literature to establish the relative importance of the various collision processes which contribute to the sustainment of the discharge. Several processes peculiar to the extremely high operating potential (not of particular importance in ordinary glow discharges) are shown to assume dominant roles. These include electron backscattering from the anode, ionization of the gas by fast ions and fast neutral atoms, and ionization in the plasma region due to the secondary electrons released from the glass walls. The necessity of a complete revision of the usual theoretical point of view regarding electron ionization in the cathode region of the discharge and in the anode (plasma) region is indicated.

I. INTRODUCTION

It has been shown by Pokrovskaia[1] that extremely intense glow discharges, conducting currents of 100 amp at applied potentials of 15 kv, exist under quasi-steady (pulsed) conditions in hydrogen gas. These discharges were observed between plane parallel electrodes 1.6 cm apart at a gas pressure of 0.21 mm Hg. Using five times larger electrode spacings and fivefold lower pressures, Pollock[2] found that 1- to 10-amp glow discharges could be sustained in deuterium for several μsec at potentials of the order of 100 kv. This discovery led to the development of compact neutron source tubes, filled with deuterium (D) or tritium (T) gas and fitted with D- or T-loaded cathodes, which produce neutron yields of $\sim 10^5$ per μsec for several tens of microseconds.[3]

The conditions of voltage, current, and pressure under which these discharges operate are apparently beyond the realm of applicability of existing glow discharge theories. Neu[4] has presented a theoretical discussion of high-current, high-voltage discharges which introduces many applicable concepts, but there is reason to believe

* Work performed under the auspices of the Atomic Energy Commission.

[1] A. S. Pokrovskaia and B. N. Kliarfel'd, Soviet Phys.—JETP 5, 812 (1957).

[2] H. C. Pollock, General Electric Research Laboratories, Schenectady, New York (private communication, April, 1957).

[3] R. D. Kelley, J. C. Hamilton, and L. C. Beavis, Rev. Sci. Instr. 32, 178 (1961).

[4] H. Neu, Z. Physik 155, 77 (1959).

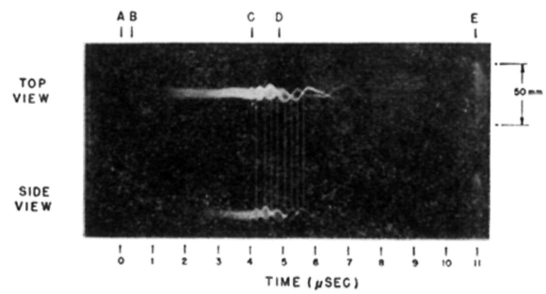

FIG. 1. Stereoscopic streak photographs of the plasma compression and $m=2$ instability mode with -3 kilogauss initial reverse field. (A) Start of discharge; (B) imploding shock reflects off axis; (C) onset of instability; (D) fission into two well-defined filaments; (E) imploding shock wave at beginning of the second half-cycle. The camera slit is imaged to the left of the photograph and defines the relative positions of the two filaments in the side and top views.