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ABSTRACT

The formative time of electric breakdown in low-pressure (0.2 to 2.0 torr) hydrogen across a strong magnetic field ($10 < \omega_b\tau < 350$; maximum B of 18 kG) has been measured in a coaxial cylindrical geometry. Attention was centered on the region of breakdown that occurs with a formative time less than the time required for an electron to drift across the electrode gap in the applied fields. This crossing time was inferred by extrapolations of previous measurements by Bernstein. These formative time measurements are compared with a simplified theory that assumes a constant number of e-folding times until breakdown, and neglects electron losses as well as secondary production at the cathode. This model predicts that the formative time is inversely proportional to the gas pressure and otherwise a function of only the ratio E/B and not of either field separately. The predicted pressure dependence is confirmed, but some deviations from the predicted functional dependence on E/B are found. These deviations are attributed to electron losses along the magnetic field. A prediction of the magnitude of the formative time based on this simplified theory must necessarily involve extrapolation of certain previously obtained results. Such a prediction is found to be in reasonable agreement with the experimental values.

I. INTRODUCTION

We describe here a study of the formative time of the electrical breakdown of hydrogen gas in an electric field perpendicular to a strong uniform static magnetic field. The magnetic field $\vec{B}$ is parallel to the electrodes, and it is of such magnitude that the electron gyrofrequency ω_b is much greater than the electron collision frequency τ^{-1}. In fact, during the period of ionization buildup the ratio of these frequencies $\omega_b\tau$ varies from 10 to 350 over the experimental range of pressure (0.2 to 2.0 torr) and magnetic field (6 to 18 kG). The voltage across the electrode gap is pulsed and rises to a value high enough above breakdown threshold that the formative time turns out to be shorter than the time required for an electron to drift from the cathode all the way to the anode (as inferred from the measurements by Bernstein[1]). This condition, which requires electric fields $\vec{E}$ of several kilovolts per centimeter, means the breakdown is accomplished essentially by the primary avalanches alone, and the effect of secondary electrons released at the cathode can in a first approximation be neglected. It is thus not surprising that our results differ markedly from those obtained by Deutsch, who investigated the formative time of cross-field breakdown involving several electron transits.[2] The simplifications under our condition permit theoretical predictions for the breakdown time as a function of the applied fields and the gas density. These predictions are based on extrapolated values of observed or computed ionization rates and drift speeds, for the authors are not aware of any direct measurements of these quantities in the phyical domain of the present experiment.

In the crossed-field configuration with $\omega_b\tau >> 1$ the motion in the plane normal to the magnetic field of a growing electron avalanche consists of a drift in the $\vec{E} \times \vec{B}$ direction and a slower drift in the $-\vec{E}$ direction, in the laboratory frame of reference. The velocity of the former drift is very nearly given by $\vec{v}_d \approx c\vec{E} \times \vec{B}/B^2$, and the speed of the latter drift, v_E, is approximately $v_E \approx (\omega_b\tau)^{-1}v_d$. The drift motion v_E, by means of which the electrons gain energy from the electric field, is a result of the collisions between the electrons and the neutral gas. For positive ions in this experiment the ratio of the gyrofrequency to the collision frequency is less than unity; therefore the positive ions drift essentially straight along $\vec{E}$ to the cathode. The speed of this ion drift motion is much greater than v_E.

II. THEORY

As an appropriate model we consider a cold neutral gas at rest with respect to the electrodes in uniform static orthogonal electric and magnetic fields with $E < B$ (in gaussian units). The treatment is limited to the case of a strong magnetic field, i.e., $\omega_b\tau >> 1$. We neglect any secondary electrons released from the cathode, and also any loss of electrons from the growing avalanche. All speeds are assumed to be nonrelativistic.

Rather than considering the problem in the laboratory frame (in which the drift motions are as described above), we choose to use a reference frame moving at velocity $\vec{v}_d$ with respect to the laboratory frame. In this "drift frame" the electric field vanishes, and the gyrating electrons are in a gas wind of velocity $-\vec{v}_d$. Note that in the drift frame the large electron drift velocity $\vec{v}_d$ is no longer present; therefore the electron velocity distribution function is much more nearly isotropic than it is in the laboratory frame. This means the usual expansion of the distribution function in spherical harmonics converges more rapidly in the drift frame.

In the drift frame, then, the isotropic part of the electron velocity distribution function f_0 can be shown[3] to depend parametrically on the ratio E/B only and not on E, B, or p (gas pressure) separately, if the collisional scattering of the electrons is independent of the azimuthal angle about the direction of the incident velocity. That is,

$$f_0 = f_0(v,t;\ v_d), \tag{1}$$

where v is the electron velocity in the drift frame. This result can of course be derived from a complete analysis of electron dynamics such as described by Allis,[4] but it can be understood with the help of a single physical argument. In the drift frame there is no electric field, so the Boltzmann equation becomes, to lowest order,

$$\frac{\partial f_0}{\partial t} = \left(\frac{\partial f_0}{\partial t}\right)_{coll}, \tag{2}$$

where $(\partial f_0/\partial t)_{coll}$ refers to the collisions with molecules only and electron-electron collisions are neglected. Thus the velocity distribution function is determined by the relative rates of all the collisional processes, and the magnetic field does not enter separately. Similarly, the gas pressure does not affect the velocity distribution function, even though τ^{-1} is admitted to be a function of velocity, because all the collision rates are proportional to the gas density.

Physically it is expected that in a "few" mean collision times after the initiation of an avalanche the electron velocity distribution function reaches an "equilibrium" shape at which the energy gain due to elastic collisions is balanced by the energy loss due to inelastic collisions. Then the distribution function is of the form

$$f = f(\vec{v};\ v_d)\ \exp\ (\beta t), \tag{3}$$

where β is the ionization rate. In what follows the distribution function is assumed to be of this form.

The ionization rate β is given, to the lowest order in the spherical harmonic expansion, by

$$\beta = 4\pi n_g \int_0^\infty \sigma_i(|\vec{v} + \vec{v}_d|)\,|\vec{v} + \vec{v}_d|\,v^2 f_0(v;\ v_d)dv, \tag{4}$$

where n_g is the gas density and $\sigma_i(|\vec{v} + \vec{v}_d|)$ is the ionization cross section. It therefore follows that the ionization rate, which is of course independent of the reference frame, is proportional to the gas density and otherwise is a function of E/B only. Symbolically, this dependence can be expressed in the form

$$\beta = n_g F\left(\frac{E}{B}\right). \tag{5}$$

As indicated in Fig. 1, for the purpose of this paper the formative time T_B is defined as the interval between the instant the applied voltage first reaches 85% of its full value and the point in time when a marked voltage decrease is discernible. For the circuit used in this experiment, the latter implies currents in the ampere range, so that electric-field distortions must always have been involved in the process of our "breakdown" regardless of its detailed mechanism and the further development of the discharge. We therefore argue that "breakdown" occurs when a certain avalance strength is reached, and we assume for simplicity that this critical amplification is the same for all fields and pressures used. Our criterion is thus analogous to that for midgap breakdown by the streamer mechanism in the absence of a magnetic field:[5]

$$T_B = (\ell n\ N/N_0)/n_g F(E/B), \tag{6}$$

where N is the number of electrons in the gap at time T_B and N_0 is the initial number present.

It is recognized by the authors that the assumption that the critical avalanche amplification is the same for all fields and pressures used is probably not strictly valid. Due to the exponential nature of the avalanche growth, however, it is evident that the formative time is much more sensitive to the ionization rate than it is to the critical avalanche amplifications, as is indicated by Eq. (6). In view of the statistical scatter in the observed formative times (see Fig. 1), we argue that this assumption is reasonable for the work at hand.

From Eq. (6), then, the observed formative time is expected to be inversely proportional to the gas pressure, and otherwise a function of E/B only.

III. APPARATUS

As in the previous work by Bernstein,[1] our experiment was performed with coaxial cylindrical electrodes having a gap much smaller than their radius so that a plane parallel configuration was closely approximated.

A cross-sectional diagram of the electrode structure is shown in Fig. 2, and further details of its construction can be found elsewhere.[6] The copper anode (outer cylindrical shell) was penetrated in one place by an array of small holes which formed a "window" through which ultraviolet light could be directed onto the cathode. This ultraviolet illumination provided a substantial but unknown number of electrons in the electrode gap. A set of interchangeable aluminum cathodes of different diameters allowed the gap spacing d to be varied. Gap spacings of 0.3, 0.5, and 0.8 cm were used. The entire electrode structure was located in a vacuum system between the poles of a large magnet.

The perturbing effects of the electric fringe fields at the ends of the gap were suppressed by covering the ends with thin glass plates coated on the outside with a resistive paint in such a manner as to simulate a gap of infinite length.[6]

The voltage V_g from a dc power supply was applied to the gap abruptly via a thyratron tube and a 100-ohm current-limiting resistor. The range of V_g was 2.4 to 12.0 kV. Taking 5% as an "observable" voltage drop on the oscilloscope traces, the current flowing at the onset of the "voltage collapse" ranged from 1.2 to 6.0 A.

IV. RESULTS

Figures 3 and 4 show two examples of a set of measured formative times plotted against 1/p for given fixed values of the electric and magnetic fields and of the gap spacing. The time $T_c \equiv d/v_E$ required for an electron to drift across the electrode gap in the applied fields is also indicated in these figures. This time is estimated by extrapolation of the drift velocity measurements by Bernstein[1] into the range of v_d of our experiment. Bernstein's results show a nearly constant value for the ratio v_E/v_d when v_d is above about 5×10^6 cm/sec, but his measurements extend to only about $v_d \doteq 9 \times 10^6$ cm/sec. We use the expression

$$v_E = 2.9 \times 10^{10} \frac{pE}{B^2} \text{ cm/sec} \qquad (7)$$

(p is gas pressure in torr at 20°C, E in V/cm, B in G) that one obtains from this constant ratio as an estimate for v_E over our range of v_d (3.3×10^7 cm/sec $< v_d < 6.6 \times 10^7$ cm/sec). When $T_c > T_B$, as in Fig. 3, the inverse pressure dependence predicted by Eq. (6) is indeed observed. When $T_B > T_c$, however, secondary processes play a significant role and a nonlinear dependence on 1/p is usually observed. Figure 4 shows an example of this latter case. In general we have restricted our attention to the case in which $T_c > T_B$, but some data (marked by the letter m), for which $T_B > T_c$, are included in Figs. 5 and 6.

According to Eq. (6) the product $n_g T_B$ should depend only on the ratio E/B and not on E, B, or d separately. Figure 5 shows a plot of $n_g T_B$ against B for a fixed gap spacing and three fixed values of E/B (or v_d). It is seen that some direct effect of B (or E) alone is observed. Figure 6 illustrates that similar deviations, notably at large values of B, occur when the gap spacing is varied. Each of the data bars in Figs. 5 and 6 is obtained from the slope of a straight line drawn through the points of a pressure-dependence graph such as Fig. 3. The length of the data bars represents the uncertainties in determining these slopes. The arrows on the right side of Fig. 5 show predictions obtained in the next section for the magnitude of $n_g T_B$ for the three cases shown.

V. DISCUSSION

A. Ionization Rate

A quantitative prediction of the time taken for the ionization to reach the critical "breakdown" value in our simplified model, according to Eq. (6), requires a knowledge of the magnitude of the growth rate β. Unfortunately, as mentioned before, no direct measurements of the ionization rate under conditions similar to those in our experiment are available. It is, however, possible to make use of extrapolations of published data on the ionization coefficient α in strong crossed fields.[7,8] The value of β then follows directly by multiplication with expression (7), since here, just as in the absence of a magnetic field,

$$\beta \equiv \alpha v_E. \tag{8}$$

The required extrapolation to large values of $\omega_b\tau$ and to the high drift speeds $v_d = cE/B$ of interest in our experiment is best accomplished by means of the equivalent-pressure concept introduced by Blevin and Haydon.[9] These authors showed that when the mean free time τ is independent of electron energy the quantity $p' \equiv p(1 + \omega_b^2\tau^2)^{1/2}$ permits an approximate analytic expression for α in the form

$$\alpha/p' = C_1 \exp(-C_2 p'/E). \tag{9}$$

In the limit of very strong magnetic fields this relation reduces to

$$\alpha/B = C_1' \exp(-C_2' B/E), \tag{10}$$

so that the product $\alpha v_E = \beta$ indeed is of the form (5).

We must not expect Eq. (9) to hold over a wide range of values E/p', of course, since τ in general will not be a constant. On the other hand, the functional form (10) for the limiting case of large $\omega_b\tau$ may be quite accurate again, because it is independent of τ. This last statement merely reflects the fact that in the strong-magnetic-field limit the distribution function for the electrons deviates much less from a drfiting Maxwellian than in the absence of a magnetic field.[3] When the values observed by Bernstein[7] (after proper reduction these values are found to differ very little from those given by Fletcher and Haydon[8]) are inserted in Eq. (9) and use is made of relation (7), we find for hydrogen the quantitative expression

$$\beta/n_g = 2.6 \times 10^{-8}(E/B)\ \exp(-0.78B/E)\ \mathrm{cm}^3/\mathrm{sec}, \qquad (11)$$

which is shown graphically in Fig. 7. It is interesting to note that this curve differs by less than a factor of 2 from the results of a numerical computation,[10] based on the theoretical considerations treated in Ref. 3, which is also displayed on Fig. 7 for comparison. It should be noted, however, that the computation had to make use of several poorly known cross sections.

B. Secondary Electrons from the Cathode

As stated in the Introduction, we believe that in this experiment the release of secondary electrons at the cathode can contribute in only a minor way to the charge accumulation in the gap. There are two very good reasons for this assumption: (a) In a strong transverse magnetic field--i.e., for large values of $\omega_b\tau$--most secondaries are returned promptly to the cathode, so that the effective value of the coefficient for release of secondaries, γ, is reduced roughly by the ratio $\omega_b\tau$. (b) Quite in general, in midgap breakdown, which must involve significant space-charge development, the secondary avalanches find themselves in regions of reduced electric field, so that their growth rate is always less than that of the first generation. This is not to say, of course, that secondary electrons can be completely absent, because in that case there would be no possibility of multiple-avalanche breakdown, contrary to observations. In fact, the absence of a sharp jump in formative time when it is longer than the electron crossing time can be explained only by the existence of secondary avalanches. The very smooth transition from single- to multiple-avalanche breakdown times is, however, most probably caused by a smearing-out of the crossing time itself. The latter must be expected, because the field distortion caused by the space charges are undoubtedly not completely uniform around the anode perimeter.

The only point we need to make is that in midgap breakdown the contributions from secondary avalanches can, in a first approximation, be neglected. It is difficult to see how otherwise the formative time could be strictly proportional to the inverse of the gas pressure.

C. Magnitude of the Critical Value of N/N_0

Prediction of the time T_B also requires knowledge of the numerical value of the critical ionization gain N/N_0 that is postulated in Eq. (6). Fortunately, an estimate only is sufficient because of the logarithmic dependence. As a first step we may argue, in analogy to the treatment of midgap breakdown at high gas density,[5] that the breakdown goes to rapid completion (see Fig. 1a) as soon as the accumulated space charge has significantly enhanced the electric field on the anode side of the electron cloud. In most of our cases this leads to a value for N in the order of

10^{12} provided the charge is spread in a reasonably uniform manner over a cylindrical shell within the gap. But even if the electric field distortion is not leading to a significantly accelerated ionization growth, the current in the gap always produces an observable voltage drop across our limiting resistor when N lies between 10^{12} and 10^{13}. Naturally this would lead to a more gradual voltage collapse, and presumably Fig. 1b is an example of such a case. Quite in general, the longer formative times tend to be correlated with this second type of "breakdown."

Unfortunately, we do not have equally good arguments for our estimate of N_0 because the resistive coating of the insulators precluded measurement of the photoemission current. That some statistical scatter of formative times remained (see Fig. 1) when strong cathode illumination was added seems to indicate that N_0 must be quite small. Presumably the net emission of photoelectrons is very much reduced by the strong magnetic field, exactly as is the release of secondaries, because at large $\omega_b\tau$ most electrons are returned to the cathode. Thus we make the guess that N_0 is larger than unity but certainly much less than 100. This means that $25 \lesssim \ln N/N_0 \lesssim 27$. The arrows on the right of Fig. 5 correspond to the upper limit of this range. The fair agreement between observed and predicted formative times indicates that our simple model may be rather good and that our extrapolation for the ionization rate β is probably quite adequate. The remaining discrepancies are readily explained, at least in a qualitative way, as resulting from our neglect of electron losses.

D. Electron Losses

The most serious oversimplification of the breakdown model discussed so far is the assumption that no electrons are lost from the system. In reality some electrons must be expected to reach one of the bounding surfaces during the buildup process. Because of this removal the net growth rate will always be smaller than $\bar{\beta}$, so that the formative time is underestimated in our calculations. A quantitative treatment of such refinement becomes quite complex[6] and is therefore not attempted here. Instead, we present the major physical arguments in order to explain at least qualitatively the observed deviations from our simplified model.

For two separate reasons some electrons can reach the anode in a time much shorter than previously assumed. Firstly, avalanches that start out in midgap of course arrive at the anode in less than a normal crossing time. Thus if, for instance, the initial electrons are distributed evenly throughout the gap the effective value for N_0 is proportional to $(d - v_E T_B)$. Although it is in the right direction, this effect is not likely to explain the dependence of $n_g T_B$ on the gap spacing d shown in Fig. 6 because it enters only logarithmically and because we expect most initial electrons to start very close to the illuminated cathode. The other possible, and probably much more significant, complication leading to "premature" electron removal at the anode involves radial drifts caused by azimuthal components of the electric field, which appear whenever the space charge has azimuthal nonuniformities. Even if they were initially absent such nonuniformities are in fact expected to develop in the later stages of the buildup as a form of the so-called "diocotron instability" of space-charge streams.[11] The subsequent behavior is equivalent to that found in a local reduction of the electrode spacing, which reduces the transit time for the participating electrons but does not lead to breakdown since only a fraction of the avalanche is involved. The final consequence is a partial removal of electrons at the anode and thus a stretching of the formative time. The effect is largest for the shortest electrode spacings, and thus might well explain the results shown in Fig. 6, but the B dependence seen in Fig. 5 is probably caused by a different loss process.

It is clear that in this experiment electrons diffuse freely along the magnetic lines, and in the later stages of the growth process they are also forced to flow in this direction by the space-charge repulsion. In fact, simple estimates indicate that in all our cases the ionization must be rather well distributed in the axial direction. It follows that electrons are deposited on the end insulators, building up a surface charge until the potential there is sufficient to suppress any further axial motion. The resulting electron-loss rate can be shown to be quite substantial, and there are two good reasons why it should be decreasing with increasing magnetic field. The area of insulator surface that is swept out by the progress of an electron avalanche is larger for lower magnetic fields. But more important is perhaps that at larger magnetic fields the flux lines in our magnet acquire a slight curvature which is concave towards the cathode. This small curvature is insufficient to affect the electron energy in any noticeable way, but it has a marked effect on the axial diffusion, particularly in the early phases of the buildup. Such a variation could easily explain the observed B (or E) dependence of $n_g T_B$ shown in Fig. 5.

Unfortunately, none of the explanations advanced in this section has yet been verified by special tests. We therefore consider our interpretation of the deviations from the simple model as merely tentative at this point. The gross agreement with the general predictions, however, is felt to be good enough to inspire confidence in the validity of our theory.

VI. CONCLUSION

Reproducible values for the formative time of breakdown in hydrogen across a strong magnetic field have been measured for conditions under which this time is less than the time required for an electron to drift across the electrode gap in the applied fields. These measurements are compared with a simplified theory which neglects electron losses from the avalanche and postulates that breakdown occurs when a certain critical space charge or a certain critical current is reached. The pressure dependence predicted by this theory is confirmed over a range of pressures spanning a decade, but deviations from the predicted functional dependence on E/B are found. These deviations are tentatively ascribed to electron losses along the magnetic field. Predictions for the magnitude of the formative time based on extrapolated values of ionization rates and drift speeds are found to be in reasonable agreement with the measured values. This agreement may be taken as a partial confirmation of our theoretical model.

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FIGURE LEGENDS

Fig. 1 Typical oscilloscope traces of the potential across the electrodes as a function of time. (a) 0.8 cm gap; 1.0 torr; 7.2 kV; 18 kG (four trace overlay). (b) 0.8 cm gap; 1.0 torr; 3.6 kV; 18 kG (five trace overlay).

Fig. 2 Schematic drawing of a cross section of the electrode structure. All parts have cylindrical symmetry about the center.

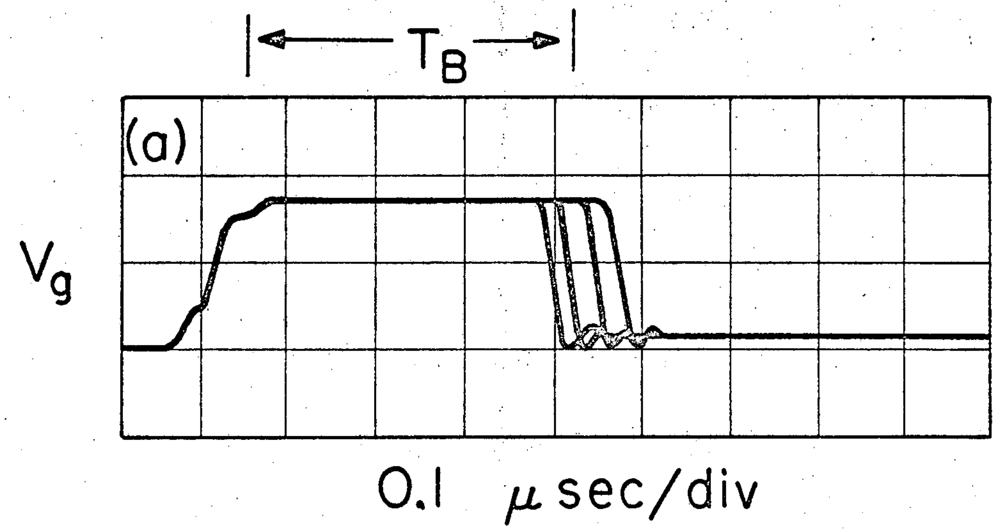
Fig. 3 The formative time as a function of 1/p for $T_B < T_c$. Each point represents a single shot.

Fig. 4 The formative time as a function of 1/p for $T_B > T_c$. The crosses represent single shots, and each bar represents a group of several shots.

Fig. 5 The product of the measured formative time and the gas density plotted against the magnetic field strength for the 0.8 cm gap. The data labeled with the letter m are cases for which $T_B > T_c$.

Fig. 6 The product of the measured formative time and the gas density for $v_d = 5.0 \times 10^7$ cm/sec. The data labeled with the letter m are cases for which $T_B > T_c$.

Fig. 7 β/n_g as a function of v_d. Solid curve based on extrapolation of Bernstein's data (Ref. 7); broken line computed by Pearson and Kunkel (Ref. 10).



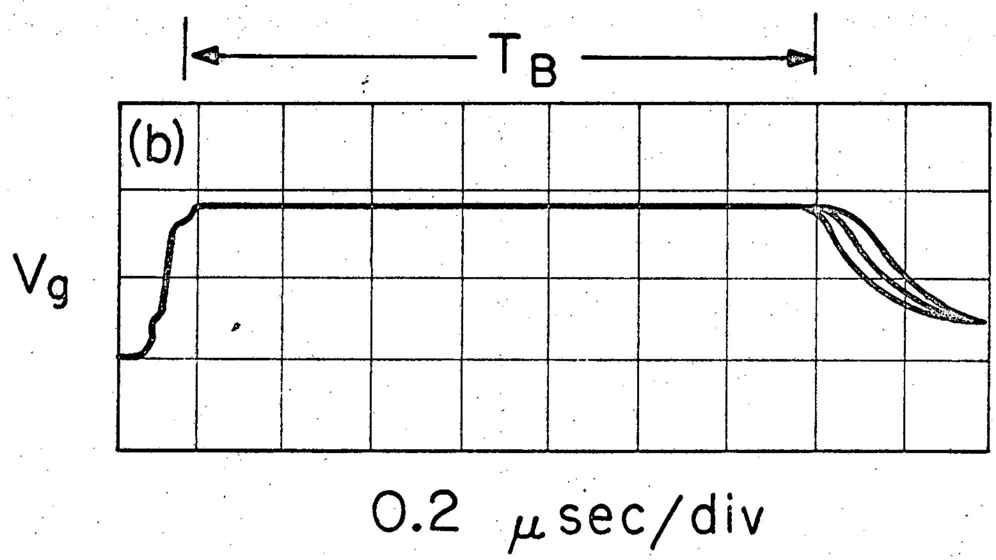
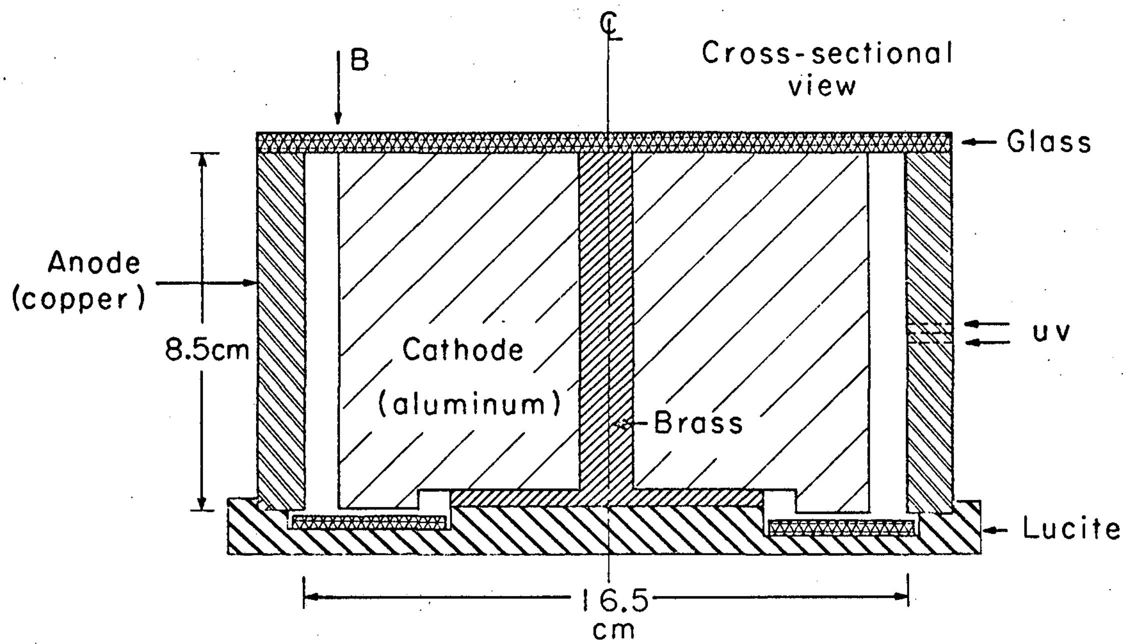


Fig. 1

XBL673-2408



Electrode structure

Fig. 2

MUB-12992

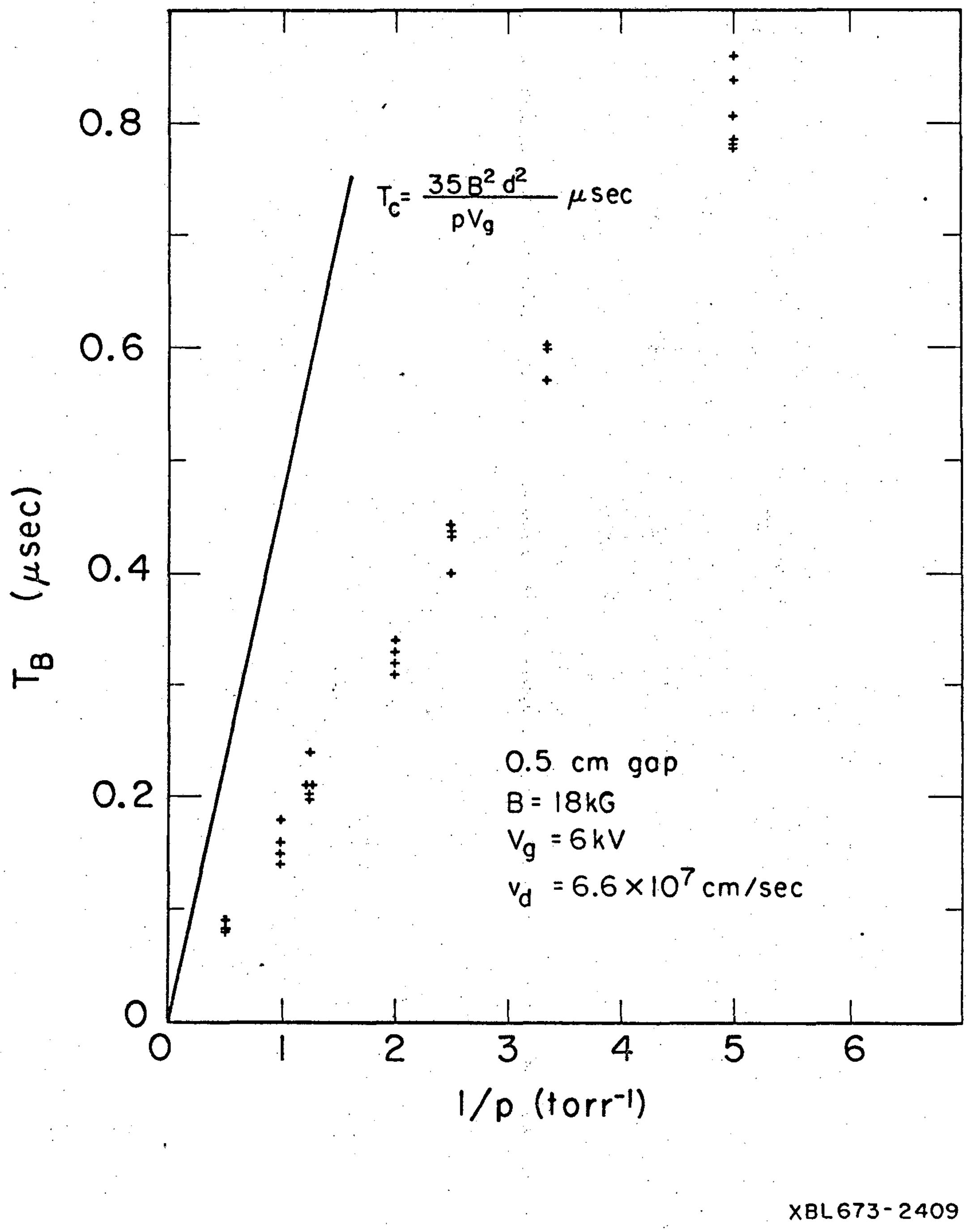

0.8
0.6
0.4
0.2
0
T_B (μsec)
$T_c = \frac{35 B^2 d^2}{p V_g}$ μsec
0.5 cm gap
B = 18kG
V_g = 6 kV
v_d = 6.6×10^7 cm/sec
0 1 2 3 4 5 6
1/p (torr^{-1})
XBL673-2409

Fig. 3

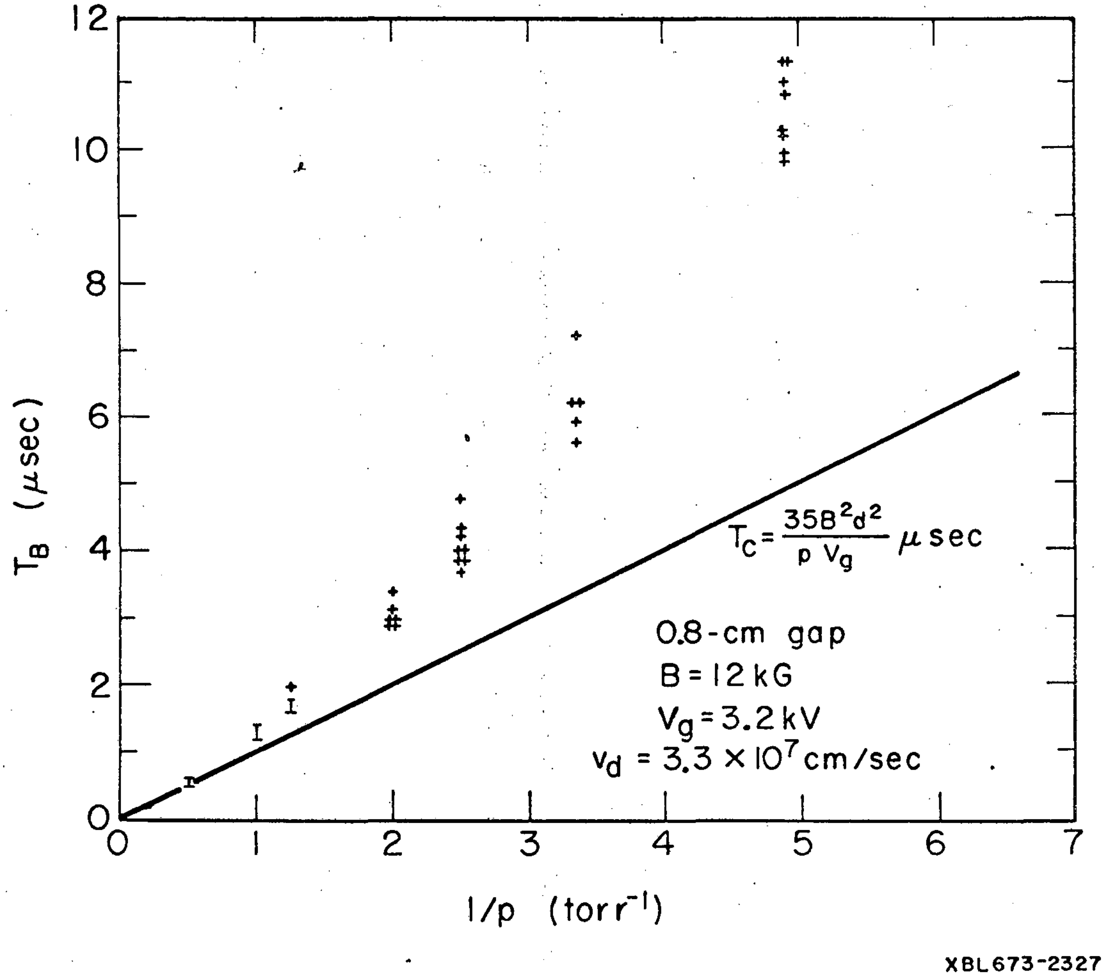
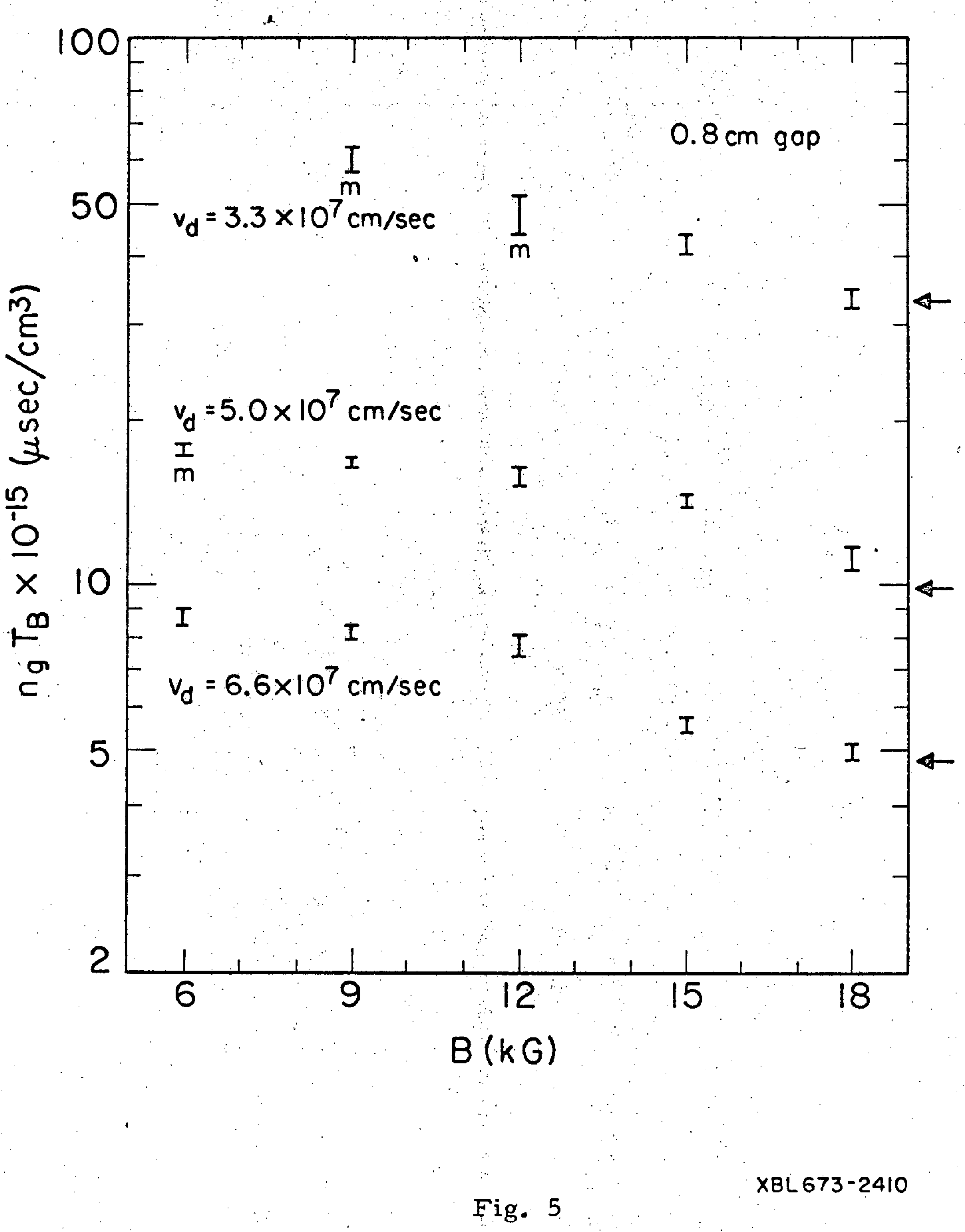
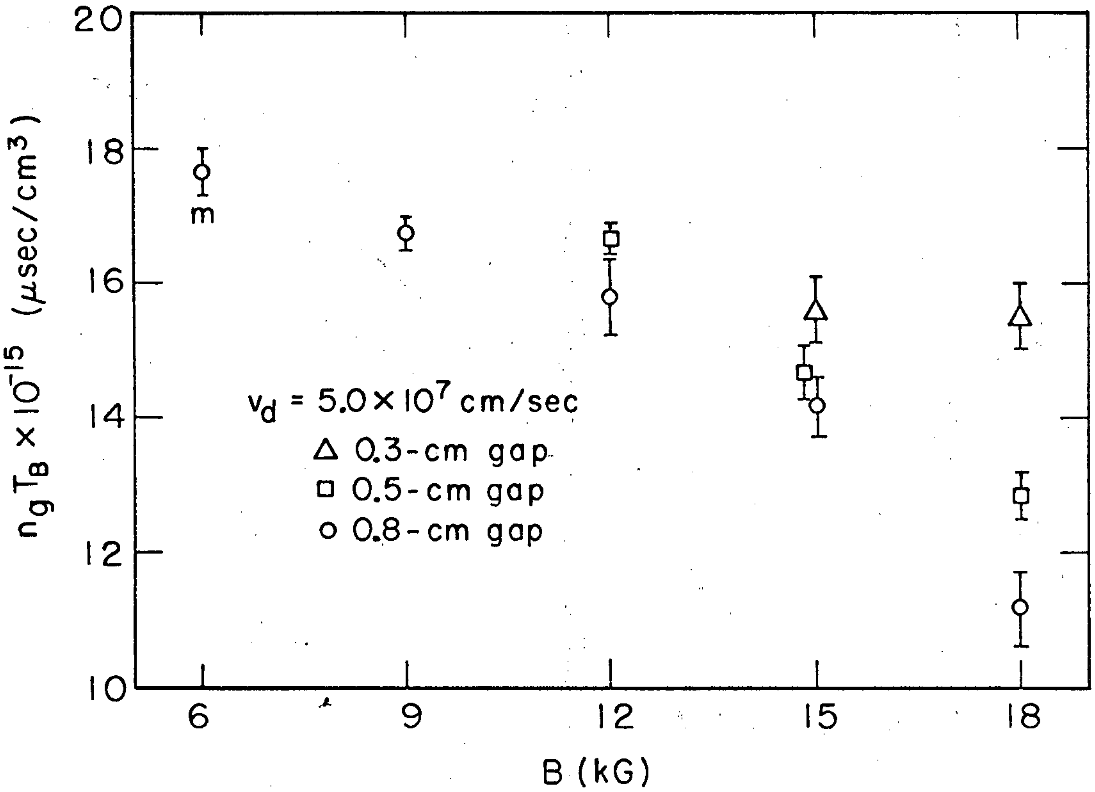

12
10
8
6
4
2
0
T_B (μsec)
0
1
2
3
4
5
6
7
1/p (torr^{-1})
$T_C = \frac{35B^2d^2}{p\,V_g}\ \mu$sec
0.8-cm gap
B = 12 kG
V_g = 3.2 kV
v_d = 3.3 × 10^7 cm/sec
XBL673-2327

Fig. 4



100
50
10
5
2
6
9
12
15
18
B (kG)
$n_g T_B \times 10^{-15}$ (μsec/cm^3)
0.8 cm gap
$v_d = 3.3 \times 10^7$ cm/sec
$v_d = 5.0 \times 10^7$ cm/sec
$v_d = 6.6 \times 10^7$ cm/sec
XBL 673-2410

Fig. 5



XBL673-2329

Fig. 6

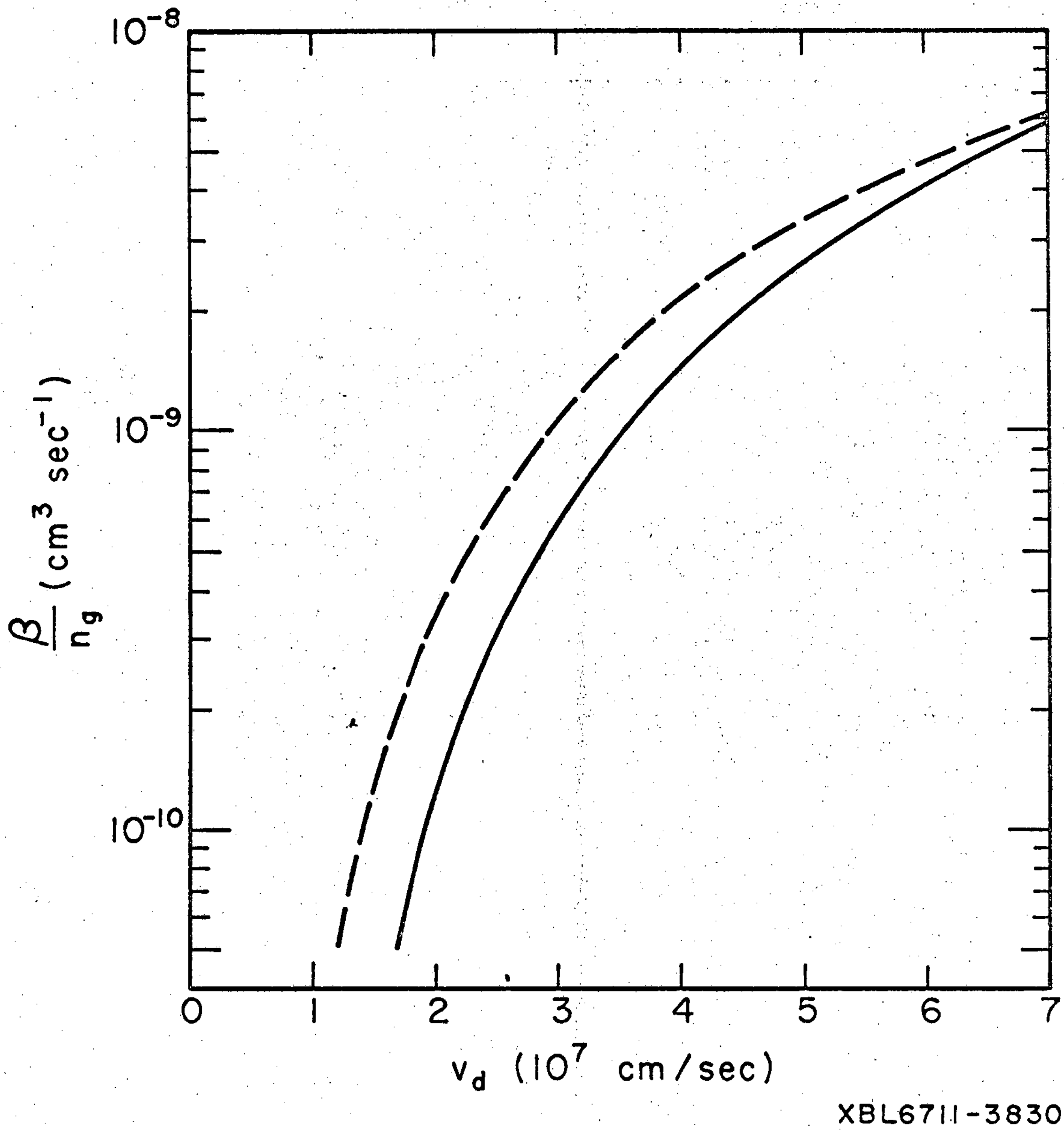

10^{-8}
10^{-9}
10^{-10}
$\frac{\beta}{n_g}$ (cm^3 sec^{-1})
0
1
2
3
4
5
6
7
v_d (10^7 cm/sec)
XBL6711-3830

Fig. 7