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Laser beam induced breakdown in helium and argon

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Abstract. Measurements of the laser radiation power density $\hat{P}_{\text{th}}$ required to cause ionization growth and breakdown in the pressure ranges $1800 \leqslant p \leqslant 16\,000$ torr and $200 \leqslant p \leqslant 20\,000$ torr have been made in helium and argon, respectively. Three Q-switched lasers (ruby, Nd, and dye) were used to investigate the dependence of $\hat{P}_{\text{th}}$ upon pressure, radiation frequency ω and flash duration 2τ. Values of $\hat{P}_{\text{th}}$ lie between 10^{11} W cm^{-2} at low pressures and 10^{10} W cm^{-2} at high pressures.

The experimental results are compared with expressions based upon an extrapolation of classical microwave breakdown theory to optical frequencies. Reasonably good agreement between measured and calculated threshold intensities is obtained.

The interaction between the laser radiation and the gas proceeds via inelastic electron–atom collisions and photoionization. Three pressure regimes may be specified: a low pressure regime in which the threshold intensity and ionization growth rate are governed principally by diffusion losses, an intermediate régime in which the laser flash duration is the dominant factor and a high pressure régime where electron–ion recombination controls the rate of ionization growth and the onset of breakdown. The ionization rate is shown to be proportional to the laser beam intensity and, in the intermediate régime, the threshold intensity $\hat{P}_{\text{th}}$ is proportional to ω^2/τ at constant pressure. A value for the electron–atom momentum transfer collision frequency of about $3\text{–}4 \times 10^9 p$ and of about 5 eV for the electron mean energy is derived for argon at the onset of breakdown.

1. Introduction

The non-linear interaction of powerful focused laser radiation with gases leading to ionization growth and plasma production has aroused much interest. Gases which are normally transparent to long wavelength radiation of low intensity may be rapidly transformed into opaque highly conducting plasmas in times of the order of nanoseconds by intense laser radiation of the same wavelength.

This transformation takes place in four stages: initiation, growth, plasma development and, finally, extinction. The initiatory stage is the time which elapses between the arrival of the laser radiation flash in the lens focal region and the initiation, by the release of an electron–ion pair, of the growth in the free electron and ion concentration in the gas. The initiatory mechanism appears to be due to electron liberation from an atom as a result of multiphoton resonance absorption of the laser light (Tozer 1965, Bebb and Gold 1966, Baravian *et al.* 1970). It occurs at a very early stage in the flash.

The formative growth stage to breakdown is the subsequent period of amplification in the number of ion pairs until the state of 'breakdown' is reached. Breakdown is arbitrarily defined as the attainment of a certain electron concentration (greater than or equal to 10^{15} cm^{-3}) or degree of fractional ionization $\delta \sim 0{\cdot}1\%$ of gas atoms in the focal region which is sufficient say to cause significant absorption and scattering of the laser radiation

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(Phelps 1966, Young and Hercher 1967). The combined duration of the initiatory and formative growth times is exceedingly small and may be only a few nanoseconds or less.

Once the conditions for the onset of breakdown are satisfied, there follows, if the laser intensity continues to increase, a rapid plasma development stage and the formation of highly conducting hot expanding gas in which further absorption and hydrodynamic effects become important (Daiber and Thompson 1967, Lampis and Brown 1968).

The final or extinction phase lasts for a time which may be three or four orders of magnitude larger than the duration of the laser flash, i.e. of the order of 50 μs compared with about 20 ns. During this phase the plasma gradually dies away as a result of radiation emission, diffusion, attachment and recombination. Brief surveys and bibliographies of some aspects of the physical processes involved in these various stages have recently been given by Morgan (1967), De Michelis (1969) and Veyrié (1970).

The present paper is concerned with the growth stage, i.e. with a study of the processes responsible for the amplification of ionization and with the conditions which must be satisfied for laser-induced breakdown to occur. In particular it describes a theoretical and experimental study of laser-induced breakdown in helium and argon over a wide pressure range, using three types of Q-switched lasers to provide radiation of various wavelengths and flashes of various durations.

2. Theory

Precise measurement of the intensity of laser radiation required to create ionization and breakdown is difficult to make. The difficulties arise on account of imprecise definition of the focal region, inaccurate knowledge of the spatial and temporal variation of the radiation intensity within the focal region, uncertainties in the absolute value of the instantaneous radiation power density and temporal changes in the mode structure and in the consequent variation of the angular divergence of the light beam. These uncertainties also give rise to theoretical difficulties, because it is not possible to specify exactly the boundary conditions which must be used to evaluate the spatio-temporal growth of ionization from the solution of the continuity equation describing the interaction of the laser light with electrons, ions and atoms. They also make inter-comparison of data published by various workers exceedingly difficult to interpret.

Nevertheless, despite these difficulties, considerable progress has been made and there is now a large body of evidence in support of the view that the growth stage proceeds as a result of the cascade or avalanche mechanism of inverse bremsstrahlung absorption (Wright 1964), which is physically equivalent to classical excitation and ionization by inelastic collisions between atoms and electrons in which the interaction energy is drawn from the electromagnetic field of the light wave (Browne 1965). This is essentially an extrapolation of microwave discharge concepts to optical frequencies.

Thus, once liberated in the focal region during the early part of the laser flash by multi-photon absorption, free electrons gain energy from the electromagnetic field in colliding with atoms and, when their energies exceed the atomic excitation and ionization potentials, cause an amplification in the electron and ion concentrations. Simultaneously there is a loss of free electrons from the focal region owing to diffusion, recombination and attachment to electronegative atoms. If the rate of production exceeds the rate of loss the concentration grows and breakdown may occur.

The net rate of change in electron concentration is described by the following continuity equation (Morgan 1965)

$$\frac{\partial n}{\partial t} = n\nu + D\nabla^2 n - \operatorname{div}(nW) - Rn^2 - \beta nN \tag{1}$$

in which n is the electron concentration at a point in the focal volume at a time t after the release of the initiatory electron, ν is the total ionization collision frequency which, in the case of laser-induced breakdown, may be the sum of the ionization frequency ν_i due to single impact collisions between electrons and atoms and of the photoionization frequency

ν_{ex} due to electronic excitation to the first excited state followed by rapid photoionization (Phelps 1965, Zeldovich and Raizer 1965). In the noble gases, resonance radiation may be confined to the focal volume for relatively long times so that photoionization of the effectively long-lived excited atoms is very probable. An estimate based on the Holstein (1947) and Biberman (1947) theory of radiation trapping shows, for example, that in helium at atmospheric pressure resonance ultraviolet radiation may be trapped in the focal volume for about 50 μs—very much longer than the laser flash duration. There is thus little effective removal of highly excited atoms by radiation losses from the focal region.

D, R, β and W are the electron diffusion, recombination and attachment coefficients and drift velocity, respectively. N is the neutral atom concentration and, for a gas at pressure p, is given by $N = N_0 p = 3{\cdot}56 \times 10^{16} p$.

In order to solve equation (1) we make assumptions and simplifications based upon the following considerations:

(*a*) Initiatory electrons are released while the radiation intensity is still relatively small.

(*b*) During the earliest stages of ionization growth there is no net loss of electrons owing to drift. This is so since the electromagnetic field is oscillatory and induces no unidirectional motion. Any anisotropic space-charge field will tend to confine the electrons in the focal volume. We thus set $W = 0$.

(*c*) We confine the discussion to non-attaching gases such as used in the present work and thus set $\beta = 0$.

(*d*) Our analysis (Evans and Morgan 1969 a, b) of the spatial distribution of the radiation intensity, produced by a single-mode laser in the focal region of simple lenses commonly used in laser breakdown studies has shown that the focal region consists of a number of closely spaced cylindrical regions of high intensity lying symmetrically about the optic axis. We have shown that the length l of the cylinders greatly exceeds their diameter $2a$. Consequently we may assume that there is negligible electron loss by diffusion through the cylinder ends compared with the lateral diffusion loss. The dimensions of the cylinder depend upon the radiation wavelength used and upon the lens focal length and aberration function as well as beam divergence.

(*e*) We assume that the intensity is uniform inside the cylinders and zero outside. This cannot be justified rigorously and is made in order to make the following analysis tractable. It implies constancy of the ionization frequency inside the cylinder. In addition the diffusion coefficient is taken as constant throughout this volume.

We may specify two pressure régimes—one in which the gas pressure is so low that electron–ion recombination is unlikely to occur, so that the dominant loss mechanism is by lateral diffusion, and one in which the pressure is so high that recombination losses are far in excess of diffusion losses. The extent of these régimes depends of course upon the values of D and R of the particular gases used.

3. Diffusion dominated losses

For the low pressure régime, equation (1) may be written in cylindrical coordinates as

$$\frac{\partial n(r,t)}{\partial t} = n\nu(t) + D\,\frac{\partial^2 n}{\partial r^2} + \frac{D}{r}\,\frac{\partial n}{\partial r}.$$

Its solution, for a cylinder of radius a, is

$$n(r,t) = A J_0\left(\frac{2{\cdot}405 r}{a}\right) \exp\left[\int_0^t \left\{\nu(t) - D\left(\frac{2{\cdot}405}{a}\right)^2\right\} \mathrm{d}t\right] \tag{2}$$

where A is a constant governed by the initiatory electron distribution in the focal volume V. We take this as equal to the reciprocal of V, i.e. $A = 1/V = 1/\pi a^2 l$ and imply that the laser radiation releases one electron therein early during the flash, i.e. $Vn(0, 0) = 1$.

In order to proceed further, it is necessary to relate $\nu(t)$ to the radiation intensity. We do this using the well-known concept of the static 'effective field' which produces the same energy transfer as the oscillatory field of the laser beam (Brown 1959, 1966, McDonald 1966). The effective field is related to the root mean square laser field E_L by $E_e = E_L\{\nu_m^2/(\nu_m^2+\omega^2)\}^{1/2}$, where ω is the angular frequency of the radiation and ν_m is the electron–atom collision frequency for momentum transfer which is assumed to be independent of electron energy. This is so in helium and neon for electron energies between 3 and 36 eV. When $\omega \gg \nu_m$, which is the case in the present work, we may write $E_e = E_L\nu_m/\omega$. The values deduced from this simplified form differ by less than 1% from those obtained using the full formula for pressures up to 10^5 torr in helium and neon.

Provided the electrons very rapidly acquire the same energy distribution (Afanas'ev *et al.* 1970) under the optical frequency field as they would have under the equivalent static field E_e, we may relate ν_i and ν_{ex} to the Townsend primary ionization coefficient α and excitation coefficient θ. Over ranges of interest θ/N, α/N and W are proportional to E_e/N so

$$\frac{\nu}{N} = \frac{\nu_i + \nu_{ex}}{N} = \frac{W(\alpha+\theta)}{N} \simeq k\left(\frac{E_e}{N}\right)^2. \tag{3}$$

Here k is a coefficient for a given gas which may be derived from a knowledge of values of θ, α and W. Figure 1 illustrates the dependence of ν/N upon E_e/N deduced from data

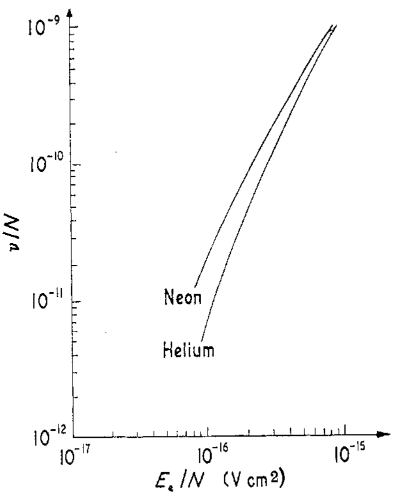

Figure 1. ν/N against E_e/N for helium and neon.

on θ/N, α/N and W obtained in our laboratories (Davies *et al.* 1962, Hughes 1967, Willis and Morgan 1968). It shows that, to a first order, equation (3) is satisfied, the graphs having a gradient of approximately two.

Table 1 gives corresponding ranges of values of the coefficient k. Those for argon are based on energy balance considerations since no data on θ/N is available for this gas, and upon the drift velocity data collected by Caren (1963).

Table 1

Gas	k ($cm^{-1}\ s^{-1}\ V^{-2}$)
Helium	7×10^{20} to $1{\cdot}3\times10^{21}$
Neon	1×10^{21} to 2×10^{21}
Argon	7×10^{20} to 1×10^{21}

Using Poynting's theorem we may relate the effective ionization rate ν/N to the laser beam intensity $P(t)$. Thus using equation (3), we find

$$\frac{\nu}{N}=\left\{\frac{377k}{\omega^2}\left(\frac{\nu_{\mathrm{m}}}{N}\right)^2\right\}P(t) \tag{4}$$

where $P(t)$ is expressed in W cm^{-2}.

In order to derive an expression for the total number of electrons N_{e} liberated during the laser flash, we note that in experiments the temporal variation of the laser beam intensity is approximately triangular, with a peak intensity $\hat{P}$ and duration 2τ, i.e.

$$\begin{aligned} P(t)&=\hat{P}t/\tau \qquad &\text{for} \quad 0\leqslant t\leqslant \tau \\ &=\hat{P}(2-t/\tau) \qquad &\text{for} \quad \tau\leqslant t\leqslant 2\tau. \end{aligned}$$

Integration of equation (2) with respect to time yields

$$N_{\mathrm{e}}(2\tau)=0{\cdot}434 \exp\left[\tau\left\{\frac{377k}{N}\left(\frac{\nu_{\mathrm{m}}}{\omega}\right)^2\hat{P}-2D\left(\frac{2{\cdot}405}{a}\right)^2\right\}\right].$$

In accordance with our definition of breakdown we let $N_{\mathrm{e}}(2\tau)=\delta NV$ and so derive an expression for the beam intensity $\hat{P}_{\mathrm{th}}$ which yields the fractional degree of ionization δ in the form

$$\hat{P}_{\mathrm{th}}=\frac{pN_0}{377k}\left(\frac{\omega}{\nu_{\mathrm{m}}}\right)^2\left\{\frac{1}{\tau}\ln\left(2{\cdot}3\delta pN_0V\right)+2D\left(\frac{2{\cdot}405}{a}\right)^2\right\}. \tag{5}$$

This is the threshold laser beam intensity required to cause breakdown at low gas pressure when diffusion losses dominate.

4. Recombination dominated losses

When free electron removal from the focal region is principally by recombination with positive ions, the continuity equation (1) becomes

$$\frac{\partial n}{\partial t}=n\nu-Rn^2$$

so that

$$n(t)=n_0 \exp\left(\int_0^t \nu(t)\,\mathrm{d}t-R\int_0 n(t)\,\mathrm{d}t\right). \tag{6}$$

In order to solve this non-linear integral equation to a first order of approximation, we assume that the value of $n(t)$ is that which obtains in the absence of recombination, i.e.

$$n(t)|_{R=0}=\frac{1}{V}\exp\int_0^t \nu(t)\,\mathrm{d}t$$

so that, on insertion in equation (6), we obtain

$$n(t)=\frac{1}{V}\exp\left\{\int_0^t \nu(t)\,\mathrm{d}t-\frac{R}{V}\int_0^t \exp\left(\int_0 \nu(t)\,\mathrm{d}t\right)\right\}. \tag{7}$$

For a triangular flash of duration 2τ and peak intensity $\hat{P}$, to which corresponds a maximum ionization collision frequency $\hat{\nu}=(377k\hat{P}/N)\,(\nu_{\mathrm{m}}/\omega)^2$, integration of equation (7) yields the electron concentration at the end of the flash as

$$n(2\tau)=\frac{1}{V}\exp\left\{\hat{\nu}\tau-\frac{R}{V}\left(\frac{\pi\tau}{2\hat{\nu}}\right)^{1/2}\exp\left(\hat{\nu}\tau\right)\operatorname{erf}\left(2\hat{\nu}\tau\right)^{1/2}\right\}. \tag{8}$$

For power densities in excess of about 10^{10} W cm^{-2}, as used in the present study, erf $(2\hat{\nu}\tau)^{1/2}$ may be taken as unity. Even for power densities as low as 10^{9} W cm^{-2}, it is within 4% of unity. Consequently, the number of electrons in the focal volume is given to a good approximation by

$$N_e(2\tau) = Vn(2\tau) \simeq \exp(\hat{\nu}\tau) \exp\left\{-\frac{R}{V}\left(\frac{\pi\tau}{2\hat{\nu}}\right)^{1/2} \exp \hat{\nu}\tau\right\}. \tag{9}$$

In order to derive an expression for the threshold power density P_{th} we first set $N_e(2\tau) = \delta N_0 p V$, i.e. define breakdown as the ionization of a fraction δ of the atoms present in V, and, once more repeat the assumption that this electron number is approximately equal to what it would be in the absence of recombination, i.e. $\simeq \exp(\hat{\nu}\tau)$.

It follows that

$$\hat{P}_{th} = \frac{pN_0}{377k}\left(\frac{\omega}{\nu_m}\right)^2 \left\{\frac{1}{\tau} \ln(\delta N_0 p V) + R\delta N_0 p \frac{\omega}{\nu_m}\left(\frac{\pi p N_0}{754k\hat{P}_{th}\tau}\right)^{1/2}\right\}. \tag{10}$$

Equations (5) and (10) for the breakdown intensity in the diffusion controlled and recombination controlled cases, respectively, are analytically similar to expressions derived by Young and Hercher, but differ in the following important respects: these disclose the sensitive dependence of $\hat{P}_{th}$ upon the radiation frequency ω and they do not require a knowledge of the inverse bremsstrahlung cross sections which are necessary in applications of their expressions.

5. Flash duration limited breakdown

Between these limiting cases of very low and very high pressures, we recognize an intermediate régime in which the duration τ of the laser flash is the dominant factor in governing the onset of breakdown. Thus when the second terms, i.e. those involving D or R inside the brackets in equations (5) and (10) are small compared with the term involving τ, an approximate expression for $\hat{P}_{th}$ is

$$\hat{P}_{th} \simeq \frac{pN_0}{377k}\left(\frac{\omega}{\nu_m}\right)^2 \left(\frac{1}{\tau} \ln(\delta N_0 p V)\right). \tag{11}$$

Thus, at a constant pressure, when neither recombination nor diffusion losses occur markedly, the breakdown intensity is expected to be roughly proportional to ω^2/τ. This expression is useful in experimental tests of the theory as it is not easy to vary ω over a wide range independently of τ. If correct it provides a method of estimating either k or ν_m at breakdown.

Inspection of equations (5), (10) and (11) shows that the slopes of graphs of $\ln \hat{P}_{th}$ against $\ln p$ are expected to be -2, -1 and $-\frac{1}{3}$ in the very low (diffusion limited), intermediate (flash duration limited), and very high (recombination limited) pressure régimes. These conclusions are in broad agreement with published data (Young and Hercher 1967).

6. Apparatus and experimental procedures

The breakdown threshold intensity of argon lies well below that of helium. Consequently measurements could be made in argon over a wider pressure range than for helium without danger of irreversible damage to the laser rods. For this reason it was necessary to use two ionization chambers. The first, made of stainless steel, was evacuated to 10^{-9} torr and slowly filled with argon via vapour traps immersed in liquid nitrogen to a predetermined pressure between 200 and 1200 torr. The pressure was measured by means of a bellows micromanometer which separated the gas from a mercury manometer.

For experiments in the range 2000 to 20 000 torr a small aluminium chamber was used. The gases, admitted after initial evacuation and repeated flushing, were dried by passage through liquid nitrogen vapour traps. In both chambers the incident laser radiation was focused by simple bi- or plano-convex Spectrasil-'A' lenses having focal lengths of 3 and 5 cm. The use of plano-convex lenses prevented ablation of the rear surface at high beam powers. Breakdown could be observed visually and recorded by means of a Mullard type 56 TVP photomultiplier fitted with a narrow band (16 Å) interference filter centred at the 4158 Å line of argon I. Both chambers were fitted with conical light dumps to absorb transmitted laser light.

7. Lasers

Three lasers were used: a plane polarized ruby laser Q-switched by means of a cryptocyanine in isopropanol solution, a plane polarized neodymium-in-glass laser, Q-switched by means of a rotating totally internal reflecting Spectrasil-'A' prism and a dye laser which was optically pumped by means of the ruby laser. The dye used was 3,3′-diethylthiatricarbocyanine iodide in methanol. The dye concentration was adjusted to obtain a maximum power output of about 10 MW with a conversion efficiency of the order of 15%. The output spectrum was measured using a monochromator and was centred at 8010 Å with a half-width of 50 Å. The flash durations were 40 ns for the ruby and dye lasers and 80 ns for the Nd laser.

In order to minimize errors due to variation in beam divergence the laser pump power was kept constant and the intensity of the beam entering the ionization chamber was smoothly adjusted by means of a set of rotatable thin glass plates set at the Brewster angle (Evans and Morgan 1968) for the ruby and Nd lasers. The output of the non-polarized dye laser was adjusted by inserting neutral density filters in the beam. In addition an infrared filter was used to eliminate stray ruby laser light from the output beam of the dye laser. The beam power was measured using TRG type 105B photodiodes.

8. Beam divergence

The beam divergence was estimated at the beginning and end of each series of measurements with a particular laser. For this purpose the attenuated beam was focused using a 118 cm focal length lens onto a photographic plate or film. A single flash was recorded and then the film displaced so that, without altering any parameter except the insertion of a 50% absorbing neutral density filter, a second flash, having half the intensity of the first, could be recorded alongside, while simultaneous measurements of beam power were made.

The intensity distributions of both spots on the developed plates were measured using a microdensitometer and the beam half-width points and divergence deduced. This method has the advantage of requiring no calibration of the film response. The beam divergences obtained by this method were: ruby laser, 3·5 mrad; dye laser, 4 mrad; Nd laser, 4 mrad.

9. Results and conclusions

Figure 2 shows the measured values of $\hat{P}_{\text{th}}$ as a function of pressure for helium at 6943 Å and for argon at 6943 Å, 8010 Å and 10 600 Å covering the pressure range 200 to 2×10^4 torr.

9.1. *Helium*

In helium $\nu_{\text{m}} = 2{\cdot}4 \times 10^9 p$ and is largely independent of electron energy over a very wide range (MacDonald 1966). The recombination coefficient R of helium is very much smaller than for other noble gases (Ferguson *et al.* 1965). Values lie between $R < 4 \times 10^{-9}$ cm^3 s^{-1} at 300 K (Oskam and Mittelstadt 1963) and $R < 2 \times 10^{-11}$ cm^3 s^{-1} at 1800 K (Collins and Robertson 1965). It is not surprising therefore that there is no evidence for an asymptotic approach to a line of slope $(-\frac{1}{3})$ in the measured values of $\hat{P}_{\text{th}}$ at high pressure. Values of $\hat{P}_{\text{th}}$ calculated using equation (5) and known values of D (Brown 1959), k (table 1) and ω give the straight line shown in figure 3. Compared with the experimental points the agreement is good and the gradient is -1, as required by the analysis. We conclude that the analysis describes the mechanism of laser-induced ionization growth to breakdown in

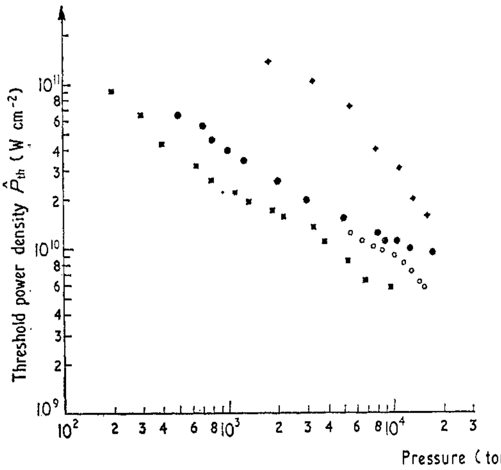

Figure 2. Measured threshold intensities in helium for $\lambda = 6943$ Å (♦) and argon for $\lambda = 6943$ Å (●), 8010 Å (○) and 10 600 Å (■).

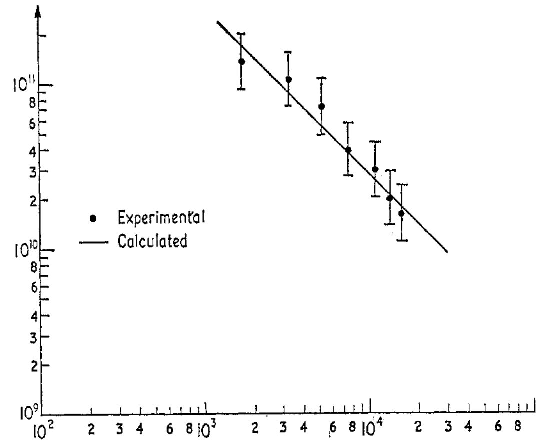

Figure 3. Comparison of measured and calculated threshold intensities for helium and ruby laser radiation with $\lambda = 6943$ Å.

helium for $2000 < p < 16 \times 10^3$ torr. The linear trend of the data with gradient -1 shows that in this pressure range neither diffusion nor recombination losses play an important part in controlling the breakdown threshold. It is the duration of the flash which principally determines whether or not breakdown occurs.

9.2. *Argon*

In contrast to helium the recombination coefficient is relatively large, of the order of 3×10^{-7} cm^{-3} s^{-1} at 300 K (Biondi 1951, Sexton and Craggs 1958, Oskam and Mittelstadt 1963). It decreases rapidly with electron temperature T as $T^{-3/2}$ and has a value of 2×10^{-8} cm^3 s^{-1} at 1800 K (Fox and Hobson 1966). However, the momentum transfer collision frequency is remarkably sensitive to electron energy and as we have *a priori* no exact value of the latter we cannot assign a well-defined value to ν_m. We note, however, from figure 2 that at the lowest pressures used the experimental data tend to lie along a line of gradient -1, while at the higher pressures there is a tendency towards asymptotic behaviour with a gradient of $-\frac{1}{3}$. These gradients are those deduced in the theory and support the view that it correctly describes laser-induced ionization growth and the onset of breakdown in argon.

In order to make a more direct comparison between predicted and measured values of $\hat{P}_{th}$ we may insert known values of R, K, δ and ω in equation (10) and adjust the value of ν_m, used as a variable parameter, to give the best fit between measured and calculated values of $\hat{P}_{th}$ as a function of pressure. This procedure was followed and the full line curve in figure 4 shows the calculated curve for $\nu_m = 3{\cdot}3 \times 10^9 p$ s^{-1} and $R = 3 \times 10^{-7}$ cm^{-3} s^{-1} which combination gave the best fit. No adequate agreement could be obtained by taking other combinations of ν_m and R—a result which is in accord with Young and Hercher. This value of ν_m is in reasonable agreement with published data (MacDonald 1966).

A further test of the analysis is based on the predicted dependence of $\hat{P}_{th}$ upon ω^2/τ. This should be a linear function with gradient $pN_0 \ln(\delta p N_0 V)/377 k \nu_m{}^2$ at constant pressure provided neither diffusion nor recombination effects are important. We have been unable to satisfy entirely the latter restriction and have succeeded in examining this dependence at relatively high pressure only ($p \gtrsim 5500$ torr) when recombination effects are not negligible. This is on account of the relatively low output obtainable from the dye laser. Despite these limitations we find the graph of measured $\hat{P}_{th}$ values is a linear function of

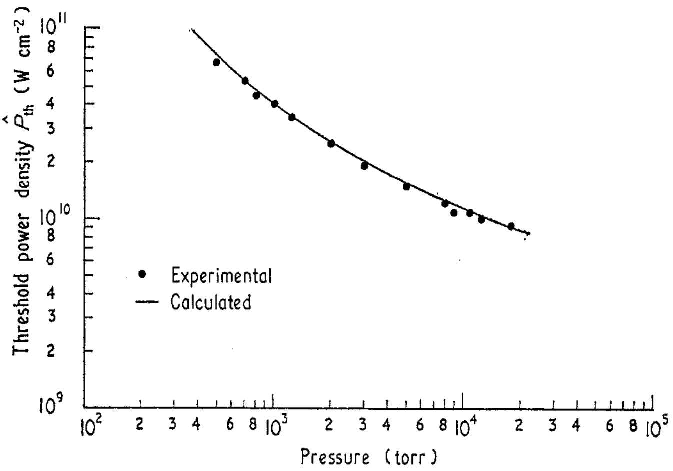

Figure 4. Comparison of measured and calculated threshold intensities for argon and ruby laser radiation with $\lambda = 6943$ Å.

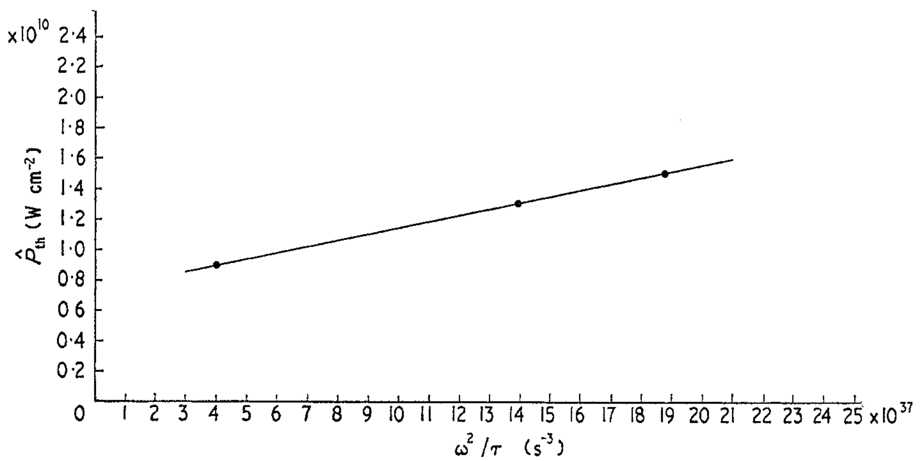

Figure 5. Threshold intensity as a function of ω^2/τ for argon at 5500 torr.

ω^2/τ for the three frequencies and two flash durations we can attain. This is shown in figure 5.

In order to examine the self-consistency of this result we note that we can use it to obtain an entirely independent estimate of ν_m for argon by equating the gradient of figure 5, which is approximately 4×10^{-29} W cm^{-2} s^{-3} to the expression given above using the known values of k, etc. The value obtained in this way is $\nu_m \simeq 3{\cdot}9 \times 10^9 p$ s^{-1}, which in view of the approximations made, is in broad agreement with the value of $3{\cdot}3 \times 10^9 p$ s^{-1} deduced earlier.

A further conclusion may be drawn from this self-consistency and from the fact that for values of R less than about 3×10^{-7} cm^{-3} s^{-1} corresponding to a temperature of about 300 K we cannot get agreement between measured and calculated values of $\hat{P}_{th}$. It follows that, at the onset of breakdown, the electron temperature is relatively low, and the electron mean energy is, from the published dependence of ν_m on electron energy (MacDonald 1966), about 4 or 5 eV. Much higher values are achieved during subsequent stages of plasma development (Lampis and Brown 1968).

These independent sets of experimental data and the reasonable agreement they show with the theoretical values lend support to the view that the cascade processes of direct impact ionization and excitation, followed immediately by photoionization, are responsible for laser-induced ionization growth and the onset of breakdown in helium and neon over the pressure range studied. Recombination is important in argon but not in helium where, in the present work, flash duration governs the growth.

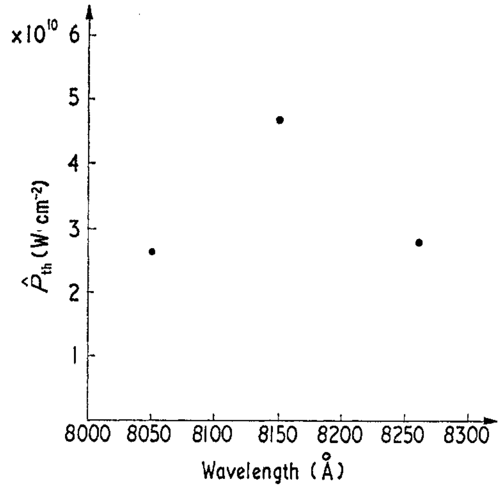

Figure 6. Threshold intensity as a function of wavelength.

It should be noted, however, that the results obtained with the dye laser at wavelengths of 8050 Å, 8150 Å and 8250 Å using argon at 10 640 torr, shown in figure 6, reveal a maximum in the $\hat{P}_{th}$ against λ curve at constant pressure and do not give a linear $\hat{P}_{th}$ against ω^2/τ dependence. They thus appear to contradict the view that laser-induced breakdown is described satisfactorily in terms of an extension of microwave theory. Similar departures have previously been reported by Buscher *et al.* (1965) and Alcock *et al.* (1969) for argon and xenon. We can, at this stage, but speculate that this experimental departure from the theoretical considerations may be connected with a substantial reduction in the probability of resonant multiphoton absorption over a narrow energy bandwidth corresponding to these wavelengths. Such reductions have been reported by Bebb and Gold in the case of hydrogen, but there are no data for argon in the relevant energy bandwidth. If this view is correct, then there could be a significant delay in the production of initiatory electrons and hence an effective reduction in the laser flash duration τ. Further work over a wider wavelength and lower pressure ranges is required to clarify this important point.

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