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BREAKDOWN AND HEATING OF GASES UNDER THE INFLUENCE OF A LASER BEAM

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INTRODUCTION

THE invention of lasers and improvements in laser techniques have uncovered extensive possibilities for the study of the interaction between matter and exceedingly intense beams of coherent radiation at optical frequencies. One of the most interesting phenomena in this region, which was observed in 1962 and immediately attracted the attention of many researchers, was the breakdown of gases and the formation of a "spark" in the focus of a laser beam. Experiments have shown that at very high light flux intensity, in gases which are usually perfectly transparent to optical frequencies (argon, helium, air, etc.), breakdown takes place, that is, strong ionization develops. The breakdown of the gas is accompanied by a light flash, as any ordinary spark discharge, and by virtue of this this phenomenon is frequently called a "spark."

Ionization occurs under these conditions if the radiation intensity exceeds a certain rather strongly pronounced threshold value. Experiments have shown that for gas breakdown it is necessary to employ very high intensities. If, as is frequently done, we characterize the radiation intensity by the intensity of the electric field in the light wave, then threshold fields are of the order of 10^6—10^7 V/cm (depending on the type and pressure of the gas).

Such high intensities of light radiation can be obtained by focusing with a lens the beam from a Q-switched laser operating in the so called giant-pulse mode, when the maximum (peak) power reaches several times 10 megawatts. The spark phenomenon was therefore observed only after methods were developed for pulsed Q switching of lasers.* The durations of the giant pulses are usually 20—40 nsec (1 nsec = 10^{-9} sec). A beam of laser light in a giant pulse has a divergence of 3—30′ ≈ 10^{-3}—10^{-2} rad, and is focused by the lens into a circle of radius 10^{-3}—10^{-2} cm. The characteristic densities of the light-energy fluxes at the focus are in the case of ruby lasers of the order of 10^{11} W/cm^2, and the photon flux densities $\sim 10^{30}$ cm^{-2} sec^{-1}.†

One can conceive of several mechanisms for gas ionization at very high intensities of light radiation. The electrons can break away from the atoms directly by the action of the radiation, as a result of a simultaneous absorption of many quanta. (The point is that the ionization potentials I of the atoms are much higher than the energies of the laser quanta; for example, I = 15.8 eV for argon, whereas a ruby laser quantum has $\hbar\omega$ = 1.78 eV, so that ionization of argon requires nine quanta.)

An electron cascade can develop in the radiation field. The multiquantum photoeffect produces in the gas the first, "priming" electrons. The electron absorbs light quanta by colliding with the atoms (in a process which is the inverse of bremsstrahlung) and

*See[46—49, 58, 59] concerning Q-switched lasers.

†We shall henceforth for brevity use simply "flux" in lieu of "flux density."

is accelerated. After accumulating an energy sufficient for ionization, the electron ionizes the atom, so that one electron is replaced by two slow ones, which begin the same process anew, etc. A combined cascade-light process is also possible, in which the electron does not ionize the atom, merely exciting it; the excited atom is then ionized by the radiation. It will be shown later that in dense gases (at pressures on the order of and higher than atmospheric) cascade ionization takes place. Only in rarefied gases, in which the electrons are emitted from the region of field action without experiencing many collisions, is the multi-quantum photoeffect the only ionization mechanism.

The spark phenomenon is also remarkable because when the generator power exceeds noticeably the value necessary for the breakdown, the gas becomes strongly ionized under the influence of the light even during the initial stage of the laser pulse, and the resultant plasma absorbs the laser beam to a considerable degree. The light energy is then released in a relatively small volume and the gas is heated in the absorption region to temperatures of tens and hundreds of thousands of degrees. This effect can be used to concentrate energy in matter and to obtain high temperatures in small volumes. The laser-beam absorption mechanism has interesting features. The light is not absorbed in a single volume with dimensions characterized by the radius of the focusing circle in which the primary gas breakdown has occurred. Owing to the ionization which occurs ahead of the layer that absorbs the light at a given instant, the new layer itself becomes capable of absorbing light. Thus, the laser-beam absorption region is continuously displaced after the primary breakdown in the direction towards the lens, producing something like a wave of light absorption and gas heating, which propagates in the gas opposite to the beam direction.

We consider below the new physical "laser spark" phenomenon, present a review of the experimental data, and describe the results of theoretical studies of the process.

The material is divided into two parts. The first contains everything pertaining to the breakdown effect, that is, to the development of ionization in the region of the focus and to the formation of the primary plasma. In the second part we consider the effects of absorption of the laser beam by the plasma and of the gas heating, and also some phenomena that are observed after the termination of the laser pulse. The last sections contain a brief report of results of an experimental study of the effect of giant pulses on solid targets and estimates of the possibility of heating hydrogen to thermonuclear temperatures. We begin each part with a summary of the obtained experimental data, as complete as possible, followed by an examination of the mechanism of the phenomena.

1. BREAKDOWN OF GASES

1. Measurement of Threshold Parameters

The first communication of an observation of breakdown of a gas (air) under the influence of a focused beam from a ruby laser, operating in the giant-pulse mode, was presented in February 1963 at the Third International Conference on Quantum Electronics in Paris.[1] This was followed by reports from several authors of measurements of threshold parameters for gas breakdown at different pressures

Meyer and Haught [2] used a ruby rod 15 cm long and 1.3 cm in diameter. A diagram of the setup is shown in Fig. 1. Pumping was by means of four xenon flash lamps, through each of which a capacitor bank was discharged. The Q was switched with a polarizer and a Kerr cell, as in the experiments of [1]. The result was a single giant pulse. The pulse energy, determined by calorimetric means, was in a typical case 1 J, the duration 30 nsec, and the peak power 30 MW.

No oscillograms of the laser pulse are published in [2] but one can be found in another paper by the same authors,[5] devoted to a study of the absorption of a laser beam at above-threshold laser power, with the measurements performed apparently by means of the same setup. We shall deal with this investigation in Sec. 6, in which an oscillogram of the pulse is shown (Fig. 5). The dimension of the focusing circle was determined from the hole burnt by the beam in thin gold foil (thickness 5×10^{-6} cm). The hole diameter was 2×10^{-2} cm. The same value was estimated from a measurement of the divergence of the laser beam 2α and the known focal distance of the lens f (radius of circle $r_0 \approx f\alpha$; the values of f and α are not given in the article). At 30 MW peak power and a focusing-circle diameter of 2×10^{-2} cm, the flux of radiant energy at the focus was 10^{18} erg/cm^{2}sec = 10^{11} W/cm^2. The occurrence of the breakdown was established by the appearance of a light flash in the focal region; the glow lasted approximately 50 μsec. In addition, a pair of electrodes, to which a low voltage of 100—200 V was applied, was used to register the breakdown. Approximately 10^{13} electron charges were drawn from the region of the focus during the breakdown.

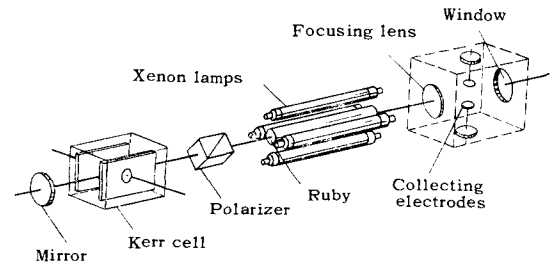

FIG. 1. Experimental setup used in[2].

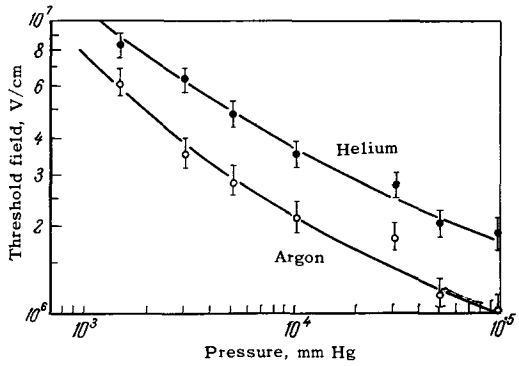

FIG. 2. Threshold fields for breakdown in argon and helium (as measured in[2]).

Breakdown in argon and helium was investigated at 1.5×10^3—10^5 mm Hg. The results of the experiments—the threshold fields for the breakdown as functions of the pressure—are shown in Fig. 2.* The authors note that the pressure for the breakdown is very critical, and a slight reduction in the power at a given pressure was sufficient to prevent occurrence of the breakdown.

Generally speaking, the threshold for the breakdown can be characterized by various quantities: the flux of radiant energy J, the flux of quanta $S = J/\hbar\omega$, the mean-square (average) intensity of the electric field in the light wave $E = \sqrt{4\pi J/c}$ (the amplitude of the field oscillations is $E_0 = \sqrt{2}E$). For convenience in conversion, we present some numerical formulas. If W is the laser power and r_0 the radius of the focusing circle, then

$$J = 3.2 \cdot 10^9 \frac{W\,[\text{MW}]}{(r_0\,[\text{cm}]/10^{-2})^2}\ \text{W/cm}^2,$$

$$S = 3.4 \cdot 10^{18}\,(J\,[\text{W/cm}^2]) = 1.1 \cdot 10^{28} \frac{W\,[\text{MW}]}{(r_0\,[\text{cm}]/10^{-2})^2}\ \frac{\text{photons}}{\text{cm}^2\ \text{sec}},$$

$$E = 10^{-8}\sqrt{S} = 19\sqrt{J\,[\text{W/cm}^2]} = 1.1 \cdot 10^6 \frac{\sqrt{W\,[\text{MW}]}}{r_0\,[\text{cm}]/10^{-2}}\ \frac{\text{volt}}{\text{cm}}.$$

Mink [3] used a Q-switched ruby laser with a rotating prism. The sawtooth pulse had a half-width of 25 nsec and a peak power of 3—5 MW. A lens with a focal distance of 2 cm was used. The breakdown was registered by noting the appearance of the glow. The plasma spectrum contained lines of singly ionized atoms and molecules; no lines of double ions were observed. At gas pressures above 15 atm the emission had a continuous spectrum. Breakdown was investigated in argon, helium, hydrogen, and nitrogen at pressures from 0.5 to 100 atm. The results of the work—the threshold laser power—are shown in Fig. 3. The diameter of the focusing circle was not measured. It is merely noted that an estimate, with allowance for the lens aberration and the characteristics of the laser beam, gives a focusing-circle diameter of 1.2×10^{-3} cm. This appears to be an underestimate. On the other hand, however, the diameter of the focusing circle in Mink's experiments was undoubtedly much smaller than the 2×10^{-2} cm obtained by Meyer and Haught,[2] since the power required for the breakdown of the argon and helium in the experiments of [2] exceeded by tens of times the threshold powers in Mink's experiments [3] (at the same gas pressures).

Nelson et al. [7] determined the threshold for breakdown in atmospheric air by a different method—by measuring the absorption of the laser beam. A Q-switched ruby laser produced pulses with energy up to 0.8 J and duration 30 nsec. The beam divergence, measured photometrically, was $2\alpha = 4 \times 10^{-3}$ rad, and the beam diameter was 1 cm. The radius of the focusing circle was defined as $r_0 = f\alpha$ (the focal distance is not indicated; judging from the maximum value of the beam flux at the focus, $\sim 8 \times 10^{10}$ W/cm^2, the radius of the circle was apparently $\sim 10^{-2}$ cm, that is, f = 5 cm). The coefficient of transmission of the laser beam through the focusing region was measured as a function of the laser power by a calorimetric method accurate to 5%. At power below threshold, the beam passed without absorption. Starting with some value of the power, the absorption of the beam increased very sharply with increasing power, indicating that an absorbing plasma was produced. At the same value of the power, a light flash—a spark—appeared. The threshold value of the energy flux was 5×10^{10} W/cm^2 $= 0.5 \times 10^{18}$ erg/cm^2 sec (field $E = 4.4 \times 10^6$ V/cm), and at this value of the flux the beam transmission

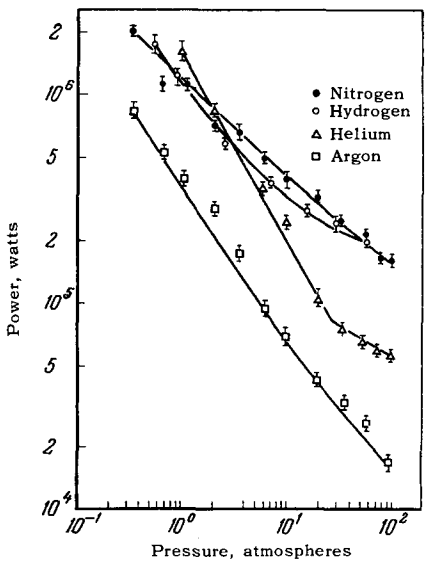

FIG. 3. Threshold laser power for breakdown in nitrogen, helium, and argon (as measured in[3]).

*The paper does not show how the field was calculated from the energy flux. It is not excluded that the values given are not the average fields, but maximum fields, which exceed the values averaged over the electromagnetic oscillation cycle by a factor of $\sqrt{2}$.

coefficient decreased very abruptly, almost jumpwise from 1 to ~0.3. With increasing incident flux to 8×10^{10} W/cm, the transmission coefficient decreased to ~0.2.

Tomlinson [15] investigated breakdown in noble gases at 400 mm Hg. The diameter of the focusing circle was determined by photographing the focal region through a microscope and defined as the distance between the points where the intensity of the beam was half the maximum value. The diameter was 3×10^{-3} cm. The laser pulse was attenuated with the aid of filters; the instant when the flash disappeared was recorded. From the corresponding value of the peak power, the threshold flux of quanta was determined. Values of 0.28, 0.33, 0.66, and 0.70×10^{30} photons/cm^{2}sec were measured for Kr, Ar, Ne, and He respectively. These data agree in order of magnitude with the results of the experiments of [2]. The transmission of light through the focal region during the time of the pulse was measured in [15]. The transmitted flux was registered with a time resolution ~1 nsec. Transmission curves are presented for argon at 1850 mm Hg and various values of excess of power above threshold. It is seen that when the power increases the breakdown and the absorption begin earlier and earlier.

To determine the frequency dependence of the threshold, Akhmanov et al.[27] measured the threshold for the breakdown of atmospheric air with the aid of a neodymium-glass laser. The measurements were made at the first and second harmonics ($\lambda_1 = 1.06\ \mu$, $\lambda_2 = 0.53\ \mu$). For the first harmonic, the diameter of the focal spot was 2×10^{-2} cm, the pulse duration 40 nsec, the threshold energy 0.3 J, and the threshold field 5×10^6 V/cm. The pulse duration at the second harmonic was smaller by a factor of 1.3, the focusing was approximately the same or about 15% better. The threshold energy, measured calorimetrically, was 1.6—1.7 times larger, that is, the threshold field at the second harmonic turns out to be 50% higher than at the first. Thus, the threshold for breakdown drops with decreasing frequency.

In all the foregoing investigations, the breakdown was produced in gases of not too low density, at pressures on the order of and above atmospheric, when the most important part in the formation of the free electrons is played by ionization by electron impact, and an electron cascade develops (see below). Unlike these experiments, Voronov and Delone [6] investigated ionization of a strongly rarefied gas—xenon—at pressures $\lesssim 10^{-2}$ mm Hg, when the ionization of the atoms by the electron impact certainly plays no role, since the electron mean free paths exceed by tens and hundreds of times the dimensions of the laser-beam focusing region. These experiments were undertaken to study the direct knock-out of electrons from atoms by an intense light wave. The ruby laser beam was split by a semitransparent mirror. One of the beams was focused in a vessel containing the gas. The positive ions were drawn by a collector with a field ~10 V/cm, the construction of the collecting device prevented charges which could be produced outside the focal region from striking the collector. The second beam was attenuated, focused, and the focal region was recorded with the aid of a micro-objective on a film. The dimension of the focusing circle was thus measured and the distribution of the illumination over the focusing cross section determined. The detector made it possible to register 4×10^3 ions (produced within a pulse time 20 nsec). At fields $(0.8—1.5) \times 10^7$ V/cm there were produced ~10^5 ions corresponding to an ionization probability of ~5×10^5 sec^{-1} under the influence of the light wave. Ionization of strongly rarefied gases (air, argon, helium) at pressures ~2×10^{-2} mm Hg was investigated earlier by Damon and Tomlinson. [4] They obtained strange results, which are subject to doubt: strong ionization (~10^7—10^8 ions) was observed at very low power ~200 kW (fields ~4×10^5 V/cm; laser without Q switching). Yet, according to the measurements of [6], no signal was observed in xenon at fields lower than 0.8×10^7 V/cm. The authors of that paper repeated the experiments of [4] and showed that the charges observed in [4] were not the result of ionization of the gas in the focus but apparently of ionization on the surface of the lens or on the vessel walls.

The question of the field distribution in the focusing cross section is of great importance in explaining the breakdown mechanisms. The point is that the radiation flux and the field are not at all uniformly distributed over the cross section, but have local maxima, corresponding to different vibrational modes of the laser (Leontovich and Veduta [8], Bradley [9]). The structure of the radiation field in the region of the focus was investigated experimentally by Barkhudarov et al. [10] The beam of a ruby laser, which was Q-switched by means of a Kerr cell, had a power ranging from several MW to several times 10 MW and was focused by lenses corrected for spherical aberration (f = 4.5 and 12 cm). An image of a selected beam cross section was obtained by means of a micro-objective on photographic film, and the distribution of the illumination over the cross section was recorded. Local inhomogeneities in the cross section of an unfocused beam are of the order of 10^{-2} cm (beam diameter ~1 cm); at the minimum cross section of the focused beam, the dimensions of the inhomogeneities are of the order of 10^{-3} cm. The measurements have shown that the cross section of the beam has a minimum not in the focal plane, but somewhat farther away from the lens. Thus, at a focal length f = 12 cm and a distance of ~1 meter from the lens to the ruby, the average dimension of the cross section (the square root of the area) was minimal and equal to 0.04 cm

at a distance 0.7 cm from the focal point. At the focus the cross section had a dimension of 0.1 cm, and at a distance of 1.2 cm behind the focus—0.06 cm. Photometric measurements of the inhomogeneity in the distribution of the illumination in the minimum cross section have shown that in this cross section there are points where the local illuminations exceed the average over the cross section by 30—50 times, and the electric fields are consequently 5—7 times larger than the average over the cross section. Of course, the dimensions of regions with very large local fields are exceedingly small.

In this connection, we point to the work of Zel'dovich and Pilipetskiĭ,[17] who examined theoretically the influence of lens aberrations on the focusing of a beam of light which has no intrinsic divergence, and carried out a diffraction calculation of the field with caustics.

2. Multiquantum Photoeffect

We proceed to consider the breakdown mechanism and to explain the results of the experiments. As noted in the introduction, one can conceive of two principal ways of gas ionization in a field of intense radiation: direct knock-out of the electrons from the atoms by the laser beam, and formation of an electron cascade.

Direct knock-out of electrons from atoms in a radiation field was considered by Bunkin and Prokhorov,[11] Keldysh,[12] and Gold and Bebb[13]; the most general results were obtained by Keldysh. The value of a quantum from a laser is much smaller than the ionization potential, so that the usual photoeffect is impossible in atoms. A multiquantum photoeffect, however, is possible in which the atomic electron is released as a result of simultaneous absorption of several photons (nine in the case of argon). The probability of simultaneous absorption of n photons by an atom is proportional to the n-th power of the quantum flux S, that is, to E^{2n}. The dependence on the field is very abrupt, and the multiquantum photoeffect occurs only in sufficiently strong fields in the light wave.

Along with the multiquantum photoeffect, one can conceive a different mechanism, similar to the tunnel effect—the knock-out of an electron from the atom by a static electric field.[18,19] The probability w_0 [sec^{-1}] of the tunnel effect in a field E is equal to

$$w_0 = \beta \exp\left(-\frac{4}{3}\frac{\sqrt{2m}I^{3/2}}{\hbar eE}\right), \qquad (2.1)$$

where the pre-exponential factor β is equal in order of magnitude to the frequency of motion of the electron in the atom $I/\hbar$.

It is obvious that the electric field of the light wave acts like a "static" field if it varies little during the time of flight of the electron through the potential barrier τ. The width of the barrier is of the order of $\Delta \sim I/eE$, and the electron velocity in the atom is of the order of $v \sim \sqrt{I/m}$, so that $\tau \sim \Delta/v \sim \sqrt{Im/eE}$. If the circular frequency of the field is ω, then the condition for the field to be quasistatic is the inequality $\omega\tau \sim \omega\sqrt{Im/eE} \ll 1$. At sufficiently high frequencies, the electron does not have time to jump through the barrier within a time on the order of the field-oscillation period, and the knock-out probability is frequency dependent; the inverse inequality $\omega\tau \gg 1$ is the condition under which the knock-out process has the character of the multiquantum photoeffect. The latter inequality is satisfied in the case of optical frequencies of interest to us: $\omega = 2.7 \times 10^{15}$ sec^{-1}* I ~ 10 eV, $E \sim 10^7$ V/cm, and $\omega\tau \sim 20$.

At first glance it may appear that the multiquantum photoeffect and the tunnel effect are different or even competing mechanisms. However, as follows from the work of Keldysh,[12] both mechanisms have a common nature and are limiting cases of a single process wherein the electron goes from a bound state in the atom into a free state under the influence of an alternating electric field. In that paper the probability is calculated of such a transition from the ground state of the hydrogen atom. Unlike the usual perturbation theory, in which stationary final states are considered, the author considered a nonstationary final state of the free electron in the oscillating electric field of the light wave $E(t) = E_0\cos\omega t$, giving rise to the transition. In calculating the matrix element, use is made of the exact wave function corresponding to this state; by the same token, the main accelerating action of the field on the free electron is taken into account.

As a result of all these calculations, a formula is obtained for the probability W [sec^{-1}] of knocking out an electron from the atom. The probability depends on the frequency ω and the field amplitude E_0, the latter entering into the formula only in the combination $\gamma = \omega\sqrt{2mI/eE_0}$, which coincides with the characteristic parameter $\omega\tau$.

In the limit of low frequencies and very strong fields, $\gamma = \omega\tau \ll 1$, the general expression for the probability is practically the same as the formula for the tunnel effect†. In the opposite limiting case of high frequencies, $\gamma = \omega\tau \gg 1$, the formula describing the multiquantum photoeffect is obtained.

This formula can be written in the following simplified form, which is convenient for numerical

*This is the frequency of the light from a ruby laser; wavelength λ = 6934 Å.

†The principal, exponential factor is exact, and the pre-exponential factor differs somewhat from the correct expression,[18] since no account was taken in the calculation of the influence of the Coulomb field of the nucleus on the wave function of the free electron.

estimates:

$$w = B\omega n^{3/2}\left(\frac{e^2E_0^2}{8m\omega^2 I}\right)^n = B\omega n^{3/2}\left(\frac{\pi e^2}{mc}\frac{\hbar S}{\omega I}\right)^n. \qquad (2.2)$$

here n is the number of quanta necessary to knock out the electrons, equal to the integer part of the quantity $(\tilde{I}/\hbar\omega) + 1$, with the effective ionization potential $\tilde{I}$ exceeding the usual ionization potential by a value equal to the average oscillation energy of the free electron in the field of the electromagnetic wave: $\tilde{I} = I + e^2E_0^2/4m\omega^2$. The values of the factor vary in a rather small range about unity. In estimates one can put B = 1. The point is that the probability w depends exceedingly strongly on the field E_0, and in practice it is necessary not to estimate w from E_0, but conversely, E_0 from w. Therefore even an appreciable spread in the values of B leads to a negligible variation in the value of the field: $E_0 \sim B^{1/2n}$.

Numerically we have for $\hbar\omega = 1.78$ eV

$$w = B\cdot 2.7\cdot 10^{15} n^{3/2}\left(4.75\cdot 10^{-34}\frac{I_H}{I}S\right)^n \text{sec}^{-1}, \qquad (2.3)$$

where $I_H = 13.6$ eV and S is in photons/cm^{2}sec.

If the energy of excitation of any state in the atom is very close to an integer number n_1 of quanta, the transition with the greater probability is that through an intermediate state, which includes resonant absorption of n_1 photons with excitation of the atom, and absorption of $n - n_1$ quanta, leading to the knock-out of the electron. The probability of such a process is calculated in[12]. Account was taken there of the Stark shift of the intermediate level; the width of the level is the sum of the natural width and the ionization probability of the level. Away from resonance, the probability of a two-step process is much smaller than the probability of direct knock-out; in the direct vicinity of the resonance, it can exceed by many times the latter, but the resonance for atoms is very narrow.

Gold and Bebb[13] calculated the probability of the multiquantum photoeffect by the usual method of perturbation theory.[16] The authors used a classical Hamiltonian of interaction between the atomic electron and the electromagnetic wave, $H = (e/mc)\mathbf{A}\cdot\mathbf{p}$, where $\mathbf{A}$ is the vector potential and $\mathbf{p}$ the electron momentum, and wrote down a general expression for the amplitude of the transition through the intermediate bound state in the lowest nonvanishing order of perturbation theory. The final, free state of the electron is described by a plane wave. As a rule, the main contribution to the sum over the intermediate states is made by one or two terms, including transitions at which the energy of several quanta turns out to be close to the excitation energy of some atomic level. Thus, for example, for the ionization of argon it is necessary to have nine quanta with $\hbar\omega = 1.78$ eV, while the energy of eight quanta differs by 0.18 eV from the excitation energy of the $5p^3S_1$ level. The ionization of helium requires fourteen quanta; the main contribution is made by transitions through the intermediate state $3p'P_1$, the energy of which differs from $13\hbar\omega$ by 0.116 eV, and the state $2s^1S_0$, the energy of which differs from $12\hbar\omega$ by 0.811 eV. The paper contains a summary table with the results of calculations for all the noble gases. The numerical formulas for the ionization probabilities of argon and helium are of the form

$$w_{Ar} = 1.4\cdot 10^{-281}S^9,\ w_{He} = 5.4\cdot 10^{-452}S^{14}\ \text{sec}^{-1}.$$

Numerically the formulas of Gold and Bebb give approximately the same results as the formulas of Keldysh (in the sense that for the same probability w they obtain very close values of the quantum fluxes), but the dependence of w on the frequency ω is somewhat different*.

Gold and Bebb[13] calculated the threshold fluxes necessary to "break down" gases by the multiphoton mechanism (to knock out 10^{13} electrons within 10^{-8} sec in the volume of the focus at a density of 10^{20} atoms/cm^3 as in the experiments of[2]). The threshold fluxes are 7×10^{31} and 4.4×10^{32} photons/cm^{2}sec for Ar and He, respectively. The experimental values are $S \sim 1-2\times 10^{29}$ for Ar and $(3-5)\times 10^{29}$ for He.

Thus, from these data for strong multiphoton ionization of the atoms, the fields required exceed by one order of magnitude the experimental averaged threshold fields. The same conclusion was reached also by Tomlinson[15] who compared the calculations based on[13] with the results of his own experiments with all the noble gases.

In drawing conclusions concerning the real role of this mechanism of ionization, it necessary to bear in mind that under the experimental conditions the local fields in very small volumes can greatly exceed the average values over the focusing volume, obtained from experiment; multiphoton ionization occurs predominantly just in the places with the large fields. In addition, no account is taken in the calculations of[12,13] of the smearing of the upper levels of the atom in the strong electric field of the light wave, which may actually cause the number of quanta necessary for the ionization of the atom to be re-

*Tozer[39] derived a semi-qualitative formula for multiquantum absorption. The probability of this process is defined as the probability of collision of n particles – photons – with an electron during an interaction time $\sim 1/\omega$, which follows from the uncertainty principle $\hbar\omega t \sim \hbar$. The formula contains an undetermined collision cross section, which is assumed in the estimate to be 10^{-16} cm^2. The multiquantum photoeffect could be analyzed with the aid of this formula. However, it is hardly meaningful to apply this formula, as was done by the author, to the multiquantum transition of an atom to the first excited state, which unlike the transition to the continuous spectrum (photoeffect) has a sharply pronounced resonant character.

duced by unity; this leads to an increase in the knock-out probability.*

However, just the same, it is improbable that the multiquantum photoeffect alone can explain the observed threshold for the breakdown in gases at high pressures. First, the gap between the experimental values of the threshold field and those calculated from [12,13] is too large; second, the assumption of the multiphoton mechanism is incompatible with the dependence of the threshold on the gas pressure. In fact, with increasing pressure, say from 1.5×10^3 to 10^5 mm Hg in helium, the threshold field decreases from 8×10^6 to 2×10^6 V/cm, that is, by a factor of 4. Even if we assume that at the lower pressure and in the stronger field the multiquantum photoeffect takes place, the photo-ionization probability at the higher pressure and the weaker field should be $\sim 4^{28} \approx 10^{17}$ times smaller. Consequently, at high pressures and low experimental threshold fields this effect certainly plays no role.†

In a strongly rarefied gas, where the possibility of formation of an electron cascade is excluded, the multiquantum photoeffect is the only ionization mechanism. Experiments [6] with rarefied xenon yield an average threshold field of 8×10^6 V/cm. From Keldysh's formula ($n = 7$) for the experimental conditions the threshold field amounts to 4.5×10^7 V/cm, and according to the formulas of [13], to 2.7×10^7 V/cm. If we take into consideration the existence of strong local fields and the unaccounted for reduction in the ionization potential, then we must recognize that these experimental data agree with the assumption of multiphoton ionization.

An important role may be assumed in the breakdown process by ionization of the excited atoms by the radiation. Zernic [14] calculated exactly the probability of the two-quantum photoeffect from the metastable 2s level of the hydrogen atom. For $\hbar\omega = 1.78$ eV, $w = 1.53 \times 10^{-47}$ S sec^{-1}. It is interesting to note that the approximate calculations of [13] give for this case a value 9 times larger, and formula (2.3) gives a value 250 times smaller (Keldysh's formula, generally speaking, pertains to the case of absorption of a larger number of quanta). We call attention to [60], in which two-quantum ionization of negative ions of iodine under the influence of a laser beam was investigated.

3. Cascade Ionization (Qualitative Picture)

Cascade ionization of a gas at the focus of a laser beam, which occurs at relatively high pressures (on the order of atmospheric and above) is the subject of theoretical papers by Ya. B. Zel'dovich and the author,[20] Wright,[21] Ryutov,[22] and Askar'yan and Rabinovich.[23] According to the general notions developed in [20], the process proceeds as follows.

In places with large local fields, multiphoton absorption (most likely by impurity atoms with low ionization potentials) causes the first "priming" electrons to appear at the beginning of the laser pulse.*

The free electron absorbs light quanta by collision with neutral atoms. Along with absorption, there occurs in the collisions under the influence of the intense radiation a stimulated bremsstrahlung of quanta of the same energy and direction. The energy of the electron in the random acts of interaction, with emission of $\hbar\omega$ batches, changes alternately in one direction and then in the other, so that the energy variation has in the main the character of a one-dimensional diffusion along the energy axis. The larger the quantum flux, the faster this process. After accumulating an energy somewhat higher than the ionization potential, the electron ionizes an atom, with high probability, as a result of which one electron gives rise to two electrons with lower energy (a new generation) which begin the entire cycle anew. At sufficiently strong fields there is not even any need for the electron to acquire the ionization energy. It is sufficient that the electron excite the atom (the excitation potentials of the first levels of the atoms, following the ground state, are approximately 3/4 of the ionization potential), since the excited atom is rapidly ionized by the radiation after absorbing two or three quanta.

On the other hand, in weak fields, insufficient for rapid ionization of the first excited states of the atoms, the energy loss by the electron to excitation hinders the development of the cascade; the electron can acquire and lose its energy many times before it succeeds in "jumping over" the energy band in which it excites the first states of the atoms above the ground state, and reaches an energy sufficient for the ionization or excitation of high-lying states that can be readily ionized by the radiation (we note that the ionization of excited atoms by electrons is at first, while the number of electrons is still small, very rare).

There are still no accurate calculations of the probability of ionization of excited atoms by radiation with account of the different effects which occur in this connection. Therefore there is considerable uncertainty in the question concerning the role of the excitation, more accurately the value of the field that

*This argument was advanced by L. V. Keldysh.

†We note that in[7] the breakdown in air at atmospheric pressure was erroneously interpreted on the basis of the notion of the multiquantum photoeffect. The apparent agreement with experiment was attained by the authors only because their formulas contained two undetermined constants, which were chosed in a suitable manner.

*Calculations by means of the formulas of Keldysh and of Gold and Bebb (the latter have been published in[15]) show that in fields $(3 - 4) \times 10^7$ V/cm, even without impurities, one electron appears in argon or helium within a time ~ 1 nsec.

demarcates the regions of positive and negative influences of excitation on the cascade development. For the time being it is necessary to make do with semiqualitative considerations. Two-photon ionization, undoubtedly, proceeds rapidly. From Zernic's formula [14] (which, to be sure, pertains to the 2s state of the hydrogen atom), we obtain even in a field of 10^6 V/cm a probability $\sim 10^9$ sec^{-1}. Under ordinary conditions, however, the first excited levels of argon, helium, etc. atoms are separated from the end point of the continuous spectrum by more than two quanta; three quanta are necessary (and even more for nitrogen or oxygen molecules). For three-quantum ionization with probability $\sim 10^9$ sec^{-1} it is necessary, according to Keldysh's formula (2.3), to have fields of 10^7 V/cm. Of course, an important role may be played here by the lowering of the end point of the continuous spectrum in the field of the wave, connected with the smearing-out of the upper levels (but possibly the first levels also drop simultaneously). The considerable broadening of the levels uncovers great possibilities for resonant transitions via intermediate states. This pertains in particular to molecules where there are many more levels and where they are more likely even to overlap. Thus, one can conceive of different possibilities of obtaining a very large probability for ionization of excited atoms by radiation at fields lower than 10^7 V/cm, say at fields equal to several times 10^6 V/cm and perhaps even 10^6 V/cm*. We emphasize once more that the excitation of the upper readily-ionized levels of atoms always contributes to the development of the cascade.

The developed cascades always slow down the electron energy loss in elastic collisions with atoms, and also the loss of the electrons themselves which diffuse from the effective region of the radiation (unlike energy loss, a loss of electrons breaks the multiplication chains in the cascade). The energy loss due to elastic collisions is the higher, the lighter the gas, and plays accordingly a relatively increasing role the slower the electrons acquire energy from the radiation, that is, the weaker the field. Under the conditions of the experiments in [2], elastic losses in argon and helium are generally speaking noticeable, and in helium at high pressures and weak fields they are even large.

The diffusion of the electrons is the slower, the higher the gas density and the larger the dimensions of the focusing region. According to estimates, the electrons under the conditions of the experiments in [2] do not have time to leave an effective radiation region with radius $\sim 10^{-2}$ cm even during the entire time of the pulse. However, from the very small volumes in which large local fields exist, the electrons drift rapidly, that is, they become distributed over the entire volume of focusing, and in general are not subject to the action of the local fields but of fields averaged over the volume.

If the diffusion drift of the electrons from the region of action of the field does not play a primary role (as in the experiments of [2]), the threshold field is determined by the condition that a sufficient number of electrons N_1 appear during the short time of the light pulse ($t_1 \approx 30$ nsec). The number of electrons in the cascade, at least at the start of the process, increases exponentially $N = N_0 e^{t/\theta} = N_0 \cdot 2^k$. The number of generations of electrons k, produced by the end of the pulse, $k_1 = t_1/\theta \times \ln 2 = 1.45 \ln (N_1/N_2)$, depends very little on the numbers N_0 and N_1. We can therefore assume approximately that the cascade begins, say, with one electron ($N_0 = 1$), and that the breakdown condition is the attainment of a certain number N_1 of electrons, for example 10^{13}, as in the experiments of [2]. Thus, the breakdown condition fixed the number of generations ($k_1 = 43$), and for a given pulse duration (30 nsec) also the cascade time constant $\theta = 1$ nsec (the time necessary for the number of electrons to double is $\theta_2 = \theta \ln 2 = 0.7$ nsec). The final number of electrons N_1 is extremely sensitive to the time constant θ, which depends on the field, and this explains the existence of an abrupt breakdown threshold.

In the case of a high degree of focusing or sufficiently low pressures (lower than atomospheric if $r_0 \approx 10^{-2}$ cm), the multiplication of electrons is greatly influenced by their diffusion drift. Under certain conditions, a quasi-stationary state can arise, when the electron production is offset by their loss (critical conditions, using nuclear reactor theory terminology; the conditions considered above, when the loss is small, are supercritical). At low pressures, the electrons leave the focusing volume rapidly and no cascade develops at all. In this case, the breakdown can result only from multiphoton processes.

4. Growth of Electron Energy in the Radiation Field

The rate of development of the cascade is determined primarily by the rate of growth of the electron energy in the radiation field. The latter is a quantity appearing in the theory of high-frequency breakdown in gases, which has features common with the breakdown at optical frequencies. In the theory of high-frequency breakdown,[24] the acceleration of the electron by the electromagnetic wave is treated classically. A free isolated electron oscillates in the alternating electric field of the wave, and its energy remains in the mean constant. The mean energy of electron oscillation is $e^2E^2/2m\omega^2$. The electron acquires energy from the wave, only when it collides with the atoms, when the abrupt change in the velocity

*Some effects of this kind were discussed by L. V. Keldysh.[12]

direction causes the oscillation energy to go over into the energy of random translational motion. The rate of growth of the electron energy ϵ under the influence of the electromagnetic wave is [25]

$$\frac{d\varepsilon}{dt}=\frac{e^2E^2}{m\omega^2}\nu_{eff}\frac{\omega^2}{\omega^2+\nu_{eff}^2}. \quad (4.1)$$

Here $\nu_{eff} = N_a v \sigma_{tr}$ is the effective frequency of collisions between the electron and the atoms, N_a the number of atoms per cm^2, v the electron velocity, $\sigma_{tr} = \sigma(1 - \overline{\cos\vartheta})$ the transport cross section, σ the elastic-scattering cross section, $\overline{\cos\vartheta}$ the average cosine of the scattering angle. If $\omega^2 >> \nu_{eff}^2$, formula (4.1) assumes the simpler form

$$\frac{d\varepsilon}{dt}=\frac{e^2E^2}{m\omega^2}\nu_{eff}. \quad (4.2)$$

The energy equal in order of magnitude to the electron oscillation energy in the field, which on the average is taken during each collision from the energy of the electromagnetic waves, consists under the high-frequency breakdown conditions of many quanta. Its order of magnitude is 10^{-3} eV, whereas $\hbar\omega \sim 10^{-4}$—10^{-5} eV ($\lambda \sim 1$—10 cm), making the classical treatment natural. In the case of breakdown at optical frequencies, the situation is reversed, the quantum energy $\hbar\omega = 1.78$ eV is much larger than the oscillation energy, $\sim 10^{-2}$ eV. Therefore in the overwhelming majority of collisions the electron will not acquire energy from the field (nor will it deliver any to the field), and only once in a large number of collisions will it acquire from the radiation instantaneously a large batch of energy $\hbar\omega$. The effect is quantum-like in character and it would appear that classical theory is not applicable here. Similar statements were made by many authors (for example [2,21]). Some authors, to the contrary, use classical theory without any stipulations. A solution of this problem is contained in [20], where the energy acquisition by the electron is treated on the basis of quantum notions, which in this case undoubtedly agree more with the nature of the process than the classical concepts. However, a clarification of the conditions of the transition to the classical limit shows that formula (4.2) can be used approximately also in the case of optical frequencies (in this case always $\omega^2 \gg \nu_{eff}^2$).

To find the rate at which energy is acquired and the time constant of the cascade development it is necessary, generally speaking, to start from the quantum kinetic equation for the distribution function of the electrons interacting with the radiation. To this end it is necessary to find first the coefficient of absorption of quanta by the electrons in the gas. This can be done by considering, on the basis of classical electrodynamics, the bremsstrahlung of an electron that experiences a collision with a neutral atom, and then determining the corresponding absorption coefficient, using the principle of detailed balance. The approximate differential cross section for bremsstrahlung of quanta was obtained in [20]*

$$d\sigma_\omega=\frac{4}{3\pi}\frac{e^2v^2\sigma_{tr}}{c^3\hbar\omega}d\omega. \quad (4.3)$$

The coefficient of absorption of light by electrons with energy $\epsilon = mv^2/2$, divided by the number of the electrons and by the number of atoms, $a_\omega(\epsilon)$ [cm^5], is

$$a_\omega(\varepsilon)=a_{\omega c}\cdot\frac{2}{3}\frac{\varepsilon+\hbar\omega}{\varepsilon}\frac{\varepsilon+\hbar\omega}{\hbar\omega}\frac{\sigma_{tr}(\varepsilon+\hbar\omega)}{\sigma_{tr}(\varepsilon)},$$
$$a_{\omega c}=\frac{4\pi e^2v\sigma_{tr}(\varepsilon)}{mc\omega^2}. \quad (4.4)$$

The quantity $a_{\omega c}$, separated in the formula, is the effective absorption coefficient obtained in the classical theory of absorption of electromagnetic waves in a weakly ionized gas [25] when $\omega^2 \gg \nu_{eff}^2$.†

The coefficient of stimulated emission of quanta $\hbar\omega$ by electrons with energy $\epsilon + \hbar\omega$, according to Einstein's relation, is‡

$$b_\omega(\varepsilon+\hbar\omega)=\sqrt{\frac{\varepsilon}{\varepsilon+\hbar\omega}}a_\omega(\varepsilon). \quad (4.5)$$

Figure 4 shows the coefficients a_ω and b_ω for argon and helium, calculated in [20] by means of formulas (4.4) and (4.5) on the basis of experimental elastic-scattering cross sections. We note that in this case spontaneous emission does not play any role.

We now set up a kinetic equation for the electrons interacting with the radiation. Let $n(\epsilon, t)d\epsilon$ be the number of electrons per cm^3, possessing energies from ϵ to $\epsilon + d\epsilon$. Leaving out for brevity the index ω, and disregarding the spatial dependence of the

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*In a more rigorous derivation with account of the correlation between the individual collisions of the electrons and the interference of the radiated waves, there appears in (4.3) the factor $\omega^2/(\omega^2 + \nu^2_{eff})$ which is contained also in (4.1). The effect of correlation is considered in a paper by the author.[26] In this case, as already noted, this factor is close to unity.

†Inasmuch as the absorbed energy of the electromagnetic wave goes into the acceleration of the electrons, the classical absorption coefficient $a_{\omega c}$ can also be obtained from (4.2). If N_e is the number of electrons per cm^3 and $J = cE^2/4\pi$ is the radiant energy flux, then

$$N_e\frac{d\varepsilon}{dt}=N_e\frac{e^2E^2}{m\omega^2}\nu_{eff}=Ja_{\omega c}N_aN_e.$$

From this we get directly expression (4.4) for $a_{\omega c}$.

‡The absorption coefficient of frequency ω by electrons with energy ϵ, corrected for the stimulated emission, which corresponds indeed to the classical coefficient of effective absorption, is

$$a'_\omega(\varepsilon)=a_\omega(\varepsilon)-b_\omega(\varepsilon)=a_\omega(\varepsilon)-\sqrt{\frac{\varepsilon-\hbar\omega}{\varepsilon}}a_\omega(\varepsilon-\hbar\omega).$$

In the limit as $\hbar\omega/\epsilon \to 0$

$$a'_\omega(\varepsilon)=(a_{\omega c}/3)\,[1+2d\ln(a_{\omega c}\varepsilon)/d\ln\varepsilon].$$

Thus, if $\nu_{eff}(\epsilon) = \text{const}$ and $a_{\omega c}(\epsilon) = \text{const}$, the limiting value of a'_ω coincides exactly with the classical value $a_{\omega c}$.

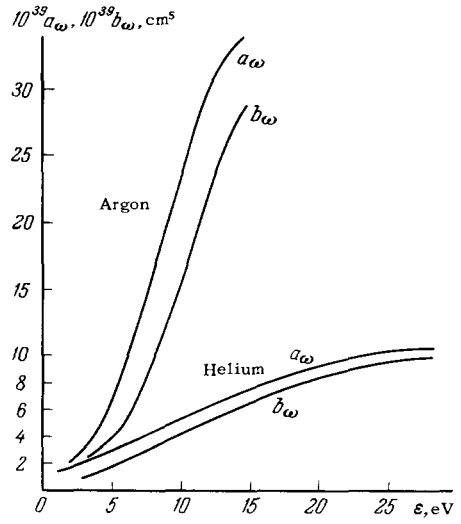

FIG. 4. Coefficients of true absorption and stimulated emission of quanta from a ruby laser in argon and helium, calculated per electron and per atom (from the data of[20]).

electron density, we write

$$\frac{\partial n(\varepsilon)}{\partial t} = SN_a\{-a(\varepsilon)n(\varepsilon) + b(\varepsilon+\hbar\omega)n(\varepsilon+\hbar\omega) - b(\varepsilon)n(\varepsilon) + a(\varepsilon-\hbar\omega)n(\varepsilon-\hbar\omega)\} + Q, \quad (4.6)$$

where Q describes processes which are not connected with radiation (elastic and inelastic energy losses, electron drift from the volume, etc.). Since we are interested in a rather appreciable interval of energies ϵ, which is of the order of the excitation energy or ionization energy of the atom, and is approximately ten times larger than $\hbar\omega$, we shall regard $\hbar\omega/\epsilon$ as a small quantity. Expanding in powers of this parameter in (4.6) (accurate to terms of third order) and in the expression (4.5) for b_ω, we obtain

$$\left.\begin{aligned} \frac{\partial n}{\partial t} &= -\frac{\partial j}{\partial \varepsilon} + Q, \\ j &= nu - D\frac{\partial n}{\partial \varepsilon}, \\ D(\varepsilon) &= SN_a(\hbar\omega)^2 a(\varepsilon), \\ u(\varepsilon) &= \frac{1}{2}\frac{D(\varepsilon)}{\varepsilon}. \end{aligned}\right\} \quad (4.7)$$

This is the equation of one-dimensional diffusion of particles in a flux with sources Q; the role of the coordinate is played by the energy ϵ. The quantity D has the meaning of a diffusion coefficient, and u is the flow velocity along the ϵ axis. Actually $SN_a \times (a + b) \approx 2SN_a a = 1/\tau'$ is the reciprocal of the lifetime of the electron relative to absorption or emission of a quantum, $\hbar\omega$ is the magnitude of the jump along the energy axis, that is, the "mean free path," $\hbar\omega/\tau'$ is the average velocity of "random motion" along the axis, and $D = (\hbar\omega/2)\hbar\omega/\tau'$ in accordance with the definition of the diffusion coefficient.

Let us trace the variation of the energy of the "average" electron produced with the low energy ($\epsilon \approx 0$) under the assumption that there are no other processes except interaction with the radiation ($Q = 0$). We put here $\nu_{eff}(\epsilon) = \text{const}$ (this is necessary for an exact transition to the classical limit), so that according to (4.7) and (4.4) we have $D \sim a(\epsilon) \sim \epsilon$ when $\hbar\omega/\epsilon \ll 1$. Multiplying the diffusion equation (4.7) by ϵ and integrating from 0 to ∞, we obtain after simple calculations an equation for the average electron energy:

$$\frac{d\bar{\varepsilon}}{dt} = u + \frac{dD}{d\varepsilon} = u + \frac{D}{\varepsilon} = \frac{3}{2}\frac{D}{\varepsilon}. \quad (4.8)$$

We substitute here expressions (4.7) for D, u, and a in accord with formula (4.4). Going over to the limit $\hbar\omega/\epsilon \to 0$ and noting that $S\hbar\omega = cE^2/4\pi$, we obtain for $d\epsilon/dt$ the classical formula (4.2). Thus, the condition for its approximate applicability is the smallness of the quantum $\hbar\omega$ compared only with the total electron energy ϵ, not compared with the energy of the oscillations in the field.

If we take into account the electron energy losses in elastic collisions with the atoms, we must write in lieu of (4.2)

$$\frac{d\varepsilon}{dt} = \left(\frac{e^2E^2}{m\omega^2} - \frac{2m}{M}\varepsilon\right)\nu_{eff}, \quad (4.9)$$

where M is the mass of the atom (in principle one could also add here a term describing inelastic losses).

This formula demonstrates clearly the shortcoming of the classical description of the processes of energy acquisition in the case when $\hbar\omega \gg e^2E^2/m\omega^2$. Judging from this formula, the electron energy can never exceed*

$$\varepsilon_m = \frac{e^2E^2M}{m\omega^2 \cdot 2m}.$$

Actually, however, in view of the quantum character of the energy acquisition there always exists a finite probability that the electron can absorb a quantum within a short time, and at the same time does not lose much energy to elastic collisions, so that its energy grows to a value exceeding ϵ_m. This pertains to an even greater degree to the case when inelastic losses are appreciable. The foregoing circumstance is automatically taken into account in the description of the process with the aid of the quantum kinetic equation.

We present a numerical example of the rate of acquisition of energy by means of formula (4.9). In argon at 1500 mm Hg $N_a = 5.3 \times 10^9$ cm^{-3} and $\nu_{eff} \approx 1.7 \times 10^{13}$ sec^{-1} (for $\epsilon \sim 10$ eV). In a threshold field $E = 6 \times 10^6$ V/cm^2 the elastic losses amount to approximately three percent of the energy acquired. The energy increase per collision is $\Delta\epsilon \approx 0.9 \times 10^{-2}$ eV and $d\epsilon/dt \approx 150$ eV/nsec.

The electron energy reaches a value on the order of excitation and ionization potentials (~15 eV)

*For example, in helium $\epsilon_m = 20$ eV at $E = 4.7 \times 10^6$ V/cm.

within ~0.1 nsec. Analogously, for helium at the same pressure we have $\nu_{eff} \approx 4 \times 10^{12}$ sec^{-1} and $E_{thr} \approx 8.5 \times 10^6$ V/cm^2; at $\epsilon = 15$ eV, the ratio of the elastic losses to the energy acquired from the radiation is 0.23; $\Delta\epsilon \approx 1.6 \times 10^{-2}$ eV and $d\epsilon/dt \approx 64$ eV/nsec. An energy of 20 eV is reached after 0.3 nsec.

5. Calculations of Cascade Ionization

The electron cascade was calculated in [20] on the basis of the quantum kinetic equation, or more accurately the diffusion equation (4.7), to which the kinetic equation reduces when $\hbar\omega/\epsilon \ll 1$. The nonstationary solution of the equation, describing the electron multiplication for supercritical conditions, is of the form $n(\epsilon, t) = n(\epsilon) \exp(t/\theta)$. The boundary conditions which must be satisfied by the electron energy distribution function $n(\epsilon)$, with account of the multiplication, give an equation for the time constant θ of the cascade. It was assumed in that paper that the excited atoms are not ionized by the radiation, that is, that the excitation acts retard the development of the cascade. In this case the constant θ was found to have an approximate value $\theta = t^*/\alpha$. Here t^* is the time required for the energy of the electron to grow as a result of absorption and stimulated emission of quanta ("diffusion along the energy axis") to a value sufficient for rapid ionization of the atom by electron impact. In accordance with the statements made in the preceding section regarding the relation between the quantum and classical treatments, this time coincides with the analogous quantity calculated from the classical formula (4.2). The quantity α is the probability that the electron can "jump" through the excitation band, that is, $1/\alpha$ is the number of cycles of motion along the energy axis, which must be executed by the electron before multiplication takes place. The probability α, which decreases from several tenths to several hundredths in the pressure range from 1 to 100 atm and at threshold fields, was calculated by solving the equation of diffusion along the energy axis. Elastic losses were not taken into account in the calculations. As a result of the calculation of the threshold fields, reasonable agreement with the measurements of [2] was obtained. However, in light of the notions concerning the rapid ionization of the excited atoms by the radiation (see Sec. 3), it is possible that it is sufficient to have even one or a small number of cycles for electron multiplication (the value of α is close to unity). It must be noted that under this assumption, in the case of low pressures, multiplication is too fast. According to the estimate made at the end of the last section, the time necessary for doubling of the number of the electrons at a pressure of 1500 mm Hg and at experimental values of the threshold fields is of the order of 0.1—0.3 nsec, whereas it follows from experiment that it is several times larger, approximately 0.7 nsec. In other words, the threshold fields should be somewhat smaller than those measured experimentally. At high pressures ~100 atm and small measured threshold fields, an important role is played by the elastic losses, which were not taken into account on the calculations. However, at such high pressures, when the spatial diffusion of the electrons in the focusing region is slow, the cascade develops predominantly in places where the local fields exceeds noticeably the average threshold fields. This makes it difficult to compare the calculations with experiment, which yields only fields averaged over the volume.

D. D. Ryutov [22] studied breakdown on the basis of the classical kinetic equation for an electron velocity distribution function $f(\mathbf{v})$

$$\frac{\partial f}{\partial t} - \frac{e\mathbf{E}_0}{m} \sin \omega t \frac{\partial f}{\partial \mathbf{v}} = I(f),$$

where $I(f)$ was the integral of collisions between the electrons and the atoms. The atoms were assumed to be infinitely heavy, that is, the elastic energy losses were not taken into account. The solution of the equation is sought in Lorentz form $f(\mathbf{v}) = f_0(\mathbf{v}) + g(\mathbf{v}) \cos \vartheta$ where ϑ is the angle between the directions of the velocity $\mathbf{v}$ and the electric field vector $\mathbf{E}_0$. As a result of several transformations, the kinetic equation reduces to an equation of the diffusion type along the energy axis, analogous to (4.7). Terms are added to the equation to describe the inelastic energy loss to excitation and the electron drift from the region of action of the field as a result of spatial diffusion. The excitation of the atoms, as in [20], is assumed to be a "harmful" effect. The equation is solved for stationary conditions, which were called above "critical," when the electron production is compensated by loss due to escape from the volume. The fields corresponding to the stationary conditions are obtained and are regarded as corresponding to threshold of the breakdown.* No estimates were made, however, of the actual number of electrons produced at the focus under such conditions. The estimated fields are in reasonable agreement with the results of the experiments of [3], in which a high degree of focusing was obtained and a "quasistationary" mode was probably realized.

Wright [2] started with quantum notions concerning the absorption of radiation, writing down a semiqualitative expression for the quantum absorption coefficient and a formula for the rate of growth of the electron energy $d\epsilon/dt$; the stimulated emission was not taken into account. He estimated the time required for the electron to attain excitation energy; the paper favors the idea of rapid ionization of the

*The same procedure is used in the theory of high-frequency breakdown.

excited atoms by the radiation. The breakdown conditions are estimated.

The calculations of Askar'yan and Rabinovich [23] pertain to strong fields, in which the excited atoms are rapidly ionized by the radiation.

The paper emphasizes the fact that owing to the large quantum absorption probability, the electron with enough energy to excite the first level of the atom, following the ground state, acquires rapidly, as a result of absorption of only one quantum, the ability of exciting the upper states, which can be readily ionized by the radiation. The kinetic equation for the number of electrons is written with allowance for the excitation and ionization of the upper levels of the atom. The effective electron energy loss per multiplication act, ϵ^*, is introduced, so that the time constant θ of the cascade is defined by the formula $1/\theta = (d\epsilon/dt)/\epsilon^*$. Notice is taken in the paper of the influence of the striction force on the electron drift from the focusing volume.

Summarizing the discussion of the cascade ionization mechanism, we must emphasize that there are still several unclear points which require additional theoretical and experimental research. The role of the excitation of the atoms at not very strong fields is not yet clear, there are no consistent calculations for breakdown in light gases at high pressures and relatively weak fields, when the elastic losses are high. There are no calculations for breakdown with allowance for spatial inhomogeneity of the field, which plays a very important role at high pressures when the cascade is localized at the points with the maximum fields. There is no theory allowing for the effects that are specific for molecular gases (air), where the excitation levels lie very low and where it is consequently very important to understand the influence exerted by the excitation of the molecules under various conditions. At the same time, the low value of the times necessary for the electrons to acquire energy in the radiation field, and the reasonable agreement between the threshold fields calculated under various assumptions with those measured experimentally, offer evidence that at pressures on the order of atmospheric and above the breakdown under threshold conditions has undoubtedly a cascade-like character. Favoring the cascade mechanism, in the opinion of the authors of [27], is also the observed lowering of the threshold fields with decreasing frequency.

II. ABSORPTION OF THE BEAM AND HEATING OF THE GAS

6. Measurements of Absorption and Scattering of a Laser Beam and Heating of the Gas in the Breakdown Region

If the intensity of the light beam exceeds noticeably the threshold intensity for breakdown, the latter occurs even before the end of the laser pulse, and an appreciable fraction of the radiation energy is absorbed by the produced plasma. At the same time the plasma is heated to high temperatures. This phenomenon was the subject of numerous experimental papers.

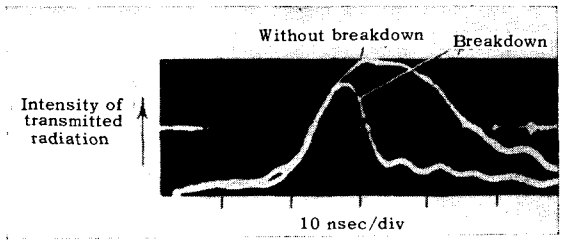

FIG. 5. Oscillograms of the total laser flux passing through the focus in the absence of breakdown and during breakdown (as measured in[5]).

Meyerand and Haught [5] measured the absorption of a laser beam at intensities above threshold. The measurements were apparently made with the aid of the generator described in the first paper of these authors. [2] Breakdown was produced in argon at atmospheric pressure and at a generator power of 40 MW, which exceeded threshold by a small amount. The diameter of the focusing circle was 2×10^{-2} cm, the length of the focusing region, according to optical calculations, was 6×10^{-2} cm. A photomultiplier was used to register the total flux through the radiation focus during the breakdown and in the absence of breakdown, with the generator power maintained the same in both measurements so as to retain fully the identical conditions. A filter was used to reduce the flux to a value below threshold. Figure 5 shows oscillograms of the total transmitted flux. We see that the absorbing plasma is produced approximately 9 nsec after the start of the pulse. By that time, approximately 10 percent of the total pulse energy has been radiated. A total of approximately 0.5 J of radiation energy was absorbed. To check that the light is indeed absorbed by the plasma, and is not scattered, the authors measured the intensity of the scattered radiation at different angles. Integration of the intensity over the solid angle has shown that only a very small fraction of the energy is scattered. The energy absorbed by the plasma goes to heating of the gas. This is confirmed, in particular, by measurements of the gas pressure after the breakdown in a vessel of fixed volume, 23 cm^3. Some time after the pulse, the gas pressure rose by 0.064 atm, corresponding to an energy of ~ 0.3 J. This agrees within the limits of accuracy with oscillograph measurements of the absorption of the radiant energy.

Interesting features of the process of absorption of the beam were observed in studying the scattering of laser light produced by the plasma and by photographing the breakdown region in the light of its own glow. This was done by Ramsden and Davis.[28] The

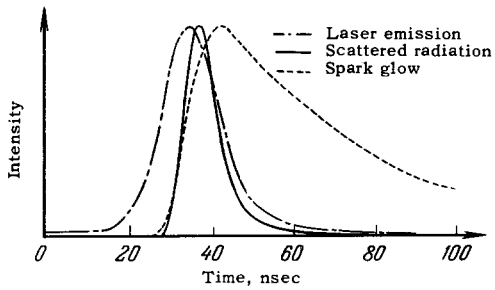

FIG. 6. Intensities of laser and scattered radiation, and also of the spark glow (as measured in[28]).

beam of a ruby laser was focused by a short-focus lens (f = 0.8 cm) into a circle of 10^{-2} cm in diameter. The breakdown in air appeared at a power exceeding 5 MW. About 60% of the pulse energy was absorbed by the plasma.

Laser radiation scattered by the plasma at right angles to the incident beam was registered. The scattered beam was gathered with a lens and focused on the slit of a spectrograph. The variation of the intensity of the scattered radiation with time was measured photoelectrically and photographically. To eliminate problems connected with the incidence of light scattered by the surfaces, the transmitted beam entered a beam stop. The scattered beam appeared only in the presence of breakdown, thus eliminating the possibility of observing scattering from unperturbed air. The scattered light was a narrow monochromatic line with wavelength close to that of the laser, but shifted up to 3 Å towards shorter wavelengths. Figure 6 shows the variation of the intensity of the scattered and laser radiations, and also the intensity of the glow of the heated air (the latter has a continuous spectrum, on which several N II lines are superimposed).

The shift of the wavelength of the scattered line is naturally interpreted as a Doppler shift connected with the motion of the scattering-plasma boundary opposite to the laser beam. The maximum shift of 3 Å corresponds to a boundary velocity of approximately 100 km/sec. This deduction was confirmed by lateral time-sweep photographs of the focal region (Fig. 7). The photograph shows how the glowing front moves opposite to the beam, and the velocity of

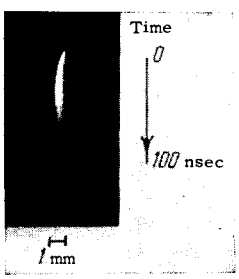

FIG. 7. Streak photograph of the process in the zone of laser-beam absorption (according to[28]).

the motion can be determined from the slope of the front line. The initial velocity is the above-mentioned 100 km/sec, and decreases with time (the Doppler shift decreases simultaneously). Sometimes one observes a weak line shifted towards longer wavelengths.

Absolute measurements of the intensity of the scattered light, with a known geometry, have made it possible to estimate the density of the electrons in the plasma, which was found to be of the order of 5×10^{19} cm^{-3}.

Within the limits of accuracy of the spectrograph resolution of 0.4 Å, the scattering line was not broadened compared with the laser line. The plasma parameters under these conditions were such that theoretically [29] the broadening of the scattering line should be determined by the ion temperature. One could therefore conclude that the ion temperatures did not exceed 10 eV.

Considerable interest, of course, should be attached to a direct determination of the temperature of the plasma produced as a result of absorption of the laser beam. Measurements of this type are described in the paper by S. L. Mandel'shtam et al.[30] The ruby-laser pulse energy, 2.5 J, was greatly in excess of the ~1 J threshold for the breakdown of air (duration 40 nsec, focusing radius 10^{-2} cm). The maximum plasma temperature was determined by measuring the intensity of the soft x radiation with $\lambda \sim 10$ Å. The electron temperature was approximately 60 eV ≈ 700,000°K. At such temperatures, the quanta with wavelength $\lambda \sim 10$ Å were in the far Wien region of the spectrum and therefore the radiation intensity was extremely sensitive to the temperature. This ensured a sufficient accuracy in the determination of the temperature, in spite of the known uncertainty in the size of the radiating volume and in the other parameters used to calculate the temperature from the measured intensity. At any rate, temperatures of 45 and 75 eV in lieu of 60 eV, are incompatible with the measured intensity for any reasonable values of the parameters.

Simultaneously with measurements of the temperature, investigations were made of the scattering of the laser beam (in a somewhat different setup than in [28]), and the results obtained were essentially the same as in [28]. The velocity of the plasma boundary moving opposite to the beam, determined from the Doppler shift, turned out to be approximately 110 km/sec.*

When the intensity of the light flux at the focus J was varied over a sufficiently wide range of values, the velocity of the plasma boundary changed slowly, approximately as $J^{1/3}$.

*Convincing proof of the Doppler nature of the shift was the fact that when observing the backward scattering towards the lens, that is, at an angle ~ 180°, the shift of the wavelength doubled compared with the shift corresponding to scattering at 90°.

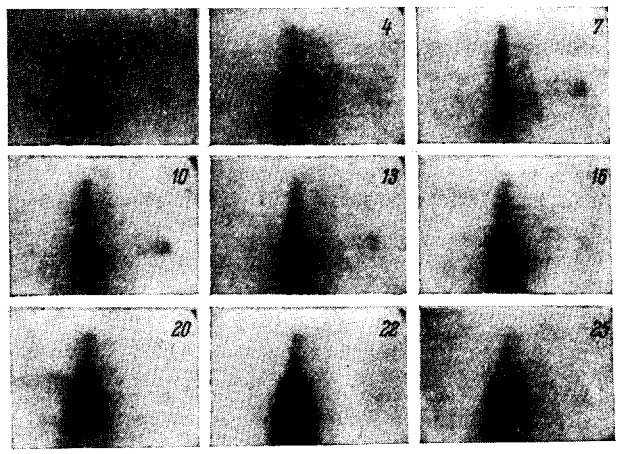

FIG. 8. Frame-by-frame photograph of the process in the region of laser-beam absorption (from[31]). The interval between successive frames is 4.4 nsec. Selected frames are shown: the numbers of the frames are indicated.

Valuable information on the dynamics of the absorption of a laser beam and the propagation of the gas-heating region is obtained by the high-speed framing photography of the phenomenon, reported by R. V. Ambartsumyan et al.[31] Photographs of the breakdown in air were obtained with the aid of a framing camera with the frames spaced 4.4 nsec apart. Figure 8 shows sample frames with serial numbers marked (negative prints). In this experiment the pulse from a Q-switched ruby laser (described in [32]) carried an energy ~3 J. The focal distance of the lens was 5 cm, the laser beam diameter 1 cm, and the divergence 10'. The pulse wave form is shown in Fig. 9. Its width at the half-peak power level was 11 nsec. The degree of blackening of the spot in the right side of the frame characterizes the laser radiation power at an instant 6 sec before the exposure of the frame. It is seen from the photographs that the breakdown develops quite unevenly in space; the determination of the boundary velocity entails certain difficulties. It can be stated, however, that it is not less than 100 km/sec. The aperture angle of the white cone in frame 7 coincides approximately with the angle of convergence

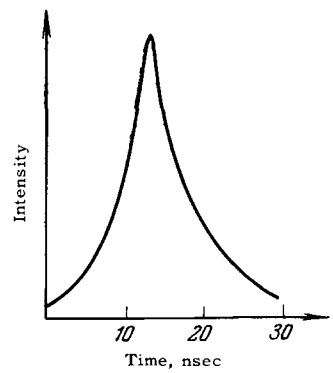

FIG. 9. Wave form of the generator pulse used in[31] (from[32]).

$10^{-5}T$, deg	Z	$10^3 l_\omega$, cm	$10^{-5}T$, deg	Z	$10^3 l_\omega$, cm
0.5	1.35	6.0	5	5.2	4.7
1	2.7	2.3	7	5.7	5.9
2.5	5.0	1.9	10	6.6	7.0

of the rays to the focus, and this, in the authors' opinion, is evidence that there is no appreciable gas-dynamic expansion of the heated gas in the transverse direction. The maximum pulse power is attained approximately in frame 8. By that time approximately 40 percent of the energy, about 1.2 J, is released.

The light cone is ~4 mm high and has a base diameter ~1 mm. Release of this energy in such a volume corresponds to an average energy of 370 eV/molecule and an average temperature of approximately 10 eV. Corresponding to the same magnitude is approximately the average velocity of the transverse expansion of the heated gas, estimated from the motion of the side boundary in frames 9—14, at a distance 4 mm from the focal point—14 km/sec. (The estimate was based on the shock-wave formula.)

7. Light Absorption and Gas Heating Wave. "Detonation" Mechanism

A general analysis of the physical process of absorption of a laser beam by the plasma resulting from the primary breakdown of air in the focal region, was presented by the author in [33]. In that paper the concept was introduced of the light-absorption and gas-heating wave, constituting the experimentally observed plasma boundary moving opposite to the beam; the shock adiabat of the wave was plotted and the wave velocities were calculated for different mechanisms of its propagation, as well as temperatures to which the gas is heated. Some problems in the theory of the phenomena described here are considered also in the article by Ramsden and Savic [34] and in papers[30,31] cited above.

When the ionization is appreciable and the temperatures reach tens and hundreds of thousands of degrees, the small quanta from the laser are absorbed in the gas as a result of free-free transitions of the electrons in the ion field. The coefficient of absorption of light by the plasma, corrected for the stimulated emission (with account of the fact that $\hbar\omega \ll kT$),[35,36]

$$\varkappa_\omega = \frac{16\pi^2}{3}\left(\frac{2\pi}{3}\right)^{1/2}\frac{Z^2e^6N_eN_+g}{(mkT)^{3/2}c\omega^2} = \frac{3{,}1\cdot10^{-31}Z^3N^2g}{T^{\circ 3/2}(\hbar\omega\,[\mathrm{eV}])^2}\ \mathrm{cm}^{-1},$$

$$g = \frac{\sqrt{3}}{\pi}\ln\left(\frac{4kT}{Ze^2N_e^{1/3}}\right) = 0.55\ln\left(\frac{2.4\cdot10^3T^\circ}{Z^{4/3}N^{1/3}}\right). \qquad (7.1)$$

Here N_e, N_+, and N are numbers of the electrons, ions, and original atoms per cubic centimeter. The ionization is assumed to be multiple so that $N_+ = N$ and $N_e = ZN$. Numerical values of the mean free path $l_\omega = 1/\kappa_\omega$ for air of normal density and for quanta with $\hbar\omega = 1.78$ eV are listed in the table. The

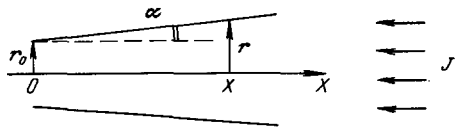

FIG. 10. Diagram of light channel in the focal region.

effective charge of the ions Z is taken from the data of [35] under the assumption of equilibrium ionization.

Let us assume that in the focal region, in the narrowest region of the converging light channel (Fig. 10), where the radiant-energy flux is maximal, breakdown took place and a high degree of ionization and a high temperature have already been attained. The light is absorbed in a very thin layer of the order of the mean free path of the quanta l_ω and heats the gas. One of the most remarkable features of this process, which is quite evident physically and which has been observed experimentally (see Sec. 6), is the motion of the absorption zone in a direction opposite to that of the light flux. In fact, the light quanta are strongly absorbed in the highly ionized medium. As soon as the degree of ionization ahead of the gas layer which absorbs at the given instant reaches for one reason or another a sufficiently high value, a new layer becomes opaque and by the same token starts absorbing. Thus, moving opposite to the beam in the light channel is an "absorption and heating wave." This effect prevents the entire pulse energy from being released in a very small volume where breakdown takes place soonest, and hampers the attainment of very high temperatures.

There are three completely different and independent mechanisms which lead to the occurrence of the absorption wave.

1. If the light flux at the focus is appreciably larger than the threshold for the breakdown, then its value exceeds the threshold also in a certain area of the light channel that expands in the direction towards the lens. The breakdown takes place also in those parts of the channel, but with a delay with respect to the narrowest spot; the larger the cross section of the channel and the smaller the flux, the longer the delay. Thus, a "breakdown wave" moves opposite to the beam. (The result of the high-speed photography of the process is interpreted in [31] on the basis of this mechanism.)

2. The heated gas in the absorbing layer expands and transmits a shock wave in all directions, including the direction along the light channel opposite to the beam. In the shock wave, the gas is heated and ionized, so that the light absorption and energy-release zone moves following the shock-wave front. This hydrodynamic mechanism is similar in many respects to the detonation of explosives.

It was noted by Ramsden and Savic,[34] who estimated the velocity of the "detonation" wave, but drew incorrect conclusions concerning the temperature to which the gas is heated, without taking into account the energy conservation law (see below).

3. The gas in front of the absorbing layer becomes ionized and acquires the ability of absorbing light because of absorption of thermal radiation from the strongly heated region of the gas (from behind the front of the absorption wave). This can be called a "radiative" mechanism.

The heating and ionization in front of the absorbing layer, connected with the electronic thermal conductivity and the diffusion of the electrons, are shown by calculation to play a small role.

The efficiency of each of the mechanisms is characterized by the rate of displacement of the absorption wave resulting from this mechanism, the real wave moving, naturally, with the largest of the possible velocities.

The absorption wave can be regarded in some sense as a hydrodynamic explosion. In the coordinate frame fixed in the wave, the process is quasi-stationary. Indeed, after a time Δt, during which the wave covers a distance on the order of its width Δx, the flux of the laser radiation and the velocity of the wave $D = \Delta x/\Delta t$ do not have time to change strongly, $\Delta x \approx l_\omega \leq 10^{-2}$ cm, $D \approx 100$ km/sec ($\Delta t \lesssim 10^{-9}$ sec).

Let us calculate the energy balance, disregarding for the time being the fact that when the gas is heated it is set in motion. The energy incident on one square centimeter of the wave surface in a time dt is Jdt, where J is the flux of radiant energy. The energy goes into heating of a mass $\rho_0 D\,dt$, which is captured by the wave during that time (ρ_0—initial gas density). Consequently, the specific internal energy $\epsilon(T)$ acquired by the gas after total absorption of the light flux, is given by the equation

$$\varrho_0 D\varepsilon(T) = J. \tag{7.2}$$

This relation expresses simply the law of energy conservation, and does not depend on the concrete mechanism of wave propagation. In a more detailed analysis of the wave as a hydrodynamic explosion, it is necessary to start from general conservation laws for the mass, momentum, and energy as the gas passes through the wave, similar to the procedure used in shock-wave theory (see [35]). As a result we obtain the equation of the "shock adiabat" of the absorption wave, which relates the pressure and the density of the gas behind the front of the wave with the initial density and the energy flux J incident on the wave. The shock adiabat of the absorption wave*,

*In spite of the common features, it differs from the shock adiabat of an explosive, in which, just as in an absorption wave, energy is released, because the energy released per gram of the explosive (q) is a constant quantity, whereas in the absorption of light the energy $J/\rho_0 D$ released per gram depends on the velocity of the wave.

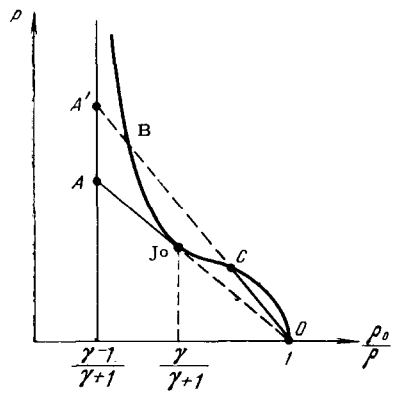

FIG. 11. Shock adiabat of the absorption wave. The straight vertical line is the shock adiabat of the strong shock wave.

calculated in [33], is shown schematically in Fig. 11.

The energy balance equation of the type (7.2) changes little when account is taken of the work of compression and of the change in the kinetic energy of the gas. The change reduces to the fact that J is replaced in the equation by a somewhat different quantity $J\beta$, with the coefficient β confined to a very narrow interval $1 < \beta \leq 2\gamma/(\gamma+1)$, where γ is the adiabatic exponent of the gas. For air at temperatures 10^5—10^6 deg, $\gamma \approx 1.33$ [35] and $1 < \beta < 1.14$, that is, the elementary energy equation (7.2) always remains valid with sufficient accuracy. This equation relates the velocity of the absorption wave and the energy of heating of the gas behind the wave, and makes it possible to estimate the temperature from the experimentally measured velocity and from the known flux J, even if the mechanism of absorption-wave propagation is not known. It must be borne in mind here, however, that an estimate based on Eq. (7.2) gives the correct result only in the case when the laser beam is practically completely absorbed by the plasma and the width of the absorption band (the width of the wave) is small compared with the diameter of the surface of the wave or of the light channel (see below).

As in the case of the shock wave, the velocity of the absorption wave D is determined by the slope of the line drawn on the pressure-specific volume diagram (Fig. 11) from the point of the initial state O to the point of the final state of the gas behind the wave. It is seen from Fig. 11 that at a given radiant energy flux J there exists a minimum possible velocity of wave propagation, corresponding to the point of the final state Jo. This is the so called Jouguet point, well known from detonation theory.[37] At this point the velocity of the wave relative to the heated matter behind the wave coincides exactly with the local velocity of sound. When all the other ionization (ignition) mechanisms are less efficient than the ionization by the shock wave, it is this hydrodynamic (detonation) mode which is realized. The gas is then compressed and heated by the shock wave to the state A, and then, after acquiring additional energy by absorption of the laser radiation, it expands along the line AJo, reaching the Jouguet point at the instant when the energy release is completed.

The maximum absorption-wave velocity is*

$$D = \left[2(\gamma^2-1)\frac{J}{\rho_0}\right]^{1/3}. \tag{7.3}$$

In this mode the heating is maximal and equal to

$$\varepsilon = \frac{\gamma}{(\gamma^2-1)(\gamma+1)}D^2 = \frac{2^{2/3}\gamma}{(\gamma^2-1)^{1/3}(\gamma+1)}\left(\frac{J}{\rho_0}\right)^{2/3}. \tag{7.4}$$

The compression behind the front is $(\gamma+1)/\gamma$ and the coefficient $\beta = 2\gamma/(\gamma+1)$.

If any of the ionization mechanisms, for example, the breakdown mechanism, produces at the given flux J a wave-propagation speed which exceeds the velocity of the "normal detonation" (7.3), then no shock wave is produced in the light channel. The gas, absorbing the light flux, goes over from the initial state O into the final state C by continuous heating and compression along the line OC, and then the increase in the pressure and density in this case are not the causes but the consequences of the wave production. In this case the wave propagates through the gas remaining behind it with supersonic velocity†.

We present a numerical example. At values $J = 2 \times 10^{18}$ erg/cm^{2}sec, $\rho_0 = 1.3 \times 10^{-3}$ g/cm^3 and $\gamma = 1.33$, corresponding to the experiments in [30], formulas (7.3) and (7.4) yield D = 133 km/sec and $\epsilon = 1.35 \times 10^{14}$ erg/g, corresponding at equilibrium to a temperature T = 910,000° (the experimental values are D = 110 km/sec and T = 700,000°).

The figures calculated in this manner agree in general with the experimental values, but they are somewhat too high, as no account was taken of the energy lost to lateral expansion of the gas. The point is that the width of the wave, that is, of the energy-release zone, is actually comparable with the radius of the surface of the wave (radius of the focusing circle), and therefore even during the stage of energy release the gas acquires radial velocities and expands through the lateral "surface" of the light channel. Taking into account the loss to lateral expansion,[33] the "effective" flux J decreases by approximately one half, so that we obtain

*Formula (7.3) can be obtained also directly from the formula for the detonation velocity $D = [2(\gamma^2-1)q]^{1/2}$ if we substitute in it the energy release $q = J/\rho_0 D$ (see the preceding footnote). This was precisely the procedure used by Ramsden and Savic,[34] but as a result of an error in the derivation of the connection between q and J, their formula which is analogous to (7.3) contains not ρ_0 but the density behind the wave.

†The hydrodynamic modes with ionization of the shock wave, but with velocity exceeding that of the detonation wave (O→A'→B) are not realized. Motion behind the wave would then be subsonic, and the expansion of the heated gas behind the wave would then attenuate the wave, bringing it into the mode of "normal detonation."

$D \approx 105$ km/sec, $\epsilon \approx 8.5 \times 10^{13}$ erg/g, and $T \approx 720{,}000°$, practically coinciding with the experimental values of [30] *.

In the paper of Ramsden and Savic,[34] in which the results of the experiments of [28] are interpreted, there is a correct estimate of the wave velocity on the basis of a formula similar to (7.3), yielding $D \sim 100$ km/sec, which agrees with experiment. However, the gas absorbing the laser beam is assigned a temperature of 40,000° which follows from a certain interpretation of the measured width of the scattering line (see Sec. 6). On the basis of the chosen value of the temperature, they estimate the absorption coefficient of the laser beam.

The authors of [34] paid no attention to the fact that the energy of gas heating and the velocity of the absorption boundary are far from independent, being connected by the energy conservation law [an expression of this is just the formula (7.4)]. The temperature 40,000° is incompatible with a velocity of ~100 km/sec, which corresponds, for an almost complete beam absorption, to a temperature on the order of a million degrees (or several hundred thousand if account is taken of the losses).

In fact, the width of the scattered line does not give the temperature to which the gas is heated after absorbing the laser beam. The question of the line width still remains unclear.[30]

It should be noted that in the experiments of [30] the authors determine from the intensity of the x radiation not the ionic but the electronic temperature. However, an estimate of the rate of energy exchange between the electrons and the ions from the known formula in [38] shows that exchange is quite rapid, under the conditions in question within a time $\sim 3 \times 10^{-10}$ sec†. At a wave velocity $D \approx 100$ km/sec, this then corresponds to an exchange length ($\sim 3 \times 10^{-3}$ cm) of the same order of magnitude as the width of the energy-release zone. Thus, there can be no great discrepancy between the electronic and ionic temperatures, such as would make the ionic temperature one order of magnitude smaller than the electronic temperature. It must be stated that the energy of the laser radiation absorbed by the electrons is transferred to the ions not only by energy exchange during collisions between these particles, but also by hydrodynamic means, via the work performed by the electron-pressure forces on the ions, since the electrons and ions are rigidly coupled by the Coulomb interaction. The latter circumstance comes clearly into play if one writes out the complete system of equations for the dynamics of a mixture of ionic and electronic gases.

*One must not have any delusions, of course, as a result of such "exact" agreement, since it undoubtedly contains a considerable element of chance.

†This value is obtained at T = 700,000°, normal air density, and five-fold ionization of the atoms. The Coulomb logarithm is taken equal to six.

8. "Breakdown" and "Radiative" Mechanism

Let us consider now the breakdown mechanism and estimate the corresponding wave velocity. As mentioned in part I of this article, the laser radiation causes an electron cascade to develop in the cold gas. The electron density increases in time in accordance with

$$N_e = N_{e0} \exp \int_0^t \frac{dt}{\theta}, \tag{8.1}$$

and at large radiation fluxes when the breakdown mechanism is the only one which plays a role, the time constant θ is determined essentially by the time required for the electron to acquire under the influence of the field an energy sufficient for excitation (or ionization) of the molecules and the atoms. In this case the rate of cascade development $1/\theta$ is approximately proportional to the light flux: $1/\theta = AJ$, $A \approx$ const.

Let us assume that the breakdown occurs when the electron density N_e reaches a certain value N_{ec}, at which the absorption of light quanta becomes sufficiently intense. This means that the instant t at which breakdown takes place in a given cross section of the light channel x (see Fig. 10) is determined the equation

$$\int_0^t \frac{dt}{\theta} = A \int_0^t J(x, t)\, dt = \ln \frac{N_{ec}}{N_{e0}} = \mu. \tag{8.2}$$

The quantity μ, which depends only logarithmically on N_{ec} and N_{e0}, will be assumed to be approximately constant.

We represent the light flux J(x, t) in the form

$$J(x, t) = \frac{W}{\pi r^2} \varphi(t),$$

where W is the peak power of the generator and $\varphi(t)$ is a dimensionless function characterizing the wave form of the pulse (see Fig. 9); $r = r(x)$ is the radius of the channel at the section x. Substituting this expression in (8.2), we get

$$A \frac{W}{\pi r^2} \int_0^t \varphi(t)\, dt = \mu. \tag{8.3}$$

If t_c is the instant of the primary breakdown at the focus, then

$$A \frac{W}{\pi r^2} \int_0^{t_c} \varphi(t)\, dt = \mu. \tag{8.4}$$

Let us assume for an estimate that the light channel in the focal region is conical and that $r = r_0 + x \tan\alpha$ (see Fig. 10).

From relations (8.3) and (8.4) follows an equation, which determines the law of motion of the breakdown

wave x(t):

$$\int_0^t \varphi(t)\,dt \Big/ \int_0^{t_c} \varphi(t)\,dt = \left[1 + \frac{x}{r_0}\operatorname{tg}\alpha\right]^2. \qquad (8.5)^*$$

When the power is appreciably in excess of threshold, breakdown usually occurs even before peak power is attained. For convenience in estimating by means of (8.5), we approximate the curve $\varphi(t)$ during the power-growth stage by a straight line. We measured the time from the point of intersection of the line with the abscissa axis, as if the pulse had a triangular form. Then $\varphi(t) = \text{const} \cdot t$. Extrapolating the line to zero power, we obtain from (8.5) the velocity of the breakdown wave during the stage of generator power growth:

$$x = D(t - t_c), \quad D = \frac{dx}{dt} = \frac{r_0}{t_c \operatorname{tg}\alpha}. \qquad (8.6)$$

In the experiments of [30] $r_0 = 10^{-2}$ cm, $t_c \approx 10$ nsec, $\tan\alpha \approx 0.1$, and $D \approx 100$ km/sec, that is, the breakdown wave velocity is close to the hydrodynamic velocity (see Sec. 7).

In the experiments of [28] $r_0 \approx 4 \times 10^{-3}$ cm and $t_c \approx 7$ nsec, $\tan\alpha = 1$ (short-focus lens). The velocity of the breakdown wave is $D \approx 6$ km/sec and is much smaller than the hydrodynamic velocity, which, as in [30], is ~ 100 km/sec, that is, in these experiments the breakdown mechanism certainly plays no role. Let us estimate the dependence of the initial velocity of the breakdown wave on the peak power of the generator W, the pulse duration Δt, and the geometry. Assuming as an estimate $\varphi(t) = \text{const} \cdot t/\Delta t$, we obtain from (6.4) that $t_c \sim \Delta t^{1/2} W^{-1/2} r_0$. Formula (6.6) then yields

$$D \sim \frac{W^{1/2}}{\Delta t^{1/2} \cdot \operatorname{tg}\alpha}. \qquad (8.7)$$

In the case of short powerful pulses and long-focus lenses (small α) the initial velocity of the breakdown wave can be very large, of the order of several hundred km/sec, and may greatly exceed the hydrodynamic velocity. A similar effect is brought about also by the difference in the form of the light channel from conical, the "sharpening" of the channel in the direct vicinity of the narrowest spot (decrease in the convergence angle α). In this case the elongated region of the channel is traversed by the breakdown wave, and then, when the channel begins to expand, the breakdown mechanism can give way to another mechanism, for example, the hydrodynamic mechanism.

In [31] a formula is derived for the breakdown propagation which is somewhat different from (8.6). The authors assume that the breakdown in this section of the channel occurs at the instant when the intensity of the light flux in this section, which grows together with the growth of the generator power, reaches a definite threshold value. The difference from the previous conclusion, consequently, consists in assuming that the cascade develops "without inertia." Although formula [31] gives a true estimate of the velocity (~ 200 km/sec for the experimental conditions of [31]), nonetheless it is more correct to take into account the accumulation of the electrons in the cascade during the time of action of the field, as was done above in the derivation of (8.6).

Let us dwell briefly on the radiative mechanism of wave propagation.

At temperatures of the order of several hundred thousand degrees, the free paths of the quanta with energy $\sim kT$, radiated by the heated gas, are $l \sim 10^{-1}$—10 cm and are considerably longer than the free paths of the small light quanta ($l_\omega \sim 10^{-3}$—10^{-2} cm), the width of the wave, or even the characteristic dimensions of the heated region. The heated gas is transparent to the thermal radiation and emits from its entire volume. This radiation is absorbed in the colder layers, where the ionization is small; the corresponding mean free paths in atmospheric air are $\sim 10^{-2}$—10^{-1} cm (for $\hbar\omega \sim 20$—200 eV). As soon as the ionization due to the absorption of the thermal radiation reaches a sufficient value (this occurs at temperatures $\sim 15{,}000$—$20{,}000°$), the new ionized layer begins to absorb the light flux intensity. The intensity of the focused flux of laser radiation is much larger than the flux of thermal radiation, so that the new layer becomes rapidly heated and the boundary of the high-temperature zone moves towards it.

To calculate the propagation velocity of this boundary it is necessary to consider the stationary conditions in a coordinate system in which the wave is at rest, as is done in the theory of the shock-wave front structure.[35] In this case doubts may arise concerning the possible existence of a stationary mode, for during the course of time, as the wave moves, the volume of the heated and emitting gas, which remains behind it increases and one might think that the flux of thermal radiation and the velocity of the wave increase. Actually this does not take place because the light channel has a finite diameter. In fact, the progress of the wave is influenced always only by the radiation produced in the layer behind the wave, the thickness of which is of the order of the channel diameter; the radiation from the remaining volume is beyond the limits of the light channel. This is what ensures a quasi-stationary behavior of the propagation process. Approximate calculations of the velocity of the radiation wave, made in [33], show that, both in magnitude and in the dependence on the laser power, the velocity of the radiative mode is close to the hydrodynamic velocity. Thus, for the experiments of [30] one obtains $D \approx 95$ km/sec. This means, that in the absence of a breakdown wave it is

*tg = tan.

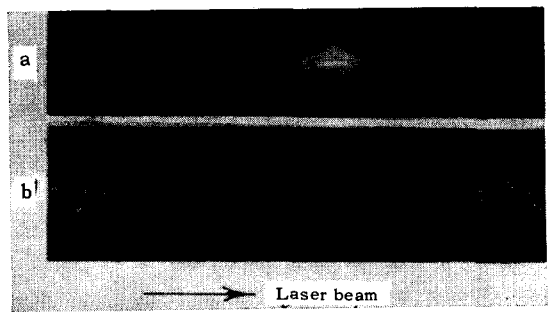

FIG. 12. Photographs of the spark (a) and of the laser beam in the absence of breakdown (b). The light channel is partially seen because of the scattering of light by particles specially introduced into the air.

most likely that the radiative and shock waves are "tied" to each other and move together, and the wave velocity can be calculated from the formulas applied to either mechanism. The radiation wave serves in this case as a "tongue" which heats the gas ahead of the front of the shock wave, as is the case in ordinary shock waves of very large amplitude.[35] The temperatures in the heating zone are relatively low, ~20,000°.

It is possible that the scattering of the laser light in this region is the cause of the experimentally observed width of the scattering line. Along with this assumption concerning the smallness of the scattering-line broadening, another assumption is advanced in [30], that the "scattered" line is the result of regular reflection from the curved shock-wave front. (Owing to the presence of lateral expansion of gas, the front of the shock wave becomes convex in the direction of motion, similar to the front of a detonation wave in cylindrical explosive charges of small diameter.) It must be stated that the question of the reasons why the scattering line has little broadening remains unclear.

In concluding the discussion of the process of absorption of the laser beam, we must emphasize that the notions presented above give apparently only the main scheme of the process; they can serve as a working model in the interpretation of the experimental data and in estimates of the expected results. The real process is much more complicated. In particular, an irregularity in the propagation of the plasma boundary was observed in the experiments.[30,31] As regards the role of the different mechanisms under different conditions, in experiments [28] with a short-focus lens the "breakdown" mechanism is certain not to be effective and the wave is most likely to be hydrodynamic. For the experiments of [30], all three mechanisms give nearly equal velocities, which does not contradict the experimental data. It is not excluded that all three mechanisms "operate" simultaneously, and that the velocities are not "additive" in this case.

In the case of long-focus lenses, which produce light channels with small convergence angle in the focusing region, and in the case of powerful short pulses, the wave is "guided" by the breakdown mechanism during the early stage of the beam absorption. However, after some time, the hydrodynamic or the radiative mechanisms must come to the forefront.

Even without knowing the exact mechanism of wave propagation, it is always possible to estimate the gas temperature from the measured wave velocity, using Eq. (7.2).

9. The Spark as a Strong Explosion. Spark in a Magnetic Field

At the instant of termination of the laser pulse, a strongly heated volume is produced in the gas, in a shape that is elongated along the laser-beam direction. The length of this region obviously is of the order of the path traversed by the absorption wave, that is, of the order of the product of the wave velocity by the duration of the pulse, equal to several millimeters. The radial dimensions are determined by the dimensions of the converging light channel and constitute several tenths of a millimeter, and by the instant of termination of the pulse the layers of the gas, heated at the start of the pulse, have time only to expand laterally, sending a shock wave into the surrounding air. The initial velocities of the radial expansion of the heated gas and the initial values of the radial component of the shock-wave velocity, which coincide with them approximately, are equal in order of magnitude to the speed of sound in the heated gas. These quantities are two or three times smaller than the velocity of the absorption wave, that is, of the order of several times 10 km/sec in those sections of the light channel where the energy release is intense. (By the end of the pulse, the absorption wave traveling through the expanding channel with decreasing light flux, is of course slowed down.) In the course of time, the radial velocities of the shock wave decrease, as is the case for a cylindrical explosion.

Figures 12 and 13 show photographs of a spark in air, taken at 625,000 frames per second by S. L. Mandel'shtam et al.[40] The time interval between frames is 1.6 μsec = 1600 nsec, that is, a prolonged

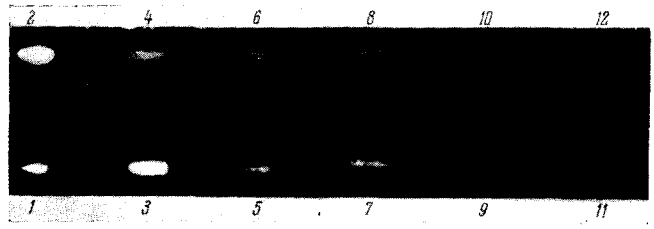

FIG. 13. Frame-by-frame photographs of the spark. The frames are numbered. The interval between the frames is 1.6 μsec.

stage has been photographed, compared to which the energy release in the gas (~ 1 J) is instantaneous.

The process during this prolonged stage has much in common with a strong explosion in air. From the point of energy release in the air, a shock wave propagates in all directions, and attenuates with time. Its surface is not fully spherical, since the energy release took place in a volume of elongated form. However, during the course of time the form of the surface of the shock wave undoubtedly becomes nearly spherical. This is not seen in the photographs since the boundary of the glowing region can coincide with the shock-wave front only at the very start of the explosion process. All that glows is the highly heated air; therefore, starting with a certain instant, when the amplitude of the shock wave decreases to a value insufficient for strong heating, the wave becomes invisible and travels forward, leaving behind it a strongly-heated glowing region of smaller dimensions. Gradually, as the air becomes cooler, this glow also attenuates. All these phenomena are well known and were considered in detail in a book by Ya. B. Zel'dovich and the author,[35] as applied to large-scale explosions.

In [40] the spectrum of the glow of the spark was plotted. Lines of singly-ionized nitrogen and oxygen atoms, which are strongly broadened, are superimposed against a strong continuous background. An estimate of the temperature and of the electron density over the spectral lines yielded values $T \sim 30{,}000$—$60{,}000°$ and $N_e \sim 3 \times 10^{18}$ cm^{-3}. These values were obtained from spectra which were summed in time, and consequently pertain to the prolonged and developed stage of the spark. As noted in Sec. 6, the final gas pressure following the production of a spark in a small closed volume was measured in [5]. The increase in pressure corresponded to the amount of energy released.

G. A. Askar'yan et al.[42] investigated the plasma cloud of an explosion produced in air as a result of absorption of a laser pulse, using microwave radio emission methods. The spark was produced between the antenna of a radio generator at 0.8-cm wavelength and the receiving horn of a detector. The signal transmitted through the plasma and reflected from it were registered simultaneously. To estimate the parameters of the plasma cloud with which the radio waves interacted, the cloud was simulated by bodies of different dimensions, made of crushed metal foil, and a body was chosen from which similar signals were obtained. It turned out that the dimensions of the plasma are of the order of a centimeter, which is much larger than the radii of the shock wave during the stage when it still leads to noticeable thermal ionization. The reflection of the radio signal lasts for hundreds of microseconds, attesting thus to the fact that during this time the electron density remains above the critical value $\sim 10^{13}$ cm^{-3}. Reflection is observed also in times shorter than a microsecond, when the shock wave has not had time to reach a radius corresponding to the measured dimensions of the reflecting region (the maximum reflection is observed at the instant ~ 5 μsec). The authors advanced the hypothesis that ionization in a volume exceeding the dimensions of the volume which can be ionized by the shock wave is connected with the action of short-wave radiation from the highly heated central zone.

In a paper by the same authors [41] they report investigations of a spark produced in the presence of an external magnetic field. The magnetic effects were predicted in a paper by G. A. Askar'yan and M. S. Rabinovich.[23] Circular induced currents were produced during the course of expansion of heated ionized air behind the shock wave propagating from the point of energy release in the presence of an external magnetic field H. The crowding out of the magnetic field from the small volume of the plasma is insignificant, according to estimates, so that the change in the external magnetic field can be neglected.

If we assume for simplicity that the plasma is scattered in spherically symmetrical fashion and if we introduce polar coordinates r, ϑ and φ with axes directed along the magnetic field, we can write for the electric induction field

$$E_\varphi(r, \vartheta, t) = \frac{1}{c} v_r(r, t) H \sin \vartheta,$$

where v_r is the radial scattering velocity. The magnetic moment of the circular current, directed opposite the external magnetic field, is

$$M = -\frac{4\pi H}{3c^2} \int_0^{R(t)} \sigma(r, t) v_r(r, t) r^3 dr,$$

where σ is the conductivity of the plasma, and R(t) the radius of the expanding plasma. Under the assumption that the scattering velocity depends linearly on the radius, $v_r = r_R r/R$, we get

$$M = -\frac{4\pi}{15c^2} v_R(t) H \bar{\sigma}(t) R^4(t),$$

where v_R is the velocity of the boundary of the plasma sphere and $\bar{\sigma}$ is some average conductivity. When $v_R \sim 30$ km/sec, $R \sim 10^{-2}$—10^{-1} cm, and $\bar{\sigma} \sim 10^{14}$ sec^{-1}, the magnetic moment is approximately 10^{-1}—10^{-2} of the moment of an ideal diamagnetic sphere having the same dimensions $M_0 = -R^3H/2$.

Owing to the time variation of the magnetic moment, the following emf should be induced in the turns surrounding the plasma sphere:

$$\mathcal{E} = \frac{2\pi}{c} \frac{n}{\rho} \frac{dM}{dt}$$

(n is the number of turns and ρ their radius). For $n \sim 100$, $H \sim 1$ kOe, and a characteristic time of the process of 10^{-8} sec, we can expect an emf on the

order of 1 volt. The recorded signal can yield information on the scattering of the plasma volume.

In the experiments of [41] the coils used had n $n \sim 2-20$ and $\rho \sim 1$ cm. The magnetic fields reached 10 kOe. A time sweep of the diamagnetic signal has shown that two pulses of opposite polarity are produced, with durations of a fraction of a microsecond separated by a time interval of several microseconds. The first signal corresponds to the appearance of a diamagnetic moment in the plasma cloud, the second corresponds to the vanishing of the currents. During the time interval between the signals, the induction currents apparently exist in the plasma. The spark glows during the entire time of existence of the diamagnetic moment. Integrating the coil voltage with respect to time, we can estimate the magnetic moment. It turns out that it reaches its magnitude $M \approx 10^{-6}$ H cm^3 even within the time of the first pulse, 0.3 μsec. The nature of the long existence of the plasma diamagnetism is still not clear.

The diamagnetism of the plasma makes it possible to act on it by means of an inhomogeneous magnetic field in order to accelerate the plasmoid [the total force acting on the plasma is $F_z = M(\partial H/\partial z \sim R^3 (\partial H^2/\partial z)]$, and use a spark and a magnetic field to obtain fast jets of pure dense plasma, to fill traps, to make plasma more compact, etc.

10. Production of Plasma from a Solid Target

It is possible to heat matter to very high temperatures with a laser beam not only by producing a spark in a gas, at a generator power noticeably exceeding the threshold value for the breakdown. A high-temperature plasma can be obtained also by focusing a giant pulse on the surface of an opaque solid (or liquid) substance. In this case, incidentally, there is no absorption threshold as in transparent gases; the light is absorbed by the opaque body at any intensity. The indicated method of obtaining a plasma has its advantages compared with the formation of the plasmoid in a gaseous medium. A solid target can be placed in a vacuum. This, in principle, makes it possible to guide to the investigated material an arbitrarily strong light pulse without the danger of screening the substance as a result of the breakdown in the surrounding gas medium. When the target is placed in a vacuum, the plasma produced under the influence of a laser pulse, expanding in vacuum, can be accelerated to high velocities. It is not our purpose here to consider in detail the mechanism of the processes which occur during the focusing of a giant pulse onto a solid target. We only mention papers on this subject for the purpose of information.

Ready [50] focused a giant pulse from a ruby laser on a carbon target placed in atmospheric air. The resultant bright flare was photographed at a frequency of 10^8 frames per second for 800 nsec. The rate of expansion of the flare, which could be determined from the photographs, was approximately 20 km/sec. The dimensions of the glowing region on the last frame were on the order of several millimeters.

Linlor [51] placed targets made of several elements, from the lightest ones to the heaviest ones, in a vacuum chamber (10^{-6} mm Hg). The pulse parameters were 0.2 J, 40 nsec, and 5.4 MW. He registered the time interval during which the plasma covered a known distance 4.3 cm from the point of beam focusing to the collector, and thus determined the velocity. The time of flight of the ions depended on the atomic number of the target material A approximately as $\sqrt{A}$. This offers evidence that the energies of ions having different masses were approximately equal. At a pulse power of 5.4 MW, the energies amount to ~ 1 kev, and at double that power—420 eV. Analogous experiments were made by Opower and Burlefinger.[55] They measured the time of flight of the ions of carbon from a carbon target placed in high vacuum. The pulse energy was five times larger than in the experiments of [51], ~ 1 J. By means of a small transverse electric field, electrons were drawn from the layers of the moving plasma farthest to the front, and consequently of lowest density (the field does not act on the denser plasma, since the removal of electrons is hampered by the resultant space charge). The forward boundary of the moving ions was registered. The ion front was very sharp. At a laser pulse of 1 J and a power of 30 MW, the velocity of the leading front was 130 km/sec, corresponding to an ion energy of ~ 1 keV, at 0.5 J and 10 MW the corresponding value was 85 km/sec. According to estimates based on the known solid angle subtended by the ion detector, assuming that the plasma propagates from the target isotropically, for a pulse of 1 J 10^{14} ions have energies from 1 to 1.1 keV. The initial plasma temperature, estimated from the ion energy, is $\sim 10^6$ deg. It must be noted that the foregoing estimate highly overvalues the temperature, since the authors, on going from the kinetic energy of the ions to the initial internal energy of the plasma, disregarded the energy lost to detachment of the electrons (only the thermal energy of the electrons was taken into account), and did not take into account the fact that the boundary of the gas that propagates in vacuum is accelerated to velocities that are larger than the average propagation velocity corresponding to the initial internal energy (see [35]).

Linlor [54] measured the absorption of a giant laser pulse, focused on gold and aluminum foils 1.8×10^{-4} and 1.6×10^{-4} g/cm^2 respectively. In the first case the plasma absorbed 94% of the incident light, and in the second 99%. The ion energy determined from the time of flight over a known distance was approximately 1 keV. He also registered the diamagnetic signal in the presence of a magnetic

field.

Archbald et al.[53] carried out a spectrographic investigation of the plasma from targets made of different metals (laser power 1 MW, pulse duration 100—500 nsec). It is noted that the craters produced on the surface of the target at the focal point had a diameter on the order of 10^{-2} and a depth of several microns.

Spectral measurements during the focusing of a neodymium-glass laser on a solid target containing lithium and placed in vacuum were made in the already cited paper[31] (pulse energy 3 J, duration 40 nsec). It was possible to distinguish against the background of the continuous plasma glow the lines of the lithium atoms and ions. The temperature, according to estimates based on the spectral-measurement data, was $\sim 230{,}000°$.

Neuman[52] measured the momentum acquired by different targets. The momentum, corresponding to the reflection of the laser beam, is only 2×10^{-4} dyne-sec, whereas the measured momentum for copper was 0.18 dyne-sec. This is connected with the recoil resulting from the evaporation and scattering of matter from the surface. This effect, which should be observed during evaporation of a surface by means of a laser beam, was called attention to by G. A. Askar'yan and E. M. Moroz[56]; it is perfectly analogous to the effect produced when very fast micrometeorites strike against a surface. There have been many studies of the action produced on the solid target by focused radiation from a laser operating in the usual (not giant-pulse) mode. In this mode, the duration of the pulse is relatively large, $\sim 10^{-3}$ sec, and the powers are much smaller than obtained with Q-switching. The material of the surface, when such a pulse is focused on solid targets, also evaporates, but no plasma is produced, since the energy release is too slow and the temperature is too low. The thermionic and ionic emissions from surfaces, as well as other effects were investigated. References to these investigations can be found in[51].

11. Conditions for Strong Heating of Hydrogen

Of great interest is the fundamental question whether hydrogen (deuterium) can be heated by laser radiation to thermonuclear temperatures $\sim 10^7$ deg. Estimates of various effects which exert an influence on the heating, and the laser power necessary for this purpose, were made by N. G. Basov and O. N. Krokhin.[43] These calculations were subsequently repeated in somewhat expanded form, by Dawson.[57] The hydrogen plasma absorbs light quanta as a result of free-free transitions of the electrons in the field of the ions, and the plasma density should not exceed 3×10^{21} cm^{-3}, which is the critical value for the reflection of light from a ruby laser. At a temperature 10^7 deg and a density $\sim 3 \times 10^{21}$ cm^{-3}, the coefficient of absorption of the light quanta (which decreases with increasing temperature) is $\sim 10^2$ cm^{-1}, the mean free path of the quanta is $\sim 10^{-2}$ cm, that is, of the order of the focusing radius, and the plasma is still opaque. Consequently, in principle heating by laser light to such high temperatures is possible. However, the heating is hampered by energy loss due to the electronic thermal conductivity (if the plasma is situated in a medium) and by hydrodynamic scatter of the plasma (the radiation losses are less significant). A particularly important role is played by expansion, as a result of which the plasma escapes from the region of action of the laser beam. For effective heating it is necessary that during the time Δt of the laser pulse, the energy of which is sufficient to heat the plasma within the focusing region to the required temperature $T \sim 10^7$ deg, the plasma be unable to expand strongly.

Let us estimate very roughly the generator power necessary for this purpose. The rate of scattering of the plasma is of the order $v \sim (kT/M)^{1/2}$, where M is the mass of the atom; the characteristic expansion time is $\tau \sim r_0/v$, where r_0 is the focusing radius. If the laser power is W and the initial density of the atoms is N, then when all the energy is released in the mass contained in the focusing volume, the plasma will be heated to a temperature $kT \sim W\Delta t/(4\pi/3)r_0^3 N$. In order for the energy actually to be released in this mass, the gas dynamic time τ should exceed the pulse duration: $\tau > \Delta t$. Thus, the generator power should exceed $W > (4\pi/3)r_0^2 N (kT)^{3/2}/m^{1/2}$. For $r_0 \sim 10^{-2}$ cm, $N \sim 3 \times 10^{21}$ cm^{-3} and $T \sim 10^7$ deg, a power on the order of 10^4 MW is necessary. These estimates concern, naturally, only the fundamental aspect of the phenomenon and do not solve all the difficulties of concrete realization of such an experiment. Of very great importance here is a search for an optimal experimental setup and effective plasma heating conditions.

One of the heating modes, in which the dimension of the region occupied by the plasma is close during all stages of the process to the mean free path of the light quanta, this being the condition for effective transformation of the radiation energy into heat, is proposed and considered by O. N. Krokhin.[44]

We make mention of an article by I. V. Nemchinov,[45] in which a problem is solved which is close in scope to the questions considered here, the problem of planar scattering of a heated gas layer. We note also a paper by Caruso and Gratton.[61] The latter paper contains an estimate of the effect of compression of a plasma produced as a result of absorption of laser light by light pressure.

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