INVESTIGATION OF THE SPARK DISCHARGE PRODUCED IN AIR BY FOCUSING LASER RADIATION. II

S. L. MANDEL'SHTAM, P. P. PASHININ, A. M. PROKHOROV, Yu. P. RAĬZER, and N. K. SUKHODREV

P. N. Lebedev Physics Institute, Academy of Sciences, U.S.S.R.

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Results are presented of an experimental investigation of the initial stage of the discharge in air induced by focusing radiation from a laser. The temperature of the plasma produced in the region of the focus is determined on the basis of the recombination radiation spectrum in the soft x-ray region, $\lambda \approx 10$ Å, and is found to be $T_e = 50-60$ eV. A study of the laser radiation scattered by the plasma indicates that the ionization front is moving toward the focusing lens. The velocity of the front, $v \approx 10^7$ cm/sec, has been measured from the Doppler shift of the scattered light. A hydrodynamic mechanism of broadening of the ionization region is considered. Values for the velocity of the detonation wavefront and the plasma temperature behind the front, estimated on the basis of this mechanism, are in satisfactory agreement with those found experimentally.

A number of recent experimental and theoretical studies[1-8] have been devoted to the study of the "spark" arising in a gas on focusing the radiation from a high-power laser. In a previous article[4] we have reported preliminary results of a study of the plasma of the spark occurring at the focus of a lens. Spectroscopic measurements on the NII lines gave values $T_e \approx (50 - 60) \times 10^3$ °K and $N_e \approx 2 \times 10^{18}$ electrons/cm^3. The spectrum was obtained without a time sweep, and these data characterize some average parameters of the plasma in the second stage of the spark, which lasts several tens of microseconds—after the action of the laser beam has stopped.

In the present communication we report results of further studies of the plasma, which characterize the plasma in the first stage of the process—during the period of action of the laser beam.

1. EXPERIMENT

A pulsed, Q-switched ruby laser was used with an energy of 2—2.5 joules and pulse length of about 40 nsec; the principal characteristics of the laser are given in previous articles.[4,9]

To estimate the maximum electron temperature of the plasma, we studied the soft x-radiation of the plasma in the wavelength region near 10 Å.

The measurements were carried out by means of photon counters with aluminum and beryllium windows of respective diameters 3 and 8 mm. The parameters for the aluminum-window counters are given by Tindo and Shurygin[10] and for the beryllium window counters by Mandel'shtam et al.[11] The counters operated in the Geiger region with a time constant of about 10^{-5} sec. Consequently, only the presence or absence of the x-ray photons could be recorded by the counters. To estimate the intensity of the radiation, aluminum and beryllium foils of known absorption[12] were added in front of the counters. With a pulse energy of 2.5 joules the aluminum-window counter still recorded photons at a distance of 25 mm from the spot with a 40 μ aluminum foil in front of it; increasing the foil thickness by 5 μ or increasing the distance from the spark resulted in disappearance of the signal. Similar conditions were observed for the beryllium-window counter on placing a 70 μ beryllium foil in front of the counter, for a distance of 12 mm.

Figure 1 shows spectral sensitivity curves of

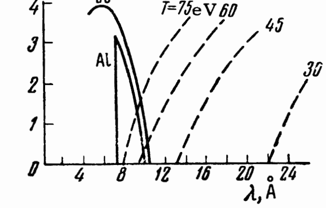

FIG. 1. Spectral sensitivity curves of counters with aluminum and beryllium windows, corresponding to the threshold for appearance of a signal under the conditions of the experiment, and theoretical curves for the intensity of recombination radiation for different temperatures (broken line).

the counters, taking into account the additional foil and the absorption of the air for the conditions given above.

According to Elwert[13] the spectral intensity of the bremsstrahlung and recombination radiation of a plasma per unit volume per unit solid angle per unit time can be described in the form

$$\varepsilon_\nu = 6.36\cdot 10^{-47} Z_i^2 \frac{N_e N_i}{(kT)^{1/2}} \exp\left(-\frac{h\nu}{kT}\right) \times\left[1+\sum_p \frac{N_{i+p}}{N_i Z_i^2}\frac{\chi_H}{kT}\sum_n\left(\frac{\chi_{i+p,\,n}}{\chi_H}\right)^2 \times\frac{\xi_n}{n}\exp\left(\frac{\chi_{i+p,\,n}}{kT}\right)\right] [\mathrm{erg/cm^3}], \tag{1}$$

where N_e and N_i are the concentration of electrons and ions with charge Z_i; χ_H is the ionization potential of hydrogen; $\chi_{i+p,n}$ is the ionization potential for an ion with charge $i + p - 1$ from a level with principal quantum number n; ξ_n is the number of free states in the n shell (the Gaunt correction factor is taken as unity). The size of the radiating volume was taken equal to $(1/4)\pi d^2 l_p$, where d is the transverse dimension of the focal region and l_p is the mean free path of a photon with the ruby laser frequency ν_p in a plasma with temperature T. According to Spitzer[14] the latter quantity is

$$l_p^{-1} = \varkappa_p = 3.69\cdot 10^8 \frac{N_e N_i Z_i^2}{T^{1/2}\nu_p^3}\left[1-\exp\left(-\frac{h\nu_p}{kT}\right)\right] = 10^{-31}\frac{Z_i N_e N_i}{T^{3/2}}\ [\mathrm{cm^{-1}}]. \tag{2}$$

Note that when the quantity l_p is used as the thickness of the radiating layer in the expression for the absolute flux of radiation, the quantities N_e and N_i, whose absolute values are not determined from the experiment, cancel out.

The main contribution to the radiation from the air in the soft x-ray region for temperatures of tens of electron volts is given by the recombination term in (1), i.e., Σ_p. In the spectral region of interest to us, where the photon energy of ~ 1 keV exceeds the ionization potentials, the principal role is played by recombination to the ground state n = 1, and in Eq. (1) we can limit ourselves to the first term in Σ_n.

Substituting (2) into (1), setting $d = 2 \times 10^{-2}$ cm, the duration of the radiation $t = 4 \times 10^{-8}$ sec, converting to the number of photons emitted per unit solid angle per unit wavelength interval (in Angstroms), and taking into account everything we have said above, we obtain finally

$$N_\lambda = 1.4\cdot 10^{13}\frac{1}{\lambda}\exp\left(-\frac{12396}{\lambda T}\right)\sum_p \frac{N_{i+p}}{N_i Z_i^2}\left(\frac{\chi_{i+p}}{\chi_H}\right)^2\frac{\xi_n}{n} \times\exp\left(\frac{\chi_{i+p}}{T}\right) \tag{3}$$

(here λ must be taken Angstroms, and T in electron volts).

The values of the concentrations N_{i+p} of nitrogen and oxygen ions in air for different values of T_e were calculated according to Zel'dovich and Raĭzer[15,16] 1) and are listed in the table together with the values of χ_{i+p}, ξ_n, and n necessary for the calculation.

Ion	n	ξ_n	χ_{i+p}, eV	Ion concentration			
				$T = 40$ eV	50 eV	60 eV	70 eV
N^{5+}	2	5	97.86	0,77	0,73	0.465	0.165
N^{6+}	1	1	551.9	0.003	0.06	0.3	0.54
N^{7+}	1	1	666.8	—	10^{-4}	0.008	0.08
O^{6+}	2	6	138.08	0.195	0.197	0.195	0.163
O^{7+}	1	1	739.1	$7.4\cdot 10^{-6}$	$3.8\cdot 10^{-4}$	0,006	0.036
O^{8+}	1	2	871.1	—	—	$5\cdot 10^{-6}$	$3\cdot 10^{-4}$

The ion concentrations were determined as the ratio of the number of ions of a given type to the number of parent atoms per cm³ for normal conditions $N_0 = 2 \times 2.7 \times 10^{19}$ cm⁻³. Figure 1 shows the results of calculation of the radiation flux according to Eq. (3) for different temperatures. The appearance of a signal in both counters under the conditions given above corresponds to $T_e \approx 60$ eV if we take as a limit the presence of one pulse.

In spite of the large uncertainty in choice of the size of the radiating volume, the criterion for appearance of a signal, etc., the evaluation of T_e turned out to be rather accurate, as the result of the very steep behavior of the intensity of radiation in the spectral region of interest to us. Thus, for example, the values $T_e = 45$ eV and $T_e = 75$ eV are inconsistent with the experimental data obtained.

To obtain further information on the state of the plasma in the first stage of the spark, we studied the width and shift of the laser line scattered by the plasma. The observation was made perpendicular to the direction of propagation of the laser beam. The focal region of the laser beam was imaged with a three-fold magnification on the spectrograph slit by the objective L_2 (see Fig. 2). We used an STÉ-1 equipment with a dispersion of 13.2 Å/mm in the region of 6900 Å.

1)The calculation was carried out by V. S. Prokudina.

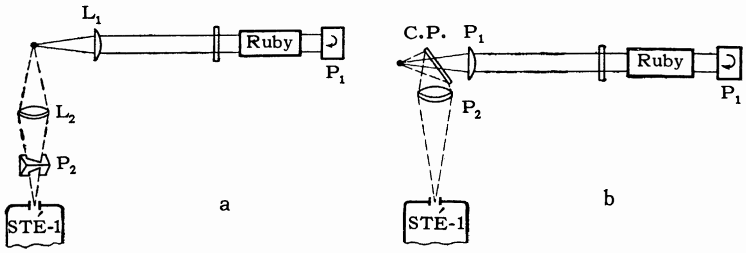

FIG. 2. Diagram of apparatus for observation of laser light scattered by a plasma: a – at an angle of 90°, b – at an angle of 180°.

The line image was photometered and the line contour in intensity units was determined from the density. The density was determined by means of an absorber placed at the slit, and the slit was uniformly illuminated by the laser beam scattered by a matte plate. The measured line width turned out to be of the order of 1.5 Å. The sum of the instrumental width of the apparatus and the line width of the radiation from the laser itself amounted under our conditions to about 0.4 Å. Thus, the broadening of the line scattered by the plasma amounts to about 1 Å for arithmetic combination of the line widths and the instrumental width and ≈ 1.4 Å for quadratic combination.

In addition to the broadening, a shift of the line is observed both in the short-wavelength and long-wavelength directions, of up to about 3 Å. This shift was also observed by Ramsden and Davies[7] and was interpreted by them as a Doppler shift. To study this phenomenon in more detail, we (and also Ramsden and Davies) imaged on the spectrograph slit a region near the focus of the lens, having rotated the spark image by 90° by means of the prism P_2 (Fig. 2a). With this arrangement the different regions along the laser beam near the focus were imaged at different heights along the spectrograph slit.

To obtain the undisplaced position of the line we used the part of the laser beam scattered by the matte plate. The characteristic form of the spectrograms obtained is shown in Fig. 3. As we can see from Fig. 3a, in the different regions of the caustic curve near the focus, the shift is different. In the regions of the caustic located close to the lens L_1, the shift occurs in the short-wavelength direction, gradually decreasing with increased distance from the center of the caustic; the shifted line is still observed at a distance of 1–2 mm from the focus. In the region beyond the focus the shift occurs in the long-wavelength direction and can be followed for 0.3–0.5 mm. The maximum value of the shift, both in the red and in the violet directions, is observed near the focus and amounts to about 3.2 Å. It is very characteristic that the displaced line is not of a continuous nature but changes its intensity noticeably along its length. In a spectrogram taken with a shorter exposure (Fig. 3b), we can see that the displaced line appears to consist of individual segments and beads of length 0.1–0.3 mm; in individual cases the segments are somewhat inclined to the average line or have a small bend.

At a distance from the focal region in the direction of the lens of 1.5–2 mm a second scattering region is observed, easily noticeable also in photographs obtained with an SFR apparatus in the form of a second region of breakdown (Fig. 4). The shifted line is also visible in this region; the magnitude of the shift is considerably smaller and amounts to ≈ 1.5 Å.

All of the data given above refer to a lens with a focal distance of 55 mm. The picture remains qualitatively the same (except for the appearance of the second luminous region) for lenses with focal lengths of 25 and 75 mm.

To verify that the shift is due to the Doppler effect, we changed the experimental setup so that it was possible to observe light scattered by the plasma in the direction of the focusing lens L_1 (Fig. 2b). The magnitude of the shift observed in this case turned out to be twice as large as in the preceding experiment, which confirms the assumption of the Doppler nature of the shift.

To determine the dependence of the maximum shift $\Delta\lambda$ on the intensity of the laser radiation J we carried out a series of experiments in which the shift of the scattered line was measured for

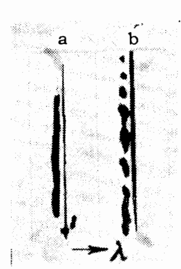

FIG. 3

FIG. 4

FIG. 3. Spectrogram of laser light scattered by a plasma: a – long exposure, with a shift visible toward the red near the focus; b – short exposure, showing the discontinuous structure of the shifted line.

FIG. 4. Photograph of the initial stage of the spark with an exposure of 1.6 μsec.

different values of laser power. The focal distance of the lens L_1 was 25 mm in these experiments. The laser operated under the same conditions, and different values of J were obtained by attenuation of the laser beam by neutral filters. With no attenuation the laser power exceeded the threshold value necessary for occurrence of a spark by approximately seven times.

The dependence of the speed of motion of the scattering region, determined from the value of the shift $\Delta\lambda$, on the intensity of the laser radiation is shown in Fig. 5. We can see from the plot that the dependence is very weak and we can state that the velocity increases with increasing J not faster than $J^{1/3}$.

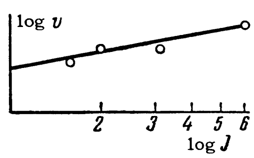

FIG. 5. Dependence of the rate of motion of the scattering region on the intensity of the laser beam.

2. INTERPRETATION OF RESULTS

The motion of the ionized region against the laser beam, which appears in the Doppler shift of the laser line, can in principle be explained by three mechanisms:

A. Hydrodynamic mechanism. The air heated in the region of the initial breakdown, expanding, sends a shock wave along the luminous channel against the beam (and also in other directions). In the shock wave the gas is heated to a high temperature, is ionized, and acquires the ability to strongly absorb light. Thus, the zone of intense absorption of the beam and of energy release moves along the luminous channel against the beam following the shock wave, much as occurs in a detonation wave.[8]

B. Luminous mechanism. The short-wavelength radiation of the air heated in the region of initial breakdown ionizes the adjacent layers so that they acquire the ability to absorb the light. This leads to heating of this region by the energy of the laser beam and as a result the region of high temperature and increased ionization moves along the luminous channel.

C. Progressive breakdown mechanism. Because of the increase in power of the laser beam with time, the region in which the threshold conditions for a discharge are attained is propagated with time to regions further from the center of the caustic curve, i.e., the region of breakdown moves along the luminous channel. A detailed analysis of all three mechanisms has been given in the work of one of the authors.[17]

Under the conditions of the experiment described above, the hydrodynamic mechanism appears more probable to us. We will assume that it applies for our conditions.

The velocity v of a detonation wave in gases is determined by the heating power q (erg/g) according to the formula[18]

$$v = \sqrt{2(\gamma^2 - 1)q}, \tag{4}$$

where γ is the exponent of the adiabatic curve. In our case the energy deposition per gram is

$$q = J / \rho_0 v, \tag{5}$$

where J is the flux of radiant energy of the laser incident per second on 1 cm^2 of the surface of the wavefront, and ρ_0 is the initial density of the air. Combining formulas (4) and (5) we obtain

$$v = [2(\gamma^2 - 1)J / \rho_0]^{1/3}. \tag{6}$$

The specific internal energy which the gas acquires as the result of the energy release (after absorption of the beam) is

$$\varepsilon = \gamma v^2 / (\gamma^2 - 1)(\gamma + 1). \tag{7}$$

If we substitute into formulas (6) and (7) the experimental values $J = 2 \times 10^{18}$ ergs/cm²-sec, $\rho_0 = 1.29 \times 10^{-3}$ g/cm³, and $\gamma = 1.33$—the effective value in the temperature region 500×10^3—1000×10^3 °K,[16] we obtain v = 133 km/sec, $\epsilon = 1.35 \times 10^{14}$ erg/g. This energy corresponds to an equilibrium temperature $T \approx 910 \times 10^3$ °K. As we see, the theoretical values of v and T are correct in order of magnitude but somewhat too large in comparison with the experimental values. This is due first of all to the fact that the "detonation" velocity (6) was computed without considering the energy loss due to lateral expansion of the gas in the zone of energy deposition. In reality the losses are important, since the width of this zone is of the order of the mean free path of the laser light $l_p \sim 10^{-2} - 10^{-3}$ cm (these values correspond to temperatures of hundreds of thousands of degrees) and is comparable with the radius of the luminous channel $\sim 10^{-2}$ cm. The "actual" value of light flux is decreased by about a factor of two because of the loss. If we substitute into formula (6) in place of $J = 2 \times 10^{18}$ a value two times smaller—10^{18} erg/cm²-sec, we obtain v = 105 km/sec, $\epsilon = 8.5 \times 10^{13}$ erg/g, $T = 720 \times 10^3$ °K, which is in very good agreement with the experimental values. The experimental dependence of v(J) also agrees well with the theoretical dependence $v \sim J^{1/3}$.

As has already been noted, Ramsden and Savic[8] utilized the hydrodynamic mechanism in interpretation of their results. To evaluate the speed of motion of the beam absorption zone, Ramsden and Savic use a formula similar to (6) [2)] and obtain a value in agreement with experiment. The plasma temperature was estimated by them from the width of the scattering line, which gives a value $T \approx 40 \times 10^3\,°K$. However, they do not note the fact that the temperature to which the gas is heated is related to the velocity of the wave by formula (7), which is essentially an expression of the law of conservation of energy. According to (7), for their conditions the gas temperature should be more than an order of magnitude higher and should amount to $\approx 700 \times 10^3\,°K$.

The question of why such a small width is obtained experimentally for the shifted line (~ 1 Å in our case and ~ 0.4 Å in the work of Ramsden and Savic[8]) is not yet completely clear.

As has been shown by Salpeter,[19] the scattering of light in a plasma is determined by the parameter $\alpha = \lambda_0/[4\pi \sin(\theta/2)D]$, where $D = (kT_e/4\pi N_e e^2)^{1/2}$ is the Debye radius. The plasma of interest here, for $T_e \approx 700 \times 10^3\,°K$, as follows from the table, has $\bar{Z}_i \approx 5$–6, i.e., $N_e \approx 3 \times 10^{20}\,cm^{-3}$; this gives $\alpha \approx 25$. For $\alpha \gg 1$ the scattering originates in the collective motions of the electrons, and the broadening of the line is determined by the motion of the ions, i.e., by the ion temperature. The time for warmup of the ions to the electron temperature[14] is

$$\tau = \frac{5.87 A_i A_e}{N_e Z_i^2 \ln \Lambda}\left(\frac{T_e}{A_i}\right)^{3/2}$$

and amounts to $\leq 10^{-10}$ sec under these conditions; consequently, in the plasma of interest to us $T_i \approx T_e \approx 60$ eV. For $\alpha \gg 1$, $T_e \approx T_i$, and $\bar{Z}_i \approx 5$–6, the scattered line has a complex shape ($\beta \approx 2$ in Salpeter's formulation[19]), and its width $\Delta\lambda \approx 3\delta\lambda_D$, where $\delta\lambda_D$ is the Doppler width corresponding to the ion temperature. In our case this gives a width $\Delta\lambda \approx 3$–4 Å. Consequently, scattering in a high-temperature plasma would give a width considerably larger than is obtained experimentally.

The small width of the observed shifted line could be explained by assuming that it is due to scattering in a colder region with a temperature $T_e \approx 20 \times 10^3$ by the forward moving front of the shock wave. Another possible explanation is that the main contribution to the intensity of the observed shifted line is given by ordinary reflection of the laser beam at the moving front of the shock wave. Since the intensity distribution over the cross section of the laser beam is very nonuniform, the wavefront can change its direction of motion somewhat in individual regions. In this case the nonuniformity mentioned above in the intensity of the shifted line can be the result of a sharp dependence of the reflection coefficient on the angle of incidence and the angle of observation. A more definite explanation of the width of the shifted line can be obtained only on the basis of further investigations.

The following considerations argue in favor of the hydrodynamic mechanism chosen by us. As has been shown by Raĭzer,[17] for the breakdown mechanism the velocity of advancement of the ionization region should depend strongly on the focal length of the lens L_1 and for $F = 25$ mm should amount to only ~ 40 km/sec. In our experiments the variation of the focus of the lens L_1 from 75 to 25 mm did not give an appreciable change in velocity. With the luminous mechanism it is difficult to explain the discontinuous structure of the scattered line. However it must be emphasized that the realization of a specific mechanism for motion of the ionization region is essentially determined by the experimental conditions and under different conditions different mechanisms can occur.

In conclusion the authors express their gratitude to V. S. Prokudina for calculating the degree of ionization of air and to A. V. Prokhindeev for assistance in the experiment.

[2)]Because of an error made in derivation of the relation between q and J, the formula similar to our Eq. (6) given by Ramsden and Savic [8] contains in place of ρ_0 the density after the wave $\rho = \rho_0(\gamma + 1)/\gamma$.

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