A criterion for the breakdown of dielectrics by a powerful optical radiation pulse

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A phenomenological criterion for the breakdown of dielectrics by a single pulse of light is proposed. An investigation has been made of the distribution of the increase in the internal energy density in the optical radiation absorption region over the electron and phonon subsystems, and of the applicability of the von Hippel criterion to the problem of optical breakdown. The validity of the Seitz criterion has been examined from the point of view of the criterion proposed. A temperature criterion for breakdown is discussed. Expressions for the absorption coefficient of powerful optical radiation in dielectrics have also been derived.

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1. INTRODUCTION

The question of a criterion for breakdown is one of the most basic and least investigated aspects of the problem of the destructive action of a powerful optical radiation pulse on dielectrics. The Seitz "40 generations"[1] criterion is most often used to determine the threshold values of the parameters for optical radiation. This was derived for breakdown in a constant electric field from energy considerations and from data on the mobility of carriers. However, apart from the satisfactory agreement in some cases of the optical breakdown thresholds derived from this criterion with the experimental values, there are no grounds for applying the Seitz criterion in the problem under examination here.

A number of authors[2–4] make use of the criteria due to Fröhlich[5] and to von Hippel[6] which, like the Seitz criterion, have been introduced from the theory of electrical breakdown. Physically these criteria amount to a requirement for an equality of the rate of acquisition of energy by the carriers in an external field $(d\varepsilon/dt)_A$ to the rate of loss of energy by them on interacting with the medium $(d\varepsilon/dt)_L$ at an energy ε equal either to the width of the forbidden band I (Fröhlich) or to the optical-phonon energy (von Hippel), or alternatively for an equality of the average value of these two quantities in terms of energy (von Hippel).[7]

The use of the concept "average" carrier, whereas the interaction of the carrier system with the field and with the phonons has an essentially stochastic nature, has the following consequence. Namely, the criteria mentioned give, on the one hand, highly inflated values of the optical breakdown thresholds (as they do for electrical breakdown thresholds!) and, on the other hand, they lead to independence of the breakdown thresholds of the duration of the radiation pulse, which is patently incorrect both theoretically and from the point of view of the existing experimental data.

Moreover, the relationship $\omega \gg \nu_{\text{eff}}$ is most frequently satisfied for the frequency ω of the optical radiation, where ν_{eff} is the effective frequency of collision of carriers with phonons. One can show that in this case the actual physical premises for the Fröhlich criterion and for the first of von Hippel's criteria lose their validity. However, in the case where $\omega \ll \nu_{\text{eff}}$, as is well known from the theory of electrical breakdown,[7] these criteria lead to the paradoxical conclusion that the breakdown threshold increases with decreasing width of the forbidden band or the optical phonon energy.

For the reasons enumerated the appropriateness of applying the above criteria is doubtful. There are also a number of reasons (which are discussed below) for regarding as erroneous the recently proposed temperature breakdown criterion[8] used subsequently by Epifanov *et al.*[9,10] On the other hand, the authors of all the papers[11–13] in which generation of shock waves in a dielectric under the influence of a single pulse of light was investigated, note that the occurrence of destruction is accompanied by the generation of a region of shock wave absorption in the vicinity. A phenomenological description of this phenomenon in[14] permits the connection of the destruction with the generation of a shock wave under the influence of a single laser pulse to be more fully investigated.

Since breakdown, considered as the appearance of any irreversible change in the dielectric, is a macroscopic phenomenon, it becomes possible on this basis to approach the determination of the optical breakdown threshold more consistently and from a viewpoint other than the microscopic approach. The breakdown criterion which emerges in doing this results directly from the physics of the phenomenon and on its basis it is possible to analyze the applicability of the Seitz and von Hippel criteria to our problem.

The general criterion for the stability of dielectrics to the action of a powerful laser pulse is discussed in Sec. 2 of this paper. In Sec. 3 the distribution of the increase in the density of internal energy $\Delta\varepsilon$ in the region of optical radiation absorption among the electron $(\Delta\varepsilon_e)$ and phonon $(\Delta\varepsilon_{\text{ph}})$ subsystems is investigated, and the limits of the applicability of the second von Hippel criterion to the problem of optical breakdown are examined. In Sec. 4 the consequences of the criterion introduced in Sec. 2 are examined and the Seitz criterion is analyzed on the basis of it. In Sec. 5 the tempera-

ture criterion for breakdown is discussed. In the Appendix expressions are given for $\Delta\varepsilon$, $\Delta\varepsilon_e$ and $\Delta\varepsilon_{ph}$, and for the absorption coefficient of optical radiation intensity near the breakdown threshold. These have been worked out on the basis of the solutions obtained in[15] for the kinetic equation. Except in cases specially noted, the notation used in[15] is employed throughout what follows.

2. THE GENERAL CRITERION FOR THE OPTICAL STABILITY OF DIELECTRICS

As shown in[14], for the laser pulse durations $\tau \lesssim 10^{-8}$ sec examined here the injection of radiation energy into the material takes place more quickly than a macroscopic response can be built up in the material (in particular, one which is manifested in the formation of compression waves and shock waves). At the same time flow of energy out of the absorption region is also effectively absent for times $t \lesssim \tau$ (the time for loss of energy due to heat conduction is $\tau_T \sim 10^{-2}$ sec, and the time for loss of carriers by diffusion $\tau_D \sim 10^{-7}$–10^{-7} sec, with $\tau_T \cdot \tau_D \gg \tau$). For the reasons given, the quantity which fully determines all the macroscopic consequences of the action of a powerful radiation pulse will be the increase in the volume density of internal energy $\Delta\varepsilon(\tau)$ in the absorption region during the time of exposure to radiation. In particular, any criterion for a macroscopic consequence such as the occurrence of irreversible changes in the material must impose limitations on the value of $\Delta\varepsilon(\tau)$.

To be specific we shall start from the strength properties.[1)] At the moment when the laser pulse ends, the material in the absorption region is under an additional pressure $\Delta p_0 = (\tilde{\gamma}-1)\Delta\varepsilon(\tau)$ ($\tilde{\gamma}$ is the effective adiabatic exponent). After the departure of the radiation this excess pressure gives rise to a compression wave which disperses in the surrounding continuum.[14] As follows from an analysis, of the resultant stressed state of the material, carried out in[16] (see Fig. 1[2)]), shear stresses arise behind the compression wave front, accompanied by tensile stresses in a direction normal to the propagation velocity of the waves (which can lead to the formation of cracks spreading out in a fan-like manner from the absorption region).

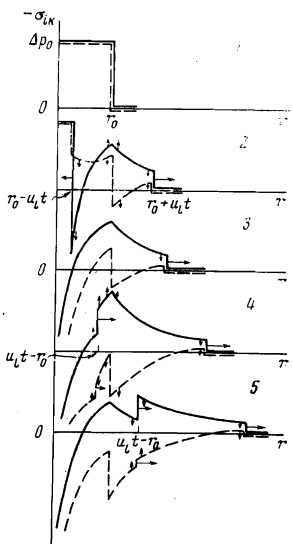

FIG. 1. Evolution of a compression and relaxation wave generated under the influence of a single pulse of optical radiation in the volume of an isotropic dielectric in the case of a spherical absorption region of radius r_0.[16] $-\sigma_{rr}(r,t)$—continuous curve, $-\sigma_{\theta\theta}(r,t)$—dotted curve. Graphs (1)–(5) correspond to the moments in time $t'=0$, $t<r_0/u_l$, $t>r_0/u_l$, $t>2\,r_0/u_l$, referred to the end of the radiation pulse. The horizontal arrows indicate the direction of movement of the waves, the vertical ones—the direction of change in stress as the waves spread out: $\sigma_{\varphi\varphi}(r,t)=\sigma_{\theta\theta}(r,t)$.

In addition, since the geometry of the problem differs from planar (the absorption region is near to spherical if the radiation is focused by a short focal-length lens, and has a cylindrical form when a lens with a long focal length is used or when the phenomenon of "traveling focus" is present), the diverging front of the compression wave drags material behind itself and tensile stresses arise behind it in a direction normal to the front. The maximum value of all the stresses is determined by the amplitude of the compression wave at the moment it is generated:

$$\Delta p_1 = (1+u_0/u_l)^{-1}(\tilde{\gamma}-1)\Delta\varepsilon(\tau),$$

where u_l is the velocity of the longitudinal elastic waves, and u_0 is the speed of sound as determined by the bulk compressibility.[14]

The critical rupture and shear stresses are respectively of the order of $E'/10$ and $\mu/10$ (E' is Young's modulus and μ is the shear modulus), i.e., they are quantities of like order. For this reason one can take, for example, the shear criterion for destruction and write the criterion for optical stability of the dielectric in the form

$$(\tilde{\gamma}-1)\Delta\varepsilon(\tau) < (1+u_0/u_l)\,p_{HEL}, \tag{1}$$

where p_{HEL} is the Hugoniot elastic limit.[3)] From (1) we obtain a breakdown criterion in the form

$$(\tilde{\gamma}-1)\Delta\varepsilon(\tau) \sim 2p_{HEL}. \tag{2}$$

3. DISTRIBUTION OF THE INCREASE IN THE DENSITY OF THE INTERNAL ENERGY IN THE ABSORPTION REGION AMONG THE ELECTRON AND PHONON SUBSYSTEMS. THE VON HIPPEL CRITERION

The quantity $\Delta\varepsilon$ is made up of the increase in the energy density of the electron ($\Delta\varepsilon_e$) and phonon ($\Delta\varepsilon_{ph}$) subsystems. One can correctly derive $\Delta\varepsilon_e(t)$ and $\Delta\varepsilon_{ph}(t)$ by starting from the kinetic equation for the carrier energy distribution function $n(\varepsilon,t)$ in the field of the light wave. Taking account of the expenditure of energy on shock ionization of the valence band, we have for $\Delta\varepsilon_e(\tau)$:

$$\Delta\varepsilon_e(\tau) = \varepsilon_e(\tau)-\varepsilon_e(0) = \int_0^I (\varepsilon+I)\,[n(\varepsilon,t)-n(\varepsilon,0)]\,d\varepsilon \sim IN_e(\tau), \tag{3}$$

$$N_e(t) = N_0 e^{\gamma t} = \int_I^\infty n(\varepsilon,t)\,d\varepsilon,$$

where N_e is the volume density of carriers (γ is the avalanche constant).

Further, from the kinetic equation and the boundary conditions given in[15] it follows that

$$\frac{d\varepsilon_{ph}}{dt}=-\int_0^I \varepsilon\frac{\partial J_2}{\partial\varepsilon}d\varepsilon=\int_0^I J_2(\varepsilon,t)\,d\varepsilon,$$

whence

$$\Delta\varepsilon_{ph}(\tau)=-\int_0^\tau dt\int_0^I d\varepsilon\, J_2(\varepsilon,t)\sim\frac{CI^{1/2}}{\gamma}N_e(\tau). \tag{4}$$

On the basis of estimates from (3) and (4) we obtain

$$\Delta\varepsilon_e(\tau)/\Delta\varepsilon_{ph}(\tau)\sim\gamma/CI^{1/2}.$$

As follows from[15], $\gamma/CI^{1/2}=\varkappa_0/\eta_0 I^{1/2}$ in the low temperature limit and $\gamma/CI^{1/2}=\varkappa_T I^{1/2}/\eta_T I$ in the high temperature limit. Since[15] in the region of intensities near the threshold $\varkappa_0, \varkappa_T I^{1/2}\ll 1$ and $2\eta_0 I^{1/2}$, $\eta_T I\gg 1$, we find that near the breakdown threshold

$$\Delta\varepsilon_e\ll\Delta\varepsilon_{ph} \tag{5}$$

(more exactly, $\Delta\varepsilon_e\sim(10^{-1}-10^{-2})\Delta\varepsilon_{ph}$).

In this way the major part of absorbed optical radiation energy (near the breakdown intensity threshold) goes, owing to the spontaneous emission of phonons by hot carriers, to heat the phonon subsystem. In this process the increasing number of carriers effectively plays the role of transmitting energy from the radiation field to the phonon subsystem. The measurements cited in the first paper by Belikova *et al.*[11] showed that the energy of the light flash accompanying the breakdown is rather small compared with the total energy absorbed. Since the flash is associated chiefly with the recombination of carriers after the departure of the laser pulse (recombination time $\tau_R\gtrsim 10^{-6}$ sec), one can consider the relationship (5) as confirmed.

Further, on the basis of the kinetic equation and the boundary conditions given in[15] we have

$$\Delta\varepsilon_e(\tau)=\int_0^\tau dt\int_0^I d\varepsilon(\varepsilon+I)\frac{\partial n}{\partial t}=\int_0^\tau dt\int_0^I d\varepsilon\left[(\varepsilon+I)\left(\frac{\partial n}{\partial t}\right)_A-\varepsilon\left(\frac{\partial n}{\partial t}\right)_L\right]$$

$$=\int_0^\tau dt\,N_e(t)\left[\left(\overline{\frac{d\varepsilon}{dt}}\right)_A-\left(\overline{\frac{d\varepsilon}{dt}}\right)_L\right],$$

$$\Delta\varepsilon_{ph}(\tau)=-\int_0^\tau dt\int_0^I d\varepsilon(\varepsilon+I)\left(\frac{\partial n}{\partial t}\right)_L=-\int_0^\tau dt\int_0^I d\varepsilon\left(\frac{\partial n}{\partial t}\right)_L\varepsilon=\int_0^\tau dt\left(\overline{\frac{d\varepsilon}{dt}}\right)_L N_e(t),$$

$$\int_0^I d\varepsilon\left(\frac{\partial n}{\partial t}\right)_L=\int_0^I d\varepsilon\frac{\partial J_2}{\partial\varepsilon}=0$$

for losses not leading to a change in the number of carriers participating in the process, the bar denoting averaging with respect to energy. Consequently,

$$\frac{\Delta\varepsilon_e(\tau)}{\Delta\varepsilon_{ph}(\tau)}=\left\{\int_0^\tau dt\left[\left(\overline{\frac{d\varepsilon}{dt}}\right)_A-\left(\overline{\frac{d\varepsilon}{dt}}\right)_L\right]N_e(t)\right\}\Big/\int_0^\tau dt\left(\overline{\frac{d\varepsilon}{dt}}\right)_L N_e(t). \tag{6}$$

From (5) and (6) it follows that the von Hippel criterion

$$\overline{(d\varepsilon/dt)}_A=\overline{(d\varepsilon/dt)}_L$$

while patently leading to independence of the threshold for optical breakdown of τ and to an incorrect dependence on I, may lead to numerical results which are not inconsistent (possibly, on the high side) within the framework of the problem being analyzed here.

4. PHENOMENOLOGICAL BREAKDOWN CRITERION AND THE SEITZ CRITERION

Taking (5) into account, one can replace $\Delta\varepsilon(\tau)$ in (2) by $\Delta\varepsilon_{ph}(\tau)$ and with equal accuracy $\bar\gamma-1$ by the Grüneisen coefficient Γ. Then according to (4) we obtain a criterion for breakdown in the form

$$\gamma\tau\sim\ln\left(\frac{2p_{HEL}}{\Gamma N_0 CI^{1/2}\tau}\gamma\tau\right). \tag{7}$$

Criterion (7) is a transcendental equation with respect to the threshold value $\gamma\tau$. It is easy to see that of the two roots of this equation only the larger has physical significance. After determining $\gamma\tau$ from (7) the determination of the breakdown threshold E_{cr0} (or $\bar W_{cr}=cE_{cr0}^2/8\pi$), as follows from expressions (12) and (14) derived for γ in the Appendix, again reduces to solving a transcendental equation. The procedure for solving this equation and for the choice of the correct root is described in[15]. However, from the general form of (7), (12), and (14) it follows, as in[15], that in the low-temperature region the breakdown threshold does not depend on temperature, and in the high temperature range $E_{cr0}\propto T^{-1/2}$ and $\bar W_{cr}\propto T^{-1}$. Then in both regions $E_{cr0}\sim\omega$ and $\bar W_{cr}\propto\omega^2(\omega^2\gg\nu_{eff}^2!)$. At the same time the breakdown threshold falls markedly with increase in τ and increases with increase in I (approximately as $W_{cr}\propto I^\beta$, where $\beta\sim\frac{1}{3}-1$).

The results of the calculation of W_{cr} for the case of breakdown induced in ruby and leucosapphire by a single noedymium laser pulse ($\tau=3\cdot10^{-8}$ sec, $\hbar\omega=1.17$ eV, $\alpha\sim1$, $m=m_e$, $p_{HEL}=180$ kbar for leucosapphire and 120 kbar for ruby,[17] $\Gamma=1.5$[18]) from the formulae (7), (12), and (14) differ comparatively little from the theoretical $\bar W_{cr}(T)$ curves given together with experimental data in[15]. On this basis one can regard the criterion (7) as reasonable from the point of view of agreement between theory and experiment.

Let us now analyze the Seitz criterion on the basis of (7). The right-hand side of (7) is of considerable magnitude and is logarithmically weakly dependent on the parameters it contains. Evidence of this is presented by the values of $\gamma\tau$ which were obtained from (7) for various values of the parameters on the right-hand side from the range of applicability of (7). From these considerations, the cause of the unsatisfactory agreement of results obtained on the basis of the Seitz criterion with experiment for a number of dielectrics is clear. However, it is also clear that it is more correct to introduce the critical value of $\gamma\tau$ not for all materials at once, as Seitz does, but for groups of materials with relatively similar parameters. This has been done in Table I, on the basis of criterion (7), for materials with $p_{HEL}\sim10^2$ kbar (ruby, leucosapphire, K-8 glass, TF-

TABLE I. Critical values of the multiplication factor for $\gamma\tau$ and the concentration of carriers $N_e(\tau)$.

No. cm^{-3}		1	10^8	10^{12}	Material	$N_e(\tau)$, cm^{-3}
$\gamma\tau$	{	45–46	27–28	18–19	Leucosapphire, ruby, K-8 glass, TF-5, ...	10^{19}–10^{20}
	{	40–42	22–24	13–15	PMMA, NaCl,...	10^{17}–10^{18}

5, ...) and for materials with a low Hugoniot limit $p_{HEL} \sim 10^{-1}$–10^{0} kbar (PMMA, NaCl, ...). After introducing the critical value of $\gamma\tau$ the breakdown threshold can be computed, as was done in[15] using the Seitz criterion (with a solution of the transcendental equation in order not to give too high a value to the threshold!). In addition, on the basis of criterion (7), an absolute value the carrier density in the absorption region $N_e(\tau) = N_0 e^{\gamma\tau}$ can be introduced for each group of materials; when this density is reached, breakdown occurs (see the table).

Let us note that the substitution of the more rigorous expressions (A.1) and (A.3) in (2) for (4) does not markedly alter the results.

5. THE TEMPERATURE CRITERION FOR BREAKDOWN

Let us now examine the temperature breakdown criterion proposed in[8]. As in[8], we shall start from the self-consistent equation for change in temperature in the absorption region obtained in the high-temperature approximation. From (A.3) it follows that

$$\frac{dT}{dt} = \frac{4}{(3\pi)^{3/2}} \frac{(eE_0)^3 \mathscr{E}_1^2 (kT)^{1/2}}{C_V \hbar^4 \rho^2 u_l^3 m^{3/2} \omega^3} N_e(t) \tag{8}$$

(C_V is the heat capacity, $\mathscr{E}_1$ is the deformation potential constant), where by virtue of taking account of the increase in temperature during the time of exposure to radiation we have on the right-hand side $T = T(t)$ and (see[16])

$$N_e(t) = N_0 \exp\left(\int_0^t \gamma_T(T(t_1))\, dt_1 \right). \tag{9}$$

From (8) and (9) it is easy to establish that the Epifanov[8] criterion amounts to an approximation in which the dependence on temperature is only taken into account in the pre-exponential multiplier ($\sim T^{3/2}$) in (8) and no substantial dependence of γ_T on T is considered. This fault must be corrected. Taking account of (14) and the equality

$$N_e(t) = \frac{1}{\gamma_T} \frac{dN_e}{dt}$$

which fixes (9), Eq. (8) can easily be integrated. In the solution obtained, the limiting transition $T \to \infty$ can formally be carried out, and having assumed, as does Epifanov[8] that $T \to \infty$ for $t \to t_0 = \tau$, one obtains the criterion for breakdown. Then, taking account of (9) and $\eta_{T_0} I \gg 1$, we have for the carrier multiplication factor

$$\int_0^\tau \gamma_T(t)\, dt = \ln\left[1 + \frac{3\alpha}{1+\alpha} \frac{u_l^2 \omega^2 m^2 \rho C_V}{(eE_0)^2 k N_0} \right]. \tag{10}$$

For breakdown in materials such as ruby and leucosapphire by a single laser pulse, criterion (10) and expression (9) give inflated values for the breakdown concentration of carriers $N_e(\tau) \approx 10^{22}$ cm^{-3}, which exceeds the value $N_p \approx 10^{21}$ cm^{-3} corresponding to the plasma frequency for neodymium and ruby laser radiation. This suggests that the approximation $t_0 = \tau$ is incorrect. In fact, strictly speaking, the limiting transition $T \to \infty$ is not valid since even at $T_1 \sim 5000$ K the approximations $\eta_T I \gg 1$ and $\varkappa_T T^{1/2} \ll 1$ used in deriving $n(\varepsilon, t)$ and (8) are violated. For the reason given γ_T rises with increase of T from T_0 to T_1, and on exceeding T_1 it begins to decrease (to zero as $T \to \infty$) although physically the decrease of γ_T with increase of T is impossible. The singularity t_0 loses physical meaning in this process and there is no reason for it not to coincide with an infinitely remote point. By virtue of the convergence of the integral in (10) with this value of $\gamma_T(t)$ as $\tau \to \infty$, it is effectively not dependent on the upper limit. The causes indicated led Epifanov *et al.*[8–10] to inflated values of the breakdown thresholds which did not depend (in contrast to experiment) on the value of τ.

At the same time, according to the data,[11] for breakdown of ruby by a single light pulse the increase in temperature in the absorption region is $\Delta T \sim 2 \cdot 10^3$ K. Since the temperature increases significantly only when $t = \tau$ is reached, the influence of this increase in T on the breakdown threshold (as shown in[16]) is effectively small. Therefore, as before, one may with sufficient accuracy not take the change in temperature into account in a self-consistent manner. Then, putting $T(t) = T_0$ on the right-hand side of (8), we obtain according to the data in the table, compared with the data in[11], $\Delta T(\tau) \sim 2 \cdot (10^2$–$10^3)$ K for $T_0 \sim 500$ K.

In the case of low temperatures $n(\varepsilon, t)$ does not depend on T and according to (A.1)

$$T(t) = T_0 + \frac{2^5}{3^2 5^2 \pi} \frac{\mathscr{E}_1^2 (eE_0)^6}{\hbar^4 \rho^2 m^2 \omega^6 u_l^3 \gamma_0 C_V(T_0)} N_e (e^{\gamma_0 t} - 1).$$

This expression does not have singularities at finite values of t, and a temperature criterion of the type examined cannot be set (allowance for the $C_V(T)$ dependence at low T only lowers the value of $\Delta T(t)$).

Thus, the temperature criterion is not general enough. Moreover, in view of the considerations presented above, it must obviously be rejected at the present stage.

6. CONCLUSION

As a result of our analysis it must be considered that criterion (7) is the most justified. At the same time, regarding the data in the table as a modification of the

Seitz criterion, the breakdown threshold should be calculated as in[15].

Let us note that at very much shorter radiation pulses one can also start from a general criterion of the type (2) if the quantity $\Delta\varepsilon(\tau)$ is calculated with allowance for the predominance of multiphoton absorption processes at such values of τ.

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APPENDIX

To calculate $\Delta\varepsilon_e(\tau)$ and $\Delta\varepsilon_{ph}(\tau)$ defined in (3) and (4) we make use of $n(\varepsilon, t)$ from[15]. In the low-temperature approximation ($mu_i^2 \ll kT \ll (mu_i^2 I)^{1/2}$) we obtain, employing a asymptotic expansion in the parameter $2\eta_0 I^{1/2} \gg 1$,

$$\Delta\varepsilon_e^{(0)}(\tau) \approx \left(\frac{12}{(2\eta_0)^2} + I\right) N_e(\tau),$$
$$\Delta\varepsilon_{ph}^{(0)}(\tau) \approx \frac{60CN_e(\tau)}{(2\eta_0)^{7/2}\gamma_0} = \frac{2^5}{3^2 5^2 \pi} \frac{\mathscr{E}_1^2 (eE_0)^6}{\hbar^4 \rho m^2 u_i^3 \omega^6 \gamma_0} N_e(\tau), \quad \text{(A.1)}$$

where[4)]

$$\gamma_0 = \frac{3^3 5^3}{4\pi} \frac{\alpha}{1+\alpha} \frac{\mathscr{E}_1^2 m^7 u_i^3 \omega^6 I^2}{\hbar^4 \rho (eE_0)^6} \exp(-2\eta_0 I^{-1}),$$
$$\eta_0 = \frac{15}{2} \frac{m^2 u_i \omega^2}{(2m)^{1/2}(eE_0)^2}. \quad \text{(A.2)}$$

On the other hand, in the high-temperature approximation ($kT \gtrsim (mu_i^2 I)^{1/2}$) we have at $\eta_T I \gg 1$

$$\Delta\varepsilon_e^{(T)}(\tau) \approx \left(\frac{3}{2\eta_T} + I\right) N_e(\tau), \quad \text{(A.3)}$$
$$\Delta\varepsilon_{ph}^{(T)}(\tau) \approx \frac{4}{3} \frac{CN_e(\tau)}{\eta_T^{5/2}\gamma_T} = \frac{4}{(3\pi)^{1/2}} \frac{(eE_0)^3 \mathscr{E}_1^2 (kT)^{1/2}}{\hbar^4 \rho u_i^3 m^{1/2} \omega^3 \gamma_T} N_e(\tau),$$
$$\gamma_T = \frac{3^{1/2} 2^{7/2}}{\pi^{1/2}} \frac{\alpha}{1+\alpha} \frac{\mathscr{E}_1^2 (2m)^{1/2} u_i^3 \omega^3 m^5 I^2}{\hbar^4 \rho (kT)^{1/2} (eE_0)^3} \exp(-\eta_T I),$$
$$\eta_T = 6 \frac{u_i^2 \omega^2 m^2}{kT (eE_0)^2}. \quad \text{(A.4)}$$

Since in computing (A.1) and (A.3) the approximations of[15] were used, the region of applicability of (A.1) and (A.3) is limited to fields near to threshold (10^5 V/cm $< E_0 < 10^7$ V/cm).

For reference we shall cite expressions for the absorption coefficient for powerful optical radiation, which is defined as

$$\xi(t) = \int_0^I \sigma_\omega(\varepsilon) n(\varepsilon, t)\, d\varepsilon.$$

With the same approximations and in the same range of fields as above we have

$$\xi_0(t) = \frac{2^{11} 7}{3^3 5^4} \frac{e^2 \mathscr{E}_1^2 \hbar^6 (eE_0)^8}{c\rho u_i^5 m^5 (\hbar\omega)^{11}} N_e(t),$$
$$\xi_T(t) = \frac{16}{9(3\pi)^{1/2}} \frac{e^2 \mathscr{E}_1^2 \hbar (kT)^{5/2} (eE_0)^3}{c\rho u_i^5 m^{5/2} (\hbar\omega)^6} N_e(t). \quad \text{(A.5)}$$

The expressions (A.5) replace the relationship

$$\xi(t) = \frac{4\pi e^2 N_e(t)}{cm\omega^2} \nu_{eff},$$

which follows from the Drude–Zener theory and which is a poor approximation for hot carriers.

From (A.5) and the data in the table it follows that for breakdown in dielectrics $\xi(t)$ varies from 10^{-11}–10^{-10} cm^{-1} for $t = 0$ ($N_0 \sim 10^8$ cm^{-3}) to 10^1–10^3 cm^{-1} for $t = \tau$. In reality the value of $\xi(0)$ is, of course, substantially larger ($\sim 10^{-4}$ cm^{-1}) since it is due to other (including linear) mechanisms of absorption which, however, are not definitive for $t \sim \tau$. At the same time the average value of the absorption coefficient (A.5) which is equal to

$$\bar{\xi} = \frac{1}{\tau} \int_0^\tau \xi(t)\, dt = \frac{1}{\gamma\tau} \xi(\tau)$$

($\sim 10^0$–10^2 cm^{-1}), no longer depends effectively on the value of $\xi(0)$.

[1)] One can start from the temperature reaching the melting point in the absorption region. However, this gives results close to those of the present approach.

[2)] Components of the stress tensor differing from zero are given in the figure for the case of a spherical absorption region. In this case the shear stresses are equal to

$$\frac{1}{2}(\sigma_{rr} - \sigma_{\theta\theta}) = \frac{1-2\sigma}{2(1-\sigma)} \sigma_{rr},$$

where σ is Poisson's coefficient.

[3)] In an approach of this kind the criterion (2) which follows from criterion (1) coincides with the criterion for the generation of a two-wave configuration consisting of a plastic shock wave with a precursor under the influence of a single laser pulse.

[4)] Since in formula (2) in[15] the number 1 ought to be absent in the second bracket of the expression under the integral sign, the numerical multipliers in (A.2) and (A.4) are slightly different from those given in[15].

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Luminescence accompanying the deformation and fracture of metals

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Experiments are described in which low intensity radiation in the visible region of the spectrum (3000-8200 Å) produced in rapid fracture and deformation of metals is observed and investigated. Two independent recording techniques are employed, one with an electron-optical image converter and the other with a photoelectric multiplier. The light is found to be emitted from fissures and from the most deformed parts. The emission duration is longer than that of the fracture process by approximately an order of magnitude. The emission has a band spectrum with a maximum of 7200–7300 Å. There are all reasons to believe that the radiation has attributes of luminescence.

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A conviction of long standing was that metals do not luminesce in the condensed state (see, e.g., [1]). We have demonstrated[2] in 1965 that the emission produced when a large-density current passes through copper conductors, under conditions when no discharges of any kind exist, is not of thermal origin. In 1969, Muradyan has excited luminescence of a metal by light to produce photoluminescence of copper and gold, while in 1971 photoluminescence came into use for the investigation of the band structure of metals.[4]

An investigation of the radiation produced upon development of neck-type magnetohydrodynamic instability in a liquid copper conductor[5–7] has shown that this radiation has all the attributes of luminescence.[1)] Two physical circumstances in the experimental conditions described earlier[5] could lead to excitation of luminescence, either an electric field or rapid destruction of the (liquid) metal by the instability.

When neck-type MHD instability arises in a cylindrical conductor, the electric field intensity in the metal and the current density are larger at the instant preceding the breaking of the current than those attained in any other experiment by many orders of magnitude. It is therefore natural to assume that electroluminescence of a metal has been registered for the first time. At the same time, the rupture of an MHD conductor by the instability takes place at rather high velocities ($v \geq 100$ m/sec) and it would not be absurd to suggest that this strongly-acting factor could lead to excitation of luminescence.[2)]

We describe here the performance and results of experiments in which we have established the production of nonthermal radiation when metals are deformed in the absence of electric and magnetic fields, as well as the results of an investigation of the temporal and spectral characteristics of the observed radiation. A preliminary communication on this subject was published earlier.[9]

Two series of experiments were performed. In the first we attempted to observe the radiation produced at a metal strain and fracture rate ~7 m/sec. A vertical ram and a reversing mechanism was used to rupture cylindrical samples of 3–4 mm diameter. The radiation detector was a photomultiplier sensitive to the 3000–8200 Å band. The maximum sensitivity of the recording system for an integral light flux was 10^{-8} lum, or in absolute power units 10^{-2} W. The experiments were performed with copper, aluminum, molybdenum, and bismuth. The radiation was not reliably registered but was fixed approximately in $\frac{1}{3}$ of the experiments, i.e., it could not essentially be reproduced in controlled fashion. One cannot exclude the possibility that this was caused by the imperfection of the detecting system (its geometry).

In the second series of experiments we attempted to