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ESTIMATE OF THE AMPLITUDES OF THE FIELDS CREATED BY AN UNSTEADY GAMMA SOURCE

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It is known [1-4] that an unsteady gamma source gives rise to an electromagnetic field in the surrounding space. Most of the studies of the characteristics of such fields have been performed in the approximation which is linear in the field [1-3]. An exception is [4] in which the slowing down of Compton electrons by the electric field is taken into account. It follows from [1, 2] that the characteristic scale of the fields created close to the source is of the order of $3 \cdot 10^4$ V/m.* Although this value is appreciably lower than the value of breakdown fields in air, electric discharges are observed [5] in the vicinity of a gamma source, indicating the presence of substantially larger fields. One effect not taken into account in the latter approximation which could lead to an increase in the field is the increase in electron termperature due to the electric field [6]. On the one hand, this decreases the electron mobility and consequently also the conductivity of the system. On the other hand, it is known that the electron attachment coefficient γ for electronegative molecules strongly affects the characteristics of electric fields and depends on the electron energy. Therefore, the electron balance equation must take account of the dependence of γ on the electric field through the electron energy, and this leads to a further change in conductivity. We take account of these effects on the shaping of electric fields in air in the vicinity of the source. It is assumed that electron lifetimes are determined solely by their attachment to molecules. This is a good approximation for air pressures near normal [1-3].

Let us consider the dependence of the electron energy and mobility on the intensity of the electric field. It is shown in [4] that if the electron thermalization time $\tau = 1/\nu\delta$, where ν is the frequency of collisions of electrons with gas molecules and δ is the average relative loss of energy of an electron in a collision, is very much shorter than the characteristic times determining the shaping of electric fields, the processes are quasistatic. In this case one can assume that the electron energy ε at a given instant is determined by the electric field E at that same instant. The relation between these quantities for δ = const and $\nu = \nu_0\sqrt{\varepsilon/\varepsilon_0}$ is derived in [6] and has the form

*A similar value is obtained also from the results of [1] for proper values of the physical constants.

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$$\varepsilon(t) = (1/2)\varepsilon_0[1 + \sqrt{1 + (E(t)/E_0)^2}], \qquad (1)$$

where $\varepsilon_0 = (3/2)\varkappa T$; $E_0^2 = 3\varkappa T\delta/8m\mu_0^2$; $\mu_0 = e/m\nu_0$; $\varkappa T$ is the temperature of the gas in energy units; and e and m are, respectively, the charge and mass of an electron.

Using (1) we obtain for the electron mobility

$$\mu = \sqrt{2}\mu_0/[\sqrt{\sqrt{1 + (E/E_0)^2} + 1}], \qquad (2)$$

where μ_0 is the weak field value of the electron mobility.

Using similar assumptions the same dependence of the electron mobility on the electron field is obtained in [7] except that it is proposed to use the experimental value of the electron mobility for μ_0 in $E_0^2 = \delta\varkappa T/\pi m\mu_0^2$. In this case Eq. (2) is in good agreement with experiment [7]. Since $\mu_0 = 10^6$ cgs units in air [8], we find $E_0 = 50 \cdot 10^2$ V/m. Henceforth we shall use this value for estimates.

The average values [8] of the physical constants determining the interaction of electrons with air molecules in the energy range $\varepsilon \leq 1.2$ eV can be written in the form

$$\nu(\varepsilon) = \nu_0(\sqrt{\varepsilon/\varepsilon_0})\xi, \quad \nu_0 = 1.75 \cdot 10^{11}\ \text{sec}^{-1}; \qquad (3)$$

$$\delta(\varepsilon) = \text{const} = 1.7 \cdot 10^{-3}; \qquad (4)$$

$$\gamma(\varepsilon) = 2.1 \cdot 10^{10}\xi^2\varepsilon^{1.66}/(1 + 470\varepsilon^{2.22}),\ \text{sec}^{-1}, \qquad (5)$$

where $\xi = p/p_0$ is the ratio of the air pressure to normal atmospheric pressure.

The value (5) characterizes the electron attachment probability in triple collisions with oxygen molecules. In the indicated energy range, and for air pressures near normal, radiative and dissociative captures of electrons by oxygen molecules can be neglected [8], and these processes are not taken into account in (5). In [8] the values of ν and δ were averaged over data from various experimental papers using total interaction cross sections of electrons with oxygen and nitrogen molecules which take account of the contributions from various elementary processes.

It follows from (3)-(5) that $\gamma/\nu\delta$ reaches a maximum value of 0.4 for $\varepsilon = 0.065$ eV and then decreases for further increases in energy. Since the electron lifetime $1/\gamma$ is the shortest of the characteristic times, the processes in the system will be quasistatic.

Thus it is clear that if the electric field heats up electrons to energies no higher than 1.2 eV, i.e., if its value is less than $4.8 \cdot 10^5$ V/m [cf (1)], then Eqs. (3) and (4) and the simple expressions for the energy (1) and mobility (2) based on them are valid. If the time dependence of the gamma source is described by a de' a function, the results obtained will be valid at those points of space where the field $E \leq 4.8 \cdot 10^5$ V/m. The results presented below show that this restriction is unneccessary if the source decays exponentially with time.

On the basis of the above considerations and the results of [1-4], the time rates of change of the field E and the secondary electron density n, taking account of the dependence on the field of the mobility, the secondary electron attachment coefficient, and the range of Compton electrons (the factor $[1 + E(t)/E_T]^{-1}$ on the right-hand sides of Eqs. (6) and (7) [4]), we write the system of equations

$$\frac{dn}{dt} + \gamma[\varepsilon(t)]n = \frac{\eta N}{\lambda^3}\frac{e^{-x}}{x^2}\frac{f(t-x)}{1 + \frac{E(t)}{E_T}}; \qquad (6)$$

$$\frac{dE}{dt} + 4\pi e\mu_0 n(t)\frac{\sqrt{2}\,E}{\sqrt{1 + \sqrt{1 + (E/E_0)^2}}} = \frac{Nel}{\lambda^3}\frac{e^{-x}}{x^2}\frac{f(t-x)}{1 + \frac{E(t)}{E_T}}, \qquad (7)$$

which we supplement by Eq. (1) and appropriate initial conditions. In Eqs. (6) and (7), η is the number of secondary electrons per MeV of absorbed energy, N is the total gamma yield, l is the range of a Compton electron, λ is the range of a gamma proton, $x = r/\lambda$, f(t) is a function characterizing the time dependence of the gamma source, and $E_T = \varepsilon_1/el \simeq 3 \cdot 10^5$ V/m, where ε_1 is the energy of a Compton electron.

The solution of system (6) and (7) is given below for a pulsed gamma source and one which decays exponentially with time.

For a pulsed gamma source

$$f(t) = \delta(t).$$

In this case the solution of system (6) and (7) can be obtained by solving the corresponding homogeneous equations with the initial conditions:

as $t-x \to 0$

$$n \to n^*_{in} = \frac{n_{in}}{1+\frac{E_{in}}{E_T}}, \qquad E \to E^*_{in} = \sqrt{\frac{1}{4}E_T^2 + E_{in}E_T} - \frac{E_T}{2}, \tag{8}$$

where $n_{in} = (\eta N/\lambda^3)e^{-x}/4\pi x^2$; $E_{in} = (eNl/\lambda^3)e^{-x}/x^2$. Using Eq. (8) we obtain from (7) as $t-x \to 0$

$$\frac{dE}{dt} \to -\frac{\sqrt{2}\,\gamma_0 \frac{E_{in}}{E_a} E^*_{in}}{\sqrt{1+\sqrt{1+\left(\frac{E^*_{in}}{E_0}\right)^2}}}, \tag{9}$$

where $\gamma_0 = 10^8\ \mathrm{sec}^{-1}$; $E_a = \gamma_0 l/\mu_0\eta \simeq 3\cdot 10^4$ V/m.

It follows from Eqs. (6), (7), and (9) that

$$\frac{dE}{dt}\frac{\sqrt{1+\sqrt{1+\left(\frac{E}{E_0}\right)^2}}}{E} + \sqrt{2}\,\gamma_0\frac{E_{in}}{E_a} + \int\limits_{t-x}^{t} \gamma(\varepsilon(t))\frac{\sqrt{1+\sqrt{1+\left(\frac{E}{E_0}\right)^2}}}{E}\frac{dE}{dt}\,dt = 0. \tag{10}$$

The integral in Eq. (10) is transformed as follows:

$$\int\limits_{t-x}^{t} \gamma(\varepsilon(t))\frac{\sqrt{1+\sqrt{1+\left(\frac{E}{E_0}\right)^2}}}{E}\frac{dE}{dt}\,dt = \int\limits_{E^*_{in}}^{E} \gamma(\varepsilon(E))\sqrt{1+\sqrt{1+\left(\frac{E}{E_0}\right)^2}}\,\frac{dE}{E}. \tag{11}$$

Using (11) the solution of Eq. (10), and, consequently, the final result, is

$$(t-x)\gamma_0 = \int\limits_{E}^{E^*_{in}} \frac{dE'\sqrt{1+\sqrt{1+\left(\frac{E'}{E_0}\right)^2}}}{E'\left[\sqrt{2}\frac{E_{in}}{E_a} + \Phi(E')\right]}, \tag{12}$$

where

$$\Phi(E) = -\int\limits_{E}^{E^*_{in}} \frac{\gamma(\varepsilon(E'))}{\gamma_0}\sqrt{1+\sqrt{1+\left(\frac{E'}{E_0}\right)^2}}\,\frac{dE'}{E'}. \tag{13}$$

In a number of cases Eqs. (12) and (13) are considerably simplified.

For example, in the approximation used in [2], where the dependence of the mobility on the electric field is neglected, slowing down of Compton electrons by the field is not taken into account, and $\gamma(\varepsilon) = \gamma_0$ = const, it follows from (8) and (13) that

$$E^*_{in} = E_{in}; \quad \Phi(E) = \sqrt{2}\ln\left(\frac{E}{E_{in}}\right). \tag{14}$$

Further we find from (12) and (14)

$$E = E_{in}\exp\left\{-\frac{E_{in}}{E_a}\left(1 - e^{-\gamma_0(t-x)}\right)\right\}, \tag{15}$$

which is identical with the corresponding expression in [2].

In a more complex case we take account of the effect for the field on the mobility and on the slowing down of Compton electrons. For $\gamma(\varepsilon) = \gamma_0$ = const the system (6), (7) can be solved in terms of elementary functions,

$$(t-x)\gamma_0 = -\ln\left\{1+\frac{E_a}{\sqrt{2}E_{in}}\left(1+\frac{E^*_{in}}{E_T}\right)\left[2\left(\sqrt{1+\sqrt{1+\left(\frac{E}{E_0}\right)^2}}-\sqrt{1+\sqrt{1+\left(\frac{E^*_{in}}{E_0}\right)^2}}\right)-\frac{1}{\sqrt{2}}\ln\left(\frac{\sqrt{1+\sqrt{1+\left(\frac{E}{E_0}\right)^2}}+\sqrt{2}}{\sqrt{1+\sqrt{1+\left(\frac{E}{E_0}\right)^2}}-\sqrt{2}}\cdot\frac{\sqrt{1+\sqrt{1+\left(\frac{E^*_{in}}{E_0}\right)^2}}-\sqrt{2}}{\sqrt{1+\sqrt{1+\left(\frac{E^*_{in}}{E_0}\right)^2}}+\sqrt{2}}\right)\right]\right\}. \quad (16)$$

When $E \gg E_0$ Eq. (16) simplifies to

$$E = E^*_{in}\left[1-\frac{E_{in}}{E_a}\sqrt{\frac{E_0}{2E_{in}}}\left(1-e^{-\gamma_0(t-x)}\right)\right]^2. \quad (17)$$

Finally, we consider the case of strong fields $E \gg E_0$. It follows from (5) that for $\varepsilon > 0.063$ eV

$$\gamma = \gamma_0(\varepsilon_2/\varepsilon)^{\omega}; \quad \omega = 0.56; \quad \varepsilon_2^{0.56} = 0.45. \quad (18)$$

Using (1) we find that (18) is valid for $E > 2.5\cdot10^4$ V/m. Since Eq. (5) is approximate, let us set $\omega = 1/2$. In this case the final result is expressed in terms of tabulated functions. Substituting (18) into (13) and the result into (12) we have

$$(t-x)\gamma_0 = \frac{1}{\sqrt{2}}\int\limits_E^{E^*_{in}}\frac{dE'}{\sqrt{E'E_0}\left(\frac{E_{in}}{E_a}+\sqrt{\frac{\varepsilon_2}{\varepsilon_0}}\ln\frac{E'}{E^*_{in}}\right)}. \quad (19)$$

Equation (19) can be reduced to the form

$$(t-x)\gamma_0 = \sqrt{\frac{E^*_{in}}{2E_0}}\sqrt{\frac{\varepsilon_0}{\varepsilon_2}}\exp\left(-\frac{E_{in}}{2E_a}\sqrt{\frac{\varepsilon_0}{\varepsilon_2}}\right)\left[E_i\left(\frac{E_{in}}{2E_a}\sqrt{\frac{\varepsilon_0}{\varepsilon_2}}\right)-E_i\left(\frac{E_{in}}{2E_a}\sqrt{\frac{\varepsilon_0}{\varepsilon_2}}+\frac{1}{2}\ln\frac{E}{E^*_{in}}\right)\right], \quad (20)$$

where $E_i(z)$ is the exponential integral [9].

From an analysis of Eq. (20) or (19) it can be shown that as $t\to\infty$

$$E \to E^*_{in}\exp\left\{-\frac{E_{in}}{E_a}\sqrt{\frac{2E_0}{E^*_{in}}\frac{\gamma_0}{\gamma(E^*_{in})}}\right\} = E^*_{in}\exp\left\{-\frac{E_{in}}{E_a}\sqrt{\frac{\varepsilon_0}{\varepsilon_2}}\right\}, \quad (21)$$

where

$$\gamma(E^*_{in}) = \gamma_0\sqrt{\frac{\varepsilon_2}{\varepsilon(E)}} = \gamma_0\sqrt{\frac{2\varepsilon_2}{\varepsilon_0}\frac{E_0}{E^*_{in}}}.$$

As $t\to\infty$ we find from Eq. (15) that $E\to E_{in}\exp\{-E_{in}/E_a\}$. Comparing this result with (21) and taking account of (17) we see that the presence of the factor $\sqrt{2E_0/E^*_{in}} \ll 1$ in the exponent of the exponential in (21) leads to an increase in the electric field as a result of the $1/\sqrt{E}$ decrease of electron mobility in strong fields. The presence in the exponent of the quantity $\gamma_0/\gamma(E^*_{in}) \gg 1$ leads to a decrease in the electric field. This is due to the fact that electrons are captured less strongly by molecules at early times, and, consequently, the conductivity in the system is larger than when $\gamma(\varepsilon)$ = const. When these factors are superimposed they largely compensate one another, since $\sqrt{\varepsilon_0/\varepsilon_2} \simeq 0.32$, and the final result does not depend on E_0.

We now consider a gamma source which decays exponentially with time

$$f(t) = be^{-bt}.$$

We assume that $b \ll \gamma$. Then the solution of Eq. (6) can be written in the form

$$n = n_{in}(b/\gamma(\varepsilon(t)))e^{-b(t-x)}[1 - E(t)/E_T].$$

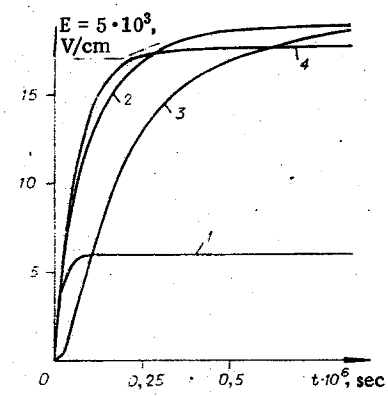

Fig. 1

In this case the time dependence of the electric field is described by the equations

$$\frac{dy}{dt} + \frac{E_{in}}{E_a}\frac{\gamma_0}{\gamma(\varepsilon(t))}\frac{\sqrt{2}\,ybe^{-b(t-x)}}{(1+\beta y)\sqrt{1+\sqrt{1+y^2}}} = \frac{E_{in}}{E_0}\frac{be^{-b(t-x)}}{(1+\beta y)}, \tag{22}$$

where $y = E/E_0$ and $\beta = E_0/E_T$. Integrating Eq. (22) with a zero initial condition, we obtain

$$(t-x)\,b = -\ln\left\{1 - \frac{E_0}{E_{in}}\int_0^{E/E_0}\frac{dy\,(1+\beta y)\sqrt{1+\sqrt{1+y^2}}}{\sqrt{1+\sqrt{1+y^2}} - \sqrt{2}\,\frac{E_0}{E_a}\frac{\gamma_0}{\gamma(\varepsilon)}\,y}\right\}. \tag{23}$$

In a number of cases Eq. (23) is considerably simplified.

We first assume that $\gamma(\varepsilon) = \gamma_0 =$ const and that there is no slowing down of Compton electrons by the electric field ($\beta = 0$). In this case by setting $dE/dt = 0$ in Eq. (22) we can determine the maximum value of the electric field

$$E_{max} = E_a E_a/2E_0. \tag{24}$$

Using the values of the physical constants given in [2, 8] we find that $E_{max} = 9 \cdot 10^4$ V/m. An electric field of this magnitude actually has little effect on the motion of Compton electrons. Setting $\beta = 0$, Eq. (23) can be expressed in terms of elementary functions

$$\begin{aligned}(t-x)\,b = -\ln\Bigg\{1 + \frac{E_a}{\sqrt{2}E_{in}}\Bigg[&2\sqrt{1+\sqrt{1+y^2}} - 2\sqrt{2} + \\ &+ 2d\ln\left(\frac{\sqrt{1+\sqrt{1+y^2}}}{\sqrt{2}} - \frac{\sqrt{\sqrt{1+y^2}+1}}{\sqrt{2}}\right) - \frac{2(1+d^2)}{\sqrt{2+d^2}} \times \\ \times \ln\left(\frac{2\sqrt{2+d^2}\sqrt{1+\sqrt{1+y^2}}}{d-\sqrt{\sqrt{1+y^2}-1}} + \frac{2(2+d^2)}{d-\sqrt{\sqrt{1+y^2}-1}} - 2d\right) &+ \frac{2(1+d^2)}{\sqrt{2+d^2}}\ln\left(\frac{2}{d}\left(\sqrt{2(2+d^2)}+2\right)\right)\Bigg]\Bigg\},\end{aligned} \tag{25}$$

where $d = E_a/\sqrt{2}E_0$.

For $y \gg 1$ (strong fields) Eq. (25) reduces to the form

$$E = E_a\frac{E_a}{2E_0}\left\{1 - \exp\left[-\frac{E_{in}E_0}{E_a^2}\left(1 - e^{-b(t-x)}\right)\right]\right\}^2, \tag{26}$$

which also follows directly from Eq. (22) for $\gamma(\varepsilon) = \gamma_0$ and $E \gg E_0$.

Neglecting the dependence of the electron mobility on the field, we obtain from (22) as $E_0 \to \infty$

$$E = E_a\{1 - \exp\left[-(E_{in}/E_a)(1 - e^{-b(t-x)})\right]\}, \tag{27}$$

which agrees exactly with the corresponding expression in [2].

Let us consider the general case. Suppose $\gamma(\varepsilon)$ is given by Eq. (5). Using (5) and (1) we obtain

$$\gamma(E) = \gamma_1 \frac{\left[1+\sqrt{1+\left(\frac{E}{E_0}\right)^2}\right]^{1.66}}{1+\alpha\left[1+\sqrt{1+\left(\frac{E}{E_0}\right)^2}\right]^{2.22}}, \tag{28}$$

where $\gamma_1 = 1.5 \cdot 10^7 \sec^{-1}$; $\alpha = 2.84 \cdot 10^{-2}$.

Substituting (28) into (23) we obtain the final result in the form

$$(t-x)b = -\ln\left\{1 - \frac{E_0}{E_{\text{in}}} \int\limits_0^{E/E_0} \frac{dy\,(1+\beta y)\,[1+\sqrt{1+y^2}]^{2.16}}{(1+\sqrt{1+y^2})^{2.16} - \frac{\sqrt{2}E_0}{E_a}\frac{\gamma_0}{\gamma_1}\,y\,[1+\alpha\,(1+\sqrt{1+y^2})^{2.22}]}\right\}. \tag{29}$$

From Eq. (22) we obtain the maximum value of the electric field determined by (29)

$$E_{\max} = E_0[(E_a/E_0)\gamma_1/\sqrt{2}\quad\alpha\gamma_0]^{1/1.06} \simeq 9.35\cdot 10^4\ \text{V/m}. \tag{30}$$

The result (30) is insensitive to the accuracy of the determination of E_0. It is obvious that for $\omega = 1/2$ [Eq. (17)] the limiting value of the electric field would not depend on E_0 at all. This is related to the fact that as the electric field increases, the decrease in electron mobility would be exactly compensated by the increase in conductivity due to the increase in the attachment coefficient. For $\gamma(\varepsilon) = \gamma_0 =$ const the limiting value of the electric field is strongly dependent on the accuracy of determining E_0 (24). There is a fundamental difference in principle in the results obtained by taking account of the dependence of the attachment coefficient on the electron energy and not taking it into account, whereas Eqs. (24) and (30) for $E_0 = 50$ V/cm lead to nearly the same result.

We note that as the height of the gamma source is increased, effects related to the influences of the electric field on the electron mobility will decrease since

$$E_a/E_0 \sim \rho/\rho_0 \sim e^{-z/h},$$

where ρ is the density of the air at the height of the source, and ρ_0 is the density of air under normal conditions.

The same conclusion can also be drawn as to the effects related to the slowing down of Compton electrons since

$$E_{\max}/E_T \sim (\rho/\rho_0)^2 \sim e^{-2z/h}.$$

At sufficiently large heights the picture is changed still further since the electron lifetime at large heights will be determined by electron–ion recombination processes [4].

Figure 1 shows the time dependence of the field for a number of cases at the point x = 1 for a gamma source which varies exponentially with time. Curve 1 is plotted from Eq. (27), curve 2 from (25), 3 from (26), and 4 from (23). An analysis of these curves shows that the effect of the field on the electron mobility leads to an appreciable increase in the maximum value of the field and its relaxation time.

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ANOMALOUS BEHAVIOR OF AN INTENSE LIGHT FLUX IN A MEDIUM WITH TWO TYPES OF ABSORPTION

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UDC 535.343

A study was made in [1] of the behavior of the spatial distribution of an intense light flux in an amorphous medium with two types of absorption – absorption by impurities with a subsequent rapid transfer of energy to the medium through radiationless processes and absorption by the medium itself. The latter type does not occur initially but as a result of heating of the zone around the impurity centers there is a temperature shift in the absorption edge and the corresponding parts of the medium start to absorb. This form of absorption eventually predominates for sufficiently large intensities. The details of the radiationless processes and of the temperature distribution around the impurities have been considered in [2, 3]. The situation studied in [1] corresponds to times by which all the populations have reached a stationary distribution and spatial zones far from the front of the light flux.

It is impossible, without making simplifications, to get an analytical solution to this problem for the initial moments of time when the nonstationary nature of the intensity distribution and of the populations are extremely important. We have therefore derived numerical solutions. The calculations showed that the behavior of the light intensity in the transient region is very peculiar. As in [1], we consider the propagation of a plane parallel monochromatic light flux which at the instant t = 0 is incident from the left on the surface of a medium which occupies the first half-space. The equations which describe the process are

$$\begin{gathered}\partial U/\partial t + c\partial U/\partial x = -N_1 c\sigma U - N_3 c\sigma_0 U + \alpha N_4 c\sigma_0 U;\\ dN_4/\partial t = c\sigma_0 U(N_3 - N_4);\ N_3 + N_4 = N_0;\\ \partial N_1/\partial t = -N_1 c\sigma U + N_2/\tau;\ N_1 + N_2 = N;\\ U(x,\ 0) = 0;\ U(0,t) = U_0;\ N_1(x,\ 0) = N;\ N_3(x,\ 0) = N_0(x,\ 0) = 0,\end{gathered} \tag{1}$$

where U is the density of quanta in the light; N is the concentration of impurities with photoabsorption cross section σ; N_1 and N_2 are the concentrations of impurities in the ground and excited states, respectively; N_0 is the concentration of absorbing molecules in the medium with photoabsorption cross section σ_0; N_3 and N_4 are the concentrations of these molecules in the ground and excited states; τ is the radiationless relaxation time of an impurity; and c is the velocity of light in the medium. The quantum density U is related to the light intensity I by $I = c\varepsilon U$, where $\varepsilon = \hbar\omega$ is the energy of one quantum. The factor α in (1) is introduced to allow for the fraction of quanta reradiated by the medium in the direction of the original flux. The diffusively scattered quanta are assumed to pass outside the zone under consideration and to play no part in (1). The equation for N_4 takes into account the absorption processes and the rapid radiationless deactivation. In order to determine N_0 we utilize the exact equation (2.3) of [1]. We get

$$N_0 = (4\pi/3)(3\varepsilon c\sigma U/8\pi v\rho_0)^{3/2}(1 + 3\varepsilon c\sigma U/8\pi v^3\rho_0(t - x/c)^2)^{-3/2}NN_c,$$

where N_c is the density of molecules in the medium, ρ_0 is the threshold oscillation-energy density at which

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