CREATION OF NONEQUILIBRIUM STATES AND CHANGE OF PLASMA ABSORPTIVITY UNDER THE ACTION OF POWERFUL LIGHT PULSES

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We examine the kinetic processes which take place in a plasma under the influence of laser radiation and lead to the creation of nonequilibrium states and variations of the plasma absorptivity. By solving the Fokker-Planck equations for the bound states we find the excited atom level population distribution under the action of the radiation. We calculate the stepwise ionization rate and then the kinetics of the variations of the electron and excited atom densities and the absorption coefficient. The results of the light absorption calculation are compared with the authors' experimental data on laser pulse propagation through a plasma. Satisfactory agreement between theory and experiment is obtained.

Experiments [1-3] have shown that plasma absorptivity depends on the radiation intensity and varies nonmonotonically; both reduction and increase of the absorptivity have been observed. A qualitative explanation of the effects was proposed in [1-3], which related them with the creation of nonequilibrium states in the plasma subjected to intense radiation. The need to understand and describe quantitatively the action of processes of laser radiation absorption by a low-temperature dense plasma arises in examining quite different phenomena, for example in estimating screening by ionized vapors of solid surfaces on which the laser ray acts, in establishing the limits of applicability of laser plasma diagnostics when the action of light on the test object cannot be tolerated, and so on.

All this makes advisable a specialized analysis of the kinetic processes which lead to the creation of nonequilibrium states and the change of plasma absorptivity under the influence of powerful light pulses. This is the objective of the present paper. For definiteness all the calculations were made in application to the conditions of [1-3], with the results of which they are compared.

1. Light Absorption in Equilibrium Plasma and Comparison of Theory with Direct Measurement Data. We shall present some known information on light absorption in gases in the first ionization region (see, for example, [4]), which is necessary for the later analysis. The absorption coefficient $\varkappa$ is made up of the coefficients corresponding to bound-free transitions and free-bound transitions of the electrons in the ion field $\varkappa = \varkappa_1 + \varkappa_2$. The coefficient of true photoionization absorption equals

$$\varkappa_1 = \sum_k N_k \sigma_{1k} \tag{1.1}$$

where N_k is the number of atoms per cm^3 in the k-th quantum state, σ_{1k} is the photoeffect cross section. The summation extends to all those atom states (realized in the plasma) from which removal of electrons by the given quanta $h\nu$ is energetically possible. The coefficient of true stopping absorption equals

$$\varkappa_2 = \frac{4}{3}\left(\frac{2\pi}{3kT}\right)^{1/2}\frac{e^6}{hcm^{3/2}\nu^3} N_e^2 g^\circ = \frac{3.7\cdot 10^8 N_e^2 g^\circ}{T^{1/2}\nu^3}\ cm^{-1} \tag{1.2}$$

where $N_e = N_+$ is the number of electrons per cm^3, T is the electron temperature, g° is a correction factor of order unity.

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For a thermodynamically equilibrium gas of hydrogen atoms with $h\nu$ and kT which are small in comparison with the ionization potential I, the summation in (1.1) can be replaced approximately by integration. If we also set $g^\circ = 1$, we obtain the well-known Unsold-Kramers equation

$$\varkappa = \varkappa^\circ = \frac{16\pi^2 ke^6}{3\sqrt{3}\,eh^4}\,\frac{TN_a}{\nu^3}\exp\frac{-I+h\nu}{kT} = 0.89\cdot 10^{24}\,\frac{TN_a}{\nu^3}\exp\frac{-I+h\nu}{kT} \tag{1.3}$$

where N_a is the number of all neutral atoms per cm^3. Then

$$\varkappa_1 / \varkappa_2 = \exp(h\nu / kT) - 1; \tag{1.4}$$

the actual attenuation of the light ray is characterized by an effective coefficient which is defined by the difference between true absorption and stimulated emission

$$\varkappa' = \varkappa\,[1 - \exp(-h\nu / kT)] \tag{1.5}$$

According to [5, 6] the absorption coefficient of a gas of complex atoms can be represented approximately in the form

$$\varkappa = \varkappa^\circ\,\frac{2g_+}{g_0}\,\xi(\nu)\exp\frac{'h\Delta\nu}{kT} \tag{1.6}$$

Here g_+ and g_0 are the statistical weights of the ion and atom; $\xi(\nu)$ are functions characteristic of each element; the last factor accounts for the lowering $h\Delta\nu$ of the continuous spectrum boundary in the plasma.

In the case of inert gases the primary correction to the Unsold-Kramers equation is that of the weighting factor $2g_\perp/g_0 = 12$;* at the ruby laser frequency $\nu = 4.3\cdot 10^{14}$ sec^{-1}; then $\xi = 1.2$-1.4. The exponential factor is usually close to one. In more exact calculations the principal terms in the sum (1.1) should be isolated, integrating only over the remaining levels [6]; for xenon, studied in the experiments of [1-3], the most complete calculations of this type were made in [7].

For the conditions of the [1-3] experiments with T = 10,000° K = 0.86 eV, $N_a = 6.1\cdot 10^{18}$ cm^{-3}, $N_e = 0.38\cdot 10^{18}$ cm^{-3}, ionization ratio 0.06, $\nu = 4.3\cdot 10^{14}$ sec^{-1}, using the Biberman-Norman equation without account for reduction ($\Delta\nu = 0$) we obtain $\varkappa' = 0.061$ cm^{-1} ($\xi = 1.35$); $\varkappa_1'/\varkappa_2' = 7$. According to Yankov's calculation $\varkappa_1' = 0.08$ cm^{-1} and the 5d level makes the dominant contribution to the sum (1.1). Calculating $\varkappa_1'$ using (1.2) and (1.5) with $\varkappa_2' = 0.0059$ cm^{-1} and using the Yankov values, we obtain $\varkappa_1'/\varkappa_2' = 13.5$; $\varkappa' = \varkappa_1' + \varkappa_2' = 0.086$ cm^{-2}.†

In the experiments of [1-3] with low light intensities, when the light does not excite the plasma and absorption is linear, the value $\varkappa' = 0.105$ cm^{-1} was obtained. However, it is difficult to be confident about the good agreement between the calculations and the direct experimental data, since the plasma parameters themselves are not known with sufficient accuracy.

2. Absorption in Nonequilibrium Plasma. In speaking of plasma nonequilibrium we will assume that the electron density N_e and the excited atom state population N_k are nonequilibrium, and we shall also have in mind the difference between the electron T and ion (atom) T_i temperatures, while retaining the assumption of Maxwellian distribution in electron (and, naturally, in ion–atom) gases. Estimates show that the Maxwellian distribution in an electron gas is usually established rapidly, for example, in the [1-3] experimental conditions after a time of $\sim 10^{-12}$-10^{-13} sec; even high intensity light $\sim 10^4$-10^5 MW/cm^2 violates this condition very little. Under these assumptions the stopping absorption coefficients $\varkappa_2$, $\varkappa_2'$ are described by the previous formulas (1.2), (1.5) and the original Eq. (1.1) for $\varkappa_1$ naturally remains valid, but the Eq. (1.5) for the effective coefficient $\varkappa_1'$ now becomes invalid.

Let us find $\varkappa_1'$. On the basis of general relations (see, for example, [4]) the partial coefficient $\varkappa_{1k}'$, associated with photoionization of the k-th atom state and induced photorecombination at this level, equals

$$\varkappa_{1k}' = \varkappa_{1k} - j_{1k}c^2 / 2h\nu^3, \qquad \varkappa_{1k} = N_k\sigma_{1k} \tag{2.1}$$

*In $\varkappa_1$ this factor appears because of the multiplicity of the terms, in $\varkappa_2$ it appears when replacing N_e^2 using the Saha equation.

† Stopping absorption during collisions of electrons with neutral atoms yields a contribution of less than 3% to the absorption.

Here j_{1k} erg · cm^{-3} (sec · ster · Hz)$^{-1}$ is the corresponding emittance. It is proportional to the rate of photoabsorption of electrons with energies $\varepsilon = h\nu - E_k$ at the k-th level (E_k is the binding energy of the k-th level). With the aid of the detailed balancing principle [4] we can express the photoabsorption cross section in terms of σ_{1k} and $N_e N_+ = N_e^2$ – in terms of the population N_k of the k-th state, equilibrium with respect to electron density and temperature

$$N_k^e = N_e^2 \frac{g_k}{2g_+}\left(\frac{h^2}{2\pi m k T}\right)^{3/2} \exp\frac{E_k}{kT} \tag{2.2}$$

where g_k is the statistical weight of the k-th level, and we obtain

$$\varkappa_1' = \sum \varkappa_{1k}', \qquad \varkappa_{1k}' = \sigma_{1k}\left[N_k - N_k^e \exp\frac{-h\nu}{kT}\right] \tag{2.3}*$$

In many real cases and specifically in the [1-3] experimental conditions, as will be shown later, the populations of the excited atom states do not differ markedly from the equilibrium states with respect to an electron gas, although the relationship between the electron N_e and all neutral atom N_a numbers may be far from equilibrium (in accordance with the terminology of [8] there exists a "block" of excited and ionized atom states). In these cases, to calculate $\varkappa'$ we can use the known data on the absorption coefficient of a thermodynamically equilibrium gas if in the corresponding expressions we convert from N_a to N_e using the Saha equation.

In the approximation in which summation over the levels is replaced by integration, on the basis of (1.6), (1.3), (1.5) we obtain for $\varkappa^{e'} = \varkappa'$ the block

$$\varkappa = \frac{2.42\cdot 10^{-37}\xi(\nu) N_e^2 [\exp(h\nu/kT) - 1]}{(h\nu)^3 T^{1/2}} \text{ cm}^{-1} \tag{2.4}$$

where $h\nu$ and T are expressed in electron volts.

In deriving this equation, corresponding to (1.6), we assumed that the reduction $h\Delta\nu$ of the continuous spectrum boundary does not differ from the decrease ΔI of the ionization potential in the plasma.

3. Excited State Populations; Rate of Stepwise Ionization of Atoms. The absorptivity is determined primarily by the numbers of atoms in the different excited states (population numbers N_k); therefore we shall examine how the atoms are distributed by levels and the degree to which the distribution differs from equilibrium with respect to the electrons. We shall do this on the basis of the Fokker-Planck equation.

In the considered case of sufficiently dense and sufficiently ionized plasma the roles of atom–atom collisions and radiative transitions are small in comparison with the role of electron collisions. We can further assume that only the very highest atom levels participate in ionization by electron collision and recombination in triple collisions. We introduce into the usual Fokker-Planck equation for these conditions [9, 10] a term describing photoionization and induced photorecombination under the action of the laser light in accordance with the equation

$$(\partial N_k/\partial t)_\nu = -S\varkappa_{1k}'$$

where S is the quantum-flux density. We have the following equation for the continuous distribution function f (E) of the atoms with respect to the binding energies E, which take the discrete numbers E_k, and $f_k(E_k) = N_k/g_k$

$$\frac{\partial f}{\partial t} = -\frac{1}{\partial(E)}\frac{\partial j}{gE} - S\sigma_1(E)\left[f - f^e \exp\frac{-h\nu}{kT}\right], \qquad f^e = \frac{N_k^e}{g_k}$$
$$j = D(E)\, g(E)\left(\frac{\partial f}{\partial E} - \frac{f}{kT}\right), \qquad g(E) = \frac{g_k}{\Delta E_k} \tag{3.1}$$

where j is the flux along the energy axis, g (E) is the state density, ΔE_k is the distance between levels, D is the diffusion coefficient, according to [10] equal to

* It is curious that if the populations of any of the "strong" levels are less than $N_e \exp(-h\nu/kT)$ and the corresponding $\varkappa_{1k}' < 0$ the situation is possible in which the overall effective absorption coefficient $\varkappa'$ is also negative, i.e., amplification takes place. We emphasize, however, that this situation cannot be accounted for per se.

$$D = D^\circ x, \qquad x = \frac{E}{kT}, \qquad D^\circ = \frac{2}{3}\sqrt{2}\,\pi e^4 \left(\frac{kT}{m}\right)^{1/2} N_e \ln\lambda$$
$$D^\circ = 1.45\cdot 10^{-6}\, T_{ev}^{1/2} N_e \ln\lambda \;\; \partial e^2/ce\kappa \tag{3.2}$$

where $T_{e\nu}$ is the temperature expressed in electron volts, and $\ln\lambda$ is the Coulomb logarithm. The expression $D = D^\circ x$ is used for convenience in converting to the dimensionless energy variable $x = E/kT$.

In the approximation of hydrogen-like excited levels $E_n = I/n^2$ (n is the principal quantum number), $g_n = 2n^2$, $\Delta E = 2I/n^3$, so that $g(E) = n^5/I$ or

$$g(E) = \langle g\rangle x^{-5/2}, \qquad \langle g\rangle = I^{3/2}(kT)^{-5/2} \tag{3.3}$$

The photoeffect cross-section [4]

$$\sigma_{1n} = 0.79\cdot 10^{-17}\,(I/h\nu)^3\, n^{-5}\ \mathrm{cm}^2$$

Hence for $E \le h\nu$ or $x \le x_\nu = h\nu/kT$

$$\sigma_1(E) = \sigma^\circ x^{5/2}, \qquad \sigma^\circ = 0.79\cdot 10^{-17}\,(I/kT)^{1/2} x_\nu^{-3}$$

For $E > h\nu$ or $x > x_\nu$

$$\sigma_1(E) = 0$$

If, as is usually the case, the number of excited atoms is far less than the number of unexcited atoms and electrons, there is rapidly established among the excited atoms a quasistationary distribution which follows the comparatively slowly varying electron density and temperature. Therefore in (3.1) we can set approximately $\partial f/\partial t = 0$ and assume that

$$f(E, t) = f[E, N_e(t), T(t)]$$

We convert, as in [8], to relative populations

$$y_k = N_k / N_k^\circ, \qquad y(E) = f(E)/f^\circ(E)$$

where the index $^\circ$ denotes the equilibrium Boltzmann quantities

$$N_k^\circ = N_a g_k \Sigma_a^{-1} \exp\frac{E_k}{kT}, \qquad f^\circ = N_a \Sigma_a^{-1} \exp\frac{E}{kT}$$
$$\Sigma_a = \sum_{k=1}^{I} g_k \exp\frac{E_k}{kT}, \qquad E \equiv I$$

Here Σ_a is the atom electron statistical sum.

We have the relation

$$N_k^e / N_k^\circ = f^e / f^\circ = (N_e/N_e^\circ)^2 = y_e^2$$

where N_e° is the equilibrium electron density, calculated using the Saha equation.

In the relative variables y,x the equation for the flux and Eq. (3.1) with expressions (3.2)-(3.4) substituted therein take the form

$$j = -\langle j\rangle x^{-3/2} e^x \frac{dy}{dx}, \qquad \langle j\rangle = D^\circ \langle g\rangle \frac{N_a}{kT}\Sigma_a \tag{3.5}$$

$$\frac{d}{dx}\left(x^{-3/2}e^x\frac{dy}{dx}\right) = \begin{cases} \gamma_0 e^x [y - y_e^2 e^{-x_\nu}] & (0 < x < x_\nu) \\ 0 & (x > x_\nu) \end{cases} \tag{3.6}$$

The parameter γ in (3.6) characterizes the direct influence of external radiation on the level populations; the other indirect influence is associated with heating of the electron gas, as a result of which the electron density becomes lower than the equilibrium value.

Let us formulate the boundary conditions. The uppermost levels are very strongly coupled with the electrons through the rapidly proceeding ionization by electron impact and absorption in triple collisions, and are therefore in equilibrium with respect to the electrons. Consequently

$$y(0) \equiv (y)_{x=0} = y_e^2 \tag{3.7}$$

and external radiation does not disturb this equilibrium, since

$$\sigma_1 \sim x^{5/2} \to 0 \quad \text{as} \quad x \to 0$$

We pose the second condition at the point $x = x_2 = E_2/kT$, isolating from the continuous energy scale the most "discrete" transition between the unexcited (k = 1) and lower excited (k = 2) states. The flux $j^2 \equiv j(x_2)$ at this point is

$$j_2 = -K_{12}N_1N_e + K_{21}N_2N_e = -K_{12}N_1^{\circ}N_e\,(y_1 - y_2) \approx -K_{12}N_aN_e\,(1 - y_2) \tag{3.8}$$

where K_{12} is the rate constant for atom excitation by electron impact; here we have used the detailed balancing principle and the fact that there are far more unexcited than excited atoms, so that $N_1 \approx N_1^{\circ} \approx N_a$, $y_1 \approx 1$. The rate constant K_{12} equals [4]

$$K_{12} = (8kT/\pi m)^{1/2}C^*\,(E^* + 2kT)\exp\frac{-E^*}{kT} \tag{3.9}$$

where $E^* = I - E_2$ is the excitation energy, C^* is the coefficient in the threshold dependence of the excitation cross section on the electron energy ε

$$\sigma_{12}(\varepsilon) = C^*(\varepsilon - E^*)$$

We substitute j from (3.5) into (3.8) and find the boundary condition

$$\left(\frac{x^{-3/2}e^x\,dy/dx}{1-y}\right)_{x=x_2} = \frac{K_{12}N_aN_e}{\langle j\rangle} \equiv \beta \tag{3.10}$$

The solution of (3.6) with the boundary conditions (3.7), (3.10) for given N_e, T, N_a, S determines not only the atom level distribution y(x) but also the rate of change of the number of electrons per cm^3, since by virtue of quasistationarity of the distribution

$$\frac{dN_e}{dt} = -\frac{dN_1}{dt} = -j_2 \approx K_{12}N_aN_e(1 - y_2) \tag{3.11}$$

Let us examine the important practical case in which external radiation is not very stong or is entirely absent: $\gamma = 0$, but the electron density is nonequilibrium for some reason or other. For $\gamma = 0$ (3.6) is easily solved analytically, since j = const [11]. Substituting the integral of (3.6) into (3.11), we find

$$\frac{dN_e}{dt} = \frac{K_{12}N_aN_e(1 - y_e^2)}{1 + \beta J(x_2)}, \quad J(x) = \int_0^x x^{3/2}e^{-x}\,dx$$
$$(J(1) = 0.21,\ J(2) = 0.61,\ J(3) = 0.93,\ J(\infty) = 1.33) \tag{3.12}$$

Equation (3.12) is similar to the known solutions of [10, 11]. In the limiting case $y_e \gg 1$, $K_{12} \to \infty$, $x_2 \to \infty$, which corresponds to the problem formulation in [10], (3.12) yields the expression obtained in [10] for the recombination rate.

For $y_e < 1$, when the kinetic process proceeds in the ionization direction, (3.12) yields the resultant rate of stepwise ionization.

If $\beta \ll 1$, we see from comparison of (3.12) and (3.11) that $y_2 \approx y_e^2$ (and in general $y(x) \approx y_e^2$), i.e., the excited atoms are in equilibrium with the electrons. In this case

$$\frac{dN_e}{dt} = K_{12}N_aN_e(1 - y_e^2) = K_{12}N_aN_e\left(1 - \frac{N_e^2}{N_e^{02}}\right) \tag{3.13}$$

Equation (3.13) automatically ensures approach of dN_e/dt to zero as the electron density approaches the equilibrium value.

The condition of $\langle j\rangle \ll 1$ for the existence of the block $\beta = K_{12}N_aN_e$ is physically understandably clear since the parameter β characterizes the binding ratio of the excited atoms with the unexcited (through the excitation rate) and with the electrons (through the diffusion coefficient D, which is proportional to the quantity $\langle j\rangle$). The quantity β equals approximately (if we set $\Sigma_a \approx g_1 \exp(I/kT)$)

$$\beta \approx 4.6\cdot 10^{13}C^*\frac{E^* + 2kT}{\ln\lambda}\frac{kT}{I}T_e^2g_1\exp\frac{E_2}{kT} \tag{3.14}$$

It is independent of the electron and atom densities and is minimal for $kT \approx E_2/3.5$, essentially for $T \approx 1$ eV, and $\beta_{min} \sim 10^{-2}$.

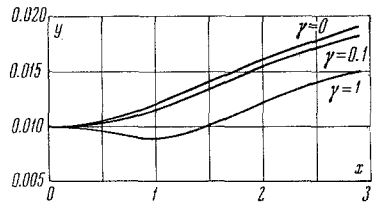

Fig. 1

Figure 1 shows examples of the numerical solution of (3.6) for $y_0 = 0.1$ and $\beta = 0.01$ (actually the numerical calculation was made only in the region $0 < x < x_\nu$, since for $x_\nu < x < x_2$ the equation can be solved analytically, and since y and dy/dx are continuous in x_ν). We shall present an example of the numerical values of the parameters. Let $N_e = 0.38 \cdot 10^{18}$ cm^{-3}, as in the initial plasma of the [1-3] experiments $N_a = 6.1 \cdot 10^{18}$ cm^{-3}; in xenon I = 12.1 eV, $E^* = 8.4$ eV, $E_2 = 3.7$ eV, we set $C^* = 3 \cdot 10^{18}$ cm^2/eV (the choice of this quantity will be discussed later), $h\nu = 1.78$ eV (ruby). Let $S = 3.3 \cdot 10^{27}$ phot/cm$^2 \cdot$ sec, which corresponds to a light intensity of 1000 MW/cm^2 and T = 1.2 eV (initial plasma temperature $T \approx 10{,}000°$ K = 0.86 eV). Setting $\ln \lambda = 0.22$ in accordance with [11], we find $\beta = 0.007$, $\gamma = 0.28$.

We see from Fig. 1 that in the high excitation region, where photoionization takes place, the populations do not differ markedly from the equilibrium populations with respect to electrons. We note that in the case of very strong radiation in a large part of the photoionization region $0 < x < x_\nu$ there is established the distribution $y = y_e^2 \exp(-h\nu/kT)$ corresponding to approximate compensation of photoionization and induced photorecombination.

4. Kinetics of Heating, Ionization, and Light Absorption. We shall formulate the system of equations describing the changes of state and absorptivity of the plasma under the action of intense radiation. We shall neglect hydrodynamic effects and also diffusion and thermal conduction, since they cannot show up during the short time of the giant laser pulse. We assume that the number of excited atoms is small in comparison with the numbers of unexcited atoms and electrons, so that the approximation of quasistationarity of the excited states, examined in the previous section, is applicable; in this case we can also neglect the excitation energy in the gas in comparison with the ionization energy. The energy of the absorbed radiation transitions into energy of the electron gas, where each free electron has the thermal energy $^3/_2$ kT and the potential energy I. Therefore

$$\frac{d}{dt} N_e (^3/_2 kT + I) = h\nu\varkappa' S - {}^3/_2 k N_e (T - T_i)/\tau_{ei}$$

where $\tau_{ei} = 3.15 \cdot 10^8 A T_{e\nu}^{3/2}/N_e \ln\lambda$ is the characteristic time for energy exchange between electrons and ions (A is atomic weight, $\ln\lambda$ is the corresponding Coulomb logarithm); the comparatively small losses to thermal radiation are neglected. Thus the equations for the temperatures of the electrons and heavy particles have the form

$$\frac{3}{2} k N_e \frac{dT}{dt} = h\nu\varkappa' S - \left(I + \frac{3}{2} kT\right)\frac{dN_e}{dt} - \frac{3}{2} k N_e \frac{T - T_i}{\tau_{ei}} \tag{4.1}$$

$$\frac{3}{2} k N_0 \frac{dT}{dt} = \frac{3}{2} k N_e \frac{T - T_i}{\tau_{ei}}, \qquad N_0 = N_a + N_e = \text{const} \tag{4.2}$$

It is worthy of note that even for very high light intensities, when ionization by the external radiation may be stronger than the ionization by electron impacts, the creation of one free electron still reduces the energy content of the electron gas by $I + {}^3/_2 h\nu$, although it would appear that the energy is derived at the expense of the external source. This is explained by the fact that splitting-out of a single excited atom during photoionization is immediately (by virtue of quasistationarity) compensated by excitation of an unexcited atom, which takes place at the expense of the electron gas energy.

The electron number rate of change dN_e/dt is expressed by the general equation (3.11), which calls for solution of the Fokker-Planck equation. We shall formulate the expression for the absorptivity in nonequilibrium conditions. On the basis of (2.3), replacing summation over the levels by integration and, in addition, using the analogous definition of the effective photoelectric absorptivity $\varkappa^{e'}$, corresponding to equilibrium of excited atoms with electrons, we obtain

$$\varkappa_1 = \varkappa_1^{e'} \frac{e^{x_\nu}}{(e^{x_\nu} - 1)^2} \int_0^{x_\nu} e^x \left[\frac{y}{y_e^2} - e^{-x}\right] dx, \qquad x_\nu = \frac{h\nu}{kT} \tag{4.3}$$

When more exact calculation of the absorption is necessary it may be advisable to substitute into this equation the value of $\varkappa_1^{e'}$ calculated with separation of the strong levels from the integral. The overall coefficient $\varkappa' = \varkappa_1 + \varkappa_2$ with integration over the levels $\varkappa_1^{e'}/\varkappa_2 = e^{x_\nu} - 1$, $\varkappa_2 = \varkappa^{e'} e^{-x_\nu}$, so that

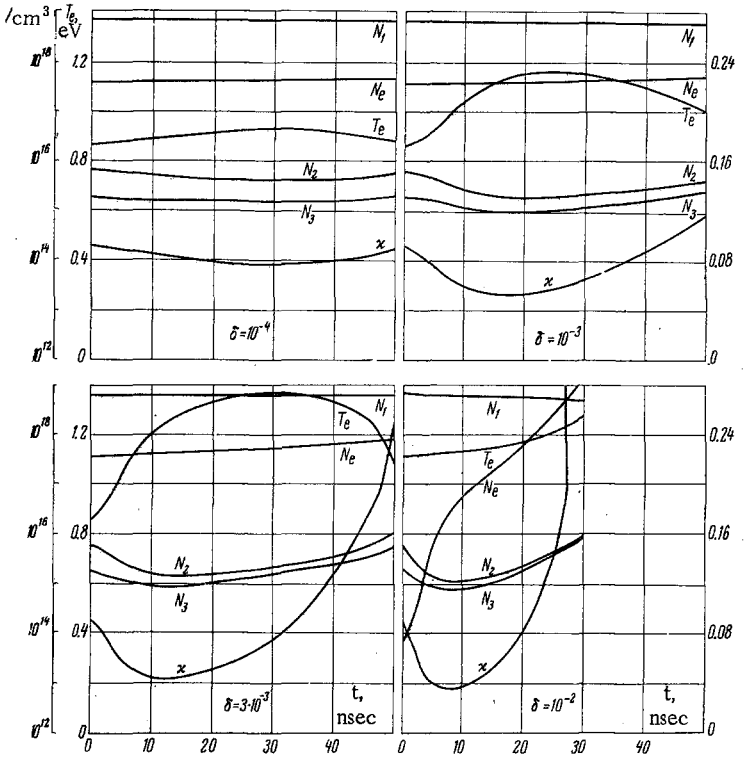

Fig. 2

$$\varkappa' = \frac{\varkappa^{e'}}{e^{x_\nu}-1}\int_0^{x_\nu} e^x\left(\frac{y}{y_e^2}-e^{-x_\nu}\right)dx + \varkappa^{e'}e^{-x_\nu} \qquad (4.4)$$

where $\varkappa^{e'}$ is given by (2.4). The use of this equation involves quite tedious calculations and it is presented only for information.

For the calculations made with the specific objective of explaining the experimental results of [1-3], it was sufficient to use the block-of-states approximation and set $\varkappa' = \varkappa^{e'}$ and also replace "exact" Eq. (3.11) for the ionization rate by the approximate Eq. (3.13).

We thus obtain a simplified system of three ordinary first-order differential equations for the time functions T, T_i, N_e – (4.1), (4.2), (3.13), (2.4) plus the Saha equation, whose solution permits finding $\varkappa'$ (t), i.e., the plasma absorptivity.

The calculation results depend markedly on the coefficient C* in (3.9) for the excitation rate constant, which defines the ionization rate (3.13). Unfortunately, there are no experimental data in the literature on C* for xenon, with which the experiments of [1-3] were conducted (the data are very incomplete for the other gases as well, since in certain cases the data relate to the overall excitation cross section at many levels, and in other cases to the excitation of definite spectral lines). It is known that for argon C* = $7 \cdot 10^{-18}$ cm²/eV, while apparently for neon C* = $1.5 \cdot 10^{-18}$ cm²/eV, and for helium C* = $4.6 \cdot 10^{-18}$ cm²/eV [4].

The calculation results presented below were obtained with the value C* = $3 \cdot 10^{-18}$ cm²/eV. We note that the larger C*, the lower the light intensities at which "clarification" and subsequent increase of the absorptivity take place.

It is interesting to see to what degree the solution of the nonstationary problem in which explicit account is made for the kinetics of the excitation of the different atom levels differs from the solution of the problem in the approximation of quasistationarity of the populations or the existence of a block. (We note that in certain cases the conditions of quasistationarity and smallness of the numbers and energy of the excited atoms may be completely invalid or questionable.) The complete system of kinetic equations for T, T_i, N_e and the ensemble of population numbers N_k is obviously quite cumbersome. Therefore it is ad-

visable to simplify this system if our objective is to clarify the question of the relationship between the nonstationary and quasistationary cases. This can be done by combining approximately all the excited states into two groups with binding energies E_2 and E_3 and population numbers N_2 and N_3. The index 1, as before, is assigned to the ground state. Assuming that only the atoms of the more highly excited third "state" are ionized by quanta and by electron impact, we write the kinetic equations

$$\begin{gathered}\frac{dN_1}{dt} = -K_{12}N_eN_1 + K_{21}N_eN_2\\ \frac{dN_2}{dt} = K_{12}N_eN_1 - K_{21}N_eN_2 - K_{23}N_eN_2 + K_{32}N_eN_3\\ \frac{dN_3}{dt} = K_{23}N_eN_2 - K_{32}N_eN_3 - K_{3e}N_eN_3 + K_{e3}N_e^3 - S\langle\sigma_1\rangle(N_3 - N_3^e\exp(-h\nu/(kT))\\ N_e + N_1 + N_2 + N_3 + N_0 = \text{const}\end{gathered} \tag{4.5}$$

The discrete transition rate constants K_{12}, K_{21}, K_{23}, K_{32} and also the ionization and recombination rate constants K_{3e}, K_{e3} are related as follows by the detailed balancing principle:

$$\begin{gathered}\frac{K_{12}}{K_{21}} = \frac{N_2^\circ}{N_1^\circ} = \frac{g_2}{g_1}\exp\frac{-E_{12}}{kT} \qquad (E_{12} = E_1 - E_2)\\ \frac{K_{23}}{K_{32}} = \frac{N_3^\circ}{N_2^\circ} = \frac{g_3}{g_2}\exp\frac{-E_{23}}{kT} \qquad (E_{23} = E_2 - E_3)\\ \frac{K_{3e}}{K_{e3}} = \frac{N_e^{\circ i}}{N_3^\circ} = \frac{2g_+}{g_3}\left(\frac{2\pi mkT}{h^2}\right)^{3/2}\exp\frac{-E_3}{kT}\end{gathered} \tag{4.6}$$

It is natural to select the average photoionization cross section $\langle\sigma_1\rangle$ so that in the equilibrium case we obtain the known true photoionization absorptivity $\langle\sigma_1\rangle N_3^\circ = \varkappa_1^\circ$. The energy of the absorbed radiation transitions to energy of the electrons and excited atoms. In place of (4.1) we obtain the new equation

$$\frac{3}{2}kN_e\frac{dT}{dt} = h\nu\varkappa' S - \left(I + \frac{3}{2}kT\right)\frac{dN_e}{dt} - E_{12}\frac{dN_2}{dt} - E_{13}\frac{dN_3}{dt} - \frac{3}{2}kN_e\frac{T - T_i}{\tau_{ei}} \tag{4.7}$$

Equation (4.2) remains in force. The effective coefficient $\varkappa'$ equals

$$\varkappa' = \langle\sigma_1\rangle\left(N_3 - N_3^e\exp\frac{-h\nu}{kT}\right) + \varkappa_2' \tag{4.8}$$

The rate constant K_{12} is found from (3.9). The rate constants K_{23} and K_{32} can be selected approximately on the basis of the known values for hydrogen-like atoms [4]; the statistical weights g_2 and g_3 of the grouped levels can be chosen by examining the atom level diagram and taking into account cutoff of the upper states in the plasma. Calculations of the nonstationary system (4.7), (4.2), (4.5), (4.8) with different values of the constants showed that the block-of-states approximation (and, consequently, the quasistationary approximation) is satisfied with adequate precision.

Figure 2 shows the results of calculations of the kinetics and absorption of the laser pulse. The pulse was approximated for ease of calculation by the sinusoidal function $S(t) = S_0 \sin(\pi t/\tau)$ with $\tau = 50$ nsec. In the figures $S_0 = S_t\delta$, where $S_t = 4.9 \cdot 10^{29}$ phot/cm^2 · sec in the peak flux corresponding to breakdown of cold xenon of the same density as the plasma.

Figure 2 shows the time variation of the electron density N_e and the excited atom densities N_2 and N_3 in the second and third groups, the electron temperature, and the light absorptivity. Also shown are the densities of all the atoms. The ion temperature does not differ markedly from the electron temperature (somewhat lower than the latter). Curves are shown for radiation intensities amounting to $\delta = 10^{-4}$, 10^{-3}, $3 \cdot 10^{-3}$, and 10^{-2} of the breakdown value S_t.

Figure 3 shows the form of the transmitted pulse for different incident radiation intensities; this value is compared with the sinusoidal form of the incident light.

Figure 4 shows the calculated ratio of the transmitted and incident light powers for different values of the incident power (dashed curve). Also shown are the experimental data (solid curve, drawn through the experimental points). The scale of the absolute values of the intensity J in MW/cm^2 is plotted with account for the slight conicity of the light beam in the plasma [1] and corresponds to the average values of the light intensity in the plasma (average along the conical channel). We see that the calculated values agree satisfactorily with the experimental data.

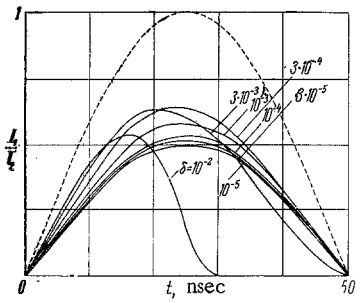

Fig. 3

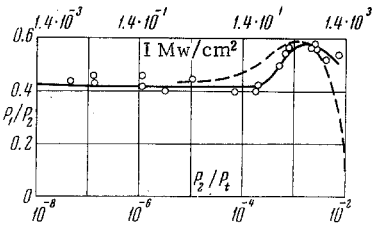

Fig. 4

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