Contribution to the theory of the absorption of a strong electromagnetic field in semiconductors

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The absorption of a strong electromagnetic field by a semiconductor is investigated in the stationary case, and also the nonstationary absorption by a field which is switched on adiabatically is examined. The investigation is carried out with the aid of the system of equations derived earlier. Approximate solutions of the equations are found, enabling one to establish that in a strong field the absorption may decrease substantially without occupation of the bottom of the conduction band by the electrons. The decrease in the absorption becomes especially strong if the field reaches a certain threshold value at which the magnitude of the gap exceeds the maximum phonon frequency. The relaxation processes turn out to be suppressed, and the rate of the removal of the electrons from the production site decreases. This leads to an additional decrease of the absorption and to the creation of an inverted population within the band and between the bands at frequencies exceeding the frequency of the external field.

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In recent years the resonant interaction of a strong electromagnetic field with semiconductors has been investigated in a number of articles[1-5] (see, for example,[6] for a more extensive review of the literature). In comparison with two-level systems, the most distinctive features of semiconductors are the continuous energy spectrum of the electrons in a band and the possibility of intraband relaxation of the electrons. The influence of a strong electromagnetic field leads to more radical changes in the electronic system of the semiconductor. Further, a strong field is to be understood as a field satisfying the condition

$$\lambda\tau_{ph}\gg 1, \tag{1}$$

where $\lambda = \mathbf{d}\cdot\mathbf{E}/\hbar$ is the frequency of electron transitions between the bands under the influence of the field $\mathbf{E}$, $\mathbf{d}$ is the dipole moment of the transition, and τ_{ph} is the time between collisions of electrons with phonons, electrons, impurities, etc. (only the interaction with phonons is investigated below).

The appearance of a gap in the energy spectrum of the electrons and holes[3] turned out to be the most important new property. The renormalization of the spectrum together with the change in the form of the electron and hole wave functions in the presence of a resonant electromagnetic field substantially alters the nature of the interaction of the electrons with phonons. The influence of the field on a collision event becomes especially strong when the magnitude 2λ of the gap exceeds the maximum phonon frequency ω_{ph}:

$$2\lambda>\omega_{ph}. \tag{2}$$

In this case the processes of energy relaxation in the conduction (valence) band turn out to be suppressed since the phonon energy is insufficient for overcoming the gap 2λ.[6,7]1)

The present article is devoted to an investigation of the absorption of a strong electromagnetic field, satisfying condition (2) in the stationary case, and in the nonstationary case—for field pulses having a duration longer than τ_{ph} (short pulses). Our investigation does not include ultrashort pulses $\tau\sim\lambda^{-1}\ll\tau_{ph}$ and the narrow region of the leading edge of the pulse. For short pulses, however, the contribution to the absorption from the region of the leading edge is negligible, and neglect of this time interval makes it possible to utilize the simpler equations derived earlier.[4]

Investigation of the absorption of light by semiconductors in the presence of large flux intensities was begun in the work by Krokhin,[1] utilizing the rate equations and assuming that $\lambda\tau_{ph}\ll 1$. It was shown[1] that the absorption coefficient in direct transitions decreases with increasing intensity, owing to occupation of the bottom of the conduction band (top of the valence band) by the electrons (holes). This mechanism is analogous to the mechanism for the Burstein-Moss effect, but for nonequilibrium electrons.

In strong fields, when condition (1) is satisfied, the absorption coefficient decreases even without occupation of the band. The point is that the direct transitions of electrons between the bands are saturated, and the processes of absorption (emission) of photons proceed via indirect transitions involving the emission (absorption) of phonons. But the probability for indirect transitions is smaller than the probability for direct transitions by $\lambda\tau_{ph}$ times. However, if condition (2) is satisfied then, as is shown below, the absorption undergoes an additional, abrupt decrease. It is due to the decrease in the rate of removal of the electrons from the point where they are produced because of the suppression of the relaxation caused by the gap. Under these conditions there is an accumulation of electrons (holes) having energies above a certain value (that is singled out by the laws of energy and momentum conservation in direct transitions) and a destruction below this value, that is, an inverted population exists inside the band. Furthermore, it is interesting to note that an inverted population also arises between the bands at frequencies greater than the frequency of the external perturbing field. We recall that usually, in stationary case, an inverted population is created at frequencies which are smaller than the frequency of the perturbing field.

In our opinion the effect of a decrease in the absorption for $2\lambda > \omega_{ph}$ is of interest, especially in connection with the interpretation of the experimental results of[8,9] in which the transmission of pulses of an electromagnetic field of large amplitude (corresponding to an intensity $P\approx 10^9$ W/cm^2) was investigated, completely satisfying condition (2), and a significant reduction of the absorption was observed.

1. SUPPRESSION BY THE FIELD OF INTRABAND RELAXATION IN SEMICONDUCTORS

In the strong field regime (1), electron transitions between the bands occur during a time interval smaller than the time between collisions with phonons and, as was shown in[3], the use of noninteracting (with the electromagnetic field) electrons as the zero order approximation loses meaning. The interband, resonant transitions is taken into account by the introduction of quasiparticles, which are superpositions of electrons a_p and holes b_p:[3]

$$\alpha_p=u_p a_p-v_p b_p^+, \quad \beta_p=u_p b_p+v_p a_p^+, \tag{3}$$

$$u_p^2,\ v_p^2=\tfrac{1}{2}(1\pm\xi_p/\varepsilon_p), \quad \xi_p=p^2/2m-p_0^2/2m,$$
$$\mu_0=p_0^2/2m=(\omega-E_g)/2>0, \quad \varepsilon_p=(\xi_p^2+\lambda^2)^{1/2}, \tag{4}$$

$$\lambda=ev_{cv}E/2\omega, \quad \hbar=c=1. \tag{5}$$

Here ϵ_p denotes the energy of the quasiparticles α_p and β_p; ω denotes the frequency of the external electromagnetic field; E_g denotes the width of the forbidden band for a semiconductor with symmetric bands and identical effective masses m of the electrons and holes; v_{cv} is the transition matrix element between the valence band and the conduction band.

Here the kinetic equations for the electrons should be replaced by the kinetic equations for quasiparticles:[4]

$$\frac{\partial}{\partial t}n_p=(1-n_p)S_p^+-n_pS_p^--n_pS_p^A+v_p^4(1-n_p)^2/\tau_R-u_p^4n_p^2/\tau_R, \tag{6}$$

$$\begin{pmatrix}S_p^+\\S_p^-\\S_p^A\end{pmatrix}=2\pi\sum_{qp'}g^2(q)\,\delta(p'+q-p)\begin{pmatrix}n'\delta(\varepsilon-\varepsilon'+\omega_q)(uu'+vv')^2\\(1-n')\,\delta(\varepsilon'-\varepsilon+\omega_q)(uu'+vv')^2\\n'\delta(\varepsilon+\varepsilon'-\omega_q)(uv'-u'v)^2\end{pmatrix},$$
$$n_p{}^\alpha=\langle\alpha_p^+\alpha_p\rangle=n_p{}^\beta=\langle\beta_p^+\beta_p\rangle\equiv n_p; \quad n_{p'},\ u_{p'},\ \varepsilon_{p'}\equiv n',\ u',\ \varepsilon'; \tag{7}$$
$$1/\tau_R=\tfrac{1}{2}e^2\omega\,|v_{cv}|^2,$$

where $g(q)$ is the matrix element of the electron-phonon interaction; $\omega_q = sq$ is the phonon frequency, s is the velocity of sound, τ_R is the radiative recombination time; the equation is written down for the case when the temperature is equal to zero. (Similar results are obtained for optical phonons.)

The first two terms in Eq. (6) describe the scattering of the quasiparticles by phonons, where S^+ describes the arrival and S^- the departure, and the third term containing S^A describes the annihilation of quasiparticles with the emission of a phonon. The fourth and fifth terms reflect the creation and destruction of quasiparticles due to recombination. (As is easy to see from Eq. (3), the destruction of an electron with $\xi < 0$ is equivalent to the creation of a quasiparticle, and the destruction of an electron with $\xi > 0$ corresponds to the destruction of a quasiparticle.)

The distribution functions of the quasiparticles allow us to determine the number of photons absorbed per unit time (due to interband transitions):

$$Q=\sum_p\frac{\xi_p}{\varepsilon_p}\frac{\partial n_p}{\partial t}+Q_{st}, \tag{8}$$

$$Q_{st}=\frac{1}{\tau_R}\sum_p[v_p^4(1-n_p)^2+u_p^4n_p^2]. \tag{9}$$

Expression (8) was previously obtained[4] in a somewhat different form and was reduced to the form (8) by Galitskiĭ and the author[10] (the expression for the steady-state absorption was found by the author[11]).

Let us analyze the solutions of Eqs. (6) and (7) for $2\lambda < \omega_{ph}$ and $2\lambda > \omega_{ph}$. As a preliminary, let us discuss the role of recombination. We shall consider the ratio τ_{ph}/τ_R to be a small parameter, since usually the relaxation time τ_R greatly exceeds the relaxation time due to collisions with phonons. However, recombination determines the total number of quasiparticles (recombination is the source of the quasiparticles) and, if $2\lambda > \omega_{ph}$, it turns out to have a singificant influence on the shape of the quasiparticle distribution function. As we shall see below, the latter is valid upon fulfillment of the condition

$$\eta=\frac{\tau_{ph}}{\tau_R}\left(\frac{\mu_0}{\lambda}\right)^2\gg 1. \tag{10}$$

In the present section we shall assume that the inequality opposite to (10) is satisfied, so that the influence of recombination on the distribution function can be neglected. Therefore, if $2\lambda < \omega_{ph}$, the solution of the steady-state Eq. (6) will be a Fermi function with the chemical potential equal to zero[4]

$$n_p=[e^{\varepsilon/T}+1]^{-1}; \quad n_p=0, \quad T=0. \tag{11}$$

As $T \to 0$ the number of quasiparticles tends to zero, but the distribution function of the electrons (holes)

$$f_p=\langle a_p^+a_p\rangle\approx v_p^2(1-n_p)+u_p^2n_p$$

has the form

$$f_p=v_p^2,$$

i.e., the electrons fill the bottom of the conduction band up to μ_0 (here and later on we shall assume that λ is small in comparison with μ_0).

However, if condition (2) is satisfied, the annihilation term S_p^A vanishes since the law of energy conservation for the annihilation process, expressed by $\delta(\epsilon+\epsilon'-\omega_q)$, cannot be satisfied. This means that the phonon energy is not sufficient for the electron to overcome the gap and emit a phonon. Allowing for $S^A = 0$, the solution of Eq. (6) will be

$$n_p=\left[\exp\frac{\varepsilon-\mu}{T}+1\right]^{-1}, \quad n_p=\begin{cases}1, & \varepsilon<\mu\\0, & \varepsilon>\mu\end{cases}, \quad T=0, \tag{12}$$

where the chemical potential μ of the quasiparticles is determined from the condition

$$\sum_p v_p^2(1-n_p)^2=\sum_p u_p^4n_p^2. \tag{13}$$

Condition (13) is obtained if Eq. (6) is integrated over the momentum. Substituting (12) into (13), we find the following result for μ

$$(1+\mu/\mu_0)^{3/2}-(1-\mu/\mu_0)^{3/2}=1, \quad \mu/\mu_0\approx 1/3. \tag{14}$$

The electrons in the conduction band are distributed in the following way:

$$f_p\approx\begin{cases}u_p^2, & \varepsilon<\mu\\v_p^2, & \varepsilon>\mu\end{cases},$$

that is, there is an inverted population within the band and between the bands. Moreover, the latter situation is attained at frequencies higher than the frequency of the perturbing field. Such a form of the distribution is due to the indirect transitions, in which electrons are created with energies in the interval from $\xi = -\mu_0$ to $\xi = \mu_0$. The gap $2\lambda > \omega_{ph}$ near $\xi = 0$ inhibits electron transitions into a state with $\xi < 0$, and they accumulate above $\xi = 0$. We note that, in contrast to the case considered in[7], no additional source is necessary for the creation of an inverted population. Recombination serves as this source.

2. STEADY-STATE ABSORPTION OF AN ELECTROMAGNETIC FIELD

In the steady-state case, the expression for the number of photons absorbed per unit time has the form (see Eq. (9))

$$Q_{st}=\frac{1}{\tau_R}\sum_p [v_p^4(1-n_p)^2+u_p^4 n_p^2]. \quad (15)$$

The absorption for $1/\tau_{ph} \ll 2\lambda \ll \omega_{ph}$ was found earlier in[11] where, by using Eq. (11) for $T = 0$, the following result for Q_{st} is obtained from Eq. (15):

$$Q_{st}^{(1)} \approx p_0^3/3\pi^2\tau_R. \quad (16)$$

This result, which agrees with[1] to within terms λ/μ_0, indicates that the number of absorbable photons does not depend on the amplitude of the field.

Now let us consider the absorption for $2\lambda > \omega_{ph}$. Since the annihilation term vanishes, it is generally necessary to take recombination into consideration. The recombination processes make $n(\xi)$ asymmetric with respect to ξ, since the creation of quasiparticles occurs for $\xi < 0$, and the destruction of quasiparticles occurs for $\xi > 0$. However, if scattering by phonons were not present, the quasiparticles would predominantly occupy the region $\xi < 0$

$$n(\xi) \approx v^2.$$

The scattering of quasiparticles by phonons tends to symmetrize the distribution function $n(\xi)$ with respect to ξ due to transitions from a state with $\xi < 0$ into a state with $\xi > 0$, which occur with a probability proportional to $\lambda^2/\tau_{ph}\xi^2$, $|\xi| \gg \lambda$ (we note that these transitions correspond to the absorption of photons involving the emission of phonons). The competition of these transitions and the recombinations determine the degree of symmetry of n. The ratio of the probability for recombination to the probability for scattering by phonons with ξ of different sign is characterized by the parameter η (see Eq. (10)).

However, both of these processes are small (for $|\xi| \gg \lambda$) in comparison with the scattering by phonons with $\xi\xi' > 0$ (corresponding to the process of intraband relaxation). This is precisely what determines the dependence of the distribution function n on $|\xi|$. Taking what has been said above into account, one can show that, to within terms of order τ_{ph}/τ_R, $\lambda^2/a^2\mu_0^2$ (where $a\mu_0$ is a quantity much larger than λ), the function $n(\xi)$ takes the form ($T = 0$)

$$n(\xi) \approx \begin{cases} 1, & -b\mu_0 < \xi < a\mu_0 \\ 0, & -\mu_0 < \xi < -b\mu_0, \quad \mu_0 a < \xi \end{cases}. \quad (17)$$

By substituting (17) into (13), and also into (6), and integrating in the interval $\xi < 0$, we obtain equations for the determination of a and b

$$1+(1-b)^{1/2}=(1+a)^{1/2}, \quad (18)$$

$$\left(\frac{1}{a}-\frac{1}{b}\right)=\eta(1-b)^{1/2}, \quad \eta=\frac{\tau_{ph}}{\tau_R}\left(\frac{\mu_0}{\lambda}\right)^2, \quad \frac{1}{\tau_{ph}}=\frac{g^2 m}{\pi}. \quad (19)$$

Here it was assumed that $g^2(q) = g^2/q$, $2\lambda \gg \omega_{ph}$ (in order to simplify the calculations), and it has been taken into consideration that $a, b \gg \lambda/\mu_0$.

If $\eta \ll 1$ then, in agreement with the results of Sec. 1, we find from Eqs. (18) and (19) that $a \approx b \approx \mu$, where μ is determined by an equation that coincides with (14). Therefore, one can neglect recombination if the condition $\eta \ll 1$ is satisfied.

In the opposite limiting case, $\eta \gg 1$, we have

$$b \approx 1-(3/2\eta)^{1/4}, \quad a \approx (2/3\eta)^{1/4}. \quad (20)$$

The electrons turn out to be localized in intervals $0 < \xi/\mu_0 < a$ and $-1 < \xi/\mu_0 < -b$ which are narrow in comparison with μ_0. We note that, just as was assumed, the quantity $a\mu_0 \approx \lambda(\tau_R/\tau_{ph})^{1/2}$ is much larger than λ.

Let us find the absorption at $\eta \gg 1$. Substituting (17) into (15), we obtain after certain calculations

$$Q_{st}^{(2)}=\frac{2}{3\pi^2}\frac{p_0^3}{\tau_R}(1-b)^{3/2} \approx \frac{p_0\lambda m}{(\tau_R\tau_{ph})^{1/2}}\frac{2}{\pi^2}\left(\frac{2}{3}\right)^{1/2} \quad (21)$$

It is interesting to compare (21) with expression (16) for the absorption for $2\lambda < \omega_{ph}$ and with the absorption Q_{weak} of a weak field. From the ratio of the absorption given by expression (21) to the weak field absorption $Q_{weak} = (4/\pi)p_0\lambda^2 m$,

$$Q_{st}^{(2)}/Q_{weak} \approx 1/\lambda(\tau_{ph}\tau_R 6)^{1/4}\pi$$

it is clear that the strong field absorption is substantially smaller. If we take $\tau_{ph} = 10^{-12}$, $\tau_R = 10^{-10}$, $\lambda = 10^{13}$ to 10^{14} ($P = 10^7$ to 10^9 W/cm^2), then $Q_{st}^{(2)}/Q_{weak} \approx 10^{-2}$ to 10^{-3}.

It follows from a comparison of expressions (21) and (16) that

$$Q_{st}^{(2)}/Q_{st}^{(1)} \approx (6/\eta)^{1/4}, \quad (22)$$

that is, the absorption is reduced by $(\eta/6)^{1/2}$ times for $2\lambda > \omega_{ph}$. With a further increase of the field, Q_{st} increases but the parameter η decreases. For certain values of λ the value of η becomes small. Then, as is indicated above, we have

$$a \approx b = \mu \approx \mu_0/3.$$

For fields $\lambda > \mu_0(\tau_{ph}/\tau_R)^{1/2}$ the absorption is given by

$$Q_{st} \approx 2^{7/2}p_0^3/3^{7/2}\pi^2\tau_R \quad (23)$$

and constitutes approximately 1.08 times the value of expression (16). Thus, the dependence of the absorption on the amplitude of the field has a nonmonotonic character, since it contains regions where $\partial Q/\partial\lambda < 0$.

3. NONSTATIONARY ABSORPTION OF AN ELECTROMAGNETIC FIELD

As already mentioned above, our investigation does not cover very early instants of time, i.e., times of order $t \sim \lambda^{-1}$. We are interested in later times $t \gg \lambda^{-1}$, omitting coherence effects of the type associated with the phenomenon of self-transparency. Such an approach makes sense for pulses having a duration $\tau \gg \lambda^{-1}$, since the contribution of the leading edge to the absorption will be negligible. It should be noted that pulses of such a duration have apparently been experimentally investigated.[8, 9]

We shall assume that the field in the leading edge of the wavefront increases adiabatically, i.e., during a time $t \gg \lambda^{-1}$. In this connection the state of the system at each instant of time is described in terms of the operators α_p and β_p, defined by the transformations (3) with the instantaneous value of λ. In the adiabatic approximation the quantum numbers are conserved quantities, that is, in our case the occupation numbers $\alpha_p^+\alpha_p$ and $\beta_p^+\beta_p$. Consequently, after switching on the interaction their values coincide with the values in the absence of a field ($\lambda = 0$). Assuming that the semiconductor is found in the ground state in the absence of the field, we have[3]

$$n_p^{\alpha}=n_p^{\beta}=n_p=\begin{cases}0, & \xi>0\\ 1, & \xi<0\end{cases}. \tag{24}$$

Thus, in the state which appears after the passage of the leading edge, the quasiparticle distribution has the form of a Fermi step-function with limiting momentum p_0. Below we shall call this instant of time the initial instant and set t = 0 in it.

The Fermi distribution for the quasiparticles is valid only in the absence of collisions. Collisions with phonons lead to a relaxation of these distributions. The relaxation is described by Eqs. (6) and (7) for the quasiparticle distribution functions. In what follows we assume that condition (2) is satisfied, so that $S^A = 0$. First of all let us calculate the absorption at t = 0. Having substituted the distribution functions given by Eq. (24) into Eq. (8), we obtain

$$Q(0)\approx 2p_0\lambda m/\tau_{ph}\pi, \tag{25}$$

considering as usual that $\tau_{ph} \ll \tau_R$ and assuming that $\omega_{ph} \ll 2\lambda$.

If expression (25) is compared with the absorption Q_{weak} in a weak field, the absorption in a strong field is reduced by $\lambda\tau_{ph}$ times, as follows from the ratio

$$Q(0)/Q_{weak}\approx {}^1/_2\lambda\tau_{ph},$$

This is due to the fact that the direct transitions are saturated in a strong field, and the absorption occurs via the emission of phonons at a rate proportional to $1/\tau_{ph}$.

The interaction with phonons and the recombinations lead to a change in the quasiparticle distribution during the course of time, and consequently lead to a change in the number of absorbable photons. In order to find the time dependence of the absorption, it is necessary to obtain the solution of the time-dependent equation (6). The problem simplifies considerably if times $t > \tau_{ph}$ are considered, which are also of fundamental interest for pulses having a duration $\tau > \tau_{ph}$.

Analysis of Eq. (6) shows that the quasiparticles participate in two relaxation processes having significantly different relaxation times. In the first place there is the scattering of quasiparticles with identical ξ and ξ' ($\xi\xi' > 0$) during a time interval of order τ_{ph}. In the second place, there is the scattering of quasiparticles with different ξ and ξ' ($\xi\xi' < 0$) during a time interval of order $\tau_{ph}(\mu_0/\lambda)^2$, and also the destruction and creation of quasiparticles due to recombination during the time τ_R. By using considerations similar to those presented in Sec. 2, one can show that the distribution function $n(\xi)$ has the form (compare with Eq. (17))

$$n(\xi, t)\approx\begin{cases}1, & -b(t)<\xi/\mu_0<a(t)\\ 0, & -\mu_0 b(t)>\xi,\ \xi>\mu_0 a(t)\end{cases} \tag{26}$$

to within terms of order τ_{ph}/τ_R.

Substituting (26) into (6) and integrating with respect to p in the ξ interval $[-\mu_0, 0]$ and then $[-\mu_0, \infty]$, we obtain two equations for a(t) and b(t):

$$(1-b)^{1/2}\frac{db}{dt}=-\frac{1}{\tau_{ph}}\left(\frac{\lambda}{\mu_0}\right)^2\left(\frac{1}{a}-\frac{1}{b}\right)+\frac{1}{\tau_R}(1-b)^{3/2}, \tag{27}$$

$$\frac{d}{dt}[(1+a)^{3/2}-(1-b)^{3/2}]=-\frac{1}{\tau_R}[(1+a)^{3/2}-(1-b)^{3/2}-1] \tag{28}$$

with the initial conditions (see Eq. (24))

$$b(0)\approx 1,\quad a(0)\approx 0.$$

Equation (28) admits the solution

$$(1+a(t))^{3/2}-(1-b(t))^{3/2}=1-C\exp\{-t/\tau_R\},$$

which takes the form

$$(1+a(t))^{3/2}-(1-b(t))^{3/2}=1. \tag{29}$$

upon taking the initial condition into account.

This equation reflects the fact that the total number of quasiparticles is conserved (in the assumed approximation), resulting from the fact that the creation and annihilation of the quasiparticles occurs via recombination. Expressing a(t) from Eq. (29) and substituting it into (27), we obtain the following equation for b(t′):

$$(1-b(t'))^{1/2}\frac{db(t')}{dt'}=-\frac{1}{[1+(1-b)^{3/2}]^{1/3}-1}+\frac{1}{b}+\eta(1-b)^{3/2}, \tag{30}$$
$$t'=t/\bar{\tau}_{ph},\quad \bar{\tau}_{ph}=\tau_{ph}(\mu_0/\lambda)^2.$$

First let us consider the case when $\eta \gg 1$. Taking into consideration that, according to Eq. (20) the function b(t) is close to unity, one can neglect the second term in (30) for all moments of time. Introducing the function z = 1 − b, we obtain the following equation for it:

$$z^{1/2}\frac{dz}{dt'}=\frac{3}{2z^{3/2}}-\eta z^{3/2},\quad z(0)=0, \tag{31}$$

whose solution is of the form

$$z(t')=\left[\frac{3}{2\eta}(1-\exp\{-3\eta t'\})\right]^{1/3}. \tag{32}$$

One can represent it in a different form by changing to the time t

$$b(t)=1-\left[\frac{3}{2\eta}\left(1-\exp\left\{-3\frac{t}{\tau_R}\right\}\right)\right]^{1/3}. \tag{33}$$

If the duration of the pulse is longer than the recombination time τ_R, the function b(t) approaches a steady-state value which coincides with expression (20). For pulses having a duration shorter than the lifetime, we have

$$z(t')\approx(9t'/4)^{1/3},\quad a(t')\approx(t')^{1/2}. \tag{34}$$

Thus, the original step of the quasiparticle distribution is shifted into the region of positive $\xi > 0$. The rate of displacement is maximal at first, but then decreases, tending to zero.

The absorption also decreases with time, reaching the minimal value (21) for $t > \tau_R$. In the time interval $t \ll \tau_R$ the absorption has the form

$$Q(t)\approx\frac{p_0\lambda m}{\pi^2\tau_{ph}}\left(\frac{\tau_{ph}}{t}\right)^{1/2}. \tag{35}$$

The reduction of the absorption with time is due to the fact that, as the states near $\xi = 0$ are filled by quasiparticles, the probability of indirect transitions with $\xi\xi' < 0$ decreases in proportion to $\lambda^2/a^2(t)\mu_0^2$. Thus, the absorption experiences an additional reduction due to the switching off of the quasiparticle annihilation for $2\lambda > \omega_{ph}$.

Let us consider the case when the parameter η is small. If $\eta \ll 1$, then in Eq. (30) the last term can be neglected since the stationary state is established during a time interval $\bar{\tau}_{ph}$, which is smaller than τ_R. Let us consider the solutions for times $t \ll \bar{\tau}_{ph}$ and $t \sim \bar{\tau}_{ph}$. In the first case the results coincide with expressions (34) and (35). When the time becomes approximately equal to $\bar{\tau}_{ph}$, the function b(t) tends to $b_0 = a_0$ according to the law

$$b(t)-b_0 \propto \exp(-\beta t'),\ \beta=\frac{(1+b_0)^{1/2}+(1-b_0)^{1/2}}{b_0^2(1-b_0^2)^{1/2}},$$

where b_0 is determined from Eq. (18), $b_0 \approx 1/3$.

The absorption in this same time

$$Q(t) \propto p_0 \lambda m \lambda \beta / \tau_{ph} \mu_0 \qquad (36)$$

decreases by approximately λ/μ_0 times in comparison with expression (25).

CONCLUSION

It follows from the cited investigation that, in the presence of strong fields, one can expect a decrease in the absorption without occupation of the bottom of the bands by the electrons. Furthermore, for $2\lambda > \omega_{ph}$ the absorption is reduced still further, which leads to a nonmonotonic dependence of the number of absorbable photons on the amplitude of the field. These effects may introduce a substantial contribution to the decrease in the absorption, which is observed experimentally.[8, 9]

In addition, if the condition $2\lambda > \omega_{ph}$ is satisfied, one can expect the following effects: an inverted population within the band, and an inverted population between the bands at frequencies higher than the frequency of the perturbing field. The last two effects also take place in the nonstationary case and can be observed in the enhancement of probing pulses of a field of the appropriate frequency.

One more effect is possible in the nonstationary case. Let a pulse with $2\lambda > \omega_{ph}$ pass through the semiconductor, creating an inverted population at frequencies above the frequency ω of the field. After the passage of the pulse, the electrons giving energy to the phonons sink to the bottom of the band. At a certain instant of time, treater than τ_{ph}, they may create an inverted population at the frequency ω. Therefore, a second pulse of the field with the same frequency, following after the first, may be intensified.

We have carried out the analysis on a model which takes the interaction with phonons into account. This is adequate for semiconductors whose interelectron collision time τ_{ee} is greater than τ_{ph} (semiconductors having a small effective mass and a large dielectric constant). If the opposite inequality is satisfied, the results remain qualitatively similar to the obtained results, except for the replacement of ω_{ph} by the plasma frequency and τ_{ph} by τ_{ee}.

Intraband absorption of an electromagnetic field leads to a heating up of the electrons and holes. Let us estimate this effect. From the condition for balancing the energy obtained by the electrons from the field and the energy given to the lattice, we obtain the following expression for the change in the temperature of the electron gas (see[6, 12]):

$$\delta T=\left(\frac{eE}{\omega}\right)^2\frac{\tau_\epsilon}{2m\tau_{im}k}.$$

Here τ_ϵ and τ_{im} denote the energy and momentum relaxation times of the electron, and k is the Boltzmann constant. In the situation under consideration the electrons possess enough energy for the emission of optical phonons, so that $\tau_\epsilon = \tau_{im}$. It is natural to compare the quantity $k\delta T$ with the energy $a\mu_0 = \lambda(\tau_R/\tau_{ph})^{1/2}$, characterizing the region of localization of the electrons:

$$\frac{k\delta T}{a\mu_0}=\left(\frac{eE}{\omega}\right)^2\frac{1}{2ma\mu_0}=\frac{4\pi e^2 P}{\omega^2 m\lambda(\tau_R/\tau_{ph})^{1/2}}.$$

Assuming P = (10^7 to 10^9) W/cm^2, $\omega = 10^{15}$ sec^{-1}, and m = 10^{-28} g, we obtain a value between 10^{-3} and 10^{-2} for the ratio $k\delta T/a\mu_0$, that is, the heating up is insignificant.

It should be noted that if the frequency ω of the field is greater than the width of the forbidden band, intraband absorption is generally absent.

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[1)]The effect of a strong electromagnetic field on the intraband motion of the electrons was considered in [6], and it was shown that the change in the electron spectrum is negligible up to fields satisfying condition (2).

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