Multiphoton intrinsic photoeffect in two-band semiconductors

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An expression different from that hetertofore known is obtained for the intrinsic-photoeffect probability within the framework of a semiconductor model with unequal electron and hole effective masses. Allowance for the effective-mass difference leads to the onset of quantum beats in the photocurrent. The possibility is demonstrated of increasing the photoeffect probability by an exponentially large amount in the presence of an electric field.

The multiphoton intrinsic photoeffect in semiconductors was first considered in Ref. 1. In this and subsequent studies, use was made of the effective-mass approximation. The validity of using this approximation to describe essentially multiphoton processes (with absorption of a large number of photons) is not obvious and raises serious doubts. Principal among them is that to calculate multiphoton matrix elements in the S-matrix approach it is necessary to know the dispersion law far from the mass shell—for large imaginary quasimomentum values. Since the effective-mass approximation is valid only for small quasimomenta, this approximation can lead in the case of multiphoton processes to incorrect results. In Ref. 2 and in similar studies the intrinsic multiphoton photoeffect was considered on the basis of the Kane dispersion law for the narrow-gap semiconductor. The electron-hole pair production probability was calculated in the adiabatic approximation, and it was assumed in addition that the effective masses of the particle and of the hole are equal. In real cases these masses are unequal, and furthermore the adiabatic approximation cannot be used if the electromagnetic field is turned on instantaneously.

We consider in the present paper the multiphoton intrinsic effect in a two-band semiconductor. The two-band semiconductor approximation makes it possible to solve the problem without assuming that the particle and hole effective masses are equal, and to forgo also the adiabatic approximation.

In Sec. I we consider a two-band semiconductor. In Sec. 2 we find the electron wave functions in electric and electromagnetic fields. In Sec. 3, using the expressions obtained in Sec. 2 for the electron final-state wave functions we calculate in the S-matrix approach the electron–hole production probability per unit time. In Sec. 4 we consider the influence of a constant electric field on the probability of the intrinsic photoeffect.

1. INTRINSIC SEMICONDUCTOR

The electron wave function in the two-band approximation

$$\Psi(x,t)=\exp\left(-\frac{i}{\hbar}Et+\frac{i}{\hbar}px\right)u(p),$$

$$u(p)=a_1(p)u_1(x)+a_2(p)u_2(x), \tag{1}$$

$u_1(x)$ and $u_2(x)$ are Bloch functions corresponding to bands 1 and 2. All the operators acting on $u(p)$ are 2×2 matrices. The Hamiltonian of an electron in a semiconductor is given in the two-band approximation by

$$\hat{H}_0=\frac{\Delta E}{2}\begin{bmatrix}1&0\\0&-1\end{bmatrix}+\frac{p^2}{2m}\begin{bmatrix}1&0\\0&1\end{bmatrix}+\frac{q}{m}\begin{bmatrix}0&p_{12}\\p_{21}&0\end{bmatrix},$$
$$p_{12}=p_{21}^{*}=\langle u_1|\hat{p}|u_2\rangle. \tag{2}$$

The dispersion law is

$$E_{1,2}(p)=\frac{p^2}{2m}\pm\left[\left(\frac{\Delta E}{2}\right)^2+\frac{(pp_{12})(p_{21}p)}{m^2}\right]^{1/2}. \tag{3}$$

We expand (3) for small quasimomentum values

$$E_{1,2}(p)\approx\frac{p^2}{2m}\pm\left[\frac{\Delta E}{2}+\frac{(pp_{12})(p_{21}p)}{m^2\Delta E}\right]=p\left(\frac{1}{2m_{1,2}^{*}}\right)p. \tag{4}$$

It must be remembered that for a three-dimensional crystal the quantities $1/m_{1,2}^{*}$ are tensors

$$\left(\frac{1}{2m_{1,2}^{*}}\right)_{ij}=\frac{1}{m}\delta_{ij}\pm\frac{2(p_{12})_i(p_{21})_j}{m^2\Delta E}. \tag{5}$$

To describe a semiconductor with a scalar mass, we must make the substitution

$$(p_{12})\otimes(p_{21})\to|p_{12}|^2\delta_{ij}\equiv q^2\delta_{ij}. \tag{6}$$

Then

$$\frac{1}{m_{1,2}^{*}}=\frac{1}{m}\pm\frac{2q}{m^2\Delta E},\qquad \frac{q}{m^2\Delta E}=\frac{1}{4}\left(\frac{1}{m_1^{*}}-\frac{1}{m_2^{*}}\right)=\frac{1}{4}\frac{1}{\mu}. \tag{7}$$

Relation (7) permits the final expressions to be written in terms of the measurable quantities m_1^{*} and m_2^{*}.

2. ELECTRON WAVE FUNCTIONS IN ELECTRIC AND ELECTROMAGNETIC FIELDS

Electric field. In the presence of a constant electric field, the Hamiltonian of an electron in a semiconductor is

$$\hat{H}=\hat{H}_0-e\mathscr{E}x, \tag{8}$$

where $\mathscr{E}$ is the electric field strength. We seek a stationary wave function of the Hamiltonian (8) in the form

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$$\Psi_a(x,t)=\exp\left(-\frac{i}{\hbar}Et\right)\int_c \exp\left(\frac{i}{\hbar}px\right)u(p)\,dp. \qquad (9)$$

The integration contour is chosen to have the integrand tend to zero at its end points. The equation for $u(p)$ with allowance for (1), (2), (8), and (9) is

$$E\begin{pmatrix}1&0\\0&1\end{pmatrix}u(p)=\left\{\frac{\Delta E}{2}\begin{pmatrix}1&0\\0&-1\end{pmatrix}+\frac{p^2}{2m}\begin{pmatrix}1&0\\0&1\end{pmatrix}+\frac{pq}{m}\begin{pmatrix}0&1\\1&0\end{pmatrix}\right\}u - i(e\mathscr{E}\hbar)\frac{d}{dp}u(p). \qquad (10)$$

Consider the unitary transformation

$$U(\varphi)=\begin{pmatrix}\cos\varphi&\sin\varphi\\-\sin\varphi&\cos\varphi\end{pmatrix}, \quad u=U(\varphi)\tilde{u}. \qquad (11)$$

We substitute (11) in (10) and multiply the result from the left by $U^{-1}(\varphi)=U(-\varphi)$:

$$E\begin{pmatrix}1&0\\0&1\end{pmatrix}\tilde{u}=\hat{H}_0\tilde{u}-i(e\mathscr{E}\hbar)\frac{d}{dp}\tilde{u} -i(e\mathscr{E}\hbar)U(-\varphi)\left(\frac{d}{dp}U(\varphi)\right)\tilde{u}, \qquad (12)$$

where $\hat{\tilde{H}}_0=U(-\varphi)\hat{H}_0U(\varphi)$,

$$U(-\varphi)\frac{d}{dp}U(\varphi)-\frac{d}{dp}\varphi(p)\begin{pmatrix}1&0\\0&-1\end{pmatrix}, \qquad (13)$$

$$(\tilde{H}_0)_{12}=\frac{\Delta E}{2}\sin 2\varphi+\frac{pq}{m}\cos 2\varphi=-(\tilde{H}_0)_{21}, \qquad (14)$$

$$(\tilde{H}_0)_{11,22}=\frac{p^2}{2m}\pm\left(\frac{\Delta E}{2}\cos 2\varphi-\frac{pq}{m}\sin 2\varphi\right)=\tilde{E}_{1,2}. \qquad (15)$$

We determine the function $\varphi(p)$ from the diagonalization condition of the right-hand side of (12):

$$\frac{\Delta E}{2}\sin 2\varphi+\frac{pq}{m}\cos 2\varphi=ie\mathscr{E}\hbar\frac{d}{dp}\varphi. \qquad (16)$$

We solve (16) by perturbation theory, assuming the right-hand side to he small:

$$\varphi_0=\frac{1}{2}\operatorname{arctg}\frac{2pq}{m\Delta E}, \qquad (17)$$

$$(\tilde{E}_{1,2})_0=E_{1,2}=\frac{p^2}{2m}\pm\left[\left(\frac{\Delta E}{2}\right)^2+\left(\frac{pq}{m}\right)^2\right]^{1/2} \qquad (18)$$

The condition that the discarded terms be small is

$$e\mathscr{E}\hbar\frac{q}{m}\frac{\Delta E}{2}\left[\left(\frac{\Delta E}{2}\right)^2+\left(\frac{pq}{m}\right)^2\right]^{-1/2}\ll 1, \qquad (19)$$

or, with allowance for (7),

$$\frac{e\mathscr{E}\hbar}{\mu^{1/2}}\frac{1}{(\Delta E)^{3/2}}\ll 1. \qquad (20)$$

The condition (20) means that the probability of tunneling from band to band is exponentially small. Use of the zero-order solution φ_0 (17) corresponds to a transition to a basis of proper wave functions of a two-band semiconductor. Substituting (17) in (11) we get

$$\tilde{u}_{1,2}(p)=\exp\left\{(i\hbar e\mathscr{E})^{-1}\int_0^p[(E-E_{1,2}(p'))]\,dp'\right\}, \qquad (21)$$

$$\Psi_a=\exp\left\{-\frac{i}{\hbar}Et\right\}\int_c\exp\left(\frac{i}{\hbar}px\right)\tilde{u}_{1,2}(p)\,dp. \qquad (22)$$

Electromagnetic field. The Hamiltonian of an electron in a semiconductor in the presence of a uniform electromagnetic field is

$$\hat{H}=\hat{H}_0(\tilde{p}), \quad \tilde{p}=p+\frac{e\mathscr{E}}{\omega}\sin\omega t, \qquad (23)$$

where $\mathscr{E}$ is the electromagnetic field intensity.

The equation for the quasistationary wave function Ψ_b is

$$i\hbar\frac{d}{dt}\Psi_b=\hat{H}\Psi_b. \qquad (24)$$

We seek Ψ_b in the form

$$\Psi_b=\exp\left(\frac{i}{\hbar}px\right)u(p,t). \qquad (25)$$

Using the transformation $U(\varphi(t))$ (11), we get

$$u_{1,2}(t)=\exp\left\{-\frac{i}{\hbar}\int_0^t E_{1,2}(\tilde{p}(t'))\,dt'\right\}. \qquad (26)$$

The condition for the validity of (26) is

$$\frac{e\mathscr{E}\hbar\cos\omega t}{\mu^{1/2}(\Delta E)^{3/2}}\ll 1. \qquad (27)$$

Electric and electromagnetic fields. The Hamiltonian of an electron in a semiconductor in the presence of an electric field and of an electromagnetic field is

$$\hat{H}=\hat{H}_0(\tilde{p})-e\mathscr{E}_2x, \quad \tilde{p}=p+\frac{e\mathscr{E}_1}{\omega}\sin\omega t. \qquad (28)$$

We seek a quasistationary wave function in the form

$$\Psi_{ab}=\exp\left(-\frac{i}{\hbar}Et\right)\int_c\exp\left(\frac{i}{\hbar}px\right)u(p,t)\,dp. \qquad (29)$$

Just as above, we use the transformation $U(\varphi(p,t))$. We obtain an equation for $\tilde{u}_{1,2}(p,t)$:

$$i\hbar\frac{\partial}{\partial t}\tilde{u}_{1,2}=\left[E_{1,2}(\tilde{p})-E-i\hbar e\mathscr{E}_2\frac{\partial}{\partial p}\right]\tilde{u}_{1,2}. \qquad (30)$$

This equation is valid under the condition

$$\frac{e\mathscr{E}_1\hbar\cos\omega t+e\mathscr{E}_2\hbar\sin\omega t}{\mu^{1/2}(\Delta E)^{3/2}}\ll 1. \qquad (31)$$

It is necessary to obtain an approximate solution of (30) such that Ψ_{ab} goes over into Ψ_a or Ψ_b when one of the fields is turned off. This requirement is satisfied by the solution

$$\tilde{u}_{1,2}=\exp\Big\{(i\hbar e\mathscr{E}_2)^{-1}\int_0^p [E_{1,2}(p')-E]\,dp'$$

$$-\frac{i}{\hbar}\int_0^t [E_{1,2}(\tilde{p})-E_{1,2}(p)]\,dt'$$

$$-\frac{i\hbar e\mathscr{E}_2}{\hbar^2}\int_0^t\frac{\partial}{\partial p}\int_0^{t'}[E_1(\tilde{p})-E_1(p)]\,dt''\,dt'\Big\}. \qquad (32)$$

The condition for the validity of (31) is that the terms discarded in the solution (30) be small:

$$\frac{e\mathscr{E}_2\hbar}{\mu^{1/2}}\frac{1}{(\Delta E)^{1/2}}\frac{\Delta E}{\hbar\omega}\ll 1. \qquad (33)$$

3. MULTIPHOTON IONIZATION

The calculation of the intrinsic multiphoton photoeffect in a semiconductor differs from the analogous problem for systems of the atomic type. In contrast to an atom, where the perturbation of the discrete state by the field is negligibly small, in a semiconductor the multiphoton transition is from a continuum into a continuum (from a band into a band). An electromagnetic field perturbs strongly the electron states (hole states) in the upper and lower bands. This circumstance must be taken into account.

We assume that the electromagnetic field is turned on instantaneously and is instantaneously turned off after a time T; the S matrix of such a process is

$$S_{21}(p)=\frac{i}{\hbar}\sum_{i=1}^{2}\langle\Psi^2(p,t=0)\,|\,\Psi_b{}^i(p,t=0)\rangle$$

$$\times\langle\Psi_b{}^i(p,t=T)\,|\,\Psi^1(p,t=T)\rangle, \qquad (34)$$

where the superscript denotes the number of the band.

The ionization probability per unit time is

$$W_b(p)=\lim_{T\to\infty}\frac{1}{T}|S_{21}|^2. \qquad (35)$$

To find the S matrix we must calculate the following integrals [see Eqs. (3) and (26)]:

$$J_{1,2}=\int_0^{2\pi}\exp\Big\{-\frac{i}{\hbar}\int_0^t [E_{1,2}(\tilde{p}(t'))-E_{2,1}(p)]\,dt'\Big\}dt. \qquad (36)$$

The integrals (36) differ from their analogs in Ref. 2. In the S-matrix approach an electromagnetic field perturbs only the final state, so that a time dependence appears in the integrals (36) only in the expression for the final-state energy. In the adiabatic approach the electromagnetic field perturbs all the states, in both the lower and upper bands. The time dependence in the analogous integrals of Ref. 2 appears correspondingly in all the expressions for the energies, of both the initial and final states. This accounts for the difference between our final expressions and those of Ref. 2.

An estimate of $J_{1,2}$ with exponential accuracy is obtained by the saddle-point method. To avoid the difficulties connected with the square-root singularity in the exponent of (36) [see Eq. (3)], and correspondingly the bypassing of the branching points along different edges of the cuts, one can replace the integrals along the cut by bypassing the branching points on another sheet of the entire Riemann surface, on which the integrand is analytic everywhere and has no singularities. The saddle points t_0 are determined from the equation

$$\frac{\tilde{p}^2(t_0)}{2m}\pm\left[\left(\frac{\Delta E}{2}\right)^2+\left(\frac{\tilde{p}(t_0)q}{m}\right)^2\right]^{1/2}$$

$$=\frac{p^2}{2m}\mp\left[\left(\frac{\Delta E}{2}\right)^2+\left(\frac{pq}{m}\right)^2\right]^{1/2},$$

$$\tilde{p}(t_0)=p+\frac{e\mathscr{E}}{\omega}\sin\omega t_0, \qquad (37)$$

the plus or minus sign corresponds to different sheets of the Riemann surface. The condition (27) coincides at the saddle point t_0 with the requirement

$$n=\frac{\Delta E}{\hbar\omega}\gg 1.$$

Calculation of the probability of the multiphoton intrinsic photoeffect for an arbitrary value of p encounters no fundamental difficulties, but in view of the unwieldy expressions we present the result for only the most interesting case—ionication from boundary to boundary of the band, i.e., for $p^2/2\mu\ll\Delta E$. We assume that $p=0$. We have then

$$\frac{\tilde{p}^2(t_0)}{2m}=\frac{2q^2}{m}\mp\Delta E,\qquad \tilde{p}(t_0)=\frac{e\mathscr{E}}{\omega}\sin\omega t_0. \qquad (38)$$

In the multiphoton limit

$$\left[\frac{e\mathscr{E}}{2\omega}\right]^2\frac{1}{2\mu}\ll n=\frac{\Delta E}{\hbar\omega}, \qquad (39)$$

so that we can assume in the calculation of $J_{1,2}$

$$\frac{e\mathscr{E}}{\omega}\sin\omega t\approx\frac{e\mathscr{E}}{2i\omega}e^{i\omega t}\equiv\alpha, \qquad (40)$$

and then (for $p=0$)

$$-\frac{i}{\hbar}\int_0^{t_0}[E_{1,2}(\tilde{p}(t'))-E_{2,1}(p)]\,dt'$$

$$\approx-\frac{1}{\hbar\omega}\Big\{\frac{\alpha^2}{4m}\pm\frac{\Delta E}{2}\ln\frac{\alpha}{\alpha_0}\pm\Big[\Big(\frac{\Delta E}{2}\Big)^2$$

$$+\Big(\frac{\alpha q}{m}\Big)^2\Big]^{1/2}\pm\frac{\Delta E}{4}\ln\frac{[(\Delta E/2)^2+(\alpha q/m)^2]^{1/2}-(\Delta E)/2}{[(\Delta E/2)^2+(\alpha q/m)^2]^{1/2}+(\Delta E)/2}\Big\}\Big|_\beta^{\alpha_0}$$

$$=-\frac{1}{\hbar\omega}\Big\{\frac{q^2}{m}\mp\frac{\Delta E}{2}-\frac{\Delta E}{2}\ln\frac{\beta^2}{2m\Delta E}\Big\}, \qquad (41)$$

where

$$\frac{\alpha_0^2}{2m}=\frac{2q^2}{m}\mp\Delta E,\qquad \beta=\frac{e\mathscr{E}}{2i\omega}\alpha_0.$$

Finally,

$$J_{1,2}=\left[\left(\frac{e\mathscr{E}}{2\omega}\right)^2\frac{1}{2m\Delta E}\right]^{n/2}\exp\left[\frac{n}{2}\left(1+\frac{m}{2\mu}\right)\right], \quad (42)$$

$$W_b(p)\approx\sum_n\frac{1}{\hbar}\left[\left(\frac{e\mathscr{E}}{2\omega}\right)^2\frac{1}{2m\Delta E}\right]^n\exp\left[(n)\left(1+\frac{m}{2\mu}\right)\right]$$

$$\times\{\delta(E_1(p)-E_2(p)+\sigma_1(p)-n\hbar\omega)$$

$$+\delta(E_1(p)-E_2(p)-\sigma_2(p)-n\hbar\omega)\}, \quad (43)$$

where

$$\sigma_{1,2}=\int_0^{2\pi/\omega}[E_{1,2}(\tilde{p}(t))-E_{1,2}(p)]\,dt\omega. \quad (44)$$

This result differs substantially from that obtained in the effective-mass approximation

$$W_{\text{e.m.}}\approx\sum_n\left[\left(\frac{e\mathscr{E}}{2\omega}\right)^2\frac{1}{2m^*\Delta E}\right]^n$$

$$\times\exp(n)\left[\delta\left(E_1(p)+E_2(p)+\left(\frac{e\mathscr{E}}{2\omega}\right)^2\frac{1}{8m^*}-n\hbar\omega\right)\right]. \quad (45)$$

The expressions for the ionization probabilities differ from one another by the exponentially large quantity

$$\frac{W_b{}^n}{W_{\text{e.m.}}{}^n}=\left(\frac{m^*}{m}\right)^n\exp\left(n\frac{m}{2\mu}\right). \quad (46)$$

The quantities $\sigma_{1,2}$ (44) have the meaning of ponderomotive potentials in the valence and conduction bands, and depend on the quasimomentum p. We use the smallness of

$$\left(\frac{e\mathscr{E}}{2\omega}\right)^2\frac{1}{2\mu}\ll\Delta E,$$

and then, retaining the first nonvanishing terms in (44), we get

$$\sigma_{1,2}=\left(\frac{e\mathscr{E}}{2\omega}\right)^2\left(\frac{1}{8m_{1,2}^*}\mp\frac{p_z^2}{\mu^2\Delta E}\right), \quad (47)$$

p_z is the quasimomentum projection on the electromagnetic-field polarization direction.

This fact means that, in contrast to the usual case, the energy spectrum of the photoelectrons in the valence band has, at a fixed electromagnetic-field frequency ω, not one δ peak whose position is determined from (45), i.e., from the relation

$$E_1(p)+\left(\frac{e\mathscr{E}}{\omega}\right)^2\frac{1}{8m_1^*}=n\hbar\omega+E_2(p), \quad (48)$$

but two peaks separated by a distance δE and having widths $\Gamma(p)$. The peak positions, δE, and $\Gamma(p)$ are given by

$$E_{1,2}(p)+\sigma_{1,2}(p)=\pm n\hbar\omega+E_{2,1}(p), \quad (49)$$

$$\delta E=\sigma_1-\sigma_2=\left(\frac{e\mathscr{E}}{\omega}\right)^2\frac{1}{8\mu}, \quad (50)$$

$$\Gamma(p)=4\left(\frac{e\mathscr{E}}{\omega}\right)^2\frac{1}{\mu}\frac{E_1(p)-E_2(p)-\Delta E}{\Delta E}\sim\frac{\delta E}{n}. \quad (51)$$

4. FRANZ–KELDYSH MULTIPHOTON EFFECT

The principal scheme for calculating the probability of the photoeffect in an electromagnetic field of intensity $\mathscr{E}_1$ in the presence of an electric field of intensity $\mathscr{E}_2$ (we assume for simplicity that the fields are parallel) is the same as for the case of multiphoton ionization (Sec. 3). We choose the initial-state wave function to be Ψ_a (21), (22) and the final states to be Ψ_{ab} (29), (32). All the peculiarities of the multiphoton processes in a semiconductor can be elucidated with the intrinsic photoeffect as the example (Sec. 3). We confine ourselves in this section therefore, after indicating the method used to calculate the entire energy spectrum, to only one problem: we consider the change of the ionization probability by application of an electric field in the case $\Delta E=n\hbar\omega$. The total ionization probability per unit time is

$$W_{\text{e.m.}}\approx\omega\int dE\sum_n|J_n|^2\delta(E_1-E_2-n\hbar\omega), \quad (52)$$

where J_n is calculated as in the preceding section. When the electric field is turned on the saddle point p_0 undergoes a small shift p_1. The electric field is weak, and p_0 is determined from the equation

$$E_1(p_0)-E_2(p_0)=n\hbar\omega=\Delta E,\quad p_0=0. \quad (53)$$

The correction to p_0-p_1 for the electric field is obtained from the equation

$$(e\mathscr{E}_2)\cdot 2\frac{p_1^2q^2}{m^2\Delta E}=-\frac{\partial}{\partial p}S_1(p_0). \quad (54)$$

Omitting the unwieldy calculations, we present directly an estimate of the exponential correction for the ionization probability per unit time following application of the electric field

$$W_{ab}=W_{ab}(\mathscr{E}_2=0)\exp(\phi), \quad (55)$$

$$W_{ab}(\mathscr{E}_2=0)=W_b, \quad (56)$$

$$\phi\sim\left[\left(\frac{e\mathscr{E}_2\hbar}{\mu^{1/2}}(\Delta E)^{-3/2}\right)\left(\frac{\Delta E}{\hbar\omega}\right)\right]^{1/2}\frac{\Delta E}{\hbar\omega}. \quad (57)$$

The value of ϕ can exceed unity noticeably. For example, its value is ~ 10 in the case $\mathscr{E}_2\approx 10^5$ eV·cm^{-1}, $\mu=0.1m$, $\Delta E=1$ eV, and $\hbar\omega=0.1$ eV.

DISCUSSION OF RESULTS

It is thus possible to obtain, within the two-band model of a semiconductor, an expression for the probability of the multiphoton intrinsic photoeffect without using the assumption that the electron and hole have equal effective masses. Nor is it necessary to assume adiabatic turning on (off) of the electromagnetic field. We have calculated here the photoeffect probability for instantaneous application of the field. The expression obtained for the probability of the intrinsic multiphoton photoeffect differs noticeably from the results obtained in the approximation of equal electron and hole masses.[2] In particular, the values of (46) and of its dependence on the effective masses differ substantially.

Further, allowance for the difference between the effective masses of the electron and hole leads to the appearance

of two maxima in the photoelectron energy spectrum. The distance between the maxima is determined by the difference between the ponderomotive potentials of the particle and hole. For an electromagnetic-field intensity $\mathscr{E} \approx 10^5$ eV·cm^{-1} and for $\mu = 0.1m$, $\Delta E = 1$ eV, and $\hbar\omega = 0.1$ eV we have $\delta E \approx 10^{-4}$ eV. The width of the peaks is $\delta E/n$. Consequently, at a fixed frequency of the perturbing electromagnetic field electron-hole states with different energies are coherently excited, and beats of frequency $\delta E/\hbar$ can be observed in the photocurrent. It must be noted that this effect can be experimentally observed only in thin films of thickness smaller than the electron mean free path, and at temperatures not higher than 10 K. the constraint on the sample temperature is determined from the condition that the temperature current be small compared with the photocurrent. From our point of view, an experimental verification of the results is of interest.

An analytic study of the multiphoton Franz–Keldysh effect in a two-band semiconductor has shown that the probability of the multiphoton photoeffect can increase exponentially in the presence of a constant electromagnetic field.

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