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On the Theory of Photon-Induced DC Hopping Conductivity in Disordered Semiconductors

By

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Based on the Kubo formalism a detailed discussion of the static hopping transport is given for the case that the electron transitions between localized gap states are caused by intense illumination of the specimen. The statistical correlations of the excited hopping system taking into account by means of the percolation theory, the resulting conductivities in the low-frequency and high-frequency limits are analyzed with respect to their spectral dependences and their orders of magnitude in relation to Mott's phonon-induced variable range hopping. Apart from its purely theoretical interest this new mechanism of photoconduction should also deserve experimental interest, in particular, because of the possibility of a frequency-controlled "reading" of the electronic density of states within the mobility gap.

Unter Verwendung des Kubo-Formalismus wird die statische Hopping-Leitfähigkeit für den Fall eingehend diskutiert, daß die elektronischen Übergänge zwischen den lokalisierten Gapzuständen durch eine intensive Beleuchtung der Probe hervorgerufen werden. Die statistischen Korrelationen des angeregten Hoppingsystems mittels der Perkolationstheorie berücksichtigend, werden die im nieder- und im hochfrequenten Grenzfall erhaltenen Leitfähigkeiten hinsichtlich ihrer spektralen Abhängigkeiten und ihrer Größenordnungen im Verhältnis zum phononindzuierten Mottschen „variable range"-Hopping analysiert. Abgesehen von seinem rein theoretischen Interesse, sollte dieser neue Mechanismus von Photoleitung auch experimentelles Interesse verdienen, insbesondere wegen der Möglichkeit eines frequenzkontrollierten „Abtastens" der elektronischen Zustandsdichte innerhalb des Beweglichkeitsgaps.

1. Introduction

Charge transport measurements in disordered semiconductors have been of considerable interest in recent years because they can provide valuable information about the electronic structure of such materials. In sufficiently disordered samples there is a significant density of spatially localized states at the Fermi level and the electrical transport is characterized by the hopping of electrons between the randomly distributed sites within the mobility gap. Today it is well known that not only thermally activated phonons or strong electric fields are able to induce this electron transfer, but also, for example, external controlled photon fields. However, only a few papers have been published up to now which investigate photon-induced hopping phenomena [1 to 7]. Photon-induced hopping was first studied by Tanaka and Fan [1]. To explain their experimental results concerning microwave absorption in p-type silicon, they assumed processes of the direct absorption of microwave photons in ionized impurity pairs without phonon participation. Later this mechanism was considered by Mott [3] in the case of the ac hopping conduction in disordered semiconductors. In contrast to such an ac field-induced hopping transport reviewed by Böttger and Bryksin [8], for example, the main concern of this paper is the analysis of that photon-induced hopping mechanism which occurs if the localized electrons get their mobility not by the direct optical absorption of photons of the ac field but by their interaction with an intensive light field within the material. The current itself is measured in a weak and static electric field. This phenomenon is obviously analogous to the usual photoconduc-

tion in the hopping region. A brief discussion of such a hopping photoconduction was first given by Bonch-Bruevich and Čápek [5]. They pointed out that one can expect a remarkably large effect although the coupling of the localized electrons with the incident photons is much smaller than, for instance, that with thermally activated phonons if only the light is intensive and the temperature is chosen to be sufficiently low.

In this paper, we present a more detailed microscopic theory of the dc hopping photoconductivity. Based on the Kubo formalism, the chosen approach takes into account both the characteristic features of the electronic structure and the statistical correlations of the excited hopping system. The investigation of the hopping photoresponse has been extended to the low-frequency optical spectral region also. The results are examined from the standpoint of their practical use.

2. General Formalism

To evaluate the photon-induced dc hopping conductivity we start from the Kubo formula for the dynamic conductivity tensor,

$$\sigma_{\mathrm{h}}(\Omega) = \frac{2\pi}{\hbar^2}\frac{e^2}{V}\left\langle \sum_{mn} \boldsymbol{R}_m \boldsymbol{R}_n \tilde{\Gamma}_{mn}(\Omega)\right\rangle_{\mathrm{c}}, \tag{1}$$

where

$$\tilde{\Gamma}_{mn}(\Omega) = \hbar\Omega \int_0^\infty \frac{\mathrm{d}\tau}{2\pi} \operatorname{Tr}\{\varrho(-\infty)\,[a_m^+(\tau)\,a_m(\tau), a_n^+ a_n]\} \exp(i\Omega^+\tau); \quad \Omega^+ \equiv \Omega + i0, \tag{2}$$

is closely related to the Fourier transform of the characteristic two-particle Green's function. The conductivity (1) represents the linear response of the coupled electron–photon system with the Hamiltonian

$$\mathcal{H} = \mathcal{H}_{\mathrm{e}} + \mathcal{H}_{\mathrm{p}} + \mathcal{H}_{\mathrm{ep}} = \sum_m \varepsilon_m a_m^+ a_m + \sum_{q\lambda} \hbar\omega_q \left(b_{q\lambda}^+ b_{q\lambda} + \frac{1}{2}\right) +$$

$$+ \sum_{mnq\lambda} g_{mn}(q\lambda)\, a_m^+ a_n (b_{q\lambda}^+ + b_{-q\lambda}) \tag{3}$$

to the weak electric ac measuring field $\boldsymbol{E}(\Omega)$. Here the Hamiltonian and, hence, the conductivity of interest are written in the so-called localization centre representation (see, for example, [9, 10]) which is based on the approximately complete orthonormal set of the localized gap states $|m\rangle$ with energies ε_m and localization centres at $\boldsymbol{R}_m$. The operators a_m^+ and a_m describe the creation and annihilation of an electron in such a state, whereas $b_{q\lambda}^+$ and $b_{q\lambda}$ are the creation and annihilation operators for a photon of momentum $\hbar\boldsymbol{q}$ and transverse polarization λ (= 1, 2). The interaction Hamiltonian $\mathcal{H}_{\mathrm{ep}}$ is characterized by the coupling function

$$g_{mn}(q\lambda) = \frac{-e}{m}\sqrt{\frac{2\pi\hbar}{\omega_q V}}\,(2\pi)^3 \int \mathrm{d}\boldsymbol{k}\, \boldsymbol{\pi}_{q\lambda} \cdot \hbar\boldsymbol{k}\psi_m^*(\boldsymbol{k})\,\psi_n(\boldsymbol{k}+\boldsymbol{q}) \tag{4}$$

in which $\boldsymbol{\pi}_{q\lambda}$ and $\psi_m(\boldsymbol{k})$ denote the unit photon polarization vector and the space Fourier transform of the one-electron wave function $\psi_m(\boldsymbol{x})$, respectively.

In deriving equation (1) it has been assumed that both the electron–phonon interaction and the perturbing ac field are adiabatically turned on at $t = -\infty$. Therefore, $\varrho(-\infty)$ is the statistical operator for the non-interacting photon and electron fields. It defines the mean occupation numbers of the unperturbed one-particle states, i.e., the well-known Fermi-function $f_{\varepsilon_m}^0$ and the photon distribution $N_{q\lambda}^0$, respectively. The latter is correlated with any realistic preparation of the external controlled radiation

field. Due to the interaction of the coupled electron–photon system with phonons treated as a part of the thermal reservoir here the initial state $\varrho(-\infty)$ is assumed to pass into a stationary distribution in the limit of $\tau \to \infty$. The phonon-induced broadening of the electronic energy levels, however, has been neglected.

The time development of the Heisenberg operators, of course, is determined by the total Hamiltonian (3), and the angular brackets in (1) indicate the configuration average reflecting the random nature of the hopping system.

Furthermore, it is to be emphasized that equation (1) was obtained by neglecting the off-diagonal elements $\delta f_{mn}(\Omega)$ of that contribution to the reduced one-electron density matrix,

$$\tilde{f}_{mn}(\Omega) = \int \frac{\mathrm{d}\tau}{2\pi} \mathrm{Tr}\,[\varrho(\tau, -\infty)\, a_m^+ a_n] \exp(i\Omega^+\tau) = f_{mn}(\Omega) + \delta f_{mn}(\Omega) \tag{5}$$

which depends on the weak electric field $\boldsymbol{E}(\Omega)$. Here $\varrho(\tau, -\infty)$ denotes the statistical operator in presence of this perturbing field. Physically such a "diagonal approximation" [10] is primarily related to the limiting case of strong localization of the effective electron states,

$$\alpha^* R^* \gg 1\,, \tag{6}$$

where α^* and R^* are the characteristic reciprocal localization radius and hopping distance of those effective electrons, respectively. Under the above condition the off-diagonal matrix elements of the coordinate operator are negligible compared with the corresponding diagonal elements. Finally, this approximation is associated with the neglect of the ac field-induced hopping transitions in the dynamic conductivity. It should be remembered that these transitions are just the direct optical absorption processes leading to Mott's familiar Ω^2-law of the ac hopping conductivity [3], for example.

On the other hand, the perturbing field-induced diagonal elements $\delta f_{mn}(\Omega)$ of the reduced one-electron density matrix (5) are connected with the spatial redistribution of the localized charges which can be described by a local change $\delta\mu_m(\Omega)$ of the chemical potential at the m-th site. Taking into account this ac field-induced population effect we ensure not only the requirement of current conservation, i.e. the total current entering each hopping site vanishes, but also the correct transition to the limit $\Omega \to 0$ with respect to the conductivity under consideration [11].

The function (2) satisfies an infinite set of coupled equations. Following Zubarev's decoupling procedure [12] this hierarchy after some straight-forward but rather lengthy calculations [13] can be reduced to a linear inhomogeneous integral equation,

$$\tilde{\Gamma}_{mn}(\Omega) = I_{mn}(\Omega) + \sum_s \int \mathrm{d}\Omega'\, \Lambda_{ms}(\Omega;\, \Omega')\, \tilde{\Gamma}_{sn}(\Omega - \Omega')\,. \tag{7}$$

If we expand here the free term $I_{mn}(\Omega)$ and the complicated kernel $\Lambda_{ms}(\Omega;\, \Omega')$ in power series in g,[1]) their leading terms are found to be of order g^2 and $g^2\Omega^{-1}$, respectively. This means, of course, that $\tilde{\Gamma}_{mn}(\Omega)$ itself cannot be developed in powers of g as Ω tends to zero because such a perturbation expansion produces divergences of the form $g^{2n}\Omega^{-n}$. Therefore, the dc conductivity of interest is given by a more general expression,

$$\sigma_\mathrm{h} = \lim_{\Omega\to 0} \sigma_\mathrm{h}(\Omega) = \left\langle \frac{1}{V} \sum_{mn} \boldsymbol{R}_{mn}\boldsymbol{R}_{mn}[(G(m,\, n) + \delta G(m,\, n)] \right\rangle_\mathrm{c}; \qquad \boldsymbol{R}_{mn} \equiv \boldsymbol{R}_m - \boldsymbol{R}_n\,, \tag{8}$$

where $\delta G(m, n)$ is nonlocal in space and time now and represents an essential correction to the linking photon-induced elementary conductances $G(m, n)$ between any two effective centres m and n. These conductances are closely related to the statistical transition

[1]) g is an abbreviation for the coupling function (4) of the electron–photon interaction.

probabilities Γ_{mn} of the photon-induced electron hops from m to n,

$$G(m, n) = \frac{e^2\beta}{4}(\Gamma_{mn} + \Gamma_{nm}) ; \qquad \beta \equiv (k_B T)^{-1} . \tag{9}$$

Here Γ_{mn} is proportional to the second power of the coupling function (4) and is separately discussed in the following section. Arising from the perturbing field-induced redistribution of the charges the correction $\delta G(m, n)$ takes into account the contribution of the diffusive currents which are trying to restore the initial distribution of the localized electrons.

We note that in $\delta G(m, n)$ is also contained the light-field-induced change of the electron population. Because of the relative high frequencies of the incident light, however, this change should be smaller than the corresponding perturbing field-induced population effect.

3. Photon-Induced Statistical Transition Probability

As mentioned above, the mobility of the localized charges is caused by their interaction with photons and, of course, is closely related to the corresponding statistical transition probability Γ_{mn} of an electron transfer between the sites m and n. In order to calculate it let us suppose that the incident light beam can be characterized by a primary photon distribution $N^0_{q\lambda}$ which is nonzero only in a given direction of propagation q^0_f and does not depend on the polarization of the light quanta here,

$$N^0_{q\lambda} \approx 4\pi N^{\,0}_{\varepsilon_q}\delta(q^0 - q^0_f) ; \qquad \varepsilon_q \equiv \hbar\omega_q , \qquad q^0 \equiv \frac{q}{q} . \tag{10}$$

Then the summation over the directions of polarization λ may be performed and the usual probability of the photon-induced hopping process $W_{mn} = \Gamma_{mn} f^0_{\varepsilon_m}(1 - f^0_{\varepsilon_n})$ becomes, with (4),

$$W_{mn} = \frac{2e^2}{m^2c^3}\int \mathrm{d}\varepsilon_q \varepsilon_q [|q^0_f \times k_{mn}(-\varepsilon_q)|^2 \delta(\varepsilon_{mn} - \varepsilon_q) + |q^0_f \times k_{mn}(+\varepsilon_q)|^2 \delta(\varepsilon_{mn} + \varepsilon_q)] N^0_{\varepsilon_q} , \tag{11}$$

where

$$k_{mn}(\mp\varepsilon_q) = (2\pi)^3 \int \mathrm{d}k\, k\psi^*_m(k)\, \psi_n\left(k \mp \frac{\varepsilon_q q^0_f}{\hbar c}\right). \tag{12}$$

Since $N^0_{\varepsilon_q}$ is assumed to be large compared to unity the contribution of the spontaneous emission of photons to the probability of interest has been neglected.

The further progress now touches the general problem of choosing the one-electron wave functions $\psi_m(x)$. We postulate here that the electrons are localized at steep and narrow potential fluctuations in hydrogen-like wave functions which will have an "envelope" falling off as $\exp(-\alpha_m |x - R_m|)$, where $\alpha_m \approx \alpha(\varepsilon_m)$ is the reciprocal localization radius. Then, far enough from the m-th site, the product $k\psi^*_m(k)$ can be approximately written as $-i\nabla_{R_m}\psi^*_m(k)$, and the evaluation of (12) is reduced to that of the integral

$$x_{mn} = \int \mathrm{d}x\, x \exp(-\alpha_m |x - R_m| - \alpha_n |x - R_n|) \tag{13}$$

which may be done exactly. For simplicity, however, we restrict ourselves to the limiting cases of a) low-energy excitation, $k_B T \approx \hbar\omega \ll \Delta_{cv}$ and b) high-energy excitation of the hopping system, $k_B T \ll \hbar\omega \approx \Delta_{cv}/4$. Here $\Delta_{cv} = \varepsilon_c - \varepsilon_v$ is the width of the mobility gap and $\hbar\omega$ denotes the peak energy of the quasi-monochromatic spectral radiation intensity $I_{\varepsilon_q} \sim \varepsilon^3_q N^0_{\varepsilon_q}$ with a finite linewidth Δ_ω at half-maximum used in the experiment. The first case (a) corresponds obviously to a transport in the imme-

diate vicinity of the Fermi level whereas the electrons with respect to the second case (b) will hop over larger energetic distances. In contrast to the situation for phonon-assisted hopping, here the most probable hopping energies, of course, are imposed on the system by the external controlled intensity distribution I_{ε_q}.

Within the framework of this limiting cases one finds the photon-induced statistical transition probability finally to be

$$\Gamma_{mn} = a \frac{\hbar}{c_0^2 m^2} \alpha_0 n_\omega^3 f^0_{\varepsilon_m}(1 - f^0_{\varepsilon_n}) \left(\frac{\varepsilon_{mn}}{\Delta_{cv}}\right)^2 (\alpha_F R_{mn})^8 \exp(-2\alpha_F R_{mn}) \cos^2 \vartheta_{f,mn} \sin^2 \vartheta_{f,mn} I_{|\varepsilon_{mn}|} \tag{14a}$$

and

$$\Gamma_{mn} = b \frac{\hbar}{c_0^2 m^2} \alpha_0 n_\omega^3 f^0_{\varepsilon_m}(1 - f^0_{\varepsilon_n}) \left(\frac{\Delta_{cv}}{\varepsilon_{mn}}\right)^4 (\alpha_1 R_{mn})^2 \exp(-2\alpha_1 R_{mn}) \cos^2 \vartheta_{f,mn} \sin^2 \vartheta_{f,mn} I_{|\varepsilon_{mn}|}\,, \tag{14b}$$

refering to the low-energy and high-energy excitation, respectively. Here α_0 is the fine-structure constant, n_ω is the index of refraction assumed to be constant in the frequency ranges under consideration, α_F and α_1 are the reciprocal localization radii of the effective bound electron states at the Fermi level and the characteristic energy ε_1, respectively, from which the electrons are predominantly excited in the high-energy case, and $\vartheta_{f,mn}$ represents the angle between the direction of propagation $\boldsymbol{q}_f^0$ of the incoming light beam and the separation vector $\boldsymbol{R}_{mn}$. The dimensionless constants are denoted by a and b.

According to our restriction on the one-photon processes the transition probabilities (14) are proportional to the spectral intensity I_{ε_q}. Because the latter is non-vanishing only around the energy $\hbar\omega$, the so-called effective subnetwork of localized states associated with the charge transport is built up of all those pairs (m, n) of sites whose energy level spacing approximately equals this peak energy, $|\varepsilon_{mn}| \approx \hbar\omega$, i.e. which are in "resonance" with the radiation field.

4. Percolation Analysis and Results

Because of the spatially and temporally nonlocal correction $\delta G(m,n)$ to the individual site-to-site conductance $G(m, n)$, each elementary electron transfer is coupled self-consistently with all other hopping transitions. Thus, the static hopping transport is characterized as a macroscopically determined phenomenon. For this reason the single microscopic processes have to be summed by means of methods starting from an overall-view of the excited hopping system. The percolation theory is such a method. Simulating the excited hopping system by an adequate random-lattice conductance network, the percolation analysis allows to find the macroscopically most favourable current paths through the effective subnetwork. Their conductances then determine the macroscopic conductivity as follows [11]

$$\sigma_h \approx R^{-1} G_c^* \,. \tag{15}$$

Here G_c^* is the so-called critical percolation conductance defined as the largest value of all conductances G^* of the effective subnetwork such that the subset of conductances with $G^* \geqq G_c^*$ still contains a cluster which spans the entire hopping system. The inverse prefactor R of the preceding expression represents a characteristic length of such a critical percolation cluster. Since the conductance spread is sufficiently large here due to the exponential dependence of the conductances (14) on the site separation, the critical path prediction (15) should be usefully accurate [14]. It is immediately apparent that the replacement of (8) by (15) is only true for an isotropic and macroscopically homogeneous hopping system. The conductivity tensor is then diagonal

Its longitudinal component $\sigma_h^{\parallel}$ with respect to the incident light beam and the two identical transverse components $\sigma_h^{\perp} \equiv \sigma_h$ are of the same order of magnitude.

In order to simplify the calculation of the critical percolation conductance G_c^*, we replace the individual intersite conductances of the effective network by their averages over all orientations relative to the direction of light propagation as well as the energetic positions of the corresponding pairs. Hence the conductances of the effective random-lattice network under consideration will merely depend on the spatial site separation, $G^* = G(R_{mn}; |\varepsilon_{mn}| \approx \hbar\omega)$, and the percolation analysis is reduced to its simplest form, namely, Pollak's so-called "R-percolation" [15]. It determines the desired conductance G_c^* in terms of a critical percolation distance R_c^*, $G_c^* = G(R_c^*; \hbar\omega)$. This situation is analogous to the case encountered not only in the impurity conduction at low compensation but also in the phonon induced hopping transport with respect to the high-temperature regime because the variable range aspect then disappears.

To obtain quantitatively computable expressions for the photon-induced dc hopping conductivities, the parameter R was approximated by the mean separation distance of the critical percolation subnetwork, where the averaging was performed with the help of a two-centre correlation function including zero probability of finding two localized states within a short distance R_0 of each other. Using the value $R_c^*/a_s^* = 0.7048$ for the critical percolation distance which is normalized by the average lattice constant a_s^* of the effective subnetwork and was obtained by Pike and Seager [16] with the aid of Monte Carlo techniques, the conductivities of interest can be finally written in the following way:

$$\sigma_h \approx \alpha_0^2 n_s^*(\hbar\omega)\, A(\hbar\omega) \left(\frac{\hbar\omega}{\Delta_{cv}}\right)^2 I_0 \exp\left[-2.819 \left(\frac{3\alpha_F^3}{8\pi N_F} \frac{1}{\hbar\omega}\right)^{1/3}\right] \tag{16a}$$

and

$$\sigma_h \approx \alpha_0^2 n_s^*(\hbar\omega)\, B(\hbar\omega) \left(\frac{\Delta_{cv}}{\hbar\omega}\right)^4 I_0 \exp\left[-2.819 \left(\frac{3\alpha_1^3}{8\pi \Delta_0 N_F} \sinh^{-1} \frac{\hbar\omega}{\Delta_0}\right)^{1/3}\right], \tag{16b}$$

relating to the limiting cases of low-energy excitation and high-energy excitation of the hopping system, respectively. Here $n_s^*(\hbar\omega)$ is the concentration of the effective sites, N_F and Δ_0 represent, respectively, the electronic density of states at the Fermi level and a measure of its tailing into the mobility gap, while I_0 denotes the total intensity of the incident light beam. The functions $A(\hbar\omega)$ and $B(\hbar\omega)$ depending on all other model parameters are too lengthy to be quoted here.

It should be noted that at low temperatures the conductivity (16b) corresponding to the limiting case of high-energy excitation is only due to absorption processes, in contrast to that of the converse limiting case (16a).

5. Discussion and Conclusions

The dc hopping photoconductivities (16) are obviously proportional to the second power of the fine-structure constant. As expected, they increase linearly with the concentration of the effective sites and the total intensity of the incident light beam. In the low-energy limit their spectral dependence is characterized by a dominant exponential $\omega^{-1/3}$-frequency dependence. Therefore the conductivity increases rapidly with the light frequency. For very low photon energies it behaves asymptotically like

$$\sigma_h \approx \text{const} \frac{\hbar^3}{c_0 m^2} \alpha_0^2 n_\omega^3 \frac{N_F^2}{\alpha_F^5} \left(\frac{\hbar\omega}{\Delta_{cv}}\right)^2 I_0 \sim \omega^2, \tag{17}$$

if all statistical correlations of the excited hopping system are neglected. Formula (17) reminds of Mott's law of the ac hopping conductivity due to direct optical absorption

processes mentioned in Section 2. That is, the spectral dependence of the photon-induced hopping conductivities should be determined foremost by the spectral dependence of the electronic mobility which is, of course, caused by the same physical elementary process here and with respect to Mott's ac hopping conduction.

In contrast to the conductivities just discussed, in the high-energy limit the spectral dependence of the dc hopping photoconductivity is markedly affected by the fine structure of the electronic density of states in a wide vicinity of the Fermi level. Thus the dominant exponential $\sinh^{-1/3}\omega$-frequency dependence presented above was only obtained by referring to the Mott CFO basic band model of disordered semiconductors.

To give a rough idea of the practical significance of our results, we have estimated the order of magnitude of the conductivities (16) in their dependence on the frequency of the incident light. Furthermore, comparing the high-energy conductivity (16b) with the phonon-induced conductivity for Mott's variable range hopping at the reference temperature $T = 81$ K, we have also evaluated the critical light intensity I_{oc} for which the photon-induced hopping mechanism begins to dominate the phonon-induced one. The results obtained are illustrated in Fig. 1 and 2. They can be summarized as follows: (i) Based on realistic variants of the Mott CFO band model, the dc hopping photoconductivities reach noticeable orders of magnitude of 10^{-15} to $10^{2}\ \Omega^{-1}\,\mathrm{cm}^{-1}$ in the spectral range under consideration when the specimens are illuminated by laser light of a relatively low intensity. (ii) Because of the larger concentration of the effective hopping sites, the high-energy conductivity was found to be always greater than the low-energy conductivity and is highly sensitive to the detailed shape of the density of states within the mobility gap. Indeed, for any more realistic and non-monotonic shape of this density of states, the conductivity is, of course, a more drastically varying function of frequency than it is demonstrated in Fig. 1. (iii) The frequency dependence of the non-exponential prefactors of all conductivities is much weaker than that of the corresponding exponential factors. It increases somewhat with increasing energy of the light quanta. (iv) The critical light intensity decreases monotonically as the frequency increases. With respect to the reference temperature chosen this intensity lies in a region of 10^{3} to 10^{-1} W cm^{-2} for higher photon energies. Of course, a decrease of the reference temperature lowers the critical intensity.

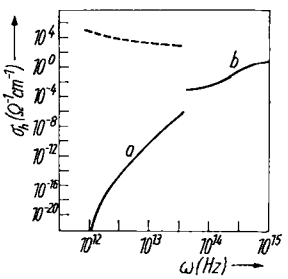

Fig. 1

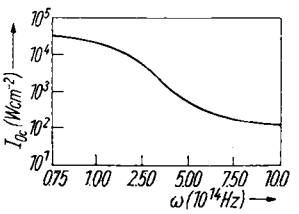

Fig. 2

Fig. 1. Typical spectral dependence of the dc hopping photoconductivity in the limiting cases of low-energy excitation (a) and high-energy excitation (b), respectively. The dashed curve represents the frequency dependence of the pre-exponential factor. $\Delta_{cv} = 1.6$ eV, $N_F = 10^{19}\ \mathrm{eV}^{-1}\,\mathrm{cm}^{-3}$, $\Delta_0 = 80$ meV, $\alpha_F = 10^{7}\ \mathrm{cm}^{-1}$, $n_\omega = 4$, $I_0 = 10^{4}\ \mathrm{W\,cm}^{-2}$

Fig. 2. The critical light intensity versus frequency at 81 K. Mott's phonon-induced hopping conductivity has been calculated with an acoustic-deformation-potential constant $C_1 = 10$ eV. The total intensity $I_0^{(\mathrm{phon})} \sim T^4$ of the assembly of the thermally activated phonons is of the order of $10^{6}\ \mathrm{W\,cm}^{-2}$

We further notice that the spectral photoresponse of the hopping system is obviously modified when both a more complicated asymptotic behaviour of the localized wave functions is used and such correlations are taken into account which go beyond the Pauli exclusion principle and the repulsion of sites of like energy, respectively, involved here. Whereas the former is only reflected in a slight modification of the weak frequency dependence of the pre-exponential factors, the latter may play a more important role in forming the spectral photoresponse. Concluding one may say that the photon-induced hopping mechanism discussed here deserves interest also, for at least two reasons, from an experimental point of view. First, under realistic conditions the hopping photoconduction can be the dominant transport mechanism leading to observable effects in such spectral ranges where the lattice absorption is weak. Secondly, it should permit a frequency-controlled "reading" of the electronic density of states within the mobility gap. Indeed, neglecting the frequency dependence of the pre-exponential factor of the high-energy conductivity the density of states of interest may be obtained by means of the relation

$$N(\varepsilon_F + \hbar\omega) \approx -3.8\alpha_F^3 \frac{d}{d\hbar\omega} \ln^{-3}[10^{-5}\sigma_h(\hbar\omega)], \tag{18}$$

where σ_h has to be expressed in units of Ω^{-1} cm^{-1}.

In our opinion technical applications of the hopping photoconduction could be expected in the field of infrared detectors, in particular, for the study of intense laser photon sources. The time constant of such detectors should be of the order of the characteristic relaxation time of the assembly of localized electrons.

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References

[1] S. TANAKA and H. Y. FAN, Phys. Rev. **132**, 1516 (1963).
[2] J. BLINOWSKI and J. MYCIELSKI, Phys. Rev. **136**, A266 (1964); **140**, A1024 (1965).
[3] N. F. MOTT, Phil. Mag. **19**, 835 (1969); **22**, 7 (1970).
[4] I. Z. KOSTADINOV, Zh. eksper. teor. Fiz., Pisma **14**, 345 (1971); Proc. VI. Internat. Conf. Amorph. Liquid Semicond., Leningrad 1975, Vol. Electronic Phenomena in Non-Crystall. Semicond., Nauka, Leningrad 1976 (p. 156).
[5] V. L. BONCH-BRUEVICH and V. ČÁPEK, Zh. eksper. teor. Fiz., Pisma **16**, 109 (1972).
[6] R. SCHUCHARDT and R. KEIPER, Proc. Conf. Amorph. Semicond., Reinhardsbrunn 1974 (p. 155); Balatonfüred 1976 (p. 89).
[7] H. BÖTTGER and V. V. BRYKSIN, Fiz. tverd. Tela **18**, 1888 (1976).
[8] H. BÖTTGER and V. V. BRYKSIN, phys. stat. sol. (b) 78, 415 (1976).
[9] I. P. ZVYAGIN and R. KEIPER, Wiss. Z. Humboldt-Univ. Berlin, Math.-Nat. R. **21**, 459 (1972).
[10] E. O. MANUCHARYANTS and I, P. ZVYAGIN, phys. stat. sol. (b) **65**, 665 (1974).
[11] V. AMBEGAOKAR, B. I. HALPERIN, and J. S. LANGER, Phys. Rev. B **4**, 2612 (1971).
[12] D. N. ZUBAREV, Fortschr. Phys. **9**, 275 (1961).
[13] R. SCHUCHARDT Dissertation, Humboldt-Univ. Berlin 1976.
[14] C. H. SEAGER and G. E. PIKE, Phys. Rev. B **10**, 1435 (1974).
[15] M. POLLAK, J. non-crystall. Solids **11**, 1 (1972).
[16] G. E. PIKE and C. H. SEAGER, Phys. Rev. B **10**, 1421 (1974).

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