

Size localization of electrons in inhomogeneous semiconductors

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A localization mechanism connected with the dependence of the percolation level on the thickness of the conducting region is considered. The corresponding dependence of the conduction activation energy makes it possible to estimate the critical exponent that determines the behavior of the correlation length near the percolation threshold.

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A transition from the activationless temperature dependence of the conduction to an activation dependence has been recently observed and corresponds to localization of the electrons as the thickness of the conducting region is decreased.^[1,2] In the cited papers, this effect is attributed to the decrease of the “minimal metallic conductivity” on going to a two-dimensional disordered system, in which the decisive role is played by small-scale fluctuations of the internal field.

We consider size-dependent localization of electrons due to the presence of a continuous random field in the sample. The model of an inhomogeneous semiconductor with a continuous random field was used earlier with success to explain certain singularities of recombination and photoconductivity of disordered semiconductors, to describe conduction over an uneven surface and for other purposes.^[3–5] We shall neglect hereafter, as is customary, the possibility of electron tunneling through the classically inaccessible regions; this assumption is justified for smooth large-scale fluctuations of the potential. Accordingly, in an inhomogeneous system with large-scale fluctuations the percolation level is close to the classical value $E_p^{(d)}$, whose position depends on the form of the characteristic functional $\mathcal{P}\{U(\mathbf{x})\}$ that determines the statistical properties of the random field $U(\mathbf{x})$, and on the dimensionality d of the space. In the two-dimensional case, for an arbitrary symmetrical functional $\mathcal{P}\{U\} = \mathcal{P}\{-U\}$ the critical fraction, corresponding to the energy $E_p^{(2)}$, of the classically accessible space is equal to $\nu_c^{(2)} = 0.5$ ^[6,7]; in the three-dimensional case calculation using a Gaussian functional of the form $\mathcal{P}\{U\} = \exp[-(1/2\psi_1) \int U^2(\mathbf{x}) d\mathbf{x}]$ yielded value $\nu_c^{(3)} \approx 0.17$.^[8] If we assume the characteristic functional to be Gaussian, then the position of the percolation level is given by

$$E_p^{(d)} = -\xi^{(d)} \psi_1^{1/2}, \quad (1)$$

and is reckoned from the average potential energy of the electron, while the parameter $\xi^{(d)}$ ($d=2,3$) is determined from the equation

$$\text{erf} \frac{\xi^{(d)}}{\sqrt{2}} = 1 - 2v_c^{(d)}. \quad (2)$$

ie values of $v_c^{(d)}$ indicated above correspond to $\xi^{(2)}=0$ and $\xi^{(3)}=0.96$. The inequality $E_p^{(2)} > E_p^{(3)}$ reflects the obvious circumstance that when the dimensionality of space increases additional possible ways of going around the "humps" of the potential relief are produced, and it is this which leads to a lowering of the percolation level. One can expect this lowering to take place also for non-Gaussian potentials and when tunneling is taken into account, although the actual expressions for the percolation level will be different.

We assume for the sake of argument that the characteristic functional is Gaussian and denote l_0 the characteristic scale of the potential. If we decrease the thickness w of the conducting region without changing the concentration and shape of the potential relief, then a transition to the one-dimensional case takes place at $w < l_0$. However, just as in the case of hopping conduction in thin films,^[9] the fact that the transverse dimensions of the system are finite comes into play already at $w > l_0$. The point is that, in accordance with the usual assumption concerning the behavior of the correlation length L_c that characterizes the shape of the classically accessible region, this length has a power-law dependence near the percolation threshold^[10]:

$$L_c = l_0 \left(\frac{|v_c - v_c^{(3)}|}{v_c^{(3)}} \right)^{-\nu}, \quad (3)$$

where ν is a critical exponent that is somewhat smaller than unity for lattice problems. There is no percolation at small thicknesses; it appears only at $L_c \lesssim w$. This means that the critical fraction of the volume, v_c , begins to depend on w —the condition $L_c \sim w$ with allowance for (3) yields

$$v_c(w) = v_c^{(3)} \left[1 + B \left(\frac{l_0}{w} \right)^{1/\nu} \right], \quad (4)$$

where B is some unknown numerical coefficient. The relation (4) is valid at $w > l_0$, when the difference between $v_c(w)$ and $v_c^{(3)}$ is small; at thickness smaller than l_0 , the value of $v_c(w)$ tends to a constant value $v_c^{(2)}=0.5$.

According to (1), (2), and (4), at a given thickness w , the critical fraction $v_c(w)$ of the volume corresponds to a percolation level

$$E_p(w) \cong E_p^{(3)} + B_1 \psi_1^{1/2} \left(\frac{l_0}{w} \right)^{1/\nu}, \quad (5)$$

where $B_1 \approx 4Bv_c^{(3)}$. If the Fermi level F in a bulky sample is located higher than $E_p^{(3)}$ but lower than $E_p^{(2)}$, then a transition from metallic conductivity to activated semiconductor conductivity takes place when w becomes smaller than a value w_0 such that $E_p(w_0)=F$. If the difference $F - E_p^{(3)} \equiv \epsilon_0$ is not very large, so that at $E_p(w) \sim F$ one can still use formula (5), then at $w < w_0$ the activation energy of the static conductivity varies like

$$\epsilon(w) = \epsilon_0 + B_1 \psi_1^{1/2} \left(\frac{l_0}{w} \right)^{1/\nu}. \quad (6)$$

When w is decreased, deviations from relation (6) should be observed—a slowing down of the increase of the activation energy and a tendency of this energy to assume the limiting value $\epsilon(0) = E_p^{(2)} - F$ at $w < l_0$.

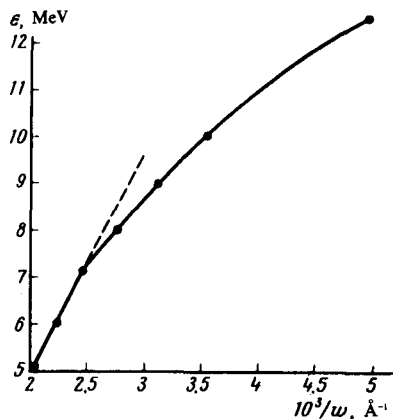


FIG. 1.

It is of interest to compare the derived relations with the experimental data on the variation of the activation energy with the thickness of the conducting region. Data of this kind, pertaining to activation in the metallized impurity band of *n*-GaAs, is contained in [2]. It should be noted that such a comparison is to some degree arbitrary, since it is not quite clear [11] to what extent the model used above suffices for the description of the system investigated in [2]. The function $\epsilon(\omega)$ calculated from the data of [2] is shown in Fig. 1. It is seen that the character of this dependence corresponds to that discussed above; at $\omega \gtrsim 400 \text{ \AA}$ the experimental points are satisfactorily described by formula (6) with a critical exponent close to unity. This value agrees with the corresponding calculations for lattice problems [10,12]. The observed slowing down of the growth of $\epsilon(\omega)$ at smaller ω also agrees with the model discussed above.

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