Coulomb term. This is close to the value 0.2 eV cited in the literature for the activation energy of indium donors. With rising temperature the maximum of this band shifts to higher energies (from 1.93 eV at 4.2 K to 1.98 eV at 300 K) and its half-width increases. The shift of the band to higher and not lower energies is apparently due to the fact that as the temperature increases the donor–acceptor pairs break up as a result of thermal dissociation and an increasing role is played by the emission resulting from the transition of electrons from the donor level to the conduction band. We point out, however, that with increasing temperature the main role in the emission is played by the band with maximum intensity at 2.374 eV at 77 K and 2.247 eV at room temperature. The maximum of this band coincides with the position of the maximum of the principal emission at room temperature for undoped crystals and, therefore, we assume that this interpretation is correct, i.e., that the emission is due mainly to electron transitions to shallow acceptor levels, most probably Li_{Zn} levels. Thus, in the case of doping with indium we observe two emission bands at room temperature, both of which are due to impurities, but the impurity band with maximum at ∿2.25 eV dominates.

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TRANSPORT OF NONEQUILIBRIUM ELECTRONS IN PHOTOEMISSION FROM SEMIBOUNDED MEDIUM

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The equation is formulated and accurately solved for the problem of the transport of photoexcited electrons which undergo scattering before reaching vacuum, with the formation of an electron–hole pair, as well as multiple scattering of elastic type. The analytical expressions obtained permit calculations of the spectral composition of the emitted-electron flux and the quantum yield of photoemission. The limiting expression for the probability that secondary electrons will reach the vacuum which is obtained when the elastic-scattering mechanism is switched off is compared with literature data.

The spectral composition of the emitted electron flux is determined by various factors, among which electron scattering processes with the formation of electron–hole pairs (e–e processes) and with the emission or absorption of a phonon (e–p scattering) are important. The transport of photoexcited electrons, taking account of such processes, has been considered in a whole series of works, both within the framework of the single-dimensional model [1-4] and in a three-dimensional propagation pattern [5-7]. Using various model assumptions, analytical expressions for the probability that the electrons will emerge into the vacuum are obtained describing electron transport in a semiinfinite medium [1, 4-7] and an emissionally active layer of finite thickness [2, 3]. The results obtained in [1-7] allow the contribution of primary electrons to the photoemission current to be calculated. If the electrons are excited by light to a state of energy exceeding the position of the vacuum level by an amount larger than the forbidden band width, however, the spectral composition of the emitted-electron flux is formed by both primary and secondary electrons.

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In [8-13], in calculating the photoemission characteristics, effects associated with the formation of secondary excitations were taken into account, but e–p scattering processes were ignored.

In events of e–p interaction, there is randomization of the directions of electron propagation (with a small change in electron energy). Therefore, for electrons able to enter repeatedly into such interactions, the pattern of transport is fundamentally changed. Since the subsequent transport of electrons undergoing e–e scattering is no different from the transport of primary electrons with the same energy, e–p scattering processes must be taken into account.

In the one-dimensional model, secondary-electron transport taking account of e–p interaction was considered for a semiinfinite medium in [14] and for a layer of material of finite thickness in [15]. In the present work, it is shown that this problem has an analytical solution for a semiinfinite medium, even with a three-dimensional pattern of electron propagation.

In describing secondary-electron transport, the approach proposed in [1] for primary electrons is followed: the propagation of a single nonequilibrium electron is traced from the event of photoexcitation to emergence into vacuum, neglecting the change in scattering parameters of the electron in e–p scattering events.

The first step is to introduce the probability $q_n(x; E, E')$ of the process in which the electron excited to a state with energy E at a depth x from the emitting surface undergoes n scattering episodes with the formation of electron–hole pairs and emerges into vacuum with energy E'. The sum of individual probabilities q_n gives the total probability that the electron will reach the vacuum: $q = \sum_{n=0} q_n$. Note that $q_0(x; E, E') = \delta(E-E')g_0(x; E)$, where $\delta(E)$ is a Dirac delta function and g_0 describes primary-electron transport [7].

Electrons undergoing two or more e–e scattering episodes are neglected, and the equation for the probability $q_1(x; E, E')$ is formulated.

Suppose that K(x, y) is the probability of electron transition from a depth x to a depth y without e–p and e–e scattering, taking account of its reflection from the boundary. At a depth y, in a small interval dy, both e–p scattering (with no change in electron energy) and e–e scattering, with a probability S(E, E') of electron transition from a state with energy E to a state with energy E' are possible. Correspondingly, the equation for q_1 takes the form

$$q_1(x; E, E') = \int_0^\infty dy K(x, y)\left[\frac{1}{l_p} q_1(y; E, E') + \frac{1}{l_e} 2S(E, E') g_0(y; E')\right], \quad (1)$$

where l_p and l_e are the mean free path lengths of electrons for e–p and e–e scattering, respectively. The factor 2 preceding the probability S appears because the photoexcited electron and the electron excited in e–e scattering make different contributions to the photoemission current.

Suppose that electron scattering by the surface is diffusional in character. Then, according to [7], the core of the integral operator in Eq. (1) takes the form

$$K(x, y) = \frac{1}{2} E_1\left(\frac{|x-y|}{l}\right) + \frac{R}{2} E_2\left(\frac{x}{l}\right) E_1\left(\frac{y}{l}\right), \quad (2)$$

where $E_n(x) = \int_1^\infty du u^{-n} e^{-xu}$ are integral exponentials; l is the effective free path length of the electrons; $l^{-1} = l_p^{-1} + l_e^{-1}$.

For convenience of exposition, the arguments E and E' are omitted, where possible, understanding that $q_1(x) = q_1(x; E, E')$, $S = S(E, E')$, $R = R(E)$, $R' = R(E')$, $l = l(E)$, etc.

Next Eq. (1) is transformed to give

$$q_1(x) = \frac{1}{2l_p} \int_0^\infty dy E_1\left(\frac{|x-y|}{l}\right)\left[q_1(y) + 2\frac{l_p}{l_e} S g_0(y; E')\right] + \frac{R}{1+R} q_1(0) E_2\left(\frac{x}{e}\right). \quad (3)$$

The only solution of this equation which is of physical interest is that satisfying the boundary condition $q_1(x) \to 0$ as $x \to \infty$.

Equation (3) was obtained in the region $x \geqslant 0$. It is solved by the Wiener–Hopf method. To this end, it is regarded as an equation on the whole x axis, predetermining the last term on the right-hand side of Eq. (3) hypothetically as zero when $x < 0$.

Complex Fourier transformation of Eq. (3) gives

$$\left(1 - \frac{\operatorname{arctg} \kappa l}{\kappa l_p}\right) \hat{q}_1^+(\kappa) + \hat{q}_1^-(\kappa) = 2S \frac{\operatorname{arctg} \kappa l}{\kappa l_e} \hat{g}_0^+(\kappa;\ E') + \frac{R}{1+R} q_1(0) \left[\frac{i}{\kappa} + \frac{1}{\kappa^2 l} \ln(1 - il\kappa)\right], \tag{4}$$

where $\hat{q}_1^{\pm}(\kappa)$ are Fourier transformations of the function $q_1(x)\theta(\pm x)$

$$\hat{q}_1^+(\kappa) = \int_0^\infty dx e^{i\kappa x} q_1(x), \quad \hat{q}_1^-(\kappa) = \int_{-\infty}^0 dx e^{i\kappa x} q_1(x).$$

Following the Wiener–Hopf method, Eq. (4) is transformed so that it contains only two functions, one of which is analytical in the upper halfplane and the other in the lower halfplane, with some common plane of overlap of their regions of analyticity.

It follows from Eq. (3), with the boundary condition that $q_1(x) \to 0$ as $x \to \infty$, that the function $\hat{q}_1^+(\kappa)$ is analytical in the region Im $\kappa \geqslant 0$ and $\hat{q}_1^-(\kappa)$ is analytical in the region Im $\kappa < l^{-1}$.

The factor preceding $\hat{q}_1^+(\kappa)$ in Eq. (4) is written in the form

$$1 - \frac{\operatorname{arctg} \kappa l}{\kappa l_p} = \frac{Y^+(\kappa)}{Y^-(\kappa)}, \tag{5}$$

where $Y^+(\kappa) = \dfrac{\kappa^2 + 1/L^2}{\kappa + i/l} e^{F^+(\kappa)}$; $Y^-(\kappa) = (\kappa - i/l) e^{F^-(\kappa)}$; L is determined from the equation $l = L \tanh(l_p/L)$. The definition and some properties of the function $F^{\pm}(\kappa)$ may be found in [7]. Here it is simply noted that the functions $F^{\pm}(\kappa)$ are analytical in the whole complex plane with cuts along the imaginary axis from $(-i/l)$ to $(-i\infty)$ and from i/l to $i\infty$, respectively; in their regions of analyticity, $F^{\pm}(\kappa) \to 0$ as $|\kappa| \to \infty$.

Using Eq. (5), Eq. (4) is reduced to the form

$$\kappa Y^+(\kappa) \hat{q}_1^+(\kappa) - i \frac{l_p}{l} \frac{2R}{1+R} q_1(0)\, Y^+(\kappa) = Y^-(\kappa) \left\{ -\kappa \hat{q}_1^-(\kappa) + \right. \tag{6}$$

$$\left. + i \frac{R}{1+R} q_1(0) \left[1 - 2\frac{l_p}{l} + \frac{1}{i\kappa l} \ln(1 + i\kappa l)\right] \right\} + 2S \frac{l}{l_e} Y^-(\kappa) \frac{\operatorname{arctg} \kappa l}{l} \hat{g}_0^+(\kappa;\ E').$$

Consider the function

$$\varphi(\kappa) = \frac{i}{\kappa - i/L} \frac{\operatorname{arctg} \kappa l}{l} Y^-(\kappa) \hat{g}_0^+(\kappa;\ E'). \tag{7}$$

It is analytical inside the strip $0 \leqslant \text{Im } \kappa < 1/L$, and decreases at its ends (as Re $\kappa \to \pm \infty$). Therefore, $\varphi(\kappa)$ may be written in the form

$$\varphi(\kappa) = \varphi^+(\kappa) - \varphi^-(\kappa), \tag{8}$$

where the functions $\varphi^{\pm}(\kappa)$ are defined by the expression

$$\varphi^{\pm}(\kappa) = \frac{1}{2\pi i} \int_{C^{\pm}} \frac{dt}{t - \kappa} \varphi(t) \tag{9}$$

and are analytical in the halfplane Im $\kappa \geqslant 0$ and Im $\kappa < 1/L$ respectively. The integration contours $C^{\pm}$ run from $|t| = \infty$, Re $t < 0$ to $|t| = \infty$, Re $t > 0$ in the region of analyticity of function $\varphi(t)$; C^+ passes below the point $t = \kappa$, and C^- above this point.

Substituting Eqs. (7) and (8) into Eq. (6) yields an equation in which the function analytical in the region Im $\kappa \geqslant 0$ is equal to the function analytical in the halfplane Im $\kappa < 1/L$, and thus its analytical continuation on the whole complex plane. The only function analytical in the whole complex plane is a polynomial.

Taking into account that both sides of Eq. (6) increase no faster than κ as $|\kappa| \to \infty$, it may be concluded that this is a first-order polynomial

$$\kappa Y^+(\kappa)\hat{q}_1^+(\kappa) - i\frac{l_p}{l}\frac{2R}{1+R}q_1(0)\,Y^+(\kappa) + 2Si\frac{l}{l_e}(\kappa - i/L)\varphi^+(\kappa) = A_0 + A_1\kappa. \tag{10}$$

It follows from this equality, when $\kappa = 0$, that

$$A_0 = -i\frac{l_p}{l}\frac{2R}{1+R}q_1(0)\,Y^+(0) + 2S\frac{l}{l_e}\frac{1}{L}\varphi^+(0). \tag{11}$$

Considering the behavior of both sides of Eq. (10) at infinity ($\kappa = iS$, $S \to \infty$), A_1 may be determined

$$A_1 = iq_1(0)\left[1 - \frac{l_p}{l}\frac{2R}{1+R}\right]. \tag{12}$$

Now let $\kappa = i/L$. Taking into account that $Y^+(0) = -(i/L)\sqrt{1 - l/l_p}$, it is found that

$$q_1(0) = 2S\frac{l}{l_e}(1+R)\frac{\varphi^+(0)}{1 - \frac{l_p}{l}R(1-\sqrt{1-l/l_p})^2}. \tag{13}$$

Expressing $\hat{q}_1^+(\kappa)$ from Eq. (10) in the form

$$\hat{q}_1^+(\kappa) = \frac{i}{\kappa}\frac{l_p}{l}\frac{2R}{1+R}q_1(0)\left[1 - \sqrt{1-l/l_p}\frac{\kappa + i/l}{\kappa + i/L}e^{-F^+(\kappa)}\right] + i2S\frac{l}{l_e}(\varphi^+(0) - \varphi^+(\kappa))\frac{\kappa + i/l}{\kappa(\kappa + i/L)}e^{-F^+(\kappa)}. \tag{14}$$

and taking account of the result obtained in [7]

$$\hat{g}_0^+(\kappa, E) = \frac{i}{\kappa}\frac{l_p}{l}\frac{1-R}{1 - \frac{l_p}{l}R(1-\sqrt{1-l/l_p})^2}\left[1 - \sqrt{1-l/l_p}\frac{\kappa + i/l}{\kappa + i/L}e^{-F^+(\kappa)}\right]. \tag{15}$$

Eq. (14) is written in the more compact form

$$\hat{q}_1^+(\kappa) = 2S\frac{l}{l_e}\left[\frac{2R}{1-R}\varphi^+(0)\hat{g}_0^+(\kappa; E) + \frac{i}{\kappa}\frac{\kappa + i/l}{\kappa + i/L}(\varphi^+(0) - \varphi^+(\kappa))e^{-F^+(\kappa)}\right]. \tag{16}$$

Now suppose that the exciting light damps exponentially in the material. In this case, there is a simple relation between $\hat{q}_1^+(\kappa)$ and the mean probability $Q_1(\alpha) = \alpha\int_0^\infty dx e^{-\alpha x}q_1(x)$ over the absorption depth of the light: $Q_1(\alpha) = \alpha\hat{q}_1^+(i\alpha)$. Thus, the expression for the probability Q_1 takes the form

$$Q_1(\alpha; E, E') = 2S(E, E')\frac{l}{l_e}\left[\frac{2R}{1-R}\varphi^+(0)G_0(\alpha; E) + \frac{\alpha + 1/l}{\alpha + 1/L}(\varphi^+(0) - \varphi^+(i\alpha))e^{-F^+(i\alpha)}\right], \tag{17}$$

where $G_0(\alpha; E) = \alpha\hat{g}_0^+(i\alpha)$ is the probability that a primary electron emerges into the vacuum. The total emergence probability $Q(\alpha; E, E') = \delta(E-E')G_0(\alpha; E) + Q_1(\alpha; E, E')$ allows the relation between the initial nonequilibrium distribution $N_0(E, \omega)$ of the excited electrons and the spectral composition of the emitted-electron flux to be established

$$N(E', \omega) = \int dE N_0(E, \omega)\,Q(\alpha; E, E') = N_0(E', \omega)G_0(\alpha; E') + \int dE N_0(E, \omega)\,Q_1(\alpha; E, E'). \tag{18}$$

The quantity $G(\alpha; E) = \int dE' Q(\alpha; E, E')$ is understood to be the mean probability of emergence into vacuum of an electron excited by light to the state with energy E. It may be used to calculate the spectral dependence of the quantum yield

$$Y(\omega) = \int dE N_0(E, \omega)G(\alpha; E). \tag{19}$$

To switch off the electron–phonon interaction, Eq. (17) is considered in the limit as $l_p \to \infty$. If the integration contour C^+ in Eq. (9) is deformed so that it passes over both sides of the cut along the imaginary axis from $(-i\infty)$ to $(-i/l')$, Eq. (17) may be reduced to the form

$$Q_1(\alpha; E, E') = 2S(E, E')\frac{1-R'}{4}\left\{R\left[1 - \frac{1}{\alpha l}\ln(1+\alpha l)\right]\times\right.$$

$$\left.\times\int_0^1 dt\, t\frac{l'}{l}\ln\left(1 + \frac{l}{tl'}\right) + \int_0^1 \frac{dt}{1 + 1/(\alpha l' t)}\left[t\frac{l'}{l}\ln\left(1 + \frac{l}{tl'}\right) + \frac{1}{\alpha l}\ln(1+\alpha l)\right]\right\}. \tag{20}$$

Equation (20) may be compared with the expression obtained in [8, 9]

$$Q_1 = 2S(E, E')\frac{1}{4}\int_{p_c/p}^{1}\frac{dt}{1+1/(\alpha l' t)}\left[t\frac{l'}{l}\ln\left(1+\frac{l}{l't}\right)+\frac{1}{\alpha l}\ln(1+\alpha l)\right]. \quad (21)$$

What is the difference between Eqs. (20) and (21)? First, penetration of the electron through the boundary was considered in [8, 9] on the basis of the emergence-cone model. In other words, it was assumed that the probability of electron transition through the surface is unity when $\theta < \theta_c$ and zero when $\theta > \theta_c$ where θ is the angle of incidence of the electron at the surface with respect to its normal and θ_c is the critical angle; $\cos\theta_c = p_c/p$. Conversion from the emergence-cone model to a diffusely scattered surface entails setting the parameter p_c/p (which characterizes the emergence cone) equal to zero, and introducing the probability $(1 - R')$ that the electron passes through the surface as a factor.

Second, the contribution to the photoemission current from electrons reflected by the surface was disregarded in [8, 9]. These electrons are described by the first term in curly brackets in Eq. (20).

Numerical analysis shows that, with a favorable relation of the scattering parameters, the electrons reflected by the surface may make a pronounced contribution to the emission current.

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