Surface photogalvanic effect in metals

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The possibility of appearance of a current in a metallic film irradiated by interband-absorbed light is demonstrated. The cases of "dirty" films at temperatures above or equal to the Debye temperature or of pure films at low temperature are considered. In the latter case the current is determined by the carriers that managed to relax to the Fermi surface but retained the direction of their momentum. The photocurrent may be as high as 0.1 A/cm^2 per W/cm^2 of the incident light intensity.

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It was demonstrated in recent studies[1-3] that a new type of photovoltaic effect (PVE) can arise in films, namely flow of current through a homogeneous film under the action of stationary and uniform illumination. This effect is due to the even anisotropy of the probability of carrier generation under the action of light, and is therefore not connected with the absence of an inversion center in the film material itself. The necessary condition for the onset of this effect is the absence of inversion in the plane of the film in the film-plus-field system. This can be attained when the polarization vector is oblique relative to the normal to the film and when the film surfaces are different,[1] or when the carriers are not uniformly generated across the film.[2]

The phenomenological expression for the current in an isotropic material is[3]

$$\mathbf{J}=\beta_1 \operatorname{Re}\{[\mathbf{e}-\mathbf{N}(\mathbf{eN})](\mathbf{Ne}^*)\}+i\beta_2[\mathbf{N}[\mathbf{e}\times\mathbf{e}^*]]. \qquad (1)$$

Here $\mathbf{N}$ is the vector normal to the plane of the film, $\mathbf{e}$ is the light-polarization vector, and β_1 and β_2 are real constants.

The surface PVE can be explained physically using as a simple example a deformed (term with β_1) or twisted (term with β_2) tennis ball that collides inelastically with a surface. Such a ball rebounds after the collision at an angle not equal to the incidence angle, and this leads to the onset of a "flux" of balls along the surface, even if all the incidence directions are equally probable. In the former case (linear PVE[4]) the anisotropy of the rebound takes place if the ball is deformed at an angle to the surface (the deformation plays the role of the polarization tensor). In the latter case (circular PVE[4]) the flux arises if the vector of the angular velocity of rotation has a component in the plane of the surface (in this case it is the analog of the vector $\mathbf{e}\times\mathbf{e}^*$).

The surface version of the PVE is apparently the principal effect in the region of opacity of the material, inasmuch as in the case of strong absorption the surface is a substantial channel for the relaxation of the photoexcited carriers. As is well known, the photocurrent is proportional to the carrier momentum mean free path and is therefore expected to be the largest in materials with large mean free path of the nonequilibrium carriers. From among all the known crystals, the longest mean free paths (up to 1 cm) are observed in pure metals at low temperatures. Metals seem therefore to be promising materials for the observation of surface PVE.

In this paper we studied the PVE in films of typical metals for interband transitions under the situation of the normal skin effect.

RELAXATION-TIME APPROXIMATION

In the relaxation-time approximation, the expression for the surface density of the current in a film thickness d, for an exponential law of damping of the generation probability in the interior of the film[1)] $g_{\mathbf{p}}=\bar{g}_{\mathbf{p}}e^{-\alpha z}$, is of the form

$$\mathbf{J}_{\parallel}=\frac{2e}{(2\pi)^3}\sum_{\mathbf{p}}\int d\mathbf{p}\mathbf{v}_{\parallel}\tau\theta(v_z)\bar{g}_{\mathbf{p}}\left\{\frac{l(1-k)}{1-\alpha^2 l^2}\right.$$
$$\left.+\frac{l(1-k')}{2(1-P_1P_2k'^2)}\left[\frac{P_1(k-k')+P_1P_2kk'-1}{1+\alpha l}-\frac{P_2(1-kk')-k+P_1P_2k'}{1-\alpha l}\right]\right\}. \qquad (2)$$

Here $l=|v_z|\tau$, τ is the momentum relaxation time, $k=e^{-\alpha d}$, $k'=e^{-d/l}$, and $P_{2,1}$ are the specularity coefficients of the illuminated and unilluminated surfaces; summation over all the bands is implied in Eq. (2). The quantity α, which determines the attenuation of the light intensity in the interior of the film, has at $\omega\ll\omega_p$ a value $2\omega_p/c$. For interband transitions, the probability of photoexcitation into the band s is given by the expression

$$\bar{g}_{\mathbf{p}}^{(s)}=\frac{\pi e^2E_0^2}{\omega^2}\sum_{s'}|\mathbf{v}_{ss'}\mathbf{e}|^2[\delta(\varepsilon_{ss'}-\omega)f_{s'}(1-f_s)-\delta(\varepsilon_{ss'}+\omega)f_s(1-f_{s'})], \qquad (3)$$

where $\varepsilon_{ss'}\equiv\varepsilon_{s\mathbf{p}}-\varepsilon_{s'\mathbf{p}}$; $\varepsilon_{s\mathbf{p}}$ is the dispersion law in the s-th band; $f_s\equiv f_0(\varepsilon_{s\mathbf{p}})$; $f_0(\varepsilon)$ is the Fermi function. The polarization and amplitude $\mathbf{e}$ and E_0 of the field inside the metal at the boundary of the film can be easily related to the value of the field of the wave incident on the surface of the metal (see, e.g., Ref. 3).

The direction of the vector $\mathbf{v}_{ss'}(p)$, in view of the absence of a preferred vector in the metal, is connected with the vector $\mathbf{p}$, and in a sufficiently isotropic medium we have $\mathbf{v}_{ss'}(\mathbf{p})=v_{ss'}(p)\mathbf{p}/p$. It is seen from (2) and (3) that the photocurrent along the surface is the result of the even anisotropy of the generation on account of unequal relaxation of the electrons moving parallel and antiparallel to the vector $\mathbf{e}$ when $\mathbf{e}$ is inclined to the normal to the film. In two limiting cases—homogeneous generation ($\alpha d\ll 1$, $\alpha l\ll 1$) and a semi-infinite

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sample ($d/l \gg 1$, $\alpha d \gg 1$)—expression (3) for the current coincides with the results of Refs. 1 and 2, respectively.

In the limit $\alpha d \gg 1$, $\alpha l \gg 1$, Eq. (2) takes the simpler form

$$J_{\parallel} = \frac{e}{(2\pi)^3\alpha} \sum (1-P_1) \int d\mathbf{p} v_{\parallel} \tau \bar{g}_{\mathbf{p}} \theta(v_z)(1-k') \frac{1+P_2k'}{1-P_1P_2k'^2}. \tag{4}$$

The τ-approximation used in (2) is valid, strictly speaking, only if the electron momentum relaxes after one collision. This is the case if the principal relaxation mechanism is scattering by impurities, or else emission of phonons with frequency of the order of the Debye frequency ω_D. The first case is realized, for example, in amorphous films and the second occurs if high temperatures ($T \gg \omega_D$) are considered or when the energy of the photoexcited carriers is high ($|\varepsilon - \varepsilon_F| \gg \omega_D$). In both cases the mean free path depends little on the energy and is of the order of several dozen angstroms (in particular, $l \sim v_F/\omega_D$ for relaxation by phonons). Since the photocurrent reaches a maximum at $d/l \gtrsim 1$, it is best to use in experiments films of thickness $d \sim l$. Films with thickness of the order of hundreds of angstroms, obtained by sputtering, are usually amorphous, therefore the film material can be regarded as isotropic. For such thicknesses $\alpha d \ll 1$. Confining ourselves to two bands, we obtain

$$\mathbf{J}_{\parallel} = \frac{e^3E_0^2p_0^2v_{12}^2(p_0)}{4\pi\omega^2\tilde{v}} \operatorname{Re}(\mathbf{e}_{\parallel}(\mathbf{e}^*\mathbf{N})) \sum_{s=1,2} (-1)^s L_s (P_2^{(s)} - P_1^{(s)}) [\Psi_0^{(s)}(d/L_s) - \Psi_1^{(s)}(d/L_s)], \tag{5}$$

where the indices 1 and 2 pertain respectively to the filled and empty bands

$$L_s = \frac{\partial \varepsilon_s(p_0)}{\partial p_0}\tau_s, \quad \tilde{v} = \frac{\partial \varepsilon_{21}(p_0)}{\partial p_0},$$

p_0 is the root of the equation $\varepsilon_{21}(p) = \omega$;

$$\Psi_n^{(s)}(x) = \int_0^1 dy\, y(1-y^2) \frac{e^{-nx/y}}{1 - P_1^{(s)}P_2^{(s)}e^{-2x/y}}.$$

According to (5), $\mathbf{J}_{\parallel} \propto L$ as $d \to \infty$ and $\mathbf{J}_{\parallel} \propto d$ as $d \to 0$. The surface density of the current is of the order of 10^{-6} A/cm at an absorbed light-flux intensity 1 W/cm^2. We note that we have neglected electron-electron collisions. A comparison of their frequency with the electron-phonon collisions shows that this neglect is valid if the energy of the photoexcited carriers is $\varepsilon < (\omega_D \omega_p)^{1/2}$.

PHOTOVOLTAIC EFFECT IN PURE METALLIC FILMS AT LOW TEMPERATURES

This case calls for a special analysis because an electron of energy ε (reckoned from the Fermi energy) lower than ω_D relaxes to the Fermi surface via emission of phonons much more rapidly than it loses momentum. The point is that when a phonon is emitted (the spontaneous emission processes are the principal ones at $\varepsilon \gg T$) the direction of the electron momentum changes by an amount $\Delta p/p_F \sim \varepsilon/\omega_D \ll 1$, whereas the energy changes by an amount equal to its own value. On reaching the temperature layer near the Fermi surface, the excitation begins to defuse along the latter. The characteristic time of the departure from the state with a given momentum is of the order of the transport mean free path time τ_T of the thermalized carriers. This means that in the case of generation in a layer $< \omega_D$ near the Fermi surface the photocurrent is determined by a relatively large transport time τ_T and is therefore large.

Two variants are possible, in which the electron is excited from a point located under and close to the Fermi surface into a higher-lying band, or from a deeply lying band into a point above the Fermi surface. For the sake of argument we shall refer to the first variant. In this case the current is determined mainly by the photoholes, since the electrons excited into the second band relax very rapidly in momentum and therefore make a negligible contribution to the photocurrent. Thus, the relaxation of the holes can be divided into two stages: rapid in energy and slow diffusion of the Fermi surface.

We consider the first stage. We assume that the field damping constant α is much smaller than the reciprocal of the mean free path of the photoexcited carriers that have not yet managed to lose energy. If the photohole has an energy $\sim\omega_D$, then this condition is equivalent to the requirement $v_F/c \ll \omega_D/\omega_p$. In the kinetic equation we can then neglect the spatial derivative. The equation for the distribution function φ_p of the holes produced under the Fermi surface takes the form $\hat{I}\varphi_p = -g_p$, where $\hat{I}$ is the integral of the collisions with the phonons. Just as before, we assume that the dispersion law is close to spherical. However, the degree of nonsphericity (corrugation) of the Fermi surface

$$\delta = (v_F p_F - \mathbf{v}_F \mathbf{p}_F)/\varepsilon_F$$

may turn out to be comparable with the parameter ω_D/ω_F. The photohole can then be produced at a distance comparable with or larger than ω_D from the Fermi surface. Therefore the nonsphericity turns out to be substantial in the generation probability. At the same time, in the collision and diffusion terms of the kinetic equation the corrugation can be neglected if $\delta \ll 1$.

Expanding the solution in spherical harmonics, we obtain

$$\varphi_{\mathbf{p}} = \sum_{lm} \Phi_{lm}(x) Y_{lm}(\mathbf{n}), \quad \mathbf{n} = \mathbf{p}/p, \quad x = (\varepsilon_F - \varepsilon_p)/2sp_F,$$

$$G_{lm} + \int_x^{x+1} dx' (x'-x)^2 P_l(x_0) \theta(x) \Phi_{lm}(x') - \frac{1}{3}\Phi_{lm}(x)[\theta(x-1) + x^3\theta(1-x)] = 0, \tag{6}$$

$$G_{lm} = g_{lm} \frac{\rho s v_F}{2\Lambda^2\pi^2 p_F^3}, \quad g_p = \sum_{lm} g_{lm} Y_{lm}(n),$$

where s is the speed of sound, ρ is the density of crystal, and $x_0 = 1 - 2(x - x')^2$. Equation (6) is valid in the region $x \gg \theta \equiv T/2sp_F$, in which spontaneous emission of phonons predominates. In the derivation of (6) we have assumed for simplicity that there is one type of acoustic phonon with a dispersion law $\omega_q = sq$, and the electron-phonon interaction does not depend on the momentum of the electron (it is described by the constant of the deformation potential Λ).

For weakly corrugated equal-energy surfaces, the main contribution to the anisotropic part of the generation is determined by the second harmonic ($l = 2$), for which the kernel of Eq. (6) takes the form

$$(x'-x)^2[1+6(x'-x)^4-6(x-x')^2]\,\theta(x). \tag{7}$$

We consider first the case when the holes are produced with energies

$$\varepsilon_F-\varepsilon_p\ll\omega_D \quad (x\ll 1).$$

In this case $P_2(x_0)\approx 1$, and the equation is easily transformed into a differential one.

$$(x^3\Phi)'''+6\Phi=3G, \quad \Phi\equiv\Phi_{2m}, \quad G\equiv G_{2m}. \tag{8}$$

Equation (8) has a solution satisfying the condition $\Phi(\infty) = 0$ in the form

$$\Phi(x)=\frac{3}{x^3}G(x)+\frac{18}{11x^4}\int\limits_x^\infty dx'G(x')-\left(\frac{2}{11}\right)^{1/2}\frac{9}{x}\int\limits_0^\infty\frac{dx'}{x'^3}G(x')$$
$$\times\sin\left(\chi+2^{1/2}\ln\frac{x'}{x}\right), \quad \sin\chi=\left(\frac{2}{11}\right)^{1/2}. \tag{9}$$

The first term in (9) represents electrons that have not yet managed to collide with the phonons. The principal term as $x\to 0$ is the second one; as $x\to 0$ we have $\Phi\propto 1/x^4$. It is easy to see that this behavior characterizes also the solution of Eq. (7). It is seen from (9) that at $x\ll 1$ the third (oscillating) term in (9) is small compared with the principal term $\propto 1/x^4$, and we shall hereafter neglect it.

For arbitrary x, the solution of Eqs. (6) can be expressed in terms of the Green's function $\mathcal{G}_l(x,x')$:

$$\Phi_{lm}(x)=\int\limits_0^\infty dx'\mathcal{G}_l(x,x')G_{lm}(x').$$

The Green's function of Eq. (7) was obtained numerically. The figure shows a plot of the asymptotic form of

$$F_2(t)=\frac{11}{18}\lim_{x\to 0}x^4\mathcal{G}_2(x,t),$$

which we shall need later on. It is seen from the figure that $F_2(t)$ is an alternating-sign function with a period ~1. The cause of the alternation of the signs is that the hole distribution function at the point x is determined mainly by transitions from the point x' into x with emission of a phonon of energy $x'-x$. If $x'-x\sim 1/2$, then the emitted phonon has a momentum $\sim p_F$, the hole is rotated on the Fermi surface through an angle $\sim\pi/2$, and this leads to opposite signs of the harmonics of the distribution function at the points x and x'.

Let us investigate the second stage of relaxation, namely diffusion of the electrons over a layer of thickness T near the Fermi surface. Following Ref. 5, we assume that the distribution function in the layer T can be represented in the form

$$\varphi=\frac{\partial f_0}{\partial x}[a(\mathbf{p}_F)+b(x,\mathbf{p}_F)],$$

where $a\gg b$; $f_0(x)=(e^{-x/\theta}+1)^{-1}$ is the Fermi function p_F is the momentum vector drawn to a point on the Fermi surface.

We introduce a certain cutoff energy $\bar\theta$, such that $1\gg\bar\theta\gg\theta$ and $\theta\ll\Delta\omega/\omega_D$, where $\Delta\omega$ is the characteristic initial energy of the photohole. We integrate the initial kinetic equation

$$v_z\frac{\partial\varphi}{\partial z}=\hat{I}\varphi+g$$

over all the energies in a given direction of the momentum (assuming that $|\varepsilon-\varepsilon_F|\ll\varepsilon_F$). The obtained collision integral can be broken up into two parts: $\mathcal{P}_T$, which contains integration of φ over the layer $x\lesssim\bar\theta$, and the part $\mathcal{P}_\infty$ which contains $\varphi(x)$ at $x\gg\theta$. After substituting the solution in the region $x>\bar\theta$, expanded in spherical harmonics, we obtain

$$v_z\frac{\partial a}{\partial z}=\mathcal{P}_T\{a\}+\sum T^{lm}Y_{lm}(\mathbf{n}),$$
$$T^{lm}=\int\limits_0^\infty dx\,g_{lm}(x)+\mathcal{P}_\infty{}^{lm}, \tag{10}$$

$$\mathcal{P}_\infty^{lm}=\frac{2\Lambda^2p_F{}^3\pi^2}{\rho s v_F}\Bigg[\int\limits_0^\infty dx\int\limits_{\max(x,\bar\theta)}^{x+1}dx'\,(x'-x)^2P_l(x_0)\,\Phi_{lm}(x')$$
$$-\frac{1}{3}\int\limits_{\bar\theta}^\infty dx\,({}^3\theta(1-x)+\theta(x-1))\,\Phi_{lm}(x)\Bigg]. \tag{11}$$

The quantity T^{lm} plays a role of the effective generation of photoholes (in the given harmonic) and the Fermi surface, while the operator $\mathcal{P}_T\{a\}$ describes the diffusion of the carriers over the Fermi surface (it is an integral with respect to x of the momentum-direction-nonconserving operator $\mathcal{L}_1$ introduced in Ref. 5).

Using Eq. (6), we can transform the expression for Φ_{lm} into

$$T^{lm}=\frac{2\Lambda^2p_F{}^3\pi^2}{\rho s v_F}\int\limits_0^{\bar\theta}dx\int\limits_{\bar\theta}^{x+1}dx'\,(x'-x)^2P_l(x_0)\,\Phi_{lm}(x'). \tag{12}$$

Substituting in (12) the expression for Φ_{lm} in terms of the Green's function, and using next the asymptotic form of the Green's function as $x\to 0$, we obtain

$$T^{lm}=\int\limits_0^\infty dx\,F_l(x)\,g_{lm}(x). \tag{13}$$

If the generation takes place at small x, then it follows from (13) that

$$T^{lm}=\int\limits_0^\infty dx\,g_{lm}(x)$$

[it is easy to show that $F_l(0) = 1$]. This means that in the case of generation in $x\ll 1$ all the carriers reach the Fermi surface and conserve their momentum direction. In the case of generation in $x\gtrsim 1$

$$|T^{lm}|<\left|\int\limits_0^\infty dx\,g_{lm}\right|,$$

which reflects the loss of momentum. Thus, the problem was reduced to a solution of Eq. (10), in which the quantities T^{lm} are defined with the aid of (13).

Since the angle of the electron scattering in the layer T is small, Eq. (10) can be transformed into a diffusion equation.[5] This solution of the obtained diffusion equation for the case of a film entails considerable difficulties. We shall therefore replace approximately (as is customarily done in the theory of conductivity of metallic films) the quantity $\mathscr{P}_T\{a\}$ by

$$-\frac{a-\langle a\rangle}{\tau_T},$$

where

$$1/\tau_T=\mathbf{v}_F\mathscr{P}_T\{\mathbf{v}_F\}/v_F^2$$

is the averaged reciprocal transport relaxation time.

We can next use Eq. (4) to find the photocurrent. In place of (4) we obtained

$$J\equiv J_x+iJ_y=\frac{2e(1-P_2)\tau_T d}{\alpha}\int\limits_0^1 dy\sum_{l=1}^{\infty}P_{2l}^1(y)\int\limits_0^{\infty}dxF_{2l}(x)(-1)^l(1-y^2)^{1/2}$$
$$\times sp_F^3\left[\frac{(4l+1)(2l-1)!}{(2l+1)!\cdot 4\pi}\right]^{1/2}(1-e^{-\xi/y})\frac{1+P_1e^{-\xi/y}}{1-P_1P_2e^{-2\xi/y}},\qquad(14)$$

where $l_T=v_F\tau_T$ is the thermal mean free path, and $\xi=d/l_T$.

Since the series in l in (14) converges rapidly, we confine ourselves only to the first term. Then the current is the product of a function that depends only on the thickness of the film and on the temperature by a function that depends only on the frequency of the light:

$$\mathbf{J}_\parallel=A(\xi)\mathbf{B}(\omega),$$
$$A(\xi)=el_T(1-P_2)(30\pi)^{1/2}(\Psi_0(\xi)+(P_1-1)\Psi_1(\xi)-P_1\Psi_2(\xi)),$$
$$\mathbf{B}(\omega)=Z(\omega)\int\limits_0^{\infty}dx\,F_2(x)\langle\!\langle\mathbf{n}_\parallel(\mathbf{nN})|\mathbf{ne}|^2\rangle\!\rangle.\qquad(15)$$

Here

$$Z(\omega)=(15/8\pi)^{1/2}e^2v_{12}^2p_F^2E_0^2/4\pi\alpha v_F\omega^2,$$
$$\langle\!\langle\ldots\rangle\!\rangle=\frac{1}{4\pi}\int d\Omega(\ldots)\delta\left(x\frac{v_F'}{v_F}-\frac{\varepsilon_F+\omega-\varepsilon_2(\mathbf{p}_F)}{2sp_F}\right)$$
$$\equiv\left\langle(\ldots)\delta\left(x\frac{v_F'}{v_F}-\frac{\varepsilon_F+\omega-\varepsilon_2(\mathbf{p}_F)}{2sp_F}\right)\right\rangle,$$
$$\mathbf{v}_F'=\frac{\partial\varepsilon_{21}(\mathbf{p}_F)}{\partial\mathbf{p}_F},\quad\mathbf{v}_F=\frac{\partial\varepsilon(\mathbf{p}_F)}{\partial\mathbf{p}_F},\quad\varepsilon(\mathbf{p})\equiv\varepsilon_1(\mathbf{p}).\quad v_{12}\equiv v_{12}(p_F).$$

In the limiting cases we have for $A(\xi)$

$$A(\xi)=\frac{3esp_F^3}{4\pi^2v_F}\left(\frac{5}{6\pi}\right)^{1/2}l_T(1-P_2)$$
$$\times\begin{cases}{}^2/_3\xi([1+P_1]/[1-P_1P_2]), & \xi\ll\max(1,[1-P_1P_2]/2P_1P_2)\\(1+P_1)/12P_1P_2, & (1-P_1P_2)/2P_1P_2\ll\xi\ll 1.\\{}^1/_4, & \xi\gg 1\end{cases}\qquad(16)$$

It is seen from (15) that the orientational and polarization dependences of the current are determined by the tensor $\langle\langle n_in_jn_kn_l\rangle\rangle$ contained in B.

Most simple metals belong to a cubic or hexagonal system. The tensor $\langle\langle n_in_jn_kn_l\rangle\rangle$ in (15) is then represented in the form

$$\langle\!\langle n_in_jn_kn_l\rangle\!\rangle=R_1(\delta_{ij}\delta_{kl}+\delta_{ik}\delta_{jl}+\delta_{il}\delta_{jk})+R_2\gamma_{ijkl}$$

(for a cubic system) and

$$\langle\!\langle n_in_jn_kn_l\rangle\!\rangle=S_1(\delta_{ij}\delta_{kl}+\delta_{ik}\delta_{jl}+\delta_{il}\delta_{jk})$$
$$+S_2c_ic_jc_kc_l+S_3(c_ic_j\delta_{kl}+c_ic_k\delta_{jl}+c_ic_l\delta_{jk}+c_jc_k\delta_{il}+c_jc_l\delta_{ik}+c_kc_l\delta_{ij})$$

(for a hexagonal system). The tensor γ_{ijkl} has along the cube axis components $\gamma_{ijkl}=1$ at $i=j=k=l$ and 0 in the remaining cases. The values of R and S are given by

$$R_1={}^1/_6\langle\!\langle 1-3n_3^4\rangle\!\rangle,\qquad R_2={}^1/_2\langle\!\langle 5n_3^4-1\rangle\!\rangle;$$
$$S_1={}^1/_3\langle\!\langle n_1^4\rangle\!\rangle,\qquad S_2={}^1/_3\langle\!\langle 3n_1^2n_3^2-n_1^4\rangle\!\rangle,$$
$$S_3=\langle\!\langle n_1^4+n_3^4-6n_1^2n_3^2\rangle\!\rangle,$$

the integrals being calculated in a coordinate system connected with the crystallographic axes (in the case of a hexagonal system the sixfold axis c coincides with the 3 axis). All the quantities R and all the quantities S, generally speaking, are of same order. Their characteristic dimension as a function of x is determined by the corrugation, which is characterized by the parameters $(\varepsilon_2''-\varepsilon_2')/2sp_F$, where ε_2'' and ε_2' are the maximum and minimum values of the function $\varepsilon_2(p_F)$.

It is seen from (15) that there is no circular photovoltaic effect. Compared with (1), Eq. (15) contains not only an isotropic term with β_1 determined by the constants R_1 and S_1, but also increments due to the anisotropy of the crystal. The structure of the tensor $\langle\langle n_in_jn_kn_l\rangle\rangle$ shows that in a cubic crystal the term with R_2 does not lead to a current if N is parallel to directions of the type [1, 0, 0]. In a hexagonal crystal at an orientation $c\parallel N$ and $c\perp N$ the contribution of the term with S_2 vanishes.

If the corrugation parameter is small, then $\varepsilon_2(\mathbf{p}_F)\approx\varepsilon_2=\text{const}$ and

$$\mathbf{B}(\omega)=\chi\left(\frac{15}{8\pi}\right)^{1/2}\frac{cE_0^2}{\omega}\cdot\frac{2}{5}F_2\left(\frac{v_F(\omega+\varepsilon_F-\varepsilon_2)}{v_F'\cdot 2sp_F}\right)\mathrm{Re}[\mathbf{e}_\parallel^*(\mathbf{Ne})],\qquad(17)$$

where

$$\chi=2e^2v_0^2p_F^2/3c\omega v_F'\theta(\omega+\varepsilon_F-\varepsilon_2)$$

is the absorption coefficient. If $\omega+\varepsilon_F-\varepsilon_2\ll 2sp_F$ the function $F_2(x)$ can be replaced by $F_2(0)=1$. At $\omega+\varepsilon_F-\varepsilon_2\gg 2sp_F$ the function $\mathbf{B}(\omega)$ oscillates and decreases rapidly in accordance with Fig. 1(a).

We proceed now to the case of large corrugation $(\varepsilon_2''-\varepsilon_2'\gg 2sp_F)$. There are three frequency regions. In the region $\omega<\varepsilon_2'-\varepsilon_F$ the absorption is equal to zero, and consequently there is no current. At $\varepsilon_2'-\varepsilon_F<\omega<\varepsilon_2''-\varepsilon_F$ the current is proportional to the number of holes produced in a layer of thickness $\sim 2sp_F$ near the Fermi surface. Since the function $F_2(x)$ is very sharp, in the expression (15) for $\mathbf{B}(\omega)$ we can take the function

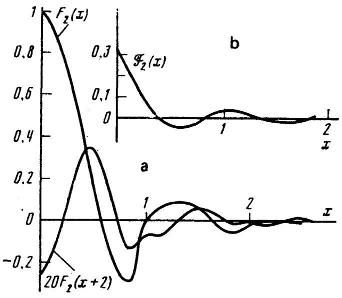

FIG. 1. a) Plot of the function $F_2(x)$. The curve $20F_2(x+2)$ gives the tail of this function at $x>2$. $F_2\to 1$ as $x\to 0$, corresponding to conservation of the electron momentum upon relaxation to the Fermi surface. b) Plot of

$$\mathscr{F}_2(x)=\int\limits_x^{\infty}dx\,F_2(x);\quad\mathscr{F}_2(0)\approx 0.32$$

$\langle\langle \mathbf{n}_\parallel(\mathbf{n}\cdot\mathbf{N})|\mathbf{n}\cdot\mathbf{e}|^2\rangle\rangle$ outside the integral with respect to x at the point $x=0$. As a result we get

$$\mathbf{B}(\omega)=Z(\omega)\mathcal{F}_2(0)\frac{s}{2\pi p_F v_F'}\oint dl\frac{\mathbf{n}_\parallel(\mathbf{nN})\,|\mathbf{ne}|^2}{[1-(\mathbf{v}_F\mathbf{v}_F')^2/v_F^2 v_F'^2]^{1/2}}, \tag{18}$$

where $\mathcal{F}_2(x)=\int_x dx' F_2(x')dx'$, $\mathcal{F}_2(0)\approx 0.32$, and $\oint dl$ stands for integration over all the intersection lines of the surfaces $\varepsilon_1(\mathbf{p}_F)=\varepsilon_F$ and $\varepsilon_2(\mathbf{p}_F)-\omega=\varepsilon_F$.

In the region $\omega>\varepsilon_2''-\varepsilon_F$, the main contribution to the PVE is determined by the holes produced closest to the Fermi surface, i.e., at the points $p_F^{(\alpha)}$ at which $\varepsilon_2(\mathbf{p}_F)$ is maximal. In this case we have for $\mathbf{B}(\omega)$

$$\mathbf{B}(\omega)=\frac{12\pi s}{p_F^5}(M_1M_2)^{1/2}\mathcal{F}_2\left(\frac{v_F(\omega+\varepsilon_F-\varepsilon_2'')}{v_F'\cdot 2sp_F}\right)\sum \mathbf{p}_{F\parallel}^{(\alpha)}(\mathbf{p}_F^{(\alpha)}\mathbf{N})\,|\mathbf{p}_F^{(\alpha)}\mathbf{e}|^2. \tag{19}$$

The summation in (19) is over the star $p_F^{(\alpha)}(\varepsilon_2(p_F^{(\alpha)})=\varepsilon_2'')$, and the quantities $1/M_{1,2}$ are the principal values of the tensor $\nabla_i\nabla_j\varepsilon_2(\mathbf{p}_F^{(\alpha)})$.

Thus, in the case of a relatively large corrugation, $\varepsilon_F\gg\varepsilon_2''-\varepsilon_2'\gg 2sp_F$, the photocurrent is a maximum and depends little on the frequency in the region where the surfaces $\varepsilon_{21}(\mathbf{p})=\omega$ and $\varepsilon_1(\mathbf{p})=\varepsilon_F$ intersect. The order of magnitude of the current, referred to the photon flux density, is determined by the product of the charge by the characteristic mean free path—the smaller of the quantities d and l_T—and by the factor $sp_F/(\varepsilon_2''-\varepsilon_2')$. The last factor decreases its value in comparison with the case of weak corrugation. In the frequency region $\omega>\varepsilon_2''-\varepsilon_F$ the photocurrent, as seen from Fig. 1(b), decreases with frequency in an oscillatory manner.

CONCLUSION

We have demonstrated in this paper the possibility of a new type of photoeffect in metals, namely the onset of a surface current in a metal illuminated by interband-absorbed light. This current is as a rule of the order 10^{-6} A/cm per W/cm^2 of the absorbed light power. In the special case of a pure metal at low temperature, the current near the absorption edge can increase all the way to 0.1 A · cm/W.

Let us discuss the assumptions used in the present paper. We regard as the most significant of them the very rarely satisfied condition of weak corrugation of the equal-energy surfaces. The weak corrugation and closing of the Fermi surface, as is well known, takes place only in alkali metals (in which the tight-binding approximation is well applicable). We do not know of any experimental data on the structure of other bands in these materials. It is natural to assume, however, that the equal-energy surfaces in high-lying bands deviate therefore little from spherical. Thus, in alkali metals the surface $\varepsilon_{21}(\mathbf{p})=\omega$ is apparently close to a sphere, and consequently the described theory is well applicable to them. One can expect it to be well applicable also to metals in which the greater part of the Fermi surface is close to a sphere, such as gold, silver, and copper. Although in these metals the Fermi surface is open, the bridges constitute a small fraction of its area. Finally, even if the equal energy surfaces are strongly nonspherical, the surface provides a qualitative description of the phenomenon. In the case described by the τ approximation is not very sensitive to the form of the dispersion laws. In the case of low temperatures, the condition for a large PVE is (just as in the case of weak corrugation) the intersection of the equal-energy surfaces $\varepsilon_{21}(\mathbf{p})=\omega$ and $\varepsilon_1(\mathbf{p})=\varepsilon_F$.

We have considered in detail the PVE in transitions of the electron from under the Fermi surface into an empty band. It is meaningful to discuss the second variant of the optical transition, namely excitation of an electron from a filled band into a point over the Fermi surface. Such a transition lies usually in the x-ray band. The frequencies of these transitions exceed the plasma frequency, and therefore α in the expressions for the current should be taken to mean the absorption coefficient. The parameter αl is always large in this case. In this situation, the effect reaches its maximum value of the surfaces $\varepsilon_{12}(\mathbf{p})=\omega$ and $\varepsilon_1(\mathbf{p})=\varepsilon_F$ intersect. The dispersion law in the deeply lying band can be regarded as close to a constant. If the width of the lower band is smaller not only than ε_F but also ω_D, then the photoexcited carriers are produced at a fixed energy distance from the Fermi surface, equal to $\varepsilon_F-\varepsilon_2-\omega$. It is then possible to use Eq. (17), in which the sign must be reversed.

It should be noted that the mechanisms considered in this paper lead to the appearance of only a linear surface PVE determined by the constant β_1 in (1). Without going into the causes of this, we note that a circular PVE can arise when account is taken of the electron spin and of the spin-flip processes in collisions with the surface.

The noted partial conservation of the electron momentum as its energy relaxes may be important in the analysis of phenomena related to the surface PVE, i.e., in situations when anisotropic production of hot carriers takes place, such as electron-photon dragging and the volume PVE. In these cases it is to be expected that if the energy of the photoexcited carriers is $\lesssim\omega_D$, then their effective mean free path will coincide with the mean free path of the thermalized carriers.

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[1)] This dependence is realized in the case of the normal skin effect, i.e., when the condition $\omega/\omega_p\gg v_F/c$ is satisfied (ω_p is the plasma frequency and v_F is the Fermi velocity). In typical metals, this region begins with optical wavelengths shorter than 10 μm.

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