Photoelectron Emission in a Ferromagnetic

A. Z. VEKSLER
Institute for the Physics of Metals, Ural Affiliate, Academy of Sciences, USSR
(Submitted to JETP editor May 10, 1954)
J. Exper. Theoret. Phys. USSR **29**, 201-208 (August, 1955)

A formula is derived which determines the velocity distribution of photoelectrons and the temperature dependence of the photocurrent in the neighborhood of the Curie point. The calculations are carried out on the basis of the $s - d$ exchange model[1], with the periodic potential of the lattice being taken account of by the method of variation of parameters. It is shown that the photocurrent depends quadratically on the magnetization, in accordance with the results of Cardwell's work[2].

1. INTRODUCTION

THE theory of photoelectronic emission in a ferromagnetic is given in reference 3, where a formula for the photocurrent at a light frequency in the neighborhood of threshold and in a temperature region in the neighborhood of the Curie point is derived. A simplified theory of the photoeffect[4,5] was used in the derivation of this formula, as was the $s - d$ exchange model, taking account of the interaction of the s electron with the d electrons, thanks to which one is able to connect the electrical, thermal and mechanical properties of a ferromagnetic with its magnetic state. Failure to take account of the interaction of the s electrons within the framework of the $s - d$ exchange model would certainly be a defect; hence, it must be considered as one of the stages in the development of a consistent theory of ferromagnetics.

The use in reference 3 of a theory of the photoeffect which does not take account of the periodic potential is inappropriate, since special features of the motion of the electrons in ferromagnetics depend upon their interaction, which is but weakly connected with the properties of the limiting boundary. Hence, the motion of the electrons is not to be described as the sum of two plane waves[5,6], but more correctly as a function of the form:

$$\psi(\mathbf{r}) = e^{i\mathbf{k}\mathbf{r}} u_{\mathbf{k}}(\mathbf{r}),$$

in which

$$u_{\mathbf{k}}(\mathbf{r}) = u_{\mathbf{k}}(\mathbf{r} + \mathbf{n}a),$$

where a is the lattice parameter. In this case there is no necessity of dividing the photoeffect into two parts: volume and surface. Moreover, the results obtained in the present work show that the photocurrent consists of two similar parts. Thus, the usual theory of the photoeffect is substantiated. Taking account of the periodic potential in a systematic way leads to the fact that the velocity distribution of the photoelectrons is different from that resulting from a theory of the photoeffect which does not take account of this potential.

2. THE WAVE FUNCTION OF THE ELECTRON

Let the metal occupy the half space $x < x_0$ and let the region $x \geq x_0$ be vacuum. Let us consider the emission of s-electrons from the metal which is produced by the absorption of light.

The wave function of the electron is determined by the equation:

$$\Delta\psi_1 + \frac{2m}{\hbar^2}[E - V(x, y, z)]\psi_1 = 0 \qquad (2.1)$$

for $x < x_0$,

$$\Delta\psi_2 + \frac{2m}{\hbar^2}\left[E + \frac{e^2}{4x}\right]\psi_2 = 0 \qquad (2.1')$$

for $x \geqslant x_0$,

where $V(x, y, z)$ is a periodic function. The solution of Eq. (2.1) has the following form:

$$\psi_1 = \sum_n \{a_n \exp[i(g_x + q_{n_1})(x - x_0)] \qquad (2.2)$$

$$+ b_n \exp[-i(g_x + q_{n_1})(x - x_0)]\}$$

$$\times \exp\{i[(g_y + q_{n_2})y + (g_z + q_{n_3})z]\},$$

where $q_{n_i} = 2\pi n_i / a$, n_i is an integer, g_x, g_y, g_z are the components of the quasi-momentum of the

[1] S. V. Vonsovskii, J. Exper. Theoret. Phys. USSR **16**, 981 (1946)

[2] A. B. Cardwell, Phys. Rev. **76**, 125 (1949)

[3] S. V. Vonsovskii and A. V. Sokolov, Dokl. Akad. Nauk SSSR **76**, 197 (1951)

[4] I. E. Tamm, Z. Phys. **68**, 97 (1931)

[5] R. H. Fowler, Phys. Rev. **38**, 45 (1931)

[6] K. Mitchell, Proc. Roy. Soc. (London) **A153**, 513 (1936)

electron.

The second equation can be solved by the method of separation of variables, setting

$$\psi_2 = f_1(x) f_2(y) f_3(z).$$

Then, on putting ψ_2 into Eq. (2.1′), we obtain

$$f_2(y) = e^{if_y y}, \quad f_3(z) = e^{if_z z}, \tag{2.3}$$

while $f_1(x)$ satisfies the equation

$$\frac{d^2 f_1}{dx^2} + \frac{2m}{h^2}\left(\frac{h^2 f_x^2}{2m} + \frac{e^2}{4x}\right) f_1 = 0, \tag{2.4}$$

and we have

$$\frac{h^2}{2m}(f_x^2 + f_y^2 + f_z^2) = E. \tag{2.5}$$

The general solution of Eq. (2.1′) has the form:

$$\psi_2(x, y, z) = \sum_{f_x, f_y, f_z} C_f f_1(x) \exp\{i(f_y y + f_z z)\}. \tag{2.6}$$

The coefficients C_f are determined by the boundary conditions

$$\psi_1|_{x=x_0} = \psi_2|_{x=x_0}, \quad \psi_1'|_{x=x_0} = \psi_1'|_{x=x_0}. \tag{2.7}$$

In the present instance the prime indicates differentiation with respect to x. These conditions can be satisfied if Eqs. (2.2) and (2.6) are compared with each other term by term. In this case identity results on fulfillment of the equations

$$f_y = g_y + q_{n_2}, \quad f_z = g_z + q_{n_3}. \tag{2.8}$$

Taking account of relation (2.5), we find

$$f_x^2 = \frac{2m}{h^2} E - (g_y + q_{n_2})^2 - (g_z + q_{n_3})^2. \tag{2.9}$$

Thus, the summation in Eq. (2.6) is to be carried out over n_2 and n_3

$$\psi_2(x, y, z) = \sum_{n_2, n_3}' C_{n_2 n_3} f_1(x) \exp\{i[(g_y + q_{n_2}) y + (g_z + q_{n_3}) z]\}. \tag{2.6′}$$

At great distances from the dividing boundary this function must have the form:

$$\psi_2(x, y, z) = \sum_{n_2 n_3} C_{n_2 n_3} \exp\{i[f_x x + (g_y + q_{n_2}) y + (g_z + q_{n_3}) z]\} \tag{2.10}$$

or

$$\psi_2(x, y, z) = \sum_{n_2 n_3} C_{n_2 n_3} \exp\{-\varkappa_n x + i[(g_y + q_{n_2}) y + (g_z + q_{n_3}) z]\}, \tag{2.10′}$$

where $\kappa_x = -if_x$ if f_x is an imaginary number.

States of the electron which are described by the function (2.10) will be referred to in what follows as "free", while those described by the function (2.10′) will be called "bound". A finite nonzero probability of finding the electron at an arbitrarily large distance from the dividing boundary corresponds to the first case. The corresponding current is also different from zero. In the second case, the current is equal to zero, and the probability of finding the electron at a very great distance from the dividing boundary is small and tends to zero with increasing x.

Calculation of the current and averaging it over the coordinates gives for the first case the expression

$$j_x = \frac{eh}{m} \sum_{n_2 n_3} |C_{n_2 n_3}|^2 f_x. \tag{2.11}$$

The photoeffect for light of frequency near the threshold will be considered below. In this case, the overwhelming majority of the photoelectrons leaving the metal have an energy of the order of thermal. As was shown in reference 7, formula (2.11) then is simplified and takes the form

$$j_x = \frac{eh}{m} |C_{0,0}|^2 f_x. \tag{2.12}$$

3. THE CURRENT OF PHOTOELECTRONS

The wave function describing the state of the photoelectrons is the solution of the equation

$$\frac{h^3}{2m} \Delta u = \frac{h}{i} \frac{\partial u}{\partial t} - W(x, y, z) u - \frac{eh}{imc} (\mathbf{A}\vec{\nabla}) u; \tag{3.1}$$

$$\mathbf{A} = 2\mathbf{a} \cos\left[\omega\left(t - \frac{x \cos\theta + y \sin\theta}{c}\right)\right], \tag{3.2}$$

where $W(x, y, z)$ is the periodic potential in the region of the metal, with $W = -e^2/4x$ for the region outside the metal.

Ordinarily the term standing at the right in Eq. (3.1) is small in comparison with the term on the

---

[7] A. V. Sokolov and A. Z. Veksler, J. Exper. Theoret. Phys. USSR **25**, 215 (1953)

left side. Hence, Eq. (3.1) may be solved by the method of perturbation theory.

We seek a solution of Eq. (3.1) in the form

$$u = u_g + v, \tag{3.3}$$

where

$$u_g = \psi(x, y, z)\exp\{-iE_g t/\hbar\}, \tag{3.4}$$

subject to the condition $|u_g| \gg |v|$.

On substituting Eq. (3.3) into Eq. (3.1) and discarding the term containing $(\mathbf{A}\vec{\nabla})v$, we obtain

$$\frac{\hbar^2}{2m}\Delta v + \frac{h}{i}\frac{\partial v}{\partial t} - W(x, y, z)u = \frac{eh}{imc}(\mathbf{A}\vec{\nabla})u_g. \tag{3.5}$$

We seek the solution of this equation in the form of an expansion in terms of the integral of the eigenfunctions of the unperturbed problem:

$$v = \iint_{-\infty}^{+\infty} dg_y dg_z \Big\{\int_0^{\infty} c_g^+(t)\,u_g^+(t)\,dg_x + \int_0^{\infty} c_g^-(t)\,u_g^-(t)\,dg_x\Big\}, \tag{3.6}$$

where

$$u^-_{g_x, g_y, g_z} = u^+_{-g_x, g_y, g_z} = u_{-g_x, g_y, g_z}.$$

On putting Eq. (3.6) into Eq. (3.5), multiplying once by u_g^{+*} and again by u_g^{-*}, and integrating, we find

$$\iint_{-\infty}^{+\infty} dg_y dg_z \Big\{\int_0^{\infty} dg_x \frac{dc_g^+(t)}{dt}\int d\tau\, u_g^+ u_g^{\alpha *} + \int_0^{\infty} dg_x \frac{dc_g^-(t)}{dt}\int d\tau u_g^- u_g^{\alpha *}\Big\} = \frac{e}{mc}\int u_g^{\alpha *}(\mathbf{A}\vec{\nabla})u_g\, d\tau, \tag{3.7}$$

where the index α takes on the values plus (+) or minus (–).

In virtue of the orthogonality of the functions u_g, we obtain

$$\int d\tau\, u_g^+ u_{g'}^{+*} = \int d\tau\, u_g^- u_{g'}^{-*} = N_g^{(1)}\delta(\mathbf{g} - \mathbf{g}'), \tag{3.8}$$

$$\int d\tau\, u_g^+ u_{g'}^{-*} = \int d\tau\, u_g^{+*} u_{g'}^- = N_g^{(2)}\delta(\mathbf{g} - \mathbf{g}'),$$

where $\delta(\mathbf{g} - \mathbf{g}')$ is the usual δ function. Hence, Eq. (3.7) can be written in the form:

$$N_g^{(1)}\frac{dc_g^+(t)}{dt} + N_g^{(2)}\frac{dc_g^-(t)}{dt} = \frac{e}{mc}\int d\tau\, u_g^{+*}(\mathbf{A}\vec{\nabla})u_g, \tag{3.9}$$

$$N_g^{(2)}\frac{dc_g^+(t)}{dt} + N_g^{(1)}\frac{dc_g^-(t)}{dt} = \frac{e}{mc}\int d\tau\, u_g^{-*}(\mathbf{A}\vec{\nabla})u_g.$$

Eliminating dc_g/dt and integrating with respect to the time, we obtain

$$c_g^{\pm}(t) = \frac{\exp[(i/\hbar)(E_g + h\omega - E_{g'})t] - 1}{(i/\hbar)(E_g + h\omega - E_{g'})}(\mathbf{g}\,|\,\mathbf{a}\,|\,\mathbf{g}')^{\pm}, \tag{3.10}$$

where

$$(\mathbf{g}\,|\,\mathbf{a}\,|\,\mathbf{g}')^{\pm} = \frac{(e/mc)}{[N_g^{(1)}]^2 - [N_g^{(2)}]^2}\int d\tau\,[(\mathbf{a}\vec{\nabla})\psi_g]\,[N_g^{(1)}\psi_g^{\pm *} - N_g^{(2)}\psi_g^{\mp *}]. \tag{3.11}$$

In the calculation, the term containing $\exp[i(E_g + h\omega - E_{g'})(t/h)]$, which corresponds to radiation, is discarded. Moreover, the factor $\exp[i\omega\frac{x\cos\theta + y\sin\theta}{c}]$ is omitted, since the case under consideration is that for which the frequency of the light is near threshold, in virtue of which $\omega/c \ll g$.

The energy connected with the quasi-momentum function[1] is

$$E = B + B'y + (C + C'_y)(\cos g_x a + \cos g_y a + \cos g_z a), \tag{3.12}$$

where γ is the magnetization; B, B', C, C' are parameters depending on the energy of interaction of the s and d electrons and the energy of exchange of the s electrons. From this it is readily found that

$$N_g^{(1)} = \frac{(2\pi)^3\hbar^2}{2m(C + C'y)\sin g_x a}\Big\{\sum_{n, n'}(2ag_x + 2\pi n_1 + 2\pi n_1')(a_n a_{n'}^* - b_n b_{n'}^*) + \sum_{n_1 n_2}|C_{n_1 n_2}|^2 afx\Big\}; \tag{3.13}$$

$$N_g^{(2)} = \frac{(2\pi)^3\hbar^2}{2m(C + C'y)\sin g_x a}\Big\{\sum_{n, n'}(2ag_x + 2\pi n_1 + 2\pi n_1')(a_n b_{n'}^* + a_{n'}^* b_n)\Big\}.$$

On putting Eq. (3.10) into Eq. (3.6) and carrying out the integration, we find

$$v = \frac{(2\pi)^3 (\mathbf{g} \mid \mathbf{a} \mid \mathbf{g}_1)^+}{(C + C'y)\, a \sin g_{1x} a} u_{g_1}(x, y, z), \quad (3.14)$$

where $\mathbf{g}_1$ is the magnitude of $\mathbf{g}'$ corresponding to the equalities:

$$E_{g_1} = E_g + \hbar\omega; \quad (3.15)$$
$$g_{1y} = g_y + \frac{2\pi n_2}{a}, \quad g_{1z} = g_z + \frac{2\pi n_3}{a}.$$

The matrix element $(\mathbf{g} \mid \mathbf{a} \mid \mathbf{g}_1)$ breaks down into two parts, one of which corresponds to the region inside the metal, the other to the region external to it. On carrying out the integration over the basic region, we find that the part of the matrix element which is connected with the wave function of the electron in the metal is equal to

$$(\mathbf{g} \mid \mathbf{a} \mid \mathbf{g}_1)_1^+ \quad (3.16)$$
$$= a_x \Bigg\{ N_g^{(1)} \sum_{n_1, n_1'} (g_{1x} + q_{n_1}) \Bigg[\frac{a_n^* a_{n'} + b_n b_{n'}^*}{g_x + q_{n_1} - g_{1x} - q_{n_1'}} - \frac{a_n^* b_{n'} + a_{n'} b_n^*}{g_x + q_{n_1} + g_{1x} + q_{n_1'}} \Bigg]$$
$$+ N_g^{(2)} \sum_{n_1, n_1'} (g_{1x} + q_{n_1'}) \Bigg[\frac{a_n^* a_{n'} + b_n^* b_{n'}}{g_x + g_{n_1} + g_{1x} + q_{n_1'}} - \frac{a_n^* b_{n'}' + b_n^* a_{n'}}{g_x + q_{n_1} - g_{1x} - q_{n_1'}} \Bigg] \Bigg\},$$

where a_x is the component of the vector $\mathbf{a}$ along the x axis.

The calculation of the matrix element for the region inside the metal is extremely difficult. Hence, we shall limit ourselves merely to a determination of its dependence on the wave function $\mathbf{f}$ and the parameter κ. In the calculation of the matrix element we can use in place of the wave function of the bound state its asymptotic representation[6]

$$f_\varkappa(x) = b_\varkappa e^{-\varkappa x} (2\varkappa x)^{s/\varkappa}, \quad (3.17)$$

where $s = me^2/4h^2$. If $\kappa \geq 10^8$ cm^{-1}, then such a substitution is valid down to $x = (3 - 4)\, 10^{-8}$ cm.

An asymptotic decomposition of the type of Eq. (3.17) is not valid for the wave functions of the free states in view of the fact that at light frequencies near threshold the component of the wave vector f_x has an order of magnitude of 10^7 cm^{-1}. Instead of the asymptotic decomposition we can make use of the result of the work of reference 6, where it is shown that for frequencies in the neighborhood of threshold,

$$f_{f_x}(x) = \lambda(x) \; ; \quad (3.18)$$

to first approximation, this does not depend on f_x. Thus, the part of the matrix element for the region outside the metal and for a light frequency in the neighborhood of the critical point can be written in the form

$$(\mathbf{g} \mid \mathbf{a} \mid \mathbf{g}_1)_2 = a_x b_x \Big[N_g^{(1)} - N_g^{(2)} \Big] \int_{x_0}^{\infty} \lambda^*(x) \frac{d}{dx} \left\{ e^{-\varkappa x} (2\varkappa x)^{s/\varkappa} \right\} dx. \quad (3.19)$$

The term corresponding to the equality

$$g_{1x} = g_x + q_{n_1} - q_{n_1'}, \quad (3.20)$$

which, because of the zonal periodicity of the energy

$$E(\mathbf{g} + \mathbf{q}_n) = E(\mathbf{g}) \quad (3.21)$$

is not in conformity with the law of conservation of energy (3.15), has been dropped in Eq. (3.16). It can occur only for an excitation of the electron such that the latter goes over into the next zone. In view of the fact that in the present work we are considering the emission of photoelectrons for a light frequency near threshold, the number of electrons going over into the next zone is small, and hence the term corresponding to the condition (3.20) may be neglected.

The member in the matrix element, the x component of the quasi-momentum of which satisfies Eq. (3.20), corresponds to the "volume" photoeffect. The selection rules for the electron here are the same as for optical transitions. Ordinarily, the "volume" photoeffect is considered separately from the "surface" photoeffect. The calculation of the photocurrent in the present work does not require such a separation. Both kinds of transitions of the electrons are obtained automatically as component parts of the matrix element.

Calculation of the current with the aid of Eqs. (2.10) and (3.14) leads to the following result:

$$j_x = \frac{(2\pi)^6 e\hbar \mid (\mathbf{g} \mid \mathbf{a} \mid \mathbf{g}_1)^+ \mid^2}{m (C + C'y)^2 a^2 \sin^2 g_{1x} a} f_x \mid C_{0,0} \mid^2. \quad (3.22)$$

As was shown in reference 7, $f_x \mid C_{0,0} \mid^2$ does not depend on f_x. Let us designate

$$f_x \mid C_{0,0} \mid^2 = A(g_{1x}). \quad (3.23)$$

Then

$$j_x = \frac{(2\pi)^6 e\hbar \mid (\mathbf{g} \mid \mathbf{a} \mid \mathbf{g}_1) \mid^2 A(g_{1x})}{m (C + C'y)^2 a^2 \sin^2 g_{1x} a} \quad (3.24)$$

Integrating over the initial states, we find the complete current

$$I_x = \frac{1}{8\pi^3} \int_{-\pi/a}^{+\pi/a} dg_y \int_{-\pi/a}^{+\pi/a} dg_y \times \int_{f_x=0}^{\infty} dg_x j_x \frac{1}{\exp[(E-\varepsilon_0)/kT]+1} \quad (3.25)$$

We expand $\frac{1}{\exp[(E-\epsilon_0)/kT]+1}$ as a series in powers of $\exp\left[\frac{E-\epsilon_0}{kT}\right]$. If $E > \epsilon_0$, then

$$\frac{1}{\exp[(E-\varepsilon_0)/kT]+1} = \sum_n (-1)^n \exp\left\{-(n+1)\frac{E-\varepsilon_0}{kT}\right\} \quad (3.26)$$

For $E < \epsilon_0$ we obtain

$$\frac{1}{\exp[(E-\varepsilon_0)/kT]+1} = \sum_n (-1)^n \exp\left\{n\frac{E-\varepsilon_0}{kT}\right\}. \quad (3.26')$$

We make use of the equality

$$E = \frac{h^2}{2m}(f_x^2 + g_y^2 + g_z^2) - h\omega$$

and make a change of the variables of integration: in place of g_x, g_y, g_z we take

$$\rho^2 = g_y^2 + g_z^2, \quad \theta = \arctan(g_y/g_z) \text{ and } f_x.$$

Taking outside the integral sign those factors which are only weakly dependent of ρ and f_x, we obtain for $E > \epsilon_0$

$$I_x = -\frac{(2\pi)^4 mek^2T^2 |(g|a|g_1)^+|^2}{h^3(C+C'y)^3 a^3 \sin^3 g_{1x}a} A(g_{1x}) \times \sum_n \frac{(-1)^n \exp\{(n+1)(h\omega+\varepsilon_0)/kT\}}{(n+1)^2} \quad (3.27)$$

If the frequency of the light is sufficient to knock out electrons with energy $E < \epsilon_0$, then in this case the region of integration can properly be divided into two: $E < \epsilon_0$ and $E > \epsilon_0$. As a result of the calculation we obtain the following expression:

$$I_x = -\frac{(2\pi)^4 mek^2T^2 |(g|a|g_1)^+|^2 A(g_{1x})}{h^3(C+C'y)^3 a^3 \sin^3 g_{1x}a} \times \left\{\frac{\pi^2}{6} + \frac{(\varepsilon_0+h\omega)^2}{2k^2T^2} + \sum_n \frac{(-1)^n \exp\{-n(h\omega+\varepsilon_0)/kT\}}{n^2}\right\}. \quad (3.27')$$

It should be remarked that $\epsilon_0 < 0$ and is equal to the work of emission with sign reversed[7]. The expression obtained for the photocurrent has the same temperature dependence as that given in reference 4. If is fully confirmed by experimental results.

Using the Eq. (2.14) of reference 7 it is not difficult to find an expression for the total photocurrent in the neighborhodd of the Curie point (in the approximation of a weakly filled zone)

$$I_x = I_x^+ + I_x^- = (A_0 + B_0 y^2) T^2 \Phi_{1,2}\left[\frac{\varepsilon(\omega)+\beta_1 y^2}{kT}\right], \quad (3.28)$$

where

$$A_0 = -\frac{(2\pi)^4 mek^2A}{h^2a^4\alpha_1^{3/2}} M_0 \sum_{n_1} |a_{n,\,00}|^2 (\gamma_1 + q_{n_1}),$$

$$B_0 = -\frac{16\pi^4 mek^2A}{h^2a^4\alpha_1^{3/2}} \times \sum_{n_1} \left\{M_0\gamma_3 + \left(M_1 - \frac{3}{2}\frac{\alpha_2}{\alpha_1}M_0\right)\gamma_2 + \left[\left(\frac{15}{8}\frac{\alpha_2^2}{\alpha_1^2} - \frac{3}{2}\frac{\alpha_3}{\alpha_1}\right)M_0 - \frac{3}{2}M_1\frac{\alpha_2}{\alpha_1}\right](\gamma_1 + q_{n_1})\right\} |a_{n,\,00}|^2;$$

$$\Phi_1(x) = \sum_n \frac{(-1)^{n+1}}{(n+1)^2} e^{(n+1)x}$$

for the case described by formula (3.27),

$$\Phi_2(x) = \left[\frac{x^2}{2} + \frac{\pi^2}{6} + \sum \frac{(-1)^n e^{-nx}}{n^2}\right]$$

for the case described by formula (3.27'); M_0, M_1 and M_2, are, respectively, the first, second and third coefficients in the expansion of the squared modulus of the matrix element in powers of y:

$$|(g|a|g_1)|^2 = M_0 + M_1 y + M_2 y^2;$$

I_x^+ and I_x^- are the total currents of electrons with right and left orientations of mechanical moment, respectively;

$$\alpha_1 = B^2 + 3C^2 + \varepsilon^2(\omega) + 4BC - 2(B+2C)\varepsilon(\omega);$$

$$\alpha_2 = 2CC' + 2\varepsilon(\omega)(B'+2C');$$

$$\alpha_3 = 3C'^2 + B'^2 + 4B'C' + \beta'(B + 2C);$$

$$\gamma_1 = \arccos \xi_0; \quad \gamma_2 = \frac{\beta_2}{(1-\xi_0^2)^{1/2}};$$

$$\gamma_3 = -\frac{\xi_0 \beta_2^2}{(1-\xi_0^2)^{3/2}} - \frac{\beta_3}{(1-\xi_0^2)^{1/2}};$$

$$\xi_0 = \frac{\varepsilon(\omega) - (B+2C)}{C}; \quad \beta_2 = \frac{B'+2C'}{C} + \xi_0 \frac{C'}{C};$$

$$\beta_3 = \xi_0 \frac{C'^2}{C^2} + \frac{B'+2C'}{C^2} C' + \frac{\beta_1}{C};$$

$$\varepsilon(\omega) = \alpha + a^2 (3\pi^2 n)^{2/3} \beta + h\omega;$$

$$\beta_1 = a^2 (3\pi^2 n)^{2/3} k_1 \left(\frac{3}{8} k_1 \beta + \frac{3}{2} \beta' \right).$$

The symbols α, α', β, β', n, k_1 are taken from the work of Vonsovskii[1]. Formula (3.28) can obviously be used for nonferromagnetic metals if $\gamma = 0$ is inserted in it.

In this case, there results a temperature dependence of the same form as in the simplified theory of the photoeffect which does not take account of the periodic potential. The distinction between the simplified view of the photoeffect and the more coherent one is especially noticeable on comparison of the expressions determining the velocity distribution of the photoelectrons [Eq. (3.24)]. This dependence is more complicated than in the simplified theory, which, apparently, is what actually occurs.

If the photoeffect produced by light of frequency much greater than threshold is considered, then the second component of the matrix element, corresponding to the so-called "volume" effect, must also be taken into account.

A comparison of Eq. (3.28) for the photocurrent with the well-known relation determining its temperature dependence shows that the work of emission must be equal to $-\epsilon_0$. This result can be substantiated not only by a formal comparison, but also on the basis of thermodynamic relations[7].

In conclusion, the author expresses his profound gratitude to A. V. Sokolov for a number of valuable suggestions and to S. V. Vonsovskii for a discussion of the present work.

Translated by Brother Simon Peter, F. S. C.
174

SOVIET PHYSICS VOLUME 2, NUMBER 2 MARCH, 1956

International Values for the Thermal Cross Section of Fissionable Isotopes

At the session of August 17* of the International Conference on the Peaceful Application of Atomic Energy, a large amount of declassified data was presented on the effective neutron cross sections of the fissionable isotopes U-233, U-235 and Pu-239. In the low energy region, for which a large number of measurements was reported, excellent agreement was obtained for these isotopes, which play important roles in reactor installations. At the initiation of the Chairman, the scientists of France, Great Britain, the USSR and the USA met after the official session to consider the effective cross sections of absorption and fission of these isotopes by thermal neutrons (with velocities of 2200 m/sec). It was decided to develop a system of international mean values for these effective cross sections. Such values would contribute to agreement of reactor calculations based on these constants. The errors in the mean international values listed here are based on the scatter of reported values and in some instances exceed the errors of particular individual measurements.

	absorption in barns	fission in barns
U-233	593 ± 8	524 ± 8
U-235	698 ± 10	590 ± 15
Pu-239	1032 ± 15	729 ± 15

* Session 17A, August 17, 1955. "The Effective Cross Section of Fissionable Isotopes." Chairman, D. Hughes (USA), Vice-chairman, D. Popovich (Yugoslavia).

Translated by R. T. Beyer
237