relatively high limit of γ-ray luminosity of the source 3C-273 is apparently connected with the large value of R for this object, for even at n = 0 (and accordingly n(95%) = 3) the value of W_γ will have the same order of magnitude.

The upper limit of the average γ-ray luminosity for infrared galaxies W_γ, can also be indirectly estimated from the intensity of the isotropic background of the hard γ rays F_γ [8] and the average density of these objects N. At the present time it is impossible to determine exactly the value of N from observations, but estimates by Burbidge [9] show that the number of powerful infrared galaxies possibly amounts to about 1% of the number of all galaxies. In order of magnitude, $N \sim R^{-3}$, where R is the distance from the nearest infrared galaxies. Then $W_\gamma \lesssim 4\pi F_\gamma H/cN \lesssim 10^{42}$ erg/sec. In this case the experimental values of W_γ do not contradict the assumption that the isotropic background of the hard γ rays is due to emission of infrared galaxies. In spite of the fact that the obtained upper limits of W_γ are relatively high, it should be noted that for all the sources these quantities do not exceed the power of the infrared radiation [1]. Such a situation does not agree with the Low hypothesis concerning the annihilation nature of infrared radiation of such objects [2]. It was assumed there that the infrared radiation is synchrotron radiation of the electrons and positrons produced upon annihilation of nucleons and antinucleons. The power of the γ radiation in the energy region ~100 MeV, in the case of such a mechanism, would exceed at least by a factor of 2 the power of the infrared radiation (see, e.g., [10]).

The authors are sincerely grateful to R.A. Syunyaev, at whose initiative this attempt to use the experimental results obtained with the "Kosmos-208" was undertaken for the determination of the upper limits of the luminosity of the indicated objects in γ rays.

[1] D.E. Kleinman and F.J. Low, Astrophys. J. 159, L165 (1970).
[2] F.J. Low, Astrophys. J. 159, L173 (1970).
[3] K.Kellermann and I.I.K. Pauliny-Toth, Ann. Rev. Astron. Astrophys. 6, 417, (1968).
[4] G.G. Fazio, Ann. Rev. Astron. Astrophys. 5, 481 (1967).
[5] G.G. Fazio, H.F. Helmken, S.J. Cavrak, and D.R. Hearn, Can. J. Phys. 46, S427 (1968).
[6] G.M. Frye and Chia Ping Wang, Astrophys. J. 158, 925 (1969).
[7] C.E. Fichtel, T.L. Cline, C.H. Ehrmann, D.A. Kniffen, and R.W. Ross, Can. J. Phys. 46, S419 (1969).
[8] L.S. Bratolyubova-Tsulukidze, N.L. Frigorov, L.F. Kalinkin, A.S. Melioranskii, E.A. Pryakhin, I.A. Savenko, and V.Ya. Yufarkin, Kosmicheskie issledovaniya (Space Research) 8, 136 (1970).
[9] G.R. Burbidge, Ann. Rev. Astron. Astrophys. 8, 360 (1970).
[10] H. Alfven and A. Elvins, Science 164, 911 (1969).

LOW TEMPERATURE ELECTRIC CONDUCTIVITY OF METALS WITH OPEN FERMI SURFACES

R.N. Gurzhi and A.I. Kopeliovich
Physico-technical Institute of Low Temperatures, Ukrainian Academy of Sciences
Submitted 5 April 1971
ZhETF Pis. Red. 13, No. 10, 569 - 572 (20 May 1971)

The present paper is devoted to the development of the theory of lattice resistance of pure metals at low temperatures, with account taken of the dragging of the phonons by the electrons.

As indicated by Peierls [1], in metals with open Fermi surfaces the Bloch law $\rho \sim T^5$ (ρ - electric resistivity, T - temperature) should be valid at arbitrary low temperatures, because the Umklapp processes in collisions between electrons with thermal phonons become possible no matter how small the momenta of these phonons may be.

It is known that electron phonon collisions can be regarded as elastic with high degree of accuracy, i.e., upon absorption (or emission) of a phonon, the electron experiences transitions within the limits of the equal-energy surface. One can therefore expect that at sufficiently low temperatures, when the thermal momentum of the phonons is $q \sim T/s \ll p_F$ (s - speed of sound, p_F - smallest characteristic dimension of the Fermi surface) the problem of electric conductivity can be formulated in terms of the diffusion of electrons in momentum space on the Fermi surface.

In a quantitative study of the question, it is natural to start from a system of kinetic equations for the interacting electrons and phonons[1], and to carry out in these equations a low-temperature expansion in the small parameter q/p_F. The diffusion equation can be obtained immediately in a relatively simple and compact form if the kinetic equation for the electrons is integrated over a certain arbitrary part of the Fermi surface prior to carrying out the low-temperature expansion, and the conservation of the number of electrons is used. The result takes the following form[2]:

$$\operatorname{div} \hat{D}_p(\nabla \chi_p - a_p) = eEn_p,$$

where

$$D_p^{ik} = T^5 \frac{30\zeta(5)}{\pi^2 \hbar^4 v_p^2} \int \frac{M_p(e)}{s^5(e)} e^i e^k \delta(e n_p)\, de$$

$$= \int D_p(e) e^i e^k \delta(e n_p)\, de,$$

$$a_p = \int \hat{A}_{pp'} \nabla \chi_{p'}\, ds_{p'}, \quad \hat{A}_{pp'} = \frac{2\hat{D}_p^{-1} \hat{D}_p(\vec{\mu}) D_{p'}(\vec{\mu})}{\sin(n_p, n_{p'})\gamma(\vec{\mu})},$$

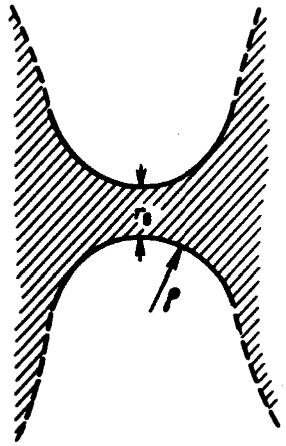
$$\gamma(\vec{\mu}) = \int D_p(\vec{\mu}) \delta(\vec{\mu} \cdot n_p)\, ds_p,$$

In these formulas, $\chi_p \partial f_0/\partial \varepsilon$ is the non-equilibrium part of the electron distribution function, $f_0(\varepsilon)$ is the Fermi function, $\vec{E}$ is the intensity of the electric field, $\vec{n}_p = \vec{v}_p/v_p$, $e = \vec{q}/q$, $\vec{\mu} = \vec{n}_p \times \vec{n}_{p'}/|\vec{n}_p \times \vec{n}_{p'}|$, $\vec{v}_p = \partial\varepsilon/\partial p$, $D_p^{ik}(\vec{e}) = e^i e^k D_p(\vec{e})$. The matrix element of the electron phonon interaction is written in the form $[qM_p(\vec{e})]^{1/2}$. The integration in the expressions for $\vec{a}_p$ and $\gamma(\vec{\mu})$ is carried out over the Fermi surface. By div and ∇ are meant two-dimensional operations carried out in the tangential plane to the Fermi surface.

---

[1] Collisions between phonons at low temperatures, as is well known, are much less probable than phonon-electron collisions, and are therefore disregarded.

[2] We note that the integral term a_p in Eq. (1) is connected with the fact that the phonons are not in equilibrium and is a reflection of the fact that each electron exchanges phonons with other electrons during the diffusion process.



The sought function χ_p should satisfy the periodic boundary conditions

$$\chi_p = \chi_{p+g}, \quad \nabla\chi_p = \nabla\chi_{p+g},$$

which in the diffusion approximation are equivalent to allowance for the Umklapp processes (g is an arbitrary reciprocal-lattice vector).

A detailed analysis of (1) will be carried out in a forthcoming article. We present here the results of its solution for a special (but far from rare) case, when the Fermi surface consists of a certain number of electron (or hole) groups randomly interconnected by narrow connecting necks. (Some of the groups can be isolated.) As can be readily shown, within the limits of the large groups, the electron distribution function has the drift form

$$\chi_p = -\mathbf{u}\mathbf{p}.$$

In the model in question, the solution of the diffusion equation (1) reduces to the problem of the flow of stationary electric current through a branched network. The necks correspond to conductors, and the large groups to the nodes of the network, χ_p has the meaning of the potential. The current flowing through the neck is determined as follows:

$$I = \oint (\hat{D}\,\nabla\chi)\, d\vec{l},$$

where the integration is over a closed line including the neck, and the vector $d\vec{l}$ is directed along the normal to this line in the tangential plane to the Fermi surface. The resistance of the neck is determined by the relation $\Delta\chi = IR$, where $\Delta\chi$ is the "potential" drop across the neck.

The currents satisfy the Kirchhoff law

$$\Sigma I = 0, \quad \Sigma IR = \mathbf{g}\mathbf{u}.$$

In the first expression, the summation is over the necks interconnected in a certain node; in the second it is over the necks forming a closed circuit. $\vec{g}$ is the summary Umklapp vector and closes the circuit ($\vec{g} = 0$ if the circuit does not intersect the boundary of the zone).

In addition, the requirement that the quasimomentum be in balance imposes the following condition on the currents crossing the boundary of the zone

$$\Sigma I\mathbf{g} = \frac{2}{h^3}\, e(n_e - n_h)\mathbf{E},$$

where $\vec{g}$ are the corresponding Umklapp vectors, and n_e and n_h are the densities of the electrons and holes.

The foregoing equations make it possible to determine the drift velocity u and consequently the electric conductivity of the metal. (The current density $\vec{j} = e(n_e - n_h)\vec{u}$.) For the simplest case, when all the necks intersect the boundaries of the zone and the metal has cubic symmetry, the electric conductivity is equal to

$$\sigma = \frac{2e^2}{h^3} \frac{(n_e - n_h)^2}{\sum_i g_i^2 \cos^2 a_i / R_i} ,$$

where α_i is the angle between the Umklapp vector $\vec{g}_i$ corresponding to the i-th neck and an arbitrary fixed direction.

The resistance of the neck R differs essentially in two limiting cases: 1) the characteristic width of the neck r_0 is of the order of the radius of curvature ρ of the neck (see the figure) and 2) $r_0 \ll \rho$. In the former case $R_1 = (1/\pi D)\ln p_0/r_0$, where D is the diagonal component of the tensor $\hat{D}$ along the generator of the neck, p_0 is the characteristic radius of the large group. In the second case $R_2 = (1/2D)\sqrt{\rho/r_0}$, and it turns out that the characteristic length of the neck is $\ell \sim \sqrt{\rho r_0}$.

We note that the presence of necks of the second type in the intermediate region of temperatures, when $\ell \gg q \sim T/s \gg r_0$, the developed approach remains in force, but the diffusion has a one-dimensional character. It turns out that in this case the resistance of the neck is $R \approx R_2 T/sr_0 \sim T^{-4}$. Accordingly, the resistance of the metal will also be proportional to T^4, thus differing from the Bloch law T^5, which is valid only when $T \ll sr_0$. This result may explain the experimentally known T^4 law for aluminum.

We note finally that the diffusion equation (1) can be solved only in the case of weak coupling, under the condition that the Harrison sphere is intersected only by one pair of Bragg planes.

[1] R. Peierls, Quantum Theory of Solids, Oxford, 1955.

POSSIBLE METHOD OF MEASURING THE ELECTROMAGNETIC FORM FACTOR OF THE NEUTRINO

V.I. Andryushin, S.M. Bilen'kii, and S.S. Gershtein
Joint Institute for Nuclear Research
Submitted 8 April 1971
ZhETF Pis. Red. 13, No. 10, 573 - 576 (20 May 1971)

A study of the electromagnetic neutrino interaction is one of the important problems of weak-interaction physics. The calculation of the electromagnetic form factor of the neutrino induced by weak interaction is dealt with in [1 - 3]. Since the weak-interaction theory is non-renormalizable, the results of the calculations depend on the cutoff parameter. The induced form factor of the neutrino depends also on whether these exists an intermediate boson, and on the electromagnetic interaction of the intermediate boson.

Some information concerning the electromagnetic radius of the neutrino can be obtained from astrophysical data. It is shown in [4] that the electromagnetic radius of ν_e and also the radius of ν_μ (in the case when the mass of ν_μ is smaller than 1 keV), does not exceed 4×10^{-14} cm.

In the present article we discuss a method that makes it possible, in principle, to obtain complete information on the electromagnetic form factor of the neutrino, including the dependence of the form factor on the square of the momentum transfer. The method is based on a comparison of the cross section of the process due to the electromagnetic interaction of the neutrino with the nucleus, with the cross section of the analogous process due to the electron.