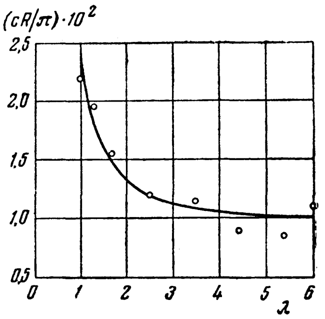

Dependence of real part of surface impedance of silver on the wavelength. Abscissa – $cR/\pi = 4n/(n^2 + \kappa^2)$, where $n - i\kappa$ is the complex index of refraction of the metal.

of the internal photoeffect, ω the frequency of the light, and ν_0 the frequency of the collisions between the electrons and the lattice. Therefore, making the important assumption that the Fermi surface is isotropic, we can determine the concentration of the conduction electrons M and the electron velocity v on the Fermi surface by using the expressions obtained for a clearly pronounced anamolous skin effect.[1] Here the reflection of the electrons from the surface of the metal is assumed to be diffuse. We have already determined N earlier, using our measurement data, (for silver $N = 5.2 \times 10^{22}\,cm^{-3}$); allowance for the term due to the interelectron collisions makes it possible to employ the quantity R_0 for the determination of v. We obtained $v = 2.4 \times 10^8$ cm/sec for silver.

Measurements of the optical constants of tin and lead in the spectral region $1 - 6\,\mu$ has shown that the contribution of the interelectron collisions to the real part of the surface impedance is substantial for these metals, too. However, the processing of the results obtained for these metals is made complicated by the circumstance that the inequality $\omega^2 \gg \nu_0^2$ is not satisfied in the above spectral region, by virtue of which it is impossible to employ the expressions obtained for the sharply-pronounced anomalous skin effect. Considerably less reliable are the calculations of the surface impedance or optical constants of the metal in the region where $\omega \sim \nu_0$. We used the expressions for the real and imaginary parts of the surface impedance, obtained for this region by Dingle.[5] Assuming that in this region the interelectron collisions lead to the appearance of a term B/λ^2 in the expression for R, we can determine for these metals not only N, but also v. Such a treatment yielded the following microscopic parameters: for tin, $N = 4.2 \times 10^{22}\,cm^{-3}$ and $v = 2.6 \times 10^8$ cm/sec (this value was obtained for $cR_0/\pi = 2.9 \times 10^{-2}$ and $cB/\pi = 24 \times 10^{-2}\,\mu^2$); for lead, $N = 3.8 \times 10^{22}\,cm^{-3}$ and $v = 4.0 \times 10^8$ cm/sec (this value was obtained for $cR_0/\pi = 5.1 \times 10^{-2}$ and $cB/\pi = 11.2 \times 10^{-2}\,\mu^2$).

In Ref. 4 we determined the upper limit of N for these metals, using a limiting formula which, as is clear from what has been said above, is not valid. Therefore, the estimate obtained there for the upper limit is too rough and is improved here. However, even these data should be considered as approximate, for we used in the calculation theoretical formulas which cannot be considered very reliable. To obtain more reliable values of these microscopic parameters it is necessary to measure the optical constants of Sn and Pb at low temperatures, at which ν_0 become substantially smaller.

In conclusion I express my gratitude to V. L. Ginzburg for discussing the results of this work.

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Translated by J. G. Adashko
147

ABSORPTION OF POLARIZED MU⁻ MESONS BY NUCLEI

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As was shown in Ref. 1, the angular distribution of neutrons produced upon capture of polarized μ^- mesons by free protons is of the form $1 + \alpha \cos\theta$. In this work we have calculated the coefficient α for the case of absorption of μ^- mesons by protons bound in the nuclei, for the scalar (s), vector (v), tensor (t) and pseudo-vector (a) versions of the four-fermion interaction of μ mesons with nucleons. The shell model is used and the recoil of the nucleus is neglected. Since the nucleons in the nucleus acquire energies on the order of several Mev by capture of a μ^- meson, the calculation is made in the nonrelativistic approximation with respect to the nucleons. The interaction of the emitted neutron with the nucleus was taken into account with the aid of a complex potential, which made it possible to take into account a definite neutron-absorption probability in the nucleus. It must be noted here that our value of α may be too low, owing to the superposition of an isotropic background of neutrons from the decay of the compound nucleus, formed as a result of absorption of the neutron, on the anisotropic angular distribution of the neutrons emitted from the nucleus directly after the μ^- capture. This effect calls for an additional evaluation.

The wave function of the proton was taken in the form

$$\psi_P = R_{njl}(r)\,\Omega_{jlj_z}(\mathbf{r}/r), \quad j = l \pm 1/2,$$

where n, j, and ℓ are the quantum numbers that characterize the subshell in the nucleus, and $\Omega_{j\ell j_z}$ is a spherical spinor.[2] If the spin-orbit interaction of the neutron with the nucleus is neglected, the wave function of the neutron becomes

$$\psi_{Ns_z} = \sum_{L=0}^{\infty} i^L (2L+1)\, a_L(r)\, P_L(\mathbf{k}_N \mathbf{r}/k_N r)\, \chi_{s_z} \qquad (1a)$$

which is asymptotically represented by a combination of a plane wave with a converging one (creation condition[3]). If the spin-orbit interaction is taken into account

$$\psi_{Ns_z} = 4\pi \sum_{ILI_z} i^L a_{IL}(r)\, (\chi_{s_z} \Omega^*_{ILI_z}(\mathbf{k}_N/k_N))\, \Omega_{ILI_z}(\mathbf{r}/r),$$

$$I = L \pm 1/2, \qquad (1b)$$

with an analogous asymptotic representation. The wave function of the neutrino was assumed plane, the dependence of the μ^--meson wave-function coordinates on the K orbit was neglected. We introduce the following notation

$$b_{L\lambda njl}(k_N) = \int_0^\infty a^*_L(r)\, j_\lambda(k_\nu r)\, R_{njl}(r)\, r^2\, dr,$$

$$b_{IL\lambda njl}(k_N) = \int_0^\infty a^*_{IL}(r)\, j_\lambda(k_\nu r)\, R_{njl}(r)\, r^2\, dr,$$

where j_λ is a spherical Bessel function and k_ν is the wave number of the neutrino

$$\varphi_{ik} = f_{ii} + f_{kk} + 2\,\mathrm{Re}\, f_{ik}, \quad \eta_{ik} = h_{ii} + h_{ik} + 2\,\mathrm{Re}\, h_{ik}$$

(i, k ≡ s, v, t, a; for a definition of f_{ik} and h_{ik} see Ref. 1);

$$A_{njl}(k_N) = \sum_{L\lambda} (2L+1)(2\lambda+1)(C^{l0}_{L0\lambda0})^2 |b_{L\lambda njl}(k_N)|^2;$$

$$B_{njl}(k_N) = \mathrm{Re}\sum_{L\lambda} [(2L+1)(2L+2)(2L+3)(2\lambda+1)$$
$$\times (2\lambda+2)(2\lambda+3)]^{1/2} \times [C^{l0}_{L0\lambda0} C^{l0}_{L+10\lambda+10} W(L+1 l 1\lambda;$$
$$\lambda+1 L)\, b_{L\lambda njl}(k_N)\, b^*_{L+1\lambda+1njl}(k_N) + C^{l0}_{L+10\lambda0} C^{l0}_{L0\lambda+10} W(Ll1\lambda;$$
$$\lambda+1 L+1)\, b_{L+1\lambda njl}(k_N)\, b^*_{L\lambda+1njl}(k_N)];$$

$$C_{njl}(k_N) = \sum_{IL\lambda} (2I+1)(2L+1)$$
$$\times (2\lambda+1)[C^{l0}_{L0\lambda0} W(Is\lambda l; Lj)]^2 |b_{IL\lambda njl}(k_N)|^2;$$

$$D_{njl}(k_N) = (2/(2l+1)) \sum_{IL\lambda} (2I+1)$$
$$\times (2\lambda+1)(C^{l0}_{L0\lambda0})^2 |b_{IL\lambda njl}(k_N)|^2;$$

$$E_{njl}(k_N) = \frac{(-)^l}{3} \sum_{IL\lambda I'L'\lambda'} (-)^{\sigma_{L'\lambda'}} (2I+1)(2I'+1)(2L+1)$$
$$\times (2L'+1)(2\lambda+1)(2\lambda'+1)\, C^{l0}_{L0\lambda0} C^{l0}_{L'0\lambda'0} C^{10}_{L0L'0} C^{10}_{\lambda0\lambda'0}$$
$$\times W(I'j1\lambda; \lambda'I)\, W(I's1L; L'I)\, W(Is\lambda l; Lj)$$
$$\times W(I's\lambda'l; L'j)\, \mathrm{Re}\,[b^*_{IL\lambda njl}(k_N)\, b_{I'L'\lambda'njl}(k_N)],$$

where

$$\sigma_{L'\lambda'} = \begin{cases} 0 \text{ for } L' = L \pm 1, \ \lambda' = \lambda \pm 1 \\ 1 \text{ for } L' = L \pm 1, \ \lambda' = \lambda \mp 1 \end{cases} \begin{matrix} I = L \pm 1/2 \\ I' = L' \pm 1/2, \quad s \equiv 1/2; \end{matrix}$$

$$F_{njl}(k_N) = 4(-)^l\, \mathrm{Re} \sum_{IL\lambda I'L'\lambda'} (-)^{I+I'}\, i^{L+\lambda-L'-\lambda'} (2I+1)$$
$$\times (2I'+1)(2L+1) \times (2L'+1)(2\lambda+1)$$
$$\times (2\lambda'+1)\, C^{10}_{\lambda0\lambda'0} C^{l0}_{L0\lambda0} C^{l0}_{L'0\lambda'0}$$
$$\times \Big\{ W(lj1s; sl) \sum_{gf} (2g+1)\, C^{10}_{g010} C^{f0}_{g010} C^{f0}_{L0L'0}\, W(L'sfI; I'L)$$
$$\times X(\lambda1\lambda'; LgL'; l1l)\, X(LgL'; IfI'; s1s)$$
$$+ 2 \sum_{gfphr} (2g+1)(2f+1)(2p+1)(2h+1)$$
$$\times C^{10}_{h0L0} C^{r1}_{L'011} C^{r1}_{h011}\, W(\lambda'1pL; fh)\, X(l\lambda L; jgI; s1s)$$
$$\times X(l\lambda'L'; jfI'; s1s)\, X(jfI'; gpL'\, ILs)$$
$$\times X(ph\lambda'; L'r1; g1\lambda) \Big\}\, b^*_{IL\lambda njl}(k_N)\, b_{I'L'\lambda'njl}(k_N),$$

where W are the Racah coefficients; for a definition of the quantities X, see Ref. 4.

Accurate to a constant factor independent of n, j, and ℓ, the probability of emission of a neutron with energy $E_N = \hbar^2 k_N^2/2m$ at a given angle θ, upon absorption of a μ^- meson by protons in a closed subshell characterized by the quantum numbers n, j, and ℓ, is given for the superposition of the s, v, t, and a versions by the following formulas. In case (1a)

$$W_{njl}(k_N; \theta) = \frac{2j+1}{2l+1} [(\varphi_{sv} + 3\varphi_{ta})\, A_{njl}(k_N)$$
$$+ (-\eta_{sv} + \eta_{ta})\, B_{njl}(k_N) \cos\theta]. \qquad (2)$$

In case (1b)

$$W_{njl}(k_N;\theta) = (2j+1)\{[(\varphi_{sv}-\varphi_{ta})C_{njl}(k_N) + \varphi_{ta}D_{njl}(k_N)] + [(-\eta_{sv}+\eta_{ta})E_{njl}(k_N) + \eta_{ta}F_{njl}(k_N)]\cos\theta\}. \quad (3)$$

The total effect due to all closed subshells in the nucleus is obtained by summation of (2) and (3) over n, j, and ℓ. Formulas (2) and (3) describe also the absorption of a μ^- mesons by a single proton located in above the closed shells (in this case the factor $(2j+1)$ must be omitted). One would expect these formulas to be a good approximation for the twice-magic nuclei (for example, ${}_{20}Ca^{40}$). The details of the calculation and numerical estimates for specific nuclei will be given in a separate article.

We are sincerely grateful to I. S. Shapiro for attention to this work and for discussion of the result.

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Translated by J. G. Adashko
148

POSSIBLE ASYMMETRY OF PARTICLES AND ANTIPARTICLES IN WEAK INTERACTIONS

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RECENT experimental data concerning the $\beta-\nu$ correlation[1] show that the positron decays of Ne^{19} and A^{35} may be explained as a mixture of A and V covariants of the β-interaction, while the electron decays of He^6 and the free neutron are dependent on T and, evidently, S covariants. Analogously, associated with the positron decays of Co^{58}, experiment gives a weak interference due to Fermi and Gamow-Teller interactions, while this interference is strong in the electron decays of Au^{198} and Sc^{46}. If we consider that all the experiments are valid, then the indicated facts sharply contradict existing theory. Indeed, if we assume that the processes $n \rightleftarrows p + e^- + \overline{\nu}'$ and $p \rightleftarrows n + e^+ + \nu$ occur as a result of different interactions, this indicates a denial of the symmetry of particles and antiparticles in weak interactions. In addition, antiparticles no longer bear an exact resemblance to the particles with opposite charge. For example, their masses may differ by the order of magnitude of g^2, the square of the weak-interaction constant.

The self-adjoint relativistic invariant Hamiltonian for β decay, in which e^-, e^+, ν, and ν' enter asymmetrically, may be written in the form

$$\begin{aligned} H = \sum_{i=1}^{5} (\overline{\Psi}_p O_i \Psi_n)\,[\Phi^{+}_{\overline{e}}\gamma_4 O_i (g_i + g'_i\gamma_5)\Phi^{+}_{\overline{\nu}} \\ + \Phi^{+}_{\overline{e}}\gamma_4 O_i (f_i + f'_i\gamma_5)\Phi_\nu + \Phi_{e^+}\gamma_4 O_i (\lambda_i + \lambda'_i\gamma_5)\Phi^{+}_{\overline{\nu}} \\ + \Phi_{e^+}\gamma_4 O_i (\mu_i + \mu'_i\gamma_5)\Phi_\nu] + \sum_{i-1}^{5} (\overline{\Psi}_n O_i \Psi_p)\cdot \\ \times [\Phi_{\overline{\nu}}\gamma^4 (g^*_i - g'^*_i\gamma_5) O_i \Phi_{\overline{e}} + \Phi^{+}_{\nu}\gamma_4 (f^*_i - f'^*_i\gamma_5) O_i \Phi_{e^-} \\ + \Phi_{\overline{\nu}}\gamma_4 (\lambda^*_i - \lambda'^*_i\gamma_5) O_i \Phi^{+}_{e^+} + \Phi^{+}_{\nu}\gamma_4 (\mu^*_i - \mu'^*_i\gamma_5) O_i \Phi^{+}_{e^+}]. \end{aligned} \quad (1)$$

In addition, in the case of the neutrino field, $\Phi_\nu(x)$ and $\Phi^{+}_{\overline{\nu}}(x)$ are respectively the positive and negative frequency components of $\Psi_\nu(x)$, so that

$$\begin{aligned} \Phi_\nu &= -i\int S^+(x-x')\gamma_4\psi_\nu(x')\,d^3x', \\ \Phi^{+}_{\overline{\nu}} &= -i\int S^-(x-x')\gamma_4\psi_\nu(x')\,d^3x' \quad (x_0 = x_0') \end{aligned} \quad (2)$$

and analogously for electrons.

In (1), g_i, g'_i and μ^*_i, μ'^*_i refer to β^- and β^+ decays (the experiments may be satisfied, letting g_T, $g_S(g_V?)$ and μ_A, μ_V be unequal to zero), f^*_i, f'^*_i describe K capture, and λ^*_i, λ'^*_i describe the absorption of an antineutrino by a proton. It is easily seen that upon reflection of the spatial coordinates (P) the Hamiltonian preserves its form if the unprimed constants remain unchanged while the primed ones change sign; with reversal of time (T) it is necessary to substitute for every constant its complex conjugate; with charge conjugation (C), $g_i \rightleftarrows \mu^*_i$, $g'_i \rightleftarrows -\mu'^*_i$, $f_i \rightleftarrows \lambda^*_i$, $f'_i \rightleftarrows -\lambda'^*_i$. From this it is evident that the Hamiltonian (1) is invariant under PTC only when $g_i = \mu_i$, $g'_i = \mu'_i$, $f_i = \lambda_i$, $f'_i = \lambda'_i$. We obtain