On the other hand, the maximum contribution to the phonon thermal conductivity is made by phonons with wavelength $\lambda_T \simeq T/\hbar v_\alpha$. Therefore the condition (6) actually denotes $\lambda_T = r_D$, apart from a numerical factor that depends, naturally, on the frequency index n of the non-electronic absorption mechanism. Relation (6) determines in essence the values of the concentration of the electrons N and of the temperature T, at which the effect considered above is possible. It is important that condition (6) does not contain the constant of the electromechanical coupling - it determines only the magnitude of the effect. In order for the effect to be large, it is necessary to satisfy near the maximum, i.e., near $x \simeq 1$, the condition

$$(a + b^\alpha)^2 < a\Delta^\alpha \qquad (7)$$

for all the acoustic modes simultaneously. Recognizing that near the maximum $a \simeq b^\alpha$, we obtain the sufficient condition $\Delta^\alpha/a > 1$. If a stronger condition is satisfied in place of (7), then the change of the thermal conductivity will be gigantic. By way of an example, let us estimate the change of the partial phonon thermal conductivity κ_α for a fixed α, say, due only to longitudinal phonons. For a crystal of the CdS type, the optimal values of the concentration and of the temperature will be $N = 10^{16}$ cm^{-3} and $T = 4°K$. If at the same time $\tau_{ph} \simeq 10^{-9}$ sec, then the partial value of the thermal conductivity is decreased by a factor of 3 compared with the value of $N = 0$ (or $N = \infty$).

The indicated effect will be large if the symmetry of the crystal and the chosen direction of heat-flow propagation (the x direction) are such that the components of the piezo tensor $\beta_{x,\,xx}$, $\beta_{x,\,xy}$, and $\beta_{x,xz}$ are simultaneously different from zero. It is precisely under these conditions that the free paths of both the longitudinal and transverse phonons will change simultaneously with changing concentration (temperature).

In conclusion we note that in a crystal with a sufficiently strong electron-phonon interaction, all the mechanisms (magnetic, electric fields, traps, etc.) influencing the electron decrement can lead to a sharp change in the phonon thermal conductivity of the crystal. This principle can be used to develop devices for automatic control and stabilization of the temperature.

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CHOICE OF OPTIMAL CONDITIONS FOR EXPERIMENTAL OBSERVATION OF COHERENT SCATTERING OF γ QUANTA BY NUCLEI

V.G. Kudryavtsev and M.I. Ryazanov
Moscow Engineering Physics Institute
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1. An experimental study of coherent scattering of γ quanta in a Coulomb field is of considerable interest, since this process is the simplest process connected with the nonlinear properties of vacuum in quantum electrodynamics [1, 2]. However, observation of coherent scattering by nuclei at not too large quantum frequencies $\omega \lesssim m$ is inhibited by the need for separating the "nuclear" scattering from the background of the much more probable process of Compton scattering. The intensity of scattering by nuclei greatly exceeds the intensity of Compton scattering only at very high frequencies, $\omega \gg m$, and small

scattering angles $\theta \lesssim m/\omega$. The purpose of the present paper is to call attention to the fact that the ratio of the cross sections for scattering by nuclei and Compton scattering is greatly increased in these processes if these processes are considered not in an amorphous matter, but in a single crystal. This circumstance can be used in choosing the optimal conditions for the observation of coherent scattering of γ quanta by nuclei.

2. Let us consider the γ quantum entering a single crystal in a direction close to the direction of the crystallographic axis x. The effective longitudinal scatterings, through which the scattering of the quantum is effected, are determined by the longitudinal transmitted wavelength, i.e., by the reciprocal of the momentum transferred in the longitudinal direction [3]. In scattering through an angle θ, the effective length $l_n \sim 1/\omega\theta^2$ for scattering by nuclei is ω/m times larger than the effective length for Compton scattering $l_C \sim (2m/\omega^2\theta^2)$ (at $\theta^2 < m/\omega$). The difference between the effective lengths leads to a different influence of the crystal structure on the two processes.

In a single crystal of length L, the ratio of the cross sections of the "nuclear" and Compton scattering can be obtained in the form

$$\frac{d\sigma_n(\omega,\theta)}{d\sigma_C(\omega,\theta)} = \frac{d\sigma_n^0(\omega,\theta)}{d\sigma_C^0(\omega,\theta)} \frac{\sin^2(L/l_n)}{\sin^2(L/l_C)} \frac{\sin^2(a/l_C)}{\sin^2(a/l_n)},$$

where $d\sigma_n^0(\omega, \theta)$ and $d\sigma_C^0(\omega, \theta)$ are the cross sections for the scattering by the nucleus and for Compton scattering by an isolated atom, and a is the lattice constant. It is obvious that in the region

$$a < l_C < L < l_n$$

the ratio of the cross sections in the crystal

$$\frac{d\sigma_n(\omega,\theta)}{d\sigma_C(\omega,\theta)} = \frac{d\sigma_n^0(\omega,\theta)}{d\sigma_C^0(\omega,\theta)} \frac{L^2}{l_C^2} \frac{1}{\sin^2(L/l_C)} \gg \frac{d\sigma_n^0(\omega,\theta)}{d\sigma_C^0(\omega,\theta)} \qquad (1)$$

is much larger than the ratio of the cross sections for the isolated atom. On the other hand, if $L \gtrsim l_n m$ then we can assume that $\sin^2(L/l_n) \sim \sin^2(L/l_C)$ and the cross section ratio becomes

$$\frac{d\sigma_n(\omega,\theta)}{d\sigma_C(\omega,\theta)} = \frac{d\sigma_n^0(\omega,\theta)}{d\sigma_C^0(\omega,\theta)} \left(\frac{l_n}{l_C}\right)^2 = \frac{d\sigma_n^0(\omega,\theta)}{d\sigma_C^0(\omega,\theta)} \left(\frac{\omega}{m}\right)^2. \qquad (2)$$

For example, for the elements of the end of the periodic system, in the case when $\omega = 25$ MeV, $\theta = 0.1^\circ$ ($\theta = 2 \times 10^{-3}$), and $(\omega/m\theta) \sim 10^{-1}$, the ratio of the cross sections of scattering by a nucleus and of Compton scattering by an individual atom is ~ 5, but the ratio of the same cross sections in a single crystal may increase to $\sim 10^4$.

3. Thermal lattice vibrations decrease the ratio (1) or (2). Allowance for the thermal vibrations leads to the appearance of the Debye-Waller factor $\exp\{-q^2\langle u^2\rangle\}$ in the expressions for the cross sections (q - momentum transferred to the lattice, $\langle u^2\rangle$ - mean squared thermal displacement of the atom). Therefore the applicability of formulas (1) - (2) is limited to the inequality $q^2\langle u^2\rangle \ll 1$, i.e.,

$$\omega^2 m^{-4} \gg \langle v^2\rangle .$$

It must be emphasized that the angular distribution has a specific form of a series of Bragg maxima, and in the cases of experimental investigation of one maximum it is necessary that the scattering by the thermal lattice vibrations not bring the quantum outside of the region of the maximum. If the observation is in a relatively broad range of angles, including several maxima, then the interference pattern becomes averaged out and the ratio of the cross sections will be smaller than (2).

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CONTRIBUTION TO THE THEORY OF ELECTRON SPIN RESONANCE IN ANISOTROPIC METALS

G.P. Alodzhants
P.N. Lebedev Physics Institute, USSR Academy of Sciences

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1. Spin waves were recently observed experimentally in nonferromagnetic metals placed in a magnetic field $\vec{B}$ [1]. The theoretical prediction of the possible existence of such waves was made by V.P. Silin [2] on the basis of the theory of a degenerate electron liquid [3]. The experiments of Dunifer and Schultz [1] are based on the use of the phenomenon, predicted in [4], of selective transparency of metallic films in the case of spin resonance of the conduction electrons. In these experiments, besides the usual transparency line at $\omega = \omega_s = 2\mu_0 B/h$ (μ_0 - magnetic moment of the electron, h - Planck's constant), there is observed a series of additional lines. The presence of the additional lines is a reflection of the fact that a possibility arises for the magnetization to propagate from the skin layer to the interior of the metal, in the form of spin waves. The experiments were performed in sodium and potassium, the Fermi surfaces of which are approximated with sufficient accuracy by a sphere. Therefore these experiments were analyzed by using the model of an isotropic electron liquid. For most other metals, a more detailed model is necessary, which takes into account the anisotropy of the Fermi surface. In [5 - 7] there were considered effects of the influence of the anisotropy of the Fermi surface on the distribution of the spin waves in the electron liquid. In these investigations [1, 2, 5 - 7], no account was taken of the difference between the magnetic moment of the electron in the crystal and the moment for the free electron, due to the weak spin-orbit interaction between the conduction-electron spin and the orbital-motion moments of the electrons of the ion lattice. On the other hand, the difference connected with the interaction between the conduction electrons was taken into account. It is known, however, that in many cases the correction to the g-factor due to the spin-orbit interaction can be appreciable and observed in experiments [8]. In particular, for metals with an anisotropic Fermi surface, this correction depends on the orientation of the magnetic field $\vec{B}$ relative to the crystal axes. In the present communication we consider alkali metals of two types: metals in which the Fermi surface lies completely within the first Brillouin zone and is sufficiently far from its faces (e.g., potassium and sodium), and metals in which the Fermi surface passes close (in particular, as tangent) to two opposite faces (e.g., cesium under slight compression [9]). In the calculation of the anisotropy of the g-factor of the conduction electrons, connected with the presence of a weak periodic pseudopotential $W(\vec{r})$, we follow the method developed in [10]. The starting point will be the previously obtained [11] expression for the shift of the g-factor of almost-free electrons in the Bloch state $\vec{k}$, which is true for alkali metals:

$$\Delta g = \langle \Delta g(k) \rangle_{av} = -\langle \mu k^2 \sin^2\theta \rangle_{av} . \qquad (1)$$