The results of the calculation were quite unexpected. At $Q < 1$, the decrease of the force constants fully offsets the change of the mass, and no local oscillations arise. To the contrary, for a heavy impurity, owing to the increase of the force constants, local oscillations do occur, but only with low frequencies. There can exist altogether two local oscillations, the first of which (odd) occurs at $Q \gtrsim 1.3$, and the second (even) at $Q > 9$. As $Q \to \infty$, the local frequencies approach from below a common limit equal to $\sim 1.0059\omega_M$ (at finite Q, the even oscillations has the lower frequency). Thus, only minute local oscillations can exist in the model under consideration, and the condition for their occurrence is the opposite of the case of a harmonic crystal.

The results are determined, naturally, by the chosen value of the parameter ρ characterizing the degree of anharmonicity. The assumed value of ρ corresponds to strong anharmonicity. With increasing ρ, at a fixed value of ω_M, the degree of anharmonicity decreases, and the local oscillations acquire the same properties as in a harmonic lattice.

Thus, in spite of the simplified character of the investigated model, it can be stated that the dynamic properties of the defects can undergo considerable changes in strongly anharmonic crystals, and an experimental study of these properties would be of definite interest.

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COULOMB EXCITATION OF NUCLEI BY HEAVY POLARIZED PARTICLES

A.V. Davydov and I.M. Narodetskii
Institute of Theoretical and Experimental Physics
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The existing methods for determining the magnetic moments (m.m.) of short-lived nuclear states are based on the study of the interaction of the m.m. with the magnetic field [1]. The spin precession under the influence of the m.m. interaction with the magnetic field leads either to a perturbation of the $\gamma\gamma$ correlation[1] or to a Zeeman splitting of the levels of the excited nucleus. The latter is observed in a number of cases with the aid of the Mossbauer effect.

[1] Resonant scattering of a γ quantum is also possible in this case.

A common shortcoming of these methods is that they apply only to nuclei with sufficiently long lifetimes. Indeed, the measured effect is proportional in any case to $\omega_L \tau$, where $\omega_L = (\mu^*/J_f)H$ is the Larmor precession frequency[2], and μ^*, J_f, and τ are respectively the m.m., spin, and lifetime of the excited level. Modern laboratory fields, including internal fields acting on nuclei in ferro- and paramagnets, do not exceed in order of magnitude $H \sim 10^5 - 10^7$ Oe, so that the m.m. can be measured only for levels with $\tau \gtrsim 10^{-11}$ sec. In addition, to determine the m.m., it is necessary to know the lifetime of the nuclear level, the measurement of which is frequently difficult.

We consider in this article a possible method of determining the m.m., based on Coulomb excitation of nuclei by polarized particles. We use in this case a charged-particle magnetic field greatly exceeding the laboratory fields, and the entire process occurs during the time of flight of the particle past the nucleus, i.e., within the nuclear time 10^{-22} sec. The effect in question is therefore independent of the lifetime of the nucleus in the excited state.

The magnetic field is determined, first, by the orbital motion, $H_{orb} \sim Z_p ev/a^2$, where Z_p is the particle charge in e units, v is the particle velocity, and 2a is the shortest-approach distance in frontal collisions. Second, it is determined by the particle magnetic moment, $H_{m.m.} \sim \mu_p/a^3$. In experiments on Coulomb excitation, these two fields are comparable in magnitude. It must be borne in mind, however, that the contribution of the M1 transitions in nonspherical nuclei is quite small and amounts to $\lesssim 10^{-2}$ of the E2 transitions [2]. The experimental errors in the determination of the quadrupole moments amounts to $\sim$20 - 50%. Therefore the contribution of H_{orb} is practically impossible to separate by studying the cross section of the Coulomb excitation or the angular correlation of the γ quanta. The $H_{m.m.}$ contribution is easier to separate, since it reverses sign when the polarization of the incident particle is changed. The difference $\sigma_\uparrow - \sigma_\downarrow$ where $\sigma_\uparrow(\sigma_\downarrow)$ is the cross section of the Coulomb excitation for a spin directed parallel (antiparallel) to the normal to the reaction plane $\vec{n} = [\vec{p}_f \times \vec{p}_i]/|\vec{p}_f \times \vec{p}_i|$ ($\vec{p}_f$ and $\vec{p}_i$ are unit vectors along the final and initial particle momenta) is determined by the m.m. of the excited state. The order of magnitude of the effect is $(\alpha/vI_f)(\lambda\!\!\!^{-}_p/a)^2$, where $\alpha = 1/137$ and $\lambda\!\!\!^{-}_p$ is the Compton wavelength of the particle. We present below quantitative estimates on the basis of the classical theory [3] of Coulomb excitation[3].

Let us consider the simplest situation, when the off-diagonal transitions in the nucleus are pure E2 transitions, and the diagonal ones can be either E2 or M1. It is easy to see that the difference $\sigma_\uparrow - \sigma_\downarrow$ is determined in this case only by one known parameter, the diagonal matrix element of the M1 transition in the final state of the nucleus $\langle J_f \| M1 \| J_f \rangle$.[4] By measuring this difference (or the left-right asymmetry in scattering) we can determine experimentally the m.m. of the excited state. We parametrize the matrix element in

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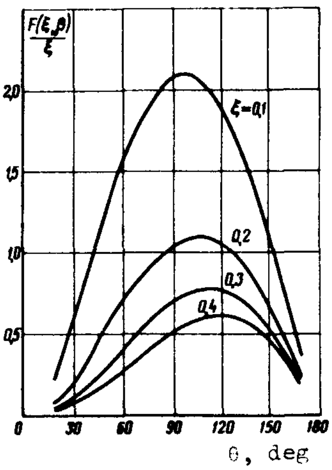
[2])Here and throughout $\hbar = c = 1$.

[3])Since we aim at obtaining estimates only, we neglect the energy lost by the incident particle and assume that the direction of the spin of this particle remains the same during the excitation process.

[4])In the general case of E2 and M1 nondiagonal transitions, the expression (1) for the asymmetry contains corrections necessitated by the magnetic moment of the ground state and the matric elements of the nondiagonal M1 transitions.

in the following manner:[5] $\langle J_f \| M1 \| J_f \rangle = \sqrt{45/8\pi}(e\mu^* \lambda\!\!\!{}^-)/2$, where $\lambda\!\!\!{}^- = 1/m$ is the Compton wavelength of the proton. We also put $\mu_p = e\mu(\lambda\!\!\!{}^-/2)$.



Using standard methods for the transitions $J_i = 0^+ \to J_f = 2^+$, we obtain accurate to terms $\sim\alpha$,

$$A = \frac{\sigma_\uparrow - \sigma_\downarrow}{1/2(\sigma_\uparrow + \sigma_\downarrow)} - K\,\frac{f(\theta,\xi)}{\xi},$$

$$K = \frac{1}{\pi}\sqrt{\frac{15}{8}}\,\mu\mu^*\,\frac{A_p(1+A_p/A_{nuc})^{-1}}{Z_p Z_{nuc}}\,\frac{\omega}{m}, \qquad (1)$$

where θ is the scattering angle, $\xi = (\alpha Z_p Z_{nuc}\omega/v^3)/A_p(1 + A_p A_{nuc})^{-1}$; Z_p, A_p, Z_{nuc}, and A_{nuc} are the charges and masses (in units of the nucleon mass m) of the particle and of the target nucleus, respectively, $\omega = E_f - E_i$ is the excitation energy, and the function $f(\theta, \xi)$ is expressed in terms of the classical orbital integrals $J_{\lambda\mu}(\theta, \xi)$ which are tabulated in [4]:

$$f(\theta,\xi) = -\frac{1}{\sqrt{30}}\left\{\sum_{\mu=\pm 2}\mu Y_{2\mu}^2\left(\frac{\pi}{2},0\right)J_{2\mu}(\theta,\xi)\,P \times\right.$$

$$\left.\times\int_{-\infty}^{\infty}\frac{J_{2\mu}(\theta,\xi+\xi')J_{20}(\theta,-\xi')\,d\xi'}{\xi'}\right\}\left\{\sum_{\mu=0,\pm 2}Y_{2\mu}^2\left(\frac{\pi}{2},0\right)J_{2\mu}^2(\theta,\xi)\right\}^{-1}. \qquad (2)$$

In the derivation of (1) and (2) we used an expression for the probability for the magnetic-transition probability (see [4], formula (2.A.42)), where it is necessary to put

$$S_{M\lambda,\mu} = \left[-\,v\,\mathrm{ctg}\frac{\theta}{2} + (\vec{\mu}_p\,\mathbf{n})\frac{\lambda}{Z_p e a}\right]\frac{1}{va\,\lambda+1}\,Y_{\lambda+1,\mu}\left(\frac{\pi}{2},0\right)\times$$

$$\times J_{\lambda+1,\mu}(\theta,\xi). \qquad (3)$$

In (3), the first term in the square brackets describes the contribution of the orbital motion of the particle [5], and the second describes the contribution of its magnetic moment. Expression (3) is valid in the case when the m.m. of the particle is parallel (antiparallel) to the normal to the reaction plane, and in this case $\vec{\mu}_p \cdot \vec{n} = \pm\mu_p$. The (-) sign in (1) is obvious beforehand, for when $\vec{\mu}_p \cdot \vec{n} = +\vec{\mu}_p$ $(-\vec{\mu}_p)$ the magnetic field at the nucleus, produced by the m.m. of the particle, is antiparallel (parallel) to the magnetic field due to its orbital motion.

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[5] In accordance with the definition:

$$\frac{e}{2m}\mu^* = \left(\frac{4\pi}{3}\right)^{1/2}\frac{(I_f I_f|0|I_f I_f)}{\sqrt{2I_f+1}} \langle I_f \| M1 \| J_f \rangle .$$

The function $f(\theta, \xi)$ has a maximum at $\theta = 90 - 120°$ and decreases monotonically with increasing ξ. Plots of the function $f(\theta, \xi)$ for $\xi = 0.1 - 0.4$ are shown in the figure. It follows from (1) that the magnitude of the effect under consideration is determined mainly by the charge of the nucleus, $K \sim 1/Z_{nuc}$. Since $\omega/m \sim 10^{-3}$ in experiments on Coulomb excitation, it follows that the maximum value of the effect is $A_{max} \sim 0.5 \times 10^{-4}$ for nuclei in the region of zinc.

The Coulomb excitation cross section $d\sigma_{Coul}/d\theta$, other conditions being equal, increases with increasing particle energy E_p and with decreasing charge of the nucleus. At large E_p and small Z_p, however, the particles can penetrate inside the nucleus, and the interpretation of the results is greatly complicated by the need for taking into account effects of the strong interaction, in which the asymmetry can reach several times ten per cent. Estimating the cross section for the production of the compound nucleus in accordance with the formula $D\sigma_{nuc}/d\theta \sim PR^2/4$, where P is the penetrability of the Coulomb barrier and $R = r_0 A_{nuc}^{1/3} + R_p$ ($r_0 = 1.4$ F), we can obtain lower bounds for the quantities $Z_p Z_{nuc}$ and $Z_p Z_{nuc}/E_p$ in the experiment. It must be borne in mind, however, that these limitations can be highly overestimated, because decay of the compound nucleus via the inelastic scattering channels is less probable than decay via all other channels.

Particle	Nucleus	ω, MeV	E_p, MeV	ξ	A/μ^*	$\frac{d\sigma_{Coul}}{d\Omega}$ cm^2/sr	$\frac{d\sigma_{nuc}}{d\Omega}$ cm^2/sr
$^7_3$Li	$^{148}_{62}$Sm	0,551	15	0,4	$2 \cdot 10^{-5}$	$0,8 \cdot 10^{-27}$	$0,7 \cdot 10^{-33}$
$^{17}_8$O	$^{64}_{30}$Zn	0,980	25	0,6	$3,7 \cdot 10^{-5}$	$5,6 \cdot 10^{-27}$	$2,5 \cdot 10^{-33}$

The table lists the values of the asymmetry A for the case of polarized $^7_3$Li and $^{17}_8$O nuclei for scattering through an angle corresponding to the maximum value of $f(\theta, \xi)/\xi$, and also the cross section for the Coulomb excitation, $d\sigma_{Coul}(E2)/d\Omega$ and the cross section $d\sigma_{nuc}/d\Omega$ for production of a compound nucleus.

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