*We observe that formulas similar to (1) and (2) obtained by Sokolov[3] are erroneous, in consequence of the fact that the components of the expression curl $\mu \mathbf{H} - \mu$ curl $\mathbf{H}$ are quantities of the first order in M'.

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[2] G. S. Krinchik, Izv. Akad. Nauk SSSR, Ser. Fiz. **21**, 1293 (1957), Columbia Tech. Transl. p. 1279. G. S. Krinchik and R. D. Nuralieva, J. Exptl. Theoret. Phys. (U.S.S.R.) **36**, 1022 (1959), Soviet Phys. JETP **9**, 724 (1959).

[3] A. V. Sokolov, Физика металлов и металловедение (Physics of Metals and Metal Research) **3**, 210 (1956).

[4] G. S. Krinchik, Физика металлов и металловедение (Physics of Metals and Metal Research) **7**, 181 (1959).

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Translated by W. F. Brown, Jr.
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ON THE OVERHAUSER EFFECT IN SATURATION OF FORBIDDEN RESONANCE

L. L. BUISHVILI

Tbilisi State University

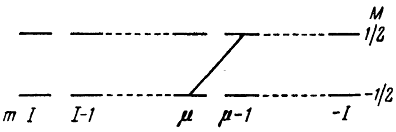
Submitted to JETP editor December 30, 1958

J. Exptl. Theoret. Phys. (U.S.S.R.) **36**, 1926-1927 (June, 1959)

KHUTSISHVILI[1] has considered the stationary Overhauser effect in saturation of an allowed transition in paramagnetic-resonance spectrum.

In the present paper we consider the stationary Overhauser effect, but in saturation of a forbidden transition in the paramagnetic-resonance spectrum (we note that Jeffries[2] has considered dynamic polarization of the nuclei, obtained in saturation of a forbidden transition, but for the case when the relaxation time of the nuclei is considerably longer than the relaxation time of the electrons).

Let us consider a system consisting of an electron shell with an effective spin S of one-half and a nucleus with spin I placed in an external magnetic field H. Considering the external field to be sufficiently strong, we neglect (in the calculation of the level population) the energies of the spin-spin interaction and the Zeeman energy of the nucleus. In such an approximation, we obtain 2I + 1 pairs of levels, the difference in the energies of the components of each pair being $g\beta H$. The level scheme is shown in the diagram (M and m are the projections of the spins of the electron and nucleus on the external field).



In the case of axial symmetry of the intra-crystalline electric field and in the case of an external field H parallel to the symmetry axis (the z axis), we obtain transitions that satisfy the selection rules $\Delta M = -\Delta m = \pm 1$ if the alternating field is parallel to H. For other directions of H relative to z we obtain also other forbidden transitions, in particular, transitions that satisfy the selection rule $\Delta M = \Delta m = \pm 1$.

We assume henceforth that only vertical relaxation (transitions $\Delta M = \pm 1$, $\Delta m = 0$) and relaxation due to the hyperfine interaction (transitions $\Delta M = -\Delta M = \pm 1$) are present.

For brevity we denote by μ the state corresponding to $M = -\frac{1}{2}$, $m = \mu$, and by μ' the state with $M = \frac{1}{2}$, $m = \mu$. We can write for relaxation transitions

$$W(\mu, \mu') = We^{-\delta}, \quad W(\mu', \mu) = We^{\delta},$$

$$W(\mu, \mu - 1') = \lambda We^{-\delta}, \quad W(\mu - 1', \mu) = \lambda We^{\delta},$$

where $2\delta = g\beta H/kT$, W is a certain function of the temperature and of the external field, and λ is a function of T, H, and μ.

Let the forbidden resonance $\mu \rightleftarrows \mu - 1'$ be saturated. We denote by $W(\mu)$ the probability of this transition, caused by an alternating field, per unit time. We introduce a resonance saturation parameter $s(\mu)$ in accordance with the formula

$$N(\mu) - N(\mu - 1') = \frac{N}{2I+1} \tanh \delta \,[1 - s(\mu)].$$

We can obtain the following expression for $s(\mu)$ and the parameters that characterize the degree of orientation of the nuclei:

$$s(\mu) = \frac{W(\mu)\,[(I + 1 - \mu)\,e^{-\delta} + (I + \mu)\,e^{\delta}]}{W(\mu)\,[(I + 1 - \mu)\,e^{-\delta} + (I + \mu)\,e^{\delta}] + \lambda W(2I + 1)},$$

$$f_1(\mu) = -s(\mu)\,\frac{[I + (I + 1) - \mu(\mu - 1)]\sinh \delta}{I\,[(I + 1 - \mu)\,e^{-\delta} + (I + \mu)\,e^{\delta}]},$$

$$f_2(\mu) = -s(\mu)\,\frac{(2\mu - 1)\,[I(I + 1) - \mu(\mu - 1)]\sinh \delta}{I(2I - 1)\,[(I + 1 - \mu)\,e^{-\delta} + (I + \mu)\,e^{\delta}]}.$$

We also calculated the parameters f_1 and f_2 in saturation of one of the transitions, satisfying the selection rules $\Delta M = \Delta m = \pm 1$. However, these expressions are too unwieldy to repeat here.

In conclusion, the author thanks G. R. Khutsishvili for help in this work.

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[2] C. D. Jeffries, Phys. Rev. **106**, 164 (1957).

Translated by J. G. Adashko
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ON THE "LARMORON" PLASMA THEORY

S. A. KAPLAN

L'vov State University

Submitted to JETP editor January 30, 1959

J. Exptl. Theoret. Phys. (U.S.S.R.) **36**, 1927-1928 (June, 1959)

In present-day quantum theory widespread use is made of the method introducing quasi-particles; it usually leads to a more effective and visualizable theory. It seems probable that a similar introduction of effective particles would also appreciably simplify plasma theory (albeit classical).

We shall use the term "larmoron" to define an effective particle situated at the guiding center of the Larmor motion of the real particle with a magnetic moment $\mu = mv_\perp^2/2H$ (where m is the particle mass, $v_\perp$ the component of the velocity perpendicular to magnetic field H) and an energy equal to the total energy of the real particle (we think this nomenclature more felicitous than the expression "Larmor circle" used elsewhere). The idea of larmorons has already been used often (but not always consistently) in many papers (see, for instance, references 1, 2, and others).

For a more consistent definition of a larmoron one must introduce its mean lifetime τ and the average translational velocities of its motion (u and v are the components perpendicular to the field H, and w the one along H). If we assume that the magnetic field is uniform, the velocity of the motion of the real particle under the action of an external acceleration a (along the x axis) in a field H (along the z axis) is of the form

$$v_x = v_x^{(0)} \cos \omega t - (v_y^{(0)} - a/\omega) \sin \omega t, \quad v_z = v_z^{(0)},$$
$$v_y = v_x^{(0)} \sin \omega t + (v_y^{(0)} - a/\omega) \cos \omega t + a/\omega, \tag{1}$$

where $\mathbf{v}^0$ is the velocity of the real particle at time $t = 0$; ω is the Larmor frequency. Since the probability that a larmoron will live through the interval of time t to t + dt is equal to $e^{-t/\tau} dt/\tau$, we have

$$u = \int_0^\infty v_x e^{-t/\tau} \frac{dt}{\tau} = \frac{v_x^{(0)} - v_y^{(0)}\omega\tau + a\tau}{1 + (\omega\tau)^2},$$
$$v = \frac{v_y^{(0)} + v_x^{(0)}\omega\tau + a\omega\tau^2}{1 + (\omega\tau)^2}, \qquad w = v_z^{(0)}. \tag{2}$$

The total larmoron energy ϵ can be expressed in terms of its velocity as follows:

$$\varepsilon = \frac{m}{2}\{v_x^{(0)2} + v_y^{(0)2} + v_z^{(0)2}\} = \frac{m}{2}\{[1 + (\omega\tau)^2](u^2 + v^2)$$
$$- 2a\tau[u + (\omega\tau)v] + (a\tau)^2 + w^2\}. \tag{3}$$

If we assume that the distribution function for the velocities $\mathbf{v}^0$ had a Maxwellian form the distribution function for the velocities u, v, w is of the form

$$f_0(u, v, w) = n(m/2\pi kT)^{3/2}[1 + (\omega\tau)^2]$$
$$\times \exp\left\{-\frac{m}{2kT}\left[\left(u\sqrt{1+(\omega\tau)^2} - \frac{a\tau}{\sqrt{1+(\omega\tau)^2}}\right)^2\right.\right.$$
$$\left.\left. + \left(v\sqrt{1+(\omega\tau)^2} - \frac{a\omega\tau^2}{\sqrt{1+(\omega\tau)^2}}\right)^2 + w^2\right]\right\}, \tag{4}$$

where n is the number of larmorons per unit volume. The transport equation for the larmorons of one kind under the conditions where there is a density and a temperature gradient along the x axis can be written in the form $u\,\partial f/\partial x = (\partial f/\partial t)_{\text{coll}}$, since the larmoron acceleration is equal to zero (if $\mathbf{a} \perp \mathbf{H}$). If we use the usual method of Lorentz's electron theory, i.e., put $f = f_0 + u\chi$ we get for the transport equation and its solution

$$\partial f_0/\partial x = -\chi/\tau, \qquad f = f_0 - \tau u\,\partial f_0/\partial x. \tag{5}$$

Substituting (4) into (5) and then evaluating, by the usual equations, the fluxes of electric charge, number of particles, and energy, we easily obtain the well known formulae for the coefficients of electrical conductivity, diffusion, and thermal conductivity of a plasma. We must note that the coefficient of electrical conductivity is evaluated with the first term of (5), i.e., with f_0 and with $a = eE/m$.

All calculations in terms of the "larmoron" theory of a plasma are very simple and suffer not at all from the unwieldiness which is so characteristic for the transport theory of a plasma. The weak point of the larmoron theory, in the formu-