certain (identical) very simple objects, or (b) the leptons and quarks are components of a single multiplet that realizes an elementary representation of the group G. Case (a) deserves a special treatment and will not be discussed here. Consider the possibility (b). If the leptons and quarks are combined in a single multiplet $\psi = (\ell, q, \ldots)$ (the dots imply the possible existence of other members of this multiplet), then the question arises of the extent to which the particles that enter initially perfectly on a par can differ ultimately so strongly in their properties. Within the framework of Weinberg's general approach, two different schemes can be advanced by way of an answer to this question.

1. We postulate the existence of short-range mass-dependent forces. This is done by introducing (in a gauge-invariant manner) massive fields with spin 2, realizing a nontrivial representation of the group G. Then the quarks, which acquire a large mass as a result of symmetry breaking, turn out to be strongly-interacting particles, whereas the leptons interact weakly. Such a theory will be renormalizable in the usual sense.

2. Another possibility that preserves the renormalizability of the theory is that there exist gluon vector and/or scalar (neutral) fields that interact strongly ($g \sim 1$) both with quarks and with leptons. On the other hand, if we assume that these fields have (or acquire) a large mass M (on the order of the quark mass) then, in spite of their strong interaction with the leptons, the latter will not form bound states, by virtue of the inequality $m_\ell \ll M$ (i.e., the effective radius $1/M$ is much less than the Compton wavelength $1/m_\ell$). At such energies, these forces reduce to a four-fermion interaction between the leptons and the neutral currents with a coupling constant $\sim g^2/M^2$. Experimental estimates of the corresponding constants [5] are $G_{h\ell} \lesssim 0.3G_F$ and $G_{\ell\ell} \lesssim 0.6G_F$. This yields for the gluon mass an estimate $M \gg 580$ mp, which is simultaneously also an estimate for the quark mass. An isolated proton remains stable, although the baryon charge, depending on the details of the model, will be conserved either rigorously or in modulo 2. The corresponding transition probability can be made small enough not to contradict the experiment.

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DYNAMIC SELF-POLARIZATION OF NUCLEI IN SOLIDS

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We consider here a system of nuclear spins interacting with electrons, which are artificially maintained in a disordered spin state. We shall show that under these conditions, at sufficiently low temperature, the disordered state of the nuclear spins is unstable (even in the absence of an external magnetic field). On the other hand, the state with practically complete ionization of the nuclei turns out to be stable. In this state, the polarized nuclei produce an effective magnetic field that acts on the spins of the electrons and is due to the hyperfine interaction. The nuclear polarization itself is maintained by the Overhauser effect in this effective field.

To understand the cause of the instability of the disordered state of the nuclear spins, let us assume that a fluctuation has occurred, consisting of the onset of a certain small polarization of the nuclei. The electron spins are situated in the magnetic field due to the hyperfine interaction, and the Overhauser effect leads to a further increase of the nuclear polarization. Such an instability (self-polarization of the nuclei) occurs only at a sufficiently low temperature, when the Overhauser value of the nuclear polarization in a magnetic field is higher than the fluctuation polarization producing this field.

The critical temperature T_c is determined by the spin energy of the electron in the effective magnetic field of the fully polarized nuclei, i.e., the hyperfine-interaction energy. The effective field of the heavy nuclei can reach several dozen kOe, and accordingly T_c can be of the order of several degrees K.

Let us consider for simplicity a metallic lattice of like nuclei with spin I = 1/2 interacting with non-oriented electrons. We denote by $N_\pm$ the number of the nuclei with spin "up" or "down." The change of N_+ is given by

$$\frac{dN_+}{dt} = - W_1 n_- N_+ + W_2 n_+ N_- , \tag{1}$$

where $n_\pm$ is the number of electrons with spin "up" or "down," and W_1 and W_2 are the probabilities of the transitions $(-+) \to (+-)$ and $(+-) \to (-+)$, respectively (the first and second signs in the parentheses correspond to the states of the electron and the nucleus, respectively). Putting $n_+ = n_- = n$ and introducing the degree of nuclear polarization $P = (N_+ - N_-)/(N_+ + N_-)$ we obtain for it from (1) the equation

$$\frac{dP}{dt} = - \gamma(P - P_0), \qquad P_0 = \mathrm{th}\left(\frac{\epsilon}{2kT}\right), \tag{2}$$

where $\gamma = (W_1 + W_2)n$ and we have used the thermodynamic ratio of the transition probabilities $W_2/W_1 = \exp(\varepsilon/kT)$. Here k is Boltzmann's constant, T the temperature, and ε the difference between the energies of the states (+-) and (-+).

Equation (3) describes the tendency of the nuclear polarization to the Overhauser value P_-. Usually ε takes into account only the electron energy in the external magnetic field. In the presence of nuclear polarization, however, it is necessary to allow also for the energy of the electron in the effective magnetic field of these nuclei[1]. Thus, in the absence of an external magnetic field we have $\varepsilon = AP/2$, where A is the hyperfine splitting factor [1]:

$$A = \frac{16\pi}{3I} \mu_B \mu_I \, |\psi(0)|^2 .$$

Here μ_B is the Bohr magneton, μ_I the magnetic moment of the nucleus, $|\psi(0)|^2$ the value of the square of the wave function of the electron (normalized in the unit-cell volume) at the location of the nucleus; in the case under consideration I = 1/2.

[1]We assume that the electron localization region contains a large number of nuclei.

According to (2), the stationary value of the nuclear polarization in the absence of an external field should satisfy the equation

$$P = \operatorname{th}\left(\frac{T_c}{T} P\right), \tag{3}$$

where

$$T_c = A/4k. \tag{4}$$

At $T > T_c$, Eq. (3) has only one solution, $P = 0$. When $T < T_c$, however, there are three solutions (see Fig. 1): $P = 0$ and $P = \pm P_s$. From (2) we obtain the condition for the stability of the stationary solutions, namely $dP_0/dP < 1$. It is easy to see that when $T < T_c$ the states with $P = \pm P_s$ are stable and the state with $P = 0$ is unstable. Figure 2 shows the temperature dependence of the polarization in the stable state. This dependence has a form typical of a second-order phase transition.

The foregoing analysis can be easily generalized to include the case when the unit cell contains several nuclei with arbitrary spins. We then obtain in place of (4) the following expression for T_c:

$$T_c = \frac{1}{3k} \sum_i A_i I_i (I_i + 1), \tag{5}$$

where the summation is over all the nuclei in the unit cell. At temperatures below T_c the degree of polarization of all types of nuclei tends to unity.

By way of an example, we calculate the critical temperature for indium antimonide. The values of $|\psi(0)|^2$ in this material were measured by Gueron [2] (see also [3]). Using these data, we get $T_c = 9$°K.

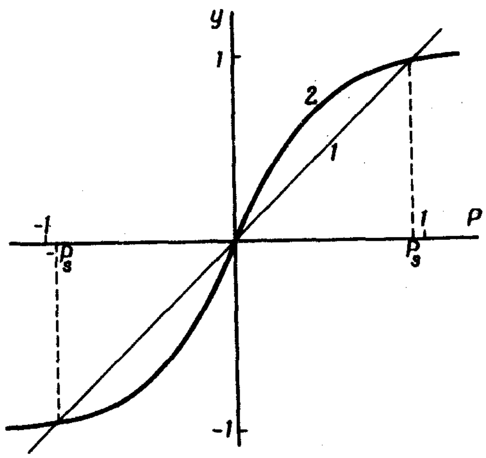

Fig. 1

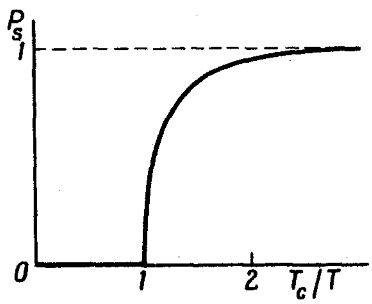

Fig. 2

Fig. 1. Graphic solution of Eq. (3) at $T = T_c/2$: 1) $y = P$, 2) $y = \tanh(T_c P/T)$.

Fig. 2. Temperature dependence of the degree of polarization of the nuclei in the stable state.

The nuclei become self-polarized because the electrons are artificially kept in a non-oriented state. Indeed, we have assumed that the fluctuation of the nuclear polarization and the appearance of an effective magnetic field are not accompanied by a corresponding change of the electron orientation. Such conditions can be created in practice, for example, in semiconductors when non-equilibrium non-oriented carriers are produced by light, or by injection of carriers through a junction, if their lifetime τ is much shorter than the spin-relaxation time τ_s. Otherwise the critical temperature becomes lower than the value given in (5), owing to the additional factor $\tau_s/(\tau + \tau_s)$. Another effect that lowers T_c may be the spin-lattice relaxation of the nuclei. The corresponding decrease of T_c is described by the factor $\gamma/(\nu + \gamma)$, (leakage factor) where ν is the reciprocal time of longitudinal nuclear relaxation.

In the absence of an external magnetic field, the direction of the macroscopic nuclear magnetization at $T < T_c$ is random (in the simple model under consideration). It can be shown that a weak external field $H << kT_c/(g\mu_B)$, without changing the general picture of the phenomenon, makes the orientation of the macroscopic magnetization of the nuclei along the magnetic field the most probable.

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SINGULARITY OF THE BEHAVIOR OF A SUPERCONDUCTING FILM IN AN ALTERNATING FIELD OF FREQUENCY CLOSE TO THRESHOLD

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We consider the behavior of a superconducting film in an external microwave field whose amplitude in a real experimental situation is small compared with the critical value. Let also the frequency of this field be close to the threshold $2\Delta_0$, where Δ_0 is the equilibrium value of the gap. The linear properties of the superconductors at such frequencies have been thoroughly investigated [1], and we are therefore interested in nonlinear effects.

To this end, we obtain the response of the ordering parameter $\Delta(t)$ which is a small increment to the equilibrium value Δ_0 to an external alternating field. It is expressed in terms of the Green's function

$$\Delta_\omega = \lambda \int \frac{d\epsilon}{4\pi i} \int \frac{d^3p}{(2\pi)^3} F_{\epsilon\epsilon-\omega}. \tag{1}$$

Expanding the function $F_{\epsilon\epsilon-\omega}$ in terms of Δ_ω and the alternating-field potentials $\vec{A}_\omega$, we rewrite (1) in an integral form

$$\int K(\omega\omega_1)\,\Delta\omega_1\, d\omega_1 = \int a(\omega\omega_1)\,\mathbf{A}_{\omega_1}\mathbf{A}_{\omega-\omega_1}\, d\omega_1 + \dots \tag{2}$$

the right-hand side of which is a series in powers of the field intensity, and K and a are certain kernels. Without allowance for the dependence of K on the