It would be most interesting to realize a condition $\omega = 2gH_0 + gH_1 = 2g\beta M_0 + gH_1$, at which the frequency of the wave that stabilizes the "overturned" state is resonant for ordinary orientation of the vector $\vec{M}$ along the field $\vec{H}$. In this case, while the amplitude of the wave h_0 decreases at the "input", it increases at a certain instant of time very sharply "at the output."

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SECOND-ORDER PHASE TRANSITIONS WITHOUT DIVERGENCES IN THE SECOND DERIVATIVES OF THE THERMODYNAMIC POTENTIAL

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It is known that the Landau theory of second-order phase transitions does not hold in the immediate vicinity of the transition point, owing to the strong increase of the long-wave fluctuations of the characteristic parameters of the transition [1, 2]. On the other hand, if the fluctuations are accompanied by the appearance of long-range fields (electric, magnetic, elastic), then their character changes significantly [3]. As a result, not all fluctuation waves increase in sufficiently anisotropic crystals near the transition point, but only waves with definite wave-vector directions. This leads to a strong decrease of the phase volume connected with the large fluctuations. An example of this type are uniaxial ferroelectrics [4, 6], in which, owing to the occurrence of the electric field, only fluctuations with wave vectors lying in the plane perpendicular to the ferroelectric axis increase. Therefore, for example, the specific heat increases on approaching the transition point in accordance with a law that differs from that in the isotropic substance.

In this article we consider another example, when the fluctuations of the characteristic parameters are even more suppressed. It turns out that if the transition parameter η and the elastic deformation u_{ik} are linearly connected in the symmetrical phase, then only fluctuations of the wave vectors parallel to definite crystallographic axes increase, i.e., the phase volume corresponding to large fluctuations decreases even more strongly than in the preceeding case. Consequently, at $T \to T_c$ (T_c is the transition temperature) the specific heat remains finite, and the thermodynamic potential can be expanded in a series in η in a manner similar to that in the Landau theory[1]. However, there is no complete agreement with the Landau theory. The singularity of the thermodynamic potential does not vanish completely at the transition point. Thus, all the derivatives of the specific heat with respect to the temperature diverge as $T \to T_c$. The same pertains, of course, to other derivatives of the thermodynamic potential above the second. It was shown in [7] that phase transitions in solids accompanied by an infinite increase of the specific heat cannot be of second order. We emphasize that the conclusions of [7] do not pertain to the transitions considered here, since the specific heat remains finite at the transition point.

Apparently, the only presently-known experimental example of the transitions of interest to us are ferroelectric transitions in crystals with piezzoeffect in the paraphase (Rochelle salt, KH_2PO_4, etc.). However, analysis shows that a linear connection between the transition parameter and the elastic deformations is possible also for certain non-ferroelectric phase transitions, including a number of transitions with change of the magnetic structure.

Let us explain now how the foregoing results are obtained. Since only the long-wave fluctuations of the characteristic parameter increase appreciably at the transition point, we can assume for the Hamiltonian of the system the continuous-medium approximation. We specify the Hamiltonian density in the form

$$n(\eta) = -\frac{a}{2}\eta^2 + b\eta^4 + c(\nabla\eta)^2 + d\eta u_{xy} + \frac{\lambda}{2} u_{ll}^2 + \mu u_{ik}^2 . \qquad (1)$$

We have neglected here the anisotropy of the elastic properties of the crystal, but its anisotropy was taken into account in the fourth term; x and y are therefore perfectly defined crystallographic directions. Starting from (1), we can obtain diagram expansions for the correlation functions and for the thermodynamic quantities, and verify that the divergences in the diagrams for the correlation functions, which usually appear in phase-transition theory, do not appear in our case. However, within the framework of the present paper, it is appropriate to calculate only the first corrections to the results of the Landau theory. It turns out that the qualitative results with respect to the temperature dependence of the thermodynamic quantities remain valid also when the higher-order approximations are taken into account. The corresponding analysis will be presented in a more detailed article. We use the method proposed in [8]. For the density of the free energy we employ an expression that differs from (1) only in that $-a$ is replaced by $\alpha = \alpha'_\theta(T - \theta)$, where θ is the transition temperature in the clamped crystal, i.e., at $u_{xy} = 0$. Actually, however, T_c is determined by the condition $\alpha - d^2/4\mu = 0$. We confine ourselves to the case of a symmetrical phase. It follows from the expression for the free energy that

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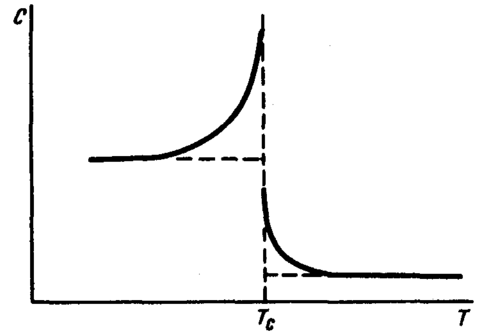
1)Before this article was sent to press, the authors have learned of a paper [9] in which it is stated, on the basis of the self-consistency criterion of the Landau theory, obtained in [2] and equivalent to the Ginzburg criterion [1], that this theory is not contradictory for ferroelectrics with the piezzoeffect in the paraphase (provided the crystal constants satisfy certain conditions). Such a statement agrees to a certain degree with one of our results. We emphasize, however, that the criteria of [1, 2] themselves do not suffice to draw definite conclusions concerning the temperature dependence of the thermodynamic quantities. It is necessary for this purpose to go beyond the limits of the first approximation of the Landau theory. The results of such investigations, reported in this paper, differ in a number of essential points from the results of [9]. The authors are grateful to D.I. Khomskii who pointed [9] out to them.

$$\langle \eta_k \eta_{-k} \rangle = \frac{T}{V\tilde{\alpha}(k)} = \frac{T}{V\left[a - \frac{d^2}{4\mu} + \frac{d^2}{4\mu}\left(\cos^2\theta + \frac{\lambda+\mu}{\lambda+2\mu}\sin^4\theta\sin^2 2\phi\right) + c k^2\right]}, \tag{2}$$

where η_k is the Fourier component of $\eta(x)$, $\langle\ \rangle$ is statistical averaging, and θ and ϕ are the angles of the spherical coordinate system with polar axis along the z axis. Confining ourselves to first-order corrections to the Landau theory theory, we obtain for the specific heat [8]:

$$C = C_0 + A \int \frac{dk}{\tilde{\alpha}^2(k)}, \tag{3}$$

where C_0 is the regular part of the specific heat, and A is a constant. It is easy to verify that the integral in (3) converges, and $\partial C/\partial T \sim |T - T_c|^{-1/2}$. The dependence of the specific heat on the temperature is shown in the figure. We recall that in the isotropic case, in the assumed approximation, we have $C \sim (T - T_c)^{-1/2}$. As shown in [5], the expression for the specific heat near T_c can be represented in the form of a series in powers of the integral contained in (3). At $T = T_c$, the terms of this series remain finite.



Temperature dependence of the specific heat. The specific heat at the transition point, calculated with allowance for the fluctuations, differs from that predicted by the Landau theory (dashed) by an amount of the order of $C_0\sqrt{T_c} \times (T_c/T_c - \theta)$.

We note, finally, that ferroelectrics with a piezzoeffect in the paraphase are at the same time uniaxial. Therefore the suppression of the fluctuations in them is due *not only* to elastic stresses, but also to the influence of the electric field. This circumstance, however, does not change the character of the temperature dependence of the thermodynamic potential, and was therefore not taken into account here.

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EXCHANGE INTERACTION OF AN ISOBAR WITH A NUCLEON, AND DYNAMIC MODEL OF TWO-BARYON RESONANCE

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This paper is devoted to multibaryon resonances, or nuclei in which one of the nucleons is replaced by a baryon resonance. The possible existence of such systems was suggested in [1, 2]. In [3], the general premises of the theory of resonance-nuclei were elucidated and the nuclei (ΔN) and (Δ2N) were considered (Δ - isobar with quantum number $T = 3/2$, $I^P = 3/2^+$, $m_\Delta = 1236$ MeV). On the basis of an analysis of the experimental data and phase-shift data, it was shown that the two-baryon resonance (ΔN) with isospin T = 1 makes a small contribution to the cross sections of the reactions $\pi d \to \pi NN$, $\pi d \to NN$, and $NN \to NN$. It was also established that induced non-exchange interaction of ΔN is not strong enough to form a bound state. In the present article we investigate the exchange forces due to the existence of the decay interaction $\Delta \to \pi N$ (see Fig. 1).

Our starting point is the Lippmann-Schwinger integral equation (formula (16) of [3]), the kernel of which is in this case the amplitude corresponding to the diagram (Fig. 1). If the kinetic energies of the baryons are small compared with the mass differences of the isobar and the nucleon $m_\Delta = 300$ MeV, then this amplitude can be regarded as dependent only on the momentum transfer $\vec{q}$, and the dependence of the isobar width Γ on the kinematic variables can be completely neglected. Then the aforementioned equation in the coordinate representation goes over into the Schrodinger equation with local energy-independent complex potential. The role of the energy in this equation is played by the quantity $E' = E + \Gamma/2$, where E is the total kinetic energy of the isobar and of the nucleon, and the operator of the potential is given by

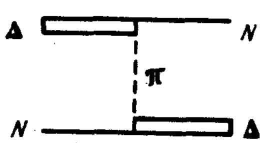

Fig. 1

$$\hat{V}(r) = \left(\frac{1}{2} + \frac{2}{3}\vec{\tau}_\Delta \vec{\tau}_N\right)\left\{-\left(\frac{1}{2} + \frac{2}{3} S\vec{\sigma}\right)\left(V_0(r) + \frac{5}{4}V_2(r)\right) + \right.$$
$$\left. + \left[2(\mathbf{S}\mathbf{n})^2 \mathbf{S}\vec{\sigma} + 2i(\mathbf{S}\mathbf{n})\mathbf{S}[\vec{\sigma}\times\mathbf{n}] + \frac{1}{2}(\mathbf{S}\mathbf{n})^2 - 3\mathbf{S}\mathbf{n}\vec{\sigma}\mathbf{n}\right] V_2(r)\right\}, \tag{1}$$

where S, $\vec{\sigma}$ and $\vec{\tau}_\Delta$, $\vec{\tau}_N$ are the spin and isospin operators of the isobar and of the nucleon, and

$$V_\Lambda(r) = -\frac{\lambda^2}{4\pi^2\mu^2}\int_0^\infty \frac{q^4\, dq}{q^2 + \mu^2 - (\Delta m - i\Gamma/2)^2}\, i^\Lambda j_\Lambda(qr)\, g^2(q^2). \tag{2}$$

Here μ is the pion mass, Γ = 120 MeV, and λ = 2. The form factor $g(q^2)$, which takes into account the departure of the pion and the nucleon from the mass shell at the vertex ΔπN, is chosen in the form $g(q^2) = (q_0^2 + c^2)/(q^2 + c^2)$, where $q_0 = 231$ MeV/c. From the data on the production of the Δ isobar in πN and NN collisions, it can be concluded (see, e.g., [4]), that the constant c