Magnetopolaron states of a spin-correlated antiferromagnet in the neighborhood of the spin-glass transition

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(Submitted September 4, 1997)
Fiz. Tverd. Tela (St. Petersburg) **40**, 310–314 (February 1998)

The problem of the spectrum of magnetopolaron states of a strongly correlated conducting canted antiferromagnet is solved. The approach used to study the spectrum is based on an atomic representation and a diagram technique for Hubbard operators. This approach makes it possible to include strong intra-ion interactions in a first-principles way, and to obtain the dispersion equation for the magnetopolaron spectrum for arbitrary values of the magnitude of the spin, temperature, and magnetic field. In the vicinity of the spin-flip transition an analytic expression is obtained for the spectrum of magnetopolaron states that goes beyond the framework of the quasiclassical approximation. © *1998 American Institute of Physics.*
[S1063-7834(98)03302-4]

In the course of their experimental studies of the de Haas-van Alphen effect in the heavy-fermion compound $CeCu_2Si_2$, Hunt *et al.*[1] observed an abrupt change in the frequency of de Haas-van Alphen oscillations as they scanned through the spin-flip point. In its normal phase, $CeCu_2Si_2$ has antiferromagnetic order with a low value of the Néel temperature. Therefore, in the range of strong magnetic fields and low temperatures, the de Haas-van Alphen effect will coexist with a spin-flip phase transition.

The role of antiferromagnetic order in the de Haas-van Alphen-effect was discussed in Refs. 2 and 3. However, these papers do not predict any change in the period of the oscillations. In Ref. 4 we used a simple model that included $s-d(f)$-exchange coupling between spin moments of the collectivized and localized electrons to demonstrate that the frequency of de Haas-van Alphen oscillations in a conducting antiferromagnet could be different on different sides of the spin-flip point. The results of Refs. 2–4 show that it is important to take into account both antiferromagnetism and the exchange coupling between collectivized and localized subsystems when attempting to understand experimental de Haas-van Alphen-oscillation data from first principles.

The latter is also important for strongly correlated systems having low concentrations of current carriers. Examples of systems that belong to this new class of strongly correlated systems are, e.g., the rare-earth pnictides (R) with the general formula RX, where X=Bi, Sb, As, P. The unusual nature of their galvanomagnetic and thermodynamic properties has led to the formulation of new concepts for constructing the ground state of these compensated RX semimetals, among them the concepts of magnetopolaron liquids and crystals.[5,6] Many-body effects play a key role in generating these scenarios for electronic structures as well as the presence of long-range antiferromagnetic order.

In order to describe the magnetopolaron states of the valence band, we can use the ideas developed in the papers by Nagaev.[7] Here, in keeping with the known peculiarities of the band structure of rare-earth monopnictides, we will assume that the width of the filled valence band is small compared with the characteristic $s-f$-interaction energy. In these narrow-band antiferromagnets, it is well known that a magnetopolaron narrowing of the band takes place.[7] This phenomenon has been well studied for the collinear geometry.

In this paper we find the spectrum of magnetopolaron states of a strongly correlated antiferromagnet under conditions of large canting of the magnetic sublattices. The use of Hubbard operators[8–10] and diagram techniques for them[11–16] allow us to derive the required dispersion equation for the spectrum for arbitrary values of temperature, magnetic field, and spin. In the low-temperature regime we obtain an analytic expression for the energies of magnetopolaron states in the neighborhood of the spin-flip transition.

1. HAMILTONIAN OF A STRONGLY CORRELATED NARROW-BAND ANTIFERROMAGNET

Let us consider a strongly correlated narrow-band antiferromagnet with hole-type current carriers. We will describe this system within the framework of the $s-d(f)$-exchange model.[17] The physics of the narrow-band variant of this model has been described in detail in Ref. 7. Strong single-site correlations are conveniently taken into account by introducing an atomic representation[8–10] and the diagram technique for the Hubbard operators.[11–16]

Let us write the Hamiltonian of the model in the form

$$H=H_h+H_{sd}+H_m. \tag{1}$$

Here the first term describes the band current carriers. In order to describe the antiferromagnetic phase we introduce two sublattices F and G. Then within the Wannier representation, and taking into account the Hubbard repulsion on a single site, we have

$$H_h=\sum_{ff'\sigma}[t_{ff'}-\delta_{ff'}(2\mu_B H\sigma-\mu)]c^+_{f\sigma}c_{f'\sigma}$$
$$+\sum_{fg\sigma}t_{fg}(c^+_{f\sigma}d_{g\sigma}+\text{h.c.})+\sum_{gg'\sigma}[t_{gg'}-\delta_{gg'}$$
$$\times(2\mu_B H\sigma-\mu)]d^+_{g\sigma}d_{g'\sigma}$$
$$+\sum_f Un_{f\uparrow}n_{f\downarrow}+\sum_g Un_{g\uparrow}n_{g\downarrow}, \quad (2)$$

where the labels f,f', and g,g' enumerate sites on the F- and G-sublattices respectively, and the operators $c_{f\sigma}(c^+_{f\sigma})$ and $d_{g\sigma}(d^+_{g\sigma})$ describe the processes of annihilation (creation) of Wannier-representation holes on the F- and G-sublattices.

The second term of the Hamiltonian (1) describes the coupling between band carriers and localized spin moments via the $s-d(f)$-exchange interaction[17]

$$H_{sd}=-A\sum_f(\mathbf{S}_f\boldsymbol{\sigma}_f)-A\sum_g(\mathbf{S}_g\boldsymbol{\sigma}_g), \quad (3)$$

where A is the $s-d(f)$-exchange parameter, $\mathbf{S}_f$ is a vector operator for a localized spin on site f, and $\boldsymbol{\sigma}_f$ is a vector operator for the spin moment of a hole in the Wannier representation for the F sublattice. An analogous definition can be made for the G sublattice.

The last term in Eq. (1) describes the Heisenberg interaction in the subsystem of localized spins, and also its Zeeman energy.

When a magnetic field is applied to an isotropic antiferromagnet, its sublattices become canted.[18,19] In order to describe this effect it is convenient to change to local coordinates so that for each sublattice the equilibrium magnetization vector is oriented along a new axis Oz. We presented a detailed derivation of the Hamiltonian in local coordinates in Ref. 4, along with rules for transforming the Fermi and spin operators. Using these rules, it is not difficult to write down the Hamiltonian H'. In this case

$$H'_h=\sum_{ff'\sigma}t_{ff'}c^+_{f\sigma}c_{f'\sigma}+\sum_{gg'\sigma}t_{gg'}d^+_{g\sigma}d_{g'\sigma}$$
$$+\sum_{fg\sigma}t_{fg}(\cos\theta c^+_{f\sigma}d_{g\sigma}+2\sigma\sin\theta c^+_{f\sigma}d_{g\sigma}+\text{h.c.})$$
$$+\sum_{f\sigma}\{\mu_B H\sin\theta c^+_{f\bar{\sigma}}c_{f\sigma}-(2\mu_B H\cos\theta\sigma-\mu)$$
$$\times c^+_{f\sigma}c_{f\sigma}\}+\sum_f Un_{f\uparrow}n_{f\downarrow}-\sum_{g\sigma}\{\mu_B H\sin\theta d^+_{g\sigma}d_{g\sigma}$$
$$+2(\mu_B H\cos\theta\sigma-\mu)d^+_{g\sigma}d_{g\sigma}\}+\sum_g Un_{g\uparrow}n_{g\downarrow}. \quad (4)$$

The Hamiltonian for the $s-d(f)$-exchange interaction remains unchanged: $H'_{sd}=H_{sd}$. For the subsystem of localized spin moments the transformed operator is given by the expression

$$H'_m=-\frac{1}{2}\sum_{ff'}I_{ff'}(\mathbf{S}_f\mathbf{S}_{f'})-\frac{1}{2}\sum_{gg'}I_{gg'}(\mathbf{S}_g\mathbf{S}_{g'})$$
$$-g\mu_B H\left\{\sum_f S^z_f(\theta)+\sum_g S^z_g(\theta)\right\}$$
$$+\sum_{fg}K_{fg}\{\cos 2\theta(S^x_fS^x_g+S^z_fS^z_g)+S^y_fS^y_g$$
$$+\sin 2\theta(S^z_fS^x_g-S^x_fS^z_g)\}. \quad (5)$$

It is clear from Eq. (4) that the noncollinear phase is characterized not only by a θ-dependent effective overlap integral but also by an additional operator term corresponding formally to a hole hopping from site to site, accompanied by a change in the projection of the spin moment.

2. THE NARROW-BAND ANTIFERROMAGNET IN ITS NONCOLLINEAR PHASE IN THE ATOMIC REPRESENTATION

In studying narrow-band antiferromagnets it is assumed that the following inequality holds between the $s-d(f)$-exchange constant and the overlap integral: $|t_{fg}|\ll|A|$. Therefore, the $s-d(f)$ coupling between localized and collectivized subsystems must be taken into account exactly.[7,12] To this end, we separate out terms from the total Hamiltonian H' that contain only single site operators:

$$H_{\text{ion}}=\sum_f\{-A(\mathbf{S}_f\boldsymbol{\sigma}_f)-\bar{H}S^z_f-h\sigma^z_f-h_\perp\sigma^x_f+Un_{f\uparrow}n_{f\downarrow}\}$$
$$+\sum_g\{-A(\mathbf{S}_g\boldsymbol{\sigma}_g)-\bar{H}S^z_g-h\sigma^z_g+h_\perp\sigma^x_g+Un_{g\uparrow}n_{g\downarrow}\}, \quad (6)$$

where the effective field $\bar{H}$ includes the self-consistent field:

$$\bar{H}=g\mu_B H\cos\theta+I_0R-K_0\cos 2\theta R,\quad R=\frac{1}{N}\sum_f\langle S^z_f\rangle. \quad (7)$$

Here I_0 and K_0 are Fourier transforms of the exchange parameters for zero quasimomentum. The longitudinal and transverse components of the magnetic field acting on a hole are determined by the expressions

$$h=2\mu_B H\cos\theta,\quad h_\perp=-2\mu_B H\sin\theta. \quad (8)$$

In what follows we will assume that the Hubbard repulsion is so strong that we can neglect states of two holes on a single site. Since in reality the inequalities $|A|\gg\bar{H}$, $|A|\gg h$, $|h_\perp|$ hold, the Schroedinger problem of finding eigenstates of the single-site operator in the noncollinear phase can be solved by following a simple kind of perturbation theory.

As is well known,[17] the $s-d(f)$-exchange operator $-A(\mathbf{S}\boldsymbol{\sigma})$ has two eigenstates with energies $E_a=-AS/2$ when the total moment $J=S+1/2$ and $E_b=A(S+1)/2$ when $J=S-1/2$. For narrow-band magnets we can limit our discussion to only those single-site states that correspond to the smaller energy.[7] For $A<0$, the states of interest are $(n=1,2,\ldots,2S)$

$$|n\rangle = \sqrt{\frac{2S-n+1}{2S+1}}\,|S-n+1,\downarrow\rangle - \sqrt{\frac{n}{2S+1}}\,|S-n,\uparrow\rangle. \tag{9}$$

If we treat the remaining terms of the single-ion Hamiltonian by perturbation theory, it is not difficult to construct a basis of single-ion states $|\psi_n\rangle$. In this case the energy spectrum of the one-hole states can be written in the form

$$E_n = \varepsilon_1 + \varepsilon(n-1) + \frac{2n-2S-1}{(2S+1)^2}\left(\frac{v^2}{\varepsilon}\right), \quad n=1, 2, \ldots, 2S,$$

$$\varepsilon_1 = \frac{A(S+1)}{2} - \overline{H}\left(S - \frac{1}{2S+1}\right) + \frac{h}{2}\frac{2S-1}{2S+1},$$

$$\varepsilon = \frac{2S+2}{2S+1}\,\overline{H} - \frac{h}{2S+1}, \quad v = \mu_B H \sin\theta. \tag{10}$$

Knowing the solution to the single-ion problem, it is not difficult to write down a representation for the Fermi operators $c_{f\sigma}$ and $c_{g\sigma}$ in terms of the Hubbard operators X_f^{α}.[11–14]

$$c_{f\sigma} = \sum_{nM} \langle M|c_{f\sigma}|\psi_n\rangle X_f^{Mn} \equiv \sum_{\alpha} \gamma_\sigma(\alpha) X_f^{\alpha}. \tag{11}$$

Here the label *alpha* identifies the transition of an ion from state $|\psi_n\rangle$ to a state $|M\rangle$ corresponding to the ion without a hole with spin moment projection M. For the G sublattice the representation can be written in the form

$$c_{g\sigma} = \sum_{\alpha} \overline{\gamma}_\sigma(\alpha) X_g^{\alpha}. \tag{12}$$

In this description we have included the fact that the single-ion problem for the G sublattice differs from the same problem for the F sublattice only by a change in the sign of the angle θ. It is obvious that the parameters of the representation $\gamma_\sigma(\alpha)$ are in reality functions of the angle θ. For brevity we will use the following written forms: $\gamma_\sigma(\theta,\alpha) \equiv \gamma_\sigma(\alpha)$, $\overline{\gamma}_\sigma(\alpha) \equiv \gamma_\sigma(-\theta,\alpha)$.

Using Eqs. (11) and (12), we find the following form for the total Hamiltonian in the atomic representation:

$$H' = H'_{\rm ion} + H_{\rm int}, \tag{13}$$

where

$$H'_{\rm ion} = \sum_{fM} E_M X_f^{MM} + \sum_{gM} E_M X_g^{MM} + \sum_{fn} (E_n - \mu) X_f^{nn} + \sum_{gn} (E_n - \mu) X_g^{nn}. \tag{14}$$

The operator $H'_{\rm ion}$ differs from $H_{\rm ion}$ in Eq. (6). First of all, $H'_{\rm ion}$ is written in the atomic representation; secondly, in $H'_{\rm ion}$ we used a truncated basis corresponding to states of the lowest multiplet.

In Eq. (14) the first two terms take into account the single-ion states without holes with energies $E_M = -\overline{H}M$, $M = S, S-1, \ldots, -S$. The remaining terms in Eq. (14) correspond to inclusion of states with a single hole.

Let us write the interaction operator in a form that is convenient for applying the diagram technique for Hubbard operators:

$$H_{\rm int} = \sum_{AA'} \sum_{l\in A, l'\in A'} \sum_{\alpha\beta\sigma} V_\sigma^{\alpha\beta}(l,l') X_l^{-\alpha} X_{l'}^{\beta}. \tag{15}$$

Here, in the interest of abbreviating the description, we have denoted summation over the sublattices as sums over A and A', where the variables A and A' can take on two values: F and G. The matrix elements of the interaction operator are defined by the expressions

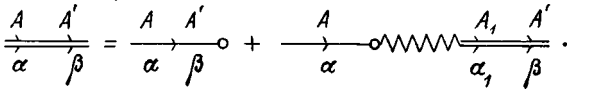
$$\begin{aligned}
V_\sigma^{\alpha\beta}(f,f') &= t_{f,f'}\gamma_\sigma(\alpha)\gamma_\sigma(\beta),\\
V_\sigma^{\alpha\beta}(g,g') &= t_{g,g'}\overline{\gamma}_\sigma(\alpha)\overline{\gamma}_\sigma(\beta),\\
V_\sigma^{\alpha\beta}(f,g) &= t_{f,g}[\cos\theta\gamma_\sigma(\alpha)\overline{\gamma}_\sigma(\beta) + 2\sigma\sin\theta\gamma_\sigma(\alpha)\overline{\gamma}_{\bar\sigma}(\beta)],\\
V_\sigma^{\alpha\beta}(g,f) &= t_{f,g}[\cos\theta\overline{\gamma}_\sigma(\alpha)\gamma_\sigma(\beta) + 2\sigma\sin\theta\overline{\gamma}_{\bar\sigma}(\alpha)\gamma_\sigma(\beta)].
\end{aligned} \tag{16}$$

3. GREEN'S FUNCTIONS AND DISPERSION EQUATION

In order to compute the spectrum of Fermi-type, we introduce into the discussion the Matsubara Green's function in its atomic representation:

$$G_{\alpha\beta}^{AA'}(l\tau, l'\tau') = -\langle T_\tau \widetilde{X}_l^{\alpha}(\tau)\widetilde{X}_{l'}^{-\beta}(\tau')\rangle, \tag{17}$$

where the lattice label $l(l')$ takes on a set of values corresponding to the $A(A')$ sublattice. The remaining notations are standard and are contained in Refs. 11 and 16. The following graphical equation can be written for the Fourier transforms $G_{\alpha\beta}AA'(\mathbf{k},\omega_n)$ in the simplest approximation:



(18)

In this graphical equation, a thin line A/α corresponds to the intra-ion propagator

$$G_\alpha(\omega_n) = [i\omega_n + \alpha E]^{-1}, \quad \omega_n = (2n+1)\pi T, \tag{19}$$

in which the scalar product of the root vector α and the vector E is determined by the equation $\alpha E \equiv \alpha(M,n)E = E_M - E_n$. Because the sublattices are identical and the values of single-ion energy levels for them are the same, $G_\alpha(\omega_n)$ does not depend on the sublattice label. When the diagrams are written out in analytic form, the wavy lines correspond to the Fourier transform of the interaction matrix element (16) summed over values of the spin moment projection. The specific values that label the interaction matrix elements are determined by the labels of the Green's function lines that join with the wavy line. For example, in Eq. (18) for $A=F$ and $A_1=G$ the interaction lines represent the expression

$$\Gamma_{\mathbf{q}}^{\alpha\alpha_1} = \sum_\sigma \Gamma_{\mathbf{q}}[\cos\theta\gamma_\sigma(\alpha)\gamma_\sigma(\alpha_1) + 2\sigma\sin\theta\gamma_\sigma(\alpha)\overline{\gamma}_\sigma(\alpha_1)],$$

$$\Gamma_{\mathbf{q}}=\frac{1}{N}\sum_{g} t_{fg}\exp\{i\mathbf{q}(\mathbf{R}_f-\mathbf{R}_g)\}. \tag{20}$$

If, however, $A=F$, $A_1=F$, the wavy line corresponds to the following analytical description:

$$t_{\mathbf{q}}^{\alpha\alpha_1}=\sum_{\sigma} t_{\mathbf{q}}\gamma_{\sigma}(\alpha)\gamma_{\sigma}(\alpha_1),$$

$$t_{\mathbf{q}}=\frac{1}{N}\sum_{f'} t_{ff'}\exp\{i\mathbf{q}(\mathbf{R}_f-\mathbf{R}_{f'})\}. \tag{21}$$

The remaining notations are standard and are discussed in detail in Refs. 11 and 15.

From Eq. (18) there follows a system of equations in analytic form

$$G_{\alpha\beta}^{FF}(\mathbf{k},\omega_n)=\delta_{\alpha\beta}G_{\alpha}(\omega_n)+G_{\alpha}(\omega_n)b(\alpha)\times\{t_{\mathbf{k}}^{\alpha\alpha_1}G_{\alpha_1\beta}^{FF}(\mathbf{k},\omega_n)+\Gamma_{\mathbf{k}}^{\alpha\alpha_1}G_{\alpha_1\beta}^{GF}(\mathbf{k},\omega_n)\}, \tag{22}$$

$$G_{\alpha\beta}^{GF}(\mathbf{k},\omega_n)=G_{\alpha}(\omega_n)b(\alpha)\Bigg\{t_{\mathbf{k}}\sum_{\sigma}\overline{\gamma}_{\sigma}(\alpha)\overline{\gamma}_{\sigma}(\alpha_1)G_{\alpha_1\beta}^{GF}(\mathbf{k},\omega_n)+\Bigg[\Gamma_{\mathbf{k}}\cos\theta\sum_{\sigma\alpha_1}\overline{\gamma}_{\sigma}(\alpha)\gamma_{\sigma}(\alpha_1)+\Gamma_{\mathbf{k}}\sin\theta\sum_{\sigma\alpha_1}2\sigma\overline{\gamma}_{\sigma}(\alpha)\gamma_{\sigma}(\alpha_1)\Bigg]G_{\alpha_1\beta}^{FF}(\mathbf{k},\omega_n)\Bigg\}. \tag{23}$$

In solving this system of equations, it is very useful to note that the nature of the interaction matrix element leads to separability with respect to the labels of the root vectors α,α_1. This separability is quite evident from Eqs. (20) and (21). Using the methods for solving equations with separable kernels,[13,14] we find that the dispersion equation can be written in the form

$$\begin{aligned}0=&[(1-t_{\mathbf{k}}L_{\uparrow\uparrow})^2-\Gamma_{\mathbf{k}}^2M_{\uparrow\uparrow}^2][(1-t_{\mathbf{k}}L_{\downarrow\downarrow})^2-\Gamma_{\mathbf{k}}^2M_{\downarrow\downarrow}^2]\\&+\Gamma_{\mathbf{k}}^4M_{\uparrow\downarrow}^2M_{\downarrow\uparrow}^2+2(\Gamma_{\mathbf{k}}^2M_{\uparrow\downarrow}M_{\downarrow\uparrow}-t_{\mathbf{k}}^2L_{\uparrow\downarrow}^2)[(1-t_{\mathbf{k}}L_{\uparrow\uparrow})\\&\times(1-t_{\mathbf{k}}L_{\downarrow\downarrow})-\Gamma_{k}^2M_{\uparrow\uparrow}M_{\downarrow\downarrow}]+2t_{\mathbf{k}}\Gamma_{\mathbf{k}}^2L_{\uparrow\downarrow}\\&\times(M_{\uparrow\downarrow}-M_{\downarrow\uparrow})(M_{\uparrow\uparrow}-M_{\downarrow\downarrow})+t_{\mathbf{k}}^2L_{\uparrow\downarrow}^2[t_{\mathbf{k}}^2L_{\uparrow\downarrow}^2\\&-\Gamma_{\mathbf{k}}^2(M_{\uparrow\downarrow}^2+M_{\downarrow\uparrow}^2)],\end{aligned} \tag{24}$$

where

$$L_{\sigma\sigma_1}(\omega)=\sum_{\alpha}\overline{\gamma}_{\sigma}(\alpha)\overline{\gamma}_{\sigma_1}(\alpha)G_{\alpha}(\omega)b(\alpha),$$

$$M_{\sigma\sigma_1}(\omega)=L_{\sigma\sigma'}(\omega)\cos\theta+(2\sigma_1)L_{\sigma\sigma_1}^{-}(\omega)\sin\theta. \tag{25}$$

In studying the de Haas-van Alphen-effect it is the low-temperature regime that is of primary interest to us. This regime is defined mathematically by the inequality $T\ll T_N$. It is well known that occupation numbers for single-ion states are distributed according to atomic statistics. Therefore, in the temperature range of interest to us, only the numbers N_1 and N_S are interesting, while the others are exponentially small. This greatly decreases the number of intra-ion transitions that participate in forming the collective quasiparticle spectrum, and simplifies the form of the dispersion equation. In fact, when $T\ll T_N$ contributions come only from those transitions in which at least one of the low-lying states of the multiplets participates, with a hole or without one.

Another factor that allows us to simplify the structure of the dispersion equation is the existence of the small parameter (ν/ε). We will carry all calculations out to quadratic accuracy in this parameter. The need to do calculations to this level of accuracy is dictated by considerations involving effects that are $\sim\theta^2$ in the neighborhood of the transition from the canted to the collinear phase.

4. SPECTRUM OF HOLES IN A NEIGHBORHOOD OF THE SPIN-FLIP TRANSITION

After calculating the functions $L_{\sigma_1\sigma_2}$ and $M_{\sigma_1\sigma_2}$ to the specified accuracy, we obtain the dispersion equation in the following form:

$$(1-\Gamma_{\mathbf{k}}^2M_{\uparrow\uparrow}^2)(1-\Gamma_{\mathbf{k}}^2M_{\downarrow\downarrow}^2)+2\Gamma_{\mathbf{k}}^2M_{\uparrow\downarrow}M_{\downarrow\uparrow}\times(1-\Gamma_{\mathbf{k}}^2M_{\uparrow\uparrow}M_{\downarrow\downarrow})=0, \tag{26}$$

where we assume that analytic continuation has already been done. For simplicity we assume that the only nonzero matrix elements are those $t_{l'}$ which correspond to hopping of a charge carrier between nearest neighbors.

Solving Eq. (26), we find the branch of Fermi excitations of interest to us:

$$\begin{aligned}E(\mathbf{k})=&\frac{A(S+1)}{2}+\frac{\overline{H}}{2S+1}+\frac{h}{2}\frac{2S-1}{2S+1}-\frac{2S}{2S+1}|\Gamma_{\mathbf{k}}|\cos\theta\\&+N_1\Bigg\{\frac{2S-1}{(2S+1)^3}\left(\frac{v}{\varepsilon}\right)^2|\Gamma_{\mathbf{k}}|\\&+\frac{2(2S-1)}{(2S+1)^2}\left(\frac{v\sin\theta}{\varepsilon}\right)|\Gamma_{\mathbf{k}}|-\frac{2S-1}{(2S+1)^2}\left(\frac{v^2}{\varepsilon}\right)\\&-\frac{\sin^2\theta}{2S+1}|\Gamma_{\mathbf{k}}|\Bigg\}.\end{aligned} \tag{27}$$

It is clear that for $H=H_c$, when $\theta=0$ this spectrum becomes the polaron spectrum obtained previously by Nagaev.[7]

These expressions for the spectrum of magnetopolaron states of an antiferromagnet in the canted phase are fundamental for studying many kinetic and galvanomagnetic phenomena in conducting antiferromagnets in a magnetic field. This statement is especially true for the de Haas-van Alphen effect in these systems. Limitations of space have not allowed us to discuss studies of the influence of magnetopolaron states on magnetic moment oscillations in a quantizing magnetic field here. The corresponding results based on derivations in this work will be published in a separate article. We note here, however, that the expressions obtained will allow us to follow the temperature evolution of magnetopolaron states. This information is necessary, for example, when temperature behavior of magnetoresistance in conducting antiferromagnets is investigated.

This work was carried out with the financial support of the Science Fund from the Krasnoyarsk region (Grant No. 5F0158).

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Translated by Frank J. Crowne