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A joint scientific session of the Physical Sciences Division of the Russian Academy of Sciences (RAS) and the Joint Physical Society of the Russian Federation was held on 9 April 2003 at the P N Lebedev Physics Institute, RAS. The following reports were presented at the session:

(1) **Kagan M Yu, Klaptsov A V, Brodskiĭ I V** (P L Kapitza Institute for Physical Problems, RAS, Moscow), **Kugel' K I, Sboĭchakov A O, Rakhmanov A L** (Institute for Theoretical and Applied Electrodynamics, Moscow) "Small-scale phase separation and electron transport in manganites";

(2) **Aksenov V L, Balagurov A M, Pomyakushin V Yu** (Joint Institute for Nuclear Research, Dubna, Moscow region) "Neutron diffraction analysis of doped manganites";

(3) **Ustinov V M** (A F Ioffe Physicotechnical Institute, RAS, St.-Petersburg) "Semiconductor quantum dot lasers";

(4) **Sergeev A M** (Institute of Applied Physics, RAS, Moscow) "Topical problems of femtosecond optics";

(5) **Lukin V P** (Institute of Atmospheric Optics, Siberian Branch of the RAS, Tomsk) "Atmospheric adaptive optics";

(6) **Kul'chin Yu N** (Far East State Technical University, Vladivostok) "Adaptive distributed optoelectronic information-measuring systems";

(7) **Bessonov E G, Vinogradov A V, Gorbunkov M V, Tur'yanskiĭ A G, Feshchenko R M, Shabalin Yu V** (P N Lebedev Physics Institute, RAS, Moscow) "Laser electron-beam X-ray source for medical applications";

(8) **Lukash V N** (Astrospace Center of the P N Lebedev Physics Institute, RAS, Moscow) "Cosmological model and universe structure formation".

An abridge version of the reports 1, 2, 5, 6, 7, and 8 is given below.

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Small-scale phase separation and electron transport in manganites

M Yu Kagan, A V Klaptsov, I V Brodskiĭ, K I Kugel', A O Sboĭchakov, A L Rakhmanov

1. Introduction

Manganites, manganese-based magnetic oxides like $LaMnO_3$, have been known for more than 50 years. In recent years they have attracted considerable interest primarily due to the discovery of the effect of colossal negative magnetoresistance in $La_{1-x}Ca_xMnO_3$ materials in 1993 [1]. The unusual physical properties and the richness of the phase diagram of manganites have generated a wealth of papers concerned with different aspects of the phase diagram and transport properties of these compounds.

Special attention to manganites arises from the possibility of forming inhomogeneous charge and spin states of various types, like lattice and magnetic polarons, droplet and stripe structures, etc. Similar phenomena are inherent in many strongly correlated structures with potential interaction energy of electrons predominating over their kinetic energy. In particular, they are widely discussed for materials exhibiting high-temperature superconductivity. We note that one of the first striking examples of inhomogeneous states of this kind is ferromagnetic droplets (ferrons) [2, 3] in an antiferromagnetic matrix under weak doping. The linkage between phase separation and the nature of colossal magnetoresistance in manganites has been the subject of wide speculation in recent years.

Phase separation is commonly investigated in the range of low densities and temperatures in the vicinity of the transition point between the antiferromagnetic state and the ferromagnetic ground state. However, there is some evidence that the state of manganites is also inhomogeneous even in the range of optimal densities and low temperatures (in the ferromagnetic phase) [4]. At high temperatures and optimal densities, experimental data once again testify to the presence of significant short-range ferromagnetic correlations (temperature ferrons [5, 6]) in the paramagnetic phase.

The fundamental results of investigations of the gross phase diagram of manganites, obtained by our group, were published in a large review paper [6]. During the last two years we have obtained a set of interesting new data concerning the compromise between the canted state and small-scale phase separation, as well as several results that are of considerable importance in describing nontrivial transport properties of manganites. Our paper is dedicated to precisely these new results.

2. Theoretical model

To account for the correlations between transport and magnetic properties in manganites, Zener [7] put forward a double-exchange model. This model assumes that a conduction electron travels in $La_{1-x}Ca_xMnO_3$ along Mn^{4+} ions and strongly interacts (in a ferromagnetic way) with the local moments of manganese ions. Since the spin of the conduction electron should be parallel with the local spin, in the classical picture this electron cannot move in the antiferromagnetic

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environment. Anderson and Hasegawa [8] solved the problem for the amplitude of atom-to-atom electron jump. They found that the hopping amplitude is of the form $t_{\rm eff} = t\cos(\chi/2)$, where t is the bare hopping amplitude, and χ is the angle between the local atomic moments. Therefore, the conduction electron tends to ferromagnetically order the nearby local spins. This gives rise to competition between the antiferromagnetism of local spins (caused by the superexchange mechanism) and the ferromagnetism of the local spins through the conduction electron (the double-exchange mechanism). De Gennes [9] hypothesized that this competition results in a homogeneous canted state: the angle between the local spins of sublattices is constant throughout the sample and varies monotonically from π (a collinear antiferromagnetic structure) to 0 (a collinear ferromagnetic structure) with increasing charge carrier concentration. Later, Nagaev [10] included the quantum fluctuations of local spins and concluded that an electron can travel even in an antiferromagnetic environment with a small hopping amplitude $t/\sqrt{2S+1}$.

Several authors independently predicted [2, 3] the existence of a self-trapped electron state in an antiferromagnetic matrix, which has come to be known as a magnetic polaron or a ferron. For an electron embedded in an antiferromagnetic environment it is quite often energetically benefited to form a ferromagnetic domain around itself and become autolocalized in it, resulting in a small-scale phase separation between ferro- and antiferromagnetic phases. One of the authors [6] showed that the double-exchange model is, in the case of low charge-carrier concentrations, unstable towards phase separation and that the energy of ferron state is lower than the energy of the homogeneous canted state [11]. However, the question arises of whether the polaron state is stable with the inclusion of quantum fluctuations of local spins.

The basic model suggested for the description of manganite properties is the ferromagnetic Kondo lattice model (the s – d model):

$$\hat{H} = -J_{\rm H}\sum_i S_i\sigma_i - t\sum_{\langle i,j\rangle} c_{i\sigma}^{+}c_{j\sigma} + J_{\rm ff}\sum_{\langle i,j\rangle} S_iS_j\,, \qquad (1)$$

where S_i is the local spin of d electrons of manganese, $\sigma_i = c_{i\alpha}^{+}\sigma_{\alpha\beta}c_{i\beta}$ is the spin of conduction electrons, and the symbol $\langle i, j\rangle$ implies summation over the nearest neighbors. We note that a strong Hund interaction prevents two electrons with different spin projections from occupying a common site. The first term in the Hamiltonian (1) describes the Hund interaction of a local spin $S = 3/2$ with a spin of a conduction electron. In real manganites, the magnitude of Hund interaction $J_{\rm H}$ is of the order of 1 eV. The second term in expression (1) corresponds to the kinetic energy of conduction electrons. We note that a strong electron – lattice interaction significantly narrows the conduction band width ($W = 2zt$), resulting in a small value of the hopping amplitude $t \approx 0.3$ eV. The last term in the Hamiltonian (1) corresponds to the weak antiferromagnetic interaction of two neighboring local spins ($J_{\rm ff} \sim 0.001$ eV).

In the case of strong Hund interaction ($J_{\rm H} \gg W \gg J_{\rm ff}$), the Hamiltonian (1) takes the form

$$\hat{H} = -\sum_{\langle i,j\rangle} t(\chi_{ij})a_i^{+}a_j + J_{\rm ff}S^2\sum_{\langle i,j\rangle}\cos\chi_{ij}\,, \qquad (2)$$

where a_i^{+} and a_j are the creation and annihilation operators for spinless fermions (conduction electrons whose spin is aligned with the local spin), $t(\chi_{ij})$ is the effective hopping amplitude, and χ_{ij} is the angle between the magnetic moments of two neighboring atoms. Interestingly, if all local spins were ferromagnetically aligned, the conduction electrons would move freely with a hopping amplitude t. Therefore, the Hamiltonian (2) describes the competition between the ferromagnetism due to double exchange (the first term) and the antiferromagnetism due to superexchange (the second term).

In Nagaev's quantum approach, the local spin and the spin of a conduction electron form, at the site occupied by the conduction electron, a state with a total spin $S + 1/2$, but with two possible spin projections $S \pm 1/2$. By this means there exist two effective bands corresponding to two different projections of the total spin:

$$t_{\pm}(\chi) = \frac{t}{2S+1}\left(\sqrt{2S+1+S^2\cos^2\frac{\chi}{2}} \pm S\cos\frac{\chi}{2}\right). \qquad (3)$$

The quantum hopping amplitude is significantly different from the classical de Gennes – Anderson – Hasegawa hopping amplitude $t\cos(\chi/2)$. Unlike the classical description, in the quantum case an electron can travel over the antiferromagnetic matrix, forming a state with $S_{\rm tot}^z = S + 1/2$ at one site and a state with $S_{\rm tot}^z = S - 1/2$ at the adjacent one:

$$\left|S_{\rm tot}^z = S+\frac{1}{2}\right\rangle \to \left|S_{\rm tot}^z = S-\frac{1}{2}\right\rangle \to \left|S_{\rm tot}^z = S+\frac{1}{2}\right\rangle \ldots \qquad (4)$$

Therefore, when one electron moves along the antiferromagnetic matrix, the widths of the two bands with $S_{\rm tot}^z = S \pm 1/2$ are equal: $t_+ = t_- = t/\sqrt{2S+1}$; in the case of motion along a ferromagnetic matrix, one band is $1/(2S+1)$ times narrower than the other: $t_+ = t$ and $t_- = t/(2S+1)$.

3. Ferron state

As discussed above, in the case of classical hopping amplitude $t\cos(\chi/2)$, an electron can autolocalize and form a ferromagnetic droplet in an antiferromagnetic matrix. In the simplest approximation we assume that the boundary between the ferromagnetic and antiferromagnetic domains is sharp, without the transition region of inhomogeneous canting. In this case, the energy of the ferron state is given by [6]

$$E_{\rm pol} = -tx\left(z - \frac{\pi^2 d^2}{R^2}\right) + \frac{1}{2}\,zJ_{\rm ff}S^2\,\frac{4\pi}{3}\,x\left(\frac{R}{d}\right)^3 - \frac{1}{2}\,zJ_{\rm ff}S^2\left[1 - \frac{4\pi}{3}\,x\left(\frac{R}{d}\right)^3\right], \qquad (5)$$

where R is the ferron radius, and d is the lattice period. The first term in expression (5) describes the gain in the kinetic energy owing to ferromagnetic domain formation. The second term corresponds to the loss in the Heisenberg energy of antiferromagnetic interaction of the local spins in the ferron. The last term corresponds to the gain in the energy of antiferromagnetic interaction outside of the ferron. The ferron radius is derived from the energy minimization condition $dE/dR = 0$. Eventually, we obtain the following relations for the energy and the radius of the

ferron (see Refs [6, 11]):

$$E_{\rm pol} = -zxt + \frac{5}{3}\,\pi^2 xt\left(\frac{2zJ_{\rm ff}S^2}{\pi t}\right)^{2/5} - \frac{1}{2}\,zJ_{\rm ff}S^2\,, \qquad (6)$$

$$R_{\rm pol} = d\left(\frac{\pi t}{2zJ_{\rm ff}S^2}\right)^{1/5}. \qquad (7)$$

We consider the case of quantum hopping amplitude (3). Since the hopping amplitude $t \neq 0$ in the antiferromagnetic domain and the conduction electron can travel along the antiferromagnetic matrix with an effective mass $m^* \sim \sqrt{2S+1}$, in this situation it is instructive to investigate the ferron state for stability. Since the analysis of the discrete case (2) is a rather complicated task, we consider the continuous limit, inferring the ferron radius to be far greater than the lattice period d (in what follows we assume that $d = 1$). Then, the total energy (2) can be written down as [12]

$$E = -\int\left[z|\Psi|^2 + \Psi^*\Delta\Psi\right]t(\chi)\,{\rm d}V + zJ_{\rm ff}S^2\int\cos^2\frac{\chi}{2}\,{\rm d}V, \quad (8)$$

$$t(\chi) = \frac{t}{2S+1}\left(\sqrt{2S+1+S^2\cos^2\frac{\chi}{2}} + S\cos\frac{\chi}{2}\right).$$

One can see from Eqn (8) that the total energy will lie between the two limiting values corresponding to electron motion over the ferromagnetic matrix (FM) or the antiferromagnetic matrix (AFM):

$$E_{\rm FM} = -zt < E < E_{\rm AFM} = -\frac{zt}{\sqrt{2S+1}}\,.$$

Since the electron wave function should be normalized to unity, $\int|\Psi|^2\,{\rm d}V = 1$, we will minimize the functional $F = E(\Psi, \chi) - t\beta\int|\Psi|^2\,{\rm d}V$ with respect to the parameters χ and Ψ, where β is the Lagrange factor. The corresponding Euler–Lagrange equations take the form

$$\left[2z\Psi + \Delta\Psi\right]t(\chi) + \Delta\left[t(\chi)\Psi\right] - 2\beta t\Psi = 0\,, \qquad (9)$$

$$\left[\left(z|\Psi|^2 + \Psi^*\Delta\Psi\right)\frac{\partial t(\chi)}{\partial\cos(\chi/2)} - 2zJ_{\rm ff}S^2\cos\frac{\chi}{2}\right]\sin\frac{\chi}{2} = 0\,. \qquad (10)$$

To solve this system of equations, we take advantage of the following iterative procedure [12]: (i) select the trial function for the canting angle $\chi(r)$; (ii) solve the first differential equation (9) to obtain the electron wave function $\Psi(r)$; (iii) employing the resultant wave function $\Psi(r)$, solve the second algebraic equation (10) in order to define the canting angle function $\chi(r)$, and (iv) revert to step (i) until the acceptable accuracy is reached.

The functions $\Psi(r)$ and $\chi(r)$ obtained through the numerical solution of Eqns (9) and (10) are plotted in Fig. 1, where they are compared with the classical solution (6). Hence, a ferron constitutes a well localized object and the transition region from the ferromagnetic phase ($\chi = 0$) to the antiferromagnetic one ($\chi = \pi$) is rather narrow. However, the ferron state may disappear for relatively small values of the parameter $\alpha = t/J_{\rm ff}S^2 \lesssim \alpha_{\rm c} \approx 75$. Indeed, one can see from Fig. 2 that there exists a threshold-type transition from the ferron state to the homogeneous antiferromagnetic state for $\alpha < \alpha_{\rm c}$. In this case, the total energy of the ferron state becomes equal to the energy of the bottom of the conduction band in the antiferromagnetic matrix. We note that a more precise calculation of the $\alpha_{\rm c}$ value calls for the solution of the variational problem in the discrete case, since the ferron radius becomes comparable to the lattice spacing for small values of the parameter α and, accordingly, the continuous approximation (8) is no longer correct. The solution of the discrete problem will be considered in a separable publication.

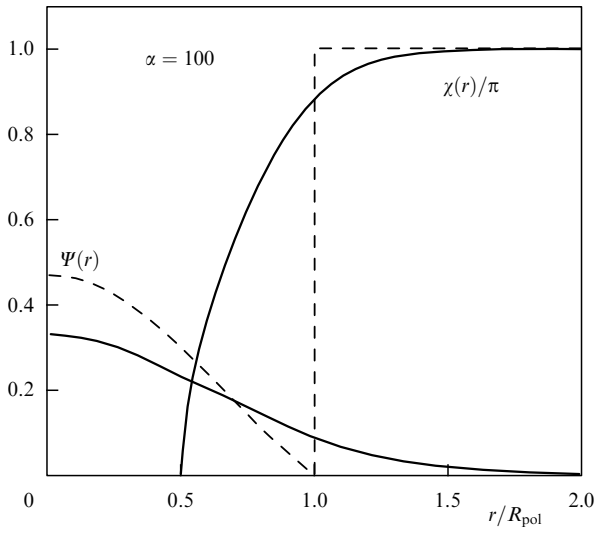

Figure 1. Electron wave function $\Psi(r)$ and canting angle $\chi(r)$ in the quantum (solid lines) and classical (dashed lines) cases.

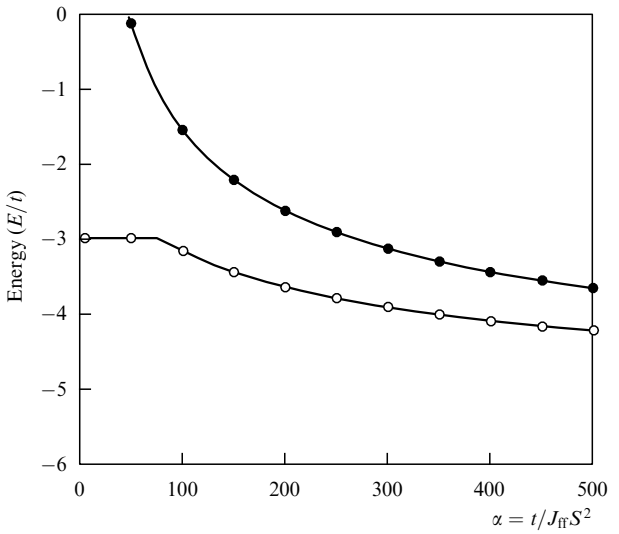

Figure 2. Ground-state energy of a ferron in the quantum (○) and classical (●) cases for a local spin $S = 3/2$. The bottom of the electronic-conduction band corresponds to $E_{\rm FM} = -6t$ in the ferromagnetic matrix, and to $E_{\rm AFM} = -3t$ in the antiferromagnetic matrix.

4. Transport properties

The transport properties of manganites, being apparent from the model of ferrons in a dielectric matrix, were first considered in Refs [13, 14]. Below we briefly recall the main points in the description of transport characteristics of manganites and outline interesting new results concerning the temperature and field dependences of magnetoresistance.

We address ourselves to a system consisting of N carriers embedded in an antiferromagnetic matrix of volume V_s. In

the ground state, the carriers form autolocalized ferron states of radius R. The ferrons and the current carriers in the sample are assumed to be equal in number. In this case, the carriers can tunnel from one ferron to a neighboring one to produce empty and doubly occupied ferrons.

To produce a doubly occupied ferron, the system has to overcome an energy barrier equal by order of magnitude to the energy V of Coulomb repulsion between the carriers. If it is assumed that the ferron radius is about 10–20 Å and the permittivity $\varepsilon \approx 10-20$, for the Coulomb repulsion energy we get $V = e^2/\varepsilon R \approx 0.1-0.2$ eV. In what follows we suppose the temperature to be low enough in comparison with the magnitude of the Coulomb barrier V, and the probability of forming a ferron containing more than two carriers can be neglected. We note that the carriers in a doubly occupied ferron can produce states with a total spin equal to 0 or 1. In the latter case, due to the Pauli exclusion principle the carriers occupy the ground and first excited levels in the ferron. The spacing of these energy levels in the ferron can be shown to exceed the Coulomb repulsion energy V, and we therefore will consider only such doubly occupied ferron states wherein the electron spins are antiparallel.

Let us introduce the following notation for the description of ferron states: $q = 0$ corresponds to an empty ferron, $q = 1\uparrow$ or $q = 1\downarrow$ corresponds to a ferron with one carrier whose spin has a projection $\sigma/2$ on the direction of the magnetic moment of the ferron, and $q = 2$ corresponds to a doubly occupied ferron. The ferron energy E_q in an external magnetic field, with the inclusion of anisotropy field, then takes the form

$$E_0 = zJ_{\rm ff}S^2\,\frac{4\pi}{3}\left(\frac{R}{d}\right)^3 - M_0(H\cos\theta + H_{\rm a}\cos^2\psi)\,, \qquad (11)$$

$$E_{1,\lambda} = E_0 + t\left(\frac{\pi d}{R}\right)^2 - \sigma\bar{J}\,, \qquad (12)$$

$$E_2 = E_0 + 2t\left(\frac{\pi d}{R}\right)^2 + \frac{e^2}{\varepsilon R}\,, \qquad (13)$$

where $H_{\rm a}$ is the anisotropy field, $M_0 = \mu_{\rm B}gS\,4\pi R^3/3d^3$ is the magnetic moment of the ferron, θ and ψ are the angles between M_0 and H and between M_0 and the easy axis, respectively, $\sigma/2$ is the carrier spin projection onto the M_0 direction, $\mu_{\rm B}$ is the Bohr magneton, and g is the Landé splitting factor. The third term in Eqn (12) describes the energy of interaction between the autolocalized carrier and the molecular magnetic field of ferromagnetically ordered local spins $H_{\rm eff} = \bar{J}/\mu_{\rm B}$ (here, $\bar{J} \sim T_{\rm Curie}$ is the effective exchange integral related to the ferromagnetic ordering temperature).

We assume that the direction of the total magnetic moment of the droplet changes relatively slowly with time and the conventional thermodynamic treatment is valid. The radius of a ferron with one carrier can then be determined by minimizing the energy (12) with respect to the ferron radius. In the linear approximation in the magnetic field value we arrive at the expression

$$R(H) \cong R_{\rm pol}\left[1 + \frac{b}{5}(H\cos\theta + H_{\rm a}\cos^2\psi)\right], \qquad b = \frac{\mu_{\rm B}g}{zJ_{\rm ff}S}\,, \qquad (14)$$

where $R_{\rm pol}$ is defined by formula (7), and the magnitude of b is measured in units of T^{-1}.

As shown in Ref. [13], the characteristic carrier tunneling time is much shorter than the relaxation time of the spin subsystem. Hence we can suppose that the radii of empty and doubly occupied ferrons coincide with the radius of a ferron containing one electron. We also assume that the total number of ferrons (empty, singly occupied, and doubly occupied) remains constant and equal to N.

In the framework of our model, the charge transfer can be effected in one of the following ways [13]:

(1) Initially there are two ferrons containing one electron each. Then, an electron tunnels from one ferron to the other, and in the final state there is one empty ferron and one doubly occupied ferron.

(2) An electron tunnels from a doubly occupied ferron to an empty ferron. In the final state there occur two singly occupied ferrons.

(3) An electron tunnels from a doubly occupied ferron to a singly occupied ferron.

(4) An electron tunnels from a singly occupied ferron to an empty ferron.

Each of the above-listed processes is defined by the corresponding tunneling probability which assumes the form [14]

$$W(q_1', q_2'; q_1, q_2) = \omega_0 f(v)\exp\left(-\frac{r}{l} + \frac{e(\mathbf{Er})}{kT} - \frac{E_{q_1'} + E_{q_2'} - E_{q_1} - E_{q_2}}{2kT}\right), \qquad (15)$$

where r and v are the distance between the ferrons and the angle between the directions of their magnetic moments, respectively, q_1, q_2 and q_1', q_2' are the initial (prior to the electron transition) and final states of the pair of ferrons, l is the tunneling length, and ω_0 is the characteristic frequency of electron motion in the potential well. The pre-exponential factor $f(v)$ in Eqn (15) describes the spin-dependent tunneling which takes into account different orientations of the electron spin in the initial and final states, and is of the form

$$f(v) = \frac{\cosh\left(\bar{J}\cos\left(v/kT\right)\right)}{\cosh\left(\bar{J}/kT\right)}\,. \qquad (16)$$

To calculate the conductivity, we should sum up the contributions to the current made by each of the possible tunnel processes. Eventually we arrive at the following expression for the conductivity

$$\sigma(H) = \frac{32\pi e^2 l^5 \omega_0}{V_s^2 kT} \times \sum_{q_1, q_2} \bar{N}_{q_1}\bar{N}_{q_2}\left\langle f(v)\exp\left(\frac{E_{q_1} + E_{q_2} - E_{q_1'} - E_{q_2'}}{2kT}\right)\right\rangle_{q_1, q_2}, \qquad (17)$$

where $\bar{N}_q$ is the average number of electrons in a ferron in the q state, and $\langle\ldots\rangle_{q_1, q_2}$ stands for averaging over the directions of the magnetic moments of two ferrons in the states q_1 and q_2. In the derivation of Eqn (17) it was assumed that the directions of the magnetic moments of the ferrons remain invariable in the course of electron tunneling from one ferron to the other.

The change of conductivity in an external magnetic field arises from the dependence of the transition probability $W(q_1', q_2'; q_1, q_2)$ and average occupation numbers $\bar{N}_q$ on the

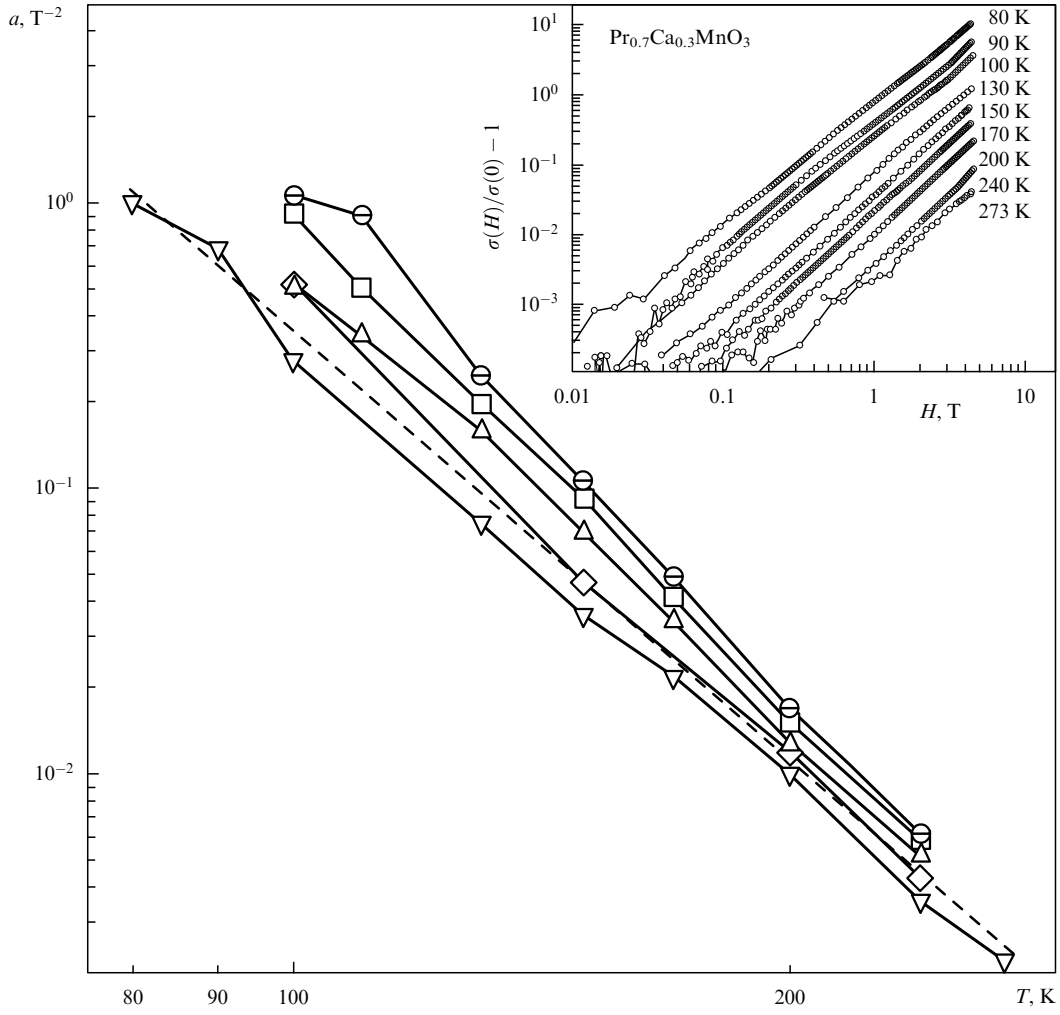

Figure 3. Experimental magnetoresistance $MR = a(T)H^2$ of different samples measured as a function of temperature and external magnetic field value (inset to figure) [16]. The dashed line corresponds to the theoretical curve obtained by formula (19): $MR(H) \sim H^2/T^5$, employing the following parameter values: $S = 2$, $g = 2$, $z = 6$, $N_m = 4\pi R^3/3d^3 = 130$, $\bar{J}/k = 15$ K, $H_a = 0.5$ T, and $\cos\beta = 1$.

applied field. It is possible to recognize several main contributions to magnetoresistance, which are possible to calculate analytically in the limiting case of small fields. First, the onset of magnetic field leads to a change in the ferron radius $R(H)$ described by Eqn (14), with a consequential change in the ferron energy (11)–(13) in the magnetic field. The corresponding contribution to the magnetoresistance will be of the form

$$MR_1(H) \approx \frac{3}{100}\,\frac{M_0^2 H^2}{(kT)^2}\,. \tag{18}$$

Second, it is possible to distinguish the contribution from spin-dependent tunneling. When a monocrystal is embedded in an external magnetic field in such a way that the angle between the easy axis and the magnetic field is β, the corresponding contribution to magnetoresistance reads

$$MR_2(H) \approx \frac{2}{225}\,\frac{M_0^3 \bar{J}^2 H_a H^2}{(kT)^5}\left(\cos^2\beta - \frac{1}{3}\right). \tag{19}$$

In all previous calculations we assumed that ferrons were spherical in shape and the field H_a was defined only by the crystal anisotropy. However, the ferron shape effect can be shown to be rather significant even for small departures from sphericity owing to the demagnetization factor. The anisotropy of a ferron shape contributes to the effective anisotropy field H_a. In this case, assuming for the sake of simplicity that the magnetic anisotropy is uniaxial, we can incorporate both the magnetic crystalline anisotropy and the shape anisotropy into the field H_a.

Hence, in small magnetic fields the magnetoresistance decreases with temperature as T^{-2} when the contribution (18) prevails or as T^{-5} when the term (19) becomes dominant. In the general case, the magnetoresistance behaves as follows

$$MR(H) = \frac{AH^2}{(kT)^2} + \frac{BH^2 H_a}{(kT)^5}\,, \tag{20}$$

where A and B are the constants. In a strong magnetic field (10–20 T and over), the magnetoresistance increases exponentially [13]:

$$MR(H) \approx \frac{\bar{J}}{kT}\coth\frac{\bar{J}}{kT}\exp\left(\frac{VbH}{10kT}\right), \tag{21}$$

and its value may range even up to several hundred percent far from the percolation threshold.

5. Conclusions

We have shown that a conduction electron produces an autolocalized state in a wide range of values of the parameter $\alpha = t/J_{\rm ff}S^2$: the electron is localized in a ferromagnetic droplet of finite radius, embedded into an antiferromagnetic matrix. Therefore, the double-exchange model exhibits a strong tendency towards phase separation. It was also proved that the ferron state becomes unstable in the quantum case for small values of the parameter α, and the electron moves freely over the antiferromagnetic matrix with an effective hopping amplitude $t_{\rm eff} = t/\sqrt{2S+1}$. Our approach to the single-electron problem corresponds to low concentrations of charge carriers (for instance, holes in $La_{1-x}Ca_xMnO_3$). It is intimately related to the recent neutron experiments of Hennion's group on the experimental detection of small ferromagnetic droplets in an antiferromagnetic or canted matrix [15].

Employing the model of ferrons in an insulating matrix [13], we calculated the temperature and field dependences of magnetoresistance in manganites. In our earlier works we also found the expression for the resistance in the framework of the ferron model [13, 14]:

$$\rho = \frac{kT}{128\pi e^2 n^2 \omega_0 l^5} \exp\left(\frac{V}{2kT}\right), \qquad (22)$$

where n is the ferron concentration. In recent experiments performed to measure the resistance in the optimal doping range [17–20] and at temperatures above $T_{\rm Curie}$(the paramagnetic phase), an exponential temperature dependence of resistance of the form $\rho \sim T\exp(V/2kT)$ was discovered, which is in qualitative agreement with formula (22) and confirms our ideas of the nature of charge transfer in the phase-separated domain.

Passing on to the analysis of magnetoresistance, we emphasize once again that it exhibits the following temperature and field dependence in small magnetic fields:

$$MR(H) = \frac{AH^2}{(kT)^2} + \frac{BH^2H_{\rm a}}{(kT)^5}.$$

We note that such a dependence of magnetoresistance on the temperature and the external field was experimentally established by Babushkina's group [16] in $(La_{1-x}Pr_x)_{0.7}Ca_{0.3}MnO_3$ samples (see Fig. 3). In particular, in the strong-anisotropy case it was possible to observe the dependence of magnetoresistance of the form $MR(H) \sim \sim H^2/T^5$ derived theoretically in our work. Note that the expression for magnetoresistance (20) is rather general, depending only slightly on the specific model and operative even in the situation with many-electron ferromagnetic droplets [4, 21].

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Neutron diffraction analysis of doped manganites

V L Aksenov, A M Balagurov, V Yu Pomyakushin

Doped manganites attract considerable attention from both theorists and experimenters in connection with the colossal magnetoresistance (CMR) effect observed in them. From the physical standpoint, manganites are of interest primarily due to an intimate connection between their electron, lattice, and spin subsystems. Recently, it has become evident that an important part in the physics of manganites is played by inhomogeneous states that show themselves, for instance, as the charge ordering of manganese cations, structural and magnetic polarons or a low-temperature phase separation. The employment of neutron diffraction analysis makes it possible to obtain detailed information on the crystal and magnetic structure of manganites, as well as to observe effects related to phase separation. The first neutron diffraction study of doped manganites was made even in the 1950s and reported in Ref. [1], which is regarded as a classical work concerned with the magnetic structure of $La_{1-x}Ca_xMnO_{3-\delta}$ compounds. It was found that there occur several different types of magnetic structures (seven types were proposed by Wollan and Koehler [1]), depending on the doping level and the oxygen stoichiometry, and, moreover, their superposition is also possible. Two mutually exclusive models were discussed for a long time to account for the simultaneous presence of diffraction peaks related to ferromagnetic (FM) and antiferromagnetic (AFM) order in the neutron diffraction patterns, which were observed in $LaMnO_{3-\delta}$ in Ref. [1] and later in other compounds (for instance, in $Pr_{0.7}Ca_{0.3}MnO_3$ [2]): the coexistence of two spatially separated magnetic (FM+AFM) phases, and a homogeneous state with a canted AFM structure. In the case of a mixed state, the question of characteristic inhomogeneity dimen-