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Tunnelling magnetoresistance and $1/f$ noise in phase-separated manganites

A O Sboychakov[1,3], A L Rakhmanov[1], K I Kugel[1], M Yu Kagan[2] and I V Brodsky[2]

[1] Institute for Theoretical and Applied Electrodynamics, Russian Academy of Sciences, Izhorskaya Street 13/19, 125412 Moscow, Russia
[2] Kapitza Institute for Physical Problems, Russian Academy of Sciences, Kosygina Street 2, 117334 Moscow, Russia

E-mail: sboycha@mail.ru

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Abstract
The magnetoresistance and the noise power of non-metallic phase-separated manganites are studied. The material is modelled by a system of small ferromagnetic metallic droplets (magnetic polarons or ferrons) in an insulating matrix. The concentration of metallic phase is assumed to be far from the percolation threshold. The electron tunnelling between ferrons causes the charge transfer in such a system. The magnetoresistance is determined both by the increase in the volume of the metallic phase and by the change in the electron hopping probability. In the framework of such a model, the low-field magnetoresistance is proportional to H^2 and decreases with temperature as T^{-n}, where n can vary from 1 to 5, depending on the parameters of the system. In the high-field limit, the tunnelling magnetoresistance grows exponentially. Different mechanisms of the voltage fluctuations in the system are analysed. The noise spectrum generated by the fluctuations of the number of droplets with extra electrons has a $1/f$ form over a wide frequency range. In the case of strong magnetic anisotropy, the $1/f$ noise can also arise due to fluctuations of the magnetic moments of ferrons. The $1/f$ noise power depends only slightly on the magnetic field in the low field range whereas it can increase as H^6 in the high-field limit.

1. Introduction

Recent theoretical and experimental studies demonstrated clearly that the tendency toward phase separation is of fundamental importance for the physics of manganites and seems to play a key role for the colossal magnetoresistance phenomenon [1–3]. The self-trapping of

[3] Author to whom any correspondence should be addressed.

charge carriers is the most widely discussed type of phase separation, first predicted in the seminal paper of Nagaev [4]. In such phase-separated states, charge carriers are confined within small ferromagnetic metallic droplets (magnetic polarons or ferrons) located in an insulating antiferromagnetic matrix. In the limit of strong Coulomb interaction, each droplet contains one charge carrier in a potential well of ferromagnetically ordered local spins S with a characteristic size of about several lattice constants.

The number of charge carriers, and consequently concentration of the metallic phase x, are related to the content of the divalent dopants. At critical metallic phase concentration $x = x_c$ (for spherical ferrons $x_c \approx 0.15$), the droplets start to overlap and the large metallic clusters arise in the system [2, 5]. We shall consider the case when x is less than x_c and the charge transport is caused by electron tunnelling from one droplet to another. Here we neglect the contribution to the current from the motion of ferrons themselves due to their large effective mass [6].

In the concentration range under study, the magnetoresistance is determined by the increase in the metallic phase volume and by the change in the probability of electron transitions. The latter mechanism gives the main contribution to the magnetoresistance when the system is far from the percolation threshold. The magnetic field dependence of the tunnelling probability is caused both by the change of the mutual orientation of magnetic moments of ferrons and by the variation in ferron size. The mechanisms of magnetoresistance are analysed in section 2.

One of the striking features of manganites in the non-metallic phase is an unusually large magnitude of $1/f$ noise [7, 8]. It is clear that the origin of this noise is closely related to the conductivity mechanisms in the phase-separated state. Several sources of the $1/f$ noise and its magnetic field dependence are discussed in the present paper. First, we generalized the result of [6] for the noise power caused by fluctuations of the number of electrons in ferrons taking into account the effect of applied magnetic field and the spin-dependent tunnelling of charge carriers (see section 3).

The fluctuations of the magnetic moment of ferrons also cause noise in the system. We consider these fluctuations in the limit of low temperatures and strong uniaxial magnetic anisotropy. In this case, each ferron can be treated as a two-level system. The transition of the magnetic moment from one state to another requires overcoming an energy barrier. The spread of the barriers can lead to the $1/f$-type noise spectrum over a wide frequency range. This mechanism is analysed in section 4. The noise power caused by fluctuations of the ferron size is considered in section 5.

The relative contributions of different mechanisms to the magnetoresistance and $1/f$ noise in different temperature and magnetic field ranges are summarized in section 6. In this section we also discuss the experimental implications of the results obtained.

2. Magnetoresistance

The model allowing analysis of the tunnelling electron transport in non-metallic phase-separated manganites was formulated in [6, 9]. In this section we briefly remind ourselves of the basic features of the model and the main results concerning magnetoresistance. Special emphasis will be placed on those details important for the consideration of the noise mechanisms.

Let us consider an insulating antiferromagnetic sample of volume V_s with total number of carriers N. In the ground state every charge carrier is self-trapped in the potential well (droplet) of radius R formed by ferromagnetically ordered local spins. The number of droplets is assumed to be equal to the number of carriers. Due to tunnelling, carriers can pass from one droplet to another. Consequently, droplets with more than one electron are created and some droplets become empty.

The creation of a two-electron droplet is associated with the energy barrier of the order of the Coulomb repulsion V between electrons in this droplet. The value of the ferron radius R can be estimated as 10–20 Å and dielectric constant $\epsilon \approx$ 10–20. Hence we have $V = e^2/\epsilon R \approx$ 0.1–0.2 eV [3]. In the following, we assume that the temperature T is low enough in comparison with this barrier and we neglect the probability of the creation of ferrons with more than two electrons. In addition, electrons in a two-electron droplet can form a state with the total spin either 0 or 1. In the latter case, electrons occupy different energy levels in the potential well. It can be easily shown that the distance between these levels is also about the Coulomb interaction energy (or even exceeds it). Therefore, we consider doubly occupied ferrons only in the state with antiparallel spins of excess electrons.

Let us introduce the notation q for the state of the ferron: $q = 0$ denotes an empty ferron, $q = 1\uparrow$ or $1\downarrow$ stand for a one-electron droplet with an electron spin projection $\sigma/2$ onto the direction of the magnetic moment $\sigma = +1(\uparrow)$ and -1 $(\downarrow)$, and $q = 2$ corresponds to a two-electron ferron. Accounting for the uniaxial magnetic anisotropy, we can write the following expressions for a part of the ferron energy E_q dependent on the droplet radius and effective magnetic fields [3, 9]:

$$E_0 = \frac{4\pi}{3} JzS^2 \frac{R^3}{d^3} - M_0(H\cos\theta + H_a\cos^2\psi), \tag{1}$$

$$E_{1,\lambda} = E_0 + t\frac{\pi^2 d^2}{R^2} - \sigma\bar{J}, \tag{2}$$

$$E_2 = E_0 + 2t\frac{\pi^2 d^2}{R^2} + \frac{e^2}{\epsilon R}, \tag{3}$$

where H_a is the anisotropy field, $M_0 = \mu_B g S 4\pi R^3/3d^3$ is the magnetic moment of a ferron, θ and ψ are the angles between $\boldsymbol{M}_0$ and $\boldsymbol{H}$ and between $\boldsymbol{M}_0$ and the easy axis, respectively, $\sigma/2$ is the projection of the electron spin onto the direction of $\boldsymbol{M}_0$, μ_B is the Bohr magneton and g is the Landé factor. In the following, we consider a single crystal in the applied magnetic field directed at an angle β with respect to the easy axis. Then we have the following relationship:

$$\cos\psi = \cos\theta\cos\beta + \sin\theta\sin\beta\cos\phi, \tag{4}$$

where ϕ is the azimuthal angle characterizing the direction of the magnetic moment $\boldsymbol{M}_0$. We assume, without loss of generality, that the corresponding azimuthal angle of the easy axis is equal to zero.

In equations (1)–(3), the term $t(\pi d/R)^2$ is the energy of an electron in a spherical potential well of radius R and $zJS^2 4\pi R^3/3d^3$ is the energy of the Heisenberg antiferromagnetic exchange between the localized spins S, where t is the hopping integral, J is the constant of the Heisenberg (antiferromagnetic) exchange, z is the number of nearest neighbours and d is the lattice constant. The last term in equation (2) is the energy of the interaction between the confined electron and the effective magnetic field $H_{eff} = \bar{J}/\mu_B$ generated by ferromagnetically ordered localized spins, where $\bar{J}$ is the normalized exchange integral related to the double exchange mechanism. Strictly speaking, in equation (2) we should write $\mu_B|\boldsymbol{H}_{eff} + \boldsymbol{B}|$ instead of $\bar{J}$, where $\boldsymbol{B}$ is the magnetic induction inside the droplet with due account taken of its demagnetization factor. However, we neglected the induction $\boldsymbol{B}$ in comparison with the effective field because H_{eff} is of the order of 100 T [10].

We assume here that the direction of the magnetic moments varies slowly enough and the usual thermodynamic approach is applicable. The radius of a one-electron droplet can be determined by the minimization of energy (2) and we find in the linear approximation with respect to the magnetic field

$$R(H) \cong R_0\left(1 + \frac{b}{5}(H\cos\theta + H_a\cos^2\psi)\right),$$
$$R_0 = d\left(\frac{\pi t}{2zJS^2}\right)^{1/5}, \qquad b = \frac{\mu_B g}{zJS} \sim 100\ \mathrm{T}^{-1}. \tag{5}$$

As was shown in [6], the characteristic tunnelling time for electrons is much smaller than the relaxation time for the local spin system. Hence, we can assume that an empty or two-electron ferron has the same radius (5) as an equilibrium ferron with one charge carrier. We also assume that the total number of ferrons (empty, one-electron and two-electron) remains constant and is equal to N.

In our model, the charge transfer can occur via one of the following four processes [6]:

(i) In the initial state we have two ferrons in the ground state, and after tunnelling in the final state we have an empty ferron and a ferron with two electrons.
(ii) An empty ferron and a two-electron ferron transform into two one-electron ferrons.
(iii) A two-electron and a one-electron ferron exchange their positions by transferring an electron from one ferron to another.
(iv) An empty ferron and a one-electron ferron exchange their positions by transferring an electron from one ferron to another.

Each process is characterized by its tunnelling probability per unit time, which is [9, 11]

$$W(q_1', q_2'; q_1, q_2) = \omega_0 f(\nu)\exp\left(-\frac{r}{l} + \frac{e(\boldsymbol{Er})}{kT} - \frac{E_{q_1'} + E_{q_2'} - E_{q_1} - E_{q_2}}{2kT}\right), \tag{6}$$

$$f(\nu) = \frac{\cosh\left(\frac{\bar{J}\cos\nu}{kT}\right)}{\cosh\left(\frac{\bar{J}}{kT}\right)}, \tag{7}$$

where l and ω_0 are the tunnelling length and the characteristic frequency for the electron motion in the potential well, r and ν are the distance between ferrons and the angle between directions of their magnetic moments; q_1, q_2 and q_1', q_2' denote initial and final states of ferrons involved in the tunnelling process. The angle ν can be expressed through polar $\theta_{1,2}$ and azimuthal $\phi_{1,2}$ angles characterizing the directions of the magnetic moments of ferrons in the states q_1 and q_2:

$$\cos\nu(\theta_1, \phi_1, \theta_2, \phi_2) = \cos\theta_1\cos\theta_2 + \sin\theta_1\sin\theta_2\cos(\phi_1 - \phi_2). \tag{8}$$

The pre-exponential factor $f(\nu)$ in equation (6) is related to the spin-dependent tunnelling and accounts for the different orientations of the magnetic moments of ferrons.

Summing the currents arising from each of the four processes, we find for the conductivity [9]

$$\sigma(H) = \frac{32\pi e^2 l^5\omega_0}{V_s^2\cosh(\bar{J}/kT)kT}\sum_{q_1,q_2}\bar{N}_{q_1}\bar{N}_{q_2} \times\left\langle\cosh\left(\frac{\bar{J}\cos\nu}{kT}\right)\exp\left(\frac{E_{q_1} + E_{q_2} - E_{q_1'} - E_{q_2'}}{2kT}\right)\right\rangle_{q_1,q_2}, \tag{9}$$

where $\bar{N}_q$ are the average occupation numbers of ferrons in state q and $\langle\cdots\rangle_{q_1,q_2}$ stands for the averaging over the direction of the magnetic moments of ferrons in states q_1 and q_2 with probability density

$$P_q(\theta, \phi) = C_q\exp\left(-\frac{E_q(\theta, \phi)}{kT}\right), \qquad \int \mathrm{d}\Omega\, P_q(\theta, \phi) = 1. \tag{10}$$

In equation (9), it is assumed that during the tunnelling process the directions of the magnetic moments of ferrons remain unchanged. The magnetic field dependence of conductivity (9) is related to the variation both in the tunnelling probability and occupation numbers $\bar{N}_q$.

Based on equation (9), we can analyse the behaviour of the magnetoresistance $MR(H) = \sigma(H)/\sigma_0 - 1$. The function $MR(H)$ in the whole magnetic field and temperature ranges could only be found numerically. However, there are several important limiting cases for which it is possible to find an explicit expression for the magnetoresistance. Later, we suppose that the temperature is not too low and obeys the inequalities $\bar{J}$, $M_0 H_a < kT$.

There are several contributions of different origin to the magnetoresistance [9]. In the low-field limit, we can write explicit expressions for all these terms. The first one originates from the magnetic field induced variation in droplet size $R(H)$ given by equation (5). This variation leads to the variation of energies (1)–(3). The corresponding contribution to the magnetoresistance is

$$MR_1(H) \approx \frac{3}{100}\frac{M_0^2 H^2}{(kT)^2}. \tag{11}$$

The second contribution stems from the spin-dependent tunnelling, and for a single crystal with an angle β between the easy axis and the magnetic field it yields

$$MR_2(H) \approx \frac{2}{225}\frac{M_0^3 \bar{J}^2 H_a H^2}{(kT)^5}(\cos^2\beta - 1/3). \tag{12}$$

We assumed above that the ferrons are spherical and the field H_a is determined only by the crystallographic anisotropy. However, it can be shown that the effect of the ferron shape becomes significant due to the demagnetization factor, even at small deviations from sphericity. The shape anisotropy affects the anisotropy field H_a. Let us assume for simplicity that the magnetic anisotropy is uniaxial and H_a is an effective anisotropy field including both the crystallographic and shape anisotropy. Note that $MR \propto H^2/T^5$ was observed over a rather wide range of fields and temperatures for the $(La_{1-x}Pr_x)_{0.7}Ca_{0.3}MnO_3$ system [12].

The third contribution to the magnetoresistance relates to the variation of the metallic phase content. When the system is far from the percolation threshold, the following evident formula is valid for the conductivity of the metal–insulator mixture, $\sigma = \sigma_d(1 + x)$, where σ_d is the conductivity of the insulating matrix and x is the volume fraction of the metallic phase. In our model, the concentration x is related to the droplet parameters by the formula

$$x = \frac{4\pi}{3V_s}\sum_q \bar{N}_q \langle R^3\rangle_q. \tag{13}$$

The corresponding contribution to the magnetoresistance is due to the variation of x with the magnetic field and can be presented in the form

$$MR_3(H) \propto \frac{H^2}{kT}.$$

Our estimations [9] show, however, that $MR_3(H)$ is much smaller than $MR_1(H)+MR_2(H)$ in most cases. Therefore, we can deduce that, in low fields, the magnetoresistance decreases with the temperature as T^{-2} if (11) is dominant or as T^{-5} if (12) is dominant. In general, it behaves as

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

where A and B are some constants.

At high fields (10–20 T and higher), the magnetoresistance grows exponentially [6]:

$$MR(H) \approx \frac{\bar{J}}{kT}\coth(\bar{J}/kT)\exp\left(\frac{VbH}{10kT}\right) - 1, \qquad (14)$$

and its value can be as high as several hundred percent, even far from the percolation threshold.

3. Fluctuations of the occupation numbers

Based on the results of the preceding section, we shall study the voltage fluctuations related to the fluctuations of the conductivity. In this section, we analyse the contribution to the noise power caused by fluctuations of the number of electrons in ferrons, taking into account the effect of applied magnetic field and the spin-dependent tunnelling of charge carriers. The analysis is a generalization of the approach of [6].

Starting from the obvious equality between the voltage and conductivity fluctuations $\delta U/U = -\delta\sigma/\sigma$ and using equation (9) we can relate the voltage noise power to fluctuations of the numbers N_q of different types of ferrons. Note that the fluctuations of N_q are not independent. They obey the simple conservation law $\delta N_2 = \delta N_0 = -\sum_\sigma \delta N_{1,\sigma}/2$. Taking into account that the average numbers $\bar{N}_0 = \bar{N}_2$ are small in comparison with $\bar{N}_{1,\sigma}$ we can write

$$\frac{\langle\delta U^2\rangle_\omega}{\bar{U}^2} = \frac{\langle\delta\sigma^2\rangle_\omega}{\bar{\sigma}^2} = \frac{\langle\delta N_2^2\rangle_\omega}{\bar{N}_2^2}, \qquad (15)$$

where $\bar{U}$ and $\bar{\sigma}$ are the average voltage and conductivity, and the symbol $\langle\ldots\rangle_\omega$ stands for the spectral density of the corresponding variables.

In the tunnelling processes, two-electron and empty ferrons are created and annihilated in pairs. The corresponding relaxation time τ associated with this process is

$$\frac{1}{\tau} = \sum_{\sigma_1,\sigma_2} W(1\sigma_1, 1\sigma_2; 0, 2). \qquad (16)$$

We should neglect the electric field term in expression (6) for probability W.

Based on the relaxation equation $\delta\dot{N}_2 = -\delta N_2/\tau$, and following the same procedure as in [6], we find

$$\langle\delta N_2^2\rangle_\omega = 4\pi\frac{\bar{N}_0\bar{N}_2}{V_s}\left\langle\int_0^\infty \frac{\tau(r)r^2\,\mathrm{d}r}{1+\omega^2\tau^2(r)}\right\rangle_{0,2}, \qquad (17)$$

where subscripts 0 and 2 denote that the averaging is performed over directions of the magnetic moments of the two-electron and empty ferrons. Substituting equation (16) into (17) and performing the integration we find that the noise spectrum has a $1/f$ form in the frequency range determined by the conditions

$$\tilde{\omega}_0\exp(-L_m/l) \ll \omega \ll \tilde{\omega}_0, \qquad \tilde{\omega}_0 = 4\omega_0 e^{V/2kT}\frac{\cosh^2(\bar{J}/2kT)}{\cosh(\bar{J}/kT)}, \qquad (18)$$

where L_m is the minimum sample size. Under conditions (18), we obtain the following expression for the noise spectral density of the voltage noise:

$$\langle\delta U^2\rangle_\omega = \frac{2\pi^2 l^3\bar{U}^2}{\omega V_s}\left\langle\ln^2\left\{\frac{\tilde{\omega}_0}{\omega}\cosh\left[\frac{\bar{J}\cos\nu}{kT}\right]\exp\left[\frac{(\frac{3}{5}+\frac{3\kappa}{10})E_m(\theta_2,\phi_2) - \frac{3}{5}E_m(\theta_0,\phi_0)}{2kT}\right]\right\}\right\rangle_{0,2}, \qquad (19)$$

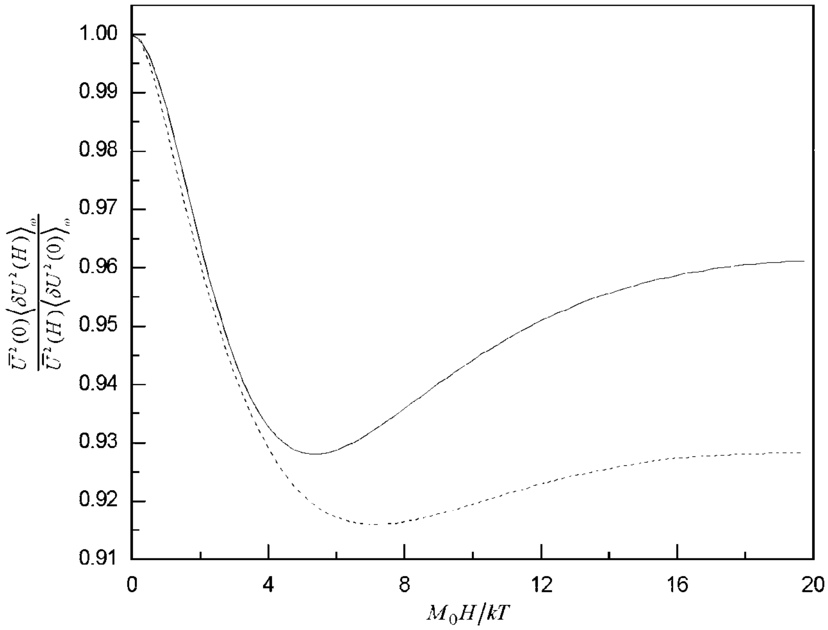

Figure 1. The magnetic field dependence of normalized spectral noise power $\langle \delta U^2 \rangle_\omega / \bar{U}^2$ generated by fluctuations of the occupation numbers N_q calculated at frequencies $\omega/\tilde{\omega}_0 = 10^{-15}$ (full curve) and $\omega/\tilde{\omega}_0 = 10^{-10}$ (broken curve). These frequencies are of the order of 10^2 and 10^7 s^{-1}, respectively. The parameter values used are $M_0H_a/kT = 3$, $\bar{J}/kT = 2$, $\beta = \pi/3$, $\kappa = 0.2$.

where $E_m(\theta, \phi) = -M_0(H\cos\theta + H_a\cos^2\psi(\theta, \phi))$ and $\kappa = VR_0^2/(\pi^2 td^2) \lesssim 1$ is the ratio of the Coulomb energy to the energy of an electron localized in a spherical potential well. We use here the relationships (1)–(3) for the ferron energies and equation (5) for the ferron radius $R(H)$.

The averaging over the directions of the magnetic moments of ferrons in equation (19) can only be performed numerically. The magnetic field dependence of the noise power (19) is shown in figure 1. We see that the normalized noise spectrum varies slowly with the field. The noise power decreases at low magnetic field owing to the alignment of the magnetic moments of ferrons related to the term with hyperbolic cosine in equation (19). The further increase in the noise power stems from the lowering of the Coulomb energy with the growth of the ferron radii in higher magnetic fields. Note that the form of normalized noise spectrum depends only slightly on the angle β.

4. Fluctuations of magnetic moments

In this section, we consider another mechanism of the conductivity noise related to fluctuations of the magnetic moments of ferrons. We suppose here that a ferron is an ellipsoid of rotation. Note that ferrons in layered manganites can have this shape [3, 13]. In this case, the magnetic moment of the ferron is $M_0 = 4\pi\mu_B gSab^2/3d^3$, where $a > b = c$ are the ellipsoid axes. We consider a single crystal and assume that the directions of the easy axes of ferrons are parallel each other. Moreover, the ratio a/b is assumed to be constant. We shall use here equation (9) for the conductivity, neglecting the effect of the droplet shape on $\sigma(H)$.

We assume that the main contribution to the anisotropy field H_a comes from the shape anisotropy of ferrons. Then, the field H_a can be written in the form

$$H_a = 2\pi \frac{\mu_B g S}{d^3}(\mathcal{N}_b - \mathcal{N}_a) = 2\pi I_s(\mathcal{N}_b - \mathcal{N}_a), \tag{20}$$

where $\mathcal{N}_b$ and $\mathcal{N}_a$ are the demagnetization factors along axes b and a, and I_s is the saturation magnetization. Let us denote $\xi = 2(\mathcal{N}_b - \mathcal{N}_a)$. This variable has the maximum value $\xi_{max} \approx 1$ at $a \gg b$, which corresponds to the anisotropy field $H_a^{max} \approx \pi g S \mu_B/d^3$. At characteristic values of the parameters $g = 2$, $S = 2$ and $d = 0.4$ nm, $H_a \approx 1.8$ kOe.

In this section, we consider the case of low temperatures when $\exp(-\pi M_0 I_s \xi/kT) \ll 1$ and we can assume that the magnetic moment of a ferron only has one or two possible directions (depending on the applied field), which are determined by the minimization of the energy of a ferromagnetic ellipsoid in magnetic field H:

$$E_m = -M_0(\pi I_s \xi \cos^2\psi + H\cos\theta), \tag{21}$$

where

$$\cos\psi = \cos\theta\cos\beta + \sin\theta\sin\beta\cos\phi$$

and β ($0 < \beta < \pi/2$) is the angle between the direction of the long axis of the ferron and the applied magnetic field. The energy E_m has two minima if $H < H_0(\xi, \beta) \approx \pi I_s \xi$ and a single minimum at higher fields.

The magnetic moment of the ferron lies evidently in the plane of the field $\boldsymbol{H}$ and the easy axis. At low fields, the deeper minimum corresponds to an azimuthal angle $\phi^+ = 0$, whereas the second minimum corresponds to $\phi^- = \pi$. The energies of these minima are

$$E_m^+(H, \xi, \beta) = -M_0(\pi I_s \xi \cos^2(\theta^+ - \beta) + H\cos\theta^+), \tag{22}$$

$$E_m^-(H, \xi, \beta) = -M_0(\pi I_s \xi \cos^2(\theta^- + \beta) + H\cos\theta^-), \tag{23}$$

where angles $\theta^\pm$ should be determined by the minimization of these expressions. At $H > H_0$, the single minimum of the energy E_m is given by equation (22). We consider the case of sufficiently low fields when ferrons have two minima of energy E_m.

At low fields, we can assume that the ferron's size does not depend on H. In this case, the probability densities (10) do not depend on the ferron state q and can be written as

$$P(\theta, \phi) = C\exp\left(-\frac{E_m(\theta, \phi)}{kT}\right). \tag{24}$$

It follows from definitions (1)–(3) that the argument of the exponential in equation (9) is independent of the angular variables and can be taken out of the angular brackets. Then

$$\left\langle \cosh\left(\frac{\bar{J}\cos\nu}{kT}\right)\right\rangle = \int \mathrm{d}\Omega_1\,\mathrm{d}\Omega_2\, P(\theta_1, \phi_1)P(\theta_2, \phi_2)\cosh\left(\frac{\bar{J}\cos\nu}{kT}\right). \tag{25}$$

In our approach, we should replace the integration over the solid angle in equation (25) by the summation over two states of the ferron with its magnetic moment directed parallel (+) and antiparallel (−) to the applied magnetic field. In the thermal equilibrium, the corresponding probabilities can be written as

$$P_S = \frac{\mathrm{e}^{-E_m^S/kT}}{\sum_{S'} \mathrm{e}^{-E_m^{S'}/kT}},\ S, S' = \pm. \tag{26}$$

The number of ferrons in state S equals $\bar{N}_S = NP_S$. When the ferron has only one minimum of the energy E_m we have $P_+ = 1$, $P_- = 0$.

The noise discussed in this section is related to the fluctuations δN_S of numbers N_S. It is convenient to introduce the notation $\delta P_S = \delta N_S/N$ since N is constant. The fluctuations of the magnetic moment and the number of electrons in the ferrons are not correlated since they are characterized by different timescales. The relaxation time of the electron numbers is $\tau_e = \exp(-V/2kT)/8\pi l^3 n\omega_0 \sim 10^{-17}$ s [6], whereas the estimation of the relaxation time for the magnetic moment gives values much higher than τ_e (see the discussion below). Therefore, we can assume that the N_q are equal to their equilibrium values.

Substituting probabilities equation (26) into (25) and taking into account that $\delta P_+ + \delta P_- = 0$, we obtain for the conductivity fluctuations

$$\frac{\delta\sigma}{\bar{\sigma}} \equiv 2\delta P_- A(H, \xi, \beta), \tag{27}$$

where

$$A(H, \xi, \beta) \cong \frac{1}{\cosh\left(\frac{\bar{J}}{kT}\right)} \sum_{S_1} P_{S_1} \left[\cosh\left(\frac{\bar{J}}{kT} \cos\nu(\theta^-, \phi^-, \theta_1^{S_1}, \phi_1^{S_1})\right) - \cosh\left(\frac{\bar{J}}{kT} \cos\nu(\theta^+, \phi^+, \theta_1^{S_1}, \phi_1^{S_1})\right)\right]. \tag{28}$$

In equations (27) and (28) we take $\bar{\sigma}(H) = \bar{\sigma}(0)$ since we discuss here the low-field limit. Note that at the sufficiently high fields $H > H_0(\xi, \beta)$, we have only one free energy minimum and the conductivity fluctuations of this type vanish.

The relaxation of the magnetic moment of the droplet from one S state to another is determined by overcoming the minimum energy barrier:

$$\Delta^{\pm}(H, \xi, \beta) = E_m^0(H, \xi, \beta) - E_m^{\pm}(H, \xi, \beta), \tag{29}$$

where E_m^0 is the energy corresponding to the lowest saddle point of the function $E_m(\theta, \phi)$. The relaxation time is then

$$\frac{1}{\tau(H, \xi, \beta)} = \Omega_0(\xi)\left[\exp\left(-\frac{\Delta^+(H, \xi, \beta)}{kT}\right) + \exp\left(-\frac{\Delta^-(H, \xi, \beta)}{kT}\right)\right], \tag{30}$$

where [15, 16]

$$\Omega_0(\xi) = \frac{\mu_B \pi g S I_s \xi}{\hbar}\sqrt{\frac{M_0 I_s \xi}{kT}}. \tag{31}$$

Substituting the characteristic values of the parameters involved in the last two equations, we find that $\tau(H, \xi, \beta) > 10^{-10}$–$10^{-11}$ s, which is much higher than the relaxation time τ_e for the electron numbers.

Assuming that there is no correlation between the fluctuations of magnetic moments of different ferrons, we obtain the following expression for the noise power from equation (28):

$$\frac{\langle\delta U^2\rangle_\omega}{\bar{U}^2} = \frac{A^2(H, \xi, \beta)}{N\cosh^2\left(\frac{E_m^+(H,\xi,\beta) - E_m^-(H,\xi,\beta)}{2kT}\right)} \frac{2\tau(H, \xi, \beta)}{1 + \omega^2\tau^2(H, \xi, \beta)}. \tag{32}$$

In the $H \ll \pi I_s \xi$ limit, we can expand equation (32) in series and find the explicit formula for the noise power:

$$\frac{\langle\delta U^2\rangle_\omega}{\bar{U}^2} = \frac{H^4 \sin^4\beta}{(\pi I_s \xi)^4} \frac{\tanh^2(\bar{J}/kT)\tanh^2(M_0 H\cos\beta/kT)}{N\cosh^2(M_0 H\cos\beta/kT)} \frac{2\tau(H, \xi, \beta)}{1 + \omega^2\tau^2(H, \xi, \beta)}, \tag{33}$$

and

$$\tau(H, \xi, \beta) = (2\Omega_0(\xi))^{-1} \frac{\exp[-M_0(\pi I_s \xi + H\sin\beta)/kT]}{\cosh(M_0 H\cos\beta/kT)}. \tag{34}$$

We see that, in contrast to equation (19), the noise spectrum does not have the $1/f$ form and depends strongly on the magnetic field and temperature. It follows from equation (33) that

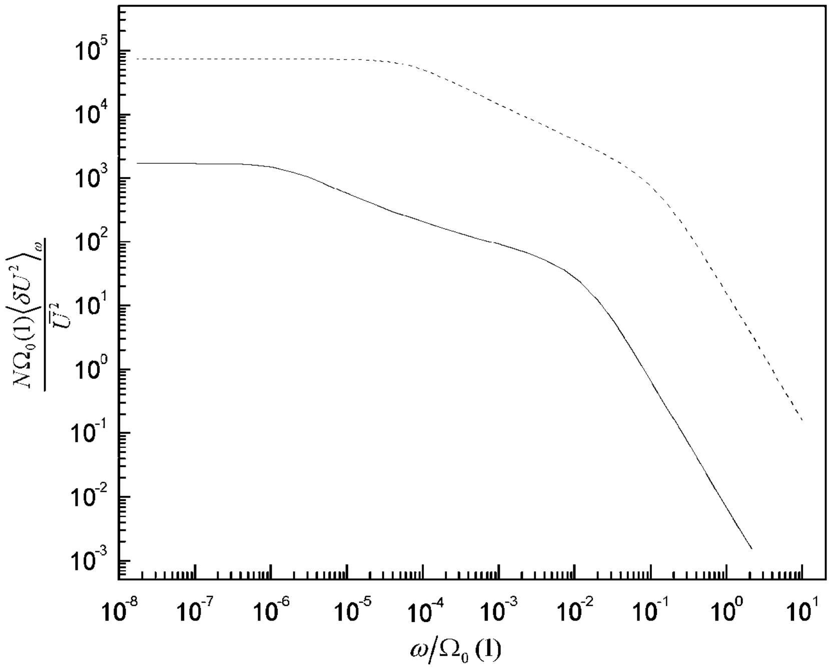

Figure 2. The spectrum of the noise generated by fluctuations of the magnetic moments of ferrons calculated at $H/(\pi I_s) = 0.05$ (full curve) and $H/(\pi I_s) = 0.5$ (broken line). The parameter values used are $M_0\pi I_s/kT = 15$, $\bar{J}/kT = 5$, $\xi_1 = 0.2$, $\xi_2 = 0.9$. At low frequencies $\omega/\Omega_0(1) \lesssim 10^{-6}$–$10^{-5}$, the noise power is constant, and at high frequencies $\omega/\Omega_0(1) \gtrsim 10^{-2}$–$10^{-1}$ it behaves as $1/\omega^2$. In the intermediate range, the frequency dependence of the noise is described approximately by the $1/\omega$ law. This region shifts to high frequencies when the field is increased.

$\delta\langle U^2\rangle_\omega \propto H^6/T^2$ if $M_0 H \cos\beta \ll kT$. In relatively high fields, when $M_0 H \cos\beta > kT$ (but still $H < \pi I_s \xi$), the noise power decreases exponentially. The steep growth of the noise at low fields is caused by the splitting between two energy minima $E_m^{\pm}$ for ferrons. The further decrease of the noise at higher fields is related to the suppression of the transitions to the state with the ferron magnetic moment antiparallel to the applied field. It also follows from equation (32) that the noise power vanishes when the applied field is parallel or perpendicular to the easy axis.

Formula (32) is valid for a single crystal. For polycrystals, we should make an averaging procedure in equation (32). We assume here that the parameter ξ is distributed in the range $\xi_1 < \xi < \xi_2$ and the directions of the easy axes of crystallites are randomly distributed. In this case, the formula for the noise power of the polycrystal can be obtained by averaging equation (32) over ξ and β with probability density $w(\xi)\sin\beta$. This procedure is adequate if the resistance of crystallites does not differ significantly from each other and the frequency ω is much less than the inverse value of the tunnelling time. The scatter of the parameters ξ and β leads to the $1/f$ form of the noise spectrum in the frequency range

$$\frac{1}{\tau(H, \xi_2, \beta_{max})} < \omega < \frac{1}{\tau(H, \xi_1, \beta_{min})}, \tag{35}$$

where β_{max}, and β_{min} are the angles corresponding to the maximum and minimum of the time τ with given ξ. The results of the numerical calculations for the case $w(\xi) = \text{constant}$ are shown in figure 2.

Note that, in addition to the noise mechanism related to the transition of the magnetic moment between two energy minima, there also exists a noise caused by the small oscillations of the ferron magnetic moment near a single equilibrium direction. However, the calculations show that the noise power related to the latter mechanism is small compared to the former one.

5. Fluctuation of ferron size

In this section we calculate the noise power generated by the fluctuations of the ferron size. In this mechanism, the main contribution to the voltage noise evidently comes from the variation of the Coulomb barrier $V = e^2/\epsilon R$. Therefore, we can neglect the effects of the magnetic field and electron spin on the ferron energy. For simplicity, we consider here spherical ferrons. In this case, the conductivity depends on a ferron radius as $\sigma \propto \exp(-e^2/(2\epsilon RkT))$ [6] and we easily obtain

$$\frac{\langle\delta\sigma^2\rangle_\omega}{\bar{\sigma}^2} = \left(\frac{V}{2kT}\right)^2 \frac{\overline{\Delta R^2}}{NR_0^2} \frac{2\tau_R}{1+\omega^2\tau_R^2}, \tag{36}$$

where τ_R is of the order of the characteristic relaxation time of magnons and $\tau_R \sim 10^{-12}$–10^{-13} s. The standard deviation $\overline{\Delta R^2}$ can be found in the following way. The energy of the ferron in the ground state $E_1(R) = t\pi^2 d^2/R^2 + 4\pi JzS^2R^3/(3d^3)$ is found from equation (2) where we should neglect the terms depending on the magnetic fields and electron spin. The mean value of R^n can be written in the form

$$\overline{R^n} = \frac{\int_0^\infty \exp\{-E_1(R)/kT\}R^{n+2}\,\mathrm{d}R}{\int_0^\infty \exp\{-E_1(R)/kT\}R^2\,\mathrm{d}R}. \tag{37}$$

The calculation of integral (37) with $n = 1$ and 2 using a saddle-point approximation gives the following relationship for the quadratic deviation:

$$\overline{\Delta R^2} = \frac{R_0^4 kT}{10\pi^2 td^2}, \tag{38}$$

and finally we obtain

$$\langle\delta U^2\rangle_\omega = \bar{U}^2 \frac{V}{kT} \frac{\kappa}{20N} \frac{\tau_R}{1+\omega^2\tau_R^2}. \tag{39}$$

In the case of one-electron ferrons, there is no reasonable mechanism of a large scatter in τ_R as well as in the ferron sizes. Therefore, the fluctuations of ferron sizes cannot be a source of the $1/f$ noise in the framework of our model. However, in real systems, the sizes of the ferromagnetic regions can undergo wide range variations, which could give an additional contribution to the $1/f$ noise.

6. Conclusions

The magnetoresistance and noise power for the non-metallic phase-separated manganites were studied in the framework of the model of small ferromagnetic metallic droplets embedded in the insulating matrix. The concentration of droplets was supposed to be far from the percolation limit. The charge transfer in the system was assumed to be due to the tunnelling of electrons from one droplet to another.

A general expression for the conductivity in an external magnetic field was derived. The low-field magnetoresistance $MR(H)$ is a quadratic function of H whereas the magnetoresistance exhibits an exponential growth with H in the high-field limit.

At low fields, the temperature dependence of the magnetoresistance has the form

$$MR(H) = \frac{AH^2}{T^2} + \frac{BH^2}{T^5},$$

where the first term is due to the variation of droplet size caused by the magnetic field and the second term originates from the spin-assisted electron tunnelling. In the case of the systems with relatively large ferromagnetic droplets and magnetic anisotropy, the second term becomes dominant and the magnetoresistance decreases with temperature as T^{-5}. Such behaviour was indeed observed for the $(La_{1-x}Pr_x)_{0.7}Ca_{0.3}MnO_3$ system [12]. The results obtained can explain the observed large (up to 100% and higher) magnetoresistance of manganites in the paramagnetic phase.

The available experiments for phase-separated manganites demonstrated strong $1/f$ noise in a wide frequency range [7, 8]. There are two main sources of $1/f$ noise in our model: the fluctuations of the number of electrons in the droplet and the fluctuations of the magnetic moments. To characterize the $1/f$ noise, it is convenient to introduce the variable $\alpha = V_s\omega\langle\delta U^2\rangle_\omega/\bar{U}^2$ [14].

First, let us estimate the parameter α_N, i.e. the contribution to the noise due to fluctuations of the electron occupation numbers. Since the Coulomb energy $V \sim 0.1$ eV and the parameter $\hbar\omega_0$ (which is about the energy of electron localization) has the same order of magnitude, we find that $\tilde{\omega}_0 \sim 10^{17}$ s^{-1}. At the frequency range 1–10^6 Hz and $H < 20$–30 T, α_N is nearly independent of the magnetic field (in agreement with the experiment [8]) and is close to $\alpha_N \approx 2\pi^2 l^3 \ln^2(\tilde{\omega}_0/\omega)$ obtained in [6]. It is natural to assume that the tunnelling length is about a ferron radius $R_0 \sim 1$ nm. As a result we get the following estimate: $\alpha_N \sim 10^{-17}$–10^{-16} cm^3. This value of $1/f$ noise power is higher by several orders of magnitude than those typical for semiconductors [14].

The parameter α_M related to the $1/f$ noise generated by fluctuations of the magnetic moments of ferrons strongly depends on the magnetic field and temperature. Using equations (33) and (34), it can be shown that, at relatively low field ($H \ll \pi I_s\xi$), $\alpha_M \propto M_0^3 H^6/T^m$, where $m = 5$ if $kT > \bar{J}$ and $m = 3$ if $kT < \bar{J}$. In the framework of our model, it turns out that $\alpha_M \ll \alpha_N$ at any realistic parameter values. However, it is clear from the analysis given in section 4 that the results for this type of noise are applicable not only to the case of one-electron ferrons but are also valid for ferromagnetic droplets of larger dimensions. Since the power of this noise is proportional to the cube of the droplet magnetic moment (that is, to the cube of the droplet volume) the discussed mechanism can give an appreciable contribution to the total $1/f$ noise power in real materials (especially at low temperatures and high fields) where ferromagnetic regions of different sizes can exist.

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