New Correlated Model of Colossal Magnetoresistive Manganese Oxides

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A new minimal model is constructed for the doped manganese oxides which exhibit colossal magnetoresistance, involving a broad spin-majority conduction band as well as nearly localized spin-minority electron states. A simple mean-field analysis yields a temperature-dependent hybridized band structure with suppressed carrier weight at the Fermi level. Spin stiffness is complex, indicating strong spin-wave damping. Further investigations are needed to verify the relevance of the proposed model.

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The unusual properties of doped manganese oxides exhibiting colossal magnetoresistance (CMR) [1] are not yet understood theoretically, and the problem of formulating a suitable microscopic model remains open. These properties include the CMR phenomenon and metal-insulator transition, which are, in turn, intimately related to the temperature-induced variation of the electron density of states (''pseudogap'' [2]) or of the effective carrier number [3]. This implies that, alongside double-exchange ferromagnetism, the effects of electron-electron interaction play a key role and, quite possibly, are responsible for CMR itself. Indeed, the band structure calculations [4] suggest that the on-site Hubbard repulsion U is the largest energy in the problem. Its effects are typically considered within the one- or two-orbital model whereby the strongly correlated behavior is induced by an interaction between the two spin-majority electronic e_g states on-site, or between the two spin components of a single e_g band [5].

It should be noted that the experimental data [6] indicate the presence of *spin-minority* electrons near the Fermi level even in the low-temperature ferromagnetic state. This agrees with the band structure calculations [4,7], suggesting that a narrow spin-minority band lies close to the Fermi energy. Since both the localized t_{2g} and itinerant e_g states originate from the same d shell of a Mn ion, and therefore are characterized by approximately the same value of Hund's rule splitting J_H, it is clear [see Fig. 1(a)] that these spin-minority electrons populate the spin-down t_{2g} (localized) states [4]. This is further corroborated by the studies of $La_{1-x}Ce_xMnO_3$ (with $1+x$ conduction electrons per formula unit) [8], which show both the spin-minority character of the carriers and a large overall increase in the resistivity (in comparison with the usual $1-x$ electron case, e.g., $La_{1-x}Ca_xMnO_3$), consistent with the extra x electrons going into the spin-down t_{2g} states. Yet, while ''two-fluid'' models involving itinerant and localized states were suggested earlier by both experimentalists and theorists [9–11], the appealing possibility (mentioned in Ref. [12]) of these states having antialigned spins has not been addressed theoretically.

Furthermore, we note that electronic properties of a model where the orbital degree of freedom is taken into account are strongly coupled to lattice dynamics via the Jahn-Teller effect [1], which results in additional splitting of both e_g and t_{2g} levels. However, CMR occurs in a broad class of bulk systems [three-dimensional (3D) perovskites and quasi-two-dimensional (quasi-2D) bilayered] and thin films of varying composition, and is presumably always due to the same physical mechanism. It is therefore worthwhile to consider a model with fewer orbitals, which still captures some of the important *intra-atomic* physics of a Mn ion, before pursuing more complicated (and probably more material-specific) options [13].

In the present Letter, we introduce such a simplified description and proceed with a simple mean-field analysis. While finer theoretical tools are required to gain a fuller picture, qualitatively our results for electron dispersion and magnetic properties appear very encouraging.

We consider a model involving a large spin $\vec{S}$ and two conduction-electron orbitals (broadened and nearly localized) at each Mn site [cf. Figure 1(b)]:

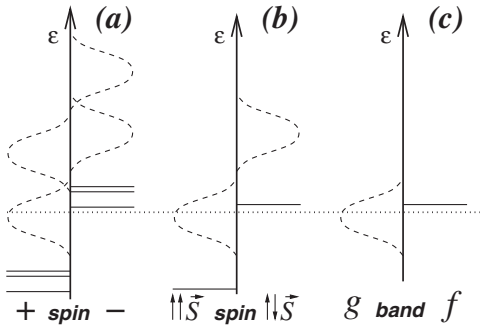

FIG. 1. (a) Crystal field splits the d-electron states of a Mn ion into t_{2g} (solid lines) and e_g (dashed lines); these split further due to a lattice distortion. Spin degeneracy is broken due to Hund's rule, and the ''spin-plus'' t_{2g} electrons form the local spin 3/2. Chemical potential (dotted line) lies within an e_g band, with a ''spin-minus'' t_{2g} level nearby. (b) A simplified model of Eq. (1), with the spin quantization axis fixed by a local spin S. (c) Relevant orbitals in the limit $J_H \to \infty$ (see text); adding an itinerant (localized) electron increases (reduces) the total on-site spin of $S + 1/2$ by $1/2$.

$$\mathcal{H} = -\frac{t}{2}\sum_{\langle i,j\rangle,\alpha}(c^{\dagger}_{i\alpha}c_{j\alpha} + c^{\dagger}_{j\alpha}c_{i\alpha}) + E^{(0)}_{d}\sum_{i,\alpha} d^{\dagger}_{i\alpha}d_{i\alpha} - \frac{J_H}{S}\sum_i \vec{S}_i\vec{\sigma}_i + \frac{J}{S^2}\sum_{\langle i,j\rangle}\vec{S}_i\vec{S}_j - \frac{H}{S}\sum_i (S^z_i + \sigma^z_i) + U\sum_i\left(c^{\dagger}_{i\uparrow}c^{\dagger}_{i\downarrow}c_{i\downarrow}c_{i\uparrow} + d^{\dagger}_{i\uparrow}d^{\dagger}_{i\downarrow}d_{i\downarrow}d_{i\uparrow} + \sum_{\alpha,\beta} c^{\dagger}_{i\alpha}d^{\dagger}_{i\beta}d_{i\beta}c_{i\alpha}\right) - \frac{V}{2}\sum_{\langle i,j\rangle,\alpha}(c^{\dagger}_{i\alpha}d_{j\alpha} + c^{\dagger}_{j\alpha}d_{i\alpha} + d^{\dagger}_{i\alpha}c_{j\alpha} + d^{\dagger}_{j\alpha}c_{i\alpha}). \quad (1)$$

Here, the operators $c_{i\alpha}$ ($d_{i\alpha}$) annihilate an e_g (t_{2g}) conduction electron of spin $\alpha = \uparrow, \downarrow$ (in the laboratory frame) at a site i of a square (or simple cubic) lattice. Localized spins S_i originate from the remaining two t_{2g} electrons; hence, in reality $S = 1$. They interact via superexchange J and are also coupled to the spins of conduction electrons on-site, $\vec{\sigma}_i = \frac{1}{2}\sum_{\alpha,\beta}\vec{\sigma}_{\alpha\beta}(c^{\dagger}_{i\alpha}c_{i\beta} + d^{\dagger}_{i\alpha}d_{i\beta})$ (where $\vec{\sigma}_{\alpha\beta}$ are the Pauli matrices) via a strong ferromagnetic Hund's rule exchange J_H; the external magnetic field H is applied along the z axis. Owing to the fact that the (electron) coordinate operator is not diagonal in the band index, there is a hybridization V between the t_{2g} and e_g states; $E^{(0)}_d$ is the (bare) energy of the t_{2g} electrons. Direct hopping between t_{2g} states on different sites is assumed to be negligible, while the nearest-neighbor hopping t between the e_g states will be used as an energy unit, $t = 1$.

We construct the spin-wave expansion for the model (1), keeping terms up to first order in $1/S$. In the spirit of Ref. [14], it is expedient to introduce a new basis of electron states on each site according to

$$c_{\uparrow} \approx g_{\uparrow} - \frac{1}{\sqrt{2S}}g_{\downarrow}\beta^{\dagger} - \frac{1}{4S}(g_{\uparrow}\beta^{\dagger}\beta + g^{\dagger}_{\downarrow}g_{\downarrow}g_{\uparrow} + f^{\dagger}_{\downarrow}g_{\downarrow}f_{\uparrow}),$$

$$c_{\downarrow} \approx g_{\downarrow} + \frac{1}{\sqrt{2S}}g_{\uparrow}\beta - \frac{1}{4S}(g_{\downarrow} + g_{\downarrow}\beta^{\dagger}\beta - g^{\dagger}_{\uparrow}g_{\uparrow}g_{\downarrow} - f^{\dagger}_{\uparrow}g_{\uparrow}f_{\downarrow})$$

(expressions for $d_{\uparrow}$ and $d_{\downarrow}$ are obtained by substituting $g_{\alpha} \leftrightarrow f_{\alpha}$). Operators $g_{i\uparrow}$ and $f_{i\uparrow}$ ($g_{i\downarrow}$ and $f_{i\downarrow}$) correspond, respectively, to the e_g and t_{2g} electrons with a spin parallel (antiparallel) to the total spin on-site, $\vec{\mathcal{T}}_i = \vec{S}_i + \vec{\sigma}_i$, whose vibrations are annihilated by a Holstein-Primakoff magnon operator β_i. The Hund's rule term takes form

$$-\frac{J_H}{S}\vec{S}_i\vec{\sigma}_i = -\frac{J_H}{2}\left\{g^{\dagger}_{i\uparrow}g_{i\uparrow} + f^{\dagger}_{i\uparrow}f_{i\uparrow} - \frac{S+1}{S}(g^{\dagger}_{i\downarrow}g_{i\downarrow} + f^{\dagger}_{i\downarrow}f_{i\downarrow}) + \frac{1}{S}[g^{\dagger}_{i\uparrow}g^{\dagger}_{i\downarrow}g_{i\downarrow}g_{i\uparrow} - g^{\dagger}_{i\uparrow}f^{\dagger}_{i\downarrow}g_{i\downarrow}f_{i\uparrow} + (f \leftrightarrow g)]\right\}.$$

Finally, the Holstein-Primakoff operators a_i of the original spins $\vec{S}_i$ are expressed as

$$a \approx \beta - \frac{g^{\dagger}_{\uparrow}g_{\downarrow} + f^{\dagger}_{\uparrow}f_{\downarrow}}{\sqrt{2S}} - \frac{g^{\dagger}_{\uparrow}g_{\uparrow} + f^{\dagger}_{\uparrow}f_{\uparrow} - g^{\dagger}_{\downarrow}g_{\downarrow} - f^{\dagger}_{\downarrow}f_{\downarrow}}{4S}\beta.$$

We next substitute these expressions into the Hamiltonian (1) and take the limit of $J_H \to \infty$ while keeping $E_d = E^{(0)}_d + J_H$ constant. Now, if the chemical potential, denoted $\mu + U - J_H/2$, lies within the spin-up e_g band, then the spin-down e_g band is completely empty, and the spin-up t_{2g} band completely filled. This is precisely the case of interest to us, containing the effects of the Coulomb repulsion between the spin-up e_g and spin-down t_{2g} electrons in the presence of a localized spin $S + 1/2$ [Fig. 1(c)]. Henceforth, we drop all the terms containing the operators $g_{\downarrow}$ and $f_{\uparrow}$, and suppress the spin index of remaining fermion operators.

Assuming the ferromagnetic ground state, the Hamiltonian takes the form $\mathcal{H}_e + \mathcal{H}_h + \mathcal{H}_m$ with the electronic and (magnon-assisted) hybridization terms,

$$\mathcal{H}_e = \sum_{\vec{k}}(\epsilon_{\vec{k}} - \mu)g^{\dagger}_{\vec{k}}g_{\vec{k}} + (E_d - \mu)\sum_j f^{\dagger}_j f_j + \frac{U}{N}\sum_{\vec{k},\vec{k}',j} e^{i(\vec{k}'-\vec{k})\vec{R}_j}g^{\dagger}_{\vec{k}}f^{\dagger}_j f_j g_{\vec{k}'}, \quad (2)$$

$$\mathcal{H}_h = -\frac{V}{\sqrt{2SN}}\sum_{\vec{k},\vec{q},j}(\epsilon_{\vec{k}} - \epsilon_{\vec{k}+\vec{q}})e^{-i(\vec{k}+\vec{q})\vec{R}_j}g^{\dagger}_{\vec{k}}f_j\beta^{\dagger}_{\vec{q}} + \text{H.c.} \quad (3)$$

Here, N is the number of lattice sites, $\epsilon_{\vec{k}} = -\cos k_x - \cos k_y(-\cos k_z)$ is the tight-binding dispersion law in two (three) dimensions, and $\vec{R}_j$ is the radius vector of site j. It is assumed that the site basis is more appropriate for describing the narrow-band fermions f_j, easily localized by fluctuations or disorder. The term $\mathcal{H}_m$, which is of order $1/S$, contains Zeeman electron energy shifts, magnon dynamics, and double-exchange band-narrowing effects:

$$\mathcal{H}_m = \frac{H}{2S}\left(\sum_j f^{\dagger}_j f_j - \sum_{\vec{k}} g^{\dagger}_{\vec{k}}g_{\vec{k}}\right) + \frac{1}{S}\sum_{\vec{k}}[H - 2J(\epsilon_{\vec{k}} + d)] \times \beta^{\dagger}_{\vec{k}}\beta_{\vec{k}} + \frac{1}{4SN}\sum_{1\div 4}{}'(2\epsilon_{2+4} - \epsilon_1 - \epsilon_2)g^{\dagger}_1 g_2\beta^{\dagger}_3\beta_4. \quad (4)$$

Here, momentum-conserving summation is denoted by Σ', and d is the dimensionality of the system (2 or 3).

The electronic term, $\mathcal{H}_e$, is the familiar Falikov-Kimball model. The rich physics contained therein [15] crucially depends on the presence (and nature) of the interband hybridization. The form of our $\mathcal{H}_h$, Eq. (3), is dictated by spin conservation: electron transfer between the two opposite-spin bands must be accompanied by magnon creation or annihilation, $\beta^{\dagger}_{\vec{q}}$ or $\beta_{\vec{q}}$. Such transfers require misalignment of spins $\vec{S}_i$ on neighboring sites; hence, the hybridization matrix element vanishes at $q \to 0$, underlining the importance of short-wavelength processes. The latter feature appears promising in the context of CMR compounds, where the unusual short-range correlations are reflected in the electronic and magnetic [16,17] properties. Importantly, magnetic field H affects the carriers both via double-exchange mechanism and by changing the energy difference between localized and itinerant states [see Eq. (4)].

We shall be interested in the regime characterized by nonzero values of *both* fermion occupancies n^g and n^f (the latter assumed independent on $\vec{R}_i$),

$$n^f = \langle f_i^\dagger f_i \rangle, \qquad n^g = \frac{1}{N}\sum_{\vec{k}} n^g_{\vec{k}}, \qquad n^g_{\vec{k}} = \langle g^\dagger_{\vec{k}} g_{\vec{k}} \rangle. \quad (5)$$

In addition, there also arises an off-diagonal average,

$$\langle f_j^\dagger g_{\vec{k}} \beta_{\vec{q}} \rangle \equiv \langle f^\dagger g_{\vec{k}} \beta_{\vec{q}} \rangle e^{-i(\vec{k}+\vec{q})\vec{R}_j} \quad (6)$$

(here and below, we omit the site indexes of operators f_j once the j-dependent exponent has been factored out).

In order to clarify the basic physics contained in our model, Eqs. (2)–(4), we will now proceed with a mean-field analysis of it. Here, we focus on the simplest self-consistent scheme, allowing only for average values (5) and (6) and for a nonzero magnon occupancy. While actual validity of this approach is probably restricted to the intermediate temperature range (on the scale of the Curie temperature) and moderate values of U (see below), it offers important guidance for future investigations. Mean-field equations can be found in a standard way by decoupling the equations of motion for the appropriate retarded Green's functions, expressing the latter as

$$\langle\langle f f^\dagger \rangle\rangle \equiv \langle\langle f_j f_j^\dagger \rangle\rangle = \left\{ \omega - \tilde{E}_d + \mu + i0 - \frac{V}{2SN^2} \sum_{\vec{k},\vec{q}} \frac{V(1+N_{\vec{q}} - n^g_{\vec{k}})(\epsilon_{\vec{k}} - \epsilon_{\vec{k}+\vec{q}})^2 + U\Phi_{\vec{q}}(\epsilon_{\vec{k}} - \epsilon_{\vec{k}+\vec{q}})}{\omega + i0 - \tilde{\epsilon}_{\vec{k}} + \mu} \right\}^{-1}, \quad (7)$$

$$\langle\langle g_{\vec{k}} g^\dagger_{\vec{k}} \rangle\rangle = \left\{ \omega - \tilde{\epsilon}_{\vec{k}} + \mu + i0 - \frac{V}{2SN} \frac{\sum_{\vec{q}} [V(N_{\vec{q}} + n^f)(\epsilon_{\vec{k}} - \epsilon_{\vec{k}+\vec{q}})^2 + U\Phi^*_{\vec{q}}(\epsilon_{\vec{k}} - \epsilon_{\vec{k}+\vec{q}})]}{\omega + i0 - \tilde{E}_d + \mu} \right\}^{-1}, \quad (8)$$

$$\langle\langle \beta_{\vec{q}} \beta^\dagger_{\vec{q}} \rangle\rangle = \left\{ \omega - \omega^0_{\vec{q}} + i0 + \frac{V^2}{2NS} \sum_{\vec{k}} \frac{(n^f - n^g_{\vec{k}})(\epsilon_{\vec{k}} - \epsilon_{\vec{k}+\vec{q}})^2}{\omega + i0 - \tilde{E}_d + \tilde{\epsilon}_{\vec{k}}} + \frac{V^2 U}{2SN^2} \frac{[\sum_{\vec{k}} \frac{(n^f - n^g_{\vec{k}})(\epsilon_{\vec{k}} - \epsilon_{\vec{k}+\vec{q}})}{\omega + i0 - \tilde{E}_d + \tilde{\epsilon}_{\vec{k}}}]^2}{1 - \frac{U}{N}\sum_{\vec{k}} \frac{n^f - n^g_{\vec{k}}}{\omega + i0 - \tilde{E}_d + \tilde{\epsilon}_{\vec{k}}}} \right\}^{-1}, \quad (9)$$

$$\Phi_{\vec{q}} = \Phi^*_{\vec{q}} \equiv \sqrt{2S} \sum_{\vec{k}} \langle g^\dagger_{\vec{k}} f \beta^\dagger_{\vec{q}} \rangle = -\frac{V}{\pi} \int \mathrm{Im} \left\{ \frac{\frac{1}{N}\sum_{\vec{k}} \frac{(n^f - n^g_{\vec{k}})(\epsilon_{\vec{k}} - \epsilon_{\vec{k}+\vec{q}})}{\omega + i0 - \tilde{E}_d + \tilde{\epsilon}_{\vec{k}}}}{1 - \frac{U}{N}\sum_{\vec{k}} \frac{n^f - n^g_{\vec{k}}}{\omega + i0 - \tilde{E}_d + \tilde{\epsilon}_{\vec{k}}}} \langle\langle \beta_{\vec{q}} \beta^\dagger_{\vec{q}} \rangle\rangle \right\} \frac{d\omega}{\exp(\frac{\omega}{T}) - 1}. \quad (10)$$

Here, the magnon occupancy is

$$\mathcal{N}_{\vec{q}} \equiv \langle \beta^\dagger_{\vec{q}} \beta_{\vec{q}} \rangle = -\frac{1}{\pi} \int \mathrm{Im} \langle\langle \beta_{\vec{q}} \beta^\dagger_{\vec{q}} \rangle\rangle \frac{d\omega}{\exp(\omega/T) - 1} \quad (11)$$

($n^g_{\vec{k}}$ and n^f are expressed in a similar way). Hartree energies of magnons and those of e_g and t_{2g} electrons read

$$\omega^0_{\vec{q}} = \frac{H}{S} - \frac{2J}{S}(\epsilon_{\vec{q}} + d) + \frac{1}{2NS}\sum_{\vec{k}} n^g_{\vec{k}}(\epsilon_{\vec{k}+\vec{q}} - \epsilon_{\vec{k}}), \quad (12)$$

$$\tilde{\epsilon}_{\vec{k}} = \epsilon_{\vec{k}} + Un^f - \frac{H}{2S} + \frac{1}{2NS}\sum_{\vec{q}} \mathcal{N}_{\vec{q}}(\epsilon_{\vec{k}+\vec{q}} - \epsilon_{\vec{k}}), \quad (13)$$

$$\tilde{E}_d = E_d + \frac{U}{N}\sum_{\vec{k}} n^g_{\vec{k}} + \frac{H}{2S}. \quad (14)$$

The last terms in Eqs. (12) and (13) contain the double-exchange physics (ferromagnetic contribution to the spin-wave spectrum and the conduction band narrowing).

In analyzing Eqs. (7)–(10) we consider a 2D system; the 3D case can be expected to be similar. We replace all factors $(\epsilon_k - \epsilon_{k-q})^2$ with their average values over the isoenergetic surfaces [18] $\epsilon = \epsilon_{\vec{k}}$ and $\epsilon = \epsilon_{\vec{q}}$. We arrive at a system of mean-field equations for μ, n_b, and four quantities $\int \Phi(\epsilon)(2+\epsilon)\nu d\epsilon$, $\int \mathcal{N}(\epsilon)\epsilon\nu d\epsilon$, $\int \mathcal{N}(\epsilon) \times \langle v^2 \rangle_\epsilon \nu d\epsilon$, and $\int \mathcal{N}(\epsilon)(4-\epsilon^2)\nu d\epsilon$, where $\nu(\epsilon)$ and $\langle v^2 \rangle_\epsilon$ are the tight-binding density of states and average velocity square at a given energy ϵ. Solving these equations numerically, we observe the following:

(i) The localized band is broadened, and a temperature-dependent gap Δ (a new small *energy scale*) opens in the spectrum of itinerant electrons (fermions $g_{\vec{k}}$). Quasiparticle weight of itinerant electrons decreases when the energy approaches the gap from either side. The Fermi level lies below the gap, and the *quasiparticle weight* at the Fermi surface is strongly *suppressed* [Figs. 2(a) and 2(b)]. This behavior, which is already reminiscent of a T-dependent (pseudo)gap found experimentally [2], will be further modified in a more exact treatment (e.g., a finite relaxation time will arise in the second order in $1/S$). Equations (7) and (8) imply that in the present model, these spectral features are directly controlled by spin dynamics.

(ii) With only the spin-majority electrons contributing to the spin stiffness D, which in 2D or 3D is given by

$$DS = -\frac{1}{4dN}\sum_{\vec{k}} \epsilon_{\vec{k}} n^g_{\vec{k}} - J - \frac{V^2}{2dN}\sum_{\vec{k}} \frac{n^f - n^g_{\vec{k}}}{\tilde{\epsilon}_{\vec{k}} - \tilde{E}_d + i0} \left(\frac{\partial \epsilon_{\vec{k}}}{\partial \vec{k}}\right)^2,$$

the latter is suppressed in comparison with the usual double-exchange case [Fig. 2(c)]. An unusual feature of the present model is that the spin-flip continuum extends down to zero energy and momentum [19]. As a result, D also develops an imaginary part. This implies *strong magnon damping*, as observed experimentally [17].

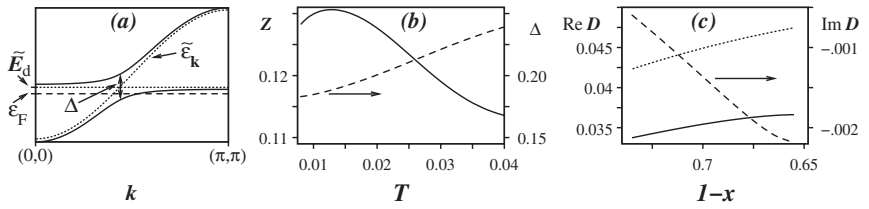

FIG. 2. Mean-field results for a 2D system with $E_d = -0.35$, $V = 0.2$, $U = 1.2$, and $J = 0$. (a) Schematic view of the itinerant fermion dispersion (solid lines); dotted lines show the unhybridized $\tilde{\epsilon}_k$ and $\tilde{E}_d$. (b) Temperature dependence of quasiparticle weight Z at the Fermi level (solid line) and the gap, $\Delta(T)$ (dashed line), for the electron density $1 - x = 0.7$. (c) Doping dependence of real (solid line) and imaginary (dashed) parts of spin stiffness D at $T = 0.0015$. Dotted line shows spin stiffness for a usual double-exchange model. t_{2g} occupancy increases from $n_f = 0.25$ at $1 - x = 0.65$ to $n_f = 0.43$ at $1 - x = 0.73$.

As for the diffusive central peak found in the inelastic neutron scattering [16], we expect it to arise once the magnon-assisted diffusive motion of t_{2g} electrons (neglected here) is taken into account. Experimentally, strongly damped magnons, central peak, and pseudogap in the density of states [2] (or optical Drude weight reduction [3]) are the key generic features of CMR manganates at the intermediate-to-high temperatures below T_C.

A relatively small value of U used in Fig. 2 is due to the reduced stability region for mean-field solutions with both n_f and n_g different from zero. This reduction is an expected artifact of a simplistic mean-field approach, mirroring, e.g., the greatly enhanced mean-field stability of ferromagnetism in the Hubbard model. This situation calls for further theoretical investigation, combining advanced mean-field schemes with numerical methods.

These future treatments will also have to address the issue of ferro- to paramagnetic transition and a possibility of charge ordering. We expect that any transition will be accompanied by a change of electron distribution between the two bands, thus changing the magnitude of the net spin $\mathcal{T}_i$ on-site. Experimentally, the relevant quantity is the average total spin $\langle \mathcal{T} \rangle = (3 + n^g - n^f)/2$ of a Mn ion, which should show temperature and magnetic field dependence, especially in the region around and above T_C. In particular, this should lead to an unconventional longitudinal spin dynamics [20] and to a renormalization of the Curie-Weiss constant (cf. Ref. [21]). In principle, the value of $\langle \mathcal{T} \rangle$ should also be accessible more directly via muon spin rotation [22] and NMR [23] measurements. We suggest that these methods (combined with electron spectroscopy) should be used to measure the value of $\langle \mathcal{T} \rangle$. Its temperature dependence, especially if it correlates with (magneto)transport properties, would imply that a successful theoretical description of CMR compounds should indeed include spin-minority localized electrons.

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