Two ferromagnetic phases in spin-Fermion systems

Naoum Karchev
Division of Material Physics, Graduate School of Engineering Science, Osaka University, Toyonaka, Osaka 560-8531, Japan and Department of Physics, University of Sofia, 1126 Sofia, Bulgaria

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We consider spin-Fermion systems which obtain their magnetic properties from a system of localized magnetic moments being coupled to conducting electrons. The dynamical degrees of freedom are spin-s operators of localized spins and spin-1/2 Fermi operators of itinerant electrons. We develop a modified spin-wave theory and obtain that the system has two ferromagnetic phases. At the characteristic temperature T^*, the magnetization of itinerant electrons becomes zero, and the high temperature ferromagnetic phase ($T^* < T < T_C$) is a phase where only localized electrons give contribution to the magnetization of the system. An anomalous increasing of magnetization below T^* is obtained in good agreement with experimental measurements of the ferromagnetic phase of UGe_2.

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This Brief Report is inspired from the experimental measurements of the ferromagnetic phase of UGe_2 which reveal the presence of an additional phase line that lies entirely within the ferromagnetic phase. The characteristic temperature of this transition T^*, which is below the Curie temperature T_C, decreases with pressure and disappears at a pressure close to the pressure at which a new phase of coexistence of superconductivity and ferromagnetism emerges.[1–3] A strong anomaly in resistivity is observed at T^*.[4] The additional phase transition demonstrates itself and through the change in the T dependence of the ordered ferromagnetic moment.[2,5–7] The magnetization shows an anomalous enhancement below T^*. An anomaly is found in the heat capacity at the characteristic temperature T^*.[8] Theoretically, it was assumed that the interplay of charge-density wave and spin-density wave fluctuations is the origin of anomalous properties.[9] Alternatively, it was proposed that the unusual phase diagram is result of novel tuning of the Fermi surface topology by the magnetization.[10]

Our objective is spin-Fermion systems, which obtain their magnetic properties from a system of localized magnetic moments and itinerant electrons. It is obtained that the true magnons in these systems, which are the transversal fluctuations corresponding to the total magnetization, are complicated mixtures of the transversal fluctuations of the spins of localized and itinerant electrons. The magnons interact with different spins in a different way, and the magnons' fluctuations suppress the ordered moments of the localized and itinerant electrons at different temperatures. As a result, the ferromagnetic phase is divided into two phases: low temperature phase $0 < T < T^*$, where all electrons contribute the ordered ferromagnetic moment, and high temperature phase $T^* < T < T_C$, where only localized spins form magnetic moments. To describe the two phases, a modified spin-wave theory is developed. We have reproduced theoretically the anomalous temperature dependence of the ordered moment, known from the experiments with UGe_2.[2,5–7]

The spin-Fermion model is known as s-d (or s-f). The model appears in the literature also as the ferromagnetic Kondo lattice model or the double exchange model. The exact results for the spin-Fermion model are reported in Ref. 11.

The dynamical degrees of freedom are spin-s operators of localized spins and spin-1/2 Fermi operators of itinerant electrons. We consider a theory with Hamiltonian

$$h = H - \mu N = -t\sum_{\langle ij\rangle} (c^+_{i\sigma}c_{j\sigma} + \text{H.c.}) - \mu\sum_i n_i - J^l\sum_{\langle ij\rangle} \mathbf{S}_i \cdot \mathbf{S}_j - J\sum_i \mathbf{S}_i \cdot \mathbf{s}_i, \tag{1}$$

where $s^\nu_i = \frac{1}{2}\sum_{\sigma\sigma'} c^+_{i\sigma}\tau^\nu_{\sigma\sigma'}c_{i\sigma'}$, with the Pauli matrices (τ^x, τ^y, τ^z), is the spin of the conduction electrons, $\mathbf{S}_i$ is the spin of the localized electrons, μ is the chemical potential, and $n_i = c^+_{i\sigma}c_{i\sigma}$. The sums are over all sites of a three-dimensional cubic lattice, and $\langle i,j\rangle$ denotes the sum over the nearest neighbors. The Heisenberg term ($J^l > 0$) describes ferromagnetic Heisenberg exchange between nearest-neighbor localized electrons. The term which describes the spin-Fermion interaction ($J > 0$) is known as a Kondo interaction ($J = J_K$) in the ferromagnetic Kondo model or as a Hund term in the double exchange model ($J = J_H$ and $J^l < 0$).

We represent the Fermi operators in terms of the Schwinger bosons $(\varphi_{i,\sigma}, \varphi^+_{i,\sigma})$ and slave Fermions (h_i, h^+_i, d_i, d^+_i). The Bose fields are doublets ($\sigma = 1, 2$) without charge, while Fermions are spinless with charges 1 (d_i) and -1 (h_i).

$$c_{i\uparrow} = h^+_i\varphi_{i1} + \varphi^+_{i2}d_i, \quad c_{i\downarrow} = h^+_i\varphi_{i2} - \varphi^+_{i1}d_i,$$

$$n_i = 1 - h^+_i h_i + d^+_i d_i, \quad s^\nu_i = \frac{1}{2}\sum_{\sigma\sigma'} \varphi^+_{i\sigma}\tau^\nu_{\sigma\sigma'}\varphi_{i\sigma'},$$

$$\varphi^+_{i1}\varphi_{i1} + \varphi^+_{i2}\varphi_{i2} + d^+_i d_i + h^+_i h_i = 1. \tag{2}$$

Next, we make a change of variables, introducing Bose doublets $\zeta_{i\sigma}$ and $\zeta^+_{i\sigma}$,[12]

$$\zeta_{i\sigma} = \varphi_{i\sigma}(1 - h^+_i h_i - d^+_i d_i)^{-1/2},$$
$$\zeta^+_{i\sigma} = \varphi^+_{i\sigma}(1 - h^+_i h_i - d^+_i d_i)^{-1/2}, \tag{3}$$

where they satisfy the constraint $\zeta^+_{i\sigma}\zeta_{i\sigma} = 1$. In terms of these fields, the spin vectors of the itinerant electrons have the form

$$s_i^\nu = \frac{1}{2}\sum_{\sigma\sigma'} \zeta_{i\sigma}^+ \tau_{\sigma\sigma'}^\nu \zeta_{i\sigma'}[1 - h_i^+ h_i - d_i^+ d_i]. \quad (4)$$

When, in the ground state, the lattice site is empty, the operator identity $h_i^+ h_i = 1$ is true. When the lattice site is doubly occupied, $d_i^+ d_i = 1$. Hence, when the lattice site is empty or doubly occupied, the spin on this site is zero. When the lattice site is neither empty nor doubly occupied ($h_i^+ h_i = d_i^+ d_i = 0$), the spin equals $\mathbf{s}_i = 1/2\mathbf{n}_i$, where the unit vector $n_i^\nu = \Sigma_{\sigma\sigma'} \zeta_{i\sigma}^+ \tau_{\sigma\sigma'}^\nu \zeta_{i\sigma'}$ ($\mathbf{n}_i^2 = 1$) identifies the local orientation of the spin of the itinerant electron. Let us average the spin of electrons in the subspace of the Fermions (d_i^+, d_i) and (h_i^+, h_i) (to integrate the Fermions out in the path integral approach) and introduce the notation

$$m = \frac{1}{2}(1 - \langle h_i^+ h_i \rangle - \langle d_i^+ d_i \rangle). \quad (5)$$

One obtains $\mathbf{s}_i = m\mathbf{n}_i$, where $\mathbf{s}_i^2 = m^2$. Hence, the amplitude of the spin vector m is an effective spin of the itinerant electrons accounting for the fact that some sites, in the ground state, are doubly occupied or empty.

It is more convenient to use the rescaled Bose fields,

$$\xi_{i\sigma} = \sqrt{2m}\zeta_{i\sigma}, \quad \xi_{i\sigma}^+ = \sqrt{2m}\zeta_{i\sigma}^+, \quad (6)$$

which satisfy the constraint $\xi_{i\sigma}^+ \xi_{i\sigma} = 2m$. Let us introduce the vector,

$$M_i^\nu = \frac{1}{2}\sum_{\sigma\sigma'} \xi_{i\sigma}^+ \tau_{\sigma\sigma'}^\nu \xi_{i\sigma'}, \quad \mathbf{M}_i^2 = m^2. \quad (7)$$

Then, the spin vector of itinerant electrons can be written in the form

$$\mathbf{s}_i = \frac{1}{2m}\mathbf{M}_i(1 - h_i^+ h_i - d_i^+ d_i). \quad (8)$$

The contribution of itinerant electrons to the total magnetization is $\langle \mathbf{s}_i^z \rangle$. Accounting for the definition of m [see Eq. (5)], one obtains $\langle \mathbf{s}_i^z \rangle = \langle \mathbf{M}_i^z \rangle$.

The Hamiltonian is quadratic with respect to the Fermions d_i, d_i^+ and h_i, h_i^+, and one can average in the subspace of these Fermions (to integrate them out in the path integral approach). As a result, we obtain an effective theory of two-spin vectors $\mathbf{S}_i$ and $\mathbf{M}_i$ with Hamiltonian

$$h_{eff} = -J^l \sum_{\langle ij \rangle} \mathbf{S}_i \cdot \mathbf{S}_j - J^{it} \sum_{\langle ij \rangle} \mathbf{M}_i \cdot \mathbf{M}_j - J\sum_i \mathbf{S}_i \cdot \mathbf{M}_i. \quad (9)$$

The first term is the term which describes the exchange of localized spins in the Hamiltonian [Eq. (1)]. The second term is obtained by integrating out the Fermions. It is calculated in the one loop approximation and in the limit when the frequency and the wave vector are small. For the effective exchange constant J^{it}, at zero temperature, we obtained

$$J^{it} = \frac{t}{6m^2}\frac{1}{N}\sum_k \left(\sum_{\nu=1}^{3} \cos k_\nu\right)[1 - \theta(-\varepsilon_k^h) + \theta(-\varepsilon_k^d)] - \frac{2t^2}{3m^2 sJ}\frac{1}{N}\sum_k \left(\sum_{\nu=1}^{3} \sin^2 k_\nu\right)[1 - \theta(-\varepsilon_k^h) - \theta(-\varepsilon_k^d)], \quad (10)$$

where N is the number of lattice sites, ε_k^h and ε_k^d are Fermions' dispersions,

$$\varepsilon_k^h = 2t(\cos k_x + \cos k_y + \cos k_z) + stJ/2 + \mu,$$
$$\varepsilon_k^d = -2t(\cos k_x + \cos k_y + \cos k_z) + stJ/2 - \mu, \quad (11)$$

and wave vector k runs over the first Brillouin zone of a cubic lattice. The third term in Eq. (9) is obtained from the last term in the Hamiltonian [Eq. (1)] using the representation [Eq. (8)] for the spin of itinerant electrons and Eq. (5).

We are going to study the ferromagnetic phase of the two-spin system [Eq. (9)] with $J^l > 0$, $J^{it} > 0$, and $J > 0$. To proceed, we use the Holstein-Primakoff representation of the spin vectors $\mathbf{S}_j(a_j^+, a_j)$ and $\mathbf{M}_j(b_j^+, b_j)$, where a_j^+, a_j and b_j^+, b_j are Bose fields. In terms of these fields and keeping only the quadratic terms, the effective Hamiltonian [Eq. (9)] adopts the form

$$h_{eff} = sJ^l \sum_{\langle ij \rangle} (a_i^+ a_i + a_j^+ a_j - a_j^+ a_i - a_i^+ a_j) + mJ^{it} \sum_{\langle ij \rangle} (b_i^+ b_i + b_j^+ b_j - b_j^+ b_i - b_i^+ b_j) - J\sum_i [\sqrt{sm}(a_i^+ b_i + b_i^+ a_i) - sb_i^+ b_i - ma_i^+ a_i]. \quad (12)$$

In momentum space representation, the Hamiltonian reads

$$h_{eff} = \sum_k [\varepsilon_k^a a_k^+ a_k + \varepsilon_k^b b_k^+ b_k - \gamma(a_k^+ b_k + b_k^+ a_k)], \quad (13)$$

where the dispersions are given by equalities,

$$\varepsilon_k^a = 2sJ^l \varepsilon_k + mJ, \quad \varepsilon_k^b = 2mJ^{it}\varepsilon_k + sJ, \quad (14)$$

$\varepsilon_k = 3 - \cos k_x - \cos k_y - \cos k_z$, and $\gamma = J\sqrt{sm}$.

To diagonalize the Hamiltonian, one introduces Bose fields $\alpha_k, \alpha_k^+, \beta_k, \beta_k^+$,

$$a_k = \cos\theta_k \alpha_k + \sin\theta_k \beta_k, \quad b_k = -\sin\theta_k \alpha_k + \cos\theta_k \beta_k, \quad (15)$$

with coefficients of transformation,

$$\cos\theta_k = \sqrt{\frac{1}{2}\left(1 + \frac{\varepsilon_k^a - \varepsilon_k^b}{\sqrt{(\varepsilon_k^a - \varepsilon_k^b)^2 + 4\gamma^2}}\right)}, \quad (16)$$

and $\sin\theta_k = (1 - \cos^2\theta_k)^{1/2}$. The transformed Hamiltonian

$$h_{eff} = \sum_k (E_k^\alpha \alpha_k^+ \alpha_k + E_k^\beta \beta_k^+ \beta_k), \quad (17)$$

where

$$E_k^\pm = \frac{1}{2}[\varepsilon_k^a + \varepsilon_k^b \pm \sqrt{(\varepsilon_k^a - \varepsilon_k^b)^2 + 4\gamma^2}], \quad (18)$$

and $E_k^\alpha = E_k^+$, $E_k^\beta = E_k^-$. With positive exchange constants $J^l > 0$, $J^{it} > 0$, and $J > 0$, the Bose fields' dispersions are positive $\varepsilon_k^a > 0$, $\varepsilon_k^b > 0$ for all wave vectors k. As a result, $E_k^\alpha > 0$, and $E_k^\beta \geqslant 0$ with $E_0^\beta = 0$. Near the zero wave vector, $E_k^\beta \approx \rho k^2$, where the spin-stiffness constant is $\rho = (s^2 J^l + m^2 J^{it})/(s+m)$. Hence, β_k is the long-range (magnon) excitation in the two-spin effective theory, while α_k is a gapped excitation with gap $E_0^\alpha = (s+m)J$.

The dimensionless magnetization per lattice site of the system M is a sum of the magnetization of the localized electrons $M^l = \langle S_i^z \rangle$ and the magnetization of the itinerant electrons $M^{it} = \langle s_i^z \rangle = \langle M_i^z \rangle$ ($M = M^l + M^{it}$). By means of the Holstein-Primakoff representation the magnetization adopts

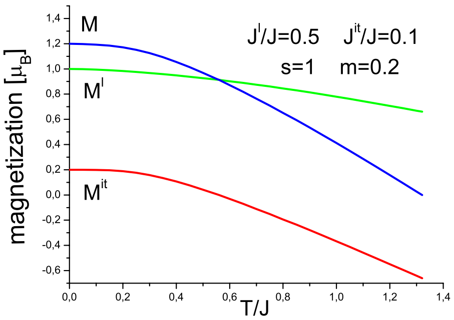

FIG. 1. (Color online) Temperature dependence of the ferromagnetic moments: M (blue line)—the magnetization of the system, M^l (green line)—contribution of the localized electrons, M^{it} (red line)—contribution of the itinerant electrons for parameters $s=1$, $m=0.2$, $J^l/J=0.5$, and $J^{it}/J=0.1$: spin-wave theory.

the form $M^l=s-1/N\Sigma_k\langle a_k^+a_k\rangle$ and $M^{it}=m-1/N\Sigma_k\langle b_k^+b_k\rangle$. Finally, by means of the transformation [Eq. (15)], one can rewrite M^l and M^{it} in terms of the Bose functions of the excitations $\alpha_k-n_k^\alpha$ and $\beta_k-n_k^\beta$,

$$M^l = s - \frac{1}{N}\sum_k [\cos^2\theta_k n_k^\alpha + \sin^2\theta_k n_k^\beta],$$

$$M^{it} = m - \frac{1}{N}\sum_k [\sin^2\theta_k n_k^\alpha + \cos^2\theta_k n_k^\beta]. \qquad (19)$$

The magnetization of the system is

$$M = s + m - \frac{1}{N}\sum_k [n_k^\alpha + n_k^\beta]. \qquad (20)$$

The magnon excitation β_k in the effective theory [Eq. (9)] is a complicated mixture of the transversal fluctuations of the spins of localized and itinerant electrons [Eq. (15)]. As a result, the magnons' fluctuations suppress in a different way the magnetic order of these electrons. Quantitatively, this depends on the coefficients $\cos\theta_k$ and $\sin\theta_k$ in Eq. (19). If the spin-Fermion interaction is very strong, $J \gg J^{it}$ and $J \gg J^l$, one can calculate the coefficients approximately using approximate expressions for dispersions [Eq. (14)]: $\varepsilon_k^a \approx mJ$ and $\varepsilon_k^a \approx sJ$. As a result, one obtains $\cos^2\theta_k \approx m/(m+s)$. For large J, the gap of the α excitation is very big, $E_0^\alpha=(s+m)J$, and one can drop this excitation in the calculations. Then, the approximate expressions for magnetization satisfy $M^l/s = M^{it}/m$, which means that the strong spin-Fermion interaction aligns the magnetic orders of the itinerant and localized electrons so strongly that they become zero at one and just the same temperature. The result is different if the spin-Fermion interaction is relatively small. The magnetization depends on the dimensionless temperature T/J and dimensionless parameters s, m, J^l/J, and J^{it}/J. We consider a theory with spin of the localized electrons $s=1$ and calculate the parameters of the effective theory [Eq. (9)] in one Fermion-loop approximation for density of Fermions $n=0.4$ and microscopic parameter $J/t=12.4$. The result is $m=0.2$ and $J^{it}/J=0.1$. Finally, we set $J^l/J=0.5$. For these effective parameters, the functions $M(T/J)$, $M^l(T/J)$, and $M^{it}(T/J)$ are depicted in Fig. 1 The green line is the magnetization of the localized electrons, the red line is the magnetization of the itinerant electrons, and the blue line is the total magnetization. The figure shows that the magnetic order of itinerant electrons (red line) is suppressed first, at temperature $T^*/J = 0.5603$. Once suppressed, the magnetic order cannot be restored at temperatures above T^* because of the increasing effect of magnon fluctuations. Hence, the magnetization of the itinerant electrons should be zero above T^*. As is evident from Fig. 1, this is not the result within customary spin-wave theory.

To solve the problem, we use the idea on the description of the paramagnetic phase of two-dimensional ferromagnets ($T>0$) by means of the modified spin-wave theory.[13,14] In the simplest version, the spin-wave theory is modified by introducing a parameter which enforces the magnetization of the system to be equal to zero in the paramagnetic phase.

In the present case, we have the two-spin system and we introduce two parameters λ^l and λ^{it} to enforce the magnetic moments both of the localized and the itinerant electrons to be equal to zero in the paramagnetic phase. To this end, we add two new terms to the effective Hamiltonian [Eq. (12)],

$$\hat{h}_{eff} = h_{eff} - \sum_i [\lambda^l S_i^z + \lambda^{it} M_i^z]. \qquad (21)$$

In momentum space, the Hamiltonian adopts the form [Eq. (13)] with new dispersions $\hat{\varepsilon}_k^a=\varepsilon_k^a+\lambda^l$ and $\hat{\varepsilon}_k^b=\varepsilon_k^b+\lambda^{it}$, where the old dispersions are given by equalities [Eq. (14)]. We utilize the same transformation [Eq. (15)] with coefficients $\cos\hat{\theta}_k$ and $\sin\hat{\theta}_k$ which depend on the new dispersions in the same way as the old ones depend on the old dispersions [Eq. (16)]. In terms of the α_k and β_k bosons, the Hamiltonian $\hat{h}_{eff}$ adopts the form [Eq. (17)] with dispersions $\hat{E}_k^\alpha$ and $\hat{E}_k^\beta$, which can be written in the form [Eq. (18)] replacing ε_k^a and ε_k^b with $\hat{\varepsilon}_k^a$ and $\hat{\varepsilon}_k^b$.

We have to do some assumptions for parameters λ^l and λ^{it} to ensure correct definition of the two-boson theory. For that purpose, it is convenient to represent the parameters λ^l and λ^{it} in the form $\lambda^l=mJ\mu^l-mJ$ and $\lambda^{it}=sJ\mu^{it}-sJ$. In terms of the parameters μ^l and μ^{it}, the dispersion reads $\hat{\varepsilon}_k^a=2sJ^l\varepsilon_k+mJ\mu^l$ and $\hat{\varepsilon}_k^b=2mJ^{it}\varepsilon_k+sJ\mu^{it}$. The conventional spin-wave theory is reproduced when $\mu^l=\mu^{it}=1$ ($\lambda^l=\lambda^{it}=0$). We assume μ^l and μ^{it} to be positive ($\mu^l>0, \mu^{it}>0$). Then, $\hat{\varepsilon}_k^a>0$, $\hat{\varepsilon}_k^b>0$, and $\hat{E}_k^\alpha>0$ for all values of the wave vector k. To explore the dispersion $\hat{E}_k^\beta=\frac{1}{2}[\hat{\varepsilon}_k^a+\hat{\varepsilon}_k^b-\sqrt{(\hat{\varepsilon}_k^a-\hat{\varepsilon}_k^b)^2+4\gamma^2}]$, we use the identity $(\hat{\varepsilon}_k^a-\hat{\varepsilon}_k^b)^2+4\gamma^2=(\hat{\varepsilon}_k^a+\hat{\varepsilon}_k^b)^2-4(\hat{\varepsilon}_k^a\hat{\varepsilon}_k^b-\gamma^2)$. It shows that $\hat{E}_k^\beta \geqslant 0$ if $\hat{\varepsilon}_k^a\hat{\varepsilon}_k^b-\gamma^2 \geqslant 0$. Since $\hat{\varepsilon}_k^a\hat{\varepsilon}_k^b \geqslant \hat{\varepsilon}_0^a\hat{\varepsilon}_0^b = smJ^2\mu^l\mu^{it}$ for all values of the wave vector k, the β_k dispersion is non-negative, $\hat{E}_k^\beta \geqslant 0$ if $\mu^l\mu^{it} \geqslant 1$. In the particular case, $\mu^l\mu^{it}=1$, $\hat{E}_0^\beta=0$, and, near the zero wave vector, $\hat{E}_k^\beta \approx \hat{\rho}k^2$, with spin-stiffness constant equals $\hat{\rho}=(s^2J^l\mu^{it}+m^2J^{it}\mu^l)/(s\mu^{it}+m\mu^l)$. Hence, in this case, the β_k boson is the long-range excitation (magnon) in the system. In the case the $\mu^l\mu^{it}>1$, both α_k boson and β_k boson are gapped excitations.

We introduced the parameters λ^l and λ^{it} (μ^l, μ^{it}) to enforce the magnetic order of localized and itinerant electrons to be equal to zero. We find out the parameters μ^l and μ^{it} solving the system of two equations $M^l=M^{it}=0$, where the ordered moments have the same representation as Eq. (19) but with coefficients $\cos\hat{\theta}_k, \sin\hat{\theta}_k$, and dispersions $\hat{E}_k^{\alpha}$, $\hat{E}_k^{\beta}$ in the expressions for the Bose functions. The numerical calculations show that for high enough temperature, $\mu^{it}>1$, $1>\mu^{l}>0$, and $\mu^{it}\mu^{l}>1$. Hence, α_k and β_k excitations are gapped. When the temperature decreases, μ^{it} decreases remaining larger than one, μ^{l} decreases too remaining positive, and the product $\mu^{l}\mu^{it}$ decreases remaining larger than one. At temperature $T_C/J=2.812$, one obtains $\mu^{it}=5.0427$, $\mu^{l}=0.1983$, and therefore $\mu^{l}\mu^{it}=1$. Hence, at T_C, long-range excitation (magnon) emerges in the spectrum which means that T_C is the Curie temperature.

Below the Curie temperature the spectrum contains magnon excitations, thereupon $\mu^{l}\mu^{it}=1$. It is convenient to represent the parameters in the following way:

$$\mu^{it}=\mu, \quad \mu^{l}=1/\mu. \tag{22}$$

In the ferromagnetic phase, magnon excitations are the origin of the suppression of magnetization. Near the zero temperature, their contribution is small, and at zero temperature, $M^{it}=m$ and $M^{l}=s$. Increasing the temperature, magnon fluctuations suppress the magnetization. For the chosen parameters, they first suppress the magnetization of the itinerant electrons at T^* $[M^{l}(T^*)>0]$. Once suppressed, the magnetic moment of itinerant electrons cannot be restored increasing the temperature above T^*. To formulate this mathematically, we modify the spin-wave theory introducing the parameter μ [Eq. (22)]. Below T^*, $\mu=1$, or in terms of λ parameters, $\lambda^{l}=\lambda^{it}=0$, which reproduces the customary spin-wave theory. Increasing the temperature above T^*, the magnetic moment of the itinerant electron should be zero. This is why we impose the condition $M^{it}(T)=0$ if $T>T^*$. For temperatures above T^*, the parameter μ is a solution of this equation. We utilize the obtained function $\mu(T)$ to calculate the magnetization of the localized electrons M^{l} as a function of the temperature. Above T^*, M^{l} is equal to the magnetization of the system. The magnetic moments of the localized and itinerant electrons as well as the magnetization of the system as a function of the temperature are depicted in Fig. 2 for parameters $s=1$, $m=0.2$, $J^{l}/J=0.5$, and $J^{it}/J=0.1$.

The figure shows an anomalous increasing of the magnetization M below T^* which is in very good agreement with the experiment [see Fig. 1 (Ref. 6)]. The present theory enables us to gain insight into the nature of the two phases. In the low temperature phase $(0,T^*)$, the localized and itinerant electrons contribute to the magnetization of the system, while in the high temperature phase (T^*,T_C), only localized electrons form ferromagnetic moments. At first sight, it seems to be counterintuitive because the local moments build an effective magnetic field, which, due to spin-Fermion interaction, leads to finite itinerant electron spin polarization. This is true in the classical limit. In the quantum case, the spin-wave fluctuations suppress the magnetic orders of the itinerant and localized electrons at different temperatures T^* and T_C as a result of different interactions of the magnon with localized and itinerant electrons. The spin-Fermion interaction increases the alignment of the local moments, and magnetic order of itinerant electrons is very strong and T^* approaches T_C.

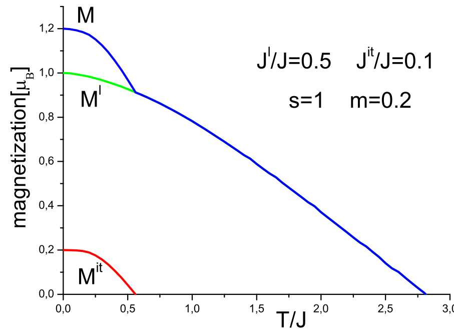

FIG. 2. (Color online) Temperature dependence of the ferromagnetic moments: M (blue line)—the magnetization of the system, M^{l} (green line)—contribution of the localized electrons, M^{it} (red line)—contribution of the itinerant electrons for parameters $s=1$, $m=0.2$, $J^{l}/J=0.5$, and $J^{it}/J=0.1$: modified spin-wave theory.

It is well known that the onset of magnetism in the itinerant systems is accompanied with a strong anomaly in resistivity.[15] This phenomenon is experimentally observed at T^* in the case of UGe_2.[4] This is another support for the theoretical interpretation of T^* as a temperature at which the itinerant electrons form ferromagnetic order.

To conclude, we note that to do more precise fitting with experimental values of the Curie temperature, one has to account for the magnon-magnon interaction. However, even the approximate calculations in the present Brief Report capture the main feature of the two-spin ferromagnetic systems and the existence of two phases.

The next step of our investigation is to understand the mechanism of decreasing the phase temperature T^*. This will help us to understand the origin of the superconductivity in these materials.

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