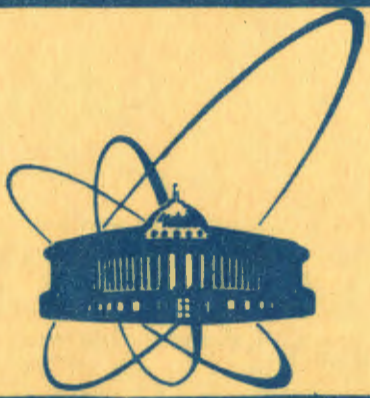

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# SUPERCONDUCTIVITY AND FERROMAGNETISM IN THE EXTENDED RANDOM ZENER MODEL

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## 1. INTRODUCTION

Up to now the problem of the simultaneous occurrence of singlet superconductivity (SC) and homogeneous ferromagnetism (FM) in the same volume of a sample has not been solved definitively. Neutron scattering experiments yielded T versus x diagrams which exhibit, e.g., mixed phases of SC with nonuniform FM in $(Er_{1-x}Ho_x)Rh_4B_4$ [1] and shortrange ferromagnetic correlations in $(Ce_{1-x}Tb_x)Ru_2$ [2]. More recently, measurements in $(Er_{1-x}Ho_x)Rh_4B_4$ [3], however, suggest the possibility of coexistence between SC and long-range FM. Theoretical studies of the coexistence question lead to different results depending on the models and approximations used.

Microscopic treatments [4-10] are usually based on the s-d exchange model combined with a BCS term [6-10]. The magnetic ions are either placed regularly on the lattice sites [6,9,10] or treated as randomly distributed impurities in the Born approximation [4,5,7]. SC and FM can coexist also at $T=0$ [4] or only at $T\neq 0$ [6,10]. A small coexistence region was found in [8] for alloys weighted in the virtualcrystal approximation.

Quasiphenomenological approaches (Ginzburg-Landau-type theories) [11-14] have been devoted to the multicritical behaviour of magnetic superconductors. In [11] it was shown that a first order phase transition between SC and FM takes place. A tetracritical point may be expected [12] due to disorder. The coexistence of SC and FM is excluded by a realistic choice of parameters in [13].

It is desirable to derive microscopically the coefficients of a Landau-type free energy functional with two coupled order parameters. Work along this line was done on the basis of a BCS or BCS-like interaction in combination with the s-d exchange [15,16], the periodic Anderson [17] and the random Hubbard [18] models. Concerning alloys with transition metals the scattering problem should be handled within the coherent potential approximation (CPA). In the present paper the scheme [18] is generalized to a random Zener-model.

## 2. DISORDERED MODEL SYSTEM

Let us model a disordered $A_xB_{1-x}$ alloy by the tight-binding Hamiltonian

$$H = H_V + H_U + H_J + H_\lambda \qquad (1)$$

composed of the one-electron part

$$H_V = \sum_{i\sigma} \epsilon_i n_{i\sigma} + \sum_{\substack{ij\sigma \\ (i\neq j)}} t_{ij}\, c^+_{i\sigma} c_{j\sigma}\,, \qquad (2)$$

the local Coulomb repulsion

$$H_U = \sum_i U_i\, n_{i\uparrow} n_{i\downarrow}\,, \qquad (3)$$

the point-like s-d exchange

$$H_J = -2 \sum_i J_i \vec{\sigma}_i \vec{S}_i\,, \qquad (4)$$

and the attractive contact term

$$H_\lambda = -\sum_{ij} \lambda_{ij}\, c^+_{i\uparrow} c^+_{i\downarrow} c_{j\downarrow} c_{j\uparrow}\,. \qquad (5)$$

Here $n^+_{i\sigma} = c^+_{i\sigma} c_{i\sigma}$ is the particle number operator, the spin operators of the itinerant electrons are defined by $\sigma_i^+ = c^+_{i\uparrow} c_{i\downarrow}$, $\sigma_i^- = c^+_{i\downarrow} c_{i\uparrow}$, and $\sigma^z_i = \frac{1}{2}(n_{i\uparrow} - n_{i\downarrow})$, and $\vec{S}_i$ denotes the operator for a spin localized at the lattice site i. The local atomic potential $\epsilon_i$, the hopping integral $t_{ij}$, the repulsion strength $U_i$, the exchange integral $J_i$, and the coupling constant $\lambda_{ij}$ are random parameters, i.e., discrete bi- or trimodal variables in a lattice of A and B atoms. As special cases of (1) one can stress disordered versions of the

(i) Hubbard model $H_V + H_U$,
(ii) s-d exchange model $H_V + H_J$,
(iii) Zener model $H_V + H_U + H_J$, and
(iv) the model $H_V + H_\lambda$, which reflects a contact-type pairing [19] in narrow-band superconductors. Note that the Hamiltonian (1) involves two non-equivalent kinds of spin, itinerant electrons and localized magnetic moments.

Introduce the thermodynamical electron Green function as the 4x4 matrix propagator

$$G_{ij}(\tau - \tau') = -\langle T \begin{pmatrix} c_{i\uparrow}(\tau) \\ c_{i\downarrow}(\tau) \\ \bar{c}_{i\uparrow}(\tau) \\ \bar{c}_{i\downarrow}(\tau) \end{pmatrix} \otimes (\bar{c}_{j\uparrow}(\tau'), \bar{c}_{j\downarrow}(\tau'), c_{j\uparrow}(\tau'), c_{j\downarrow}(\tau'))\rangle, \qquad (6)$$

where $\langle \ldots \rangle$ means the thermal average. In the Hartree-Fock-Bogolubov approximation of (2) to (5) we get the Fourier transformed inverse Greenian



$$(G^{-1}(z))_{ij} = \begin{pmatrix} (z+\frac{U_i}{2}m_i + J_i\langle S_i^z\rangle)\delta_{ij} - [H_0]_{ij} & 0 & 0 & \Delta_i\delta_{ij} \\ 0 & (z-\frac{U_i}{2}m_i - J_i\langle S_i^z\rangle)\delta_{ij} - [H_0]_{ij} & -\Delta_i\delta_{ij} & 0 \\ 0 & -\Delta_i^*\delta_{ij} & (z-\frac{U_i}{2}m_i - J_i\langle S_i^z\rangle)\delta_{ij} + [H_0]_{ij} & 0 \\ \Delta_i^*\delta_{ij} & 0 & 0 & (z+\frac{U_i}{2}m_i + J_i\langle S_i^z\rangle)\delta_{ij} + [H_0]_{ij} \end{pmatrix} \qquad (7)$$

defined at $z = z_n = i(2n+1)\pi T$. Here the notation $[H_0]_{ij} = (\epsilon_i + \frac{U_i}{2}n_i)\delta_{ij} + t_{ij}(1-\delta_{ij})$ with $n_i = \langle n_{i\uparrow}\rangle + \langle n_{i\downarrow}\rangle$, and $m_i = \langle n_{i\uparrow}\rangle - \langle n_{i\downarrow}\rangle$ is used. The zeros in (7) reflect that triplet superconducting states and transverse spin arrangements have been excluded a priori; in other words, only singlet pairing and FM are taken into account.

The free energy corresponding to the mean field level of (1) (in particular to (7)) is found to be

$$F = -\frac{T}{2}\sum_n \mathrm{Sp\,tr} \ln(-G^{-1}(z_n)) - \frac{1}{4}\sum_i U_i n_i^2 + \frac{1}{4}\sum_i U_i m_i^2 + \sum_{ij} \Delta_i^* (\lambda^{-1})_{ij} \Delta_j - T\sum_i \ln\left(2\cosh\left(\frac{J_i m_i}{2T}\right)\right) + \sum_i J_i m_i \langle S_i^z\rangle\,, \qquad (8)$$

where Sp refers to the trace over the lattice space, while tr is restricted to the 4x4 (spin) matrix. According to the variational principle one derives via $\frac{\delta F}{\delta m_i} = 0$, $\frac{\delta F}{\delta \langle S_i^z\rangle} = 0$, and $\frac{\delta F}{\delta \Delta_i^*} = 0$ the self-consistency conditions

$$\langle S_i^z\rangle = \frac{1}{2}\tanh\left(\frac{J_i m_i}{2T}\right), \qquad (9)$$

$$m_i = \frac{T}{2}\sum_n \left(G_{ii}^{11}(z_n) - G_{ii}^{22}(z_n) - G_{ii}^{33}(z_n) + G_{ii}^{44}(z_n)\right), \qquad (10)$$

$$\Delta_i = \frac{T}{2}\sum_n \sum_j{}' \lambda_{ij}\left(G_{jj}^{14}(z_n) - G_{jj}^{23}(z_n)\right), \qquad (11)$$

respectively. Analogously, the mean electron number $n_i$ at site i is determined self-consistently. Note that (9) includes the spin number of 1/2 for the localized spins. The Coulomb renormalization entering into $\Delta_i$ is competitive with $\lambda_{ii}$ [19]; the net effect of these sitediagonal elements is neglected in (11). It is pointed out that all quantities $\langle S_i^z\rangle$, $m_i$, and $\Delta_i$ are still

configuration-dependent. In the sense of a terminal point approximation they depend at least on the site i due to the lack of the translational symmetry.

## 3. LOOP EXPANSION OF THE FREE ENERGY

In order to expand the first term of (8) we write (cf. /18/)

$$G^{-1} = G_0^{-1}(1 - G_0 W) \tag{12}$$

with the definitions

$$G_0 = \begin{pmatrix} G_0^{11} & & & 0 \\ & G_0^{11} & & \\ & & G_0^{33} & \\ 0 & & & G_0^{33} \end{pmatrix} \qquad W = \begin{pmatrix} \tilde{\Phi} & 0 & 0 & -\Delta \\ 0 & -\tilde{\Phi} & \Delta & 0 \\ 0 & \Delta^* & -\tilde{\Phi} & 0 \\ -\Delta^* & 0 & 0 & \tilde{\Phi} \end{pmatrix} \tag{13}$$

$G_0$ stands for the unperturbed normal-state random Green function. More explicitly, one has

$$((G_0^{33}(z))^{-1})_{ij} = z\delta_{ij} \mp [H_0]_{ij} . \tag{14}$$

The perturbation W is a local one because of $[\tilde{\Phi}]_{ij} = -(\frac{U_i}{2} m_i + J_i \langle S_i^z \rangle)\delta_{ij}$ and $[\Delta]_{ij} = \Delta_i \delta_{ij}$. Moreover, we expand (9) in the form $\langle S_i^z \rangle = \frac{J_i m_i}{4T} - \frac{1}{48}(\frac{J_i m_i}{T})^3 + \ldots$ justified at finite temperature near the magnetic phase transition. This gives rise to the definition

$$\Phi_i = -\frac{1}{2}(U_i + \frac{J_i^2}{2T}) m_i = -\frac{\tilde{U}_i}{2} m_i \tag{15}$$

as the local magnetic order parameter.

Using (12) to (15) one can decompose (8) by

$$F[\{\Phi\}, \{\Delta\}] = F_0 + F_2[\{\Phi\}, \{\Delta\}] + F_4[\{\Phi\}, \{\Delta\}], \tag{16}$$

where $F_0$ is the irrelevant normal-state contribution. The loop expansion (16) of the free energy is carried out up to the fourth order in $\Phi_i$ (or magnetization $m_i$) and the superconducting order parameter $\Delta_i$ to give

$$F_2[\{\Phi\}, \{\Delta\}] = \sum_{ij} \Phi_i (\chi_{0ij}^{11} + \frac{1}{\tilde{U}_i}\delta_{ij}) \Phi_j + \sum_{ij} \Delta_i^* (\chi_{0ij}^{13} + (\lambda^{-1})_{ij}) \Delta_j ,$$

$$F_4[\{\Phi\}, \{\Delta\}] = -\frac{1}{3T^3} \sum_{ij} \Phi_i \chi_{0ij}^{11} \frac{J_j^4}{\tilde{U}_i^3} \Phi_j^3 - \frac{1}{4T^3} \sum_i (\frac{J_i}{\tilde{U}_i})^4 \Phi_i^4 + \tag{17}$$



$$+ \sum_{ijk\ell} \{ \frac{1}{2} \Lambda_{ijk\ell}^{1111} \Phi_j \Phi_k \Phi_\ell \Phi_i + \frac{1}{2} \Lambda_{ijk\ell}^{1313} \Delta_j \Delta_k^* \Delta_\ell \Delta_i^* + \tag{18}$$

$$+ \Lambda_{ijk\ell}^{1113} \Phi_j \Phi_k \Delta_\ell \Delta_i^* + \Lambda_{ijk\ell}^{3331} \Phi_j \Phi_k \Delta_\ell^* \Delta_i + \Lambda_{ijk\ell}^{1133} \Phi_j \Delta_k \Phi_\ell \Delta_i^* \},$$

where

$$\chi_{0ij}^{\alpha\beta} = T \sum_n G_{0ij}^{\alpha\alpha}(z_n) G_{0ji}^{\beta\beta}(z_n), \tag{19}$$

$$\Lambda_{ijk\ell}^{\alpha\beta\gamma\delta} = T \sum_n G_{0ij}^{\alpha\alpha}(z_n) G_{0jk}^{\beta\beta}(z_n) G_{0k\ell}^{\gamma\gamma}(z_n) G_{0\ell i}^{\delta\delta}(z_n) . \tag{20}$$

The quadratic expression $F_2$ describes FM and SC separately, whereas the quartic term $F_4$ involves their interaction by biquadratically coupled order parameters. The odd terms $F_{2p+1}[\{\Phi\}, \{\Delta\}]$ are traceless in the spin subspace. The result (17) and (18) is given in the static limit, but we are left with the spatially inhomogeneous case due to disorder up to now not to be specified. The present scheme holds also for a pure system.

## 4. CPA AVERAGING. COEXISTENCE CRITERION

We derive the Ginzburg-Landau free energy functional in the static and homogeneous limit on the basis of the single-site CPA. For this purpose, all the local quantities as $n_i$, $\langle S_i^z \rangle$, $m_i$, $\Phi_i$, and $\Delta_i$ depending in principle on the whole lattice configuration are taken in the terminal point approximation, i.e., they are replaced by conditional (restricted) averages. Thus one can perform a normal self-consistent multiple scattering approach at site i to get the effective medium in CPA. The presence of random hopping of the additive (ODCPA) or multiplicative type is permitted, too. Furthermore, the averaging procedure is characterized by the concept /18,19/ of Ward identities and coherent order parameters working near the phase transitions.

The effective free energy defined by $F_{eff} = \langle F_2 + F_4 \rangle_c$ ($\langle \ldots \rangle_c$: configuration average) results in

$$F_{eff}[\Phi, \Delta] = N\{(\chi_0^{11} + \frac{1}{\tilde{U}_{eff}})\Phi^2 + (\chi_0^{13} + \frac{1}{\lambda_{eff}})|\Delta|^2 + \frac{1}{2}\tilde{\Lambda}^{1111}\Phi^4 + \tag{21}$$

$$+ \frac{1}{2}\Lambda^{1313}|\Delta|^4 + 2\Lambda^{1313}\Phi^2|\Delta|^2 \},$$

where $\Phi$ and $\Delta$ are now coherent order parameters being site-independent. N is the number of lattice sites. Having determined the noninteracting susceptibilities $\chi_0^{\alpha\beta} = \frac{1}{N}\langle \sum_{ij} \chi_{0ij}^{\alpha\beta} \rangle_c$ and

the fourthorder vertex couplings $\Lambda^{\alpha\beta\gamma\delta} = \frac{1}{N} < \sum_{ijk\ell} \Lambda^{\alpha\beta\gamma\delta}_{ijk\ell} >_c$, the coefficients of the Landau-type expansion (21) are evaluated on microscopic footing as follows:

$$\chi_0^{11} = -\int_{-\infty}^{\infty} (-\frac{\partial f}{\partial E}) \rho(E)\, dE = -\rho_T(\mu)\,, \tag{22}$$

$$\chi_0^{13} = -\rho(\mu) \ln(1.13 \frac{T_D}{T})\,, \tag{23}$$

$$\tilde{U}_{eff} = <\tilde{U}^{\nu} (g_T^{\nu})^2>_c = x\tilde{U}^A (g_T^A)^2 + (1-x)\tilde{U}^B (g_T^B)^2\,, \tag{24}$$

$$\lambda_{eff} = N' <g^{\bar{\nu}} \lambda^{\bar{\nu}\bar{\nu}'} g^{\bar{\nu}'}>_c = N' \{x^2 \lambda^{AA} (g^A)^2 + (1-x)^2 \lambda^{BB} (g^B)^2 + 2x(1-x)\lambda^{AB} g^A g^B\}\,, \tag{25}$$

$$\tilde{\Lambda}^{1111} = \frac{1}{6T^3} \frac{J_{eff}^4}{\tilde{U}_{eff}^3} \rho_T(\mu) + \Lambda^{1111}\,, \tag{26}$$

$$\Lambda^{1111} = -\frac{1}{6} \rho_T''(\mu)\,, \tag{27}$$

$$\Lambda^{1313} = \frac{7}{8} \frac{\zeta(3)\, \rho(\mu)}{\pi^2 T^2} \tag{28}$$

with the Fermi function $f(E) = (e^{(E-\mu)/T} + 1)^{-1}$, the chemical potential $\mu$, the Debye temperature $T_D$ of the averaged system, $g_T^{\nu} = (\rho_T^{\nu}(\mu))/(\rho_T(\mu))$, $g^{\nu} = (\rho^{\nu}(\mu))/(\rho(\mu))$, and $N'$ being the coordination number. $J_{eff}$ is defined like (24). The partially and totally averaged densities of states $\rho^{\nu}(E) = -\frac{1}{\pi} \mathrm{Im} <i|G_i^{\nu}(E+i0)|i>$ and $\rho(E) = -\frac{1}{\pi N} \mathrm{Im}\, \mathrm{Sp} <G_0^{11}(E+i0)>_c$, resp., are available from the CPA, where $G_i^{\nu}$ denotes the conditionally (except the site i occupied by a $\nu$-atom) averaged Green function. By symmetry arguments one proves that $\Lambda^{1113} = \Lambda^{3331} = \frac{1}{2}\Lambda^{1313}$ and $\Lambda^{1133} = \Lambda^{1313}$ entering into the coupling term in (21). Note that in deriving (22), (23), (27), and (28) the Ward identity $<G_0^{11}(z_1) G_0^{11}(z_2)>_c = -(<G_0^{11}(z_1)>_c - <G_0^{11}(z_2)>_c)/(z_1 - z_2)$ taken at $z_1 = z_2$ and $z_1 = -z_2$ in the FM and SC cases, resp., plays a key role.

Using (22) and (23) the zeros of the quadratic part of (21) correspond to the generalized Stoner criterion ($T_M$: Cure temperature)

$$\tilde{U}_{eff}\, \rho_{T_M}(\mu) = 1 \tag{29}$$

and the superconducting transition temperature

$$T_S = 1.13\, T_D\, e^{-1/(\lambda_{eff}\, \rho(\mu))}\,. \tag{30}$$

Minimization of (21) through the conditions $(\delta F_{eff})/(\delta\Phi) = 0 = 0$ and $(\delta F_{eff})/(\delta\Delta^*) = 0$ yields the coupled equations



$$(\chi_0^{11} + \frac{1}{\tilde{U}_{eff}} + 2\Lambda^{1313} |\Delta|^2)\Phi + \tilde{\Lambda}^{1111} \Phi^3 = 0\,, \tag{31}$$

$$(\chi_0^{13} + \frac{1}{\lambda_{eff}} + 2\Lambda^{1313} \Phi^2)\Delta + \Lambda^{1313} |\Delta|^2 \Lambda = 0\,. \tag{32}$$

To investigate coexistence, we are looking for nontrivial solutions of (31) and (32), so that $\Delta \neq 0$ and $\Phi \neq 0$ hold simultaneously. Such solutions become stable in the case of weak coupling [20,21] between the two order parameters. This implies the coexistence criterion

$$\tilde{\Lambda}^{1111} > 4\Lambda^{1313} \tag{33}$$

which is connected with tetracritical behaviour.

On the other hand, if the condition (33) is violated one has only a bicritical point referring to a first order transition between SC and FM.

## 5. CONCLUSIONS

Now we are dealing with two special cases of the present model which can be treated analytically.

First, we restrict ourselves to itinerant FM by putting $J^{\nu} = 0$ [18]. Moreover, we are going through $\epsilon^A \to \infty$ to the independent band limit. Then the Stoner criterion (29) with (15) and (24) is reduced to

$$U^B \rho^B{}_{T_M}(\mu) = U^B [\rho^B(\mu_0) + \frac{\pi^2}{6} (\rho^{B''}(\mu_0) - \frac{(\rho^{B'}(\mu_0))^2}{\rho^B(\mu_0)}) T_M^2] = 1\,, \tag{34}$$

where $\rho^B(E) = \rho(E)/(1-x)$, $\mu_0 = \mu(T=0)$, and the temperature dependence has been approximated up to $O(T^2)$. Choosing a semielliptic shape for the unperturbed B-band with the half-band width w we obtain

$$\rho^B(E) = \frac{2}{\pi w_x} \sqrt{1 - (\frac{E - \epsilon^B - \frac{U^B}{2} n^B}{w_x})^2}\,, \tag{35}$$

where $w = w\sqrt{1-x}$ is affected by disorder. From $U^B \rho^B|_{x=x_c}(\mu_0) = 1$ one finds the critical or percolation concentration $x_c$ at T=0; for a half-filled band ($n^B = 1$) one has immediately $x_c = 1 - (4(U^B)^2)/\pi^2 w^2$. Employing (34) and (35) for $n^B = 1$ as well as (30) on the basis of (25), which simply becomes $\lambda_{eff} = N' \lambda^{BB}$, one arrives at

$$T_M(x) \propto (x - x_c)^{1/2}\,, \quad \frac{dT_s}{dx} < 0\,. \tag{36}$$

The tendencies in (36) make an intersection point $T_M = T_S = T_c$ possible. The condition (33) can be rewritten by invoking (26), (27) ($\rho''_T(\mu) < 0$ is established in view of the concavity of (35)), and (28) in the form

$$\frac{21}{\pi^2}\zeta(3) < \left(\frac{T_c}{w_x}\right)^2 + \frac{\pi^2}{2}\left(\frac{T_c}{w_x}\right)^4 \qquad (37)$$

never to be fulfilled for realistic parameters due to $T_c/w \sim O(10^{-3})$.

Secondly, we consider only the s-d exchange in the simplest way, provided that $U^\nu = 0$, $J^A \neq 0$, $J^B = 0$, $\lambda^{AA} = \lambda^{AB} = 0$, and $\lambda^{BB} \neq 0$. For simplicity $\rho$ is assumed to be independent of disorder ($\epsilon_i = 0$, $t_{ij} = \text{const}$), i.e., $\rho(\mu) = 2/\pi w$ (cf. (35)) for a half-filled band. Then (29) with (15) and (24) gives

$$T_M = \frac{1}{2} x (J^A)^2 \rho(\mu) . \qquad (38)$$

Besides, the decreasing $\frac{dT_S}{dx} < 0$ is retained because of $\lambda_{eff} = N'(1-x)^2 \lambda^{BB}$. The coexistence criterion (33) now reads with (26) and (28):

$$\frac{21}{8}\zeta(3) < \left(\frac{J^A}{w}\right)^2 , \qquad (39)$$

where the coupling $\Lambda^{1111}$ from (27) can be neglected. Then the coefficients (26) and (28) are comparable with those in /16/. In real substances it turns out that $J^A < w$, so that the inequality (39) cannot be satisfied.

A coexistence criterion for SC and FM has been derived from the free energy functional on the mean field level without fluctuations for the augmented random Zener model. To be realistic, we have considered itinerant FM in the percolation limit and the case of localized spins. There coexistence related to a tetracritical behaviour could not be verified. This result agrees qualitatively with conclusions in /11,13/, and partly with /8/ if the component exchange integrals $J^A$ and $J^B$ have the same sign.

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REFERENCES

1. MacKay H.B. et al. Phys.Rev.Lett., 1979, 42, p. 918.
2. Fernandez-Baca J.A., Lynn J.W. J.Appl.Phys., 1981, 52, p. 2183.
3. Lynn J.W. et al. Yamada Conf. VI on Neutron Scatt. of Cond.Matt., Hakone, Jap., 1982, PA-35.
4. Gorkov L.P., Rusinov A.I. Zh.Eksp.Teor.Fiz., 1964, 46, p.1363.
5. Izyumov Yu.A., Skryabin Yu.N. phys.stat.sol., 1974, B61, p.9.
6. Smit W.A., Vertogen G., Kraak J. Physica, 1974, 74, p. 97.
7. Sakurai A. Sol.St.Comm., 1978, 25, p. 867.
8. Balseiro C.A., Falicov L.M. Phys.Rev., 1979, B19, p. 2548.
9. Bogolubov N.N. Jr., Ermilov A.N., Kurbatov A.M. Physica, 1980, 101A, p. 613.
10. Tonchev N.S., Bramkov I.G. phys.stat.sol., 1980, B102, p. 179.
11. Andreychenko V.B., Burmistrov S.N. Pisma v. Zh.E.T.F. 1978, 28, p. 530.
12. Hornreich R.M., Schuster H.G. Phys.Lett., 1979, 70A, p.143.
13. Grewe N., Schuh B. Phys.Rev., 1980, B22, p. 3183.
14. Laptev V.M., Skryabin Yu.M. Fiz. Metallov i Metallovedenie, 1981, 51, p. 69.
15. Suhl H. J.Less-Common Metals, 1978, 62, p. 225.
16. Brankov J.G., Tonchev N.S. Physica, 1981, 108A, p. 459.
17. Weithege W., Entel P., Mühlschlegel B. Z.Phys., 1982, B47, p. 35.
18. Kolley E., Kolley W. phys.stat.sol., 1982, B114, p. 221.
19. Beckmann A., Kolley E., Kolley W. phys.stat.sol., 1982, B110, p. 557.
20. Lyuksyutov I.F., Pokrovskii V.L., Khmelnitskii D.E. Zh.Eksp. Teor. Fiz., 1975, 69, p. 1817.
21. Imry Y. J.Phys.C: Solid State Phys., 1975, 8, p. 567.

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Сверхпроводимость и ферромагнетизм в обобщенной неупорядоченной модели Зенера

Исследуется взаимное влияние сверхпроводимости и ферромагнетизма в случайных бинарных сплавах с учетом сильно связанных электронов и локализованных магнитных моментов. Получено петлевое разложение свободной энергии для комбинированного притяжения контактного типа и модели Зенера. Вычислены коэффициенты функционала Гинзбурга-Ландау с двумя биквадратичными связанными параметрами порядка с помощью усреднения по конфигурациям в рамках приближения когерентного потенциала. Сформулирован критерий сосуществования. Получены оценки, показывающие, что в предельных случаях зонной и локализованной спиновой модели имеется только бикритическое поведение фазовой диаграммы.

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Kolley E., Kolley W. E17-82-806

Superconductivity and Ferromagnetism in the Extended Random Zener Model

The mutual influence of superconductivity and ferromagnetism in random binary alloys is studied in the presence of tightly bound electrons and localized magnetic moments. The loop expansion of the free energy is derived microscopically by combining a contact-type attraction and the Zener model. The coefficients of the Ginzburg-Landau functional with two biquadratically coupled order parameters are obtained after configuration averaging within the coherent potential approximation. A criterion for coexistence is formulated. Estimations in the limits of itinerant and localized spin models give rise only to bicritical behaviour.

The investigation has been performed at the Laboratory of Theoretical Physics, JINR.

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