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ON THE PHASE DIAGRAM OF HIGH-T_c SUPERCONDUCTIVE GLASS MODEL

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1. Introduction

Glassy behaviour of high-T_c superconductors (HTSC) first reported by Müller et al. /1/ was then studied by various methods both in the ceramic samples /2,3/ and single crystals /4,5/. Morgenstern et al. /6/ presented numerical simulations for a 2-D system of clusters coupled by Josephson tunnelling which reproduced main features found in high-T_c superconductor experiments. Important is the proposal made of a phase diagram in the H-T plane showing possible states of the system. Analogous to the experiments of Müller et al. /1/ the numerical simulations led to the existence of a quasi-de Almeida-Thouless(AT) line separating the superconducting-glass (SCG) phase and the normal conducting regime at $H < H_c^u$ ($H_c^u \simeq 0.5$ the magnetic field in units of $2\pi/\Phi_0$). But at $H > H_c^u$ together with the SCG phase the Josephson spin-glass (JSG) phase exists.

The numerical simulations /6/ were performed for the magnetic field up to H = 0.4 since at higher fields equilibrium is hard to attain. Therefore, phase boundaries between the JSG phase and the SCG phase and normal conductor phase are to be refined.

A phase transition line $T_c(H)$ was calculated in /7/ analytically for an analogous system described with an XY model. The results of paper /7/ are valid at strong fields $H > H_c^u$ leading to strong frustration effects, i.e transition to the JSG phase. Following the method exploited by Vinokur et al. /7/ we study here the quasiequilibrium (static) properties of the model for weak coupling superconducting clusters /6/ in an arbitrary magnetic field. Phase boundary is derived separating between the normal conductor phase and Meissner phase at $H < H_c^\ell$, SCG phase at $H_c^\ell < H < H_c^u$ and JSG phase at $H > H_c^u$.

2. The model

Following Morgenstern et al /6/, let us consider an array of N superconducting (SC) clusters coupled by Josephson tunnelling. A cluster is defined as a region of coherent phases in the superconductor. The considered model allows clusters <u>inside</u> the physical grains in contrast to Ebner and Stroud /8/ and John and Lubensky /9/ . The i-th cluster is located at $\vec{x}_i$ and is descibed by a complex order parameter

$$\Psi_i = \Delta_i e^{i\phi_i}. \tag{1}$$

Here Δ_i is the complex energy gap and ϕ_i is the phase of the Cooper pairs in each cluster i . The ensemble of SC clusters in magnetic field is then governed by a 2-D XY-model Hamiltonian

$$H = - \sum_{ij} K_{ij} \cos (\phi_i - \phi_j - A_{ij}) , \tag{2}$$

where

$$A_{ij} = \frac{2\pi}{\phi_o} \int_i^j \vec{A}\cdot d\vec{l} . \tag{3}$$

We consider the case when coupling energies between domains of uniform phases i and j $K_{ij} = J$ for nearest neighbours only and otherwise zero. For Josephson tunnelling we have

$$K_{ij}(T) = \frac{\Delta(T)}{R_{ij}} \tanh\left(\frac{\Delta(T)}{2R_{ij}T}\right) , \tag{4}$$

where R_{ij} is the resistance between domains i and j in their normal state.

In contrast to the ordinary XY-SG model /10/ frustration in our system (2) is due to the external magnetic field (3). Applying the magnetic field in the Z-direction we have

$$A_{ij} = \frac{2\pi}{\phi_o} H \frac{x_i + x_j}{2} (y_j - y_i), \tag{5}$$

where x_i, y_i are the actual coordinates of i-th cluster and $\phi_o = \frac{hc}{2e}$ is an elementary flux quantum. As in /6/ the 2-D case for site disorder will be considered. We will describe this type of disorder by the Gauss distribution function over the cluster coordinates (and not by random J distribution as in the SG theory /10/)

$$P(x_i,y_i) = \frac{1}{2\pi\sigma} e^{-\frac{x_i^2}{2\sigma} - \frac{y_i^2}{2\sigma}} \tag{6}$$

with the mean-square area of SC cluster S being in the form $S = \pi\sigma$. So, the mean area of the cluster on the random Josephson lattice is, in fact, the only parameter of our theory.

A justification of the model (2) can be found in /6/. Here we should only remember the most important limitations of the model. They are :

(i) the Josephson energy J is smaller than the transition temperature for a single cluster (otherwise the dependence of J on temperature need to be taken into account),

(ii) the London penetration depth is large compared to the cluster size and their separation (otherwise one ought to discriminate between the local and applied magnetic fields in (2)).

In a recent paper /7/ a similar to (2) model as a model of Josephson spin glass was considered. Although the authors of that paper dealt with the region of very strong fields only (in fact, they were in the regime of complete frustration, where the SCG phase was not existent already), their analytical results were very interesting. We take the method they used as a guide line in our own research.

3. Density of States

Following Vinokur et al. /7/ let us rewrite the Hamiltonian (2) in the form of a 2-D XY-model for Josephson spins S_i

$$H = - \mathrm{Re} \sum_{ij} S_i^* J_{ij} S_j , \tag{7}$$

where

$$S_i = e^{i\phi_i} \quad , \quad J_{ij} = J e^{iA_{ij}} . \tag{8}$$

It is important to stress that in contrast to Vinokur et al /7/, our exchange energies J_{ij} automatically permit the symmetry $J_{ij} = J_{ji}^*$ (due to the antisymmetry of A_{ij} (5)). Let us consider the spectrum of the random (via A_{ij}) matrix J_{ij} , because this spectrum (as we are to see below) is connected with the critical behaviour of the model (7). The density of states for

J_{ij} is defined in a standard manner by averaging the one-point Green function $g_{ii}(E)$ according to (6)

$$\rho(E) = -\frac{1}{\pi}\,\mathrm{Im}\,\overline{g_{ii}(E + i0)} \quad , \tag{9}$$

where

$$g_{ij}(E) = (\delta_{ij}E - J_{ij})^{-1}. \tag{10}$$

To find the locator $g(E) \equiv \overline{g_{ii}(E)}$, we need to solve (in the virtual crystal approximation) the Dyson equation

$$g(E) = g_o(E) + g_o(E)\cdot\Sigma\cdot g(E)\,. \tag{11}$$

Here, $g_o(E) = E^{-1}$ is the "bare" Green function, and the self-energy Σ has the form

$$\Sigma(H) = \sum_{n=1}^{\infty} K_n (Ng)^{n-1} \tag{12}$$

A set of irreducible correlators K_n is defined as follows (for odd n)

$$K_{2n-1} \equiv \overline{J_{ij_1}J_{j_1j_2}\cdots J_{j_{2n-2}i}} = \frac{J^{2n-1}}{(1 + v^2)^{(2n-1)/2}} \tag{13}$$

$$v = H/H_o, \quad H_o = \Phi_o/2S, \quad S = \pi\sigma. \tag{14}$$

On Fig. 1 the diagrammatic version of Dyson equation (11) is presented. The set of one-loop diagrams can be summarized exactly and for the self-energy $\Sigma(H)$ in arbitrary field one gets

$$\Sigma(H) = \frac{J^2Ng + JU}{1 - J^2N^2U^2g^2} \quad , \quad U = (1 + v^2)^{-1/2}. \tag{15}$$

By solving the Dyson equation (11) we have the cubic equation on $g(E)$

$$Eg(E) = 1 + \Sigma(H)g(E)\,. \tag{16}$$

An upper boundary of the spectrum E_o in arbitrary field is defined by the appearance of an imaginary part in the solution of equation (16). When $E \to E_o$ such solutions have the root singularity and, hence, $g(E_o)$ is the solution of equation $\partial E(g)/\partial g = 0$, i.e.

$$J^2Ng^2(E_o)\frac{1 + J^2NU^2g^2(E_o) + 2JNU^2g(E_o)}{\left[1 - J^2N^2U^2g^2(E_o)\right]^2} = 1\,. \tag{17}$$

As will be shown in Sect. 4, the quantity $g(E_o)$ determines the critical temperature in the system (7).

In the limiting cases of zero and infinite magnetic fields equations (15) and (16) yield

a) $H = 0$ $(v = 0,\ U = 1)$

$$\Sigma = \frac{J}{1 - JNg} \quad , \quad g(E) \simeq E^{-1} \quad (N \gg 1),$$

$$\rho_o(E) = \delta(E) \quad , \quad E_o = JN \tag{18a}$$

b) $H = \infty$ $(v = \infty,\ U = 0)$

$$\Sigma = J^2Ng, \quad g(E) = 1/2\left(\frac{E}{J^2N} - i\sqrt{\frac{4}{J^2N} - \frac{E^2}{J^4N^2}}\right),$$

$$\rho_\infty(E) = \frac{1}{2\pi NJ^2}\sqrt{4NJ^2 - E^2}, \qquad E_o = 2J\sqrt{N}\,. \tag{18b}$$

In the case of intermediate fields one may distinguish three characteristic regions

I. Reversible diamagnetic region $(H < H_{c1})$

$$\rho(E) \simeq \delta(E)\cdot\left(1 - \frac{v^2}{2}\right), \quad E_o \simeq JN\left(1 - \frac{2}{\sqrt{3}}v^2\right) \tag{19}$$

$$H_{c1} = 3/4\cdot H_o \equiv H_c^{\ell}\,. \tag{20}$$

II. Irreversible region of superconductive glass $(H_{c2} > H > H_{c1})$

$$\rho(E) \simeq \delta(E)\cdot\left(1 - \frac{v^2}{2}\right) + \frac{1}{2\pi NJ^2}\sqrt{E_o^2 - E^2}\,,$$

$$E_o = -\frac{JN}{\sqrt{3}}\,(6v^2)^{1/3} + JN \tag{21}$$

$$H_{c2} = 15\cdot H_o \equiv H_c^{u}\,. \tag{22}$$

III. Irreversible region of frozen Josephson spin glass $(H > H_{c2})$

$$\rho(E) \simeq \frac{1}{2\pi N J^2} \sqrt{E_0^2 - E^2} \quad , \quad E_0 = 2J\sqrt{N}\cdot\left(1 + \frac{N}{2v^2}\right) \cdot \tag{23}$$

At $H \gg H_{c2}$ $E_0 = 2J\sqrt{N}$, i.e we deal with the regime of complete frustration /7/.

4. Phase Diagram

The transition temperature $T_c(H)$ is defined as usual by the appearance of singularity in the behaviour of the generalized susceptibility. The mean value $m_i = \langle S_i \rangle$ is described by the Thoules-Anderson-Palmer equation /11/ in a fictive field h_i

$$m_i = \frac{1}{2T}\left(\sum_j J_{ij} m_j + h_i\right) - \alpha m_i . \tag{24}$$

The correlative to the mean field theory correction $\alpha(H)$ was first introduced in /11/ for the correct description of thermodynamic properties of the SG-model. In the case of arbitrary fields this constant is determined in full analogy to the one-point Green function by the substitution of $\overline{g_{ii}(E)}$ for $\chi_{ii} = 1/2T$. By taking into account equations (13)-(16) one obtains

$$\alpha(H) = \frac{J}{2T}\left(\frac{JN}{2T} + U\right)\left(1 - \frac{J^2N^2U^2}{4T^2}\right)^{-1} ,$$

where

$$U(H) = \left(1 + \frac{H^2}{H_0^2}\right)^{-1/2} .$$

Equation (24) for the generalized susceptibility $\chi_{ij} = \frac{\partial m_i}{\partial h_j}$ yields

$$\chi_{ij} = \left[2T(1+\alpha)\cdot\delta_{ij} - J_{ij}\right]^{-1} \equiv g_{ij}(E = 2T(1+\alpha)) . \tag{25}$$

So the equation on $T_c(H)$ has the form /7/

$$2T_c(1 + \alpha(T_c)) = E_0 , \tag{26}$$

where the upper boundary E_0 for some limiting cases is given by (19)-(23). Due to (25) and (26) $\chi_{ii}(T_c) = \overline{g_{ii}(E_0)} = g(E_0)$. On the other hand in the model of "hard spins" ($|S_i| = 1$) in the high temperature phase, from linear response relation $2T\cdot\chi_{ij} = \langle\langle S_i S_j^* \rangle\rangle$, we have /7/

$$\chi_{ii} = 1/2T . \tag{27}$$

Finally the equation on T_c takes the form

$$g(E_0) = 1/2T_c . \tag{28}$$

In general the dependence of T_c on H in arbitrary fields is determined by (17) and (28).

In the regions I-II-III this dependence is described by the following explicit formulae

I. $H < H_{c1}$ (region of nearly ideal diamagnetism)

$$\frac{T_c(0) - T_c(H)}{T_c(0)} = \frac{2}{3}\left(\frac{H}{H_0}\right)^2, \quad H_{c1} = (3/4)H_0 \tag{29}$$

$$T_c(0) = \frac{JN}{2} . \tag{30}$$

II. $H_{c2} > H > H_{c1}$ (region of SCG phase : AT line)

$$\frac{T_c(0) - T_c(H)}{T_c(0)} = \sqrt[3]{\frac{2}{\sqrt{3}\cdot H_0^2}}\, H^{2/3}, \quad H_{c2} = 15 H_0 . \tag{31}$$

III. $H > H_{c2}$ (region of JSG phase : strong frustration)

$$T_c(H) = T_c(\infty)\left(1 + \frac{3NH_0^2}{2H^2}\right), \tag{32}$$

$$T_c(\infty) = \frac{J\sqrt{N}}{2} . \tag{33}$$

At moderate fields $(H \simeq H_0)$

$$T_c = \frac{JN}{\sqrt{8}} . \tag{34}$$

III. Irreversible region of frozen Josephson spin glass $(H > H_{c2})$

$$\rho(E) \simeq \frac{1}{2\pi N J^2} \sqrt{E_0^2 - E^2} \quad , \quad E_0 = 2 J \sqrt{N} \cdot (1 + \frac{N}{2v^2}) \cdot \qquad (23)$$

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$$\alpha(H) = \frac{J}{2T} \left(\frac{JN}{2T} + U \right) \left(1 - \frac{J^2 N^2 U^2}{4T^2} \right)^{-1} ,$$

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At moderate fields $(H \simeq H_0)$

$$T_c = \frac{JN}{\sqrt{8}} . \qquad (34)$$

The analytical results (29)-(33) are reflected on the phase diagram (Fig. 2). To compare it with the experimental one for the La-Ba-Cu-O system we have to translate our units into the units of paper /6/. This yields, in particular, $H_o = 0.05T$. In view of (14) for the actual value of cluster area S we have

$$S = \frac{\Phi_o}{2H_o} \simeq 0.02\ \mu^2 . \tag{35}$$

This is in reasonable agreement with the commonly used experimental estimates $S = 0.01 \div 0.1\ \mu^2$ /1-5, 12, 13/.

5. Diamagnetic Response

For the more correct identification of the phases I-II-III shown in Fig. 2 let us consider the behaviour of isothermal (field-cooled) magnetization for these phases versus applied magnetic field.

In our model the role of magnetization plays an induced (by Josephson supercurrents I_{ij}) magnetic moment of the SC cluster ensemble /7/

$$\mu_z = \frac{1}{2c} \sum_{ij} I_{ij}(x_i y_j - x_j y_i) , \tag{36}$$

where

$$I_{ij} = \frac{2e}{\hbar} J \sin\Phi_{ij} \tag{37}$$

$$\Phi_{ij} = \phi_i - \phi_j - A_{ij} . \tag{38}$$

Here A_{ij} is defined by (5).

By performing thermal averaging via (24) and random averaging via (6) we find the magnetization per cluster area S as

$$M \equiv \overline{\langle \mu_z \rangle}/S = -\chi(t,T,H)H , \tag{39}$$

where

$$\chi(t,T,H) = \chi_o(t,T)\left(1 + \frac{H^2}{H_o^2}\right)^{-3/2} , \tag{40}$$

$$\chi_o(t,T) = \frac{\chi_o(T)}{N} \sum_i D_{ii}(t) , \tag{41}$$

$$D_{ij}(t) = \overline{\langle\langle S_i(t) S_j^* \rangle\rangle} , \qquad \chi_o(T) = \frac{4JSN}{\Phi_o^2} . \tag{42}$$

In the mean field approximation from (27) follows that in the static limit

$$D_{ii}(0) = 2T \cdot \chi_{ii}(T) = 1 . \tag{43}$$

Thus in relatively small fields ($H \lesssim H_{c1}$) the relations (39)-(43) lead to a linear diamagnetic susceptibility in the form

$$\chi_{ac}^{DM} = \left.\frac{\partial M}{\partial H}\right|_T = -\chi_o(T) . \tag{44}$$

According to (39) - (43) the nonlinear effects in the behaviour of susceptibility become more essential with increasing field. This, in fact, was observed experimentally for the region of SCG ($H_{c2} > H > H_{c1}$). When $H \to H_{c2}$ (the JSG phase on Fig.2) the magnetization rather rapidly tends to zero. The behaviour of equilibrium (static) magnetization versus magnetic field is shown in detail in Fig. 3.

Conforming with the units of paper /1/ we have in accordance with (35)

$$M_o = \frac{2JN}{H_o \Phi_o} \simeq 1.2 \times 10^{-2}\ \text{emu/g} . \tag{45}$$

This, by the way, allows estimation of the coupling energy between clusters J . Namely, from (45) and (39) one obtains for N = 16 /6/ $J \simeq 3.7$ K. On the other hand as follows from (30) for the experimenal value of $T_c(0) = 28$ K /6/ the Josephson energy is

$$J = \frac{2T_c(0)}{N} \simeq 3.5\ \text{K} . \tag{46}$$

So we may say that there is a correlation between the superconductive and magnetic (glassy) properties in the high-T_c SCG model.

6. Conclusions

The phase transition temperature as a function of applied magnetic field is derived in this paper for a model of superconducting clusters coupled by Josephson tunnelling. According to the estimates reported in Sect.4 of this paper the mean area of the cluster $S \simeq 0.02\ \mu^2$ and their coupling energy $J = 3.6$ K. The $T_c(H)$ line obtained reproduces the results of the numerical simulations /6/ at $H \leq H_c^\ell$ and extends to the fields $H_c^u \geq H > H_c^\ell$ giving additional information.

The temperature $T_c(H)$ is defined by the apperance of irreversible (glassy) features in a system of random Josephson contacts. Strictly speaking this temperature is different from the critical temperature at which occurs the superconducting transition and which should be determined from the temperature dependence of resistivity. However, according to the experiments /1/ and numerical simulations /6/ these two temperatures are hardly distinguishable. The latter showed this difference to be proportional to $N^{-1/2}$, where N is the number of superconducting clusters in the model. Since the results of this paper are obtained in the thermodynamic limit of $N \gg 1$, we apparently cannot distinguish these two critical temperatures.

Data reported in papers /1,6/ have not only indicated close values for these temperatures, but also pointed to their similar behaviour following the external field variation. At fields below the upper critical field H_c^u (i.e. in the SGG region) a $T_c(H)$ line of the AT type separates between the superconducting phase and normal conductor region, i.e. the transition temperature follows the magnetic field variation (in the magnetic field range $H_c^u \gtrsim H > H_c^\ell$) as $H^{2/3}$. Only at $H \to H_c^u$ the above temperatures start to behave differently, the difference being rather essential. At higher $H > H_c^u$ the $T_c(H)$ approaches saturation corresponding to the JSG phase with a frozen disorder and zero induced magnetic moment, while the SC transition temperature on the contrary must suffer strong suppression at $H \to H_c^u$ (see also /7/). And so, if a number of features of the lower critical field H_c^ℓ allows its identification with the first critical field H_{c1} of the second order superconductor, the situation with the upper critical field H_c^u appears more complicate. In fact, the H_c^ℓ and H_c^u critical fields are introduced in connection with the magnetization behaviour which plays the role of a critical current for the model under discussion (see Fig. 3). Therefore, experimental study of the magnetic field dependence of T_c under transition from the SGG to the JSG phase is of interest.

We are indebted to Prof. K.K.Likharev for useful discussions.

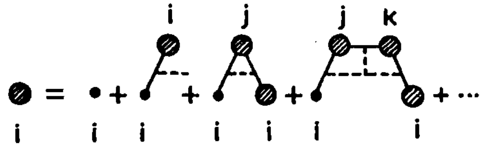

Fig. 1. Dyson equation on one-point Green function $\overline{g_{ii}(E)}$ (shaded circles). Thick points for "bare" Green function $\overline{g_{ii}^{o}}$, straight line for J_{ij}, dotted line for averaging over disorder.

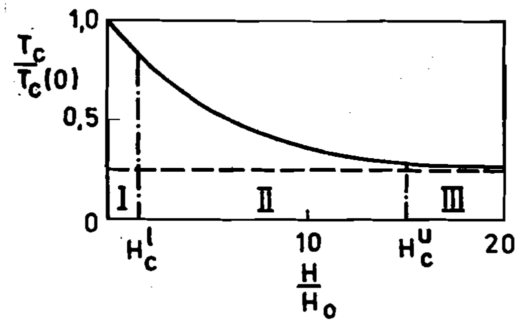

Fig. 2. Phase diagram in the H-T plane. Reversible diamagnetic phase (I). Between the H_c^ℓ and H_c^u superconductive glass phase (II). Above H_c^u the Josephson spin glass phase (III). Irreversible effects are separated from reversible ones by an AT line.

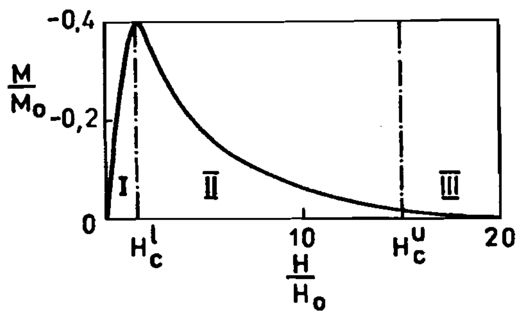

Fig. 3. Equilibrium (field-cooled) magnetization M versus the magnetic field H (cf. Fig.2).

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О фазовой диаграмме в модели
высокотемпературного сверхпроводящего стекла

В двумерной XY модели джозефсоновских спинов вычислена зависимость температуры фазового перехода T_c и изотермической намагниченности от внешнего магнитного поля H при $0 \leq H < \infty$. Показано, что на плоскости (T,H) имеется три области, различающиеся характером зависимости T_c от H: диамагнитная область, область сверхпроводящего стекла и область джозефсоновского спинового стекла. Результаты качественно согласуются с данными экспериментов и численного моделирования для "новых" сверхпроводников.

Работа выполнена в Лаборатории нейтронной физики ОИЯИ.

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On the Phase Diagram of High-T_c
Superconductive Glass Model

The transition temperature T_c and isothermic magnetization are calculated as functions of applied magnetic field in the frame of the 2-D XY Josephson glass model. Three characteristic regions are shown to be distinguishable in the H-T plane: the diamagnetic region, region of superconducting glass and region of Josephson spin glass. The results are in qualitative agreement with experimental data and the results of numerical simulations for "new" superconductors.

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

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