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NONEQUILIBRIUM BEHAVIOUR AND NONERGODICITY OF HIGH-T$_c$ SUPERCONDUCTIVE GLASS MODEL

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1. INTRODUCTION

One of the most exciting questions high-T_c superconductivity physics deals with is that of the relation between superconducting and structural properties. The glassy behaviour of new superconductive materials first reported by Müller et al.[1] and afterwards confirmed in a number of papers[2–8, 12] attracts special attention of experimentalists and theorists alike. Morgenstern et al.[9] by numerical simulations reproduced the main observable features (both equilibrium and nonequilibrium) in the behaviour of new superconductors and established the phase diagram in the plane H-T for applied magnetic fields up to $H \leq H_c^u$. An analogous model of frustrated Josephson spins on 2-D disordered lattice was treated by Vinokur et al.[10] at high fields ($H > H_c^u$).

In our previous paper[11] a phase line $T_c(H)$ in an arbitrary magnetic field was calculated analytically. One may distinguish three characteristic regions on the phase diagram: the quasireversible Meissner phase ($H < H_c^l$), superconductive glass phase ($H_c^u > H > H_c^l$), and Josephson spin glass phase ($H > H_c^u$). The obtained in[11] phase diagram serves as basis for exploration of the dynamical behaviour of glass-like superconductors. At present there are available a number of papers on experimental studies of the nonergodic and relaxation properties of magnetization including the muon spin relaxation method[12] in oxide superconductors (see, e.g.[1–8]). In these experiments the difference between field cooled (FC) and zero field cooled (ZFC) magnetic measurements due to the nonergodicity is observed below $T_c(H)$.

Numerous investigations of the long-time relaxation behaviour of the remanent magnetization indicate the nonexponential (logarithmic, power, or Kohlrausch-like) decay law with essential temperature and applied magnetic field dependence of α- and β-relaxation exponents.

The present paper is devoted to theoretical treatment of the dynamical properties of high-T superconductive glass (SCG) model with the aim to describe the above mentioned experimental data. The paper is organized as follows. In Sect.2 we consider the nonergodic behaviour of the model. The comparison of the theoretically obtained temperature field dependence of nonergodicity parameter is given versus the experimental FC and ZFC magnetization data for La-Ba-Cu-O [1], Y-Ba-Cu-O [4], and Gd-Ba-Cu-O [7] systems. In Sect.3 via the mode-coupling approximation the critical low-frequency dynamics of the model is discussed. The dynamical scaling as well as the critical crossover in behaviour of Josephson spin correlation function are present. In Sect. 4 we study the non-exponential (logarithmic and power) behaviour of the long-time relaxation of the remanent magnetization. The temperature and applied magnetic field dependence of calculated magnetic viscosity coefficient and of the power relaxation exponent are given in comparison with the corresponding experimental data [5,6,8]. In conclusion we discuss the obtained results as well as some predictions of the SCG model and experimental possibilities for their confirmation.

2. NONERGODIC BEHAVIOUR OF SCG MODEL

Let us consider the SCG model [9] with a Hamiltonian [9-11]

$$H = -J\sum_{ij}\cos\phi_{ij} = -\mathrm{Re}\sum_{ij} J_{ij}\, S_i^* S_j\,, \tag{1}$$

where

$$\phi_{ij} = \phi_i - \phi_j - A_{ij}\,, \tag{2}$$

$$A_{ij} = \frac{\pi H}{\phi_0}(x_i + x_j)(y_j - y_i)\,, \tag{3}$$

$$S_i = \exp(i\phi_i)\,,\quad J_{ij} = J\exp(iA_{ij})\,. \tag{4}$$

Here ϕ_i is the phase of the Cooper pairs in clusters i which is described by random coordinates (x_i, y_i). An ensemble of Josephson spins S_i is in an applied magnetic field $\mathbf{H} = (0,0,H)$.

To describe the dynamical behaviour of the model (1) we need an equation of motion for the spin S_i. According to [10] the superconductive current through the Josephson juction between clusters i and j is defined as

$$I_{ij} = \frac{h}{2eR}\frac{d\phi_{ij}}{dt} + \frac{2eJ}{h}\sin\phi_{ij}\,. \tag{5}$$

Here R is the resistance between domains i and j in their normal state, $J(T) = \frac{\Delta(T)}{R}\,\mathrm{th}\left(\frac{\Delta(T)}{2RT}\right)$ is the temperature dependent Josephson energy.

The equation of motion for the superconducting phase of the i-th cluster ϕ_i follows from the conservation laws [10]

$$\sum_i I_{ij} = 0\,,\quad \sum_j I_{ij} = 0\,. \tag{6}$$

In the approximation $N \gg 1$, eq.(6) yields

$$\frac{hN}{2eR}\frac{d\phi_i}{dt} + \frac{2eJ}{h}\sum_j \sin\phi_{ij} = 0\,. \tag{7}$$

Let us introduce a dimensionless time parameter

$$\tilde{t} = \frac{4e^2RT}{h^2N}\,t\,. \tag{8}$$

Then, in view of eqs. (7) and (8), the equation of motion for Josephson spin S_i takes the form

$$\dot{S}_i = \frac{1}{2T}\sum_j (J_{ij}S_j - J_{ij}^* S_j^* S_i^2)\,. \tag{9}$$

Hereafter we put $\tilde{t} \equiv t$ keeping in mind the definition (8).

To study the nonergodic properties of the system (1) let us consider a low-frequency behaviour of the correlation function [11]

$$D_{ij}(t) = \overline{\langle S_i^*(t) S_j \rangle}\,. \tag{10}$$

Here the bar denotes the Gauss-like averaging over random cluster coordinates (x_i, y_i) with a mean square deviation σ.

According to paper [11] the correlator (10) determines all nonequilibrium properties of the remanent magnetization M(t) which in the model (1) has the form

$$M(t) = -\chi(t,T,H)H\,, \tag{11}$$

$$\chi(t,T,H) = \chi_0(t,T)\left(1 + \frac{H^2}{H_0^2}\right)^{-3/2}, \tag{12}$$

$$\chi_0(t,T) = \frac{1}{N}\sum_i D_{ii}(t)\,\chi_0(T)\,, \tag{13}$$

$$\chi_0(T) = 4SNJ(T)/\phi_0^2\,,\quad \phi_0 = hc/2e\,. \tag{14}$$

In paper [11] the field behaviour of the isothermal magnetization was obtained in static limit, when $D_{ij}(0) = 2T\chi_{ij}(T)$ and $\chi_{ij} = 1/2T$ (in the high-temperature approximation). To calculate the isothermal relaxation function $D_{ij}(t)$ in low-frequency limit we use the Mori-like projection technique [15] which is well known in the theory of spin glasses [13] and proton glasses [14]. In the mode-coupling approximation the self-consistent system on $D_{ij}(t)$ can be constructed. Let us introduce the Laplace transform

$$D_{ij}(z) \equiv \overline{((S_i^*|S_j))} = i\int_0^\infty dt^{izt} D_{ij}(t). \tag{15}$$

Then the continued fraction expansion for $D_{ij}(z)$ leads to the expression

$$D_q(z) = -(z-(z+K_q(z))^{-1})^{-1}, \tag{16}$$

where

$$D_q(z) = \frac{1}{N}\sum_{jk} e^{iq(j-k)} D_{jk}(t), \quad D_q(t=0) = 1. \tag{17}$$

In eq. (16) the current-current correlator $K_{ij}(z)$ is determined as a second momentum (see [14, 15]) according to the equation of motion (9), and in the pair spin correlation approximation has the form

$$K_{ij}(z) \equiv \overline{((\ddot{S}_i^*|\ddot{S}_j))}_2 \simeq f(T,H)(D_{ij}(z) + i\int_0^\infty dt\, e^{izt} \sum_{k\ell} D_{ik}(t)D_{k\ell}(t)D_{\ell j}(t)), \tag{18}$$

where

$$f(T,H) = \frac{J^2N}{4T^2}\left(1 + \frac{J^2N^2U^2}{4T^2} + \frac{JNU^3}{T}\right)\left(1 - \frac{J^2N^2U^2}{4T^2}\right)^{-2}. \tag{19}$$

$$U(H) = (1 + H^2/H_o^2)^{-1/2}, \quad H_o = \phi_o/2S. \tag{20}$$

Here $S = \pi\sigma$ is the mean square area of the superconductive cluster [11].

As is known [16], the nonergodic behaviour is characterized by a nonzero value of the correlator $D_{ij}(t)$ when $t\to\infty$, i.e. by a parameter

$$L_{ij} = \lim_{t\to\infty} D_{ij}(t) \neq 0. \tag{21}$$

And it is the nonergodicity parameter L that is connected with the experimentally observable difference between FC and ZFC susceptibilities [13], namely

$$L \sim T(\chi_{FC} - \chi_{ZFC}). \tag{22}$$

The Fourier transform of L_{ij} may be extracted from the relaxation function $D_q(z)$ in the way

$$L_q = \lim_{z\to io}(-zD_q(z)) = (1 + m_q^{-1})^{-1}, \tag{23}$$

where

$$m_q \equiv \lim_{z\to io}(-zK_q(z)). \tag{24}$$

According to the fluctuation-dissipation theorem the appearance of the pole for $D_q(z)$ leads to a singular behaviour of the spectral function

$$D''_{ij}(\omega) = \pi L_{ij}\delta(\omega) + \text{reg}. \tag{25}$$

To simplify further calculation we neglect the dispersion of the relaxation kernel by adopting the approximation $m_{ij} = \delta_{ij}m$ which with account for eqs. (18) and (24) yields

$$m = f(T,H)(L + L^3), \tag{26}$$

where $L \equiv L_{ii}$ denotes the local nonergodicity parameter. Thus, from eqs.(18)-(26) we obtain the self-consitent equation for L

$$L = \frac{1}{N}\sum_q L_q = \frac{f(T,H)(L+L^3)}{1 + f(T,H)(L+L^3)}. \tag{27}$$

The transition temperature $T_c(H)$ to the nonergodic state is determined by the equation

$$L(T_c, H) = 0, \tag{28}$$

or equivalently by the equation $f(T_c,H) = 1$. It is worth noting that the critical temperature $T_c(H)$ from (28) coincides exactly with the glassy temperature obtained in the SCG model [11]. It is easy to check that at $T \geq T_c(H)$ the equation (27) has the only stable solution, namely a trivial one $L = 0$. Taking into account an explicit form of f(T,H) (see (19)), let us rewrite eq.(27) in the form

$$L^3 - L^2 + L - (1 - T^2/T_c^2(H)) = 0. \tag{29}$$

The general form of $T_c(H)$ is given by (28). In what follows we shall consider some special cases for three regions on the phase diagram T-H, namely[11]

(I) $H < H_c^{\ell}$ (the Meissner phase)

$$\frac{T_c(0) - T_c(H)}{T_c(0)} = \frac{2}{\sqrt{3}} \left(\frac{H}{H_0}\right)^2 , \quad H_c^{\ell} = 3H_0/4 \tag{30}$$

$$T_c(0) = JN/2 . \tag{31}$$

(II) $H_c^u > H > H_c^{\ell}$ (the superconductive glass phase)

$$\frac{T_c(0) - T_c(H)}{T_c(0)} = \frac{1}{\sqrt{3}} (6/H_o^2)^{1/3} H^{2/3} , \tag{32}$$

$$H_c^u = 15H_o . \tag{33}$$

(III) $H > H_c^u$ (the Josephson spin glass phase)

$$T_c(H) = T_c(\infty)(1 + 3NH_o^2/2H^2) , \tag{34}$$

$$T_c(\infty) = J\sqrt{N/2} . \tag{35}$$

Here N is the number of superconductive clusters, $H_c^{\ell}(H_c^u)$ is the lower (upper) critical magnetic field (see[9,11]).

In Fig.1 the calculated by (29) values of L versus $T/T_c(0)$ are shown (solid lines) in the regions (I)-(III). For comparison, experimental data on FC and ZFC susceptibility differences are plotted (dotted lines) for La-Ba-Cu-O[1], Y-Ba-Cu-O[4], and Gd-Ba-Cu-O[7], respectively. As we see, at least in the region of not very low temperatures (where ordinary activation processes may play an essential role) there is a good qualitative agreement of theoretical calculations and experimental results. It is interesting to note that as follows from (29) the dependence of L on $T/T_c(H)$ has a universal shape in the whole region of applied magnetic field variation (see Fig.2). As we are to see in Sect. 3 this temperature-field dependence is presented also in the critical low-frequency region.

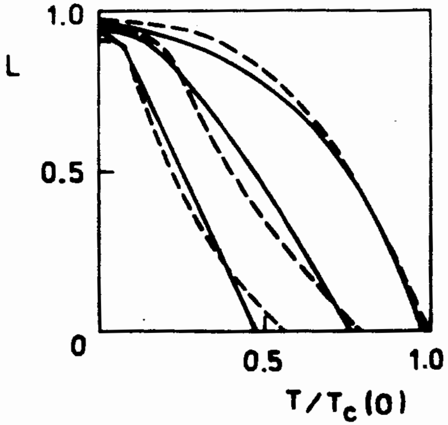

Fig.1. The dependence of the nonergodicity parameter L (29) on temperature (solid lines). The dotted lines indicate the experimental data on differences $\chi_{FC} - \chi_{ZFC}$ for oxide superconductors (from right to left): (1) $GdBa_2Cu_3O_x$, H=10 G[7], (2) La_2BaCuO_4, H=0.03 T[1], (3) $YBa_2Cu_3O_y$, H=0.5 T[4].

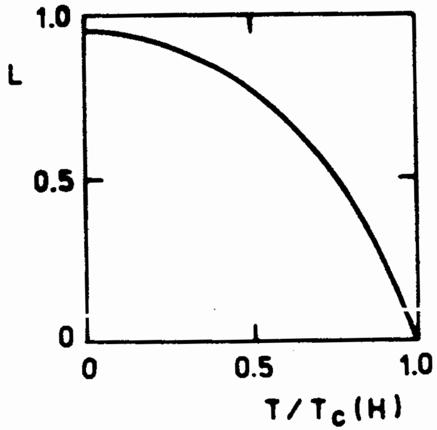

Fig.2. The universal dependence of the nonergodicity parameter L on $T/T_c(H)$ (cf. Fig.1).

3. LOW-FREQUENCY CRITICAL DYNAMICS

Following[13] let us consider in more detail the low-frequency behaviour of relaxation function $D_{ij}(z)$ (15) and, consequently (see (11)-(13)) the long-time behaviour of magnetization M(t) in a critical region, where $|\epsilon| \ll 1$, $\epsilon = (T - T_c)/T_c$. In the single-site approximation for the relaxation kernel $K_{ij}(t)$ (18) one gets

$$K(t) \equiv K_{ii}(t) = f(T, H)D(t) + i\mu . \tag{36}$$

The term $\mu = K''(0)$ corresponds to a regular contribution to the diffusive behaviour of relaxation function at $T \gtrsim T_c(H)$. In the hydrodynamic regime ($q \to 0$, $z \to 0$) the relaxation function is defined by the spin diffusivity $d \sim 1/\mu$ in a manner

$$D(q,z) \sim -(z + iq^2 d)^{-1}. \tag{37}$$

Thus, at $T \gtrsim T_c(H)$ the Josephson spin (superconductive) fluctuations propagate according to a diffusion equation with coefficient d. In this case the motion is ergodic. At $T < T_c(H)$ the situation changes drastically. The correlation function now has to exhibit a zero-frequency pole (see (23)) leading to a singular behaviour of the spectral function

$$D''(q,\omega) \sim \pi\,\delta(\omega)(1 + r_0^2 q^2)^{-1} \tag{38}$$

with a strong reactive character. Here the correlation radius of spatial spin disturbances $r_0 \sim \lim_{\omega\to 0}\lim_{q\to 0}(\omega K''(q,\omega))^{-1/2}$ in the hydrodynamic approximation.

So, below $T_c(H)$ the system is in the nonergodic state. According to (16) and (36), when $|z/\mu| \ll 1$ and $|\epsilon| \ll 1$, the equation for $D(z)$ has the form

$$D(z)(1 - zD(z)) = i\mu + f(T,H)D(z). \tag{39}$$

From (19) and (20) it follows that near $T_c(H)$ $f(T,H) \simeq T_c^2(H)/T^2$, where the form of $T_c(H)$ is given by (29). The solution of eq.(39) is (cf. [13,21])

$$D(z) = \frac{\mu}{|\epsilon|}\,g(\xi), \tag{40}$$

where

$$g(\xi) = \begin{cases} \dfrac{i}{1 + (1 - i\xi)^{1/2}}, & T \ge T_c(H) \\[2ex] -\dfrac{2}{\xi} + \dfrac{i}{1 + (1 - i\xi)^{1/2}}, & T < T_c(H), \end{cases} \tag{41}$$

$$\xi = z/\omega_c, \quad \omega_c = \epsilon^2/\mu. \tag{42}$$

In view of (15) for the time evolution of the critical relaxation function we have

$$D(t) = \begin{cases} |\epsilon|\left(\dfrac{e^{-\tau}}{\sqrt{\pi\tau}} - \operatorname{erfc}(\sqrt{\tau})\right), & T > T_c(H) \\[2ex] (\mu/\pi t)^{1/2}, & T = T_c(H) \\[2ex] L_\epsilon + |\epsilon|\left(\dfrac{e^{-\tau}}{\sqrt{\pi\tau}} - \operatorname{erfc}(\sqrt{\tau})\right), & T < T_c(H), \end{cases} \tag{43}$$

where $\tau = t\omega_c$, $L_\epsilon = -2|\epsilon|\cdot\theta(T_c - T)$ is a critical nonergodicity parameter.

At last, the behaviour of the spectral function $D''(\omega)$ below $T_c(H)$ is governed by the expression

$$D''(\omega) = \pi L_\epsilon\cdot\delta(\omega) + \frac{|\epsilon|}{\omega_c}(1 + \omega^2/\omega_c^2)^{-1/4}. \tag{44}$$

From eqs.(40)-(42) one obtains the well-known in the theory of glasses [13,21] dynamic scaling law

$$\lambda\cdot D(\lambda^2 z, \lambda\omega) = D(z,\omega). \tag{45}$$

Moreover, according to (43) the critical crossover from the experimental (at $T < T_c$ and $T > T_c$) to power law (at $T = T_c$) behaviour is reproduced (cf. [13,21]). The critical frequency ω_c separates the hydrodynamical regime ($|z| \ll \omega_c$) from the critical one ($|z| \gg \omega_c$) which in turn leads to the following asymptotic behaviour of spectral function

$$D''(\omega) \simeq \begin{cases} \dfrac{2\mu}{|\epsilon|}(1 - \omega^2/2\omega_c^2), & \omega \ll \omega_c \\[2ex] \sqrt{\dfrac{\mu}{\omega}}\,(1 - \omega_c^2/2\omega^2), & \omega \gg \omega_c. \end{cases} \tag{46}$$

Due to the dynamic character of the model (1) the scaling behaviour in time (frequency) leads near $T_c(H)$ to the corresponding scaling behaviour in space (wave-vector) (cf. [13]). As a result we have found that when approaching the critical point from the ergodic phase ($T \to T_c^+(H)$), the Josephson spin diffusivity decreased to zero like $d \sim |\epsilon|$, while when approaching it from the nonergodic SCG phase ($T \to T_c^-(H)$), the characteristic length for the spatial spread of superconductive fluctuations diverged like $r_0 \sim |\epsilon|^{-1/2}$. So we may say that the obtained in this Section results confirm in a sense the arguments of Giovannello [22] concerning the universality class of spin glass and superconductive glass models.

In paper [10] the critical low-frequency dynamics for an analogous to (1) model has been considered, but in the region of high fields ($H > H_c^u$) only. In this region (the JSG phase, see (34)) we agree with them.

4. LONG-TIME RELAXATION OF MAGNETIZATION

In this Section we are dealing with a long-time relaxation of remanent magnetization M(t) to its equilibrium value M_∞. The long-time behaviour of some physical quantities in the glass phase is due to the nonergodicity of the latter. Really, as is shown in Sect. 2, below $T_c(H)$ the equilibrium value of magnetization is of the form $M_\infty = \lim_{t\to\infty} M(t) = M_0 L(T,H)$, where the nonergodicity parameter L(T,H) is given by (29). In the model under consideration according to (11)-(14) the time evolution of the magnetization is governed by the time-dependent correlation function of superconductive fluctuations $D_{ij}(t)$ (see (10)). The latter in turn is defined by the system (15)-(18). In general this system of equations permits only numerical investigation. That is why we consider some of its solutions using the well-known empirical relaxation laws as ansatz.

As is known [20], in the glass-like state the exponential relaxation is substituted by the logarithmic one, and further (with increasing observation time) by the Kohlrausch-like law (α-relaxation), which in terms of the correlator D(t) has the form

$$D(t) = D_c \exp(-t^{\beta}) . \tag{47}$$

Considering (47) an ansatz one gets from (15)-(18) a coupled system for parameters D_c and β as

$$D_c^3 - D_c^2 + D_c - (1 - T^2/T_c^2(H)) = 0, \tag{48}$$

$$\beta = - \ln 3 / \ln(D_c^2 f/(f - 1)) , \tag{49}$$

where the function f(T,H) is defined by (19). Comparing (48) and (29) we conclude that $D_c = L$. According to the experimental data the glassy properties of oxide superconductors show the strongest peculiarities near $T_c(H)$. From eqs. (19), (29) and

(49) for the region $T \lesssim T_c(H)$ we have within a proper accuracy

$$f(T, H) \simeq T_c^2(H)/T^2, \quad L(T, H) \simeq (f - 1)/f , \tag{50}$$

$$\beta(T, H) \simeq -\ln 3/\ln L(T, H) . \tag{51}$$

Note that the form of $T_c(H)$ is given by (28) in the general case and by (30)-(35) in some special cases.

For the observation times t which satisfy a requirement $\ln t \lesssim 1$ for $0 < \beta < 1$ (i.e. when $T \lesssim T_c(H)$) from (47) one obtains within a logarithmic accuracy that

$$M(t) = M_0 D(t) \simeq M_0 (L - s \ln t) , \tag{52}$$

where the magnetic viscosity coefficient s has the form

$$s(T, H) \equiv - \frac{1}{M_0} \frac{dM(t)}{d \ln t} = L(T, H) \cdot \beta(T, H). \tag{53}$$

Thus, we have found the relation between Kohlrausch exponent β and commonly used magnetic viscosity coefficient s. Using (51) one gets within the same accuracy.

$$s(T, H) = - L(T, H) \cdot \ln 3/\ln L(T, H) . \tag{54}$$

Fig.3 shows (solid line) the calculated by (54) behaviour of s(T) versus $T/T_c(H)$. The points represent the experimental data [6] on Y-Ba-Cu-O after cooling in the field H = 500 G. As is seen a rather good qualitative agreement in the region of not very low temperatures is obtained (for discussion see Sect. 5).

Let us now consider a power long-time relaxation of the trapped flux which is observed at longer times. We look for the solution of the system (15)-(18) at $t \gg 1$ (see (8)) in the form

$$D(t) = D_\infty + L(T, H)t^{-\alpha} , \quad 0 < \alpha < 1 , \quad D_\infty = L. \tag{55}$$

As a result for the dependence on temperature and applied magnetic field of the exponent α (T,H) in the region $T \lesssim T_c(H)$ one gets

$$\alpha(T, H) = \frac{1 - L^2(T,H)}{3 - L^2(T,H)} , \tag{56}$$

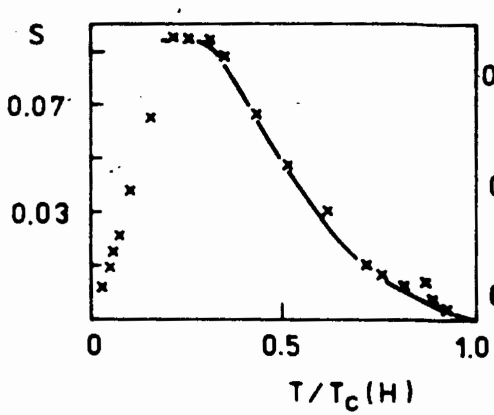

Fig.3. The temperature behaviour of the magnetic viscosity coefficient (54) (solid line) in comparison with the experimental data (points) for $YBa_2Cu_3O_y$ in the field $H = 500$ G[8].

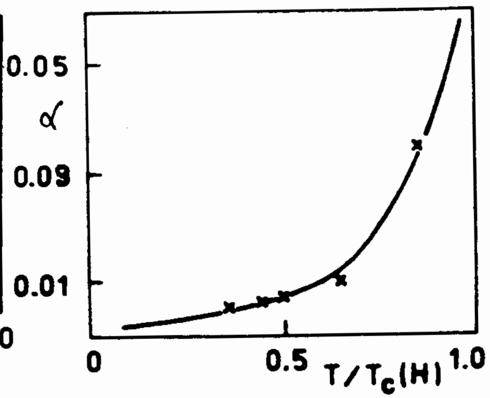

Fig.4. The comparison of the theoretical curve (solid line) for the exponent of power relaxation law (56) with the experimental data (points) for $YBa_2Cu_3O_x$[8].

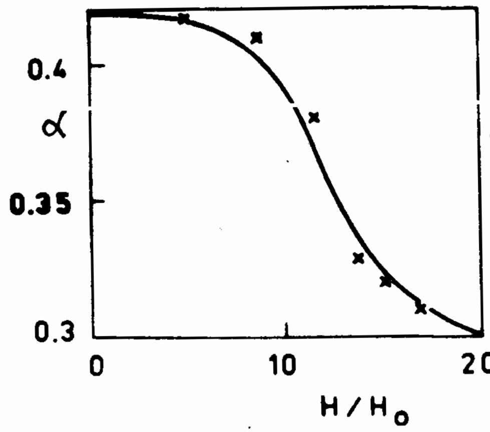

Fig.5. The theoretical curve (solid line) for applied magnetic field dependence of the power relaxation law exponent (56) and experimental data (points) for $La_{2-x}Sr_xCuO_4$ in the region of superconductive glass[5].

Fig. 4 shows the temperature dependence of $\alpha(T)$ at a fixed value of magnetic field calculated by (56) (solid line). The points represent the experimental data[8] for the Y-Ba-Cu-O system.

The behaviour of α versus applied magnetic field H below T_c in comparison with the experimental data[5] for La-Sr-Cu-O in the region of SCG phase (see (33)) is illustrated in Fig.5.

To summarize, we have obtained the following picture for the long-time relaxation of remanent magnetization in the SCG model. The slowest part of the time decay of M(t) is ruled by the logarithmic law (52) which (when the observation time increases) is substituted by the Kohlrausch-like law (47) with exponent $\beta(T,H)$ (see (51)). At longer times of observation the α-relaxation is distorted by a new β-relaxation (power law) mechanism of the magnetization time decay. At last, for extremely achieved experimental times the system decays according to the power law (55) with a nontrivial temperature-field dependence of exponent $\alpha(T,H)$ (see (56)).

5. DISCUSSION

The main result of this paper is as follows. By a relatively simple model (1) one may quite satisfactorily describe nonequilibrium properties of oxide superconductors[1-8,12]. The nonequilibrium behaviour of the system is shown to be in close relation to the nonergodicity parameter (27) which plays the role of the dynamical order parameter for the glass-like state. Its temperature behaviour can be determined from experimental data on FC and ZFC according to (22). The experimental check of the critical crossover in the behaviour of the time-dependent relaxation function (43) near $T_c(H)$ is of interest. It should be emphasized that there exists the temperature-field scaling both for the nonergodicity parameter and for the low-frequency relaxation function (see discussion in the end of Sects. 2 and 3).

Most important in understanding the glass-like properties is the investigation of their long-time relaxation behaviour. The calculated in Sect. 4 parameters of nonexponential decay laws and the discussed in the end of Sect. 4 the hierarchy of relaxation regimes are in qualitative agreement with the known experimental data. Note that the slow dynamics of oxide superconductors is more complicated than in the ordinary spin glasses[18]. For example, in the low-temperature region (below $T_c/2$) there is the well-known Anderson relaxation mechanism[19] due to the flux creep, in which the trapped flux is also known to decay in a logarithmic fashion[17] but with a linear dependence of magnetization rate on temperature (cf.[2,3]). However, when $T \to T_c(H)$ slow relaxation becomes accompanied by a nonlinear temperature-field behaviour of magnetic viscosity[6] in agreement with the predictions of the SCG model (1). Thus, to check the adequacy of the considered model (1) one needs more detailed experimental studies to be carried out on long-time relaxation dynamics of oxide superconductors. The obtained in Sect. 4 system of equations for the correlation function of superconductive fluctuations allows one to get complete information on the time dependnece of remanent magnetization without empirical laws (47) and (55).

The most interesting region on the phase diagram H-T is surely the region of weak magnetic fields ($H < H_c^{\ell}$), where the frustration is still weak and, thus, a more complicated behaviour is expected. Indeed, the recent experiments of Giovannello et al.[23] indicate the existence of fine structure on the phase diagram at $H < H_c^{\ell}$. A more detailed experimental study in this region may provide important information to improve the SCG model.

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Аксенов В.Л., Сергеенков С.А. E17-88-354

Неравновесное поведение и неэргодичность в модели высокотемпературного сверхпроводящего стекла

В рамках двумерной XY-модели фрустрированных джозефсоновских спинов исследованы неравновесные свойства оксидных сверхпроводников. Получена зависимость параметра неэргодичности от температуры и внешнего магнитного поля, показано, что она имеет универсальный характер в переменных $T/T_c(H)$. В приближении взаимодействующих мод изучена критическая низкочастотная динамика модели. Исследовано неэкспоненциальное поведение долговременной релаксации остаточной намагниченности. Получены зависимости показателей временных законов распада от температуры и магнитного поля. Показано, что модель дает качественное описание имеющимся экспериментальным данным для лантановых и иттриевых керамик.

Работа выполнена в Лаборатории нейтронной физики ОИЯИ.

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Aksenov V.L., Sergeenkov S.A. E17-88-354

Nonequilibrium Behaviour and Nonergodicity of High-T_c Superconductive Glass Model

Nonequilibrium properties of oxide superconductors are investigated in the frame of the 2-D XY Josephson glass model. The dependence on temperature and applied magnetic field of the nonergodicity parameter is calculated and shown to have a universal behaviour versus $T/T_c(H)$. In the mode-coupling approximation the critical low-frequency dynamics of the model is studied. The model long-time nonexponential relaxation behaviour of the remanent magnetization is discussed as well. The obtained temperature-field dependences of the magnetization decay rate are in qualitative agreement with the recent experimental data for the lanthanium and yttrium ceramics.

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

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