THE STUDY OF THE LOW-FIELD ELECTRODYNAMICS OF A GRANULAR AND POLYCRYSTALLINE SUPERCONDUCTORS BY METHODS OF LINEARLY AND ROTATING OSCILLATING FIELDS

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It is well known that magnetic fields much less than the lower critical field do not penetrate in superconductors. However, this statement becomes incorrect, if the superconductor is a multiply connected system, e.g., a polycrystal in which crystallites are connected through Josephson junctions. Since very weak magnetic fields can easily penetrate via Josephson junctions, it can be inferred that such fields will penetrate into polycrystalline superconductors too. A similar situation is observed in ceramic HTSC in which granules are connected through weak Josephson junctions. In this case, a wide range of irreversible and nonlinear effects is observed experimentally, which can be explained on the basis of low-field electrodynamics [1,02]. However, to our knowledge, the penetration of very weak magnetic fields into polycrystalline superconductors has not yet been investigated.

It should be recalled that low-field electrodynamics was formulated initially for granular HTSC materials and is based on the following two assumptions.

1. Bean's model of the critical state [3] is applicable to ceramic superconductors and the equation of critical state has the form

$$\left|\frac{dH}{dx}\right| = 4\pi j_c(H) . \tag{1}$$

In Eq. (1), $j_c(H) = \alpha(H)/H$, where $\alpha(H)$ is the pinning force and $j_c(H)$ has the meaning of the critical current density and is a phenomenological function of H.

It should be noted that there exist several models leading to different dependences $j_c(H)$. For instance, our experiments with various YBaCuO ceramic samples [2,4] revealed that

$$j_c(H) = \frac{j_0 H_0}{|h| + H_0}; \qquad j_c(H) = \frac{j_0 H_0^2}{h^2 + H_0^2} \tag{2}$$

Here H_0 is a certain characteristic field, approximately equal to 3 Oe.

The case when $H_0 \to \infty$ leads to Bean's model [3] in which j_c is independent on H. The case $H_0 = 0$ and $j_c H_0 =$ const in the first formula of (2) corresponds to the Kim-Anderson model in which the pinning force $\alpha(H)$ does not depend on H. Generally speaking, there is no theoretically substantiated choice of the function $j_c(H)$.

2. A Josephson medium can exhibit two qualitatively different patterns of magnetic field penetration, whose realization is determined by the dimensionless parameter β

$$\frac{\lambda_{eff}^2}{a^2} = \frac{\Phi_0}{8\pi^2 \mu_{eff} j_c a^3} = \frac{1}{\beta}, \quad \beta = 8\pi^2 \mu_{eff} \frac{j_c a^3}{\Phi_0}, \tag{3}$$

with β characterizing the number of magnetic flux quanta per elementary loop formed by adjacent granules. Here a is the characteristic size of granules, Φ_0 is the magnetic flux quantum and λ_{eff} is effective penetration depth, μ_{eff} is the effective permeability of the ceramic material, taking into account of the fact that the field does not penetrate into granules.

When $\beta \ll 1$, a granular superconductor behaves as a classical type II superconductor in which the field penetrates in the form of vortices, and anisotropy induced by the field becomes an important factor. It was shown in [4] that in another limiting case, when $\beta \gg 1$, the discretness of the Josephson medium, which is now described by equations equivalent to those for systems with self-organized critical state [5], becomes significant.

An approach that can be used to distinguish unambiguously between these two limiting cases with completely different physical patterns corresponding to them was considered in [6].

This approach is directly connected with a very important though sparsely studied problem of longitudinal currents in hard type II superconductors. The essence of the problem can be briefly formulated as follows. It is well known that in the presence of uncut vortices and pinning, the critical current $j_{c||}$, which is longitudinal relative to the magnetic field, is equal to infinity, while the longitudinal electric field $E_{||}$ is always equal to zero. However, experiments proved that $j_{c||}$ and $j_{c\perp}$ ($j_{c\perp}$ is the transverse critical current density) are of the same order of magnitude, and $E_{||}$ differs from zero. In order to explain these phenomena, a model of flux-line cutting was proposed. According to this model, nonparallel external magnetic fields penetrate into a superconductor through mutual cutting of the flux lines formed by these fields, followed by their cross restoration. As a result, finite $j_{c||}$ and $E_{||}$ are formed. In such a case, the local current-voltage characteristic (CVC) connecting the electric field $\vec{E}$ and the current density $\vec{j}$ is strongly anisotropic relative to the magnetic induction vector $\vec{B}$. All what has been said above is also applicable to a Josephson medium for $\beta \ll 1$ ($\lambda_{eff} \gg a$).

On the other hand, if $\beta \gg 1$, the local CVC does not depend on the angle between the current and the magnetic field and is isotropic (see [6]). Thus, if we can establish experimentally whether or not the CVC is isotropic, we can distinguish between the continual ($\beta \ll 1$) and discrete ($\beta \gg 1$) cases.

Let us consider an infinitely large slab in the yz plane, having a thickness d along the x-axis. Further we consider the penetration of a linearly polarized and a rotating alternating current (ac) fields of amplitude h_0 in the presence of a direct current (dc) field H for the anisotropic and isotropic models. In all the cases, we assume that, first, $h_0 \ll H$, and second, the dc field is directed strictly along the z axis, while the ac field $h(t) = h_0 cos\omega t$ lies in the zy plane at an angle γ to H. Thus, for $x = 0$, we have

$$\vec{H}(0,t) = H\vec{e}_z + \vec{h}(t), \qquad \frac{h_0}{H} \ll 1$$

$$\vec{h} = h_z\vec{e}_z + h_y\vec{e}_y,$$

$$h_z = h_0 cos\gamma cos\omega t,$$

$$h_y = h_0 sin\gamma cos(\omega t - \phi) \tag{4}$$

For $\phi = 0$ we have a linearly polarized field directed at an angle γ to the constant field; for $\phi = \pi\ /4$ and $\phi = = \pi/2$, a circularly polarized field with magnitude $h_1 = h_0/\sqrt{2}$ and for

$\gamma = \pi/4$ and arbitrary ϕ, a rotating field with an arbitrary phase shift between h_{0y} and h_{0z}. Thus, (4) includes the externally applied ac field in the most general form.

Now we consider the expressions for the magnetic induction $\vec{h}(t) =< \vec{h}(x,t) >_x$ obtained in [6] for the anisotropic and isotropic models.

1. Rotating and linearly polarized ac magnetic field in the flux-line cutting model (the anisotropic model)

Since the FLC model describes the cases of linearly polarized or rotating ac field in the presence of a dc field in a very similar way, we consider them here in parallel.

$$\vec{h}(t) = h_z(t)\vec{e}_z + h_y(t)\vec{e}_y,$$

$$h_z(t) = \sum_k a_{kz} cos(k\omega t) + b_{kz}\sin(k\omega t),$$

$$h_y(t) = \sum_k a_{ky} cos(k\omega t - \phi) + b_{ky} sin(k\omega t - \phi),$$

$$a_{1z} = \frac{h_0^2 cos^2\gamma}{4\pi j_{c\perp}(H)d}, \qquad a_{1y} = \frac{h_0^2 sin^2\gamma}{4\pi j_{c||}(H)d},$$

$$a_{2k+1,z} = a_{2k+1,y} = 0, \quad k \geq 1$$

$$b_{2k+1,z} = -\frac{h_0^2 cos^2\gamma}{8\pi^2 j_{c\perp}(H)d}\frac{1}{(k^2-1/4)(k+3/2)},$$

$$b_{2k+1,y} = -\frac{h_0^2 sin^2\gamma}{8\pi^2 j_{c||}(H)d}\frac{1}{(k^2-1/4)(k+3/2)},$$

$$a_{2z} = \frac{h_0^3 cos^3\gamma}{32\pi d}\left(\frac{\partial}{\partial H}\frac{1}{j_{c\perp}(H)}\right)$$

$$a_{2k,z} = 0; \qquad k \geq 2,$$

$$a_{2k,y} = 0,$$

$$b_{2k,z} = -\frac{h_0^3 cos^3\gamma}{16\pi^2 d}\frac{k}{(k^2-1/4)(k^2-9/4)}\left(\frac{\partial}{\partial H}\frac{1}{j_{c\perp}(H)}\right), \qquad (5)$$

$$b_{2k,y} = 0$$

It can be seen from these equations that the z and y components oscillate independently with their j_c, and odd harmonics are proportional to $h_0^2 cos^2\gamma$ and $h_0^2 sin^2\gamma$, while even harmonics differ from zero only for the z component and are proportional to $h_0^3 cos^3\gamma$.

2. Rotating or linearly polarized ac magnetic field in the isotropic model

In the case of a rotating field the spectrum of garmonics includes the first garmonic and the second one only, all others being zero.

For a linearly polarized field the expressions for the z and y components are given by:

$$a_{1z} = \frac{h_0^2}{4\pi j_c(H)d} \cos\gamma, \qquad a_{1y} = \frac{h_0^2}{4\pi j_c(H)d} \sin\gamma,$$

$$a_{2k+1,z} = a_{2k+1,y} = 0; \qquad k \geq 1,$$

$$b_{2k+1,z} = -\frac{h_0^2 cos\gamma}{8\pi^2 j_c(H)d} \frac{1}{(k^2-1/4)(k+3/2)},$$

$$b_{2k+1,y} = -\frac{h_0^2 sin\gamma}{8\pi^2 j_c(H)d} \frac{1}{(k^2-1/4)(k+3/2)},$$

$$a_{2z} = \frac{h_0^3 cos^2\gamma}{32\pi d} \left(\frac{\partial}{\partial H} \frac{1}{j_c(H)} \right),$$

$$a_{2y} = \frac{h_0^3 \sin\gamma\cos\gamma}{32\pi d} \left(\frac{\partial}{\partial H} \frac{1}{j_c(H)} \right),$$

$$a_{2k,z} = a_{2k,y} = 0; \qquad k \geq 2,$$

$$b_{2k,z} = -\frac{h_0^3 \cos^2\gamma}{16\pi^2 d} \frac{k}{(k^2-1/4)(k^2-9/4)} \left(\frac{\partial}{\partial H} \frac{1}{j_c(H)} \right),$$

$$b_{2k,y} = -\frac{h_0^3 \sin\gamma\cos\gamma}{16\pi^2 d} \frac{k}{(k^2-1/4)(k^2-9/4)} \left(\frac{\partial}{\partial H} \frac{1}{j_c(H)} \right). \tag{6}$$

Thus, it can be seen that the isotropic model strongly differs from the flux-line cutting model, first, by the form of the angular dependence of harmonics, and second, by that the even harmonics $a_{2k,y}, b_{2k,y}$, which are equal to zero in the flux-line cutting model, now differ from zero.

3. Experiment

The experiments [7] were made on polycrystalline samples of $SnMo_6S_8$ and PbM_6S_8. This choice was dictated by a very large value of H_{c2}, and accordingly, by very small coherence length $\xi = 23$ A for molybdenum chalcogenides. This makes these materials similar to HTSC for which ξ is of the order of a few Angstroms. Hence, we can expect that any (even small) defect will play the role of a Josephson junction for such values of ξ.

The linear and nonlinear responses were measured in the temperature range from 4.2 K to $T > T_c$ at frequencies varying from 10^3 to 10^5 Hz in the field interval $10^{-2} \leq h_0 \leq 1$ Oe and for $H \leq 20$ Oe. We studied the field and temperature dependences of the real and imaginary components of linear susceptibility (χ' and χ'') as well as the moduli of the higher harmonic amplitudes $c_n = (a_n^2 + b_n^2)^{1/2}$ (n = 2, 3, ...).

Consider first the results of experiments when the ac and dc fields are collinear ($\gamma = 0$). Then, as seen from (5) and (6), we obtained for both models at $h_0 \ll H$ one and the same result for the harmonic magnitudes.

The temperature dependences of the susceptibility $\chi_3 = c_3/h_0$ for various values of the amplitude of the ac field h_0 are identical to the corresponding dependences for HTSC ceramics (see, e.g., [2]). A typical feature of these curves is the presence of two peaks. The high-temperature peak can be associated with the penetration of the field into crystallites, while the low-frequency peak is associated with the penetration only into the Josephson medium formed by weak links between the crystallites. The position of the low-temperature peak is displaced towards lower temperatures depending on the increase of h_0, which is associated with a transition from the mode in which the field penetrates to the middle of the sample to a mode in which the field penetrates only to a small depth of the sample. In the latter case, the condition of a weak field ($h_0 \ll H$) can be easily realized. Henceforth, we consider the experimental results only for this limiting case.

The dependences of a_1, b_1, c_3, and c_2 (for $H \neq 0$) on h_0 are also in accordance with the results of the theory: odd harmonics exhibit a quadratic dependence on h_0, while the second harmonic exhibits a cubic dependence. Note that the scale of magnetic induction induced by the ac field (from 0.01 mOe to 1 Oe) varies from 1 μG to 0.1 G.

Analyzing the behavior of even harmonics proportional to $\frac{\partial}{\partial H}(1/j_c(H))$ and odd harmonics proportional to $1/j_c(H)$ in a dc field for $h_0 =$ const and using expressions (5) (or (6)), we can determine self-consistently the form of the function $j_c(H)$. This is how the explicit form of this function was determined for granular HTSC [see (4)]. However, the situation turned out to be more complicated for the dependence of higher harmonics on the dc magnetic field in $PbMo_6S_8$ (Fig. 1). It can be seen from the figure that starting from $H = 1$ Oe, c_3 is practically independent on H. Proceeding from formulas (5) or (6) (at $\gamma = 0$), we should expect the absence of the second harmonic (as in the case of $SnMo_6S_8$). However, in the experiments with $PbMo_6S_8$ we observed the second harmonic whose magnitude was considerably larger than the value predicted by (5) or (6). Using these relations, we can write the following expression for the ratio c_2/c_3:

$$c_2/c_3 \approx 4h_0 j_c(H)(\frac{\partial}{\partial H}\frac{1}{j_c(H)}), \tag{7}$$

which does not depend on the specific form of function $j_c(H)$, i.e., is independent of the model. If we calculate the derivative in (7) proceeding from the experimental data on the third harmonic ($c_3 \propto 1/j_c(H)$), it turns out that the ratio c_2/c_3 must be two or three orders of magnitude smaller than the observed value equal approximately to 0.1. In this case, our sensitivity (not worse than 0.1 μV) is insufficient for observing the second harmonic (precisely this situation was observed for $SnMo_6S_8$). However, in spite of such a discrepancy, all the remaining (field and angular) dependences are in excellent agreement with the theory.

The next series of experiments was carried out with a linearly polarized ac field forming an angle γ with the dc field. These experiments were aimed at determinition of what version (the flux-line cutting model or the isotropic model) is realized in the polycrystalline superconductors under investigation. It can be seen from formulas (5) and (6) that in the former case $c_{3z} \propto cos^2\gamma$, $c_{3y} \propto sin^2\gamma$ and $c_{2z} \propto cos^3\gamma$, $c_{2y} = 0$, while in the latter case $c_{3z} \propto cos\gamma$, $c_{3y} \propto sin\gamma$ and $c_{2z} \propto cos^2\gamma$ $c_{2y} \propto sin\gamma cos\gamma$.

Fig. 2 shows that the angular dependences for harmonics obtained by us unambiguously correspond to the isotropic model. Besides, the fact that the second harmonic c_{2y} differs from zero for an arbitrary γ (except $\gamma = 0$ and $\gamma = \pi/2$) also speaks in favor of this model.

The same angular dependencies of a higher harmonics were obtained by us when studying a HTSC-ceramics [6]. Besides, in these investigations the method of rotating ac field in the presence of a dc field was used. Firstly, it was found that the third harmonic was not observed

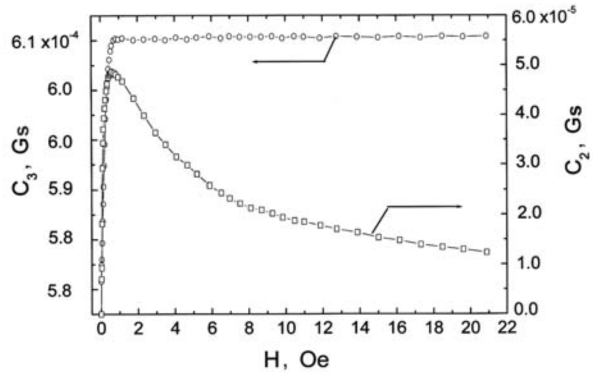

Figure 1: Dependence of c_3 and c_2 on H for $h_0 = 0.4$ Oe for $PbMo_6S_8$ at $T = 10.5$ K

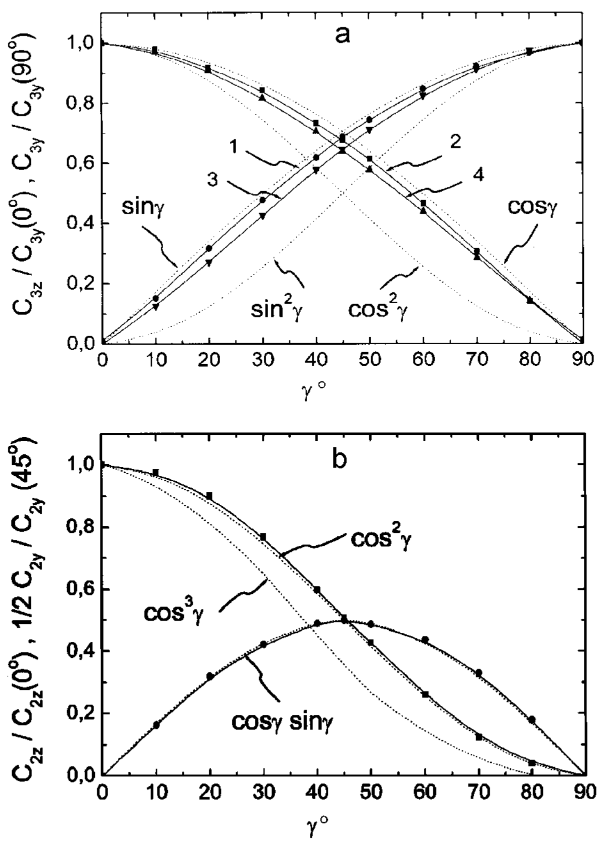

Figure 2: (a) Angular dependences of c_{3z} and c_{3y} normalized to their values for $\gamma = 0$ and $\gamma = 90$, respectively. Curves 1 and 2 correspond to $PbMo_6S_8$, $H = 8$ Oe, $h_0 = 0.6$ Oe, and $T = 10.5$ K, while curves 3 and 4 correspond to $SnMo_6S_8$, $H = 20$ Oe, $h_0 = 0.6$ Oe, and $T = 11.5$ K. (b) Angular dependences $c_{2z}/c_{2z}(0)$ and $(1/2)c_{2y}/c_{2y}(45)$ for $PhMo_6S_8$

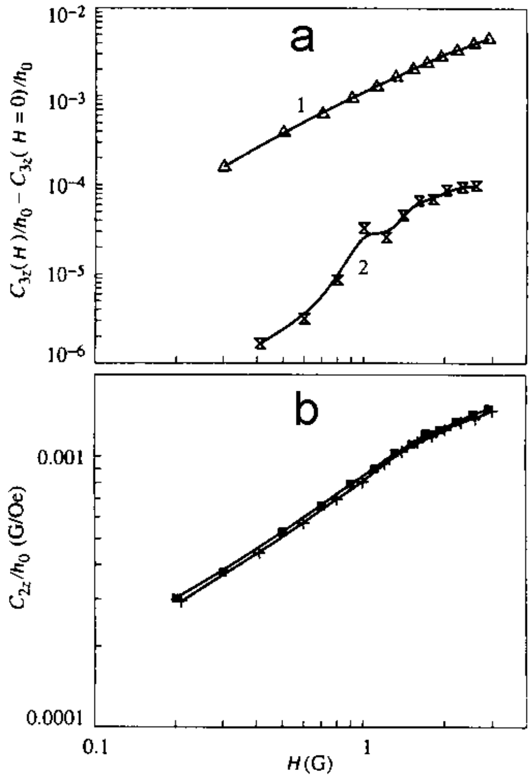

Figure 3: (a) Difference $c_{3z}(H)/h_0 - c_{3z}(0)/h_0$ as a function of H for linearly polarized (curve 1) and rotating (curve 2) fields. (b) c_z/h_0 as a function of H for the rotating and linearly polarized oscillating fields

in these experiments. For clarity Fig. 3(a) shows $c_{3z,y}$ as a function of H for both the circularly (curve 1) and linearly polarized (curve 2) fields. It can be seen that in the former case the third harmonic is practically zero, whereas in the latter $c_{3z,y}$ grows with increasing dc field, which is possibly due to a reduction on j_c.

The behaviour of the second harmonic with increasing H was one and the same for linearly polarized and rotating fields (Fig. 3(b)). We can see that in a ceramic HTSC the behaviour of higher harmonics in the rotating field also agrees well with the isotropic model.

Thus, the results of our experiments show that a polycrystalline superconductor, as well as an HTSC ceramics material, behaves as a standard Josephson medium in weak fields. Such systems display irreversible and nonlinear effects typical of low-field electrodynamics in Josephson media.

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