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Critical Behaviour and Anisotropy of High Temperature Superconductors

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ABSTRACT: The weak pinning and relatively large flux creep in extended regions of the mixed state and even at low temperatures are intimately related to the vortex lattice anisotropy. This vortex lattice anisotropy itself is a result of the layered structure of the oxidic cuprates and is very large for the Bi and Tl compounds. In this contribution we discuss some of the characteristics of the field-temperature phase diagram as appropriate for $Bi_2Sr_2CaCu_2O_8$. Collective flux pinning and flux creep in case oxygen vacancies are the dominant pinning centers, are briefly reviewed. Finally, magnetization measurements versus temperature or field in the reversible vortex liquid regime about T_c are presented. They show clear evidence for strong two-dimensional diamagnetic fluctuations and behaviour deviating from the well-known mean-field (Abrikosov) characteristics.

1. INTRODUCTION

The Tl and Bi-compound high temperature superconductors (HTS) have extremely anisotropic properties. For instance, in the normal state the c-axis resistivity is larger by a factor of 10^5 than the resistivity in the ab plane [1]. In the superconducting state the anisotropy is expressed by the large value of the anisotropy parameter Γ as determined by the angular dependence of the magnetic torque [2]. For large Γ, or equivalently for large effective mass ratios m_c/m_{ab}, the vortex lattice (VL) has properties quite different then the conventional Abrikosov VL. The reason is that superconductivity is concentrated in the (super)conducting planes. The coupling between these planes is via Josephson tunneling [3]. The conducting planes consist of double or triple CuO_2 layers (with Ca layers in between) seperated by non-conducting BiO, SrO, or TlO and BaO layers. It is appealing to contribute the much smaller anisotropy of $YBa_2Cu_3O_7$ (YBCO) to the conducting CuO chain layers. The former two compounds therefore actually form a stacking of superconducting-insulating layers (SISI) with Josephson coupling, whereas the latter compound rather is a stacking of superconducting-normal layers (SNSN) with proximity coupling.

In the following we consider the situation in which the field is applied along the c-direction. The superconducting screening currents around the vortex cores are

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restricted to the S layers. The vortex lines thus break up into stacked two-dimensional (2D) pancake vortices [4]. The short range coupling between the pancake vortices in adjacent S layers is weak. This leads to a reduction of the non-local tilt modulus by a factor Γ^{-1} as was predicted by Houghton et al. [5] and verified by experiments on the conventional layered superconductor H-$NbSe_2$ by Koorevaar et al. [6]. For the HTS the parameter Γ can be as large as 3×10^3 for Bi:2212 and 10^5 for Tl:2212. The interaction between the pancake vortices in adjacent S layers drops exponentially at a distance equal to the Josephson length $R_J = \Gamma^{1/2}s$ along the layers [7], where s is the periodicity of the SI multilayer, i.e. s = 1.5 nm for Bi and Tl:2212. Consequently, positional fluctuations of the pancake VL within a length scale R_J are effectively decoupled which leads to quasi two-dimensional (2D) behavior. When $a_o \simeq (\Phi_o/B)^{1/2} < R_J$, i.e. $B > B_{2D}$ with $B_{2D} \equiv \Phi_o/\Gamma s^2$, the VL attains 2D properties comparable to a description valid for a thin film superconductor of thickness s [7,8,9]. We compute $B_{2D} \simeq 0.3$ T for Bi:2212 and 10 mT for Tl:2212, while it is 50 T for YBCO.

The large anisotropy is responsible for an unconventional (B,T) phase-diagram. The main features in case flux pinning is important, are given in Fig. 1. At low B and T pinning plays a dominant role, it destroys the long range order in the VL and yields the 3D vortex-glass state. At low T and $B > B_{2D}$ the short-range disorder of the pancake vortex lattice is of 2D nature, i.e. determined only by the intralayer pinning and vortex interactions. Thermal fluctuations become dominant at high temperature and give rise to both thermal depinning [10] and VL melting [11, 12]. It can be shown [13, 7] that thermal depinning is preceded by melting in case of a 2D VL with weak pinning. The melting line is then given by the Kosterlitz-Thouless temperature T_M^{2D} for unbinding of vortex-dislocation pairs. The value of T_M^{2D} for Bi:2212 is estimated to be about 30 K. Melting of the 3D VL in presence of pinning has not yet been worked out in detail [9], but it is believed that it will closely follow the predictions of the Lindemann criterion $\langle u^2\rangle_T \approx (a_o/6)^2$ without pinning [11,5,12], where $\langle u^2\rangle_T$ is the squared averaged thermal displacement of the VL. In the 3D regime thermal depinning and melting are also closely related. Raising the temperature leads to a transition to a vortex liquid state which is expected to show reversible magnetic properties away from the transition line. This vortex liquid regime as sketched in Fig. 1, is very extended in Bi:2212.

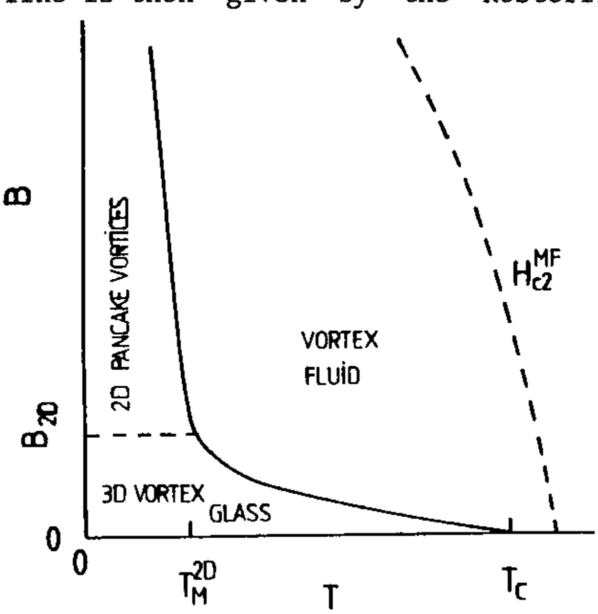

Fig. 1 Phase diagram of a layered quasi two dimensional superconductor [9].

We will now briefly address the consequences for pinning, creep and reversible magnetic properties.

2. FLUX PINNING IN EXTREMELY ANISOTROPIC VORTEX LATTICES

In the 2D regime estimates for the pinning force density F_p and the critical current density $j_c = F_p/B$ are relatively easy to give, since only the disorder of the VL along the S layers and the shear modulus are important [7]. The concept of collective pinning has recently been worked out assuming that the predominant pinning centers were oxygen vacancies [14]. The main results computed for Bi:2212 can be summarized as follows:

1. The maximum elementary pinning force at $T = 0$ below about 10 T is $f_p(0) \simeq 6.5\times10^{-14}$N.
2. The collective average effect of vacancies in the vortex core can be computed assuming an areal density n of 3.5×10^{17} vacancies per double layer per m^2. The net pinning force on each vortex is $F_v \simeq 1.7\times10^{-13}$N and the resulting macroscopic pinning force is given by $F_p = F_v/(a_o^2 s)$ which yields $j_c = F_v/\Phi_o s = 5\times10^{10}$ Am^{-2}.
3. The pinning energy of each vortex U_v is given by $U_v = F_v\ \eta(0)$ with η the Ginzburg-Landau coherence length. We obtaub $U_v = 34$ K. It is very important to realize that the quasi-2D nature of the VL causes both a large j_c and a small U_v. The reason is very simple: $j_c \propto (a_o^2 s)^{-1/2}$, whereas $U_v \propto (a_o^2 s)^{1/2}$. The volume $a_o^2 s$ arises here, because at low fields the shear modulus is so small that each vortex is pinned independently. It is precisely the smallness of $a_o^2 s$ which leads to the above paradoxial result.

3. FLUX CREEP IN EXTREMELY ANISOTROPIC VORTEX LATTICES

The current densities as obtained from magnetization measurements on Bi:2212 single crystals show a low-temperature value of about 2×10^9 Am^{-2} which decreases strongly between 15 and 30 K depending on field (4 T and 0.2 T, respectively) [15]. Even at 1.5 K a clear magnetic decay with time is observed which is related to thermally activated flux creep. That creep effects are so evidently present is not surprising with a pin energy as small as 34 K. Therefore, magnetic experiments do not probe j_c but rather a much smaller current density j that depends on the ramp speed of the magnetic field during the measurements.

Creep processes can be studied with a variety of experimental techniques. One of these techniques, ac susceptibility measurements in very small driving fields, provided data which could be well explained in terms of thermally activated creep of VL defects, especially interstitials [14]. Together with the results of the preceding section the defect-mediated creep picture yields a self consistent description of collective pinning and creep. However, it is not the only possible explanation and more decisive experiments are certainly needed to create a better understanding of the irreversible and metastable VL behavior in the HTS.

4. REVERSIBLE MAGNETIC BEHAVIOR ABOUT T_c

We now turn to the extraordinary results of very recent magnetization measurements on Bi:2212 single crystals. The characteristics of the single crystal are as follows [17]. The shape is almost of a square platelet of 4x4x0.11 mm with the c axis along the smallest dimension. The nominal composition as determined from EPMA is $Bi_{2.2}Sr_{1.9}CaCu_2O_x$. HREM only revealed minor traces of the 2201 phase in a concentration of 2 in 1000 CuO_2 double layers. The transition temperature has been determined from ac-susceptibility measurements in a μ-metal shielded environment with a background field less than 0.3 μT and an applied ac field of 2.8 μT parallel to the crystallographic ab planes. A linear extrapolation to $\chi' = 0$ gave$T_c(0) = 88.1$ K and a transition width of 1.5 K between $\chi' = 0$ and $\chi' = -1$. The above observations give us some confidence that we are dealing with a high-quality crystal.

If M is measured versus T in various constant fields, it is expected according to the G-L mean field (Abrikosov) theory that M decreases linearly at $T_c(B)$. Here $T_c(B)$ is defined as the temperature for which $B = B_{c2}(T)$. Neglecting the small temperature dependence of κ, the slope dM/dT would be given by $(2.32\kappa^2)^{-1}dH_{c2}/dT$. Close to $T_c(0)$ this would be a constant, but at lower temperatures dM/dT should slightly decrease for two reasons. In the first place $H_{c2}(T)$-H increases, so that M(H) begins to deviate from the linear (H_{c2}-H) behavior. And in the second place the $H_{c2}(T)$ curve starts to deviate from the linear dependence near $T_c(0)$ [16]. It is further expected that $T_c(B)$ determined by extrapolating the linear M(T) behaviorback to M = 0, should decrease with increasing field. Such behavior has been observed on $YBa_2Cu_3O_7$ single crystals [18]. We checked our equipment on a grain-aligned YBCO sample and found excellent agreement with Ref. 18. However, carrying out the same experiment on a Bi:2212 single crystal gave the results plotted in Fig. 2 for various fields denoted in the caption. It clearly shows that the conventional construction gives rise to $T_c(B)$ increasing with field, as displayed in the inset of Fig. 2. The slopes of the linear M(T) curves are plotted versus field in Fig. 3 on a semi-log scale. A clear kink in the data is seen at about 1 T. It has been suggested by Kogan et al. [19] to use the B dependence of dM/dT to determine the penetration depth $\lambda(0)$. In the regime $H_{c1} \ll H \ll H_{c2}$ one can show that according to the G-L theory for high κ superconductors

$$\frac{dM}{dT} = \frac{\Phi_o}{8\pi\mu_o\lambda^2(0)T_c}\left\{\ln\left(\frac{\beta\, H_{c2}}{H}\right) + 1\right\} \tag{1}$$

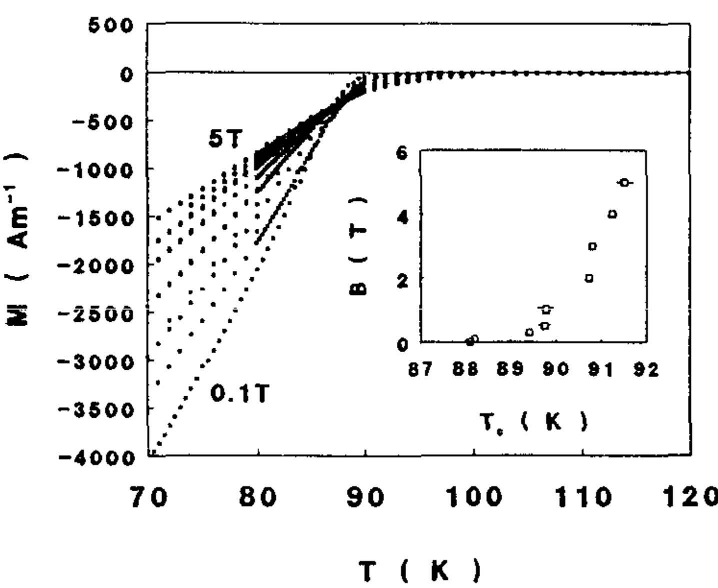

Fig. 2 Compilation of magnetization data as a function of temperature of a Bi:2212 single crystal in fields of 0.1, 0.3, 0.5, 1, 2, 3, 4, and 5 T (bottom to top) measured in both increasing and decreasing T runs. The inset shows $T_c(B)$ as determined from a linear extrapolation of M(T) back to M = 0.

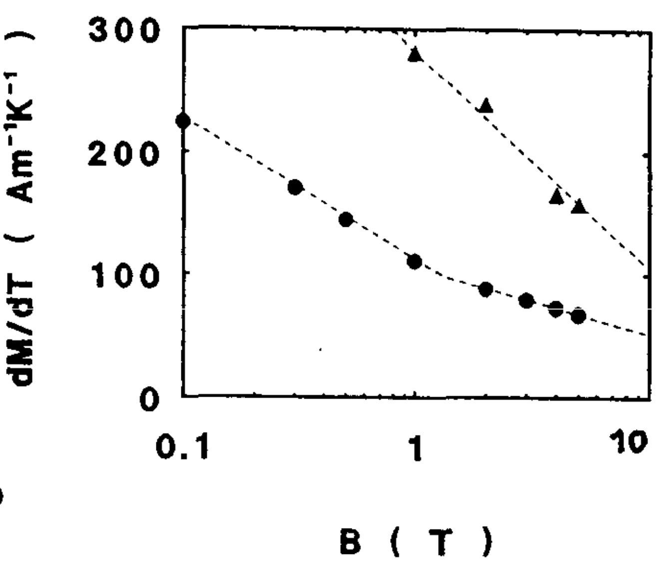

Fig. 3 The slopes of the linear parts of the M(T) curves in Fig. 2 as a function of field for Bi:2212 (circles). For comparison the data of $YBa_2Cu_3O_7$ [18] are shown as well (triangles). Note the kink at about 1 T, possibly indicating a dimensional crossover from 3D to 2D behavior. The lines are guide to the eye.

where β is a constant or order unity. From the slopes in Fig. 3 we obtain $\lambda(0)$ = 120 nm for B < 1 T and $\lambda(0)$ = 180 nm for B > 1 T. In the clean limit the low-temperature London penetration depth is computed from $\lambda_L(0) = \sqrt{2}\lambda(0)$ which yields $\lambda_L(0)$ = 170 nm and 260 nm for B < 1 T and > 1 T, respectively. Repeating this with the data for YBaCuO [18] depicted by the triangles, gives $\lambda(0)$ = 102 ±6 nm and $\lambda_L(0)$ = 141 ±0 nm in good agreement with the values obtained with a variety of techniques [1]. At this point we may conclude that this method seems to work for YBaCuO, but because the idea of two penetration depths for different field regimes doesn't seem to be physically sound, this approach doesn't work in case of Bi:2212.

In addition, the unphysical increase of T_c with field seems to indicate as well that the conventional mean-field theory doesn't describe the magnetic behavior close to $T_c(B)$. Since this is expected to be the flux-liquid regime of Fig. 1, we propose that the Abrikosov theory doesn't apply for a flux-liquid and that the mean-field transition line $H_{c2}^{MF}(T)$ just indicates a crossover from a normal state with superconducting fluctuations above the line to a "superconducting" state in which fluctuations play a dominant role as well. The question arises whether the name superconducting is still correct for this phase, because the resistance is non-zero. The superconductivity only manifests itself by virtue of the diamagnetic signal. Such a destruction of long-range superconducting order in quasi-2D superconductors has been recently proposed by De Jongh [20].
Very recently, Hao and Clem [21] extended their previous work [16] and studied the field dependence of the magnetization below the linar regime near H_{c2}. They argue correctly that because of the large increase of H_{c2} with decreasing temperature, Abrikosov's high-field result is valid only for temperatures very close to T_c and is usually not applicable due to fluctuation effects. They also obtain an analytical approximation for M(H) in the intermediate field regime which in their units is

$$-4\pi M = \frac{\alpha}{4\kappa} \ln\left(\frac{\beta\kappa}{H}\right) \qquad (2)$$

The constants α and β are determined by fitting this expression to the computed M(H) curves. Their most accurate result is $\alpha = 0.84$, $\beta = 1.08$ for $0.02 N h < 0.1$ and $\alpha = 0.70$, $\beta = 1.74$ for $0.1 < h < 0.3$, where h is H/H_{c2}. We used this result to compute M(T) at constant fields $\mu_0 H = 0.1$ T, 1 T, and 5 T assuming $\mu_0 H_{c2} = S(T_c-T)$ with S = 2 T/K. It is seen from Fig. 4 that a linear interpretation of the resulting M(T) points still yields a T_c which decreases when the field is enlarged, in contradiction to our measurements. It therefore seems that the theory of Ref. 21cannot fully explain the anomalous magnetic behaviour of our sample indicating that fluctuations lead to deviations from the mean-field theory.

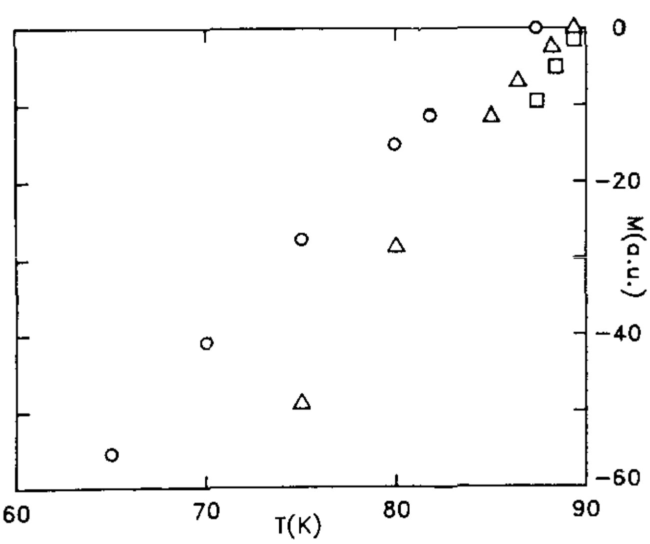

Fig. 4 M vs T computed using Eq.(2) and the α,β values of Ref. [21] for fields of 0.1 T (□), 1 T (Δ), and 5 T (O) assuming $-d\mu_0 H_{c2}/dT = 2$ T/K.

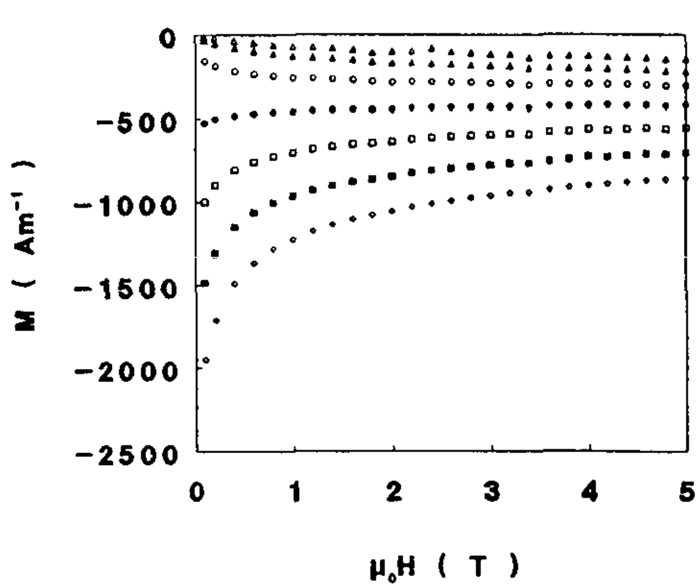

Fig. 5 Compilation of M vs H data above 0.1 T at 80 K (◆), 82 K (■), 84 K (□), 86 K (●), 88 K (O), 90 K (▲) and 92 K (Δ). Note the diamagnetic behavior in fields up to 5 T and temperatures up to 92 K.

Another way of studying the magnetic behavior is to measure M as a function of H for several constant temperatures. The results are displayed in Fig. 5 for T ranging between 80 and 92 K and B > 0.1 T. Low-field results are less reliable as far as field is concerned (trapped flux in magnet), but can be described, e.g. for 80 K, as follows: increasing the field from zero a sharp decrease of M to a value -7200 Am^{-1} is followed by a steep increase to -1950 Am^{-1} at 0.1 T, Fig. 5. This behavior resembles that of a high-κ superconductor with very small H_{c1} which is even more reduced by the demagnetization effect in the H||c configuration. The value $-\mu_0 M = 9.0$ mT may thus be close to $\mu_0 H_{c1}$ at 80 K. Assuming $H_{c1}(T) = H_{c1}(0)(1-t^2)$ we obtain $\mu_0 H_{c1}(0) = 51$ mT which is a factor of 2.5 larger than reported by Biggs et al. [22]. The high-field measurements at 80 K show a gradual increase of M with field which would give $\mu_0 H_{c2} = 26$ T when linearly extrapolated to M = 0. Thus, assuming the Abrikosov theory to be valid, we estimate $-\mu_0 dH_{c2}/dT$ to be about 3.3 T/K and $\xi(0) \approx 1.1$ nm. The latter would be much smaller than reported till now [1] and would support a local pairing mechanism for the superconductivity in these materials.

Going up in temperature, the data in Fig. 5 increasingly deviate from the Abrikosov prediction, e.g. the data at 86 K are almost independent of H, but M is still negative. Considerable negative values of M are still observed above 92 K. It turns out that these results can be well explained by the theory for diamagnetic fluctuations in layered superconductors [23]. For a detailed discussion we refer to a forthcoming publication [17]. The nature of the fluctuations above 0.2 T is strongly 2D. A mean-field transition temperature of 92.2 ±0.5 K is obtained. The regime of critical (or strong) fluctuations is about 2 K around a GL transition temperature of 88.0 ±0.4 K. It is also concluded that a theoretical description of the fluctuation regime below the mean-field transition line has still to be formulated. The kink in Fig. 3 points to a change in dimensionality from 3D to 2D at about 1 T.

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