Local Hall-probe-based susceptometry of $Tl_2Ba_2CaCu_2O_8$ epitaxial films: Critical state and flux dynamics in collinear ac and dc magnetic fields

Leena K. Sahoo, S. Patnaik, and R. C. Budhani
Department of Physics, Indian Institute of Technology Kanpur, Kanpur-208016, India

W. L. Holstein
DuPont Central Research and Development, P.O. Box 80262, Wilmington, Delaware 19880-0262
(Received 24 October 2000; published 27 April 2001)

The ac screening response of a superconducting epitaxial $Tl_2Ba_2CaCu_2O_8$ film in static and time-varying magnetic fields collinear with the film normal is reported. The complex susceptibility components measured in terms of the fundamental- ($T_{H'}$ and $T_{H''}$) and third-harmonic ($|T_{H3}|$) transmittivities of a local Hall-probe susceptometer are compared with the models of flux penetration in a disc-shaped thin superconductor. In the absence of a dc field, the low-amplitude ($h_0 \leqslant 13.2$ G) and low-frequency ($f \leqslant 1073$ Hz) response is well described by the Bean critical-state model applied to thin discs in a transverse ac field. The J_c calculated from the model compares well with the J_c measured in a four-probe transport geometry. When a dc field is superimposed over the ac field, signatures of a creeplike motion of vortices are seen in a temperature regime where $|T_{H3}|$ is nonzero. Vortex-glass/collective-pinning theories for a three-dimensional vortex system are considered to understand this behavior in a high-temperature ($T \geqslant 60$ K) low-field ($\mu_0 H \leqslant 0.5$ T) regime. The irreversibility field derived from the onset temperature of $|T_{H3}|$ varies as $H(T) \sim (1-t)^n$ on the H-T plane, with $n=1.2$. At low fields, this line also represents the variation of the vortex-glass temperature T_g with field inferred from the tail of the resistive transition.

DOI: 10.1103/PhysRevB.63.214501 PACS number(s): 74.72.Fq, 74.76.−w

I. INTRODUCTION

The theory of magnetic-flux penetration and ac losses in disc- and slab-shaped thin superconducting films subjected to perpendicular ac and dc magnetic fields has attracted much attention in recent years[1–8] It has been argued that for low-amplitude quasistatic fields, the response of a disc-shaped superconductor can be described under the framework of Bean's critical-state (CS) model.[9] Accordingly, in a transverse-field geometry the flux front penetrates the disc radially inwards up to a distance over which the local current density J is equal to the critical current density J_c. However, unlike the ideal case of CS in an infinite cylinder subjected to an axial field, the current density in the flux-free central region of the disc is not zero, but decreases smoothly from $J = J_c$ at the edge of the flux front to $J=0$ at the center of the disc. While the static nonlinear response as embodied in the Bean model is followed in the strong-pinning regime of a superconductor subjected to low-frequency low-amplitude ac field, deviations from this ideal behavior are expected at the higher values of temperature, frequency, and field strength. These deviations from the Bean model occur due to thermally activated flux-creep and flux-flow processes, which are quite dominant in high-T_c cuprates. While there are a large number of experiments that have established the existence of thermally activated processes,[10–13] and the general theories of diffusive and nonlinear vortex motion are available in detail,[14–17] the specific case of a high-T_c thin-film superconductor subjected to a transverse ac field is being tackled only recently.[6–8]

Unlike the case of $YBa_2Cu_3O_7$,[6,7] where the low-frequency ac response is in agreement with the predictions of the modified Bean model,[1–3] $Tl_2Ba_2CaCu_2O_8$ (Tl-2212) and $Bi_2Si_2CaCu_2O_8$ (Bi-2212) are expected to show the effects of thermally activated flux creep even at small driving forces due to their two-dimensional character. In fact, the vortex phase diagram of a system like Tl-2212 is rather complicated due to many competing interactions of nearly the same strength. The dynamical response of vortices in such a system will have elements of self organization and criticality, hysteretic and history effects, and plastic/elastic motion.[18,19] Measuring these effects needs probes of varying strengths and time scales. Measurements of ac susceptibility and its higher harmonics provide a powerful method to study different regimes of vortex dynamics. For example, the third harmonic susceptibility is very sensitive to the onset of a nonlinear response. Similarly, by changing the strength of ac perturbation, a transition from elastic to plastic motion can be triggered in the vortex phase. However, the studies of vortex dynamics in the Tl-2212 system have been limited primarily to standard dc transport and magnetization measurement.[20–24] In this paper we report measurements of vortex dynamics in Tl-2212 films using a local Hall-probe ac susceptometer. We first apply transverse-field theories of ac-field penetration to square-shaped thin-film samples and establish the applicability of the static Bean model. Detailed measurements of low-field high-temperature vortex dynamics and the applicability of vortex-glass/collective-creep models for dissipation are discussed subsequently.

II. SCREENING RESPONSE OF A THIN-FILM SUPERCONDUCTOR IN TRANSVERSE-FIELD GEOMETRY

The applicability of the Bean critical-state model to a thin circular disc of a superconductor subjected to a low-frequency oscillating magnetic field perpendicular to its plane was examined first by Mikheenko and Kuzovlev[1] and subsequently by Zhu *et al.*[2] This work has been extended by Clem and Sanchez[3] to calculate hysteretic-magnetization curves and ac-susceptibility components (χ_n' and χ_n''). Brandt and coworkers, in a series of publications,[4–6] have discussed the applicability of the static (without creep) and dynamic (with creep and flow) critical state to thin superconductors of rectangular, strip, and ring geometries. Following Clem and Sanchez,[3] the components of the complex susceptibility under the static-CS case for a disc of radius R and thickness d placed in a transverse ac field of the type $h=h_0\exp(-i\omega t)$ are given as

$$\chi_n'=\frac{2\chi_0}{\pi}\int_0^\pi(1-\cos\theta)S[(x/2)(1-\cos\theta)]\cos n\theta\,d\theta \tag{1}$$

and

$$\chi_n''=\frac{2\chi_0}{\pi}\int_0^\pi\{-S(x)+(1-\cos\theta)S\times[(x/2)(l-\cos\theta)]\}\sin n\theta\,d\theta \tag{2}$$

where $x=h_0/H_d$, $H_d=J_cd/2$, and $\chi_0=8R/3\pi d$, and

$$S(x)=\frac{1}{2x}\left(\cos^{-1}\left[\frac{1}{\cosh x}\right]+\frac{\sinh x}{\cosh^2 x}\right).$$

In the limit of very small ac-field amplitudes ($x\ll 1$), the components of the fundamental susceptibility are

$$\chi'=-\chi_0(1-\tfrac{15}{32}x^2), \tag{3}$$

$$\chi''=\chi_0x^2/\pi. \tag{4}$$

Numerical calculations of Clem and Sanchez further show that the peak in χ'' occurs at $x=1.942$, which implies that the ac-field amplitude corresponding to the peak is $h_m=0.97J_cd$. Calculations of Brandt for square- and disk-shaped thin samples suggest that the difference in h_m for the two cases is only $\sim$0.2%.

Deviations from this quasistatic description of the mixed state, in which the flux front moves towards the center of the thin disc, occur only when the current J exceeds the critical current density J_c. These deviations are due to thermally activated flux-creep and flux-flow processes. A macroscopic description of the diffusive and creeplike motions of the flux into the virgin areas of the sample when it is subjected simultaneously to dc and collinear low-amplitude ac magnetic fields has been given by many workers.[25–27] van der Beek, Geshkenbein, and Vinokur[26] have derived a diffusionlike equation for flux motion by making use of the Maxwell's equations and the constitutive relation $\mathbf{E}=\rho(J,\omega,B,T)\mathbf{J}$, where ρ is the complex resistivity and $\mathbf{E}$ and $\mathbf{J}$ are the electric field and current density, respectively. In the flux-flow regime, the resistivity ρ is independent of $\mathbf{J}$, and the solution of the diffusion equation directly gives the ac penetration depth. The linear-response theory can be extended to the regime of temperature and field where the current-voltage curves are nonlinear by assuming a thermally activated current-dependent resistivity of the form

$$\rho(J,\omega,B,T)=\rho_0\exp[-U(J,\omega,B,T)/k_BT], \tag{5}$$

where ρ_0 is a preexponential term and k_B is Boltzman's constant. The activation energy $U(J,\omega,B,T)$ is given as

$$U(J,\omega,B,T)=k_BT\ln(1/\omega\tau), \tag{6}$$

where the relaxation time $\tau=\mu_0h_0^2/\rho_0J^2$. The barrier for vortex motion depends on the current density J. In the vortex solid phase with a weak zero-dimensional pinning disorder, the vortex-glass[28] or collective-creep[29] type of scenario best describes the motion of flux lines. These models predict that for a given temperature, flux density, and frequency, the activation energy $U(J)$ diverges algebraically with the decreasing current as $U(J)=U_c(J_c/J)^\mu$, where the power μ is a universal constant of the order unity in the vortex-glass picture. A similar type of behavior is also expected in the presence of correlated pinning disorder, albeit with different values of U_c, J_c, and μ.[30,31] Combining this form of the current-dependent barrier height with Eqs. (5) and (6), we can write the following explicit form for the screening current induced in the sample by the ac field;

$$J=J_c\left[\frac{k_BT}{U_c}\ln\left(\frac{1}{\omega\tau}\right)\right]^{-1/\mu}. \tag{7}$$

This expression shows that measurements of the frequency dependence of ac screening currents provide important information about the mechanism of flux penetration in a superconductor with quenched-in pinning disorder.

III. EXPERIMENTAL DETAIL

Thin films of $Tl_2Ba_2CaCu_2O_8$ were prepared on (100) $LaAlO_3$ substrates by a two-step process that involved rf magnetron sputtering of a BaCaCuO target followed by annealing in an atmosphere of controlled Tl_2O and O_2 partial pressures. Details of film preparation and x-ray-based characterization are described elsewhere.[32] Transport critical-current-density measurements were carried out on a photolithographically fabricated bridge ($100\times1000\,\mu$m^2) of the film in a four-probe geometry. The ac response of the sample in a perpendicular field configuration was measured using a local Hall-probe-based susceptometer, which consists of a GaAs/Ga$_x$Al$_{1-x}$As Hall sensor of effective area $25\times25\,\mu$m^2 mounted at the center of a 850-turn copper coil of inner diameter $\sim$0.7 cm. This coil, when driven with a sinusoidal voltage, provides an ac magnetic field of desired amplitude ($\leqslant$13.2 G) and frequency ($\leqslant$1073 KHz) perpendicular to the plane of the Hall probe. A square piece (5×5 mm^2) of Tl-2212 film is placed on the top of the Hall sensor such that the latter is located at the center of the film. This assembly is placed inside a stainless steel dip-stick cryostat for measurements down to 60 K in a subatmospheric-pressure liquid-nitrogen bath. The cryostat is placed between the pole pieces of a 1-T electromagnet. In all our measurements the dc field was always perpendicular to the plane of the sample. The Hall probe was biased with a highly stable

30 mA dc current from a battery source. In the present geometry, the Hall sensor measures the vector sum of the magnetic fields of all external sources and of the induced screening currents in the sample. The voltage output therefore has both ac and dc components. In this experiment we have measured only the fundamental- and third-harmonic components of the ac Hall voltage using a two-phase digital lockin amplifier. Following the procedure of Gilchrist and Konczykowski,[33] we calculate the fundamental- and third-harmonic transmittivities, $T_{H'}$ and $T_{H''}$, and $|T_{H3}|$, respectively, from the measured voltage at the fundamental- and third-harmonic frequency. The transmittivities are defined as

$$T_{H'}=\frac{[V'(f,T)-V'(f,T\ll T_c)]}{[V'(f,T\gg T_c)-V'(f,T\ll T_c)]}, \tag{8}$$

$$T_{H''}=\frac{[V''(f,T)]}{[V'(f,T\gg T_c)-V'(f,T\ll T_c)]}, \tag{9}$$

and

$$|T_{H3}|=\frac{V(3f,T)}{[V'(f,T\gg T_c)-V'(f,T\ll T_c)]} \tag{10}$$

where $V'(f,T)$ and $V''(f,T)$ are the inphase and quadrature components of the lockin voltage at a given frequency and temperature and $V(3f,T)$ is the rms value of the third-harmonic signal. The inphase and quadrature components of the fundamental transmittivities are related to the real and imaginary parts of the ac susceptibility χ as $\chi'=T_{H'}-1$ and $\chi''=T_{H''}$.

IV. RESULTS

We have measured the frequency and ac-field-amplitude dependence of $T_{H'}$, $T_{H''}$, and $|T_{H3}|$. In the absence of a dc field, the screening response of the sample does not show any frequency dependence over a frequency range of 73 Hz–1 kHz. This suggests that the static critical-state model can be applied to understand the magnetic state of the thin film. When a dc magnetic field is superimposed over the small ac field, the fundamental- and third-harmonic transmittivities show frequency as well as amplitude dependence. This behavior is indicative of thermally activated flux-creep and flux-flow processes. For a thin film in a transverse geometry, even a small dc field would lead to uniform distribution of flux lines in the plane of the sample. The screening currents induced by the ac field over the length scale of the order of the Bean length from the periphery of the sample exert Lorentz force on these flux lines. Depending on the temperature and flux density, the flux lines move towards the center of the sample during one half of the ac cycle and towards the edges during the other half. In a local Hall-probe geometry we can define a normalized shielding current density J^*, which is related to the broad-band transmittivity ($T_{H'}$) as[31,33]

$$J^*=1/\pi\cos^{-1}(2T_{H'}-1). \tag{11}$$

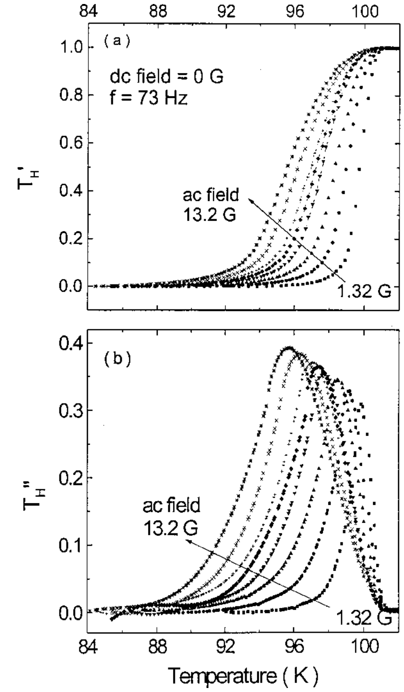

FIG. 1. (a) Temperature dependence of $T_{H'}$ measured at 1.32, 2.64, 3.96, 5.28, 6.6, 7.92, 10.6, and 13.2 G amplitudes of a 73-Hz ac field. (b) The corresponding variations of $T_{H'}$.

The frequency-dependent J^* computed from the measured $T_{H'}$ data can be compared with various models for the creep-like motion of various in high-T_c superconductors.

A. Ac response in the absence of a dc field

First we present results of screening measurements performed when only a small sinusoidally time-varying magnetic field is applied perpendicular to the plane of the sample. Figure 1(a) shows the inphase transmittivity ($T_{H'}$) measured as a function of temperature at different amplitudes of a 73-Hz ac field. The temperature at which $T_{H'}$ reaches unity has been identified as the critical temperature T_c of the sample. As the ac field amplitude is increased from 1.32 to 13.2 G, the superconducting transition broadens considerably. This large broadening of the transition is a characteristic feature of the layered cuprates such as Tl-2212 and Bi-2212, and it indicates a highly dynamic vortex state in these superconductors. In Fig. 1(b), we show the out-of-phase component of the transmittivity ($T_{H''}$). As expected for the imaginary part of the fundamental susceptibility, the $T_{H''}$ goes through a peak value on decreasing the temperature. The peak broadens and its position shifts to lower temperature at the higher values of h_0. We also observe a slight increase in the amplitude of the peak with increasing h_0. We have measured the broadband transmittivities ($T_{H'}$ and $T_{H''}$) at several temperatures as a function of the ac field amplitude. The static critical-state (SCS) model[1–3] shows that for a thin disc-shaped sample, if the $T_{H''}$ is measured as a function of h_0, the field h_m corresponding to the peak in $T_{H''}$ is related to J_c as $h_m=J_cd/1.03$. The $T_{H''}$ vs h_0 data therefore allow determination of the temperature dependence of J_c.

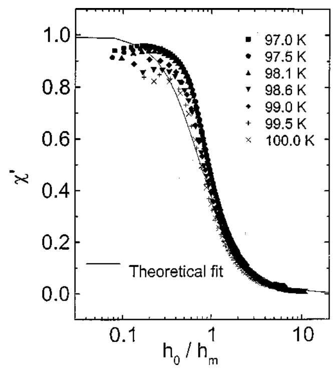

FIG. 2. Real part of the fundamental susceptibility (χ') plotted as a function of the normalized field (h_0/h_m) at different temperatures. The solid line is a calculation of χ' using Eq. (1) of the text.

We have calculated the critical current density using the data of Fig. 1(b). The J_c follows a temperature dependence of the type $J_c \sim J_{c0}(1-T/T_c)^{\beta}$ with J_{c0} and β equal to 2.5 $\times 10^7$ A/cm^2 and 1.7, respectively. The J_c deduced from the transport measurements using a 1-μV/cm field criterion also shows a similar temperature dependence albeit with a marginally lower value, which is perhaps due to some degradation of the film during lithographic processing.

The applicability of the SCS model to the present case of low-frequency and low-amplitude ac response is valid if it can be shown that the susceptibilities χ' and χ'' at all temperatures and frequencies scale with the normalized field h_0/h_m, as predicted by Eqs. (1) and (2). In Fig. 2, we show χ' vs h_0/h_m at seven different temperatures. A good scaling of the data with the normalized field is evident and the calculated χ' from Eq. (1) also agrees with the experimental data. The fact that normalized susceptibility is independent of temperature indicates that thermally activated processes are not operational in these low-ac-field amplitude measurements, and the SCS model gives a correct description of the ac response. We, however, could not establish the upper bound of the ac field to which the theory is applicable because of the limited field generated by the Hall-probe coil.

In the presence of a collinear dc magnetic field, the fundamental- and third-harmonic transmittivities become frequency dependent. In Figs. 3(a), 3(b), and 3(c) we show the behavior of $T_{H'}$, $T_{H''}$, and $|T_{H3}|$ as a function of temperature when both a 1-kG dc field and a 1.32-G ac field of variable frequency were applied to the sample. A progressive shift of the superconducting transition to higher temperatures is seen in both $T_{H'}$ and $T_{H''}$ data with the increasing frequency. The peak in third-harmonic transmittivity $|T_{H3}|$ becomes sharper, its amplitude decreases, and it shifts to higher temperatures as the frequency of the ac field is raised. The temperature at which the third-harmonic susceptibility becomes detectable on cooling, marks the onset of irreversible behavior in the sample. The third-harmonic susceptibility has also been measured for different values of the dc field while keeping the ac field frequency and amplitude fixed at 73 Hz and 1.32 G, respectively. The irreversibility line extracted from these measurements is shown in Fig. 4. In the inset of Fig. 4, we show the frequency dependence of the irreversibility line at 1 kG field. Extrapolation of the frequency dependence to zero frequency yields the true irreversibility temperature at 1 kG field. The irreversibility field $H_{\mathrm{irr}}(T)$ shown in Fig. 4 has a temperature dependence of the type $H_{\mathrm{irr}} \sim (1-T/T_c)^n$ with $n=1.2$. It is instructive to compare the behavior of $H_{\mathrm{irr}}(T)$ with the vortex-glass field $H(T_g)$ of similar Tl-2212 films deduced by Deak *et al.*[21] from the form of the resistivity in the critical regime of the vortex-glass transition. The power "n" in the work of Deak *et al.*[21] is 2.6 ± 0.5. For a vortex-glass to vortex-liquid phase boundary, however, the expected field dependence is of the type $H(T_g) \sim (T_c - T_g)^{4/3}$.[28]

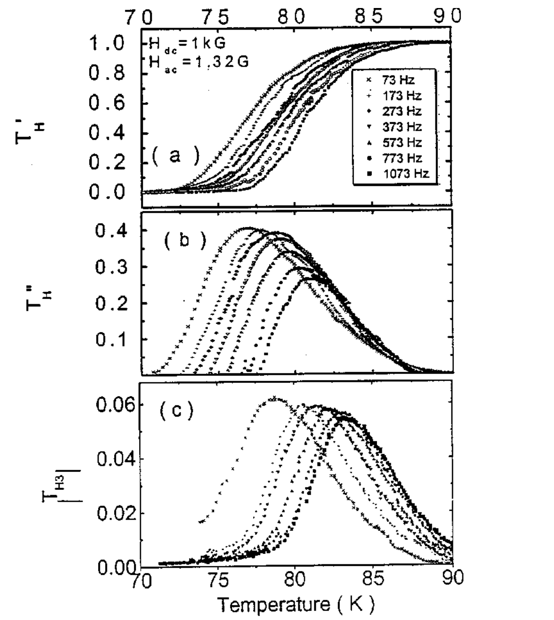

FIG. 3. Frequency dependences of $T_{H'}$, $T_{H''}$, and $|T_{H3}|$ when a 1-kG dc field is applied parallel to a 1.32-G, 73-Hz ac field are shown in panels (a), (b), and (c), respectively.

In the case of high-T_c cuprates of general formula $Bi_2Sr_2Ca_{n-1}Cu_nO_{2n+4}$ and $Tl_2Ba_2Ca_{n-1}Cu_nO_{2n+4}$ (where $n=2$ and 3), the copper-oxide planes are separated by insu-

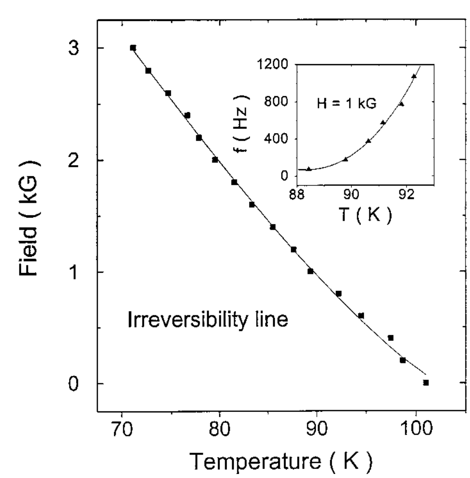

FIG. 4. Irreversibility line for the $Tl_2Ba_2CaCu_2O_8$ thin film. Solid line is a fit to equation $H_{\mathrm{irr}} \sim (1-T/T_c)^n$. Inset shows the frequency dependence of the irreversibility temperature at a 1-kG dc field.

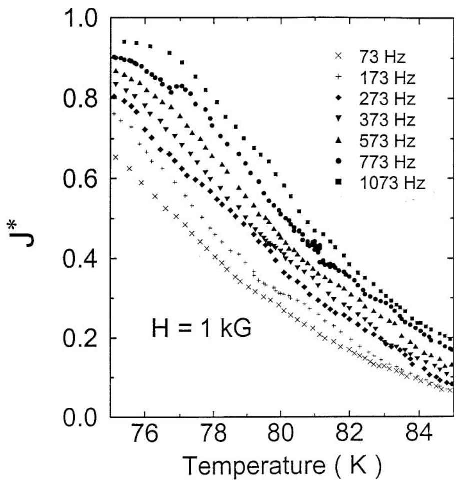

FIG. 5. Temperature dependence of the normalized screening current density J derived from the $T_{H'}$ data of Fig. 4(a).

lating (Bi/Tl)-O and (Sr/Ca)-O layers of total thickness ~15 Å. This large separation between the superconducting planes affects the mixed state significantly. A vortex in this case is not a stringlike object having uniform azimuthal supercurrents. For magnetic-field direction perpendicular to the CuO_2 planes, it is a string of planar current loops on the superconducting planes. These so-called pancake vortices couple along the field direction via Josephson interlayer coupling. The physics of a clean system is governed by two characteristic lengths: the vortex spacing $a=(\phi_0/B)^{1/2}$ and the Josephson length γs. Here ϕ_0 is the flux quantum, and s and γ are the interlayer spacing and anisotropy parameter, respectively.[18] The factor γs defines the crossover field $B_{cr} \sim \phi_0/(\gamma s)^2$ above which the intraplane interaction of pancakes exceeds the interlayer coupling. The vortex solid for a clean system at $B<B_{cr}$ forms the Abrikosov lattice as seen through neutron scattering[34] and muon spin-rotation experiments.[35] In the presence of quenched-in disorder, a vortex-glass phase is expected to form for $B<B_{cr}$ and show highly nonlinear dissipation.

For temperatures below the onset of third-harmonic susceptibility, the ac response of the system is indeed highly nonlinear. The behavior of the induced currents in the nonlinear regime is governed by nucleation and subsequent growth of vortex-loop excitations under the influence of the Lorentz force. The critical loop size becomes vanishingly small as J goes to zero. In Fig. 5, we show the temperature dependence of normalized screening current density deduced from the $T_{H'}$ data of Fig. 3 as calculated from Eq. (11). These measurements allow us to plot the normalized screening current density as a function of the drive frequency. The $J(\omega)$ in the vortex-glass/collective-creep scenario is expected to follow Eq. (7) in the temperature range where the third-harmonic susceptibility is nonzero. In Fig. 6 we plot the normalized current density as a function of the log of frequency. Solid lines in the figure are the fits to the vortex-glass/collective-creep model. The parameters deduced from these fits are U_c, J_c, τ, and the creep exponent μ. In Fig. 7(a) we show the variation of U_c and μ with temperature for the 1 kG field measurements. Similarly, variations for the relaxation time τ and J_c are shown in Fig. 7(b).

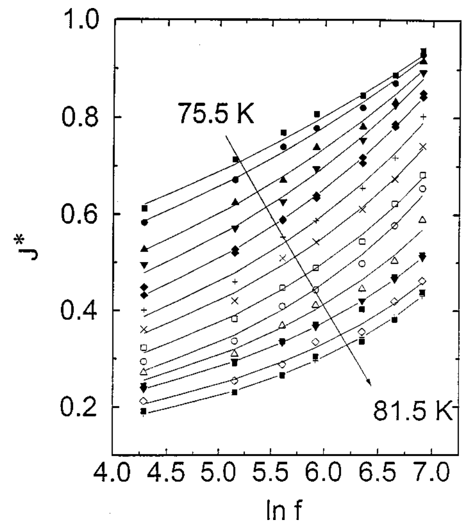

FIG. 6. Frequency dependence of the normalized J at temperatures ranging from 75.5 to 81.5 K in 0.5 K increments when a 1-kG dc field is applied normal to the film plane (parallel to c axis). Solid lines in the figure are fits to Eq. (7).

V. DISCUSSION

The dynamic behavior of vortices in a highly anisotropic superconductor such as Tl-2212 depends on the relative strengths of Josephson coupling across the planes, in-plane repulsion between pancakes, and pinning disorder vis-a-vis the thermal energy k_BT. As the vortices come closer with the increasing field, the in-plane interaction far exceeds the Josephson coupling, and the vortices in each plane become

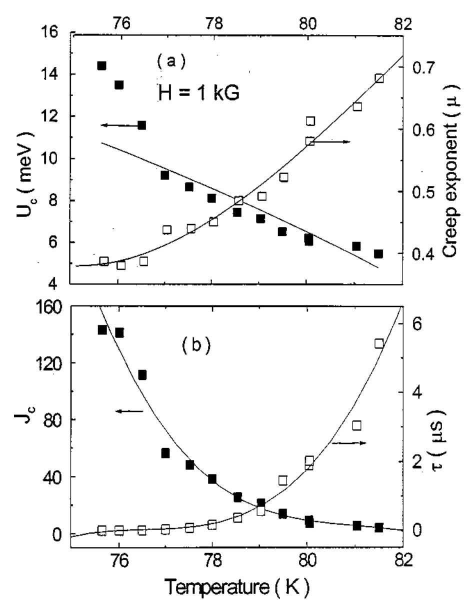

FIG. 7. (a) Temperature dependence of U_0 and μ at 1 kG field. (b) Temperature dependence of J_c and τ at the same field. Solid lines in the figure are guides to the eyes.

decoupled. In a clean system these decoupled vortices will freeze into a two-dimensional (2D) lattice at a temperature close to the solidification temperature of vortices in a superconducting layer.[36] For a single layer, this temperature is independent of field. In the presence of a weak pinning disorder, the layers of pancakes can behave at low temperatures as a stack of independent 2D vortex glasses with a diverging glass-correlation length.[28] At fields $B < B_{cr} (= 2\pi[\phi_0 \ln(\gamma K_{max} d)]/\gamma^2 d^2)$, where K_{max} determines the size of the vortex core and can be approximated as $\sim 1/\xi_{ab}$, where ξ_{ab} is the ab-plane coherence length,[37] the 2D vortices are correlated along the field direction (c axis), and a true 3D vortex-glass phase showing zero dissipation in the limit of infinitesimal driving force is formed. It has been argued that the 3D glass in highly layered superconductors undergoes a two-step transition on heating beyond T_g.[37] First, a liquid of stringlike vortices appears at T_g that eventually goes to a liquid of uncorrelated 2D vortices. The crossover field B_{cr} for Tl-2212 compound is ~400 G.

It has been argued by Glazman and Koshelev[37] that the 3D liquid in fact is also pinned by disorder, and just as in the case of the vortex-glass phase the linear resistance is zero here as well. The linear dissipation appears only when the system goes to the 2D liquid state. The field at which this happens is given as,[37]

$$\tilde{B}_m(T) = \frac{\phi_0^3(1-t)}{(4\pi\lambda_c)^2 dT}\alpha_m, \qquad (12)$$

where t is the reduced temperature, α_m is a universal constant ~0.1, and λ_c is the c-axis penetration depth. Since the present experiments focus only on the low-field behavior, the ac response of the system on cooling below T_c should, in principle, show this two-step transition.

Measurements of Deak *et al.*[21] on similar samples show two distinct forms of the temperature dependences for resistivity in the tail of the transition. In both these regimes, the resistivity follows the vortex-glass type of behavior $\rho(T) = (T-T^*)^{v(2-1)}$, albeit with different T^* and critical exponents (v and z). These authors[21] have attributed the higher and lower values of T^* to the decoupling transition in the liquid phase and the true vortex-glass transition T_g, respectively. The contrast between the model of Glazman and Koshelev and the measurements of Deak *et al.* is evident here. The latter's observation of linear resistivity and reversible magnetization above T_g contradicts the argument that in the vortex-line liquid-state disorder effectively pins each line and the linear resistance is zero. The issue whether the first-order melting transition and the loss of c-axis correlation in the liquid phase of quasi-2D superconductors occur at the same temperature or different temperatures has gained more attention recently.[38–40] The Josephson plasma resonance experiments[40] seem to suggest that the melting and loss of c-axis correlation in the liquid occur at the same temperature. This issue in clean superconductors, as well as in those with weak pinning disorder, is not yet fully settled.

In order to establish a correspondence between the loss of linear resistivity and the onset of third-harmonic susceptibility, we show in Fig. 8 the resistivity of a sample at 0, 1, and 2 kG fields along with its third-harmonic transmittivity. The shape of the ρ vs T curve in the figure is similar to the earlier reports on Tl-2212 films.[21,22,41] It is clear from the figure that $|T_{H3}|$ starts appearing at the temperature where $\dot{\rho}(T)$ vanishes. In fact the glass transition T_g deduced from a linear fit of the type $\rho(T) \sim (T-T_g)^x$, where $x = v(z-1)$ is the same as the onset temperature of $|T_{H3}|$. The value of x in our measurement is 1.8. This observation suggests that T_g is same as the irreversibility temperature. The resistivity for two fields reported here also indicates a change of slope at a slightly higher temperature, which is perhaps due to a 3D-to-2D vortex liquid transition identified by Deak and co-workers.[21]

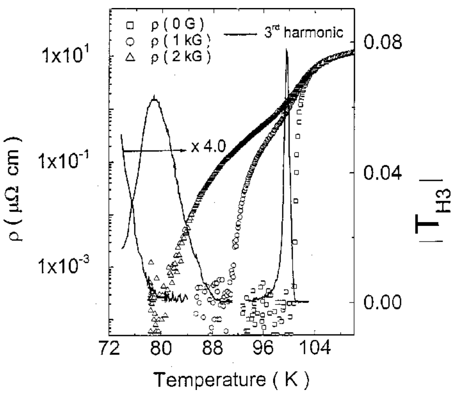

FIG. 8. Temperature dependence of the resistivity at 0-G, 1-kG, 2-kG dc fields applied perpendicular to the film plane. Figure also shows the third-harmonic transmittivity of the same film measured at 73 Hz.

We now discuss the behavior of the system in the regime of nonlinear response. The screening current J at low fields ($< B_{cr}$), where vortices are indeed rectilinear, can be fitted to the vortex-glass/collective-creep formalism [Eq. (7)]. The fitting procedure allows calculations of the current-independent activation energy $U_c(T,H)$ and the vortex-glass/collective-creep exponent μ. The latter is related to the critical size of the vortex loop, which grows under the action of the Lorentz force and leads to dissipation. The vortex-glass theory[28] predicts μ to be a universal constant of the order unity. The exponent μ in the collective-creep theory[29] depends on magnetic field, temperature, and current density, and varies from $\frac{1}{7}$ at low temperatures to $\frac{1}{2}$ close to the irreversibility line with a peak value of $\frac{3}{2}$ in between. For an anisotropic superconductor like Tl-2212, 2D collective creep may represent the correct form of dissipation in the vortex solid phase. Here also the exponent varies between $\frac{1}{3}$ and $\frac{9}{8}$.[18] The μ in the present study varies from 0.4 to 0.7 on increasing the temperature. Other workers have also reported a temperature- and field-dependent μ.[23,24,42–44] Magnetization measurements on $YBa_2Cu_3O_7$ (YBCO) crystals,[42] for example, reveal that the variation of μ with T at 10 kG has a domelike shape with a peak value of 1.8 at ~30 K. A systematic dependence of μ on the applied field H and temperature T has been deduced by Wen *et al.*[23] for ring-shaped Tl-2212 films using torque magnetometry. At low fields (<4 kG), they observe a monotonic increase in μ from ~0.05 to ~0.2 with temperature in the range 4.2 to ~60 K. This is followed by a sudden drop of μ to zero as the irreversibility temperature is reached. For fields between 4 and 7 kG, μ changes sign from positive to negative in the temperature range of 4.2 K and $T_{\rm irr}$. For $H>7$ kG however, μ is always negative for $T>4.2$ K and its extrapolated value at $T=0$ is zero. Wen *et al.*[23] associate the $\mu=0$ line on the *H-T* plane to a vortex-liquid to vortex-glass transition. Since for a 2D vortex glass $T_g=0$,[28] Wen *et al.* conclude that a dimensional crossover from a 3D to a 2D vortex-glass phase occurs around $H \sim 7$ kG in this system. A further study of magnetization relaxation on similar samples by Wen *et al.*[24] shows a 3D to 2D vortex-glass crossover at a field as low as 1 kG. These authors attribute the difference in the crossover field reported in the two studies[23,24] to a varying degree of structural defects in their films.

Recalling the relation $U(J)=U(T,H)(J_c/J)^{\mu}$, $U(T,H)$ is the effective barrier height when $J=J_c$. In 3D collective creep, it varies as

$$U_c \cong \left(\frac{T_c}{\sqrt{G_i}}\right)\left(\frac{J_c}{J_0}\right)^{1/2}$$

where J_0 is the depairing current density and G_i is the Ginzberg number.[18] Clearly, $U_c(T,H)$ is expected to go as $\sqrt{J_c}$ in temperature. Its value at 1 kG field is in agreement with the activation energy for flux motion in Bi-2212 deduced from the Arrhenious plots of the frequency and the peak temperature χ'',[45] but it is smaller than the reported values for YBCO (Ref. 46) and $Nd_{1.85}Ce_{0.15}CuO_4$ (Ref. 47) systems. One important parameter that comes out after fitting the data to Eq. (7) is the relaxation time τ. Blatter *et al.*[18] have shown that τ is a microscopic time of the order of $\sim 10^{-6}$ sec, which depends on sample size and geometry. The value of τ in our case varies from $\sim$0.1 to $\sim$5.5 μsec as the temperature is raised from 76 to 82 K. Similar values of τ have been reported in the case of Bi-2212 and YBCO crystals.[48,49] Furthermore, it is shown that $\tau \sim U_c^{-1}$ under the collective-creep model. This dependence is evident in Fig. 7.

VI. CONCLUSIONS

We have made use of a local Hall-probe-based susceptometer to measure the screening response of $Tl_2Ba_2CaCu_2O_8$ epitaxial films when an ac field of varying amplitude and frequency was applied perpendicular to the plane of the film. We have compared our measurements with a recent model for screening response of a disc-shaped thin superconductor subjected to an ac field normal to its plane. Our observation of a temperature-independent susceptibility, which scales with the normalized field, and the calculated J_c, which compares well with the measured transport J_c, lends support to the applicability of the model. The ac response when a dc field is applied collinear with the ac perturbation has the signature of thermally activated creep in the regime of temperature where the third-harmonic susceptibility is nonzero. The *H-T* phase boundary, below which the response is irreversible, varies as $H(T) \sim (1-t)^n$, with $n=1.2$. We find that the frequency dependence of screening currents in the irreversible regime is consistent with the vortex-glass/collective-creep picture. The critical parameters of the collective-creep theory such as U_c, μ, and τ have been inferred form the screening data. We have also extracted the low-field T_g from the temperature dependence of flux-flow resistivity at very low levels of dissipation. The glass-transition temperature is same as the irreversibility temperature deduced from the onset of the third-harmonic component to the susceptibility. While identification of T_g with the irreversibility temperature ($T_{\rm irr}$) over the entire phase diagram would require measurements of T_g for large values of the dc field, our data over a limited dc field range suggest that the response of 3D flux liquid is not irreversible.

ACKNOWLEDGMENTS

We thank Dr. Kanwaljeet Singh for assistance in the early stages of this work. Financial support for this research has come from a U.S.-India fund monitored by the Office of Naval Research through Grant No. N00014-97-1-0854. We also thank Dr. C. S. Pande, Dr. D. O. Welch, and Dr. M. Suenaga for helpful discussions.

[1] P. N. Mikheenko and Yu. E. Kuzovlev, Physica C **204**, 229 (1993).
[2] J. Zhu, J. Mester, J. Lochhart, and J. Turneaure, Physica C **212**, 216 (1993).
[3] John R. Clem and Alvaro Sanchez, Phys. Rev. B **50**, 9355 (1994).
[4] E. H. Brandt, Phys. Rev. Lett. **71**, 2821 (1993).
[5] E. H. Brandt, Phys. Rev. Lett. **74**, 3025 (1995).
[6] Th. Herzog, H. A. Radovan, P. Ziemann, and E. H. Brandt, Phys. Rev. B **56**, 2871 (1997).
[7] M. Wurlitzer, M. Lorenz, K. Zimmer, and P. Esquinazi, Phys. Rev. B **55**, 11 816 (1997).
[8] B. J. Jonsson, K. V. Rao, S. H. Yun, and U. O. Karlsson, Phys. Rev. B **58**, 5862 (1998).
[9] Charles P. Bean, Rev. Mod. Phys. **36**, 31 (1964).
[10] Y. Iye, T. Tamegai, H. Takeya, and H. Takei, Jpn. J. Appl. Phys., Part 2 **26**, L1057 (1998).
[11] T. T. M. Palstra, B. Batlogg, R. B. Van Dover, L. F. Schneemeyer, and J. V. Waszczak, Phys. Rev. B **41**, 6621 (1990).
[12] P. H. Kes, J. Aarts, J. van den Berg, C. J. van der Beek, and J. A. Mydosh, Supercond. Sci. Technol. **1**, 242 (1989).
[13] D. H. Kim, K. E. Grey, R. T. Kampwirth, and D. M. McKay, Phys. Rev. B **42**, 6249 (1990).
[14] M. W. Coffey and J. R. Clem, Phys. Rev. Lett. **67**, 386 (1991).
[15] V. B. Geshkenbein, V. M. Vinokur, and R. Fehrenbacher, Phys. Rev. B **43**, 3748 (1991).
[16] N. C. Yeh, Phys. Rev. B **43**, 523 (1991).
[17] A. E. Campbell and J. E. Evetts, Adv. Phys. **21**, 199 (1972).
[18] G. Blatter, M. V. Feigelman, V. B. Geshkenbein, A. I. Larkin, and V. M. Vinokur, Rev. Mod. Phys. **66**, 1125 (1994).
[19] G. W. Crabtree and D. R. Nelson, Phys. Today **50**(4), 38 (1997).
[20] J. D. Hettinger, D. H. Kim, K. E. Grey, U. Welp, R. T. Kampwirth, and M. Eddy, Appl. Phys. Lett. **60**, 2153 (1992).
[21] J. Deak, M. McElfresh, D. W. Face, and W. L. Holstein, Phys.

Rev. B **52**, 3880 (1995).
[22] R. C. Budhani, B. D. Weaver, and W. L. Holstein, Phys. Rev. B **50**, 3499 (1994).
[23] H. H. Wen, A. F. Th. Hoekstra, R. Griessen, S. L. Yan, L. Fang, and M. S. Si, Phys. Rev. Lett. **79**, 1559 (1997).
[24] Hai-hu Wen, Paul Ziemann, Henri A. Radovan, and Thomas Herzog, Physica C **305**, 185 (1998); Hai-hu Wen, H. A. Radovan, F.-M. Kamm, P. Ziemann, S. L. Yan, L. Fang, and M. S. Si, Phys. Rev. Lett. **80**, 3859 (1998).
[25] E. H. Brandt, Phys. Rev. Lett. **67**, 2219 (1991).
[26] C. J. van der Beek, V. B. Geshkenbein, and V. M. Vinokur, Phys. Rev. B **48**, 3393 (1993).
[27] A. E. Koshelev and V. M. Vinokur, Physica C **173**, 465 (1991).
[28] Daniel S. Fisher, Matthew P. A. Fisher, and David A. Huse, Phys. Rev. B **43**, 130 (1991).
[29] M. V. Feigelman, V. B. Geshkenbein, A. I. Larkin, and V. M. Vinokur, Phys. Rev. Lett. **63**, 2303 (1989).
[30] D. R. Nelson and V. M. Vinokur, Phys. Rev. Lett. **68**, 2398 (1992); Phys. Rev. B **48**, 13 060 (1993).
[31] C. J. van der Beek, M. Konczykowski, V. M. Vinokur, G. W. Crabtree, T. W. Li, and P. H. Kes, Phys. Rev. B **51**, 15 492 (1995).
[32] W. L. Holstein, L. A. Parisi, C. Wilker, and R. B. Flippen, Appl. Phys. Lett. **60**, 2014 (1992).
[33] J. Gilchrist and M. Konczykowski, Physica C **212**, 43 (1993).
[34] R. Cubitt, F. M. Forgan, G. Yang, S. L. Lee, D. McK. Paul, H. A. Mook, M. Yethitaj, P. H. Kes, T. W. Li, A. A. Menovsky, Z. Tarnawski, and K. Mortensen, Nature (London) **365**, 407 (1993).
[35] S. H. Lee, P. Zimmermann, H. Keller, M. Warden, I. M. Savic, R. Schauwecker, D. Zech, R. Cubitt, E. M. Forgan, P. H. Kes, T. W. Li, A. A. Menovsky, and Z. Tarnawski, Phys. Rev. Lett. **71**, 3862 (1993).
[36] D. S. Fisher, Phys. Rev. B **22**, 1190 (1980).
[37] L. I. Glazman and A. E. Koshelev, Phys. Rev. B **43**, 2835 (1991).
[38] L. N. Bulaevskii, M. P. Maley, and M. Tachiki, Phys. Rev. Lett. **74**, 801 (1995).
[39] D. Lopez, E. F. Righi, G. Nieva, F. de la Cruz, W. K. Kwok, J. A. Fendrich, G. W. Crabtree, and L. Paulius, Phys. Rev. B **53**, R8895 (1996).
[40] Y. Matsuda, M. B. Gaifullin, K. Kumagai, M. Kosugi, and K. Hirata, Phys. Rev. Lett. **78**, 1972 (1997).
[41] D. H. Kim, K. E. Grey, R. T. Kampwirth, J. C. Smith, D. S. Richeson, T. J. Marks, J. H. Kang, J. Talvacchio, and M. Eddy, Physica C **177**, 431 (1991).
[42] J. R. Thompson, Y. R. Sun, and F. Holtzberg, Phys. Rev. B **44**, 458 (1991).
[43] C. Dekker, W. Eidelloth, and R. H. Koch, Phys. Rev. Lett. **68**, 3347 (1992).
[44] C. J. van der Beek, P. H. Kes, M. P. Maley, M. J. V. Menken, and A. A. Menovsky, Physica C **195**, 307 (1992).
[45] C. J. van der Beek and P. H. Kes, Phys. Rev. B **43**, 13 032 (1991).
[46] Y. Yeshurun, A. P. Malozemoff, and F. Holtzberg, J. Appl. Phys. **64**, 5797 (1988).
[47] C. C. Almasen, C. L. Seaman, Y. Dalichaouch, and M. B. Maple, Physica C **174**, 93 (1991).
[48] D. A. Brawner, N. P. Ong, and Z. Z. Wang, Phys. Rev. B **47**, 1156 (1993).
[49] P. Svedlindh, C. Rossel, K. Niskanne, P. Norling, P. Nordblad, L. Lundgren, and G. V. Chandrashekhar, Physica C **176**, 336 (1991).