# Effects of Vortex Pinning and Thermal Fluctuations on the Josephson Plasma Resonance in $Tl_2Ba_2CaCu_2O_8$ and $YBa_2Cu_3O_{6.5}$

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We investigated the temperature dependence and $c$-axis magnetic field dependence of the Josephson plasma resonance in optimally doped $Tl_2Ba_2CaCu_2O_8$ thin films and underdoped $YBa_2Cu_3O_{6.5}$ (YBCO) ortho-II single crystals using infrared spectroscopy. We obtained the $c$-axis penetration depths, at low temperature, in zero fields of about 20 and 7 $\mu$m, respectively. While the temperature dependencies of the resonances in the two compounds are very similar, the magnetic field dependence in YBCO is much weaker. We attribute this weak magnetic field dependence to the lower anisotropy of YBCO and interpret the observed behaviors in terms of thermal fluctuations and pinning of pancake vortices.

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The Josephson plasma resonance (JPR) phenomenon [1–3] provides direct information about the Josephson coupling between the $CuO_2$ planes in the highly anisotropic high $T_c$ superconductors [4]. Studies of the temperature and magnetic field dependence of the JPR provide useful insights into the pinning and thermal fluctuations of vortices in the cuprate family [5–10]. It is natural to anticipate unusual magnetic field behavior in these highly anisotropic materials. Indeed, just as in any conventional type-II superconductor, if a dc magnetic field is applied perpendicular to the 2D planes it penetrates the sample in the form of vortices which, when the anisotropy is large, are not vortex lines but pancake vortices. For low magnetic fields and low temperatures, these pancake vortices *may* form a nearly perfect vortex lattice. Under these conditions there should be no magnetic field dependence of the JPR since $\omega_p^2(B,T) = \omega_p^2(0)\langle\cos(\phi_{n,n+1})(r)\rangle$, where $\omega_p^2(0)$ is the plasma frequency in zero field and $\phi_{n,n+1}(r)$ is the gauge-invariant phase difference between the neighboring layers $n$ and $n+1$ [8]. In zero field, $\phi_n(r)$ is independent of $r$ and independent of $n$ for zero $c$-axis current. In nonzero field, $\phi_n(r)$ becomes position dependent but $\phi_{n,n+1}(r)$ is everywhere zero, provided the vortex cores are lined up, as in a perfect lattice. However, due to the effects of pinning (at low temperatures) and thermal fluctuations (at higher temperatures), pancake vortices do not in general line up from layer to layer. Both regimes have been extensively studied theoretically by Bulaevskii and one of us (A. E. K.) [9–11], but there is very little data on the JPR in the low magnetic field regime [7]. In this paper we present our studies of the temperature and field dependence of the JPR in two members of the high $T_c$ family with differing anisotropies.

Two types of reflectivity spectra were measured, each optimized to the type of sample: First normal incidence reflectivity spectra were recorded of a mosaic made of six $ac$-oriented single crystals of $YBa_2Cu_3O_{6.5}$ with ortho-II oxygen ordering, and a superconducting transition temperature of 58 K. The mosaic was held together with epoxy and had an area of 6 mm$^2$. Second, grazing incidence reflectivity measurements were made on epitaxially grown $Tl_2Ba_2CaCu_2O_8$ thin films, with a superconducting transition temperature of 98 K. Dimensions of the films in the $ab$ plane were 64 mm$^2$, and the thickness was 600–700 nm. Details of the sample preparation and growth methods have been reported elsewhere [12,13]. Here we report data in the far infrared region (17–700 cm$^{-1}$) obtained using a Fourier transform spectrometer. A grazing angle of 80° was chosen to probe predominantly the $c$-axis response for $p$ polarization of the light. Absolute reflectivities were obtained by calibrating the reflectivities against the reflectivity of a gold film deposited *in situ* on the sample.

The reflection experiments were carried out in magnetic fields oriented along the $c$ axis, which could be varied continuously between 0 and 0.8 T using a sliding permanent magnet, and were calibrated using a GaAs 2D electron gas placed in the sample holder block, 1 mm underneath the sample. Details of the magnet-cryostat design are described elsewhere [14].

In Fig. 1 we show the grazing incidence reflectivity of $Tl_2Ba_2CaCu_2O_8$ (Tl2212) for a range of temperatures in zero field and for various magnetic fields at 4 K under field cooled (FC) conditions: The sample temperature was first increased above the phase transition, next the field was set at the desired value, the sample temperature was lowered to 4 K, and the spectra were recorded. This sequence was then repeated for each external field value. In Fig. 2 we present the FC normal incidence reflectivity of YBCO, with the electric field along the $c$ axis in the frequency region from 18–100 cm$^{-1}$ for various magnetic fields at 4 K. In the figures we show only the low frequency part

 

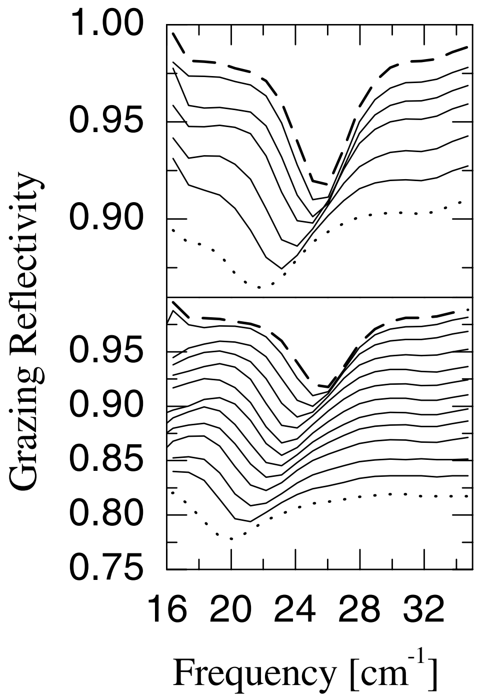


FIG. 1. Top panel: Grazing incidence reflectivity of Tl2212 without external magnetic field at (from top to bottom) 10, 20, 30, 40, 50, 60, and 70 K. Bottom panel: Field cooled grazing incidence reflectivity of Tl2212 at 4 K in an external magnetic field of (from top to bottom) 0, 0.03, 0.06, 0.11, 0.17, 0.21, 0.23, 0.26, 0.29, 0.32, 0.36, 0.4, and 0.46 T, perpendicular to the surface. Except for the zero field data, all curves have been given vertical offsets.

of the spectra since the only features in the spectra below 50 $\mathrm{cm}^{-1}$ correspond to the JPR. However, for the spectral analysis and accurate determination of the plasma resonance frequency, the entire measured spectral range was used. Figure 1 reveals that the JPR shifts down quasilinearly as a function of magnetic field for all temperatures (about 10 $\mathrm{cm}^{-1}$ per tesla). The gradual redshift without a change of amplitude upon increasing the magnetic field indicates that the resonance is a Josephson plasma resonance for all fields considered, and not a vortex mode [15]. The FC magnetic field dependence of the JPR of Tl2212 is displayed in Fig. 4 for each temperature. For YBCO the JPR has a much lower field dependence of about 0.5 $\mathrm{cm}^{-1}$ per tesla (see Fig. 2).

The $c$-axis penetration depths extracted from the zero field plasma frequency at 4 K are 19.8 $\mu$m in Tl2212 and 7.1 $\mu$m in YBCO, using $\epsilon_S = 9.1$ [16] and $\epsilon_S = 28$ [17], respectively. Previously we have addressed the issue of the temperature dependence of the resonance [16] in Tl2212 and Tl2201, from which we obtained the temperature dependence of the $c$-axis penetration depth, $\lambda_c^2(0)/\lambda_c^2(T) = 1 - (T/T_0)^\eta$ with $\eta \sim 2$. From the present data on YBCO we obtained the same temperature dependence.

In addition we measured reflectivity data under zero field cooled (ZFC) conditions. These data are displayed in

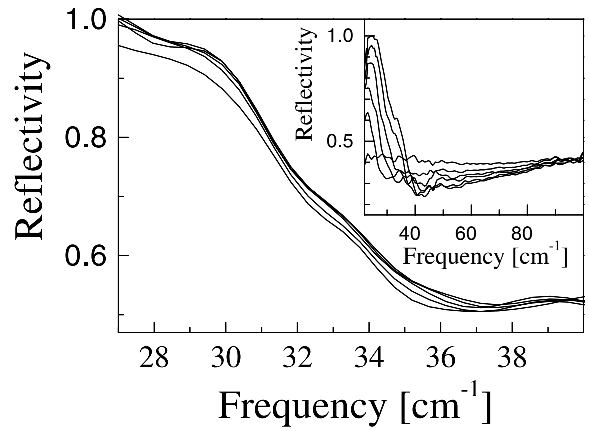


FIG. 2. Field-cooled normal incidence reflectivity of YBCO. At 4 K in an external magnetic field of (from top to bottom) 0, 0.2, 0.4, 0.6, and 0.8 T, perpendicular to the surface. Inset: Reflectivity of YBCO without external magnetic field at (from top to bottom) 4, 30, 40, 50, and 70 K.

Fig. 3, showing that under ZFC conditions the resonance stays at the same frequency, whereas the *intensity* of the resonance drops gradually upon raising the field. The intensity also exhibits hysteresis during cycling of the magnetic field (not shown). Briefly, this behavior is caused by the self-screening of the sample: As the field is applied perpendicular to the film, the outer parts of the film are penetrated by the field, while the central part of the film remains at $B = 0$. Assuming that the JPR is absent from the field penetrated region of the film, the spectra are just linear combinations of the zero field spectrum and the featureless high field spectrum. One can make this quantitative using the result of Mikheenko and Kuzovlev [18] for the distance $a$ from the center of the disk that is penetrated by the magnetic field: $a = R/\cosh[H/(2\pi j_c^{ab} d)]$, where $R$ is the radius of the disk, $H$ is the applied field, $j_c^{ab}$ is the $ab$-plane critical current, and $d$ is the thickness of the film. From comparison to Fig. 3 we obtained

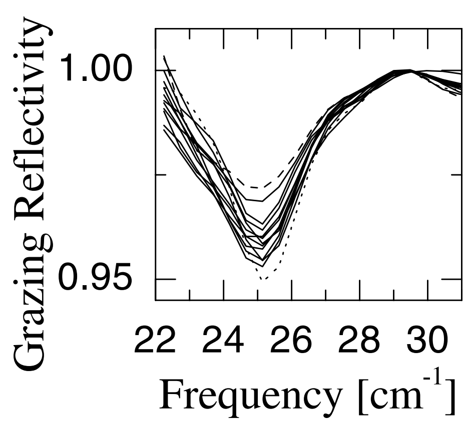


FIG. 3. Zero field cooled grazing incidence reflectivity of Tl2212 at 10 K in an external magnetic field of (from top to bottom) 0.46, 0.39, 0.32, 0.25, 0.18, 0.12, 0.06, and 0 T, perpendicular to the surface.

$j_c^{ab} = 6.7 \times 10^{10}$ A/m$^2$, which is a reasonable value for a high quality film.

The resonance is due to the zero crossing of $\mathrm{Re}\epsilon_c(\omega)$, for which we adopt the two-fluid expression appropriate for the superconducting state

$$\epsilon_c(\omega) = \epsilon_S(\omega) - \frac{c^2}{\lambda_c^2(B,T)\omega^2} + \frac{4\pi i \sigma_{qp}(\omega)}{\omega}, \quad (1)$$

where $\lambda_c$ is the $c$-axis penetration depth and $1/\lambda_c(B,T)^2$ represents the superfluid density appropriate for the $c$ direction. The pole strength of the second (London) term, coming from the condensate, provides directly, within a Josephson coupled layer model, the critical current in the $c$ direction. Oscillations of the condensate are electromagnetically coupled to interband transitions and lattice vibrations, which are represented by the dielectric function $\epsilon_S$. The third term arises from dissipative currents due to quasiparticles and determines the width of the plasma resonance in our spectra [16]. The condition for propagation of longitudinal modes in the medium along the $c$ direction in the long wavelength limit is that $\epsilon_c(\omega) = 0$, which occurs at the JPR frequency $\omega_J = c/\{\lambda_c\sqrt{\mathrm{Re}\epsilon_S(\omega_J)}\}$. By fitting Eq. (1) to the measured spectra, using the full Fresnel expression for the grazing/normal incidence reflectivity of uniaxial crystals [16], we obtain accurate values of the $c$-axis superfluid density, $\lambda_c^2(0)/\lambda_c^2(B,T)$, as a function of temperature and magnetic field.

As mentioned above, a dc magnetic field applied along the $c$ axis penetrates the cuprate superconductors in the form of pancake vortices. For sufficiently low fields and temperatures they form a vortex lattice, which for an ideal material would be composed of straight lines along the $c$ axis, and as a consequence have no influence on the interlayer Josephson coupling. If the coupling between the layers is very weak, pinning and thermal fluctuations may be sufficient to destroy the alignment of the pancake vortices along the field direction. The resulting spatial fluctuations of the relative phases between the planes cause a suppression of the Josephson coupling between the layers. Below the irreversibility line, the vortex phase diagram is very rich [7,19], and a similar richness is required of theories aimed at explaining all parts of the diagram. In the low temperature regime the main cause of misalignment of the pancakes is pinning due to impurities in the materials. At higher temperatures, thermal fluctuations take over and become the main cause for suppression of the Josephson current. Suppression of the Josepson coupling due to a magnetic field applied along the $c$ axis in the vortex lattice state was extensively studied theoretically by one of us (A. E. K.) in both the low temperature (where the pinning is the most important contribution) and the high temperature (where thermal fluctuations play the most important role) regimes. In Ref. [11] a model was developed where the influence of weak pinning on the Josephson coupling between layers was considered. A characteristic "decoupling" field $B_w$ was introduced, above which the JPR frequency saturates at a certain value. The vortex lines are destroyed by pinning, and $B_w$ is determined by the wandering length of the vortex lines. The following expression was derived for the effective Josephson coupling

$$E_J^{\mathrm{eff}} - E_J = -\frac{E_J B}{B_w}. \quad (2)$$

Here $E_J$ is Josephson energy per unit area, which is related to the $c$-axis penetration depth via $E_J = \Phi_0^2/(16\pi^3\lambda_c^2 d)$, where $\Phi_0$ is the elementary flux quantum and $d$ is the interlayer distance. With this expression we can calculate the Josephson energies per unit area with the result $E_J = 1.6 \times 10^{-4}$ erg/cm$^2$ for Tl2212, and $E_J = 1.5 \times 10^{-3}$ erg/cm$^2$ in YBCO. In the inset of Fig. 4 we show the fit to Eq. (2) at 4 K from which we obtain values for $B_w = 1.40 \pm 0.02T$ in Tl2212. In YBCO the field dependence is much weaker (Fig. 2), corresponding to $B_w = 30 \pm 10T$. Now we can pose the question, "Why does YBCO have a weaker field dependence than Tl2212?" From the discussion above it is clear that it can be due to either the higher purity of the sample or the more three dimensional nature of YBCO. The theoretical prediction for $B_w \sim E_J\Phi_0/U_p$ from Ref. [11], where $U_p$ represents the pinning potential can help us resolve this problem. The ratios of the decoupling fields for Tl2212 and YBCO should scale as $B_w^{Tl}/B_w^{Y} = [E_J^{Tl}/(E_J^{Y}) \cdot (U_p^{Y}/U_p^{Tl})]$. Since we have already determined the Josephson coupling constants $E_J$ and the decoupling fields $B_w$, we have sufficient experimental information to estimate the ratio of the pinning potentials, providing $U_p^{Y}/U_p^{Tl} \approx 0.44$. In the strongly anisotropic material Bi2212 $\lambda_c(0) = 110$ $\mu$m [7], corresponding to $E_J = 6.0 \times 10^{-6}$ erg/cm$^2$, while $B_w \approx 0.12$ T [7]. In the same way as for YBCO

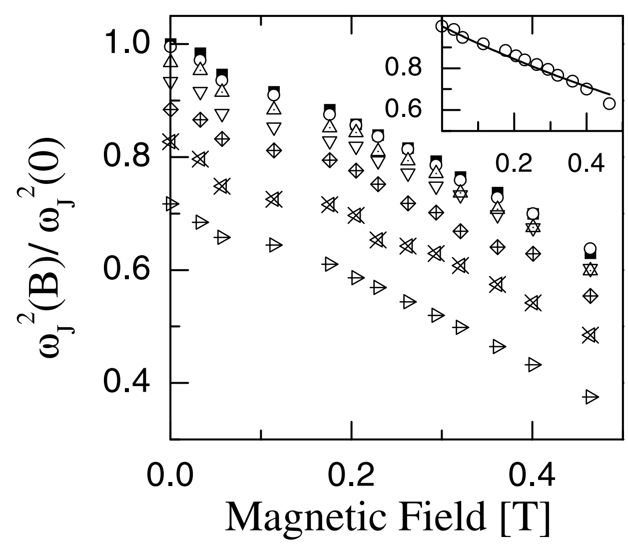


FIG. 4. Magnetic field dependence of $\omega(B,T)^2$ normalized to the value at zero field at 4 K (field-cooled) in Tl2212 for temperatures from top to bottom: 10, 20, 30, 40, 50, 60, and 70 K. Inset: 4 K data, with a fit to Eq. (2), using $B_w = 1.40$ T.

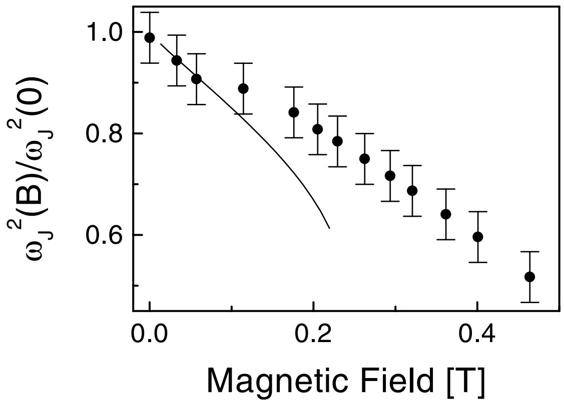


FIG. 5. Comparison between the experimental values of $\omega(B,T)^2/\omega(0,T)^2$ of Tl2212 at 70 K and the theoretical prediction based on Ref. [10], with the parameters $\lambda_{ab} = 245$ nm and $\lambda_c = 24.2$ $\mu$m for the in-plane and $c$-axis penetration depth, respectively.

we can estimate that $U_p^{Bi}/U_p^{Tl} = 0.43$. The ratios of the decoupling fields of these three compounds are $B_w^{Bi}{:}B_w^{Tl}{:}B_w^{Y} = 1{:}12{:}250$. The ratios of the pinning potentials are $U_p^{Bi}{:}U_p^{Tl}{:}U_p^{Y} = 1{:}2{:}1$, which does not correlate with $B_w$ in Bi2212, Tl2212, and YBCO. On the other hand, the relative values of the Josephson coupling energies are $E_J^{Bi}{:}E_J^{Tl}{:}E_J^{Y} = 1{:}27{:}250$, from which we can conclude that the spread of $B_w$'s over 2 orders of magnitude is a consequence of the strong material dependence of the interlayer Josephson coupling constant.

As temperature is increased towards $T_c$, a crossover may take place to a different type of magnetic field suppression of the JPR, where the pancake vortices become progressively misaligned due to thermal fluctuations. In Ref. [10] it was shown for low fields in the thermal fluctuation regime that $\omega_J^2(0,T) - \omega_J^2(B,T) \propto B$ in the low field region, with a gradual downward curvature at higher field values, similar to our experimental data [10]. The numerical calculation for the present case is shown in Fig. 5, for $T = 70$ K, using the experimental values for $\lambda_{ab}(T) = 245$ nm [12,20,21] and $\lambda_c(T) = 24.2$ $\mu$m [16] as the (only, experimentally fixed) parameters. We see, that, although the field dependence predicted by the thermal fluctuation model is close to the observations at 70 K, still the experimental field dependence is weaker than in the calculation. This is probably due to the fact that there is still a significant contribution from pinning even at 70 K. This is also borne out by the fact that in Fig. 4 there is no indication of a sudden change of the field dependence as a function of temperature. A quantitative theoretical framework of these observations would require a unified theoretical model, which takes into account both the influence of pinning and of thermal fluctuations on the Josephson coupling.

In conclusion, we have studied the $c$-axis Josephson plasma resonance in Tl2212 and YBCO as a function of temperature and low magnetic field. In the low temperature regime the observed behavior of Tl2212 fits very well the model of vortex disorder due to weak pinning [11] while above 50 K, good quantitative agreement with model calculations provides evidence for the presence of disorder due to thermally fluctuating pancakes [10]. We attribute the weaker magnetic field dependence of the resonance in YBCO to the lower anisotropy.

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[1] P. W. Anderson, in *Progress in Low-Temperature Physics,* edited by C. J. Gorter (North-Holland, Amsterdam, 1967), Vol. V.
[2] A. J. Dahm *et al.,* Phys. Rev. Lett. **20**, 859 (1968).
[3] V. N. Gubankov *et al.,* Supercond. Sci. Technol. **5**, 168 (1992).
[4] D. van der Marel *et al.,* Physica (Amsterdam) **341C–348C**, 1531 (2000).
[5] O. K. C. Tsui *et al.,* Phys. Rev. Lett. **73**, 724 (1994).
[6] Y. Matsuda *et al.,* Phys. Rev. Lett. **75**, 4512 (1995).
[7] M. B. Gaifullin *et al.,* Phys. Rev. Lett. **84**, 2945 (2000).
[8] L. N. Bulaevskii *et al.,* Phys. Rev. Lett. **74**, 801 (1995).
[9] L. N. Bulaevskii *et al.,* Phys. Rev. B **61**, R3819 (2000).
[10] A. E. Koshelev and L. N. Bulaevskii, Physica (Amsterdam) **341C–348C**, 1503 (2000).
[11] A. E. Koshelev *et al.,* Phys. Rev. B **53**, 2786 (1996).
[12] B. A. Willemsen *et al.,* Phys. Rev. B **58**, 6650 (1998).
[13] M. M. Eddy, U.S. Patent No. 5,476,836 (1995); U.S. Patent No. 5,508,255 (1996).
[14] D. Dulic, Ph.D. thesis, University of Groningen, 2000.
[15] N. B. Kopnin *et al.,* Phys. Rev. Lett. **74**, 4527 (1995).
[16] D. Dulić *et al.,* Phys. Rev. B **60**, R15051 (1999).
[17] C. C. Homes *et al.,* Physica (Amsterdam) **296C**, 230 (1998).
[18] P. N. Mikheenko and Yu. E. Kuzovlev, Physica (Amsterdam) **204C**, 229 (1993).
[19] M. F. Goffman *et al.,* Phys. Rev. B **57**, 3663 (1998).
[20] T. Jacobs *et al.,* Phys. Rev. Lett. **75**, 4516 (1995).
[21] Shih-Fu Lee *et al.,* Phys. Rev. Lett. **77**, 735 (1996).