Excitation of the Josephson Plasma Mode in $Bi_2Sr_2CaCu_2O_{8+\delta}$ in an Oblique Field

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We present evidence identifying the resonance in microwave absorption observed in $Bi_2Sr_2CaCu_2O_{8+\delta}$ in a field **H** as the Josephson plasma mode. The resonance field displays an unusual reentrant cusp when **H** is very close to alignment with the layers. By fitting to a model of Bulaevskii *et al.*, we obtain an anisotropy parameter $\gamma = 400-420$ at 5 K, and a Josephson plasma frequency ω_J equal to 5.3 cm^{-1} (at $H = 0$). We also describe the dependence of ω_J on H and T when the applied field $\mathbf{H} \parallel \mathbf{c}$.

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When the superconductor $Bi_2Sr_2CaCu_2O_{8+\delta}$ (Bi 2212) is exposed to microwave radiation in the mixed state, an unexpectedly sharp resonance in the absorption is observed as a function of the applied field [1]. Because the resonance field B_0 varies significantly with radiation frequency (ω) and temperature (T), the resonant absorption was tentatively described in terms of a collective mode of the vortex lattice hybridized with the cyclotron mode of quasiparticles [1]. However, several researchers [2] have pointed out that the resonance may correspond to the Josephson plasma oscillations. In $La_{2-x}Sr_xCuO_4$, a sharp Drude edge has been observed [3] at 60 cm^{-1} (for $x = 0.16$) in the c-axis reflectivity (in zero field) and has been attributed to the Josephson plasma mode [4]. The larger anisotropy in Bi 2212 implies a lower plasma frequency $\nu_J(0)$. A search down to 30 cm^{-1} did not detect its presence [5]. Torque experiments obtain estimates for the anisotropy parameter γ varying from >200 [6] to ~ 900 [7] [equivalent to frequencies $\nu_J(0) < 12$ cm^{-1}]. An intense magnetic field **H** may further suppress the mode to within the range of our microwave experiment (1–3 cm^{-1}). A particularly decisive test of whether the resonances are Josephson plasma excitations has been proposed by Bulaevskii, Maley, Safar, and Domînguez (BMSD) [8]: When **H** is rotated into alignment with the CuO_2 layers, the resonance field should display a reentrant cusp analogous to that observed in the c-axis resistivity [8,9]. We present measurements confirming the reentrant behavior of B_0 in Bi 2212. Moreover, using BMSD's model, we show that the angular dependence near alignment provides a novel, sensitive way to determine the value of γ (400–420) and $\nu_J(0)$ (~ 5.3 cm^{-1}).

We recall that, in a junction, a fluctuation of the relative phase ϕ from zero lowers the Josephson coupling energy and induces a supercurrent which leads to charging of the electrodes. The Josephson plasma mode [10] corresponds to the charging energy oscillating 90° out of phase with the Josephson coupling energy at the frequency $\omega_J = [(2\pi/\Phi_0)I_0/C]^{1/2}$, where I_0 is the critical current, Φ_0 the flux quantum, and C the junction capacitance. For a layered superconductor in zero field, the c-axis critical current density J_0 is related to γ by $J_0 = (\Phi_0/2\pi)/\mu_0 s\gamma^2\lambda_{ab}^2$, where λ_{ab} is the in-plane penetration depth, s the layer spacing, and μ_0 the permeability constant. In the mixed state, the presence of vortices strongly suppresses the Josephson current $J_m(B)$ (B is the induction) [11]. At $T = 0$, the reduction is expressed by the factor $J_m(B)/J_0 = \langle\cos\phi_{n,n+1}\rangle$, where $\phi_{n,n+1}$ is the gauge-invariant phase difference between adjacent layers indexed by n (where $\langle\cdots\rangle$ denotes spatial and disorder averaging) [11]. We may express the field-dependent Josephson plasma frequency as

$$\omega_J^2(B) = (2\pi/\Phi_0)(sJ_0/\varepsilon_0\varepsilon_r)\langle\cos\phi_{n,n+1}\rangle \qquad (T \to 0) \tag{1}$$

(ε_0 is the permittivity constant and ε_r the dielectric constant). The field affects ω_J through the phase factor $\langle\cos\phi_{n,n+1}\rangle$ only.

In our experiment, modulation of the microwave power produces a slight oscillation in the sample's temperature, which is measured with a cernox bolometer by lock-in detection [1,12]. The oscillation amplitude ΔT is proportional to the surface resistance R_s in the mixed state. As shown in Fig. 1, the sample platform (inside a cylindrical waveguide of radius 3.5 mm) may be rotated to vary the angle θ between **H** and the CuO_2 layers. In all runs, ω, θ, and T are held fixed, and R_s is recorded as H is swept from 0 to 7 T and back to 0 T.

Figure 1 displays a subset of the 32 traces taken at $T = 5$ K. Starting at large, negative tilt angle (lowest trace) the resonance field B_0 (arrow) moves rapidly to

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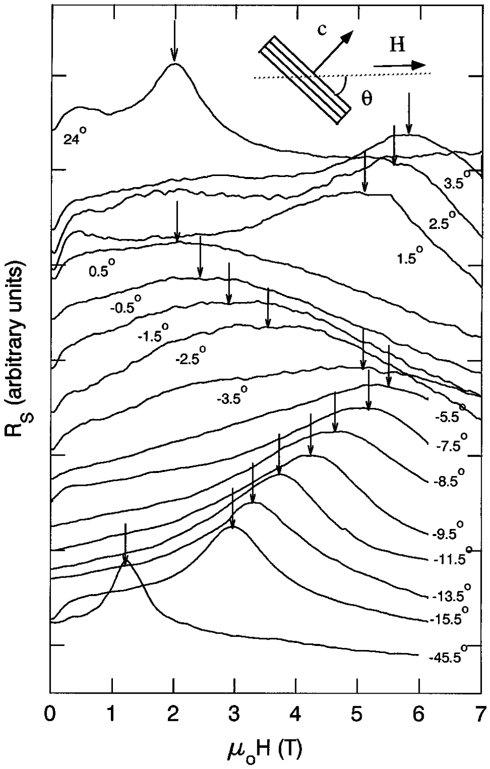

FIG. 1. Variation of the resonance field B_0 (arrows) with tilt angle θ in $Bi_2Sr_2CaCu_2O_{8+\delta}$ (sample 1), measured at 5.0 K at $\omega/2\pi = 47.3$ GHz (θ is defined in inset). **H** is fixed parallel to the cylindrical waveguide axis. The microwave E field $\mathbf{E}_\omega$ lies in the plane of **H** and **c**. Each trace represents the surface resistance R_s vs the applied field H (only sweep-up traces shown). As θ increases from –5.5° and 0, B_0 falls to a deep minimum. At small, positive θ (1.5°, 2.5°), we observe a discontinuous jump of B_0 from 2 to 5.5 T. Further increase of θ to 24° moves the peak back to low fields. At $\theta = 90°$ (not shown), the coupling to the radiation is greatly reduced, but B_0 at 0.8 T may still be resolved.

higher fields as we tilt **H** away from the c axis. While for large $|\theta|$, $B_0(\theta)$ scales approximately as $H\sin\theta$, deviation from this simple form is apparent at $|\theta| \sim 10°$. At $\theta = -5.5°$, B_0 attains a local maximum value of 5.6 T, while the resonance broadens noticeably. Remarkably, as **H** is brought into closer alignment with the layers, B_0 rapidly decreases to a deep minimum at 1.8 T. At positive θ, B_0 attains a second maximum (6.2 T) and subsequently approaches the approximate form $H\sin\theta$ (top trace). A clearer perspective is obtained by plotting the parallel component of the resonance field $B_0\cos\theta$ in each of the 32 traces vs $B_0\sin\theta$ [Fig. 2(a)]. Starting at the left ($\theta = -45.5°$), the trajectory of B_0 rises almost vertically. As θ approaches zero, the trajectory goes through a maximum before falling into a deep cusp. With further increase in θ, the curve goes through another maximum before falling to the right (we discuss the slight asymmetry below). The reentrant behavior of B_0 vs θ is a striking confirmation of Bulaevskii's prediction.

In parallel field ($\theta = 0$), flux lines exist as Josephson vortices trapped in the potential trough between CuO_2 bilayers in triangular coordination. The Josephson vortex array leads to a strong reduction of the Josephson energy [8]. At finite tilt angle θ, pancake vortices appear. BMSD [8] find that at $T = 0$ the total Josephson energy is optimized when the pancake vortex array assumes a zigzag distortion. The distortion leads to a recovery of a fraction of the loss in Josephson energy engendered by the parallel array (pinning neglected). Thus, $J_m(B)$ increase steeply as **H** is tilted out of the plane, in qualitative agreement with the angular dependence of the c-axis resistivity [8,9]. We may express ω_J^2 in an oblique field as (c is the velocity of light) [8]

$$\omega_J^2 = 2(c/s)^2(\Phi_0/Bs^2)(\varepsilon_r\gamma^4)^{-1}\{(2\pi)^{-2}(\Phi_0/\lambda_{ab}^2 B) + \theta[\ln(\lambda_{ab}^2/\pi s^2) + 2(\pi\lambda_{ab}sB/\Phi_0)^2]^{-1}\} \quad (T \to 0,\ \theta \ll 1). \tag{2}$$

The term in θ represents the compensation effect of the pancake vortices. Setting $\omega_J/2\pi = 47.3$ GHz, we solve

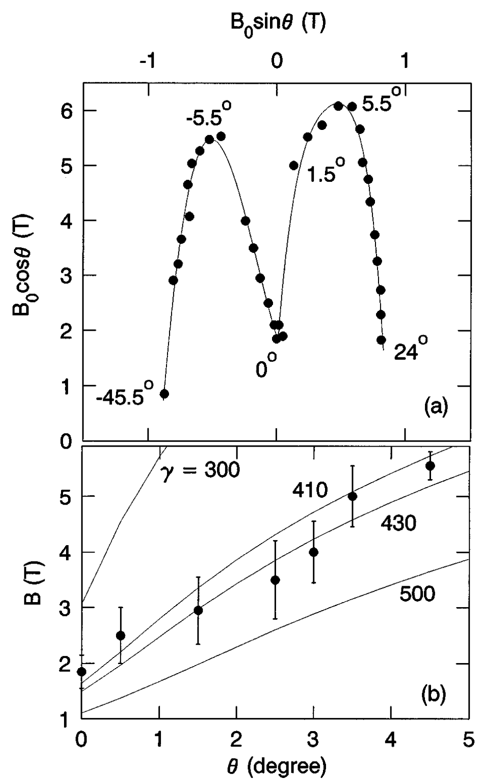

FIG. 2. (a) Plot of the parallel component of the resonance field $B_0\cos\theta$ vs the perpendicular component $B_0\sin\theta$ with tilt angle as a parameter (same conditions as in Fig. 1). The trajectory shows a striking reentrant cusp at small $|\theta|$. The uncertainty in determining the $\theta = 0$ position is $\pm 0.3°$. The data were taken in the sequence $-45° \to 0 \to 24°$. The line is a guide to the eye. (b) Comparison of the resonance field B_0 vs 0 with the prediction of BMSD's model [8] [Eq. (2)]. The lines are curves for $\gamma = 300$, 410, 430, and 500 (with $s = 15.6$ Å, $\lambda_{ab} = 1400$ Å, and $\varepsilon_r = 25$). See explanation of fits in text. The data are restricted to the negative branch of the cusp in (a).

Eq. (2) for the resonance field B_0 in the range $0° < \theta < 5°$, using γ as an adjustable parameter (we take $s = 15.6$ Å, $\lambda_{ab} = 1400$ Å, and $\varepsilon_r = 25$ [Ref. [5]]). The best fit to B_0 (for negative θ) is obtained with $\gamma = 420 \pm 10$. [See Fig. 2(b). For comparison, we also show curves with $\gamma = 300$ and 500.] With $\gamma = 420$, we calculate $J_0 = 4870$ A/cm^2 and $\nu_J(0) = 163$ GHz (If the fit is repeated with $\lambda_{ab} = 1700$ Å, we obtain the slightly different values $\gamma = 400$ and $\nu_J(0) = 171$ GHz.) Previous estimates of γ in Bi 2212 range from ~55 [13] to 200 [6]. A recent torque measurement at 77 K by Steinmeyer *et al.* [7] obtained $\gamma \sim 900$ in Ar-annealed Bi 2212 ($T_c = 85$ K). We have examined several crystals annealed under different conditions and found that oxygen reduction severely decreases the resonance field (by as much as a factor of 4 in a crystal in which T_c remains above 80 K). In optimally doped Bi 2212, our value for $\gamma \sim 400$ implies that the coupling between adjacent *bilayers* remains very weak. A useful way to express this is by the characteristic field scale $H_0 = \Phi_0/\gamma s^2$ which equals ~2.0 T in Bi 2212. It is significant that H_0 (280–370 T) is over 100 times larger in optimally doped LaSrCuO [14]. Oxygen reduction should decrease this field scale to more accessible values.

We discuss briefly the effect of vortex pinning in our experiment. The scans in Fig. 1 display considerable hysteresis between sweep-up and sweep-down traces when θ is nonzero. However, the hysteresis becomes negligible at $\theta = 0$, suggesting that the pinning force on vortex pancakes is much greater than that on the Josephson vortices. The disparity accounts for two observed features, namely, the asymmetry relative to $\pm\theta$ in Fig. 2(a) and the discontinuous jump observed in B_0 as **H** is rotated out of alignment (at $\theta = 1.5°$ and $2.5°$ in Fig. 1, the peak is observed to stick at 2 T and then jump discontinuously to 5.5 T). At finite, negative θ, a large number of vortex pancakes remain trapped in the sample as H is cycled from 0 to 7 T and back (finite hysteresis). At $\theta = 0$, however, it is possible to expel almost all vortex pancakes by field cycling (negligible hysteresis). Next, as **H** is tilted slightly out of alignment ($\theta > 0$), the sample remains free of vortex pancakes until the normal component $H \sin\theta$ exceeds the "perpendicular critical field" $H_{c1\perp}$ for pancake formation [15]. Thus B_0 sticks at 2 T until the field is large enough for pancakes to reenter the sample. From the jump at 5.5 T in the curve at 1.5°, we estimate $H_{c1\perp} \sim 800$–1300 Oe at 5 K. The jump is not seen for $\theta < 0$ because a finite number of pancakes is always trapped. The lock-in phenomenon has been investigated by magnetization studies in Bi 2212 [7,16], $Tl_2Ba_2Ca_2Cu_3O_{10}$ [7], and (BEDT-TTF)$_2$Cu(SCN)$_2$ [17].

We turn next to the field-perpendicular experiments (**H** ∥ **c**), focusing on the vortex solid phase below the irreversibility line T_m where the measurements are fairly complete. With T fixed, we tracked the resonance frequency ω_J as a function of field B. In the two crystals examined, we find that ω_J varies as a fractional power of the field, viz. $\omega_J^2 \sim B^{-\nu}$ (main panel of Fig. 3) [18]. The exponent ν is not sensitive to T and has the value 0.7–0.8 in both crystals. If ω is held constant, the resonance field B_0 varies roughly exponentially with T (inset, Fig. 3). Thus we may write $\omega_J(B,T)$ as

$$\omega_J^2(B,T) = (B_J/B)^{\nu}\omega_0^2 D(T) \qquad (\mathbf{H} \parallel \mathbf{c},\ T < T_m) \tag{3}$$

(ω_0 and B_J are constants independent of T). Below its cusp at T_m, the factor $D(T)$ is a monotonically rising function of T, well approximated by the form $\exp(T/T_0)$, with the characteristic temperature $T_0 \sim 12$ K [above T_m, $D(T)$ decreases slowly with increasing T [1]].

At present, our understanding of the distinctive behavior of ω_J in perpendicular field is much poorer than for the field-parallel case. In an extension of the model described above, Bulaevskii, Pokrovsky, and Maley (BPM)

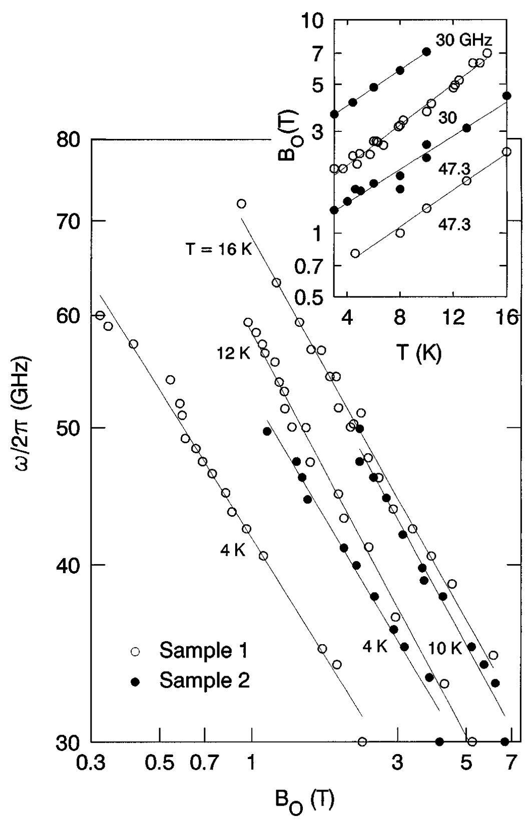

FIG. 3. (Main panel) The variation of the resonance field B_0 (in **H** ∥ **c** geometry) with ω in log-log plot, in samples 1 and 2 with T fixed at the values indicated (sample in a K-band rectangular waveguide). Each data point is the average of four peaks detected in a complete hysteresis cycle (e.g., -2 T → $+2$ → -2 T). The straight lines correspond to $\omega^2 \sim B_0^{-\nu}$. In sample 1, $\nu = 0.70$, 0.82, and 0.78 at 4, 12, and 16 K, respectively. In sample 2, $\nu = 0.72$ and 0.80 at 4 and 10 K. The inset shows semilogarithmic plot of B_0 vs T for the same crystals at the two ω indicated. Straight lines are fits by $B_0 \sim \exp(T/T_1)$ [in $D(T)$, $T_0 \equiv T_1/\nu$]. In sample 1, $T_1 = 9.1$ K (with $\omega = 30$ GHz) and 10 K (47.3 GHz). In sample 2, $T_1 = 10$ and 9 K at 30 and 47.3 GHz, respectively.

[19] propose that, in perpendicular field, disorder in the vortex pancake array is caused by strong pinning (at $T = 0$). BPM show that this can produce the power-law behavior $\omega_J^2 \sim B_0^{-\nu}$. However, this $T = 0$ result cannot be meaningfully confronted with experiment, since the most dramatic effects involve the behavior of $D(T)$, which increases rapidly to a cusp at T_m and then decreases slowly above [1]. Recalling that $\omega_J^2 \sim J_m(B)$, the behavior of $D(T)$ implies that $J_m(B)$ increases rapidly as we approach T_m from below. The experimental finding is that increasing thermal fluctuations counter the disordering effects of the random pins, so that the *vortices become better aligned with increasing* T. This reflects an interesting competition between strong pinning and thermal fluctuations in the lattice that is poorly understood. The Debye-Waller form for $D(T)$ suggests that the characteristic temperature (for $\mathbf{H} \parallel \mathbf{c}$) is $T_0 = 12.5$ K. A second interesting feature is that ω_J does not decrease abruptly to zero (as would be expected if we assume that J_0 is zero in the liquid state). Instead $D(T)$ decreases monotonically [1]. This remains to be explained.

In summary, we have investigated the absorption of microwave radiation in Bi 2212 in the presence of an oblique field. As $\mathbf{H}$ is tilted away from $\mathbf{c}$, B_0 varies rapidly and displays a narrow, reentrant cusp at $\theta = 0$. We remark that the reentrant behavior shown in Figs. 1 and 2 is unexpected in a highly anisotropic but *homogeneous* superconductor. The opposite effects the Josephson and pancake vortices have on the critical current J_m (the latter partially canceling the effect of the former) seem to be essential in producing the behavior. Thus the tilt experiment provides very strong evidence that we are exciting the Josephson plasma oscillation in the presence of Josephson vortices. Using the BMSD model, we have estimated γ (400–420) and $\nu_J(0)$ (160–170 GHz). In contrast to the field-parallel geometry, the behavior of the plasma mode in perpendicular field is poorly understood. Concurrent with our work, Matsuda *et al.* [20] have recently shown that the mode is strongest when $\mathbf{E}_\omega \parallel \mathbf{c}$ and absent when $\mathbf{E}_\omega \perp \mathbf{c}$, consistent with excitation of the Josephson plasma mode.

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