Excitation of the Josephson Plasma Mode in Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ in an Oblique Field

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We present evidence identifying the resonance in microwave absorption observed in Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ 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 $\omega_f$ equal to 5.3 cm$^{-1}$ (at $H = 0$). We also describe the dependence of $\omega_f$ on $H$ and $T$ when the applied field $H \parallel c$.

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When the superconductor Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ (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 ($\omega$) 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 mode. The reentrant behavior of $B_0$ 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 offer estimates for the anisotropy parameter $\gamma$ varying from $>200$ [6] to $\approx 900$ [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 Domínguez (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 $\gamma$ (400–420) and $\nu_J(0)$ ($\approx 5.3$ cm$^{-1}$).

We recall that, in a junction, a fluctuation of the relative phase $\phi$ 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 $\omega_f = [(2\pi/\Phi_0)I_0/C]^{1/2}$, where $I_0$ is the critical current, $\Phi_0$ the flux quantum, and $C$ the junction capacitance. For a layered superconductor in zero field, the $c$-axis critical current density $J_c$ is related to $\gamma$ by $J_c = (\Phi_0/2\pi)/\mu_0s\gamma^2\lambda_{ab}^2$, where $\lambda_{ab}$ is the in-plane penetration depth, $s$ the layer spacing, and $\mu_0$ the permeability constant. In the mixed state, the presence of vortices strongly supresses 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

$$
\omega_f^2(B) = (2\pi/\Phi_0)(sJ_c/\varepsilon_0\varepsilon_r)(\cos \phi_{n,n+1}) \quad (T \to 0)
$$

(1)

($\varepsilon_0$ is the permittivity constant and $\varepsilon_r$ the dielectric constant). The field affects $\omega_f$ 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 $\Delta T$ 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 $\theta$ between $H$ and the CuO$_2$ layers. In all runs, $\omega$, $\theta$, 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
higher fields as we tilt $\mathbf{H}$ 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 $|\theta| \sim 10^\circ$. At $\theta = -5.5^\circ$, $B_0$ attains a local maximum value of 5.6 T, while the resonance broadens noticeably. Remarkably, as $\mathbf{H}$ is brought into closer alignment with the layers, $B_0$ rapidly decreases to a deep minimum at 1.8 T. At positive $\theta$, $B_0$ attains a second maximum (6.2 T) and subsequently approaches the 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 ($\theta = -45.5^\circ$), the trajectory of $B_0$ rises almost vertically. As $\theta$ approaches zero, the trajectory goes through a maximum before falling into a deep cusp. With further increase in $\theta$, the curve goes through another maximum before falling to the right (we discuss the slight asymmetry below). The reentrant behavior of $B_0$ vs $\theta$ 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 $\theta$, 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, $J_s(B)$ increase steeply as $\mathbf{H}$ is tilted out of the plane, in qualitative agreement with the angular dependence of the $c$-axis resistivity [8,9]. We may express $\omega_J^2$ in an oblique field as ($c$ is the velocity of light) [8]

$$\omega_J^2 = 2(c/s)^2(\Phi_0/Bs^2)(e_r \gamma^4)^{-1}(2\pi)^{-2}(\Phi_0/\lambda_{ab}^2 B)$$

$$+ \theta[\ln(\lambda_{ab}^2/\pi s^2) + 2(\pi \lambda_{ab} s B/\Phi_0)^2]^{-1}$$

(T $\rightarrow 0$, $\theta \ll 1$).

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

![Figure 1](image1.png)

**FIG. 1.** Variation of the resonance field $B_0$ (arrows) with tilt angle $\theta$ in Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ (sample 1), measured at 5.0 K at $\omega/2\pi = 47.3$ GHz ($\theta$ is defined in inset). $\mathbf{H}$ is fixed parallel to the cylindrical waveguide axis. The microwave $E$ field $E_y$ lies in the plane of $\mathbf{H}$ and $c$. Each trace represents the surface resistance $R_s$ vs the applied field $H$ (only sweep-up traces shown). As $\theta$ increases from $-5.5^\circ$ and 0, $B_0$ falls to a deep minimum. At small, positive $\theta$ (1.5°, 2.5°), we observe a discontinuous jump of $B_0$ from 2 to 5.5 T. Further increase of $\theta$ to 24° moves the peak back to low fields. At $\theta = 90^\circ$ (not shown), the coupling to the radiation is greatly reduced, but $B_0$ at 0.8 T may still be resolved.

![Figure 2](image2.png)

**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 $\pm 0.3^\circ$. The data were taken in the sequence $-45^\circ \rightarrow 0 \rightarrow 24^\circ$. The line is a guide to the eye. (b) Comparison of the resonance field $B_0$ vs $\theta$ with the prediction of BMSD's model [8] [Eq. (2)]. The lines are curves for $\gamma = 300, 410, 430$, and 500 (with $c = 15.6$ Å, $\lambda_{ab} = 1400$ Å, and $s_r = 25$). 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 $0^\circ < \theta < 5^\circ$, using $\gamma$ 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 $\theta$) 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 $\gamma$ in Bi 2212 range from $\sim 55$ [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 $\sim 2.0$ 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 $\theta$ 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 $\theta = 1.5^\circ$ and $2.5^\circ$ in Fig. 1, the peak is observed to stick to 2 T and then jump discontinuously to 5.5 T). At finite, negative $\theta$, 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 to 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 $1.5^\circ$, 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], Ti$_2$Ba$_2$Cu$_4$O$_{10}$ [7], and (BEDT-TTF)$_2$Cu(SCN)$_2$ [17].

We turn next to the field-perpendicular experiments ($H \parallel c$), 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 $\omega_J$ as a function of field $B$. In the two crystals examined, we find that $\omega_J$ varies as a fractional power of the field, viz. $\omega_J^2 \sim B^{-\nu}$ (main panel of Fig. 3) [18]. The exponent $\nu$ is not sensitive to $T$ and has the value 0.7–0.8 in both crystals. If $\omega$ 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

$$\omega_J^2(B,T) = (B_J/B)^{\nu_0}D(T) \quad (H \parallel c, \; T < T_m)$$

($\omega_0$ 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 $T$ [11]].

At present, our understanding of the distinctive behavior of $\omega_J$ 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)
[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 \( \omega^2 \sim B_0^{-\gamma} \). 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 \( \omega^2 \sim J_m(B) \), 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 \( H \parallel c \)) is \( T_0 = 12.5 \) K. A second interesting feature is that \( \omega^2 \) 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 \( H \) is tilted away from \( c \), \( B_0 \) 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 \( \gamma \) (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 \( E_\theta \parallel c \) and absent when \( E_\theta \perp c \), consistent with excitation of the Josephson plasma mode.

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[14] In \( c \)-axis reflectivity measurements in \( La_{2-x}Sr_xCuO_4 \) (\( x = 0.15 \)) [A.M. Gerrits et al., Phys. Rev. B 51, 12 049 (1995)] did not resolve any shift in the Drude edge in fields up to \( 25 \) T (\( H \perp c \)). The large field scale \( H_0 = 280–370 \) T (estimated with \( \gamma = 15–20 \) and \( s = 6 \) Å) implies the shift may be too small to observe for \( x = 0.15 \). Recent reflectivity measurements on undoped \( LaSrCuO \) have succeeded in detecting the shift in \( H \perp c \) [S. Uchida (private communication)].