# Excitation of Josephson plasma and vortex oscillation modes in $Bi_2Sr_2CaCu_2O_{8+\delta}$ in parallel magnetic fields

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The Josephson plasma resonance of $Bi_2Sr_2CaCu_2O_{8+\delta}$ in parallel magnetic fields has been measured by angular alignment with ~0.02 degree precision to reveal the phase collective modes in the Josephson vortex lattice. When the angle $\theta$ between the field and the $CuO_2$ plane is tilted from the $c$ axis, the resonance field scales as $H \sin\theta$ in the range $\theta \gtrsim 4°$. At lower angles, the resonance field shows anomalous decreases with decreasing $\theta$. A disappearance of the resonance has been unexpectedly found for angles within $\theta \sim 0.2°$, which may correspond to the angle where the lock-in transition takes place. The origin of the phenomenon is discussed in terms of the Josephson plasma oscillation and Josephson vortex oscillation modes. [S0163-1829(97)52214-4]

It is well established that strongly anisotropic high-$T_c$ superconductors (HTSC) behave as stacks of superconductor-insulator-superconductor (S-I-S) Josephson tunnel junctions. This peculiar electronic structure is responsible for the unusual magnetic and electrodynamic properties of these materials in the superconducting state. Recently, Josephson plasma, which is a phase collective mode representing the Cooper pair oscillation through the insulating barriers in the crystal, has attracted much interest.[1–10] Since the plasma modes of HTSC lie well below the superconducting gap, plasma damping processes such as Landau damping and optical phonon damping are prohibited from occurring and a very stable mode appears in the superconducting state.[6] The Josephson plasma excitation mode as a resonance in microwave absorption has been observed recently in $Bi_2Sr_2CaCu_2O_{8+\delta}$ (Refs. 1–3) and $Bi_2Sr_2CuO_{6+\delta}$ (Ref. 4) by applying a magnetic field perpendicular to the conducting layers. Quite recently, similar resonance has been reported in the organic superconductor $\kappa$-(BEDT-TTF)$_2$Cu(NCS)$_2$ with extremely large anisotropy.[5] It has been shown that in these Josephson coupled superconductors the pancake vortices induced by a perpendicular field play an important role in determining the plasma modes, yielding the reduction of the interlayer Josephson coupling and plasma frequency.[7,8,10,11]

In a parallel field, on the other hand, Josephson vortices, which do not have the usual normal core, are known to play an important role in determining the electromagnetic properties of layered superconductors.[12,13] However, the plasma modes in a parallel field have so far never been examined experimentally, not only in the layered superconductors but also in the single S-I-S junction. Recently, Tsui, Ong, and Peterson measured Josephson plasma resonance in an oblique field for $Bi_2Sr_2CaCu_2O_{8+\delta}$ and found that the resonance field displays an unusual reentrant cusp when $H$ is very close to alignment with the layers.[3] This phenomenon has been discussed by Bulaevskii, Maley, Safar, and Domĭnguez (BMSD) in light of the reduction of the plasma frequency due to the formation of Josephson vortices.[9] In this paper we present a very precise angular dependence of the resonance of $Bi_2Sr_2CaCu_2O_{8+\delta}$, particularly in fields very close to the $ab$ plane. We demonstrate that the phenomenon in the parallel field is much more drastic than the result of Ref. 3. When the angle $\theta$ between the field and layers is tilted from the $c$ axis, the resonance scales as $H \sin\theta$ in the wide angle range $\theta \gtrsim 4°$. At lower angles ($\theta \lesssim 0.6°$), however, the resonance field $B_0$ shows an anomalous decrease with decreasing $\theta$. A disappearance of the resonance has been unexpectedly found for angles within $\theta \sim 0.2°$, below which vortex lines may be locked in a parallel orientation. The origin of the phenomenon is discussed in terms of the phase collective modes in the Josephson vortex lattice.[14–16] It is revealed that a collective mode stemming from the vortex lattice oscillation that has not been considered in the previous studies plays an important role in determining the plasma resonance in the parallel fields.

The single crystals of $Bi_2Sr_2CaCu_2O_{8+\delta}$ were grown by the traveling-solvent–floating-zone technique. In this study, we used two crystals with different oxygen stoichiometry: optimally doped ($T_c = 89.5$ K) and underdoped ($T_c = 87$ K) crystals. Both samples were cut and cleaved into square plates with approximate dimensions of $0.9 \times 0.9 \times 0.07$ mm$^3$. Both were obtained by annealing the as-grown crystals

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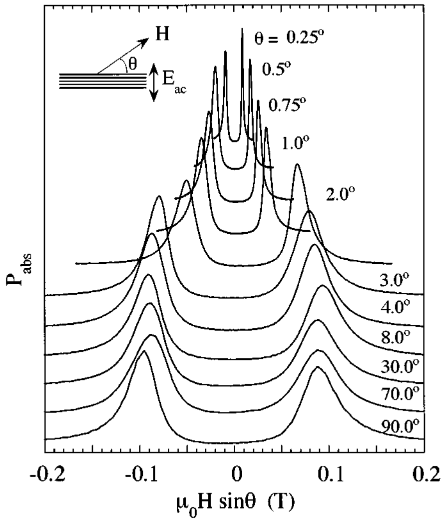


FIG. 1. Resonance in oblique fields for optimally doped crystal measured at $T=30.2$ K at $\omega=45$ GHz. The resonance is plotted as a function of the field component perpendicular to the layers, $H$ sin $\theta$.

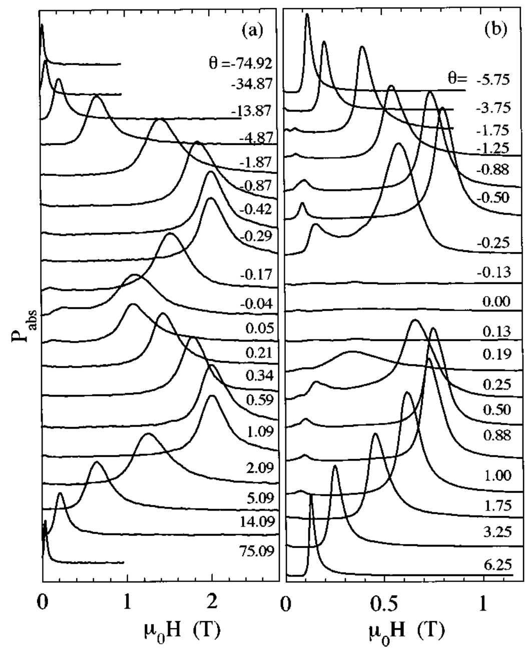


FIG. 2. (a) The angular dependence of the resonance in a magnetic field nearly parallel to the layers for optimally doped crystal at 40.2 K. (b) Same data for an underdoped crystal with a lower resonance field at 36.0 K. In both cases, the angle is rotated from − to +.

in a vacuum or nitrogen atmosphere. The crystal is placed in a rectangular copper cavity with $TE_{102}$ mode ($\omega=45$ GHz) and is placed at the maximum microwave electric field $\mathbf{E}_{ac}$ and oriented such that $\mathbf{E}_{ac}$ is parallel to the $c$ axis.[17] This configuration generates the Josephson plasma mode in which the supercurrent oscillates normal to the layers.[2] The measurements of the angular dependence of the resonance were carried out in the split pair magnet providing a transverse field. The microwave cavity was rotated in the magnet with a precision of 0.02°, keeping $\mathbf{E}_{ac}$ parallel to the $c$ axis.

Figure 1 shows the angular dependence of the resonance for the optimally doped sample taken at $T=30.2$ K. The magnetic fields are swept from $+6$ T to $-6$ T through the zero field. The resonance in positive and negative fields is almost symmetric. In Fig. 1, the resonance is plotted as a function of the field component parallel to the $c$ axis, $H$ sin $\theta$. While the resonance is scaled well as $H$ sin $\theta$ at large angles ($\theta \gtrsim 4°$), deviation from this simple form is apparent at lower angles. Figure 2(a) displays the traces near $\theta=0$ for the same crystal taken at $T=40.2$ K. The resonance is plotted as a function of $H$. Starting at a negative tilt angle (top trace) the resonance field moves rapidly to higher fields, reaches maximum at 0.4°–0.6°, and then decreases rapidly as $\theta$ approaches zero. The resonance behavior in positive and negative angles is almost symmetric. The traces for the underdoped sample are shown in Fig. 2(b).[18] Unexpectedly, the resonance vanishes within the critical angle $\theta_c \sim 0.13°$ for this crystal. The angular dependence of $B_0$ for both samples is plotted in Fig. 3(a). Although the reentrant behavior of $B_0$ vs $\theta$ resembles that of Ref. 3, the present results are quite different in the following respects. First, the resonance broadens noticeably in Ref. 3 near $\theta \sim 0$, but the present resonance does not show such a broadening.[19] Second, and more importantly, *the vanishing of the resonance* has not been observed in previous studies.

We first discuss the resonance at the higher angles. Josephson plasma frequency, $\omega_{pl}=c/\sqrt{\varepsilon_0}\lambda_c$ ($\varepsilon_0$ is the dielectric constant of the crystal and $\lambda_c$ the out-of-plane penetration length) in the presence of the field is written as

$$\omega_{pl}^2(B,T)=\omega_{pl}^2(0,T)\langle\cos\varphi_{n,n+1}\rangle, \tag{1}$$

where $\langle\cos\varphi_{n,n+1}\rangle$ represents the thermal and disorder average of the cosine of the gauge-invariant phase difference between layer $n$ and $n+1$. The field affects $\omega_{pl}$ through the phase factor only. The reduction of the plasmon frequency is caused by the Josephson strings that are created by the deviation from the straight alignment of the pancakes along the $c$ axis.[7,8,10,11] This deviation is caused by thermal fluctuation and pinning. In the perpendicular field, the plasma frequency decreases monotonically with $H$. Consequently, resonant absorption occurs when $\omega_{pl}$ coincides with the microwave frequency $\omega$ as a function of $H$, if the zero-field plasma frequency is larger than $\omega$.[2] According to the recent theory based on the high-temperature expansion, the plasma frequency is expressed as[10]

$$\omega_{pl}^2=\frac{2\pi s\Phi_0 j_J^2}{\varepsilon_0 k_B T H_\perp}\exp\left(-\frac{\pi s^2 H_\parallel^2}{H_\perp \Phi_0}\right), \tag{2}$$

in the vortex liquid phase, provided

$$\frac{E_J\Phi_0}{H_\perp k_B T}\ll 1. \tag{3}$$

Here $j_J$ is the Josephson critical current in the zero field which is related to the anisotropy factor $\gamma$ by $j_J = c\Phi_0/8\pi^2\gamma^2\lambda_{ab}^2 s$ ($\lambda_{ab}$ is the in-plane penetration depth), $H_\parallel$ and $H_\perp$ are the field components parallel and perpendicular to the layers, respectively, and $E_J(=\Phi_0 j_J/2\pi c)$ is the Josephson coupling energy. Equation (3) assumes the almost decoupled layers. The exponential factor in Eq. (2) arises

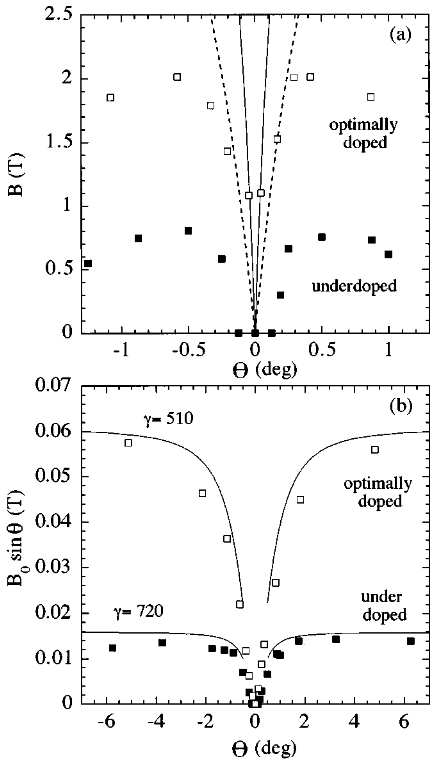


FIG. 3. (a) Angular dependence of the resonance field $B_0$ for optimally doped (open squares) and underdoped (filled squares) crystals in $H$ very close to the layers. Solid and dashed lines represent $B_0$ obtained by Eq. (5) assuming $\gamma$=250 and 300, respectively. (b) Comparison of $B_0$ vs $\theta$ with the prediction of Eq. (2) for optimally doped (open squares) and underdoped (filled squares) crystals. The lines are curves for $\gamma$=510 and 720. See explanation of the fits in the text. The curves are restricted to the region of Eq. (3).

from the extra phase factor added to $\varphi_{n,n+1}$ due to the presence of $H_\parallel$. Since $\Phi_0/\pi s^2 \sim 100$ T, the exponential factor is not important at high angles and resonance fields are well scaled by $H_\perp = H \sin\theta$. Although we do not show it here, the plasma frequencies in both the optimally doped and underdoped crystals in the perpendicular field ($H_\parallel = 0$) are well described by Eq. (2), assuming $\gamma$=510 for optimally doped and 720 for underdoped crystals. (We take $\epsilon_0 = 25$ and $\lambda_{ab} = 2000$ Å.) Zero-field plasma frequencies at $T=0$ are estimated to be 113 and 81 GHz for optimally doped and underdoped crystals, respectively. A detailed discussion of these results will be presented elsewhere.[20] We therefore analyze the angular dependence of $B_0$ in accordance with Eq. (2). Setting $\omega_{\rm pl} = 45$ GHz, we solve Eq. (2) for $B_0$ using these $\gamma$ values. The inequality Eq. (3) holds for $H_\perp > 210$ Oe ($\theta > 0.5°$) in the optimally doped and $H_\perp > 105$ Oe ($\theta > 1°$) in the underdoped sample. The result is depicted by solid lines in Fig. 3(b) as a function of $B_0 \sin\theta$. The deviation from the scaling relation $B_0 \sin\theta$ is appreciable below a few degrees for both samples. The results of the calculation well reproduce the experimental data. Thus the resonance at high angles is found to be well described by Eq. (2).

We next discuss the vanishing of the resonance observed in the underdoped crystal below $\theta_c \sim 0.13°$. This phenomenon cannot be accounted for by the suppression of the interlayer coupling; hence it calls for a new theory. The vanishing admits of two interpretations; a discontinuous jump of the resonance mode to a different frequency or a strong smearing of the resonance. We show below that the jump can be accounted for naturally by considering the electromagnetic excitation mode of the Josephson vortex lattice. We assume that $\theta_c$ corresponds to the lock-in angle below which no pancake segment exists in the vortex line. This transition takes place when $H_\perp$ is less than the lower critical field perpendicular to the layers $H_{c1}^\perp$. The lock-in angle $\theta_L$ is given as

$$H \sin\theta_L = H_{c1}^\perp (1-N), \tag{4}$$

where $N$ is the demagnetization factor.[12] Assuming $H_{c1}^\perp \simeq 100$ Oe and $N \simeq 0.9$, $\theta_L$ at $H=0.5$ T is estimated to be 0.12°. This value is close to $\theta_c$. Thus, a Josephson vortex lattice with no pancake segment would be formed below $\theta_c$. We consider the discontinuous jump of the plasma resonance mode on the analogy of the phase collective modes in the Josephson vortex lattice that have been discussed in the long single Josephson junction.[14] In the Meissner phase, there is only one collective mode corresponding to a plasma wave with a dispersion, $\omega^2(\mathbf{k}) = \omega_{\rm pl}^2 + c_0^2 \mathbf{k}^2$ ($c_0$ is the Swihart velocity). As a result, the uniform ac electric field causes the resonance only at $\omega = \omega_{\rm pl}$. However, a parallel field changes the situation drastically. In the presence of $H_\parallel$ larger than the lower critical field parallel to the junction $H_{c1}^\parallel = 2\Phi_0/\pi^2 d\lambda_J$ ($d$ is the thickness of the junction and $\lambda_J$ is the Josephson length), two different phase collective modes appear that can be referred to as the vortex oscillations and the plasma oscillations. Both modes accompany the charge (Cooper pair) oscillation through the insulating layers. The former mode corresponds to sound-wave-like collective oscillation associated with the sliding mode of the Josephson vortex lattice. Since the Josephson vortex lattice can move freely parallel to the $ab$ plane, this mode has no gap at $\mathbf{k} = 0$. On the other hand, the latter modes are analogous to the optical branch of the phonon spectra and have gaps which are always larger than the zero-field plasma frequency. As a result the zero-field plasma resonance at $\omega = \omega_{\rm pl}$ splits into an infinite number of resonances with frequencies larger than the zero-field plasma frequency.[14] Quite recently, Bulaevskii, Domĭnguez, Maley, and Bishop (BDMB) found that the optical weight of the plasma modes shifts to that of the vortex oscillation mode with the field and finally vanishes in the limit of strong field $H \gg H_{c1}^\parallel$.[16]

The importance of the mode arising from the Josephson vortex lattice oscillation in layered superconductors has also been pointed out by many authors.[15,16] BDMB predicted that the behavior of the plasma modes in a layered superconductor may be similar to those in a single junction, but the characteristic field scale $H_0 = \Phi_0/\gamma s^2$ replaces $H_{c1}^\parallel$ in a single junction.[16] $H_0$ is several T in $Bi_2Sr_2CaCu_2O_{8+\delta}$ with $\gamma$=500. The net optical weight of plasma resonance peaks near the zero-field plasma frequency in low fields $H \ll H_0$, but it shifts to the vortex lattice oscillation mode with the field and vanishes at $H \gg H_0$. Thus the resonance modes of the layered superconductors in parallel fields exist but they are always above the zero-field plasma frequency, similar to those in single junctions. Consequently, *if $\omega$ is less than the zero-field plasma frequency, no resonance is observed*. This is in contrast to the perpendicular and oblique field cases where plasma frequency decreases monotonically with the field. Thus the vanishing of the resonance below the lock-in angle can be accounted for by considering the excitation mode of the Josephson vortex lattice. The vanishing of the resonance has not been observed in the optimally doped crystal [Fig. 2(a)]. The reason for this may be that the high resonance fields for this sample reduce the lock-in angle to an inaccessible low angle range that may be within the mosaic spread inside the crystal.

We finally discuss the angular dependence of the resonance slightly tilted away from the parallel configuration. In this orientation, the vortex lines with very long intervals parallel to the layers separated by pancake vortices are strongly pinned by the pancake segments. Recently, BMSD found that the pancakes adjust to the lattice of Josephson vortices by forming a zigzag structure along the $c$ axis, giving rise to the periodic pinning for Josephson vortices.[9] As a result, the sliding mode of the Josephson vortex lattice becomes an oscillating mode with nonzero frequency. This theory successfully explains $c$-axis resistance that shows a narrow maximum as a function of $\theta$ near parallel orientation with a width of $\sim 1°$.[9] According to BMSD, the plasmon frequency at a very low angle ($|\theta| \ll 1$) may be written as[21]

$$\omega_{\mathrm{pl}}^2 = \frac{8\Phi_0 c^2 \theta}{\pi \gamma^4 s^4 \varepsilon_0 B}[\ln(\lambda_{ab}^2/\pi s^2) + 2(\pi \lambda_{ab} s B/\Phi_0)^2]^{-1}. \qquad (5)$$

Setting $\omega_{\mathrm{pl}} = 45$ GHz, we solve Eq. (5) for the resonance field $B_0$, using $\gamma$ as an adjustable parameter. In Fig. 3(a), we plot the result of the fit for the optimally doped crystal. For comparison, we show curves with $\gamma = 250$ and 300. The range we can fit is very narrow ($\theta < 0.2°$). The $\gamma$ value obtained by this method is about half the value obtained by the resonance in the perpendicular fields.

In summary, we have investigated the Josephson plasma resonance of $Bi_2Sr_2CaCu_2O_{8+\delta}$ in parallel magnetic fields. The angular dependence of the resonance obeys a scaling relation $H \sin\theta$ for $\theta \gtrsim 4°$, but shows significant deviations at lower angles. A disappearance of the resonance has been found within angles where the lock-in transition takes place. We attribute the phenomenon to the phase collective modes of the Josephson vortex lattice.

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[18] In the underdoped sample, a weaker second resonance occurs at fields of about 0.1 to 0.2 T for small angles. A similar but much weaker feature is present in the optimally doped sample. The shape and position of this second resonance depend on the field sweep ratio, while those of the main resonance line do not. Therefore the second resonance lines may be related to the nonequilibrium properties of the vortex state.
[19] We note that the experiment of the angular dependence in Ref. 3 was performed by rotating the sample in a waveguide, not in the cavity resonator. In such a condition, the electric field is not perpendicular to the layers, and both transverse and longitudinal plasma modes are excited simultaneously. The transverse plasma mode is known to be much broader than the longitudinal one. See Ref. 6.
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[21] This result was first obtained in Eq. (9) of Ref. 9, but in the present equation the first term, $J_m(0)$, is dropped. This revised result is given in Ref. 16.