Reduction of the Superfluid Density in the Vortex-Liquid Phase of Bi$_2$Sr$_2$CaCu$_2$O$_y$

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In-plane complex surface impedance of a Bi$_2$Sr$_2$CaCu$_2$O$_y$ single crystal was measured in the mixed state at 40.8 GHz. The surface reactivity, which is proportional to the real part of the effective penetration depth, increased rapidly just above the first-order vortex-lattice melting transition field and the second magnetization peak field. This increase is ascribed to the decrease in the superfluid density rather than the loss of pinning. This result indicates that the vortex melting transition changes the electronic structure as well as the vortex structure.

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The electronic structure of high-$T_c$ superconductors (HTSC’s) in the mixed state attracts much attention. In conventional superconductors (CSC’s) with $s$-wave gap, quasiparticle (QP) states in the mixed state are localized in the vortex cores where the superconducting gap is suppressed. The QP excitation spectrum inside the vortex core consists of quantized energy levels separated by $\Delta E \sim \Delta_0/k_F\xi$, where $\Delta_0$ is the bulk gap, $k_F$ is the Fermi momentum and $\xi$ is the coherence length [1]. Since $\Delta E$ in CSC’s is much smaller than the scattering energies, the vortex core can be regarded as a normal metal. In HTSC’s, above picture can not be applied because of the following reasons. First, the symmetry of the superconducting gap in HTSC’s is not an $s$-wave but is most likely a $d$-wave. Since the amplitude of the $d$-wave gap is zero at the node, QP’s are not localized in the vortex core but extend along the node direction [2,3]. Calculations based on the Bogoliubov-de Gennes equations suggest that there are no truly localized states in the vortex core of pure $d$-wave superconductors [4,5,6]. Secondly, even if there are localized QP states inside the vortex core owing to the mixing of the different gap symmetry [7], $\Delta E$ may exceed any other energy scales, since $k_F\xi$ of HTSC’s is very small. Therefore the vortex core in HTSC’s should be different from that of CSC’s in either cases. Finally, in HTSC’s, the greater part of the mixed state is a vortex liquid phase which is practically lost in CSC’s. It may be possible that the difference in the vortex structure brings the difference in the electronic state.

Experimentally, scanning tunneling spectroscopies [8,9] and thermal conductivity measurements [10] have been used to study the QP state of HTSC’s in the mixed states. High frequency surface impedance $Z_s = R_s + iX_s$ also provides useful information on the electronic state. Using high enough frequencies, one can be free from the vortex pinning and the information not only on QP’s but also on superfluids can be deduced from $R_s$ and $X_s$, respectively. So far, a vortex flow resistivity and a QP lifetime inside the vortex core were discussed from the $R_s$ measurements in YBa$_2$Cu$_3$O$_y$ [11-13]. Recently, Mallozzi et al. measured the complex resistivity in the mixed state of Bi$_2$Sr$_2$CaCu$_2$O$_y$ and argued a possible $d$-wave effect [14]. However, measurements of $Z_s$ near the vortex-melting transition are still lacking. In this paper, we report both $R_s$ and $X_s$ measurements on a Bi$_2$Sr$_2$CaCu$_2$O$_y$ single crystal and show that the high-frequency response can not be explained in terms of simple vortex flow. Moreover, we succeeded in measuring $Z_s$ across the first-order vortex melting transition line and found that the reactive part, $X_s$, increases rapidly above the transition while there was little change in $R_s$. This result indicates that the superfluid density or the amplitude of the order parameter decreases in the vortex liquid phase.

A Bi$_2$Sr$_2$CaCu$_2$O$_y$ single crystal was grown by the floating zone method. The as grown crystal was annealed in air at 800 °C for 3 days and was quenched to room temperature to achieve an optimum oxygen content. A superconducting transition temperature $T_c$ defined at zero resistivity was 91 K. Prior to the surface impedance measurements, we measured the magnetic-field dependence of the local magnetization of the same crystal using the micro Hall probe magnetometry and determined the first-order vortex melting transition field from the position of the magnetization jump. Surface impedance was measured by the cavity perturbation technique with a cylindrical Cu cavity operated at 40.8 GHz in the TE$_{011}$ mode. The sample was located at the center of the cavity which is the anti-node of the microwave magnetic field $H_{rf}$ being parallel to the $c$-axis of the sample. The dimensions of the sample were $0.5 \times 0.5 \times 0.02$ mm$^3$, which is appreciably larger than the normal-state skin depth (~1 $\mu$m at 40.8 GHz). The surface resistance $R_s$ and the surface reactance $X_s$ can be obtained from the changes in the quality factor of the cavity and the resonance frequency, respectively. We determined absolute values of $R_s$ and $X_s$ from comparison with the dc resistivity and assuming that $R_s = R_s$ (Hagen-Rubens limit) in the normal state. Using this procedure, we obtained the reasonable zero-temperature penetration depth of $\sim 2 \times 10^2$ nm. Surface impedance measurements were performed with swept temperature $T$ under field cooled conditions to
avoid any extrinsic effects associated with pinning e.g. giant magnetostriction [4]. In all the measurements, the static magnetic fields $H_{dc}$ were applied along the $c$-axis.

First, we briefly introduce the general behavior of the surface impedance of type-II superconductors [3]. The surface impedance $Z_s$ is related to the complex effective penetration depth $\lambda$ as $Z_s = i\mu_0 \omega \lambda$, where $\mu_0$ is the vacuum permeability and $\omega$ is the angular frequency. The complex resistivity $\rho$ is also expressed by $\lambda$ as $\rho = i\mu_0 \omega \lambda^2$. In the Meissner state, the response is reactive and $\lambda$ is purely real. Therefore, $R_s \approx 0$ and $X_s = \mu_0 \omega \lambda_L$, where $\lambda_L$ is the London depth. In the normal state, the response is dissipative and $\lambda^2$ is purely imaginary. Therefore, $R_s = X_s = \mu_0 \omega \delta/2$ where $\delta = (2\rho_p/\mu_0 \omega)^{1/2}$ is the skin depth and $\rho_p$ is the normal resistivity. In the mixed state, vortex dynamics affects $\lambda$. If the frequency is low, vortices are effectively pinned and a response is similar to that of the Meissner state except that $\rho$ is replaced by $\lambda^2 \sim \lambda_L^2 + B \Phi_0/\mu_0 \kappa_p$, where $B$ is the magnetic induction, $\Phi_0$ is the flux quantum and $\kappa_p$ is the Labusch parameter which denotes the pinning strength. On the other hand, if the frequency is high enough, the viscous loss becomes dominant and a response is similar to that of the normal state except that $\delta$ is replaced by the vortex flow skin depth $\delta_f \sim (2B \Phi_0/\mu_0 \omega \eta)^{1/2}$, where $\eta$ is the viscosity of the vortex motion being related to the QP excitation inside the vortex core. Within the vortex flow theory of Bardeen and Stephen [4], $\eta$ is independent of $B$ and $R_s \approx X_s \propto B^{1/2}$ at high enough fields where $\delta_f \gg \lambda_L$ and $R_s \propto B$, $X_s \sim \mu_0 \omega \lambda_L$ at low fields. A crossover from the pinned regime to the dissipative regime occurs at the pinning frequency $\omega_p = \kappa_p/\eta$.

The field dependence of $Z_s$ of the Bi$_2$Sr$_2$CaCu$_2$O$_y$ single crystal is shown in Fig. 1. At high fields above 0.2 T, $R_s$ is almost proportional to $B^{1/2}$. However, $X_s$ is larger than $R_s$ up to the highest field we used. This result means that the contribution from the reactive parts, superfluid and/or pinned vortices, can not be neglected. Therefore, the observed $B^{1/2}$ dependence of $Z_s$ can not be attributed to the high-field region of the Bardeen-Stephan type vortex flow.

At low fields, a prominent anomaly can be seen. Above a certain field, $X_s$ increases rapidly and saturates at a higher field while there is little change in $R_s$ in the same field region. To examine this anomaly in more detail, we plot the field induced changes in the real part of the effective penetration depth $\text{Re}\Delta\tilde{\lambda}(H_{dc}) \equiv (X_s(H_{dc}) - X_s(0))/\mu_0 \omega$ in Fig. 2. At low fields and low temperatures, $\text{Re}\Delta\lambda$ is independent of $T$ and linear in $H_{dc}$. At high temperatures, $\text{Re}\Delta\lambda$ begins to deviate from the linear $H_{dc}$ dependence and rapidly increases above a certain field $H_{kink}$. In Fig. 3, $H_{kink}$ is plotted on the magnetic phase diagram together with the vortex melting transition field $H_m$ and the second magnetization peak field $B_{sp}$. Here, $H_{kink}$ is defined at the onset of the rapid increase in $\text{Re}\Delta\lambda(T)$ as shown in the inset of Fig. 3. The agreement between $H_m$ and $H_{kink}$ is excellent, indicating that the vortex melting affects $\text{Re}\Delta\lambda$. In the present crystal, the “critical point” is located at 45 K and 33 mT. Above this field, the sharp feature in $\text{Re}\Delta\lambda(T)$ disappears as shown in the inset of Fig. 3. However, as shown in Fig. 2, the small deviation from the linear $H_{dc}$ dependence is still observed in $\text{Re}\Delta\lambda(H_{dc})$ around 40 mT, somewhat higher than $B_{sp}$. Considering the difference between the applied field and the magnetic induction inside the sample and the different processes of the field applications in the local magnetization and the microwave measurements, it is reasonable to regard that the increase in $\text{Re}\Delta\lambda(H_{dc})$ below 45 K and the second magnetization peak have the same origin, most likely the field induced vortex decoupling [7]. To sum up, $\text{Re}\Delta\lambda(H_{dc})$ increases when the 3-dimensional vortex lattice model is no longer valid. Note that this anomaly only appears in the reactive response (Re$\Delta\lambda$ or $X_s$) and almost no anomaly appears in the dissipative response ($R_s$).
Here, we introduce the ratio of these two terms: \[ \frac{\rho_1}{\mu_0 \omega L_L^2} = \frac{s}{1 + s^2} + \frac{r^2 + sr}{(1 + s^2)(1 + r^2)} b, \] (1a)

\[ \frac{\rho_2}{\mu_0 \omega L_L^2} = \frac{1}{1 + s^2} + \frac{r - sr^2}{(1 + s^2)(1 + r^2)} b. \] (1b)

Here, \( s = (2 \lambda_L^2/\delta_{nf}^2) \), \( \delta_{nf} \) is the normal-fluid skin depth, \( r = \omega/\omega_p \) and \( b = B_d \phi_0 / \mu_0 \omega_L \lambda_L^2 \). The parameters \( s \) and \( r \) denote the normal-fluid fraction relative to the superfluid density and the weakness of the pinning, respectively. Note that all the parameters related to the vortex motion \( (r \text{ and } b) \) are included only in the second terms of the right hand sides of Eq. (1). Therefore, if only the vortex motion is considered, magnetic field dependent parts of the complex resistivity \( \Delta \rho_1, \Delta \rho_2 \) are nothing but these terms. Here, we introduce the ratio of these two terms: \[ \Delta \rho_2 / \Delta \rho_1 = (r - sr^2)/(r^2 + sr). \] Let us consider the behavior of this ratio at the vortex melting transition. If the vortices melt, the effective pinning strength must be weakened. Accordingly, \( r \propto 1/\kappa_p \) should increase above the transition. Since \( \Delta \rho_2 / \Delta \rho_1 \) monotonically decreases with increasing \( r \), \( \Delta \rho_2 \) should decrease relative to \( \Delta \rho_1 \) at the transition. This behavior is expected to be irrespective of the nature of pinning, surface or bulk, because any depinning processes cause energy losses which mainly affect \( \rho_1 \). To compare the above arguments with the experimental results, we calculated \( \bar{\rho} \) from \( Z_s \). As shown in Fig. 4, the experimentally obtained \( \Delta \rho_2 \) increases relative to \( \Delta \rho_1 \) at the vortex melting transition and the second magnetization peak field. Therefore, the increase in \( \Delta \rho_1 \) can not be attributed to the loss of pinning. Instead, it should be originated from the increase in \( L_L \) itself. Since \( 1/L_L^2 \) is proportional to the superfluid density, this result strongly suggest that the superfluid density decreases in the vortex liquid phase of \( \text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_y \) due to the additional pair breaking. If all the \( \text{QP’s} \) are localized in the vortex core as in \( s \)-wave superconductors, the electronic state might be insensitive to the change in the vortex structure. Therefore, we speculate that the reduction of the superfluid density at the vortex melting transition is related to the \( d \)-wave superconductivity.

Next we discuss the behavior of \( \text{Re} \Delta \lambda \) in the vortex solid phase. Since \( R_s \sim 0 \) in this regime, we can recognize that vortices are effectively pinned so that \( r \ll 1, \) and \( \lambda \sim (\lambda_L^2 + B \Phi_0 / \mu_0 \delta_{nf})^{1/2} \sim \lambda_L + B \Phi_0 / 2 \lambda_L \mu_0 \kappa_F. \) If \( \lambda_L \) is independent of \( H_{dc} \), \( \kappa_F \) should also be independent of \( H_{dc} \) and \( T \), since \( \text{Re} \Delta \lambda \propto B \) and independent of \( T \) as shown in Fig. 2. Field independent \( \kappa_F \) is only expected when the vortices are individually pinned by strong pinning.
ning centers. This is unreasonable because we observed the sharp first-order magnetization jump in our sample. Therefore, $\lambda_L$, as well as $\kappa_p$, should be $H_{dc}$ dependent even in the vortex solid phase. Let us assume that the field dependence of $\lambda_L$ is dominant, as has been argued by Mallozzi et al. in the vortex liquid phase [13]. The field dependent $\lambda_L$ in the mixed state of $d$-wave superconductors is discussed in terms of the magnetic-field effect on the extended QP state by introducing the Doppler energy shift due to the circulating current around the vortices [13][14][20]. According to these normal theories, the kinetic density, which is proportional to the change in energy shift due to the circulating current around the vortices $\lambda_L$. We calculated this quantity at 10 K and found that $\lambda_L(H_{dc})$ increases linearly with $H_{dc}$ when $H_{dc} < 1 \text{T}$. Even if $H_{dc}$ is given by $f_p(H_{dc}) \sim 1 - (H_{dc}/H_{c2})^{1/2}$ [20] and $H_{c2}$ is considered to be around 100 T. Therefore, both the increase in $\lambda_L$ and the contribution from the vortex motion should be considered in the vortex liquid phase. We stress here again that the behavior of $\Delta \rho_2/\Delta \rho_1$ at the transition can not be explained in terms of the vortex motion alone, even though the two effects contribute to Re$\Delta \lambda$. To separate both contributions, detailed frequency dependence measurements are indispensable. This experiment is now underway.

In conclusion, we have measured the complex surface impedance $Z_s = R_s + iX_s$ of a Bi$_2$Sr$_2$CaCu$_2$O$_y$ single crystal in the mixed state. We succeeded in detecting the change in $Z_s$ at the first-order vortex melting transition and the second magnetization peak field for the first time. Above the transition, $X_s$ which is proportional to the real part of the effective penetration depth increases while there was little change in $R_s$. From the analysis of the complex resistivity deduced from $Z_s$, we showed that the increase in $X_s$ can not be ascribed to the loss of pinning but arises from the reduction of the superfluid density. Namely, the additional pair breaking mechanism may exist in the vortex liquid phase. Our results indicate that not only the phase but also the amplitude of the order parameter take different values in different vortex phases. We speculate that this effect is related to the $d$-wave superconductivity.

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