

Non-quasiparticle microwave absorption in $Bi_2Sr_2CaCu_2O_{8+\delta}$ *

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We show that a non-quasiparticle charge collective mode, in parallel and coincident with the d -wave pair conductivity, leads to a quantitative understanding of microwave surface impedance measurements on superconducting $Bi_2Sr_2CaCu_2O_{8+\delta}$. The analysis suggests an inhomogeneous charge ground state in $Bi_2Sr_2CaCu_2O_{8+\delta}$ and other HTS.

A complete understanding of the microwave response of high T_c superconductors has proved to be as elusive as the mechanism of superconductivity itself. Two principal outstanding issues are:

(i) The basic mechanism of linear microwave absorption is still not understood. While one would have thought that an identification of the order parameter symmetry as $d_{x^2-y^2}$ would have provided quantitative description of the linear microwave properties, this has not turned out to be the case. Calculations of the absorption based upon d -wave (or even mixed $s+d$ symmetry) are actually lower by orders of magnitude (as we show below), even though they can account for the penetration depth measurements [1,2]. A unified picture of the microwave loss of single crystals, thin films and ceramics has not yet been achieved.

(ii) The HTS display a surprisingly high level of nonlinear response which is not properly understood, despite several attempts at modelling it. The measured nonlinear response (which is a major limitation of the use of the cuprate superconductors in microwave applications) is significantly higher than estimates based upon d -wave calculations [3]. Particularly intriguing manifestations of nonlinearity are the Josephson-like response in single crystals [4], the so-called magnetic recovery effect[5](a) and 2nd harmonic generation[5](b).

In this paper we discuss a new analysis of the microwave response in HTS. In the superconducting state, electromagnetic relaxation is shown to occur predominantly via a charge mode even at frequencies in the GHz ranges. This model quan-

titatively describes the available microwave data for $Bi_2Sr_2CaCu_2O_{8+\delta}$ ($Bi : 2212$) and also for several other cuprate HTS.

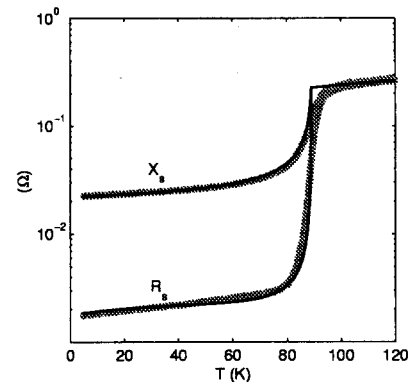


Figure 1. R_s and X_s vs. T for $Bi : 2212$ single crystal (data: thick gray line). The thin solid line represents the calculations using Eq. 1 and 2.

Microwave data for the surface impedance $Z_s = R_s - iX_s$ are presented in Fig.1 (thick gray line). The data are obtained from Nb superconducting cavity resonator measurements, which is operated in TE_{011} mode at $10GHz$, with $H_w \parallel \hat{c}$. The surface impedance Z_s and the conductivity $\tilde{\sigma} = \sigma_1 + i\sigma_2$ are related by $Z_s = \sqrt{-i\mu_0\omega/\tilde{\sigma}}$. $\sigma_1 + i\sigma_2 = -i\mu_0\omega/Z_s^2$ extracted from the Z_s data are shown in Fig.2 (thick gray line). σ_2 rises from zero at and above T_c to a large value at low T . In the conventional picture of microwave superconductivity, σ_2 is regarded as a measure of the superfluid den-

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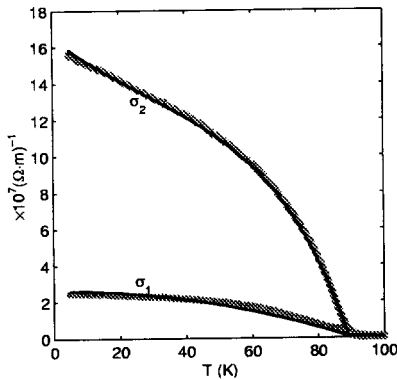


Figure 2. σ_1 and σ_2 vs. T (data: thick gray line). The thin solid line represents the calculations of $\text{Re}[\tilde{\sigma}_{total}]$ and $\text{Im}[\tilde{\sigma}_{total}]$ using Eq. 2.

sity, since in the standard Mattis-Bardeen picture, $\sigma_2(T) = [\mu_0\omega\lambda^2(T)]^{-1} = n_s(T)e^2/\omega m$.

In ref.[1] the T -dependence of σ_2 was shown to be consistent with d -wave calculations [1], assuming a BCS temperature dependence for the $d_{x^2-y^2}$ gap parameter $\Delta(T, \phi) = \Delta_d(T) \cos(2\phi)$. The d -wave model correctly describes the low T behavior of $\lambda(T) \propto T$, a feature which is taken to be the hallmark of d -wave superconductivity [6](a). Similar data and conclusions were subsequently presented in ref.[7] and [8].

Despite the excellent agreement for X_s , and hence $\lambda(T)$ and $\sigma_2(T)$, the understanding of the absorptive part represented by the surface resistance R_s and the normal conductivity σ_1 has remained elusive. This is demonstrated in Fig.3. If the surface resistance is calculated using a d -wave gap and the scattering rate $\Gamma_\sigma(T)$ extrapolated from $T > T_c$ (Fig.3 (a)), then the expected R_s would be $\sim 3 \times 10^2$ times smaller than that measured (Fig.3 (b)). The calculation method is described in ref.[2]. Conversely, in order to explain the large R_s , the scattering rate $\Gamma_\sigma(T)$ would have to drop abruptly at T_c by about the order of 10^3 . This highly unusual behavior would require theoretical explanation which is lacking.

Also the behavior of σ_1 is very anomalous. As is evident from Fig.2 (thick gray line), σ_1 rises

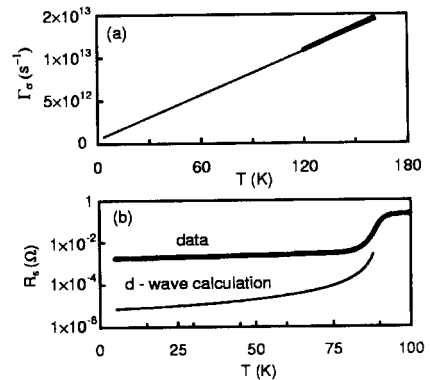


Figure 3. (a) Scattering rate $\Gamma_\sigma(T < T_c) \approx 9 \times 10^{10}T$ (thin line) linearly extrapolated from normal state $\Gamma_\sigma(T > T_c)$ (thick dots). (b) Calculated $R_s(T < T_c)$ using d -wave model and $\Gamma_\sigma(T < T_c)$. The calculated R_s (thin line) is $\sim 3 \times 10^2$ smaller than experimental data (thick dots).

monotonically from very low values, and shows no indication of a downturn. At low T , $\sigma_1(4K)$ is much larger than the residual conductivity $\sigma_{res} \sim 10^{3-4}(\Omega \cdot m)^{-1}$ predicted by quasiparticle localization. Surprisingly, essentially similar behavior of σ_1 in $BSCCO$ thin films is observed at THz frequencies [9]. This is completely different from that theoretically expected for any superconductor, where ultimately one expects that the presence of a gap would lead to freeze out of quasiparticles. Indeed a peak is a prominent feature of σ_1 in other high T_c superconductors such as $YBa_2Cu_3O_{7-\delta}$ – the current explanation for conductivity peaks is that this can be understood from the expression $\sigma_1 = n_{qp}(T)e^2/m\Gamma_\sigma(T)$, as arising from the competition between a decreasing $\Gamma_\sigma(T)$ (increasing $\tau_\sigma(T) = \Gamma_\sigma^{-1}(T)$) and a decreasing $n_{qp}(T)$ (due to quasiparticle freeze-out) [6](b).

These disagreements clearly indicate that an additional mechanism is likely to be responsible for the large microwave loss in this material. We have achieved excellent agreement with a remarkably simple model in which we postulate

that in the superconducting state, in addition to the usual Mattis-Bardeen complex conductivity $\tilde{\sigma} = \sigma_{s1} + i\sigma_{s2}$, a non-quasiparticle polarization contribution also appears at T_c .

In conventional analyses of the microwave response of metals, the displacement current is always ignored. This holds very well in homogeneous metals. Thus Maxwell's equation $\nabla \times \vec{H} = \vec{J} + \partial\vec{D}/\partial t$, is always approximated as $\nabla \times \vec{H} = \vec{J} = \tilde{\sigma}\vec{E}$. This only holds true if $\omega\epsilon' \ll \sigma$. This assumption may not be valid in inhomogeneous metals such as the cuprates. We have recently thoroughly analyzed the cavity measurement technique for the general case of a conducting (or even superconducting) dielectric[10]. Inclusion of the displacement current term then yields a modified equation for the surface impedance as:

$$Z_s = \sqrt{-i\omega\mu_0/\tilde{\sigma}_{total}} \quad (1)$$

where $\tilde{\sigma}_{total} \equiv \tilde{\sigma} - i\omega\tilde{\epsilon}$ is the effective conductivity. For $T < T_c$, $\tilde{\sigma}_{total}$ is represented as:

$$\tilde{\sigma}_{total} = \frac{\sigma_{s1}(T) + i\sigma_{s2}(T) - \frac{i\omega\epsilon_0\epsilon(T)}{1 - \omega^2/\omega_0^2 - i\omega\tau_\epsilon(T)}}{\quad} \quad (2)$$

For $T > T_c$, $\sigma(T > T_c) = \sigma_{10}/t$ where $t = T/T_c$. We use simple forms such as the 2-fluid expressions for the superconducting σ_{s1} and σ_{s2} . We use $\sigma_{s2}(T) = \sigma_{20}n_s(t)$, where $n_s(t)$ is a d -wave superfluid density calculated as discussed in ref.[2] using a gap ratio $\Delta(0)/kT_c = 2.8$, and a monotonic temperature dependence $\sigma_{s1} = \sigma_{10}t^2$ for $\sigma_{s1}(T)$.

For $\tilde{\epsilon}$ we have used a collective mode response like a CDW which turns on at T_c , viz. $\epsilon(T) = \epsilon(0)(1 - t^2)$, where $t = T/T_c$ ($T_c = 89K$). In the limit $\omega \ll \omega_0$, where ω_0 is the oscillator resonant frequency, we get a pinned CDW $\epsilon(T)/(1 - i\omega\tau_\epsilon(T))$, where $\tau_\epsilon(T)$ is the pinning relaxation time. If $\omega \gg \omega_0$, the response is a Drude-like conductivity $\epsilon(T)/(i\omega\tau_\epsilon)/(1 - i\omega\tau_\sigma(T))$, where here $\tau_\sigma = 1/\omega_0^2\tau_\epsilon$. Because we are considering fixed frequency data varying T , either of these assumptions are possible - they only differ in sign for the real part of $\tilde{\epsilon}$. Further frequency dependent measurements are needed to distinguish between these possibilities.

We find that excellent fits to the data are obtained in either limit. The parameters used for the fit in the limit $\omega \ll \omega_0$ are: $\sigma_{20} = 2.3 \times 10^8(\Omega \cdot m)^{-1}$, $\sigma_{10} = 7.7 \times 10^5(\Omega \cdot m)^{-1}$, $\epsilon(0) = 1.4 \times 10^8$, and a temperature independent scattering time $\tau_\epsilon = 6.2 \times 10^{-12}$ sec. Thus in this case $\omega\tau_\epsilon = 0.39$. In the $\omega \gg \omega_0$ Drude conductivity limit, the corresponding parameters are: $\sigma_{20} = 1.6 \times 10^8(\Omega \cdot m)^{-1}$, $\sigma_{10} = 7.7 \times 10^5(\Omega \cdot m)^{-1}$, $\epsilon(0)/\omega\tau_\epsilon = 4.8 \times 10^7$, and a temperature independent scattering time $\tau_\sigma = 2.2 \times 10^{-13}$ sec. Clearly the model describes the data extremely well (see Fig.1 and 2 solid thin lines). The model shows that although the reactive response is dominated by the pair conductivity σ_2 , the absorptive part is completely determined by the non-quasiparticle channel and not by quasiparticles.

A key outcome is that this approach yields an alternative explanation of the "conductivity" peaks mentioned previously. We now identify these peaks as relaxation loss peaks which occur at a peak temperature T_p where $\omega\tau_\epsilon(T_p) = 1$. The peak corresponds to a crossover from $\omega\tau_\epsilon \ll 1$ at high T to a regime $\omega\tau_\epsilon \gg 1$ at low T . We have observed such peaks in non-superconducting oxides, such as $Sr_{14}Cu_{24}O_{41}$ and insulating $YBa_2Cu_3O_{6.0}$ [10,17]. In $Bi : 2212$ the peak is not observed because $\omega\tau_\epsilon(T) < 1$ at all temperatures. In contrast in almost all the other HTS, the scattering rate $\Gamma_\epsilon(T) = \tau_\epsilon^{-1}(T)$ is significantly smaller or $\tau_\epsilon(T)$ is larger, and the relaxation rate crosses the probe frequency even at microwave frequencies, leading to a peak, as seen in superconducting $YBa_2Cu_3O_{6.95}$ [2].

The dielectric strength $\epsilon(0) \sim 10^8$ is typical of values in low dimensional CDW systems [11]. The presence of such appreciable polarization may be indicative of spatial modulation of charge. This is in clear contrast with LTS, and may well be associated with the presence of stripes in HTS [12]. Note that the present results for $H_\omega \parallel \hat{c}$ are distinct from the c-axis polarization measured by ref.[13]. It should also be noted that large dielectric constants both in-plane and along the c-axis have been measured in the parent compound $Bi_2Sr_2(Dy, Er)Cu_2O_8$ [14].

The onset of strong polarization at or near T_c may not be surprising in view of the many re-

ports of structural or lattice distortions reported at T_c in $YBCO$, Hg : 1201 and Tl : 2212 [15]. It is quite reasonable that these structural distortions are accompanied by charge density instabilities resulting in large changes in polarization. It has long been recognized that the oxide superconductors are also incipient ferroelectrics, since ferroelectricity in perovskite oxides is well known. Thus it is possible that the results reported here are observing strong polarization modes associated with the $Bi - O$ and $(Ca, Sr) - O$ layers. Theories where dielectricity and superconductivity are both present have already been presented [16]. In $YBa_2Cu_3O_{6.02}$ we have found evidence for dielectric transitions at 110K and 55K [17].

There are several important implications of the work presented here. The displacement current channel, which is always neglected in analysis of microwave response of conventional metals and superconductors, and in previous analysis of HTS, cannot be ignored in HTS. Consequently, a charge collective mode is essential to understand the microwave response of the cuprate superconductors. The charge mode dynamics manifests itself via the presence of relaxation loss peaks, which are the correct explanation of the microwave absorption peaks even in the superconducting state.

The Mattis-Bardeen conductivity, so successful in describing the electrodynamics of low T_c superconductors, is clearly inadequate for HTS. Quasiparticle contributions alone are clearly insufficient to describe the microwave absorption, and do not account for the most prominent feature of the microwave response of HTS – the “conductivity” peaks which we instead show here are more appropriately called microwave absorption peaks or relaxation loss peaks.

The analysis presented here is not restricted to Bi : 2212. Indeed, for the first time, by including non-quasiparticle contributions, we are able to quantitatively describe the microwave response of other cuprate superconductors also [17].

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