Intrinsic Tunneling Spectra of Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ near Optimal Doping

S. P. Zhao, X. B. Zhu, and Y. F. Wei
Beijing National Laboratory for Condensed Matter Physics,
Institute of Physics, Chinese Academy of Sciences, Beijing 100080, China

We report tunneling spectra of near optimally doped Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ intrinsic Josephson junctions with area of 0.09 $\mu$m$^2$, which avoid some fundamental difficulties in the previous tunneling experiments and allow a stable temperature-dependent measurement. A $d$-wave Eliashberg analysis shows that the spectrum at 4.2 K can be well fitted by considering electron couplings to a bosonic magnetic resonance mode and a broad high-energy continuum. Above $T_c$, the spectra show a clear pseudogap that persists up to 230 K, and a crossover can be seen indicating two different pseudogap phases existing above $T_c$. The intrinsic electron tunneling nature is discussed in the analysis.

PACS numbers: 74.50.+r, 74.25.Jb, 74.72.Hs

The pairing mechanism and pseudogap phenomena are known as the two key issues in the studies of high-$T_c$ superconductors. In addition to the views based on resonating valence bond (RVB) state [1], one possible mechanism that is of much interest is boson-mediated pairing of the electrons. Recent angle-resolved photoemission spectroscopy experiments have focused on a "kink" in single-particle dispersion and both phonon [2, 3] and magnetic resonance [4, 5] are considered as the bosonic candidate. In the tunneling [6] and optical [7] experiments, a broad high-energy continuum, in addition to the resonance mode, has also been considered.

Tunneling has traditionally been a useful tool in revealing the material’s superconducting properties. For the high-$T_c$ superconductors, mostly Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ (Bi-2212), the scanning tunneling microscope (STM) [8], break junctions (BJs) [9], and intrinsic Josephson junctions (IJJs) [10, 11, 12, 13, 14, 15] are used, yielding many important informations. In particular, Zasadzinski et al. have correlated the tunneling dip structure with the magnetic resonance mode [8]. In Ref. 8, the pseudogap opening temperature $T^*$, rather than the superconducting transition temperature $T_c$, is suggested to be the mean-field critical temperature, thus supporting precursor pairing as the origin of the pseudogap phase.

IJJs, being intrinsic, have the advantage of avoiding the problems like the surface chemical deterioration and unstable junction structure, and can offer a convenient temperature-dependent measurement. In the previous studies, however, the micro-meter-sized IJJs are found to suffer from significant self-heating, which severely distorts the tunneling spectra [13, 14].

Recently, we have demonstrated that self-heating in IJJs can be reduced considerably when junction sizes decrease down to the sub-micron level [14]. In this Letter, we report and present a detailed analysis on the tunneling spectra of an IJJ mesa with further-reduced area $S = 0.3 \times 0.3 \mu$m$^2$ on a near optimally doped Bi-2212 crystal [16], which are shown in Fig. 1. Comparison with the previous reports (including those from STM and BJ experiments) is presented in Table I in two important spectroscopic aspects: the relation between the superconducting gap (SG) and pseudogap (PG), and whether there is a clear dip at low temperatures. It can be seen that self-heating can lead to considerably different results in the IJJ experiments.

The spectra in Fig. 1 show a well-defined, almost temperature-independent normal-state resistance $R_N$, which allow a straightforward tunneling analysis. Below we analyze the data, along with a discussion on the tunneling nature in IJJs, of which a consensus is still lacking. We will show that the spectrum at 4.2 K can be explained within the $d$-wave Eliashberg description by considering...
TABLE I: Some differences and similarities of the tunneling spectra for (near) optimally doped Bi-2212. Results for STM, BJ, and IJJ are from Refs. [8] [19] and [11], respectively.

|       | STM | BJ   | IJJ | This work |
|-------|-----|------|-----|-----------|
| SG vs. PG | yes | yes | no  | yes       |
| Clear dip? | yes | yes | no  | yes       |

aSG→PG and PG denote a smooth evolution from SG to PG at Tc, and that they coexist independently, respectively.

Discussion on the tunneling nature in IJJs. — A usual way to fit the tunneling spectra is to build an angle-averaged DOS: \( N_d(\omega) = (1/2\pi) \int \frac{1}{\sqrt{\omega^2 - \Delta^2(\omega)\cos^2(2\theta)}} d\omega \), from which the \( I(V) \) curve can be calculated straightforwardly considering an incoherent tunneling process \([6, 20, 21]\). The basic features from such calculation at 4.2 K using a set of parameters described below can be seen in Fig. 2(a) and (b) (dotted lines). While the lines follow the experimental data closely above the gap voltage, they deviate considerably below it. From our calculation, we find that a better fit below the gap voltage can be obtained by considering coherent electron tunneling with a directional matrix element of \( t \sim \cos^2(2\theta) \), but this leads to a worse agreement near and above the gap voltage \([20]\).

Coherent electron tunneling has been discussed in a number of publications \([20, 22, 23]\) and used to interpret the c-axis transport data in IJJs \([24, 25]\). The present results may imply that there is a crossover from coherent to incoherent tunneling as the bias voltage increases. Difficulty arises, however, that when a \( t \sim \sin^2(\theta) \) is used, which is theoretically more strict for the Bi-2212 material in which the Cu atoms in CuO2 planes across the BiO and SrO barrier layers do not lie collinearly \([23]\), the computed results give a worse fit for voltages both below and above the gap edge. To avoid the conceptual difficulty, we assume that intrinsic electron tunneling in IJJs is basically incoherent for voltages away from zero. At low bias, tunneling can be described by an effective DOS: \( N_{eff}(\omega) = (1/\pi) \int \frac{1}{\sqrt{\omega^2 - \Delta^2(\omega)\cos^2(2\theta)}} d\omega \), which produces similar I-V and \( dI/dV \) characteristics (dashed

FIG. 2: (color online). Fit of the experimental (a) I-V and (b) \( dI/dV \) curves at 4.2 K using the d-wave Eliashberg theory. The calculated DOS and self-energy \( \Sigma \) are shown in (c) and (d). Note the horizontal shifts of 150 and 300 meV for \( N_{eff} \) and \( N(\omega, 0) \) in (c) for clarity. \( \alpha^2F \) used in the calculation is also shown in (d). See text for detailed fitting parameters.
lines in Fig. 2(a) and (b)) as coherent tunneling plus \( t \sim \cos^2(2\theta) \) is considered. Physically, the \( \cos^2(2\theta) \) factor in \( N_{\text{eff}} \) can result from the wave-function overlap (having \( d_{x^2-y^2} \) symmetry) between the O-2p and Cu-4s orbitals while the latter plays an important role in assisting a hole hopping to and out of the O-2p orbital \[ 22 \]. At high bias, tunneling is described by \( N_d(\omega) \).

Switch from \( N_{\text{eff}} \) to \( N_d \) implies a lesser role the Cu-4s orbital plays in assisting hopping as bias increases and more high-energy quasiparticles with reduced lifetimes are involved in tunneling. Phenomenologically, the tunneling current can be written as \( I = IS_d + (1 - P)I_{\text{eff}} \), where \( I_d \) and \( I_{\text{eff}} \) represent those obtained using \( N_d \) and \( N_{\text{eff}} \), and \( P \) is a voltage-dependent probability distribution. In our fitting, an exponential function \( P(V) = (e^{V/V_1} - 1)/(e^{V_2/V_1} - 1) \) with \( 0 \leq V \leq V_2 \) is used, in which \( V_1 \) characterizes the voltage scale for \( I_{\text{eff}} \) to \( I_d \) crossover and \( V_2 \) sets a point above which \( I = I_d \). The solid lines in Fig. 2(a) and (b) are obtained from the dotted and dashed results with \( V_1 \) and \( V_2 \) equal to 20 and 75 mV, respectively. The fit appears satisfactory.

**Tunneling spectrum at 4.2 K.** — In the Eliashberg analysis of the 4.2-K data in Fig. 2, we follow Zasadzinski \textit{et al.} and use two Lorentzians for \( \alpha^2F \) to model the magnetic resonance peak (around 35–40 meV) and the high-energy continuum \[ 6 \]. The fitting parameters are: \( c_3 = 0.18, c_2 = 1; \alpha_1^2 = 20, \Omega_0 = \Omega_{11} = 35 \text{ meV}, \eta_1 = 5 \text{ meV}; \alpha_2^2 = 46.5, \Omega_{02} = \Omega_{12} = 150 \text{ meV}, \eta_2 = 2000 \text{ meV}; R_N = 1.92 \text{ k}\Omega \). The parameters are adjusted so as to have a fair fit of the voltage positions of the \( dI/dV \) peak and dip and also the depth of the dip, which are therefore judged primarily by \( I_d \). The resulting antinodal DOS \( N(\omega, 0) \), together with \( N_d \) and \( N_{\text{eff}} \), is shown in Fig. 2(c), and \( \alpha^2F \) is depicted in (d).

\( \alpha^2F \) is found to change the shape of the \( dI/dV \) curve in the following way. While the resonance peak creates the dip, the high-energy continuum tends to reduce its depth. In addition, as the continuum extends to higher energy, the width of the dip decreases. The latter change is more obvious when the upper cutoff of the continuum reduces from 150 to 80 meV, where the resulting dip width in both the \( dI/dV \) and DOS curves becomes too wide to fit the experimental data. This means that a continuum extending above 150 meV must be considered.

We have chosen an upper cutoff of 300 meV for the continuum following a recent analysis of the optical data \[ 17 \]. As can be seen in Fig. 2(d), this results in a considerable high-energy contribution to \( -2\text{Re}\Sigma \), which may need further experimental testing. Using a smaller cutoff of \( \sim 160 \text{ meV} \) \[ 6 \] will reduce such contribution and slightly increase the dip width. The overall size of \( -2\text{Re}\Sigma \), on the other hand, is mostly affected by \( c_5 \). Reducing \( c_5 \) will reduce \( -2\text{Re}\Sigma \) and make the peak-dip separation smaller.

Compared with the BJ results \[ 6 \], the present spectra show a shallower dip, which leads to a larger contribution of the continuum in \( \alpha^2F \) relative to the resonance mode. Presently, it seems unclear as to which part in \( \alpha^2F \) should be more essential for the pairing \[ 6, 7 \].

**Spectra near and above \( T_c \).** — As temperature increases, the dip gradually disappears at \( \sim 75 \text{ K} \) (see Fig. 1(b) and (c)). In this case, it is possible to fit the data using \( N_{\text{eff}} \) in which \( N(\omega, 0) \) has a Dynes’ form \( \text{Re}[\omega - i\Gamma]/\sqrt{\omega - i\Gamma}^2 - \Delta^2\cos^2(2\theta) \), where \( \Delta \) and \( \Gamma \) are the energy-independent parameters \[ 20 \]. The fitting can be satisfactory using \( N_{\text{eff}} \) in the entire voltage range since difference of using \( N_{\text{eff}} \) and \( N_d \) above the gap voltage is smaller at higher temperatures and in the absence of the dip. Comparison between experiment and theory can be seen in Fig. 3 and the \( \Delta \) and \( \Gamma \) parameters are plotted in Fig. 4. It is interesting to note the linear temperature dependence of the scattering rate \( \Gamma \), a signature predicted in the marginal Fermi-liquid theory \[ 27 \].

It should be pointed out that using \( \Delta \) to describe the data above \( T_c \) implies that we are taking the pre-pairing view and \( \Delta \) is understood as the thermodynamic-averaged value \[ 28 \]. The pre-pairing idea has been suggested previously in the STM experiments to account for the pseudogap phase \[ 3 \]. From our experiments, the \( dI/dV \) spectra show a gapped structure for temperatures as high as 230 K. To look at this in more detail, we plot in Fig. 4 the temperature-dependence of half the peak energy in the \( dI/dV \) curves (\( E_P/2 \) solid squares). It can be seen that \( E_P/2 \) decreases as the temperature increases up to \( T_c \). It then increases, reaches a maximum at \( \sim 140 \text{ K} \), and starts decreasing again after that. The open squares in the figure are the results taken from the \( dI/dV \) curves that are normalized to the 250-K curve. In this case, as temperature decreases from 230 K, \( E_P/2 \) is
almost flat until \( \sim 170 \) K, below which it decreases at a relatively fast speed.

These results point to a change within the pseudogap phase, which can be characterized by two temperatures \( T_1^* = 230 \) K and \( T_2^* = 140\sim 170 \) K. Considering our near optimally doped sample, this is consistent in the phase diagram with the recent Nernst experiment \[29\]. From the RVB scenario \[1\], the results can be viewed as the opening of a spin gap below \( T_1^* \). Below \( T_2^* \), holons start to Bose-condense and the Nernst state results. A global superconducting state is finally developed below \( T_c \) through the Berezinskii-Kosterlitz-Thouless transition. In the view of boson-mediated pairing, the change around \( T_2^* \) should correspond to the point below which the superconducting order starts to form out of (with \( \Gamma \) comparable to \( \Delta \) at initial stage), and compete with the spin-gap state.

**Summary.** — Tunneling spectra of near optimally doped Bi-2212 IJJs in the temperature range of 4.2 to 250 K are presented. We have analyzed intrinsic tunneling process and used \( d \)-wave Eliashberg theory and Dynes’ DOS to fit the tunneling data. The fit between theory and experiment is satisfactory. We have shown that the electronic coupling to a magnetic resonance mode plus a broad continuum can well explain the spectrum at 4.2 K, thus providing an evidence for the electronic-originated boson-mediated pairing. However, a doubt remains as to which part in the boson spectrum is more essential for superconductivity. We have also found that there exist two pseudogap phases above \( T_c \).

We acknowledge many valuable discussions with T. Xiang, N. L. Wang, and Q. H. Wang. This work was supported by the Ministry of Science and Technology of China (2006CB601007), the Knowledge Innovation Project of the Chinese Academy of Sciences, and the National Natural Science Foundation of China (10604064).

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