In-plane optical spectral weight transfer in optimally doped Bi$_2$Sr$_2$Ca$_2$Cu$_3$O$_{10}$

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We examine the redistribution of the in-plane optical spectral weight in the normal and superconducting state in trilayer Bi$_2$Sr$_2$Ca$_2$Cu$_3$O$_{10}$ (Bi2223) near optimal doping ($T_c=110$ K) on a single crystal via infrared reflectivity and spectroscopic ellipsometry. We report the temperature dependence of the low-frequency integrated spectral weight $W(\Omega_c)$ for different values of the cutoff energy $\Omega_c$. Two different model-independent analyses consistently show that for $\Omega_c=1$ eV, which is below the charge transfer gap, $W(\Omega_c)$ increases below $T_c$, implying the lowering of the kinetic energy of the holes. This is opposite to the BCS scenario, but it follows the same trend observed in the bilayer compound Bi$_2$Sr$_2$CaCu$_2$O$_8$ (Bi2212). The size of this effect is larger in Bi2223 than in Bi2212, approximately scaling with the critical temperature. In the normal state, the temperature dependence of $W(\Omega_c)$ is close to $T^2$ up to 300 K.

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I. INTRODUCTION

The discussion as to what extent the superconductivity in high-$T_c$ cuprates can be understood as a BCS-like pairing instability of a standard Fermi liquid is open for almost two decades. The dominating role of the Coulomb interaction in the formation of the non-Fermi-liquid electronic state in the underdoped region of the phase diagram questions the ability of the BCS theory, where these interactions are treated only as a perturbation, to universally explain the mechanism of the formation of the non-Fermi-liquid electronic state in the theoretical implications related to the choice of the cutoff frequency. We present the temperature dependence of $W(\Omega_c)$ for different relevant values of $\Omega_c$.

While the physical meaning of the temperature dependent $W(\Omega_c)$ is a matter of theoretical interpretations, it can be, in fact, experimentally determined without model assumptions. The latter is nontrivial, since the integration in Eq. (1) requires, at first glance, the knowledge of $\sigma_1(\omega)$ down to zero frequency and a separate determination of the superfluid density.

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determine accurately $W(\Omega_c)$ and its temperature dependence without low-frequency data extrapolations. Although this point was emphasized in our previous publications, we reiterate in this article the details of the experimental determination of $W(\Omega_c)$ due to its topmost experimental importance. Two different numerical approaches were used which gave us consistent results: (i) the calculation of $W(\Omega_c)$ at each temperature and (ii) the temperature-modulation analysis of superconductivity induced changes of the optical properties.

II. EXPERIMENT AND RESULTS

Two large single crystals of Bi$_2$Sr$_2$Ca$_2$Cu$_3$O$_{10}$ with $T_c = 110$ K and transition width $\Delta T_c \sim 1$ K were prepared as described in Ref. 19. The samples had dimensions $(a \times b \times c)$ of $4 \times 1.5 \times 0.2$ mm$^3$ and of $3 \times 0.8 \times 3$ mm$^3$, respectively. The first crystal has been used to measure the in-plane optical properties and was cleaved within minutes before being inserted into the cryostat. We measured the normal-incidence reflectivity from 100 and 7000 cm$^{-1}$ (12.5 meV–0.87 eV) using a Fourier transform spectrometer, evaporating gold in situ on the crystal surface as a reference. The reflectivity curves for selected temperatures are displayed in Fig. 1. The ellipsometric measurements were made on the same sample surface in the frequency range between 6000 and 36 000 cm$^{-1}$ (0.75–4.5 eV) at an angle of incidence of 74°. The ellipsometrically measured pseudodielectric function was numerically corrected for the admixture of the $c$-axis component which provided the true $ab$-plane dielectric function, whose real and imaginary parts are shown in Fig. 2. The $c$-axis dielectric function that is required for this correction was measured independently on the $ac$-oriented surface of the second crystal as described in the Appendix. It is shown in Fig. 3.

The superconductivity induced changes of the optical properties at photon energies above the superconducting gap (in the midinfrared and higher frequencies), are rather small but their reliable detection is crucial to determine to correct sign and magnitude of the spectral weight transfer. We used homemade optical cryostats, whose special design preserves the sample alignment during thermal cycling. In the visible-ultraviolet (UV) region, in order to avoid spurious temperature dependencies of the optical constants due to adsorbed gases at the sample surface, an ultra-high-vacuum cryostat was used, operating at a pressure in the $10^{-10}$ mbar range. All data were acquired in the regime of continuous temperature scans at a rate of about 1 K/min between 20 and 300 K with

![FIG. 1. In-plane reflectivity spectra of optimally doped Bi2223 for selected temperatures.](image1)

![FIG. 2. Bi2223 ab-plane dielectric function at selected temperature.](image2)

![FIG. 3. c-axis optical spectra of Bi2223 at selected temperatures. Top panel: normal incidence reflectivity, middle panel: $\sigma_1(\omega)$, bottom panel: $\varepsilon_1(\omega)$.](image3)
TABLE I. Fit of the measured reflectivity and ellipsometry data with one Drude and four Lorentz oscillators: $\varepsilon(\omega) = \varepsilon_{\infty} + \sum_i \frac{\varepsilon_i \omega_i^2}{\omega^2 - \omega_i^2 - i\gamma_i \omega}$ at selected temperatures. All parameters, except $\varepsilon_{\infty}$, are given in cm$^{-1}$, i.e., they should be multiplied by $2\pi c$ to convert to angular frequencies in units of s$^{-1}$.

<table>
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<th>Temperature</th>
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<th>110 K</th>
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<td>1179</td>
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A resolution of 1 K. The signal to noise ratio of the temperature dependent reflectivity in the midinfrared is about 2000.

In order to obtain the optical conductivity $\sigma(\omega)$ in the whole frequency range, we used a variational procedure described in Ref. 20. In the first stage, the infrared reflectivity and the ellipsometrically measured complex dielectric function in the visible and UV range were fitted simultaneously with a Drude-Lorentz model. The corresponding parameters at selected temperatures are listed in Table I. We found that one Drude and four Lorentz terms form a minimal set of oscillators fitting data well enough at all temperatures. The Drude peak narrows with cooling down and transforms to a condensed $\delta$-peak below $T_c$. The two lowest frequency oscillators which mostly describe the mid infrared absorption, show dramatic changes below $T_c$, mimicking the formation of the superconducting gap. The high-frequency Lorentzians corresponding to the interband transitions, show very little temperature dependence. In the second stage, the fitting was refined with a variational dielectric function added on top of the Drude-Lorentz model. The former is essentially a set of a large number of narrow oscillators, each corresponding to one or two spectral data points. This yields the Kramers-Kronig consistent dielectric function which reproduces all the fine details of the infrared reflectivity spectra while simultaneously fitting to the measured complex dielectric function at high frequencies. In contrast to the “conventional” KK reflectivity transformation this procedure anchors the phase of the complex reflectivity to the one at high energies measured directly with ellipsometry.21

FIG. 4. In-plane optical conductivity of Bi2223 at selected temperatures. The inset displays the low energy part of the spectrum.

In Fig. 4 we show the optical conductivity at selected temperatures. The spectral and temperature dependence of $\sigma(\omega)$ of Bi2223 is very similar to the one of Bi2212,5,22 although the conductivity of Bi2223 is slightly larger, likely due to a higher volume density of the CuO2 planes in the trilayer compound. The strongest changes as a function of temperature occur at low frequencies. In the normal state the dominant trend is the narrowing of the Drude peak. The onset of superconductivity is marked by the opening of the superconducting gap which suppresses $\sigma(\omega)$ below about 120–140 meV, slightly higher than in Bi2212. Such a large scale is apparently caused by a large gap value in Bi2223, which amounts up to 60 meV, as shown by tunnelling measurements.23

The much smaller absolute conductivity changes at higher energies, which are not discernible at this scale, can be better seen in Fig. 5 where we show the temperature dependent optical constants taken at selected photon energies. The change induced by superconductivity in the optical constant is clearly visible as a kink at $T_c$ for energies up to at least 2 eV, which tells that the energy range where the redistribution of spectral weight takes place is very large.

III. INTEGRATED SPECTRAL WEIGHT

A. Experimental determination of $W(\Omega_c)$

The extraction of the spectral weight $W(\Omega_c)$ from the measured spectra is a delicate issue. Formally, one has to integrate the optical conductivity over a broad frequency range, including the region below the low-frequency experimental cutoff $\Omega_c$ (in our case about 100 cm$^{-1}$) containing the condensed $\delta$ peak (below $T_c$) at $\omega$=0 and a narrow quasiparticle peak. According to a frequently occurring misconception the existence of such a cutoff inhibits the calculation of this integral. Indeed if only the real part of the optical conductivity in some finite frequency interval was available, clearly an essential piece of information needed to calculate $W(\Omega_c)$ would be missing, namely, $\sigma(\omega)$ below $\Omega_c$. However, due to the fact that the real and imaginary part of the dielectric constant are related nonlocally via the Kramers-Kronig transformation, any change in one of them will affect the other in a broad region of the spectrum. In particular, any
The leading contribution of $\sigma_1(\omega)$ below $\Omega_{\text{min}}$ must influence $\epsilon_1(\omega)$ at higher frequencies. Since the latter is measured independently (directly by the ellipsometry above 0.75 eV cm$^{-1}$ and indirectly via the reflectivity in the infrared), it puts constraints on the possible values of $\sigma_1(\omega)$ below $\Omega_{\text{min}}$ and $W(\Omega_c)$. Obviously, these constraints are going to be the more tight the more accurately the optical constants are determined in the accessible interval.

According to the KK relation

$$\epsilon_1(\omega) = 1 + 8\rho \int_0^{\Omega} \frac{\sigma_1(\omega')d\omega'}{\omega'^2 - \omega^2}$$

the real part of the dielectric function can be done via the KK transform.

In the superconducting state, the model contains a narrow quasiparticle peak (provided that its width $\gamma \ll \Omega_{\text{min}}$ to $\epsilon_1$ (and thus to the reflectivity) at high frequencies are almost indistinguishable. Therefore the value of the integral of $W(\Omega_c)$ can be model independently determined from our experimental data, while resolving the details of $\sigma_1(\omega)$ below 100 cm$^{-1}$, for example the separation of the superfluid density and quasiparticle spectral weight, is not possible.

In practice, the realization of the extra bounds on the value of $W(\Omega_c)$ using the additional information contained in the real part of the dielectric function can be done via the aforementioned procedure of variational KK constrained fitting. Essentially, this is a modeling of the data with a very large number of narrow oscillators, which are added to the model dielectric function until all the fine details of the measured spectra are reproduced. Importantly, the model function always satisfies the KK relations. Once a satisfactory fit of both reflectivity in the infrared region and $\epsilon_1(\omega)$ and $\epsilon_2(\omega)$ above 0.75 eV is obtained, we use the integral of $\sigma_1(\omega)$ generated analytically by this multioscillator model as an estimate of the spectral weight. Since the number of oscillators is very large and their parameters are all automatically adjustable, this procedure is essentially model independent. The spectral weight as a function of temperature is shown in Fig. 6, for different cutoff frequencies.

In order to illustrate that the measured spectra can indeed anchor the value of the low-frequency spectral weight without the need of the low-frequency extrapolations of reflectivity (as it would be the case in the conventional KK transform of reflectivity), we performed the following numerical test. In the superconducting state, the model contains a narrow oscillator centered at zero frequency which accounts for the spectral weight of the condensate and a narrow quasiparticle peak. We artificially displaced this oscillator from the origin to a finite frequency, yet below $\Omega_{\text{min}}$. The least mean square routine then readjusted all the other parameters in order to produce a new best fit. After that the overall fitting quality (the mean-squared error $\chi^2$) and the low frequency spectral

![Image](https://example.com/image.png)

**FIG. 5.** Optical constants of Bi2223 at selected photon energies as a function of temperature. Left panel: normal incidence reflectivity at 0.027, 0.14, and 0.56 eV. Right panel: real and imaginary parts of the complex dielectric function at 0.8, 1.24, and 2.05 eV. The photon energies are chosen close to the borders and in the middle of the experimental range.

![Image](https://example.com/image.png)

**FIG. 6.** Integrated spectral weight as a function of temperature (plotted vs $T^2$ for the reasons described in the text). The results corresponding to different ranges of integration are shown: (a) from 0 to 0.3 eV, (b) from 0 to 1 eV, (c) from 0 to 2.5, (d) from 0.3 to 1, and (e) from 1 to 2.5 eV.
weight remained almost unchanged with respect to the initial values, which is indicated by the errorbar in Fig. 7, while the reflectivity below 100 cm\(^{-1}\) now shows a very different behavior, as it is shown in Fig. 8. This demonstrates that the details of \(\sigma_1(\omega)\) and \(R(\omega)\) below 100 cm\(^{-1}\) are not essential for the determination of \(W(\Omega_c)\).

### B. Superfluid density

Even though the described above procedure does not involve a separate determination of the condensate spectral weight, it is nevertheless interesting to estimate the superfluid density and the penetration depth of Bi2223 as compared to other cuprates. This is approximately given by the spectral weight of the \(\delta\)-peak in the Drude-Lorentz model (see Table I). The condensate plasma frequency \(\omega_{ps}\) is about 10 300 cm\(^{-1}\) (\(\sim\)1.3 eV), which corresponds to the penetration depth \(\lambda\sim0.15\sim0.16\) \(\mu\)m. For comparison, the same procedure applied to our previous data on Bi2212\(^5\) gives \(\omega_{ps}=9500\) cm\(^{-1}\) and \(\lambda\sim0.17\) \(\mu\)m. The stronger superfluid component of Bi2223 as compared to Bi2212 correlates with a slightly higher plasma frequency of the former. One should note that in Ref. 22 a smaller value \(\omega_{ps}=9250\) (8890) cm\(^{-1}\) was found along the a (b) axis.

Another way to extract the superfluid density is to plot the value of \((\pi/2)\sigma_2(\omega)\omega\) plotted as a function of frequency (Fig. 9). The zero energy extrapolation of this value represents the spectral weight of the condensate. This value is about \(2.7\pm0.2\times10^6\) \(\Omega\)\(^{-1}\) cm\(^{-2}\). After multiplying with the usual factor \(120/\pi\) to make the conversion to plasma-frequency squared, we obtain \(\omega_{ps}/2\pi c=1030\pm40\) cm\(^{-1}\), in good agreement with the value extracted from the Drude-Lorentz fitting. Importantly, the reactive response of the low-lying quasiparticle peak cannot be separated from the condensate using only optical data down to a few meV.

### C. Temperature dependent redistribution of spectral weight

It is important to establish the relevant cutoff energy in the spectral weight integral. The optical conductivity of Bi2223 (see Fig. 4) as well as the one of Bi2212 (Refs. 5 and 22) shows a minimum around 1 eV separating the Drude peak from the lowest-energy interband transition, which is believed to be a charge transfer excitation. It is interesting that such a separation is even more pronounced if we plot the difference between conductivity spectra at \(T_c\) and a high temperature: \(\sigma_1(\omega, 280\) K\()\sim\sigma_1(\omega, 110\) K\) (see upper panel of Fig. 10). A sharp upturn below 20 meV, a dip at around 50 meV and a slow approaching of this difference to zero (while being negative) as the energy increases can be attributed to the narrowing of the Drude peak\(^{19}\). However, above the same characteristic energy of about 1 eV this monotonic trend totally disappears. Instead, \(\sigma_1(\omega, 280\) K\()\sim\sigma_1(\omega, 110\) K\) shows a dip at 1.5–2 eV, which corresponds to the removal of spectral weight from this region, as the system is cooled down. The difference \(\sigma_1(\omega, 25\) K\()\sim\sigma_1(\omega, 110\) K\) between the two spectra in the superconducting state also shows that the effect of the narrowing of the Drude peak does not noticeably extend above 1 eV (Fig. 10).

We can learn more from the corresponding differences of the integrated spectral weight \(W(\omega, 280\) K\()\sim W(\omega, 110\) K\) and \(W(\omega, 25\) K\()\sim W(\omega, 110\) K\) which are displayed in the lower panel of Fig. 10. In order to separate the effect of the superconducting transition from the temperature dependence already present in the normal state, we additionally plot the “normal-state corrected” spectral weight difference of the superconducting state relative to the normal state \(W_N-W_S\) calculated according to the procedure described in Sec. III D. Not surprisingly, all curves show an intense spectral structure below \(~0.3\) eV as a result of strong changes of the shape of...
the Drude peak with temperature. However, between 0.3–0.5 and 1.0–1.5 eV the spectral variation is weak and, importantly, both normal- and superconducting-state differences remain positive in this “plateau” region. This indicates an increase of the intraband spectral weight as the sample is cooled down and an extra increase in the superconducting state. Above 1.5–2 eV, the normal-state difference \( W(\omega,110K) - W(\omega,280K) \) decreases rapidly and becomes negative, which suggests that spectral weight is transferred between the charge-transfer and the intraband regions. In contrast, the superconducting state differences (both normal-state corrected and not) continue decreasing slowly and remain positive up to at least 2.5 eV, which means that the Ferrell-Glover-Tinkham sum rule is not yet recovered at this energy.

Another way to visualize the spectral weight transfer is to plot \( W(\Omega_c) \) as a function of temperature for different cutoff energies. In Figs. 6(a)–6(c) we present such curves for \( \Omega_c = 0.3, 1 \) and 2.5 eV.

One can immediately notice that the curves \( W(0.3 \text{ eV},T) \) and \( W(1 \text{ eV},T) \), apart from different absolute values, have almost identical temperature dependencies; accordingly, the integrated spectral weight between 0.3 and 1 eV [Fig. 6(d)] shows a very little variation with temperature. This is, of course, a manifestation of the existence of the discussed above “plateau” region in the frequency dependent spectral weight differences (Fig. 10). This observation is in line with the theoretical findings of Wrobel et al.\(^{16}\) who pointed out that spectral weight integrated to the hopping parameter \( t \) \(-0.3–0.4 \text{ eV} \) is representative of the kinetic energy of the \( t-J \) model.

Above \( T_c \), the spectral weight \( W(T) \) for the cutoff of 1 eV increases gradually with cooling down in a virtually \( T^2 \) fashion, which is most clearly seen when \( W(\Omega_c) \) is plotted versus \( T^2 \) [Figs. 6(a) and 6(b)]. A similar normal state behavior was observed in optimally and underdoped Bi2212 (Ref. 5) and in \( \text{La}_2\text{CuO}_4 \).\(^{24}\) Although a \( T^2 \) term follows trivially from the Sommerfeld expansion for the temperature broadening of the Fermi-Dirac distribution, the absolute value of this term turns out to be several times larger than what one expects from this expansion.\(^{25,26}\) The dynamical mean filed theory (DMFT) calculations within the Hubbard model\(^{29}\) showed that this may be caused by strong correlation effects. On the other hand, it was recently pointed out\(^{27,28}\) that the temperature dependence of the one-electron spectral function due to inelastic electron-boson scattering contributes to the overall temperature dependence of the optical sum rule much stronger than the Sommerfeld term. The extra contribution though is predominantly \( T \) linear if the boson energy is small.\(^{28}\)

At the superconducting transition the curve \( W(T) \) shows a sharp upward kink (slope change) close to \( T_c = 110 \) K. The same effect was observed in optimally and underdoped Bi2212.\(^{3}\) In order to directly compare these data to the results of Ref. 5, we plot in Fig. 7 the \( W(\Omega_c,T) \) for \( \Omega = 1.25 \text{ eV} \) together with the temperature derivative. One should stress that the corresponding kinks are already observed on the temperature dependence of directly measured optical constants (see Fig. 5). By extrapolating the normal state trend to \( T = 0 \) K, we can estimate the size of the superconductivity induced spectral weight transfer in the intraband region \( \Delta W = 0.8 \times 10^5 \text{ cm}^{-2} \), which is about 1% of the total intraband spectral weight (as shown by Fig. 6).

Remarkably, the upward kink of \( W(\Omega_c,T) \) at \( T_c \) is still observed for the cutoff of 2.5 eV [Fig. 6(c)] suggesting that superconductivity induced spectral weight transfer involves energies above the charge transfer gap, as we could already see from Fig. 10. In the context of the Hubbard model, the integrated spectral weight corresponds to the kinetic energy of the Hubbard Hamiltonian\(^{16}\) when the cutoff frequency is set much higher than the Sommerfeld term. The extra contribution though is predominantly \( T \) linear if the boson energy is small.\(^{28}\)

D. Temperature modulation analysis at \( T_c \)

The existence of relatively sharp kinks (slope discontinuities) at \( T_c \) on the curves of the temperature dependence of various optical functions (see Fig. 5) enables an alternative way to quantify the superconductivity-induced spectral weight changes which gives perhaps a better feeling of the error bars involved.\(^{30}\) We recently applied this procedure, which is essentially a temperature-modulation method, to a similar set of data on optimally doped Bi2212.\(^{18}\)
In order to separate the superconductivity-induced effect from the large temperature dependence in the normal state not related to the onset of superconductivity we apply the “slopecdifferencetion operator” \( \Delta_s \) defined as\(^{30} \)

\[
\Delta_s f(\omega) = \left. \frac{\partial f(\omega, T)}{\partial T} \right|_{T_0+\delta} - \left. \frac{\partial f(\omega, T)}{\partial T} \right|_{T_0-\delta},
\]

(3)

where \( f \) stands for any temperature dependent function. In Fig. 11 the slope-difference spectra \( \Delta_s R(\omega) \), \( \Delta_s \epsilon_1(\omega) \), and \( \Delta_s \sigma_1(\omega) \) are displayed with the error bars, which we determined from the temperature dependent curves such as shown in Fig. 5, using a numerical procedure, described in Ref. 18. Since \( \Delta_s \) is a linear operator, the KK relation between \( \epsilon_1 \) and \( \epsilon_2 \) holds also for the slope-difference spectra \( \Delta_s \epsilon_1(\omega) \) and \( \Delta_s \epsilon_2(\omega) \). Thus we can fit the latter spectra with a multisdelectric Lorentz model, which automatically satisfies the KK relations. If the number of oscillators is large enough, the procedure becomes essentially model independent. Using the same dielectric function we can additionally calculate the slope-difference spectrum of reflectivity as it is related to \( \Delta_s \epsilon_1(\omega) \) and \( \Delta_s \epsilon_2(\omega) \):

\[
\Delta_s R(\omega) = \frac{\partial R}{\partial \epsilon_1}(\omega, T_c) \Delta_s \epsilon_1(\omega) + \frac{\partial R}{\partial \epsilon_2}(\omega, T_c) \Delta_s \epsilon_2(\omega).
\]

(4)

The functions \( \frac{\partial R}{\partial \epsilon_1}(\omega, T_c) \) and \( \frac{\partial R}{\partial \epsilon_2}(\omega, T_c) \) can be derived from the analysis of optical spectra at \( T_c \).\(^{18} \)

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**FIG. 12.** The demonstration that \( \Delta_s W(\Omega_c) \) is well defined by the available data \((\Omega_c=1 \text{ eV})\). The mean squared error \( \chi^2 \) for the total fit of \( \Delta_s R(\omega), \Delta_s \epsilon_1(\omega), \) and \( \Delta_s \sigma_1(\omega) \) (shown in Fig. 11) as a function of \( \Delta_s W(\Omega_c) \) imposed. More description is given in the text.

The best fit of the slope-difference optical constants for Bi2223 is shown by the solid curves in Fig. 11. From this we also calculate the slope-difference integrated spectral weight \( \int_0^\infty \Delta_s \epsilon(\omega) d\omega \) as shown in Fig. 11(a). At \( \Omega_c =1 \text{ eV} \) we obtain the value of \( \Delta_s W = +1100 \text{ \Omega}^{-1} \text{ cm}^{-2} \text{ K}^{-1} \) which corresponds to the superconductivity-induced increase of spectral weight, in agreement with the previous analysis.

To check that the value of \( \Delta_s W \) is well defined by the present set of experimental spectra we repeated the fitting routine while forcing \( \Delta_s W(\Omega_c=1 \text{ eV}) = \Omega_c^2 \) to be equal to some imposed value \( \Delta_s W(\Omega_c) \) imposed. We did this for different values of \( \Omega_c \) imposed and monitored the quality of the fit, as expressed by the mean-squared error \( \chi^2 \). In Fig. 12 we plot \( \chi^2 \) as a function of \( \Delta_s W(\Omega_c) \) imposed. The best fit quality is, of course, achieved for the mentioned value \( \Delta_s W(\Omega_c) \) imposed = 1100 \text{ \Omega}^{-1} \text{ cm}^{-2} \text{ K}^{-1} \). It is evident that the fit quality deteriorates rapidly as \( \Delta_s W(\Omega_c) \) imposed is dragged away from this value. For example, the value of \( \chi^2 \) for the case \( \Delta_s W(\Omega_c) =0 \) (which would be the full recovery of the sum rule at 1 eV) is about 10 times larger than the optimal value; the corresponding data fit should be regarded as unacceptable. As it was discussed in Ref. 18 and in this paper, this is due to the fact that any change of the low-frequency spectral weight inevitably affects the value of \( \epsilon_1(\omega) \) and \( R(\omega) \) at higher frequencies.

To compare the approach described in this section and the one from the previous section, we derive from the curve \( W(T) \), plotted in Fig. 6 \( \Delta_s W(\Omega_c=1 \text{ eV}) \approx 1500 \text{ \Omega}^{-1} \text{ cm}^{-2} \text{ K}^{-1} \). This is not far from the aforementioned value of 1100 \text{ \Omega}^{-1} \text{ cm}^{-2} \text{ K}^{-1} which quantitatively confirms the central assertion of this paper, namely, that the low frequency spectral weight increases below \( T_c \).

As it was discussed in Ref. 18, the “normal-state corrected” spectral weight difference of the superconducting state relative to the normal state \( W_{SC}-W_N \) can be estimated from the slope-difference conductivity spectra

\[
W_{SC}-W_N = \alpha T_c \Delta_s W,
\]

(5)

where \( \alpha \) is a dimensionless coefficient. The choice of \( \alpha \) is suggested by the temperature dependence of \( W(\Omega_c) \).\(^{18} \) Since
we observe in both normal and superconducting state a temperature dependence close to \( T^2 \), then we choose \( \alpha = 1/2 \). The corresponding curve \( W_{SC} - W_N \) is shown as a dotted line in Fig. 10. One can see that it is slightly smaller than the direct difference \( W(\omega, 25 \text{ K}) - W(\omega, 7 = 110 \text{ K}) \). The latter fact is not surprising since the spectral weight in the normal state is increasing as a function of temperature.

IV. DISCUSSION AND CONCLUSIONS

The superconductivity-induced increase of low-frequency spectral weight, which implies the opposite to the BCS type lowering of the electronic kinetic energy, was previously observed in the underdoped and optimally doped bilayer \( \text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8 \). Recent studies\(^{11,31} \) show that this effect changes sign for strongly overdoped samples of the \( \text{Bi}2212 \), following the trend of the BCS model. A temperature dependence of the low energy spectral weight in accordance to the BCS theory has also been claimed for optimally doped \( \text{Bi}2223 \)\(^{5,6,18} \), although the latter result is controversial.\(^{18,33} \) Hence the “unconventional” superconductivity induced increase of low energy spectral weight appears to be a property characteristic of the Bi-based multilayer cuprates at and below optimal doping, but not of the entire cuprate family for all doping levels.

Recently an intriguing connection has been pointed out between the superconductivity induced increase of \( W(T) \) on the one hand, and the drop of scattering rate on the other hand.\(^{34} \) Since the former involves a spectral weight integral over 1 eV, whereas the latter is measured at microwave frequencies, these two experimental observations are seemingly unrelated. However, decreasing the scattering results in sharpening of the occupation distribution in \( k \) space. Hence a decrease of scattering automatically implies a decrease of the average kinetic energy, which in turn is observed as an increased \( W(T) \) when the sample turns from normal to superconducting. Since the standard BCS model (without a change in scattering) already predicts a superconductivity induced increase of the kinetic energy, the net result of both effects (scattering change and intrinsic BCS effect) will depend on the relative magnitude of these two effects. This explanation successfully relates two classes of experiments, without directly relating either one of these experiments to a particular pairing mechanism. The starting assumption of an anomalous drop in scattering is at this stage a phenomenological one, and still requires a microscopic explanation.

We observe that the increase of spectral weight below \( T_c \) in \( \text{Bi}2223 \) is larger than in \( \text{Bi}2212 \).\(^{5} \) We believe that this can be quite generally understood. It is well known that in the BCS theory

\[
\Delta E_{\text{kin}} \sim \Delta_{\text{SC}}^2 \sim T_c^2.
\]

The first equality holds generally for situations where the electronic occupation numbers \( n_k \) are redistributed in an energy range \( \Delta \) around \( E_F \), whereas the second is suggested simply by a dimensional analysis. Recent STM studies\(^{25} \) evidence that the SC gap is indeed larger in \( \text{Bi}2223 \) than in \( \text{Bi}2212 \) at a similar doping level. Assuming that Eq. (6) is also valid in \( \text{Bi}2212 \) and \( \text{Bi}2223 \) at optimal doping, and Eq. (1) is exact, we obtain

\[
\frac{\Delta W_{\text{Bi}2223}}{\Delta W_{\text{Bi}2212}} = \frac{T_c^2}{T_c^2} \approx 1.6.
\]

This ratio value matches the experimental observation, supporting the idea that the spectral weight transfer is intimately related to the SC-induced redistribution of the occupation numbers.

In summary, we have observed a superconductivity-induced increase of the in-plane low-frequency optical spectral weight \( W(\Omega_\perp) \) in trilayer \( \text{Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_8 \) near optimal doping from the reflectivity and ellipsometric measurements. Comparison to the case of optimally doped \( \text{Bi}2212 \) suggests that the size of the spectral weight transfer scales with the value of the critical temperature. In the normal state, the temperature dependence of the \( W(\Omega_\perp) \) is essentially proportional to \( T^2 \).

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APPENDIX A. DETERMINATION OF THE \( c \)-AXIS DIELECTRIC FUNCTION

In order to properly convert the pseudodielectric function measured ellipsometrically on the \( ab \)-crystal surface to the true dielectric function along the \( ab \) plane, we additionally measured the \( c \)-axis dielectric function using another crystal of \( \text{Bi}2223 \) grown under the same conditions.

We did spectroscopic ellipsometry from 6000 and 36 000 cm\(^{-1} \) on an \( ac \) surface of dimensions \( (a \times c) 3 \times 0.8 \text{ mm}^2 \) which we cut and polished with a diamond paper of 0.1 \( \mu \text{m} \) grain size. The surface image is shown in Fig. 13. Two orthogonal orientations of the sample were used, designated as \( (ac) \) and \( (ca) \), as shown in Fig. 13 which provided
and the real part of the dielectric function
played in Fig. 3. We observe a weak wavelength dependence
for
simultaneously in order to improve the accuracy of the in-

The main advance compared to these earlier results is that
the dielectric tensor and the angle of incidence \( \theta \): \( \epsilon_{\text{pseudo}} = f(\epsilon_{ab}, \epsilon_c, \theta) \). It was shown by Aspnes\(^{36} \) that in this case the pseudodielectric function should be much more sensitive to the \( ab \)-plane component, which lies along the crossing line of the plane of incidence and the sample surface, than to the \( c \)-axis one. In order to verify that this is the case, we show in Fig. 14 the “sensitivity functions” \( \partial \epsilon_{\text{pseudo}} / \partial \epsilon_{ab} \) and \( \partial \epsilon_{\text{pseudo}} / \partial \epsilon_c \) for the actual angle of incidence (in our case 74\(^\circ \)) on the base of described above \( ac \)-plane ellipsometry results.

One can see that the pseudodielectric function is indeed much less sensitive to the \( c \)-axis component since \( \partial \epsilon_{\text{pseudo}} / \partial \epsilon_c \) is about 4–5 times smaller than \( \partial \epsilon_{\text{pseudo}} / \partial \epsilon_{ab} \). For this reason, the temperature dependence of the \( c \)-axis dielectric function is expected to have only a minor effect on that of the measured pseudodielectric function. In order to verify this, we performed the \( c \)-axis correction of the \( ab \)-plane pseudodielectric function and calculated the spectral weight integral \( W(\Omega_c = 1 \text{ eV}) \) using in the first case the temperature dependent \( c \)-axis dielectric function and in the second case a constant, temperature-averaged one. The resulting temperature dependence of \( W(\Omega_c) \) in the former case is shown in Fig. 15, while the one in the latter case is shown in Fig. 6. One can see that accounting for the temperature dependence of the \( c \)-axis dielectric function does not have any significant influence on temperature dependence of \( W(\Omega_c) \), except for a stronger scatter of the datapoints as a result of the inevitable noise introduced by extra measurement on a small crystal surface. Therefore we used the temperature averaged \( c \)-axis data to correct the \( ab \)-plane pseudodielectric function in the main part of this paper.

FIG. 14. Real and imaginary parts of the sensitivity functions \( \partial \epsilon_{\text{pseudo}} / \partial \epsilon_{ab} \) and \( \partial \epsilon_{\text{pseudo}} / \partial \epsilon_c \) defined in the text.

FIG. 15. Spectral weight as a function of temperature (for the cutoff of 1 eV) obtained when a temperature dependent \( c \)-axis data were used to correct the \( ab \)-plane pseudodielectric function for the admixture of the \( c \)-axis component. The temperature resolution in this plot is lower than in Fig. 6 because the \( c \)-axis temperature dependent data have a resolution of 5 K.