I. INTRODUCTION

The study of the time evolution of complex systems through symmetry breaking transitions (SBT) is of great fundamental interest in different areas of physics. An SBT of particular general interest is the ultrafast normal-to-superconducting (N→S) state transition. Due to the small heat capacity of the electronic system, an optical pulse can efficiently suppress the SC state without heating the low-frequency phonon heat bath, which remains well below the critical temperature (Tc). This enables us to perform an ultrafast effective electron temperature quench across Tc with an ultrashort laser pulse, which is then followed by an ultrafast non-equilibrium N→S transition.

The ultrafast S→N→S transition in the cuprate superconductors has been initially studied by all-optical pump-probe technique followed by laser ARPES10–13. While the laser ARPES can directly resolve the momentum dependent quasiparticle (QP) distribution function, all-optical techniques offer better bulk sensitivity and greater flexibility. The lack of momentum resolution of an optical probe can be partially compensated by use of the optical dipole transition selection rules that depend on the probe-photon polarization15,16 and energy8,17 and enable selection of different parts of the Brillouin zone (BZ).

The electronic Raman-scattering tensor analyses have shown20 that the dielectric tensor fluctuations of different symmetries can be linked to charge excitations in different parts of the BZ. In particular, in a D4h point-symmetry corresponding to the ideal CuO2-plane symmetry, the dielectric tensor fluctuations with the B1g and B2g symmetries are linked to the anti-nodal and nodal BZ charge excitations, respectively, while the totally symmetric A1g fluctuations do not discriminate between the regions. The transient reflectivity, ΔR, is related to the Raman tensor and in Bi2Sr2CaCu2O8+δ (Bi2212) the B1g-like transient reflectivity component shows sensitivity to the SC state only, while A1g-like and B2g-like transient reflectivity components couple to both the SC and pseudogap (PG) order.20

While the all-optical transient response in the cuprates under weak excitation can be well described in terms of the photoinduced absorption of the photoexcited quasiparticles20, the response function in highly nonequilibrium states is unclear due to unknown relative contributions of collective and single-particle degrees of freedom to the transient optical reflectivity. To overcome this problem the standard two-pulse all-optical pump-probe technique was extended to a multi-pulse technique, which was shown to be instrumental in extracting the order parameter dynamics in a charge density wave compound20 as well as in the prototypical cuprate superconductor Bi2Sr2CaCu2O8+δ.
La$_{1.9}$Sr$_{0.1}$CuO$_4$\textsuperscript{21}.

Here we extend our previous study\textsuperscript{21} of an ultrafast S→N→S transition in La$_{1.9}$Sr$_{0.1}$CuO$_4$ to Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ in search of universality, and also to uncover potential important differences in the two materials with substantially different critical temperatures and pseudogap/SC gap ratios. By means of the all-optical multi-pulse technique combined with the polarization selective optical probe we were able to separate the SC state recovery dynamics from the previously studied\textsuperscript{22} PG state recovery dynamic and enable discrimination between relaxation in the nodal and anti-nodal BZ regions. The material has been studied previously by time-resolved techniques\textsuperscript{7,8,10–12,23}, but thus far there has been no systematic study of the the non-equilibrium transitions in this material, especially by the 3-pulse technique.

While we found that in La$_{1.9}$Sr$_{0.1}$CuO$_4$ the time dependent Ginzburg-Landau (TDGL) theory can provide a fair quantitative description of the SC order parameter recovery, only a qualitative description of the data is possible in Bi2212, which we attribute to the large SC order fluctuations in the PG state near time of the transition. In addition, when only a partial SC state suppression is achieved, the polarization resolved optical probe enables us to detect anisotropic SC-order recovery timescales, revealing a faster SC gap recovery in the anti-nodal direction in comparison with the nodal BZ regions.

II. EXPERIMENTAL

The sample used in this work was underdoped Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ (Bi2212) single crystal with $T_c \approx 78$ K ($\delta = 0.14$) grown by means of the traveling solvent floating zone method. Before mounting into a liquid-He flow cryostat the sample was freshly cleaved using sticky tape.

The pulse train from a 250-KHz 1.55-eV Ti:Sapphire regenerative amplifier was split into 50 fs destruction (D), pump (P) and probe (pr) pulse trains that were independently delayed with respect to each other. The result-

![Figure 2. The influence of the destruction pulse on the transient reflectivity in the superconducting state, when the D pulse arrives after the P pulse using the PDS. (a) and (b) show dependence on the D-pulse fluence and $t_{D-P}$ delay at $T = 15$ K, respectively. (c) $T$-dependence of the transient reflectivity suppression.](image)

In order to suppress the P beam scattering contribution to the signal in the BDS the P-beam frequency was doubled (3.1-eV P-photon energy) and a long-pass filter in front of the detector was used while the 1.55-eV D-photon energy was the same as in the first scheme.\textsuperscript{25}
Figure 3. (a) Transient reflectivity using the PDS at $T = 40$ K with $F_D = 68 \mu$J/cm$^2$. For comparison a transient measured in the PG state ($T = 120$ K) in the absence of the D pulse is shown vertically shifted below the main data. The vertical shaded areas indicates the read-out $t_{P-pr}$ delays (see text). (b) The same data set as in (a) shown as a colormap. At $t_{D-P} = 0$ both the PG and SC signal are suppressed. With increasing $t_{D-P}$ one can observe a sequential recovery of the negative PG response followed by the positive SC response.

Figure 4. (a) Recovery of the transient-reflectivity $B_{1g}$ component at $T = 40$ K and $F_D = 56 \mu$J/cm$^2$. The vertical shaded area represents the interval used to determine $A_{SC}$. Inset: black squares (left axis) - normalized amplitude of the response, red circles - relaxation rate extracted from single-exponential fits to the traces. Error bars are the standard errors of the regression analysis. (b) The same data as in (a) shown as a color map in $t_{D-P} - t_{P-pr}$. In the absence of the PG response recovery of the SC signal is evident already at $t_{D-P} \sim 1$ ps.

III. RESULTS

A. SC state destruction

To illustrate the destruction of the SC state, in Fig. 2 we plot the transient reflectivity for the case when the D pulse arrives after the P pulse using the PDS. Depending on the D pulse fluence, $F_D$, the transient reflectivity is suppressed to different degrees. Above $F_D \sim 70 \mu$J/cm$^2$, SC order is completely suppressed on a 200-fs timescale after the D-pulse arrival. Above $F_D \sim 70 \mu$J/cm$^2$ we observe also a small negative overshot lasting a few hundred femtoseconds followed by a weak recovery of the signal on a picosecond timescale. Both features vanish at the highest fluence of $\sim 400 \mu$J/cm$^2$. The suppression timescale does not depend on the D-pulse arrival time [Fig. 2(b)] nor temperature [Fig. 2(c)].

B. SC state recovery

In Fig. 3 we show a typical transient reflectivity data set measured in the PDS for the case when the P pulse arrives after the D pulse. After a complete suppression for $t_{D-P} \lesssim 0.5$ ps we first observe a recovery of the negative PG component on a 1-\sim-ps timescale followed by the
Figure 5. (a)-(c) The transient reflectivity at \( t_{P-pr} = 0.15 \) ps, corresponding to the delay at which the PG response peaks, as a function of \( t_{D-P} \) for different D-pulse fluences at different temperatures. (d)-(f) Evolution of the normalized \( \Delta R/R \) amplitude averaged in the \( 0.5 \text{ ps} \leq t_{P-pr} \leq 0.7 \text{ ps} \) range as a function of \( t_{D-P} \) for different D-pulse fluences at different temperatures. (g)-(i) the same for \( \Delta R_{B1g}/R \).

recovery of the positive SC component.

As shown previously,\(^{16}\) the PG response does not contribute to \( \Delta R_{B1g} \) so recovery of the SC component on the short \( t_{D-P} \) timescale can be observed more clearly in the BDS. In Fig. 4 we show a typical transient reflectivity data set measured using the BDS. As expected, the PG component is suppressed, but the signal-to-noise ratio is reduced due to a smaller \( \Delta R_{B1g}/R \) amplitude.

### IV. Analysis and Discussion

#### A. SC state destruction

The destruction timescale of \( \sim 200 \text{ fs} \) is \( T \) and \( F \) independent and faster than \( \sim 700 \text{ fs} \) in LSCO.\(^{6,9}\) In LSCO it was suggested\(^{6,9}\) that the high energy optical phonons created during the relaxation of the primary photo electron-hole pair are the dominating pair breaking excitation setting the destruction timescale.

The faster destruction timescale in Bi2212 does not exclude the same phonon mediated destruction mechanism since one polar optical phonon can be generated by a photoexcited electron/hole every \( \sim 5 \text{ fs} \).\(^{9}\) Taking the initial photo electron/hole energy of \( \sim 1 \text{ eV} \) and optical phonon energy of 50 meV leads to \( \sim 100 \text{ fs} \) photoelectron/hole energy relaxation time that is fast enough to be compatible with the experimental data.\(^{27}\) The phonon dominated pair-breaking destruction of the SC state is supported also by the large optical SC state destruction energy that is 5 times larger than the SC condensation energy.\(^{24}\)

#### B. Analysis of the SC state recovery

To analyze the recovery we first fit a finite-rise-time single-exponential relaxation model to the transient reflectivity in Fig. 4 to obtain the \( t_{D-P} \) dependent relaxation rate \( \gamma \). In the inset to Fig. 4 (a) we compare the relaxation rate \( \gamma \) from the fit to the amplitude of the \( B1g \) SC response \( A_{B1g} = \Delta R_{B1g}/\Delta R_{B1g, no-D} \),
where \( \gamma \) corresponds to the average of \( y \) in the interval\(^{28}\) \( t_{p-pr} = 0.5 - 0.7 \) ps and \( \Delta R_{B_{\perp},no-D} \) to the transient reflectivity in the absence of the D pulse.

\( \gamma \) and \( A_{B_{\perp}} \) initially recover on a similar time-scale of \( \sim 4 \) ps followed by slower dynamics extending towards the nanosecond timescale. As in the case of (La,Sr)CuO\(_{4+\delta}\) (LSCO), we attribute the suppression of \( \gamma \) during the first part of the recovery to the critical slowing down of the SC fluctuations in the vicinity of the transition.\(^{28}\) Upon the initial increase \( \gamma \) decreases on the nanosecond timescale indicating cooling of the probed volume: Since the effective temperature on longer timescales is far from the critical temperature the \( T \)-dependence of \( \gamma \) is no longer critical but determined by the Rothwarf-Taylor bottleneck dynamics.\(^{29}\)

Contrary to LSCO, where the PG component shows no suppression up to a rather high excitation fluence\(^{30}\), the PG component in Bi2212 shows suppression already below\(^{22}\) \( \sim 100 \) \( \mu \)J/cm\(^2\) so also the PG component is affected by the D pulse. To extract the SC component recovery dynamics in the PDS it is therefore necessary to take the PG dynamics into account.

The PG component peaks at \( t_{D-pr} = 0.15 \) ps. Traces of \( \Delta R/R \) at this \( t_{D-pr} \) as function of \( t_{D-pr} \) are shown in Fig. 6 (a-c). It is evident that at higher \( F_D \) the PG recovery leads to non-monotonous traces due to the sub-ps recovery timescale\(^{22}\) of the negative PG component preceding the recovery of the positive SC-state component. Due to the rather fast PG-component relaxation time\(^{30}\) [see Fig. 6 (a)] the contribution of the PG component to \( \Delta R/R \) should diminish with increasing \( t_{D-pr} \). Taking \( t_{p-pr} \) in the interval \( 0.5 \) ps \(< t_{p-pr} < 0.7 \) ps where \( \Delta R_{B_{\perp}} / R \) has a peak\(^{31}\) in the absence of the D pulse and the PG response is already suppressed, we calculate the normalized average, \( \gamma = \Delta R / \Delta R_{no-D} \). Indeed, \( A(t_{D-pr}) \) traces presented in Fig. 6 (d-f) show significantly less PG-component recovery and appear very similar to the equivalent \( A_{B_{\perp}}(t_{D-pr}) \) traces shown in Fig. 6 (g-i).

At \( T > T_c \) the PG recovery on the \( \sim 0.7 \) ps timescale\(^{22}\) to check whether the amplitude of the PG component is modified during the slower SC state recovery\(^{17}\) we compare in Fig. 6 the two readouts with different PG contribution taken at \( t_{p-pr} = 0.15 \) ps and the average in the interval \( 0.5 \) ps \(< t_{p-pr} < 0.7 \) ps \( \langle A \rangle \). At the highest \( F_D = 132 \) \( \mu \)J/cm\(^2\) it is possible to overlap the traces beyond \( t_{D-pr} \) \( \geq 1 \) ps at all measured temperatures when plot as a function of \( t_{D-pr} \) by vertically shifting\(^{31}\) and rescaling. At intermediate \( F_D \)s the complete overlap is not possible. The shifted and rescaled readouts at \( t_{p-pr} = 0.15 \) ps show slightly higher values in the \( \sim 2 - \sim 10 \)-ps delay range. This could indicate that the negative PG response at 1.55-eV probe-photon energy is transiently suppressed\(^{24}\) by the appearance of the SC order.

A possibly related suppression of the PG component in the SC state at 1.08-eV probe-photon energy was suggested recently\(^{17}\). Considering an earlier report\(^{17}\), however, where by selecting a particular polarization and probe-photon energy no suppression of the PG component in the SC state was observed, we attribute the difference between readouts in our experiment to the SC-gap dependent pre-bottleneck SC-state dynamics, which influences the readouts at \( t_{D-pr} = 0.15 \) ps.

\[ A_S = A_T - A_e^{\gamma - t_{D-pr}}/t_{rec}, \]

Figure 6. Comparison of \( A_S \) at two different \( t_{p-pr} \) as a function of \( t_{D-pr} \). The traces are vertically shifted for clarity as indicated by the horizontal thin lines. Full and open symbols correspond to \( t_{p-pr} = 0.15 \) ps and \( t_{p-pr} = 0.5 - 0.7 \) ps, respectively. The strongly PG-affected \( (t_{p-pr} = 0.15 \) ps) traces (full symbols) are vertically shifted and scaled to achieve the best match for \( t_{D-pr} \) \( \geq 1 \) ps.

C. SC state recovery timescale in nodal and anti-nodal response

In Fig. 6 (b) we compare the fluence dependencies of the SC recovery time, \( t_{rec} \), for both symmetries to the standard 2-pulse transient-reflectivity relaxation time, \( t_{2p} \). We estimate \( t_{rec} \) using a phenomenological exponential fit:

\[ A_S = A_T - A_e^{\gamma - t_{D-pr}}/t_{rec}, \]
Figure 7. (a) Fits of Eq. (1) to the $A_{B1g}$ trajectories at $T = 15$ K. (b) The recovery time of the superconducting response from the fits at $T = 15$ K as a function of fluence for $A_{B1g}$ and $A$ trajectories [Fig. 5 (d) and (g)]. For comparison the $\Delta R/R$ relaxation time from a two-pulse experiment at 15 K in Bi2212 is shown by open squares. The corresponding relaxation times in LSCO are at $T = 4$ K are shown by stars. (c) Temperature dependence of $\tau_{rec}$ of $A_{B1g}$ (full symbols) and $A$ (open symbols) trajectories.

Due to the sensitivity of the the $B_{1g}$ configuration to the anti-nodal Brillouin-zone (BZ) region this is consistent with a faster quasiparticle relaxation around anti-nodes either by recombination or by scattering into the nodal BZ region, which contributes to the $A_{B1g}$-dominated channel showing a slower decay.

The effect of the faster antinodal quasiparticle relaxation is also evident in our 3-pulse experiment, but only, when the SC order is not completely suppressed and $T$ is below $\sim 40$ K. From the 3-pulse data it appears that upon a modest suppression the SC gap recovers faster at the anti-nodes than near the nodes.

At higher $F_D$ upon a complete suppression of the SC gap our data suggest the recovery that is more homogeneous across the Fermi surface. This could be attributed to two factors. First, during the initial part of the recovery the suppression of the Rothwarf-Taylor phonon bottleneck and lifting of the SC-gap-imposed QP-relaxation phase space restrictions enable efficient transfer of the excess QP energy to the phonon bath together with efficient diffusion of excitations across all of the BZ. Second, at higher $F_D$ the lattice bath is heated closer to $T_c$ so the QP-relaxation phase space restrictions can be easier overcome by the phonon assisted QP scattering.

In LSCO [Fig. 7 (b)] $\tau_{rec}$ is similarly to $\tau_{2p}$, significantly longer than in Bi2212. The generally slower $\tau_{2p}$ and $\tau_{rec}$ in LSCO could be attributed to the smaller SC gap enhancing the quasiparticle relaxation bottleneck. Moreover, in LSCO $\tau_{rec}$ increases monotonically above the destruction threshold fluence $F_{Dth} = 4.2 \mu J/cm^2$, while in Bi2212 with $F_{Dth} = \sim 16 \mu J/cm^2$, the increase starts only above $\sim 4F_{Dth}$. This could be attributed to the lattice temperature after the quench being close to $T_c$ in LSCO than in Bi2212 resulting in a stronger critical slowing down of the SC order parameter dynamics.

D. Time dependent Ginzburg-Landau model

We proceed by analyzing the trajectory of the SC amplitude through the transition in the framework of the time-dependent Ginzburg-Landau (TDGL) theory. In previous study of the SC-order recovery in LSCO we have shown that the TDGL theory fails to consistently describe the ultrafast optical destruction of the SC condensate. On the other hand, the SC condensate recovery can be quantitatively modeled using a phenomenological response function and the Ginzburg-Landau time, $\tau_{GL}$, as the only free fit parameter assuming a finite magnitude of the initial depth-dependent order parameter (Fig. 3c in Ref. [21]). The magnitude of the initial order parameter corresponds to the magnitude of the frozen SC fluctuations after the quench from the normal/PG to the SC state which is a function of the depth-dependent quench-rate (Eq. (4) in Ref. [21]). In LSCO even better fit is possible using a phenomenological depth-dependent initial order parameter $\psi_{BC}(z)$:

$$\psi_{BC}(z) = \begin{cases} c z \left( \frac{U_D(z)}{U_{th}} > 1 - T/T_c \right) & ; U_D(z) > U_{th} \\ \sqrt{1 - \frac{T}{T_c}} \left( \frac{U_D(z)}{U_{th}} \right) & ; U_D(z) \leq U_{th} \end{cases}$$

where $c$ is an additional $F_D$-dependent free parameter, $U_D(z)$ the depth-dependent absorbed optical-energy density and $U_{th}$ the SC-destruction-threshold optical-energy density.

In the following we apply a similar TDGL approach to the SC state recovery dynamics in Bi2212.

1. Response function

As a starting point one needs to establish the relation between the superconducting order parameter magnitude, $|\psi_{GL}|$, and the transient optical response amplitude. This relation was in the case of LSCO established phenomenologically from the temperature dependence of the normalized weak-excitation $\Delta R/R$ amplitude, $A_S$.

In Bi2212 $A_S$ does not go to zero at $T_c$ due to the large pairing fluctuations above $T_c$ as shown in Fig. [8]
ing) order parameter, $\psi$ phase coherence is established at still apply the GL description assuming that only the SC amplitude $K$. Implying the standard GL $T$ to the mean-field pairing critical temperature, $T$.

A superconducting order parameter evolution presented in Fig. 8. (a) The amplitude of the normalized transient superconducting response as a function of temperature. The response function for the theoretical calculation of the superconducting order parameter presented in Fig. 8(b). In Fig. 8(c) and (d) we show comparison of our data to both, $A_{\text{AR}}$ and the TR-ARPES gap, $\Delta_{30g}$.

For the $B_{1g}$ configuration we find a surprisingly good match between $A_{\text{AR}}$ and $A_{\text{B1g}}$ in the low fluence region, where the SC gap is only partially suppressed. At higher $F$, where the gap is completely suppressed, the dynamics appears significantly different below $\sim 3$ ps, unless we compare curves with very different fluences. Ignoring the response function a direct comparison of $\Delta_{30g}$ at $F = 23 \mu J/cm^2$ to $A_{\text{B1g}}$ at 4.4 times higher $F = 102 \mu J/cm^2$ gives a good match in the region of the strong suppression of the gap.

For the $A_{1g}$ dominated configuration a better match is observed when we compare $A$ to $\Delta_{30g}$ directly [Fig. 8(d)] while $A_{\text{AR}}$ shows consistently higher magnitude than $A$. Similarly to the $B_{1g}$ configuration, a good match is observed at a complete SC gap suppression between the TR-ARPES trajectory at $F = 15 \mu J/cm^2$ and $A$ at 4.5 times higher $F = 68 \mu J/cm^2$.

Assuming that the TR-ARPES SC gap dynamics is identical to the bulk gap dynamics the difference between the fluences of the corresponding-timescales data can be, at least partially, attributed to the smearing of the optical-probe dynamics due to the depth-dependent excitation density and SC gap suppression. This is corroborated by the convergence of the optical and TR-ARPES trajectories with similar fluence on longer timescales, when the spatial inhomogeneity is expected to decrease.

The inaccuracy of the empirical response function [Fig. 8(b)] can further contribute to the difference, especially in the region of small SC order parameter. Contrary to LSCO where the response function is linear up to $A_0 \sim 0.8$ the s-shape of the response function in the present case suggests that $A_{\text{AR}}$ might be underestimated for low values of the gap.

Importantly, taking into account the inherent differences between the techniques we can conclude that the TR-ARPES Fermi-arc SC gap and the antinodal SC gap inferred from the $B_{1g}$ channel multi-pulse optical probe show qualitatively identical suppression and recovery dynamics.

2. Simulations

As in the case of LSCO we simulate the evolution of the order parameter through the transition by solving the dimensionless form of the first of the two TDGL equations:

$$\frac{\partial \psi}{\partial t} = \alpha_r(t, z) \psi - \psi|\psi|^2 + \nabla^2 \psi,$$

where time and length are measured in units of $\tau_{\text{GL}}$ (fit parameter) and the coherence length, respectively.
\[ \alpha_r(t, z) \] is a time- and depth-dependent reduced temperature which is the solution of the three temperature model \[ \text{Eq. (2)} \] combined with the heat diffusion equation \[ \text{Eq. (21)} \].

We neglect the second TDGL equation and any lateral variation of the order parameter, assuming that all the Kibble-Zurek (KZ) physics \([\text{Eq. (21)}]\) can be phenomenologically absorbed into the initial order parameter \( \psi_{BC}(z) \) using \[ \text{Eq. (2)} \], and the phase dynamics, i.e. the dynamics of vortices, does not significantly modify the order parameter amplitude.

The two fitting parameters \( c \) and \( \tau_{\text{GL}} \) are rather independent. While the first defines the Kibble-Zurek-physics-related amplitude of the response at \( t = 0 \), the second defines the time-scale of the recovery. In Fig. 9 we present typical results of the simulations for two different values of the \( \tau_{\text{GL}} \) optimized to fit the highest-\( F_D \) and the lowest-\( F_D \) trajectories at 15 K, respectively. One can see that while a decent agreement for a targeted curve can be achieved, one needs to significantly vary \( \tau_{\text{GL}} \) to fit the complete data set. Since such variation is unphysical, we can state that the presented TDGL approach is only sufficient to describe the present data qualitatively, contrary to what was found in LSCO, where a more quantitative description is possible.

The lack of quantitative description can not be attributed to the omission of the second TDGL equation and the resulting vortex dynamics it describes. While at a partial SC-state order parameter suppression no KZ-vortices formation is expected more vortices would be created with further suppression. The presence of vortices at increased order parameter suppression is expected to further slow down the SC-state recovery. Looking at Fig. 9 one can clearly see, that even without the vortex dynamics the TDGL solutions display a stronger recovery-timescale slowdown with increased order parameter suppression than the experimental data, so inclusion of the vortex dynamics into modeling is expected to only increase the discrepancy.

On the other hand, the lack of quantitative description is not very surprising due to the large pairing fluctuations contribution to the transient reflectivity above \( T_c \) that prevents strict applicability of the TDGL theory and undermine the phenomenological link between the order parameter magnitude and the experimentally observable \( A_S \).

\section{V. SUMMARY AN CONCLUSIONS}

Our systematic investigation of the ultrafast optical suppression and recovery of the superconducting state in Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ by means of polarization-selective multi-pulse optical time-resolved spectroscopy leads to some interesting, and somewhat surprising new findings. We found that the SC order is suppressed on the 200-fs timescale, comparable to the recent laser TR-ARPES results. The destruction timescale is independent of the temperature and optical destruction pulse energy and is consistent with a photoexcited carrier energy-transfer to the high-energy pair breaking phonons.

The recovery of the SC order is slower appearing on the 2-8 ps timescale showing non-monotonous dependence on the destruction pulse energy. At low \( T \) and a partial SC-state suppression the data shows that the SC gap in the antinodal region recovers faster than near the nodes. Perhaps surprisingly, the recovery also slows down with decreasing \( T \) highlighting the importance of thermal fluctuations in the recovery mechanism. When the SC state is strongly suppressed, the recovery becomes non-exponential with the recovery timescale slowing down, becoming \( T \)-independent.

The fact that the antinodal SC order parameter recovery dynamics inferred from the B$_{1g}$ channel and the TR-ARPES Fermi-arc SC gap dynamics show qualitatively identical recovery dynamics gives us confidence in the significance of the multipulse technique.

Despite strong SC fluctuations above \( T_c \) and the anisotropic SC-gap recovery the time dependent Ginzburg-Landau model qualitatively describes the SC-order temporal dynamics reasonably well, considering its limitations.

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The optical penetration depth in Bi2212 is of the order of 100 nm in comparison to a nm scale photoelectron escape depth. The Bi_{1g} optical response is sensitive to a broad region near the anti node while the A_{1g} response samples both the nodal and antinodal regions.

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