Carrier relaxation dynamics in intra-gap states: the case of the superconductor $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ and the charge-density-wave semiconductor $\text{K}_{0.3}\text{MoO}_3$.

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(January 14, 2018)

The unusual slow carrier relaxation dynamics - observed in femtosecond pump-probe experiments on high-temperature superconductors and recently also in a charge-density-wave system - is analyzed in terms of a model for relaxation of carriers in intra-gap states. The data on $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ near optimum doping and $\text{K}_{0.3}\text{MoO}_3$ are found to be described very well with the model using a BCS-like gap which closes at $T_C$. From the analysis of the data we conclude that a significant intra-gap density of localized states exists in these materials, which can be clearly distinguished from quasiparticle states by the time-resolved optical experiments because of the different time- and temperature-dependences of the photoinduced transmission or reflection. Localized charges are suggested to be the most likely origin of the intra-gap states, while the similarity of the response in the two materials appear to exclude spin and vortex excitations.
I. INTRODUCTION

Photoinduced absorption or reflection spectroscopy using femtosecond lasers is potentially a very powerful tool for the study of the electronic structure of superconductors and related materials. Recent pump-probe experiments on cuprate superconductors and some other materials have shown that a photoinduced change in absorption or reflection can be observed at low temperatures and especially for $T < T_c$. The effects are believed to be caused by excited state absorption of the probe pulse from photoexcited quasiparticle (QP) states and theoretical analysis of the response was found to be in good agreement with experimental data on YBa$_2$Cu$_3$O$_{7-\delta}$ over a wide range of doping. However, in addition to the QP response which occurs on the picosecond and subpicosecond timescale, a distinct slower response was also consistently observed in YBa$_2$Cu$_3$O$_{7-\delta}$ (YBCO) and Bi$_2$Sr$_2$CaCu$_2$O$_8$ (BISCO) and Bi$_2$Y$_x$Ca$_{1-x}$SrCu$_2$O$_8$ and more recently in the charge-density-wave (CDW) quasi one-dimensional Peierls insulator K$_{0.3}$MoO$_3$. It was thought to be of non-thermal origin (detailed analysis is given in Ref. 3) and occurs on a timescale of $10^{-8}$ s or longer (see Fig. 1 for details). Its anomalous $T$-dependence - which is qualitatively different from the $T$-dependence of the fast QP recombination dynamics - lead the authors to the suggestion that the signal is due to localized states near the Fermi energy. However, the processes involved were so far not discussed in any detail.

In this paper we examine quantitatively the photoinduced absorption (or reflection) from localized intra-gap states. We develop a theoretical model for two different cases: (i) the case of a material with a BCS-like collective gap and (ii) the case of a $T$-independent gap (sometimes called a pseudogap) which exists above $T_c$. We compare the model predictions with the available data on the temperature-dependence of the photoinduced relaxation on the nanosecond timescale in two different materials, both of which are generally thought to have a low-temperature gap: YBa$_2$Cu$_3$O$_{6.9}$ ($T_c = 90$K) and K$_{0.3}$MoO$_3$ ($T_c = 183$K).

II. THEORETICAL MODEL

As the name implies, pump-probe spectroscopy involves the excitation of the material by an ultrashort pump laser pulse and the subsequent measurement of the resulting change in optical absorption, transmission or reflection of the sample caused by photoexcited charge carriers. As photons from the pump laser pulse are absorbed, they excite electrons and holes in the material (see schematic diagram in Fig.2). These particles release their extra kinetic energy by scattering amongst themselves and with phonons (step 2 in Fig. 2a)). This energy relaxation process is very rapid and the particles end up in QP states near the Fermi energy ($E_F$) within $\tau_e = 10^{-10}$ fs. Subsequent relaxation is slowed down by the presence of the gap and a relaxation bottleneck is formed. From pump-probe photoinduced transmission experiments in YBCO and BISCO, the relaxation times of the photoexcited QPs were found to be in the range $\tau_{QP} = 0.3-3$ ps.

In addition to the picosecond transient, the signal on nanosecond timescale has been consistently observed in HTSC and recently also in quasi 1D CDW insulator K$_{0.3}$MoO$_3$. In Fig. 1 the photoinduced transient taken on YBCO at $T = 80$ K is shown (squares). After photoexcitation ($t = 0$) the signal relaxes within 10 ps to some non-zero value, that can be represented by the constant on the 100 ps timescale. The lifetime of the slow component, $\tau_L$, cannot be directly measured, since it appears to be longer than the inter-pulse separation of $t_r \approx 10 - 12$ ns. This results in the signal pile-up due to accumulation of the response from many pulses. The magnitude of this pile-up, $A$, given by the difference in the signal amplitude when the pump pulse is unblocked (2), and zero signal, when the pump is blocked (1) at negative time delays (see Fig. 1), is several times larger than single pulse contribution, $a_0$, given by the difference between (4) and (2).

Assuming - for simplicity - that relaxation of the signal is exponential, we can write an equation for the steady state amplitude $A$:

$$A = a_0 \sum_{n=0}^{\infty} \exp(-nt_r/\tau_L) = a_0/(1 - \exp(-t_r/\tau_L)) \quad (1)$$

Here $a_0$ is the amplitude of the signal from a single pulse and $\tau_L$ is the relaxation time. Usually, $A \gg a_0$ and we can expand the exponent in the denominator of Eq.(1), to obtain $\tau_L = \tau_r (A/a_0)$. From the experiments it appears that $\tau_L > 10^{-7}$ s. This is long in comparison with the phonon relaxation time and with the phonon escape time from the excitation volume into the bulk or thin film substrate, which is typically $10^{-10}$s, so we can ignore phonon escape effects and discuss only intrinsic relaxation processes.

The process giving rise to the actual photoinduced optical signal from intra-gap states is shown as step 3 in Fig.2a). Here we do not discuss the optical probe process in detail, but make the very general assumption that the photoinduced change in sample transmission $\Delta T/T$ or reflection $\Delta R/R$ is proportional to the photoinduced density of localized states populated by the laser pump pulse. The photoinduced signal $\Delta T/T$ (or $\Delta R/R$) is then proportional to the number of filled localized states and the $T$-dependence is mainly determined by the occupation of the intra-gap states.

Since the relaxation time $\tau_L$ is long in comparison
with \( \tau_{QP} \) and phonon relaxation times, we assume that phonons and quasiparticles can be described by equilibrium densities \( N_\omega \) and \( N \) respectively. For the relaxation of the localized carriers we apply arguments similar to those originally proposed by Rothwarf and Taylor [3]. The rate equation for the total density of localized excitations \( N_L \) is then given by:

\[
\frac{dN_L}{dt} = -RN_L^2 - \gamma N_L + \gamma N + \beta N_\omega.
\] (2)

The first term in Eq.(2) describes the recombination of two localized excitations to a Cooper pair with a recombination rate \( R \). The second and the third terms describe the exchange of an electron or a hole between the localized and quasiparticle states with density \( N_L \) and \( N \) respectively and with a rates \( \gamma \) proportional to \( \Delta \). (Above \( \alpha \) disappear as \( \Delta \) at low temperatures. The term proportional to the square of the order parameter [9,10]:

\[
\text{behavior at low temperatures.}
\]

Given in the Appendix. Next we consider the limiting case of pre-formed pairs with a level splitting between paired and unpaired states, and the latter case, the gap is better considered as an energy-gap formed by a CDW) and b) a Bose condensate of pairs.) In this paper we consider a) a superconductor case of a CDW gap, the Cooper pairs are replaced by pre-formed pairs with a BCS-like gap (which can also be used in the case of a gap formed by a CDW) and b) a Bose condensate of pre-formed pairs with a T-independent gap. In the latter case, the gap is better considered as an energy-level splitting between paired and unpaired states, and the main difference is that the gap is T-independent and exists above \( T_c \). A general exact solution to Eq.(3) is given in the Appendix. Next we consider the limiting behavior at low \( T \) and near \( T_c \).

In the case of a collective BCS-like gap, \( \Delta_s(T) \), the recombination rate below \( T_c \), is to the lowest order in \( \Delta_s \) proportional to the square of the order parameter [9,10]:

\[
R \approx \alpha(\Delta_s(T)/\Omega_c)^2
\] (4)

where \( \alpha \) is a constant, and \( \Omega_c \) is the phonon spectrum cutoff frequency. (Above \( T_c \), all recombination processes disappear as \( \Delta_s(T) \to 0 \).) To relax the carriers in localized intra-gap states via quasiparticle states (the term \( \gamma n_L \) in Eq.(3)), an energy of the order of \( \Delta_s \) is required and so this process is exponentially suppressed at low temperatures. The term proportional to \( N_{LO} \) is also small at low temperatures, because the number of thermally excited localized excitations is small as \( (k_B T/\Delta_s(0))^\mu \), where \( \mu \) depends on the density of localized states. Therefore, for analysis of the relaxation of localized excitations at low temperatures \( T \ll T_c \), we retain only the first term in Eq.(3) giving a solution of the form:

\[
n_L(t) = n_L(0)/(n_L(0)Rt + 1)
\] (5)

To obtain the stationary solution for a repetitive laser pump pulse train excitation, we use the condition that the total number of localized excitations that recombine between two laser pulses should equal the number of localized excitations created by each laser pulse:

\[
n_L(0) - n_L(0)/(n_L(0)Rt + 1) = \eta n_{ph}(T)
\] (6)

where \( \eta \propto \gamma \Delta \) is the probability of trapping a QF into a localized state and \( n_{QP}(T) \) is the number of photoinduced QPs at temperature \( T \) created by each laser pulse. Since the number of photoexcited carriers is typically small compared to the overall carrier density \( \eta n_{QP} \ll n_L(0) \), we can estimate \( n_L(0) \) as

\[
n_L(0) = \sqrt{\frac{\eta n_{QP}(T)}{Rt}}
\] (7)

As a result, combining (4) and (7) we get an expression for the \( T \)-dependence of the photoinduced transmission (or reflection) amplitude at low temperatures:

\[
|\Delta T/T| \propto n_L(0) = \sqrt{\frac{\eta n_{QP}(T)}{\Omega_c \Delta_s(T) \alpha^2}}
\] (8)

where \( n_{QP} \) for a BCS-like case is given by [9,10]:

\[
n_{QP}(T) = \frac{\mathcal{E}_I}{(\Delta_s(T) + k_B T/2)} \frac{(2\pi)^{1/2}}{\sqrt{\Gamma_0}} \sqrt{\frac{2k_B T}{\Delta_s(T) \exp(-\Delta_s(T)/k_B T)}}
\] (9)

Here \( \mathcal{E}_I \) is the energy density deposited per pulse, \( \nu \) is the effective number of phonons per unit cell involved in the relaxation process and \( N(0) \) is the density of states at the Fermi energy in units eV\(^{-1}\)cell\(^{-1}\)spin\(^{-1}\). Near \( T_c \), when the number of thermally excited localized carriers becomes comparable or larger than number of nonequilibrium carriers, the relaxation terms \( 2RN_{LO} n_L \) and \( \gamma n_L \) become dominant, and the solution to Eq. (3) has an exponential rather than a power law time dependence of the form \( n_L(t) = n_L(0) \exp(-t/\tau) \), where \( 1/\tau = (2R(T)N_{LO} + \gamma) \). For \( T \to T_c \), the temperature dependence of the photoinduced signal amplitude is then given by:

\[
|\Delta T/T| \approx n_L(0) = \frac{\eta n_{QP}(T)}{(2N_{LO} R(T) + \gamma)t_c}
\] (10)
Thus the predicted amplitude of the signal increases with increasing $T$ up to $T_c$ and drops to zero above $T_c$.

In the case of a $T$-independent gap (or "pseudogap"), $\Delta_p$, we assume that the gap exists at all temperatures. Since the gap does not close at $T_c$, the recombination rate does not go to zero at $T_c$ and instead of (4) we have a constant, $\Gamma$.

However, below $T_c$ the presence of the condensate may also have an effect on the recombination of localized excitations. In general, the relaxation rate is a function of the order parameter. To take this into account, we can expand it in terms of even powers of $\Delta$. Near $T_c$ the order parameter is small and we can keep only the lowest power in $\Delta^2$ [4]. If we assume that the order parameter exhibits mean-field behavior ($\Delta \propto \sqrt{1 - T/T_c}$), then:

$$R \simeq \alpha(1 - T/T_c) + \Gamma$$

(11)

where $\alpha$ is phenomenological constant which describes the dependence of the relaxation rate on $T$ below $T_c$, which in general is not equal to 0.

A slightly different expression for $R$ is obtained if we assume that the recombination rate is dependent on the pair momentum in the condensate through the kinetic energy. To illustrate this, we write the recombination rate as $\Gamma = \Gamma_0 + \tilde{\Gamma}$ where $\tilde{\Gamma}$ is the momentum averaged recombination rate, and $\Gamma_0$ is the recombination rate for pairs with $k = 0$. The recombination rate is generally proportional to the number of pairs in the condensate $n_p$. For Bose condensation this is given by $n_p \propto (1 - (T/T_c)^{3/2})$, and we thus obtain a formula for the total relaxation rate which is similar to Eq. (11), but with a different temperature dependence in the first term:

$$R = (\Gamma_0 - \tilde{\Gamma})(1 - (T/T_c)^{3/2}) + \tilde{\Gamma}$$

(12)

In principle the two cases (11) or (12) can be distinguished by measurements of the $T$-dependence of the photoinduced transmission or reflection below $T_c$, although the difference will be very small and very high quality data is needed to do this.

To obtain the photoinduced signal amplitude, we substitute Eq.(11) into eq. (3):

$$|\Delta T / T| \propto n_L(0) = \sqrt{\frac{\nu n_{QP}^\prime}{\alpha(1 - (T/T_c)^\beta) + \Gamma} \frac{1}{R}}$$

(13)

where $\beta$ is a constant (generally 1 or 3/2) and $n_{QP}^\prime$ is the equivalent of formula (9) for the case of a temperature independent gap [4];

$$n_{QP}^\prime = \frac{E_T / \Delta^p}{1 + \frac{2\nu}{\alpha(1/T_c)} \exp(-\Delta^p / k_B T)}$$

(14)

In contrast to the BCS case, expression (13) is non-zero above $T_c$ and reduces to:

$$|\Delta T / T| \propto n_L(0) = \sqrt{\frac{\nu n_{QP}}{\Gamma T}}$$

(15)

which implies that the photoinduced absorption signal should remain observable well above $T_c$ and should reveal the presence of a pseudogap if it exists.

Just as before, a crossover to exponential time relaxation takes place when the number of thermally excited excitations become large and the second term in the Eq.(3) becomes dominant, i.e. $T \sim \Delta_p$, leading to a linear intensity dependence of $|\Delta T / T|$ given by modified Eq.(10) with $n_{QP}$ and $R$ from Eq.(12). Note that this crossover may occur at higher temperatures than experimentally measured since $T$-independent gap is typically $\Delta_p \gg 300 K$.

### III. COMPARISON WITH EXPERIMENTAL DATA

The temperature dependence of $|\Delta T / T|$ or $|\Delta R / R|$ for optimally doped YBa$_2$Cu$_3$O$_{7-\delta}$ [11] and for K$_{0.3}$MoO$_3$ [4] is plotted in Fig. 3a) and b) respectively. In both compounds the signal amplitude increases with increasing temperature followed by an abrupt drop above $T_c$. Now we compare the predicted $T$-dependences of $|\Delta T / T|$ (Eqs. (8) and (10)) with the data. The parameters used in the fits are the same as previously used in the analysis of the fast relaxation component [2-11] and for values of dimensionless constant $\frac{2\nu}{\alpha(1/T_c)} \simeq 30$ for YBa$_2$Cu$_3$O$_{7-\delta}$ [8] and $\simeq 10$ for K$_{0.3}$MoO$_3$ [4]. Since the magnitude of the gap is of the order of $\Delta_c(0) \sim 5k_T$, we expect that Eq. (8) is valid up to temperatures close to $T_c$. In Fig. 3a) and b) we show the calculated temperature dependences of $|\Delta T / T|$ using Eq.(8) in comparison with experimental data for YBa$_2$Cu$_3$O$_{7-\delta}$ and K$_{0.3}$MoO$_3$ respectively. We would like to stress that Eq. (8) is independent of the shape of DOS of localized states within the gap and shows a universal temperature dependence. It can be seen from the Fig.3b) that at low temperature there is deviation of the calculated curve from the experimental points. To explain this effect we should remember that Eq.(4) for constant $R$ is valid near $T_c$ where $\Delta_c(T)$ is small. (If we add next - fourth order - term in the expansion of $R$ in powers of $\Delta_c(T)$ we can account for this discrepancy.)

Near $T_c$ Eq.(8) fails and Eq. (10) should be used (see also Appendix). It leads to a crossover from square root intensity dependence of $|\Delta T / T|$ at low $T$ described by Eq.(8) to linear intensity dependence near $T_c$ predicted by Eq.(10).

Finally let us discuss the effect of $\tilde{\gamma}$. In Fig. 3a) and b) fits to the data using the general solution (Eq.(19)) are shown with dashed lines. In these fits we have also added the fourth order term in the expansion of $R$ in powers of $\Delta_c(T)$. As can be seen from these fits the effect of $\tilde{\gamma}$
becomes important near $T_c$ by cutting the divergence of $|\Delta T/T|$ as $T \to T_c$.

In Fig.3c) we have plotted calculated values of $|\Delta T/T| \propto n_L(0)$ for the case of temperature independent pseudogap (which might be applicable in underdoped cuprates for example) as a function of temperature for different values of parameter $\alpha$ using Eq.(13). As can be seen from the Fig.3c) in this case slow relaxation via localized states is present also above $T_c$. This effect is due to temperature dependence of $n_QP$ controlled by $T$-independent pseudogap above $T_c$.

In cuprate superconductors there is spectroscopic evidence suggesting that there is a significant density of states in the gap possibly extending to the Fermi level, which is often attributed to a $d$-wave gap symmetry. However, by normal spectroscopies it is difficult to determine if the states in the gap are QP states or, for example, localized states. Time-resolved spectroscopy can answer this question rather effectively because of the different time- and temperature-dependences of the QPs and localized carrier relaxations. It was argued that in the presence of impurity scatteringQP DOS in the $d$-wave state remains finite at zero energy $\Delta T/T_L$ [4]. Recently it was proposed [5] that the quasiparticles in the superconducting state may become strongly localized for short coherence length $d$-wave superconductors. However this statement has been questioned by Balatsky and Salkola [6] and remains controversial. On the basis of available experimental data we cannot make any definite conclusion about origin of intragap localized states.

We can, however estimate the density of the intra-gap states from the available data by assuming that the optical probe process (step 3) is similar for excited state absorption from localized states and for QPs. Both optical probe processes involve allowed transitions to the same final state $E_F$ and so this assumption is not unreasonable. From typical photoinduced reflection data for YBCO (as in Figure 1), we find that approximately $|\Delta T/T|_L \simeq |\Delta T/T|_{QP}$, implying that also $n_L(0) \simeq n_{QP}$. From this we can conclude that the density of intra-gap states is comparable with the density of QP states. This observation has important implications for the interpretation of frequency-domain spectroscopies, since it suggests that the spectra should show a very significant intra-gap spectral density due to localized states, irrespective of the gap symmetry.

Assuming that the optical transition probability of the probe pulse is the same for QPs as for the intra-gap states, from Eqs. (8) and (9) we obtain:

$$\frac{|\Delta T/T|_L}{|\Delta T/T|_{QP}} \sim \frac{n_L}{n_{QP}} = \frac{\eta \gamma L}{t_r}$$

(16)

where $|\Delta T/T|_{QP}$ is the photoinduced transmission due to the QPs. From Fig. (1) typically $n_L/n_{QP} = 0.1 - 1$ and using a pulse repetition rate $t_r = 12$ ns and with $\gamma L = 100$ ns, we obtain an estimate of the trapping probability for carriers by localized states of $\eta = 0.1 - 1$.

A detailed discussion of the origin of the localized intra-gap states in the cuprates should be deferred until more systematic data as a function of doping is available, and we only mention some of the most likely possibilities: (i) localized states associated with the inhomogeneous ground state of the cuprates (stripes) [7], (ii) intrinsic defect states, (iii) localized QP states in $d$-wave superconductor [8] and possibly (iv) holons [10]. In $K_{0.3}MoO_3$, the nature intra-gap excitations has been a subject of extensive study over the years and the reader is referred to ref. [2] for a review. However, the fact that the signals in $K_{0.3}MoO_3$ and $YBa_2Cu_3O_7-\delta$ are very similar appears to rule out both spin excitations and vortex states, leaving localized charges as the most plausible origin of the intra-gap states.

**IV. CONCLUSIONS**

To conclude, the calculated time- and temperature-depence of the photoinduced absorption for the case of a BCS-like gap is found to be in good agreement with experimental data from femtosecond time-resolved spectroscopy on the cuprate superconductor $YBa_2Cu_3O_7-\delta$ near optimum doping and the charge-density wave insulator $K_{0.3}MoO_3$. We find that time-resolved spectroscopy can very effectively distinguish between QP states and localized states in the gap. A rather surprising feature of the data is the remarkable separation of the QP response on the femtosecond timescale and the slow response of intra-gap state relaxation on the scale of 100s of nanoseconds. In both materials we find a significant intra-gap density of states, which display very different time-dynamics and $T$-dependence than the QP states above the gap.

The authors wish to acknowledge the ULTRAFAST network and the Ministry for Science and Technology of Slovenia for supporting part of this work.

**V. APPENDIX**

Analytic solution of Eq. (3) has the form:

$$n_L(t) = \frac{2N_{L0}C \exp(-t/\tau)}{1 - C \exp(-t/\tau)}.$$

(17)

Here $1/\tau = 2N_{L0}R + \tilde{\gamma}$. Constant $C$ can be found from the following equation (see also Eq.(6)):

$$n_L(0) - n_L(t_r) = \eta n_{QP}.$$ 

(18)

Combining this two equations one obtains the following form for $n_L(0)$:
This solution reduces to Eq.(8) if \( n_L(0) = N_{LO}(1 + \frac{\eta}{2RN_L}) \)

\[
\left[ 1 + \frac{\eta Q P}{N_{LO}^2(1 + \frac{\eta}{2RN_L})^2 R_T} - 1 \right]
\]

(19)

This solution reduces to Eq.(8) if \( \frac{\eta Q P}{N_{LO}(1+\frac{\eta}{2RN_L})^2 R_T} \gg 1 \) and to Eq.(10) in the opposite limit.

**Figure 1.** A photoinduced transmission signal \( \Delta T/T \) as a function of time \( t \) after photoexcitation in YBa\(_2\)Cu\(_{3}\)O\(_{7-\delta} \) (\( T_c = 90\) K) taken at \( T = 80 \) K (points), together with the fit (solid line). (1) is the baseline signal with no pump applied. (2) is the long-lived signal pile-up remaining from previous pulses, (3) (dashed line) is the signal due to QP recombination and (4) (dotted line) is the long-lived signal remaining after all the QP signal has decayed. Signal pile-up contribution and the single pulse contribution to the slow photoinduced signal are given by \( A \) and \( \alpha_0 \) respectively.

**Figure 2.** a) The pump-probe optical diagram. 1 represents the pump pulse exciting charge carriers into a higher-lying band, 2 the carriers rapidly relax their energy to states near the Fermi energy. 3 represents the probe pulse. b) A schematic diagram of the terms contributing to the relaxation of intra-gap states in a superconductor with a gap \( 2\Delta \). In the case of a CDW gap, the \( e - h \) pairs take the place of Cooper pairs, so the rate equation remains the same.

**Figure 3.** a) The temperature dependence of the photoinduced absorption from localized states in YBa\(_2\)Cu\(_{3}\)O\(_{6.9}\) taken from Ref. [3] (open circles) and Ref. [11] (solid circles). The solid line is a plot of expression (8) with \( \Delta_c(0)/k_BT_c = 5 \) whereas the dashed line represents a general solution (Eq. 19) with non-zero \( \gamma \) term and additional fourth order term in the expansion \( R \) in powers of \( \Delta_c \). b) The temperature dependence of the photoinduced reflection in K\(_{0.3}\)MoO\(_3\) from Ref. [12] (open circles) compared with the model fit using Eq.(8) with \( \Delta_c(0)/k_BT_c = 4.8 \) (solid line). The general solution (Eq.(19)) is represented by the dashed line. c) The calculated temperature dependence of the photoinduced absorption from localized states in case of \( T \)-independent pseudogap Eq.(13) with \( \Delta^0/k_BT_c = 8 \) and different \( \alpha/T \) ratios.

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Figure 1

Time [ps]

- $\Delta T/T [10^{-4}]$

(1) $A$

(2) $A + a_o$

(3)

(4)
Fig. 2

Diagram a)

Diagram b)
Fig. 3

(a) $\Delta R/R$ [norm.] vs. Temperature [K]

(b) $\Delta R/R$ [norm.] vs. Temperature [K]

(c) $\Delta R/R$ [norm.] vs. $T/T_c$ for different $\alpha/\Gamma$ values:
- $\alpha/\Gamma = 0$
- $\alpha/\Gamma = 1$
- $\alpha/\Gamma = 10$
- $\alpha/\Gamma = 100$