Comments on a gapless high-T_c superconductivity

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The zero-bias conductance peaks observed on the tunneling current-voltage characteristics of high-T_c superconductors may stem from the onset of a gapless region and a strong gap dispersion near the Fermi surface. These effects would in turn stem from an inelastic relaxation of quasiparticles due to phonons or low-frequency collective excitations of the charge density (acoustic plasmons).

1. The differential current-voltage characteristics of tunnel junctions and point contacts formed from the cuprate metal oxides indicate that the quasiparticle density of states is nonzero at any voltage V and temperatures T below the superconducting transition temperature T_c of the compound.[1,2] In addition, narrow peaks in the conductance dI/dV are often observed[3–6] near the point $V = 0$ at $T \ll T_c$. These peaks are independent of the magnetic field and thus unrelated to the flow of a Josephson current.

In the present study we have shown that the reason for these "zero-bias" peaks in the density of states (a "second gap") may be a strong dispersion of the complex gap parameter near the Fermi surface. This dispersion would result from an inelastic relaxation of quasiparticles due to phonons or damped collective low-frequency excitations of the charge density in a system of nearly localized carriers: acoustic plasmons. These plasmons exist in metals with overlapping wide and narrow bands and may play an important role in high-T_c superconductivity.[7]

2. As was shown in Ref. 8 on the basis of the Éliashberg equations[9] for superconductors with a strong electron–phonon interaction, the damping of quasiparticles near the Fermi surface due to a relaxation involving acoustic phonons with a damping rate $\gamma_{ph} \simeq \lambda_{ph} T^3/\Theta_D^2$ (λ_{ph} is the constant of the electron–phonon coupling, and Θ_D is the Debye temperature) leads to the vanishing of the real and imaginary parts of the renormalized complex gap parameter $\Delta(\omega)$ as $\omega \to 0$:

$$\mathrm{Re}\Delta(\omega) = \frac{\omega^2}{\gamma_{ph}^2}\Delta_0; \qquad \mathrm{Im}\Delta(\omega) = -\frac{\omega}{\gamma_{ph}}\Delta_1, \tag{1}$$

where Δ_0 and Δ_1 are parameters that depend on T (more on this below).

In the new high-T_c superconductors based on layered cuprate metal oxides, the resistivity is found[10,11] to be an approximately linear function of the temperature T, $\rho(T) \approx \rho_0 + \rho_1 T$, over a wide temperature range. The implication is a carrier relaxation mechanism with a reciprocal lifetime (damping rate) which is a nearly linear function of T:

$$\tau^{-1}(T) = \gamma(T) \simeq A + BT. \tag{2}$$

According to Ref. 7, this is the T dependence which would be characteristic of the quasiparticle damping rate $\gamma_{pl}(\omega,T)$ in the limit $\omega \to 0$ if the damping were due to an inelastic relaxation involving acoustic plasmons, and if the acoustic plasmons were damped by the Landau mechanism, in a process involving nearly localized, nondegenerate heavy (h) carriers in a narrow band, and by Drude damping, as a result of an elastic scattering of degenerate, light (l) carriers in a wide band by lattice defects and impurity centers with a time scale $\tau_l = \text{const}$. The elastic scattering of the h carriers is suppressed because these carriers localize near lattice sites ($\tau_h \gg \tau_l$), and the quantum Landau damping by the l carriers is slight in the low-frequency region $|\omega| \ll \Omega_{pl}$, where Ω_{pl} is their plasma frequency ($\Omega_{pl} \approx 1$ eV).

As was shown in Ref. 7, an exchange of virtual acoustic plasmons, which hybridize with optical phonons, may lead to a Cooper pairing of the l carriers. An electron–plasmon coupling of this sort is particularly effective if the strong-coupling approximation is valid for the h carriers and if the spectrum of the acoustic plasmons is, like the phonon spectrum, a periodic function of the quasimomentum.

Since the imaginary part of the kernel of the electron–plasmon coupling in the Éliashberg equations,[9] which stems from the Landau and Drude damping of virtual acoustic plasmons, is, as in the case of the electron–phonon coupling,[8] proportional to ω, the complex gap parameter $\Delta(\omega)$ in the l-carrier spectrum satisfies conditions (1) at $T < T_c$ as $\omega \to 0$ (with γ_{ph} replaced by γ_{pl}).

3. The quasiparticle density of states of a superconductor[12] of a complex gap $\Delta(\omega)$ is ($\omega > 0$)

$$N_s(\omega) = N_n(0)\mathrm{Re}\left[\frac{\omega}{\sqrt{\omega^2 - \Delta^2(\omega)}}\right] = N_n(0)\omega\left\{\frac{\sqrt{K^2(\omega) + L^2(\omega)} + K(\omega)}{2[K^2(\omega) + L^2(\omega)]}\right\}^{1/2}, \tag{3}$$

where

$$K(\omega) = \omega^2 - [\mathrm{Re}\Delta(\omega)]^2 + [\mathrm{Im}\Delta(\omega)]^2; \quad L(\omega) = 2\mathrm{Re}\Delta(\omega)\mathrm{Im}\Delta(\omega); \tag{4}$$

and $N_n(0)$ is the density of states of the l carriers in the normal state on the Fermi surface. Substituting (1) into (3), and using (4), we find an expression for the density of states near the Fermi surface

$$N_s(\omega) \approx N_n(0) \left\{ \frac{\sqrt{(1-\omega^2\tilde{\Delta}_0^2+\tilde{\Delta}_1^2)^2+4\omega^2\tilde{\Delta}_0^2\tilde{\Delta}_1^2}+(1-\omega^2\tilde{\Delta}_0^2+\tilde{\Delta}_1^2)}{2[(1-\omega^2\tilde{\Delta}_0^2+\tilde{\Delta}_1^2)^2+4\omega^2\tilde{\Delta}_0^2\tilde{\Delta}_1^2]} \right\}^{1/2}, \tag{5}$$

where $\Delta_0 = \Delta_0/\gamma_0^2$, $\Delta_1 = \Delta_1/\gamma_0$, and either $\gamma_0 = \gamma_{ph}$, in the case of the electron–phonon coupling, or $\gamma_0 = \gamma_{pl}$ in the case in which the electron–plasmon coupling is predominant as $\omega \to 0$. It follows from (5) that in the gapless region $\omega < \gamma_0$, in which the gap dispersion is determined by (1), the quasiparticle density of states has a maximum at the point

$$\omega_0(T) = \frac{\gamma_0(T)}{\Delta_0(T)}[\Delta_1^2(T)+\gamma_0^2(T)]^{1/2}, \tag{6}$$

if $\Delta_0^2(T) > \Delta_1^2(T) + \gamma_0^2(T)$. Outside the gapless region, where the real and imaginary parts of the gap take on constant finite values $\mathrm{Re}\,\Delta(\omega) \approx \Delta_0$ and $\mathrm{Im}\,\Delta(\omega) \approx \Delta_1$, according to (3) and (4), we find

$$N_s(\omega) \approx N_n(0)\omega \left\{ \frac{\sqrt{(\omega^2-\Delta_0^2+\Delta_1^2)^2+4\Delta_0^2\Delta_1^2}+(\omega^2-\Delta_0^2+\Delta_1^2)}{2[(\omega^2-\Delta_0^2+\Delta_1^2)^2+4\Delta_0^2\Delta_1^2]} \right\}^{1/2}. \tag{7}$$

It follows that the density of states has a second maximum at the point

$$\omega_1(T) = [\Delta_1^0(T) - \Delta_1^2(T)]^{1/2}, \tag{8}$$

which corresponds to the primary gap feature of the density of states $(\omega_2 > \omega_0)$.

Figure 1 (a–c) shows curves of the ratio $N_s(\omega)/N_n(0)$ versus ω/T_c, which are symmetric with respect to the point $\omega = 0$, for various temperatures $T < T_c$ and for the following approximation of the quasiparticle damping rate:

$$\gamma_{pl}(\omega, T) = A + BT + \frac{C^2\omega^2}{(Q + C|\omega|)}. \tag{9}$$

This approximation uses a γ_{pl} quadratic in ω in the low-frequency region and a γ_{pl} linear in $|\omega|$ at high frequencies.[13] It also incorporates a linear dependence of γ_{pl} on T at $\omega = 0$ [see Eq. (2)]. The constants A, B, C, and Q in (9) were selected through a comparison with experimental data for the static resistivity[10] $\rho(T)$ and the optical conductivity[13] $\sigma(\omega)$. The constant A is related to the residual resistivity ρ_0 and describes the contribution to $\gamma_{pl}(\omega,T)$ from Drude damping by l carriers $(A \simeq \tau_l^{-1})$.

In calculating $N_s(\omega)$ from (3), with the help of (4), we used model functions for the real and imaginary parts of $\Delta(\omega)$. In the limit $\omega \to 0$, these functions agree with the asymptotic behavior in (1), while in the region $\omega > \gamma_{pl}$ they reconstruct a finite gap in the quasiparticle spectrum:

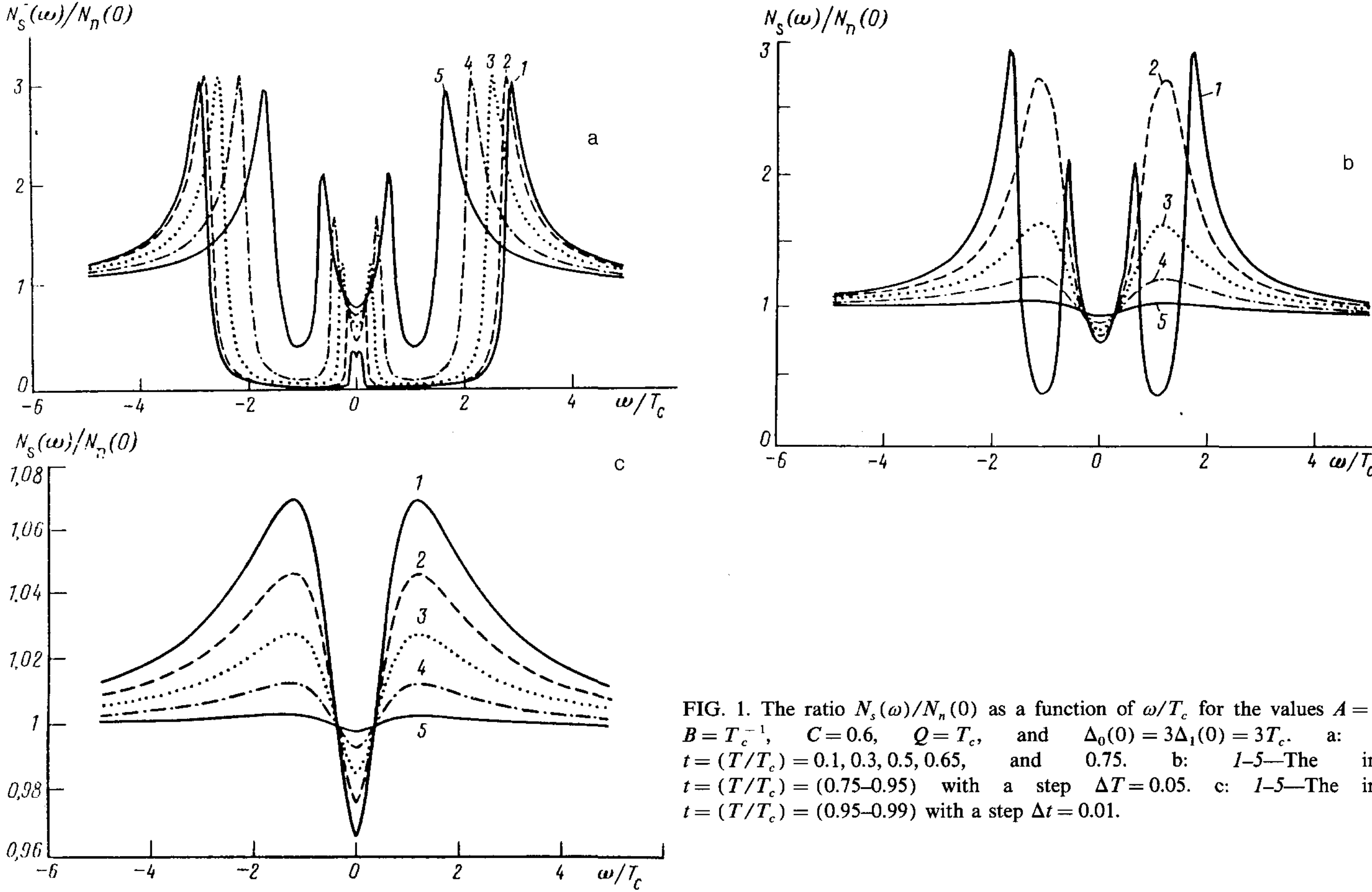

FIG. 1. The ratio $N_s(\omega)/N_n(0)$ as a function of ω/T_c for the values $A = 0.2T_c$, $B = T_c^{-1}$, $C = 0.6$, $Q = T_c$, and $\Delta_0(0) = 3\Delta_1(0) = 3T_c$. a: *1–5*—$t = (T/T_c) = 0.1, 0.3, 0.5, 0.65$, and 0.75. b: *1–5*—The interval $t = (T/T_c) = (0.75–0.95)$ with a step $\Delta T = 0.05$. c: *1–5*—The interval $t = (T/T_c) = (0.95–0.99)$ with a step $\Delta t = 0.01$.

$$\mathrm{Re}\Delta(\omega,T) = \Delta_0(T)\left[1-\exp\left\{-\frac{\omega^2}{\gamma_{pl}^2(\omega,T)}\right\}\right]; \tag{10}$$

$$\mathrm{Im}\Delta(\omega,T) = -\Delta(T)\left[1-\exp\left\{-\frac{|\omega|}{\gamma_{pl}(\omega,T)}\right\}\right]\mathrm{sign}\ \omega. \tag{11}$$

The T dependence of Δ_0 and Δ_1 was approximated by the function

$$f(t) = [1 - t^{4<1+t^4>}], \qquad t = T/T_c, \tag{12}$$

which is approximately the same as the BCS curve[12] for a gap. The values of the parameters $\Delta_0(0)$, $\Delta_1(0)$, and γ_{pl} were chosen to satisfy the condition for the existence of a low-frequency maximum in the density of states at $T=0$: $\omega_0(0) < \gamma_{pl}(0,0)$. The value of $\omega_0(0)$ was taken to be $3/T_c$, with allowance for strong-coupling effects.

It can be seen from Fig. 1a that there are two symmetric peaks in the density of states inside the gap at $T \ll T_c$. As T is increased, they shift in the direction opposite that of the primary gap features and simulate the appearance of a second gap. The size of this second gap increases with increasing T, in agreement with the experiments of Ref. 4. The existence of additional low-frequency peaks in $N_s(\omega)$ near the point $\omega = 0$ as $T \to 0$ stems from the gap dispersion near the Fermi surface [see Eqs. (10) and (11)] if $\gamma_{pl}(0,0) = A \neq 0$. Their position is determined by the ω dependence of the damping rate $\gamma_{pl}(\omega,0)$ [see Eq. (9)]. This dependence is manifested in the region $|\omega| > \sqrt{AQ}/C$, in which there is a crossover from Drude damping by l carriers to Landau damping by h carriers. In ideal single crystals, in which there is essentially no Drude damping ($\tau_l \to \infty$), the damping rate exhibits a behavior $\gamma_{pl}(\omega,0) \simeq \omega^2$ at $T=0$

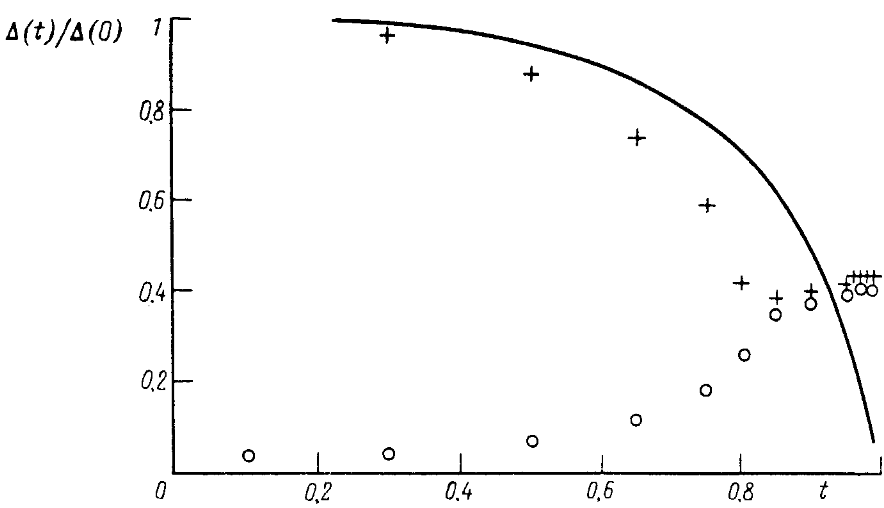

FIG. 2. Positions of the $N_s(\omega)$ peaks versus $t = T/T_c$. Solid line—The function $f(t) = [1 - t^{4(1+t^4)}]$; plus signs —primary gap feature; circles—zero-bias peak in the density of states.

as $\omega \to 0$. The low-frequency peaks in the density of states thus shift toward the point $\omega = 0$ as $T \to 0$ according to (6), and they merge to form a single narrow peak. The height of the latter peak increases as $\omega^{-1/2}$, while its width vanishes, $\Delta\omega \to 0$, since the gapless region on the Fermi surface becomes covered, and we have $\mathrm{Re}\,\Delta(0,0) \neq 0$.

These arguments also hold in the case of an inelastic relaxation of quasiparticles by phonons with a damping rate $\gamma_{ph}(\omega,T) = aT^3 + b\omega^2$, but in this case the splitting of the zero-bias peak in the density of states and the shift of the low-frequency peaks with increasing T occur in a T^3 fashion, according to (6). This prediction is at odds with the experiments of Ref. 4. On the other hand, a peak in the density of state has been observed[3,6] on the differential current-voltage characteristics at $V = 0$ and low values of T. That peak apparently corresponds to a relaxation of carriers due to acoustic plasmons with a very weak Drude damping.

Figure 2 shows the positions of the $N_s(\omega)$ peaks versus T for the curves in Fig. 1 (a–c). We see that the frequency $\omega_1(T)$, which corresponds to the primary gap features in the density of states, falls off with increasing T slightly more rapidly than the BCS curve does.[12] Near T_c, the gap feature merges with the zero-bias peak, $\omega_0(T)$. Its position becomes nearly independent of T, all the way up to T_c, because of the increase in $\gamma_{pl}(\omega,T)$ with increasing T. This anomaly in the temperature dependence of the gap has been observed in optical experiments, on IR reflection.[13]

A more detailed analysis of the numerous experimental data and a comparison of these data with the theory will make it possible to identify the quasiparticle relaxation mechanism in the cuprate metal oxides. It will also shed light on the nature of high-T_c superconductivity.

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