LA-UR 93-1354

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TITLE: EFFECT OF SPONTANEOUS DECAY OF SUPERCONDUCTOR PARTICLES IN THE TUNNELING DENSITY OF STATES

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SUBMITTED TO: Conference on Spectroscopies in Novel Superconductors, Santa Fe, NM, March 17-19, 1993

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Effect of Spontaneous Decay of Superconductor Quasiparticles in the Tunneling Density of States

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ABSTRACT - Superconductivity has been successfully described with either the Landau-Ginsburg theory of second order phase transitions or with strong-coupling versions of the original BCS theory for almost fifty years. Recent tunnelling and photoemission data on the cuprate oxide superconductors may now provide evidence of corrections to the mean field approximation. It has been shown by Zasadzinski et al.[a][1992] that there is a dip at $eV \simeq 3\Delta_0$ in the SIS tunneling conductance, which is the derivative of the current across a superconductor-insulator-superconductor junction with respect to the applied voltage, for a set of cuprate superconductors whose T_c's range from 5.5K to 100K. Recently L. Coffey and I[Coffey and Coffey 1993] proposed an explanation of this feature in terms of the spontaneous decay of mean field quasiparticles. We showed that corrections to the mean field approximation for a superconductor lead to different frequency thresholds for spontaneous quasiparticle decay with different superconductor order parameter symmetries. These effects lead to features in the superconductor density of states and in the SIS tunneling conductance and provide experimental evidence of d-wave symmetry for the superconductor order parameter in the cuprates. I discuss this model and also evidence of quasiparticle decay in ARPES data on $Bi_2Sr_2CaCu_2O_8$.

Keywords: Superconductivity, Tunneling, Quasiparticles

I. Introduction

One of the earliest experiments which supported the predictions of the BCS theory were the tunneling experiments of Giaever[1960]. These experiments showed the piling up of the density of states, as measured in SIN tunneling conductance, at Δ_0 for $T << T_c$. Later developments verified the results of more sophicated treatment of the electron-phonon interactions within the Eliashberg formalism and lead to a consistent account of the phonon density of states together with the electron-phonon coupling and the tunneling conductance.[Schrieffer et al. 1963] This analysis of the electron-phonon coupling relies on features in $g_{SIN}(eV) = dI/dV$ at bias voltages such that $eV = \Delta_0 + \omega_{ph}$ where Δ_0 is the position of the peak in $g_{SIN}(eV)$ and ω_{ph} are characteristic phonon frequencies in the superconductors.[McMillan and Rowell 1969] These features are dips whose depth is a few % of the BCS value of $g_{SIN}(eV)$ at that bias and are about four times larger in the tunneling conductance of superconductor-insulator-superconductor junction, $g_{SIS}(eV)$. Zasadzinski et al.[a][1992] found a dip feature in the $g_{SIS}(eV)$ curves for a number of cuprate superconductors with values of T_c ranging from 5K to 100K which scales with Δ_0, other than the peak at $2\Delta_0$ expected from the mean field approximation. This also shows up when g_{SIS} is generated numerically using measured tunneling conductances on superconducting-insulator-normal metal (S-I-N) junctions, g_{SIN}. The analysis of Zasadzinski et al.[a][1992], shown in Figure 1, is the first to identify the feature in g_{SIS} with a superconductor energy scale. The same feature is present in the data of other groups on the BSCCO system[Wnuk et al. 1991, Chen and Ng 1992] and has been shown to go away as the temperature increase above T_c.[Mandrus et al. 1991, Mandrus et al. 1993]

This feature at $eV \simeq 3\Delta_0$ in $g_{SIS}(eV)$ is consistent with the dip feature seen in the angle resolved photoemission data(ARPES) on $Bi_2Sr_2CaCu_2O_8$.[Dessau et al. 1991, Hwu et al. 1991, Shen et al. 1993] This consistency between $g_{SIN}(eV)$ and ARPES data has been pointed out before by Zasadzinski et al.[b][1992] and by Mandrus et al [1993] ARPES measures the spectral density for energies below the Fermi Surface and from it the density of states, $g_{SIN}(eV)$, may be calculated. In the ARPES data there is a dip along one direction at $\omega \simeq 2\Delta_0$. The anisotropy of the dip leads to it being reduced by averaging when the $g_{SIN}(eV)$ is calculated from the spectral density. The resulting feature at $eV \simeq 2\Delta_0$ in $g_{SIN}(eV)$ leads to the dip at $eV \simeq 3\Delta_0$

in the $g_{SIS}(eV)$ generated from the $g_{SIN}(eV)$ curves. In our calculation[Coffey and Coffey 1993], the dip in g_{SIS} is a consequence of deviations from weak coupling mean field behavior of the superconductivity in these materials. Here we argue that the dip seen in g_{SIS} is a consequence of these effects and that the value of the biasing across the junction at which it occurs points to the conclusion that the superconducting order parameter in the cuprates is d-wave.

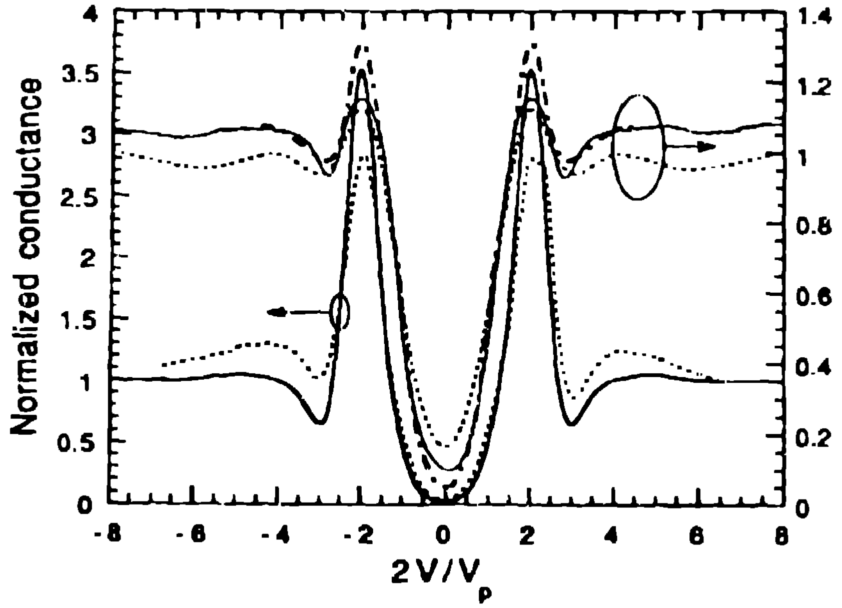

Figure 1 Scaling of the Dip Feature(from Zasadzinski et al.[8][1992]). g_{SIS} curves for various cuprate superconductors (T_c from 5.5K to 100K) plotted on a voltage axis scaled in units of Δ_0. Left scale: $Nd_{2-x}Ce_xCuO_{4-y}$(T_c =23K, solid line), $Bi_2Sr_2CaCu_2O_8$(T_c =86K, dashed line). Right scale: $Tl_2Ba_2CaCuO_8$(T_c =100K, dashed-dot line), $Bi_2Sr_2Cu_1O_6$(T_c =5.5K, solid line), BSCCO film(dashed line).

One consequence of deviations from a weak coupling mean field treatment of superconductivity is that the mean field quasiparticles spontaneously decay which leads to changes in the density of states.[D. Coffey 1990] These decay processes are characterised solely by the gap function, $\Delta(k)$, and the quasiparticle excitation spectrum. Consequently they are a probe of the nature of the order parameter.

II. Spontaneous Quasiparticle Decay

The starting point of the present analysis is a Hamiltonian, Eq.(1), describing a system of fermions interacting via a potential, $U(\vec{q})$.

$$H = \sum_{k,\sigma} \xi_k c^\dagger_{k\sigma} c_{k\sigma} + \sum_{k,k',q} U(q) c^\dagger_{k+q,\sigma} c^\dagger_{k'-q,\sigma'} c_{k',\sigma'} c_{k,\sigma} \quad (1)$$

where $\xi_k = -2t(\cos(k_x a) + \cos(k_y a)) - \mu$. a is the lattice spacing, μ is the chemical potential and t is the hopping matrix element. One assumes that the ground state of the system at low temperatures is superconducting, the nature of which depends on $U(\vec{q})$. $\Delta(\vec{k})$ and the quasiparticle operators, $\gamma_{k\sigma}$ and $\gamma^\dagger_{k\sigma}$, are determined by the weak coupling approximation for the gap equation and the Hamiltonian is written in terms of these operators. The Hamiltonian becomes

$$
\begin{aligned}
H = \sum_{\vec{k}\sigma} E_{\vec{k}} \gamma^{\dagger}_{\vec{k}\sigma} \gamma_{\vec{k}\sigma} + \sum_{\vec{k},\vec{k}',\vec{q},\alpha,\beta} U(\vec{q}) \Big[& H_A(\vec{k}_1,\vec{k}_2,\vec{k}_3,\vec{k}_4,\alpha,\beta) \gamma^{\dagger}_{\vec{k}_1\alpha} \gamma^{\dagger}_{\vec{k}_2\beta} \gamma^{\dagger}_{\vec{k}_3\alpha} \gamma_{\vec{k}_4\alpha} \\
& + H_B(\vec{k}_1,\vec{k}_2,\vec{k}_3,\vec{k}_4,\alpha,\beta) \gamma^{\dagger}_{\vec{k}_1\alpha} \gamma^{\dagger}_{\vec{k}_2\beta} \gamma^{\dagger}_{-\vec{k}_3\alpha} \gamma^{\dagger}_{\vec{k}_4\beta} \\
& + H_C(\vec{k}_1,\vec{k}_2,\vec{k}_3,\vec{k}_4,\alpha,\beta) \gamma^{\dagger}_{\vec{k}_1\alpha} \gamma^{\dagger}_{\vec{k}_2\beta} \gamma_{\vec{k}_3\beta} \gamma_{\vec{k}_4\alpha} + h.c. \Big]
\end{aligned}
\tag{2}
$$

where $\vec{k}_1 = \vec{k} - \vec{q}$, $\vec{k}_2 = -\vec{k}' + \vec{q}$, $\vec{k}_3 = -\vec{k}'$, and $\vec{k}_4 = \vec{k}$. H_A, H_B and H_C are

$$
H_A(\vec{k}_1,\vec{k}_2,\vec{k}_3,\vec{k}_4,\alpha,\beta) = (u_{\vec{k}_1} u_{\vec{k}_4} + v_{\vec{k}_1} v_{\vec{k}_4}) u_{\vec{k}_2} v_{\vec{k}_3} \delta_{\alpha,-\beta}, \tag{3a}
$$

$$
H_B(\vec{k}_1,\vec{k}_2,\vec{k}_3,\vec{k}_4,\alpha,\beta) = u_{\vec{k}_1} u_{\vec{k}_2} v_{\vec{k}_3} v_{\vec{k}_4} \delta_{\alpha,-\beta}, \tag{3b}
$$

and

$$
\begin{aligned}
H_C(\vec{k}_1,\vec{k}_2,\vec{k}_3,\vec{k}_4,\alpha,\beta) = & (u_{\vec{k}_1} u_{\vec{k}_3} + v_{\vec{k}_1} v_{\vec{k}_3})(u_{\vec{k}_2} u_{\vec{k}_4} + v_{\vec{k}_2} v_{\vec{k}_4}) \delta_{\alpha,-\beta} \\
& + u_{\vec{k}_1} v_{\vec{k}_2} v_{\vec{k}_3} u_{\vec{k}_4} \delta_{\alpha,\beta}
\end{aligned}
\tag{3c}
$$

The first term in Equation(1) has been replaced by the pairing Hamiltonian which incorporates all the weak coupling mean field contributions from the interaction term. Consequently in calculationg the effect of the interactions in Equation(2) all contributions which include mean field contributions should be dropped to avoid double-counting. The residual interactions lead to corrections to the weak-coupling mean field approximation. Here these deviations from this mean field approximation are calculated to second order in the interactions. The contributions to the self-energy, $\Sigma(\vec{k},\omega)$, of the $\gamma^{\dagger}_{\vec{k}\sigma}\gamma_{\vec{k}\sigma}$ propagator are shown in Figure 2

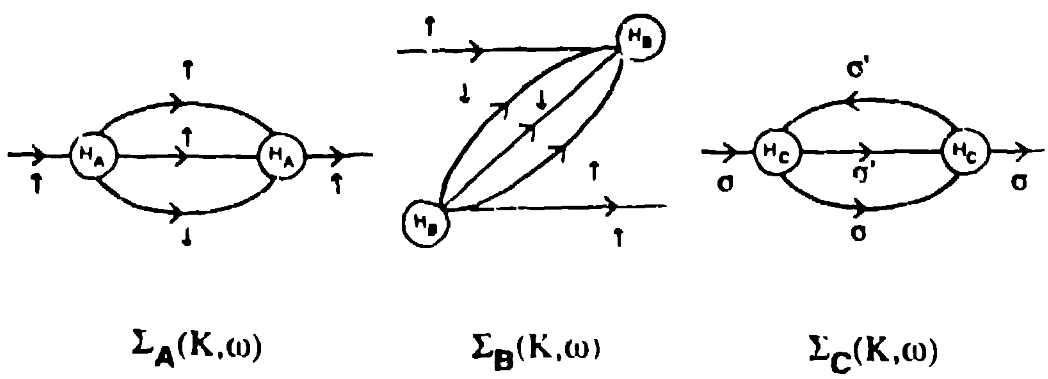

Figure 2 Second order contributions to the self-energy of the $\gamma^{\dagger}_{\vec{k}\sigma}\gamma_{\vec{k}\sigma}$ propagator

Taking $U(q) = U_o$ the gap equation gives $\Delta(\vec{k}) = \Delta_0$, an s wave superconductor. If $U(q)$ is assumed to be sharply peaked at $\vec{Q} = (\pi/a, \pi/a)$,

$$
U(q) = \frac{U_0}{1 + \xi^2 |q - (\pm\pi/a, \pm\pi/a)|^2} \tag{5}
$$

an approximate solution of the gap equation is $\Delta(\vec{k}) = \frac{\Delta_0}{2}(\cos(k_x) - \cos(k_y))$, a superconductor with $d_{x^2-y^2}$ symmetry. $\Sigma_C(\vec{k},\omega)$ is negligible at low temperatures for both s and d wave cases since it relies on the presence of thermal quasiparticles and I drop this contribution in the present discussion. $\Sigma_A(\vec{k},\omega)$ and $\Sigma_B(\vec{k},\omega)$ are,

$$
\begin{aligned}
\Sigma_A(\vec{k},\omega) &= \frac{1}{2} \sum_{\vec{k}',\vec{q}} U^2(q) \frac{(u_{\vec{k}} u_{-\vec{k}'-\vec{q}} + v_{\vec{k}} v_{-\vec{k}'-\vec{q}})^2 (u_{\vec{k}'} v_{\vec{k}'+\vec{q}} - v_{\vec{k}'} u_{\vec{k}'+\vec{q}})^2}{E_{\vec{k}+\vec{q}} + E_{\vec{k}'} + E_{-\vec{k}'-\vec{q}} - \omega} \\
\Sigma_B(\vec{k},\omega) &= \frac{1}{2} \sum_{\vec{k}',\vec{q}} U^2(q) \frac{(u_{-\vec{k}'} v_{-\vec{k}'-\vec{q}} - v_{-\vec{k}'} u_{-\vec{k}'-\vec{q}})^2 (u_{\vec{k}} v_{\vec{k}+\vec{q}} - v_{\vec{k}} u_{\vec{k}+\vec{q}})^2}{E_{-\vec{k}-\vec{q}} + E_{-\vec{k}'} + E_{\vec{k}'+\vec{q}} + \omega}
\end{aligned}
\tag{6}
$$

When the imaginary parts of these self-energies become non sero the mean field quasiparticles spontaneously decay.

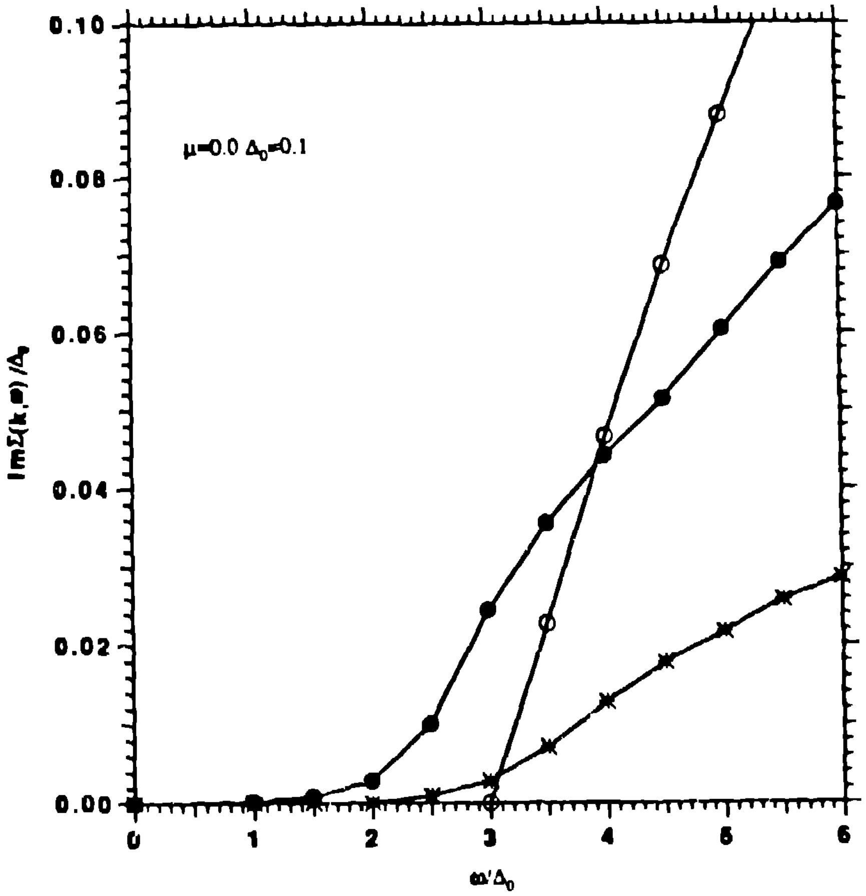

Figure 2 $\mathrm{Im}\Sigma_A(\vec{k},\omega)$ for s-wave at $\vec{k} = (\frac{k}{\sqrt{2}}, \frac{k}{\sqrt{2}})$(open circles) and for d-wave at $\vec{k} = (\frac{k}{\sqrt{2}}, \frac{k}{\sqrt{2}})$(solid circles) and at $\vec{k} = (k,0)$(stars) $\mu = 0$ and $\Delta_0 = 0.1$. In the s-wave case $U(\vec{q}) \sim -1.22$ was used and for the d-wave case $U(\vec{q})$ in Eq.(3) was used with $U_0 \sim 14.707$ and $\xi = 1$. With these parameters $\Delta_0 = 0.1$ when $\mu = 0$.

In Figure 3 $Im\Sigma_A(\vec{k}_F,\omega)$ is shown for s-wave and d-wave with $\Delta_0 = 0.1$ when $\mu = 0$ for both. From Eqs.(5) one can easily see that $Im\Sigma_A(\vec{k}_F,\omega) = Im\Sigma_B(\vec{k}_F, -\omega)$. In the s wave case $\Sigma_A(\vec{k},\omega)$ and $\Sigma_B(\vec{k}, -\omega)$ are real until $|\omega| > 3\Delta_o$ where spontaneous decay is possible [D. Coffey 1990] There are two points to be made about the d-wave results: the threshold for spontaneous quasiparticle decay is different than for the s wave case and there can be strong anisotropy in momentum space. For the $d_{x^2-y^2}$ case, $\Delta(\vec{k}) = \frac{\Delta_0}{2}(\cos(k_x) - \cos(k_y))$ has nodes and this results in a lower threshold for quasiparticle decay. As a result both $Im\Sigma_A(\vec{k},\omega)$ and $Im\Sigma_B(\vec{k},\omega)$ increase rapidly for $|\omega| > 2\Delta_0$ and go over to a linear dependence on ω at this level of approximation. Monthoux and Pines[1992] have used an interaction very similar to Eq. (4) and solved the Eliashberg equations to get $d_{x^2-y^2}$ superconductivity. In their calculations this linear dependence on ω is suppressed and $Im\Sigma(\vec{k},\omega)$ eventually falls as ω increases [Monthoux and Pines 1992] The magnitudes of $Im\Sigma_A(\vec{k},\omega)$ with these parameters are $\sim 0.03\Delta_0$ at $\omega = 3\Delta_0$. The magnitude of $Im\Sigma(\vec{k},\omega)$ increases with increasing Δ_0. It also depends on μ and ξ. The more sharply peaked $U(\vec{q})$ is the more $Im\Sigma(\vec{k},\omega)$ increases as $\mu \to 0$. For the value of ξ used here $Im\Sigma_A(\vec{k}_F,\omega)$ has a sharper turn on at $\omega \sim 2\Delta_0$ but the magnitude is practically unchanged as μ goes from 0 to -0.1. From Figure 3 one sees that there is an anisotropy in the d wave case between the direction in which $\Delta(\vec{k}) = 0$ and the direction in which $\Delta(\vec{k}_F) = \Delta_0$ I will discuss this below in the context of angle resolved photoemission spectroscopy. The most important point here is that the microscopic calculations clearly show that spontaneous quasiparticle decay becomes appreciable at ω's smaller by $\sim \Delta_0$ in the d-wave case than in the s-wave case. We now show that the switching on of this decay channel may be seen in the tunneling conductance.

III. Phenomenological Model

The current across an insulating barrier is given by[Schrieffer 1963]

$$I(eV) \propto \sum_{\vec{k}_L,\vec{k}_R} \int_{-eV}^{0} d\omega A^R(\vec{k}_R,\omega) A^L(\vec{k}_L,\omega + eV) |T_{\vec{k}_L,\vec{k}_R}|^2 \tag{6}$$

where $A^{R,L}(\vec{k},\omega)$ is the spectral density on either side of the barrier. Assuming that the tunneling matrix element, $|T_{\vec{k}_L,\vec{k}_R}|^2$, does not depend on momenta on the left and right hand sides of the barrier, $\vec{k}_L$ and $\vec{k}_R$, the current becomes

$$I(eV) \propto \int_{-eV}^{0} d\omega N^R(\omega) N^L(\omega + eV). \tag{7}$$

$N^{(R,L)}(\omega)$ are the densities of states of the charge carriers on the left and right hand sides of the insulating barrier, which for $\omega > 0$ is,

$$N^L(\omega) = \sum_{\vec{k}} A^L(\vec{k},\omega) \simeq \frac{1}{\pi} \sum_{\vec{k}} \frac{|Im\Sigma_A(\vec{k},\omega)|}{(\omega - E_{\vec{k}} - Re\Sigma_A(\vec{k},\omega))^2 + (Im\Sigma_A(\vec{k},\omega))^2} \tag{8}$$

The effects of spontaneous quasiparticle decay are more easily seen in $g_{SIS}(eV)$ curves than in $g_{SIN}(eV)$ curves. This is because two superconductor densities of states are convoluted with each other in $g_{SIS}(eV)$, whereas a superconducting and normal density of states are convoluted with each other in $g_{SIN}(eV)$. Examining Eq.(4) one sees that there is a big contribution to the current across the junction when $\omega = -\Delta_0$ and $\omega + eV = \Delta_0$, i.e. when the bias across the junction is $eV \simeq 2\Delta_0$. In the same way the effect of the decay processes can be seen in a d-wave superconductor when $\omega \simeq -\Delta_0$ and $\omega + eV \simeq 2\Delta_0$, the frequency at which the quasiparticle decay process starts to become appreciable in a d-wave superconductor. The calculated current is a monotonic function of eV with changes in slope at $eV = 2\Delta_0$ and at $eV = 3\Delta_0$ for d-wave. Calculating $g_{SIS}(eV)$ picks out these values of eV at which the slope changes and it is the rapid increase in $\Sigma_A(\vec{k},\omega)$ at $\omega \simeq 2\Delta_0$ which is responsible for the features at $eV \simeq 3\Delta_0$. The corresponding value of eV for an s-wave superconductor is $\simeq 4\Delta_0$
In order to generate the g_{SIN} and g_{SIS} curves for a d-wave superconductor we introduce a phenomenological model in which $ImG(\vec{k},\omega)$ is replaced by a Lorentzian of the form

$$ImG(\vec{k},\omega) = \frac{1}{2\pi} \frac{\Gamma(\vec{k},\omega)}{(\omega - E_{\vec{k}})^2 + \Gamma(\vec{k},\omega)^2} \tag{9}$$

where $E_{\vec{k}}$ is a renormalised quasiparticle spectrum which is assumed to have the same form as that given by the mean field approximation but in which the parameters have been renormalised by the interactions. The qualitative features of the $\vec{k}$ and ω dependence of $\Gamma(\vec{k},\omega)$ are determined by our calculations of $Im\Sigma_A(\vec{k},\omega)$. Our model for $\Gamma(\vec{k},\omega)$, Eq. (7), has the rapid increase at $\omega \sim 2\Delta_0$ and the anisotropy in $\vec{k}$ space seen in Fig. 2. These features are present for all parameters in the microscopic calculations and survive the effects of repeated scattering

$$\Gamma(\vec{k},\omega) = \Gamma_0 + \frac{\Gamma_1}{2}\left[1 + tanh(\frac{\omega - 2.5\Delta_0}{5\Delta_0})\right]\left[1 - \frac{1}{2}(\frac{\Delta(\vec{k})}{\Delta_0})^2\right] \tag{10}$$

Γ_0 and Γ_1 are taken to be free parameters. In Figure 4 we compare (A) $g_{SIN}(eV)$ and (B) $g_{SIS}(eV)$ curves for the case where the frequency dependent decay process is taken into account, $\Gamma_0 = 0.05\Delta_0$ and $\Gamma_1 = 0.5\Delta_0$ (full line), with the case where this frequency dependence is ignored $\Gamma_0 = 0.05\Delta_0$ and $\Gamma_1 = 0$ (broken line). One sees that, as a consequence of the frequency dependent damping due to the decay of Bogoluibov quasiparticles, there is a dip in $g_{SIS}(eV)$ for $3\Delta_0 \leq |eV| \leq 4\Delta_0$. This feature has been clearly identified in the experimental data by the work of Zasadzinski et al.[a](1992). This dip is completely missing from the $g_{SIS}(eV)$ where the frequency dependent decay process is ignored (broken line). Looking at the $g_{SIN}(eV)$ curves, Figure 4(A), there is no strong feature and it would be difficult to identify the effect in experimental data.

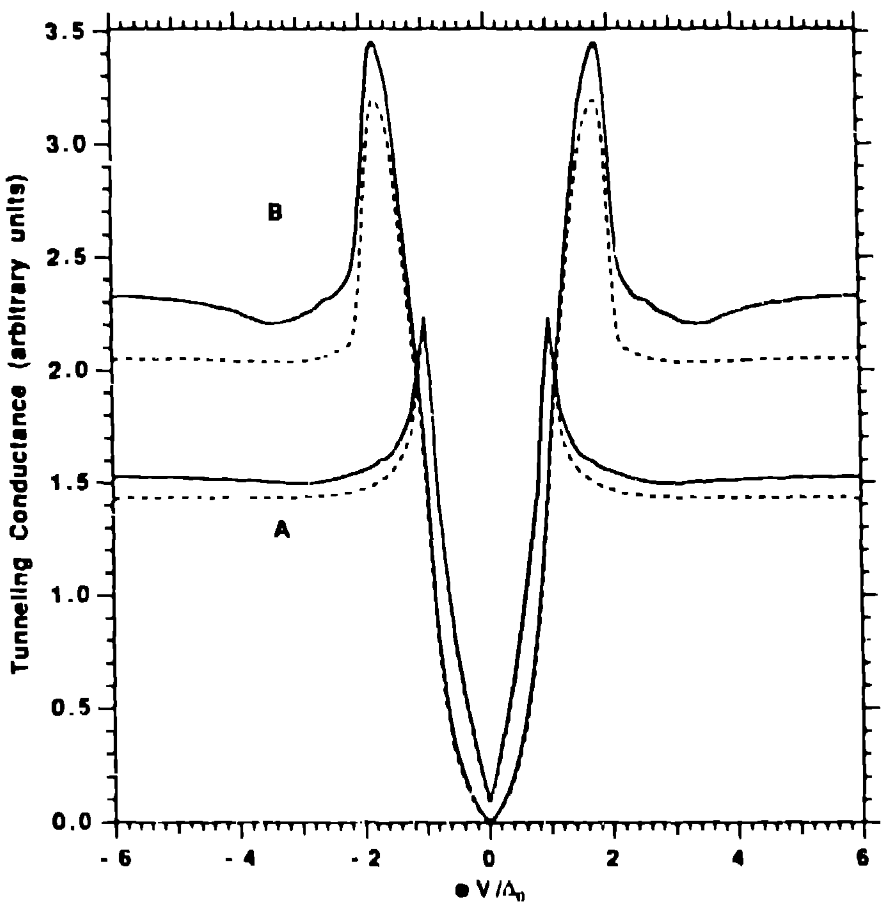

Figure 4 Tunneling conductances, (A) g_{SIN} and (B) g_{SIS}, as a function of the voltage across the junction, eV. The full line is the result with frequency dependent damping given in Eq.(5), where $\Gamma_0 = 0.05\Delta_0$ and $\Gamma_1 = 0.5\Delta_0$. The broken line is the result for a frequency independent damping, $\Gamma_1 = 0$. Here $\mu = 0.2$ and $\Delta = 0.01$.

The assumption of an s wave order parameter will always produce a dip in the g_{SIS} curves starting at $eV = 4\Delta_0$ which is not in agreement with the experimental data.(see Figure 1) The size of Γ_1 in Fig. 4 is an order of magnitude larger than the size of $Im\Sigma(\vec{k},\omega)$ found in the microscopic calculations and so our microscopic calculations are used here only to motivate our phenomenological model for $\Gamma(\vec{k},\omega)$. However this difference between the value of Γ_1 and $Im\Sigma_A(\vec{k}_F,\omega)$ is partially due to our crude approximation for the tunneling matrix element which leads to the density of states entering into the expression for the tunneling current. This has the effect of averaging over the anisotropic effect of spontaneous decay in our model which reduces its effect for a given value of Γ_1. The magnitude of $Im\Sigma(\vec{k},\omega)$ is enhanced by antiferromagnetic spin fluctuations which require going beyond second order perturbation theory and allowing for repeated scattering through H_C.[Coffey and Coffey, unpublished]

IV. Angle Resolved Photoemission

The treatment of the tunneling conductance above ignored any momentum dependence of the tunneling matrix element. As a result the density of states, an average of the spectral density, entered the expression for the current. In reality the tunneling matrix element will pick out different $\vec{k}$ states for tunneling depending on the geometry of the experiment. In making this approximation I have averaged over the anisotropy present in the microscopic calculations which shows up in the spectral density. One experiment in which the spectral density is measured directly is angle resolved photoemission(ARPES) and the data of Shen et al.[1993] presented here at this conference clearly shows this anisotropy. They have concentrated on the anisotropy of the piling up of states and so conclude that $\Delta(\vec{k})$ is very anisotropic. However their data also clearly shows the anisotropy of the dip feature. Along one direction in momentum space there is a piling up of the density of states at $\sim \Delta_0$ with a dip which appears at $\sim 2\Delta_0$. This feature in the spectral density is the signature of the spontaneous quasiparticle decay for a d-wave superconductor. Along a direction at 45° there is no evidence for a gap in $\Delta(\vec{k})$ and there is no feature at $2\Delta_0$. This is not quite what is seen in Figure 2. There, the largest values of $Im\Sigma(\vec{k},\omega)$ appear along the direction in which $\Delta(\vec{k}) = 0$ and not along the other direction. This discrepancy between our model and the experiments on $Bi_2Sr_2CaCu_2O_{8+\delta}$ turns out to depend on band structure. In the microscopic calculations discussed above, a simple tight-binding band structure was used. If next nearest neighbor hopping is put in, the $t - t'$ band structure, the self-energy becomes appreciable along both directions at $\omega \simeq 2\Delta_0$. This sensitivity to the Fermi surface is understandable given that our analysis is based on conservation of energy and momentum and consequently depends on where the low energy states close to the Fermi surface are in momentum space. So it is important to carry out our microscopic calculations with a Fermi surface which better approximates the one found by Shen et al.[1993]. The results of this calculation will be reported elsewhere.

V. Conclusions

The scaling of the dip feature in the cuprates found by Zasadzinski et al.[a][1992] clearly points to its origin being intrinsic to the superconductor state. It would be of interest to see if there is evidence for these effects in systems other than the cuprates. Since the work of MacMillan and Rowell[1968] weak features in the tunneling conductance have provided information on the electron-phonon coupling consistent with the known phonon density of states and so one would look for discrepancies between the phonon density of states and $g_{SIS}(eV)$ or for features close to eV equal to multiples of Δ_0 outside the phonon range in a survey of g_{SIS} curves for different superconductors.

In conclusion we have postulated that features seen in the tunneling conductance of S-I-S junctions at bias voltages $eV \sim 3\Delta_0$ arise from quasiparticle decay. This quasiparticle decay is a correction to the weak coupling mean field treatment of the quasiparticle states. The voltages at which these features are seen in $g_{SIS}(eV)$ have been identified with $3\Delta_0$ for a range of cuprate superconductors. This provides strong evidence for $d_{x^2-y^2}$ order parameter in the cuprate superconductors or at the minimum a very anisotropic one, where the degree of anisotropy has to be essentially the same over the range of the superconductors shown in Figure 1. This feature is unambiguous because of the high energy at which the dip is seen and is a clear signature of the nature of the order parameter. Attempts to identify the symmetry of the order parameter from data up to now have relied on power law fits to low temperature and low energy behavior of various experiments such as penetration depth, spin lattice relaxation. These attempts are plagued by extrinsic effects, such as the effects of impurities, which limit the range of the fits and can make the conclusions controversial. Even the absence of a Hebel Slichter bump in the spin lattice relaxation time of the cuprates can be accounted for with an s wave order parameter and sufficient inelastic scattering [L. Coffey 1990, Allen and Rainer 1991]. By comparison the scaling property of the dip feature discovered by Zasadzinski et al.[a][1992] is present even in films of cuprate superconductors with low T_c's and seems to be a robust feature of the tunneling conductance in the cuprates [Zasadzinski et al. 1993].

In our calculation we have concentrated on spontaneous quasiparticle decay and have taken a simple model for the normal state spectrum. In particular we have ignored the effects of interactions on this spectrum which have the effect of smoothing out sharp features associate with the tight-binding structure in the density of states. We consider a more realistic treatment of the normal state spectrum and interaction effects in a forthcoming publication.[Coffey and Coffey unpublished]

I am grateful to J. Zasadzinski et al. for permission to reproduce their data prior to publication and to L. Coffey, D. Mandrus, and J. Zasadzinski for useful conversations This work was supported by the US Department of Energy.

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