Mechanisms of heat conductivity in high-T_c superconductors

Yury Pogorelov, Miguel Angel Arranz, Raúl Villar, and Sebastián Vieira
Departamento de Física de la Materia Condensada, Universidad Autónoma de Madrid, Cantoblanco, 28049 Madrid, Spain

(Received 11 November 1994)

A phenomenological approach to the heat conductivity κ in the mixed state of a high-T_c superconductor is proposed, which permits resolution between the electronic and phononic contributions, κ_{el} and κ_{ph}. The scattering amplitude for bandlike quasiparticles vs bound states in the vortex cores has been microscopically estimated. The measurements of the in-plane thermal conductivity $\kappa(B,T)$ of single crystals $YBa_2Cu_3O_7$ and $Bi_2Sr_2CaCu_2O_8$ have been performed at several temperatures T from 4 to 110 K and magnetic fields B (along the c axis) up to 11 T. From comparison of experimental data with the phenomenologic model, the ratio κ_{el}/κ_{ph} (at $B=0$) has been extracted. This analysis shows that only κ_{el} is field sensitive and it is mostly responsible for the maximum of κ vs T, observed at $B=0$ in both materials.

The mechanism of heat conductivity κ below the superconducting transition temperature T_c in high-T_c superconductors (HTSC's) is, at the present, an object of discussion. Some controversial points of view exist about the relative importance of phononic and electronic contributions: $\kappa=\kappa_{ph}+\kappa_{el}$, mainly in order to explain the broad maximum in $\kappa(T)$ observed below T_c.[1] Widely accepted is the concept of the phononic origin of this maximum,[2] in similarity with the classical theory of heat transport in superconductors.[3] However, this model cannot explain some properties peculiar to HTSC's, such as a strong in-plane/out-of-plane anisotropy and absence of the low-temperature maximum in the out-of-plane conductivity.[4,5] Therefore an alternative approach was recently set forth,[6] claiming electronic origin of the maximum.

An additional insight into the transport properties can be obtained by scattering the heat carriers by flux lines in the mixed state of type-II superconductors[7,8] so that κ becomes a function of T and magnetic induction B. This is of particular interest for HTSC's, known as "extreme type-II" superconductors. In this paper, we present the measurements of a-b plane heat conductivity of single crystals $YBa_2Cu_3O_7$ (YBCO) and $Bi_2Sr_2CaCu_8$ (BiSCCO), as functions of T and B (normal to the plane). The field dependence $\kappa(B)$ is found to be noticeably nonlinear already at low fields, like the data known from the literature.[7,9] It is worth mentioning that the previous attempts to describe this nonlinearity[9] involved complicated empirical laws, as $\sim\exp(-dB^{1/4})$, bearing no clear physical meaning. Below, we present a simple phenomenological analysis,[10] which shows how the nonlinear $\kappa(B)$ transforms into a certain *linear* function with a slope related to the ratio $\nu=\kappa_{el}/\kappa_{ph}$ at zero field. Having estimated this ratio from our experimental data, we are able to separate the two contributions in the zero-field $\kappa(T)$. In particular, it turns out that the above mentioned maximum is mainly due to κ_{el}.[10,11]

The field effect on heat conductivity is well studied for traditional type-II materials (TT2M's), as *Nb* and its compounds and alloys, where κ first decreases linearly with B, then saturates at some low level, and is restored to its normal state value κ_n for $B\to B_{c2}$.[12–14] The heat carriers in a superconductor are the phonons and normal electronic quasiparticles and, in principle, both of them may be influenced by the magnetic flux. But it is commonly accepted that the main field effect in TT2M's is related to phonons[12,17] which dominate in the total $\kappa(T)$ at temperatures of order of or less than $T_c/2$.

The electronic structure of a vortex consists of bound quasiparticle states localized near the vortex core,[15] such that the corresponding energy levels fill the superconducting gap Δ with a separation between them of the order of $\hbar$ times ω_c, the cyclotronic frequency within the core. In TT2M's, the quasiparticle lifetime τ usually refers to the "dirty" limit;[16] $\omega_c\tau\ll 1$, so that the discrete levels overlap and the core spectrum becomes similar to that of normal state quasiparticles. On the other hand, the size of the "normal metal" region near the vortex line is about the superconducting coherence length $\xi_0\sim E_F/(k_F\Delta)$, and it exceeds the thermal phonon wavelength $\lambda_T\sim a\Theta_D/T$ at $T\gg T_d=\Theta_D\Delta/E_F$ (Θ_D is the Debye temperature, and the lattice parameters $a\sim k_F^{-1}$). Typically for TT2M's, T_d is as low as ~ 0.1 K so, within a core, the phonons can always be well characterized by their lifetime against the electron-phonon interaction in a normal metal, which provides the main field effect.[17]

The heat conductance in HTSC's differs from that in TT2M's already at $B=0$, by the above mentioned maximum in $\kappa(T)$ at some $T_m<T_c$. The field effect mainly consists in a monotonic lowering of this maximum, up to the highest accessible fields in experiment.[7,8] In contrast to TT2M's, the "superclean" limit $\omega_c\tau\gg 1$ is more likely for HTSC's at lower T,[18] and then the discrete spectrum of few bound states hardly resembles that of a normal metal. Also, T_d is as high as ~ 30 K here, so that λ_T is comparable to or greater than ξ_0 and the vortices act on the phonons rather as pointlike (inelastic) scatterers. A standard method to calculate $\kappa_{ph}(T)$ uses the Eliashberg equations[2] which, unfortunately, are too complicated to include the bound vortex states. However, the qualitative field effect in this situation can be expected to be similar

for both types of heat carriers, consisting in that the respective inverse lifetimes grow linearly with B (unless the bound states from adjacent vortices overlap, that is, for $B \ll B_{c2}$). So, phenomenologically, we can write

$$\kappa(B,T)=N_{\rm ph}(\tau_{\rm ph}^{-1}+\alpha_{\rm ph}B)^{-1}+N_{\rm el}(\tau_{\rm el}^{-1}+\alpha_{\rm el}B)^{-1} , \quad (1)$$

where N_i,τ_i,α_i are temperature dependent (the α's are proportional to the scattering amplitudes of respective heat carriers by an individual vortex). Now let us define a function $f(B,T)=B[1-\kappa(B,T)/\kappa(0,T)]^{-1}$ and insert Eq. (1) into it. This generally gives $f(B,T)=\alpha_0+\alpha_1 B+\alpha_2 B^2+\cdots$, with the temperature-dependent coefficients α_n, which are proportional to the product $\alpha_{\rm ph}\alpha_{\rm el}$ for all $n \geq 2$. Remarkably, in experiment $f(B,T)$ occurs as a *linear* function of B (see below). This suggests that either $\alpha_{\rm el}$ or $\alpha_{\rm ph}$ is zero. Supposing $\alpha_{\rm ph}=0,\alpha_{\rm el}\neq 0$, we get immediately

$$f(B,T)=(1+1/\nu)(B+B_0) , \quad (2)$$

where $B_0^{-1}=\alpha_{\rm el}\tau_{\rm el}$ and ν is the before mentioned zero-field ratio $\kappa_{\rm el}/\kappa_{\rm ph}$. Thus the electron-phonon ratio ν, as a function of temperature, can be simply inferred from the experimental slope of f vs B at given T. Typically, the estimates from the Wiedemann-Franz law at $T>T_c$ give $\nu<1$ (see, e.g., Ref. 6), and this agrees well with the values obtained below. The alternative possibility for linear f vs T, $\alpha_{\rm ph}\neq 0,\alpha_{\rm el}=0$, would change $\nu\rightarrow\nu^{-1}$ in Eq. (2) and thus would lead to an inconsistency. Simultaneously with our previous communication,[10] an expression equivalent to Eq. (1), was used by Yu *et al.*,[11] who also concluded that only the electronic part of κ is field sensitive.

Unlike the phonon channel, the field-dependent electronic scattering rate $\alpha_{\rm el}$ admits an approximate microscopic calculation. Considering the bandlike quasiparticles already coherent with the vortex system at $T=0$ (when the vortices form a regular lattice and there are no excitations at any vortex), the field-induced effect on quasiparticle lifetime can only result from thermally excited bound states.[19] Though scarcely occupied at $T \ll T_c$, they can produce a considerable effect, due to strong Coulomb interaction between the band and bound quasiparticle states (this interaction is supposed unscreened and realized through electronlike and holelike components of the corresponding Nambu spinors). We take the relevant band states as plane waves with wave vectors $\mathbf{k},k\approx k_F$, lying within the a-b plane (that is, neglecting the electronic transport along the c axis). The spinor components of Caroli-Matricon bound states[15] with energies $\varepsilon_j \approx j\hbar\omega_c, j=1,2,\ldots$, respective to the Fermi energy E_F, are modeled as sines and cosines of $k_F r$(r is the distance from the vortex axis), normalized within circles $r<r_j=\hbar v_F/\sqrt{\Delta^2-\varepsilon_j^2}$. Then a standard second order perturbation theory gives finally

$$\alpha_{\rm el}=(e^5/\hbar^2 c E_F^2)\sum_{\varepsilon_j<\Delta} n_F(\varepsilon_j)\sqrt{\Delta^2-\varepsilon_j^2} , \quad (3)$$

where n_F is the Fermi function. The reported lowest excitation energy for vortex states in YBCO corresponds to $\hbar\omega_c/k_B \approx 38$ K,[23] so the higher levels can be neglected for $T\sim T_m$, and from Eq. (3) we estimate $\alpha_{\rm el}(T_m)\sim 10^8$ G^{-1}s^{-1}. The electronic lifetime $\tau_{\rm el}$ can be roughly estimated by its normal state value τ_n (from the Drude law, extrapolated to $T<T_c$). With typical normal electric conductivity $\sigma_n\sim(10^6\ \Omega^{-1}\ {\rm K})/T$, we have $\tau_n\sim 10^{-13}$ s at $T\sim T_m$, which results in an order of magnitude estimate for the characteristic field in Eq. (2): $B_0\sim 10$ T. The data presented below give a direct check for these theory predictions and permit a conclusion about the heat transport in HTSC's.

The thermal conductivity was measured using the steady-state method with a temperature gradient applied along the a-b plane (see inset to Fig. 1). The typical sample sizes were $1\times 3\times 0.03$ mm^2 for YBCO, and $2\times 5\times 0.3$ mm^3 for BiSCCO. The sample holder was made with several pieces of copper coated substrate, joined only with two thin tubes of Pyrex glass. In each case, the heat flow due to conduction by these tubes amounted to less than 2% of the flow through the sample. Samples are attached to the copper plates with a little amount of silver epoxy to ensure a good thermal contact; for measuring the transverse gradient through the crystals, two little Speer resistors were also fixed with cryogenic varnish on copper strips and connected at intermediate positions of the sample. In order to minimize heat losses from the sample holder, low conductivity wires (Manganin) were used for all connections. Radiation losses are also avoided by mounting a copper shield around the sample. Simultaneous resistance and thermoelectric coefficient measurements indicate superconducting transitions at 92 K for YBCO and at 80 K for BiSCCO (transition width $\sim$1 K). In the absence of field,

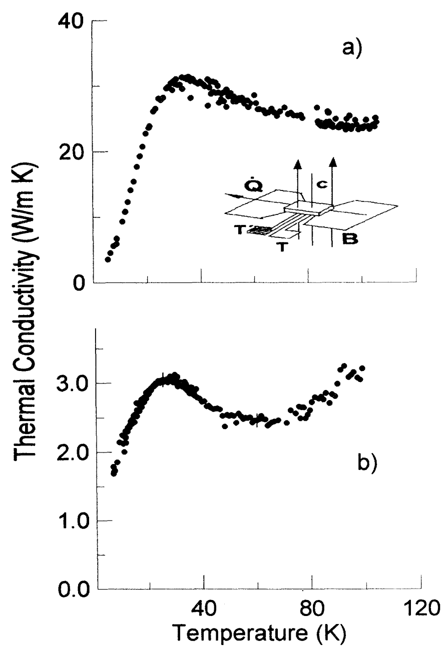

FIG. 1. Total zero-field conductivity of (a) YBCO and (b) BiSCCO specimens; the inset shows schematic of experimental setup; the experimental error bars are also marked.

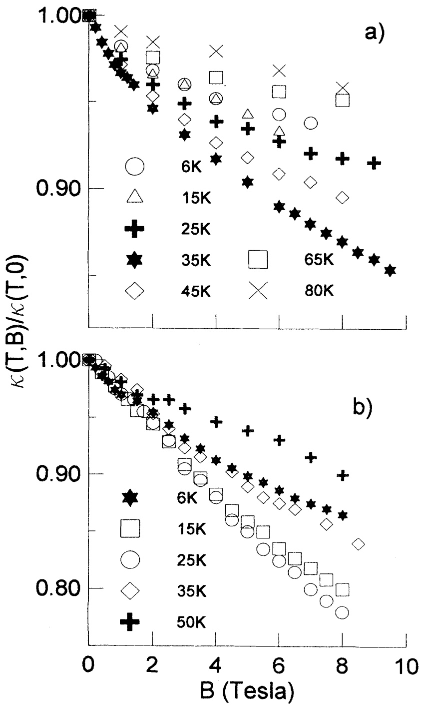

FIG. 2. The field effect on κ for the two samples of Fig. 1 at various temperatures (note smaller nonlinearity for BiSCCO, related to a greater zero-field peak).

the data on $\kappa(0,T)$ for both materials are presented in Fig. 1; they show a maximum at $T_m \sim T_c/3$. Decreasing the temperature below T_c, a sudden increase of $\kappa(0,T)$ is observed. The twinned nature of our crystals explains

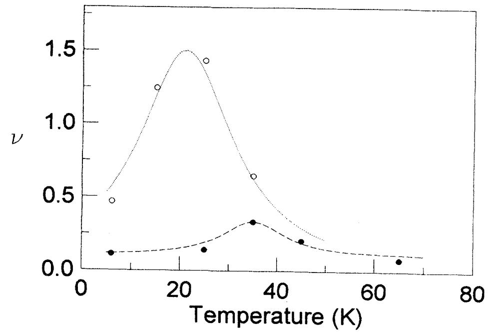

FIG. 4. Electron-to-phonon ratio ν vs temperature for YBCO (bold circles) and BiSCCO (light circles). The broken lines are guides to the eye.

their smaller relative maximum [$\kappa(0,T_m)/\kappa(0,T_c) \sim 1.25$ for YBCO vs ~ 2 for untwinned YBCO samples[8,6]]. The magnetic field was applied along the c axis. The ratio $\kappa(B,T)/\kappa(0,T)$ for several temperatures is shown in Fig. 2; it decreases continuously with B at all temperatures, revealing marked nonlinear behavior (especially for YBCO). The corresponding function $f(B,T)$ is presented in Fig. 3. Within experimental errors, it is linear in B so that, according to Eq. (2), ν can be readily obtained from its slope, and B_0 from its initial value $f(0,T)$. The values of ν, obtained from the best linear fits to experimental f for both materials, are plotted in Fig. 4. The next natural step is to use these $\nu(T)$ to "cut out" the electronic contribution from the total zero-field $\kappa(T)$. The separated κ_{ph} for both materials is traced by dashed lines in Fig. 5. They clearly demonstrate that *the zero-field peak is com-*

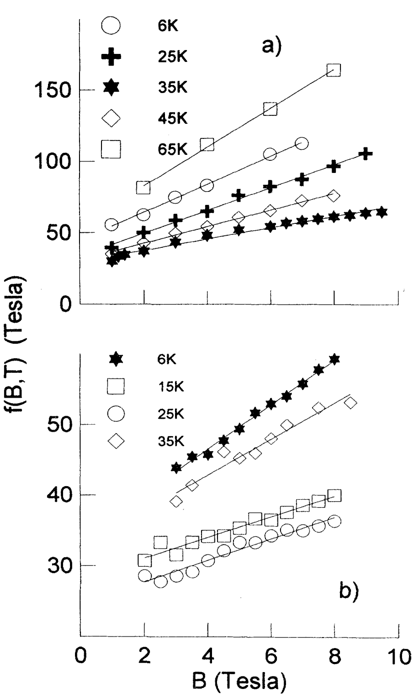

FIG. 3. Linear field dependences $f(B,T)$ obtained from the data of Fig. 2 at various temperatures.

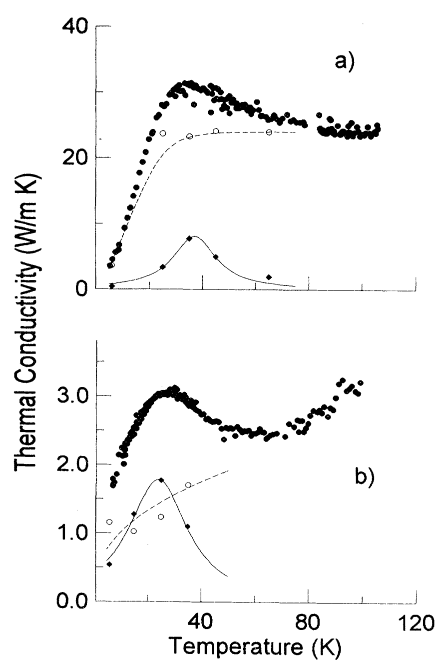

FIG. 5. Separated phononic (light circles and broken lines) and electronic (bold diamonds and solid lines) contributions to the heat conductivity of the two materials under study, as resulted from Figs. 1 and 4.

pletely due to the electronic contribution K_{el}. This central result of our study agrees well with the bulk of available experimental data on heat conductivity in HTSC's (see below).[23] Finally, the values of B_0, obtained from Fig. 3 for YBCO, are between 4 and 6 T, which is the same order of magnitude as the theoretical estimate [see after Eq. (3)]. This suggests that the proposed Coulomb scattering by bound vortex states can be a plausible mechanism of field-dependent quasiparticle relaxation in HTSC's.

In conclusion, the presented phenomenological procedure applies well to all the data on field-dependent κ in HTSC's, known by us up to the moment. Thus, all the nonlinear curves $\kappa(B)$ from the experiments[7–9] on YBCO and BiSCCO are found to result in linear $f(B)$. The following separation of particular contributions shows the low-temperature maximum to be mainly related to the electronic part κ_{el}. On the other hand, the monotonic dependence $\kappa_{ph}(T)$ obtained in Fig. 5(b) is quite similar to that concluded in a recent zero-field experiment on BiSCCO and its insulating isomorph $Bi_2Sr_2YCu_2O_8$.[24] At least, our results are in agreement with the theoretical analysis of zero-field heat conductivity,[25] which ascribes the low-temperature maximum to the electronic contribution. All these facts are in favor of the electronic mechanism of Ref. 6.

We are grateful to A. V. Inyushkin and A. N. Taldenkov for communicating to us about the results of this phenomenological procedure applied to their recent experimental data.

[1] C. Uher, J. Supercond. **3**, 337 (1990).
[2] L. Tewordt and Th. Wolkhausen, Solid State Commun. **75**, 515 (1990).
[3] J. Bardeen, G. Rickayzen, and L. Tewordt, Phys. Rev. **113**, 982 (1959).
[4] M. F. Crommie and A. Zettl, Phys. Rev. B **43**, 408 (1991).
[5] S. J. Hagen, Z. Z. Wang, and N. P. Ong, Phys. Rev. B **40**, 9389 (1989).
[6] R. C. Yu, M. B. Salomon, J. P. Lu, and W. C. Lee, Phys. Rev. Lett. **69**, 1431 (1992).
[7] V. V. Florentiev, A. V. Inyushkin, A. N. Taldenkov, O. K. Mel'nikov, and A. B. Bykov, Superconductivity **3**, 378 (1990).
[8] S. D. Peacor, J. L. Cohn, and C. Uher, Phys. Rev. B **43**, 8721 (1991).
[9] R. A. Richardson, S. D. Peacor, F. Nori, and C. Uher, Phys. Rev. Lett. **67**, 3856 (1991).
[10] The previous report of the present approach can be found in M. A. Arranz, Yu. Pogorelov, R. Villar, and S. Vieira, in *Advances on Science and Technology,* Proceedings of 8th CIMTEC, Florence, Italy, 1994, edited by P. Vincenzini (Techna, Faenza, 1995), Vol. 8, p. 422.
[11] F. Yu, M. B. Salomon, V. Kopylov, M. N. Kolesnikov, and H. M. Duan, Physica C **235-240**, 1489 (1994).
[12] W. F. Vinen, E. M. Forgan, C. E. Gough, and M. J. Hood, Physica **55**, 94 (1971).
[13] J. B. Sousa, Physica **55**, 507 (1971).
[14] P. H. Kes, J. P. M. van der Veeken, and D. de Kler, J. Low Temp. Phys. **18**, 355 (1975).
[15] C. Caroli and J. Matricon, Phys. Kondens. Mater. **3**, 380 (1965).
[16] J. Bardeen and M. J. Stephen, Phys. Rev. **140**, 1197 (1965).
[17] E. Canel (unpublished).
[18] J. Matsuda, N. P. Ong, Y. F. Yan, J. M. Harris, and J. B. Peterson, Phys. Rev. B **49**, 4380 (1994).
[19] In fact, these states may not be strictly localized in the case of *d*-wave superconducting coupling; in this case a consequent treatment may include some specific microscopic interactions (see, e.g., Refs. 20, 21). But, nevertheless, the simple microscopic model used to obtain Eq. (3) does not contradict existing theoretic approaches to vortex states in *d*-wave systems (Ref. 22).
[20] T. Dahm, C. T. Rieck, L. Tewordt, and S. Wermbter, Physica C **223**, 179 (1994).
[21] S. Wermbster and L. Tewordt, Physica C **183**, 365 (1991).
[22] K. Maki, N. Schopohl, and H. Won, Physica B **204**, 214 (1995).
[23] K. Karraï, E. J. Choi, F. Gunmore, S. Liu, H. D. Drew, Qi Li, D. B. Fenner, Y. D. Zhu, and Fu-Chun Zhang, Phys. Rev. Lett. **69**, 152 (1992).
[24] P. B. Allen, X. Du, L. Mihaly, and L. Forro, Phys. Rev. B **49**, 9073 (1994).
[25] M. Houssa and M. Ausloos, Physica C **235-240**, 1483 (1995).