PHYSICAL REVIEW B VOLUME 43, NUMBER 4 1 FEBRUARY 1991

Vortex-motion dissipation in high-T_c superconductors at microwave frequencies

R. Marcon, R. Fastampa, M. Giura, and E. Silva
Dipartimento di Fisica, Università di Roma La Sapienza, Roma, Italy
and Gruppo Nazionale di Struttura della Materia del CNR, Roma,
Centro Interuniversitario di Struttura della Materia del Ministero U.R.S.T., Roma, Italy
(Received 29 June 1990)

Measurements of the microwave surface resistance on samples of Y-Ba-Cu-O and Bi-Sr-Ca-Cu-O at different frequencies (23 and 48 GHz) indicate a universal behavior of the dissipation due to the vortex motion as a function of temperature, at low magnetic field. A temperature $T_0=(0.97\pm0.01)T_c$ separates two regimes such that for $T<T_0$ the dissipation increases with T while, for $T_0<T<T_c$ the dissipation decreases with increasing T. The role of fluctuations is expedient for explaining the experimental data near T_c.

I. INTRODUCTION

The appearance of a finite resistivity in the presence of an external magnetic field is a property of high T_c as well as of the conventional type-II superconductors. This feature, which is also important from a practical point of view, has been the subject of a number of works, both experimental and theoretical, on type-II superconductors. It is known that the finite resistivity is a consequence of the motion of the vortices, which are produced at a sufficiently high magnetic field (mixed state). When a transport supercurrent flows around a vortex or fluxon, a Lorentz-like force works between the current and the vortex. If the current is sufficiently high to suppress the action of pinning centers, a fluxon motion results that induces an electric field directed along the transport current, so that a finite resistivity appears and power is dissipated: flux-flow regime.[1]

The microscopic theory of superconductivity is able, in principle, to calculate the resistivity in the mixed state for any temperature T and magnetic field H. However, the resistivity has been calculated for the entire interval of temperatures both at low and high magnetic fields, only for alloys with a high concentration of paramagnetic impurities.[2,3] In general, relatively simple analytic expressions can be obtained in limited ranges of temperature ($T\rightarrow0$ or $T\rightarrow T_c$, T_c being the transition temperature) and of magnetic field ($H\rightarrow0$ or $H\rightarrow H_{c2}$). A review of calculations and comparison with experimental results for type-II superconductors can be found in the paper of Gor'kov and Kopnin[4] and in the work of Huebener.[5]

Two methods have generally been used for the measurement of the flow resistivity: the direct-current (dc) method and surface-resistance method. In the first method, a transport dc density $j_{\rm tr}$ is made to flow along the major side of a flat rectangular-shaped sample placed in a magnetic field H (with $H_{c1}<H<H_{c2}$, H_{c1} and H_{c2} being the lower and the upper critical fields, respectively) orthogonal to $j_{\rm tr}$.[6] A potential difference V along the current intensity $I_{\rm tr}$ is then measured and the slope of the linear section of the voltage-current ($V-I$) characteristic allows us to deduce the flow resistivity ρ_f. The behavior of $\rho_f(H,T)$ at low temperature and low magnetic field, as shown by a number of experiments, is in agreement with the expression $\rho_f=\rho_n B/H_{c2}$, where ρ_n is the normal-state resistivity and B is the magnetic induction ($B=n\phi_0$, n being the density of vortices and ϕ_0 the flux quantum). A review of the results can be found in the paper of Kim.[7] The dc method presents some disadvantages, namely the large transport current used to overcome the strength of the pinning centers ($j_{\rm tr}>j_c$, with j_c the critical current density for depinning) produces a heating of the sample and a self-field which bends the vortex filaments. In the surface resistance method an ac electric field produces flux motion at current densities several orders of magnitude below the critical values, provided that its frequency is higher than the pinning frequency ω_p.[8,9] For type-II superconductors, ω_p lies below the microwave range[9] and the microwave surface-resistance method has been frequently used to measure the flux-flow dissipation.

At microwave frequencies a phenomenological approach, based on the London electrodynamics and on the vortex-motion equation, which does not contain the restrictions present in the microscopic formulas, has often been employed.[9–12] The calculations are in good agreement with experiments[11,13] provided that the external magnetic field is not close to H_{c2}.

In the microwave-frequency range, the flux-flow regime in high-T_c superconductors has been first found in Y-Ba-Cu-O samples.[14] In the present paper, we apply the phenomenological approach to high-T_c superconductors for the calculation of the microwave surface resistance R_f under flux-flow conditions (Sec. II) and find that the measurements of R_f show a universal behavior as predicted by the theory (Sec. III).

II. CALCULATION OF THE MICROWAVE SURFACE RESISTANCE IN THE FLUX-FLOW REGIME

As well known, in the mixed state, for a spatially constant harmonic time-dependent $\exp(-i\omega t)$ electromagnetic field, the supercurrent **J** is given by[15]

$$J=-\frac{c^2}{i4\pi\omega\mu\lambda_L^2}\left[\mathbf{E}-\frac{\mathbf{v}\times f\mathbf{B}}{c}\right], \tag{1}$$

where the second term in parentheses is connected to vortex motion. In Eq. (1), λ_L is the London penetration length, **E** the electric field, **v** the vortex velocity, **B** the magnetic induction ($\mathbf{B}=n\boldsymbol{\phi}_0$, with n the two-dimensional vortex density and ϕ_0 the flux quantum) and f (≤ 1) the fraction of the weakly pinned vortices. The equation of motion for the unity of length of a vortex is

$$m_v\frac{d^2\mathbf{X}}{dt^2}+\eta\frac{d\mathbf{X}}{dt}+p\mathbf{X}=\frac{\mathbf{J}\times\boldsymbol{\phi}_0}{c}, \tag{2}$$

where m_v is the vortex mass per unit length, η the coefficient of viscosity, and p a constant which restrains the vortex displacement **X** (i.e., $p\mathbf{X}$ is the pinning force). At microwave frequencies, both the vortex mass and the pinning force can usually be neglected[8,16] (a discussion of this point will be made in Sec. III B), so that, in a pure viscous regime, we have ($\mathbf{v}=d\mathbf{X}/dt$)

$$\eta\mathbf{v}=\frac{\mathbf{J}\times\boldsymbol{\phi}_0}{c}. \tag{3}$$

Assuming $\mathbf{J}\perp\mathbf{B}$ and eliminating the vortex velocity from Eqs. (1) and (3), one obtains, for the flow resistivity $\rho_f=E/J$, the expression

$$\rho_f=\frac{\phi_a fB}{c^2\eta}-i\frac{4\pi\omega\mu\lambda_L^2}{c^2}, \tag{4}$$

which is valid when the vortex lattice is sufficiently rarefied, i.e., for the Ginzburg-Landau parameter $\kappa\gg 1$ and $B<H_{c2}$, conditions normally satisfied in high-T_c superconductors. By inserting Eq. (4) into the expression of the surface impedance

$$Z=(-4\pi i\omega\mu c^{-2}\rho_f)^{1/2},$$

one obtains

$$Z=\frac{4\pi\omega\mu\lambda_L}{c^2}\left[-1-i\frac{\phi_a fB}{4\pi\omega\mu\lambda_L^2\eta}\right]^{1/2}, \tag{5}$$

which yields the surface resistance

$$R_f=\mathrm{Re}(Z)=\frac{4\pi\omega\mu\lambda_L}{\sqrt{2}c^2}\left\{-1+\left[1+\left[\frac{\phi_0 fB}{4\pi\omega\mu\lambda_L^2\eta}\right]^2\right]^{1/2}\right\}^{1/2}, \tag{6}$$

which coincides with the analogous expression found by Gilchrist[11] by means of a different procedure. The viscosity coefficient is given by[4]

$$\eta=\gamma\frac{\phi_0}{c^2\rho_n}H_{c2}, \tag{7}$$

where ρ_n is the normal-state resistivity and γ a coefficient which, in the most general case, depends both on the field and the temperature (in the case $\gamma=1$ one obtains the viscosity coefficient given by Bardeen and Stephen.[17] Using Eq. (7), Eq. (6) can be cast in the more expressive form

$$R_f=aR_n^2\left\{-1+\left[1+\left[\frac{bB}{H_{c2}}\right]^2\right]^{1/2}\right\}^{1/2}, \tag{8}$$

where R_n is the surface resistance in the normal state $a=\mu\lambda_L\sqrt{2}/\rho_n$, and $b=\mu f/R_n^2a^2\gamma$ is a quantity much larger than unity.

For $bB\ll H_{c2}$, we obtain a behavior linear with B; that is,

$$R_f(t,B)\cong\beta(t)B=\frac{\rho_n f(t)}{2\gamma(t,B)\lambda_L(t)H_{c2}(t)}B, \tag{9}$$

where we have put into evidence the dependence of the various quantities from the reduced temperature $t=T/T_c$.

For $bB\gg H_{c2}$ we have

$$\frac{R_f}{R_n}=\left[\frac{\mu f}{\gamma H_{c2}}B\right]^{1/2}, \tag{10}$$

so that, after writing the resistivity in the general form $\rho_f=\rho_n B/\gamma H_{c2}$, Eq. (10) becomes

$$\frac{R_f}{R_n}=\left[\mu f\frac{\rho_f}{\rho_n}\right]^{1/2}, \tag{11}$$

which coincides (for $f=1$) with the analogous formulas for the resistance obtained by other authors.[4,13,18] One usually refers to Eq. (11) to describe the flow-dissipation data in a number of type-II superconductors (e.g., Pb-In and Nb-Ta alloys).[11,13,18] However, the experimental behavior of R_f is a linear function of B for small B, this linear behavior persisting for about the first 5% of the range of variation of B/H_{c2}. Thus, in the superconducting alloys as Pb-In and Nb-Ta for which H_{c2} is of order of a few kG, the dominant behavior is given by Eq. (10), while in high-T_c superconductors, as Y-Ba-Cu-O for which H_{c2} is of order 1 MG, the linear behavior described by Eq. (9) prevails up to a few kG (provided the temperature is not close to T_c). It is to be noted that the London penetration length λ_L is present at low fields [Eq. (9)] but not at high fields [Eq. (10)]. Therefore, the description of the behavior of R_f at low fields cannot leave the superconducting electrons out of consideration. In the following we will rely on the linear behavior (9) in order to fit our experimental results.

III. EXPERIMENTAL RESULTS AND DISCUSSION

We have made use of the method of microwave surface resistance to detect the flux-flow motion in high-T_c granular superconductor samples of Y-Ba-Cu-O and Bi-Sr-Ca-Cu-O at low magnetic fields ($H\ll H_{c2}$) and in the range of temperatures between 40 K and T_c. Measurements were made by means of two microwave apparatus, one working in the K band (23 GHz), the other in the Q-band (48 GHz). The disk-shaped samples were the bottoms of cylindrical cavities tuned in the TE_{011} mode. The microwave apparatus was made sensitive to the absorption of the cavity and the power reflected from it was detected. In this condition, the reflected power is proportional to surface resistance of the sample and the apparatus allows us to measure the variations of the surface resistance with the external magnetic field. In Fig. 1, a sketch of the experimental setup is shown.

For granular superconductors, the absorbed power exhibits a rather complicated behavior, due to presence of various effects depending on the range of temperatures and/or of magnetic fields. Part of the absorption is due to the decoupling of weak links (Josephson junctions or proximity effect links, among grains or twins) and increases with the increasing of the magnetic field H. This absorption, which has an exponential-like behavior, is present at low magnetic fields ($H < 100-300$ G) and in the entire temperature range below T_{cj} (the temperature T_{cj}, located a few K below T_c, is defined the temperature above which all the junctions are decoupled at $H=0$). This phenomenon has been quantitatively studied in previous works.[19,20] Furthermore, at low temperatures (about liquid-helium temperature), a sufficiently high magnetic field (of order of some kG) is able to drive the set of junctions in a frozen state.[20,21] Another cause of the microwave absorption, which shows a hysteretic behavior, is connected to flux-creep effects and, as measurements in progress seem to show, is present at sufficiently high magnetic fields and at temperatures lower than $0.5T_c$.

At higher temperatures ($T > 0.5T_c$), the measurements show that the microwave absorption is reversible and connected to motion of free or weakly pinned vortices, as we will describe below; in other words, we are in the presence of a flux-flow regime. This regime has the feature to be a general property of the mixed state of a superconductor. The dissipation should be, apart from scale factors, a universal function of the temperature and of B/H_{c2}; more precisely, the phenomenon should be magnetically reversible, independent from the type of material and from the microwave frequency (at least, within certain limits). Thus, to test the model, the resistance R_f as a function of temperature T has to be measured with different types of samples at different frequencies.

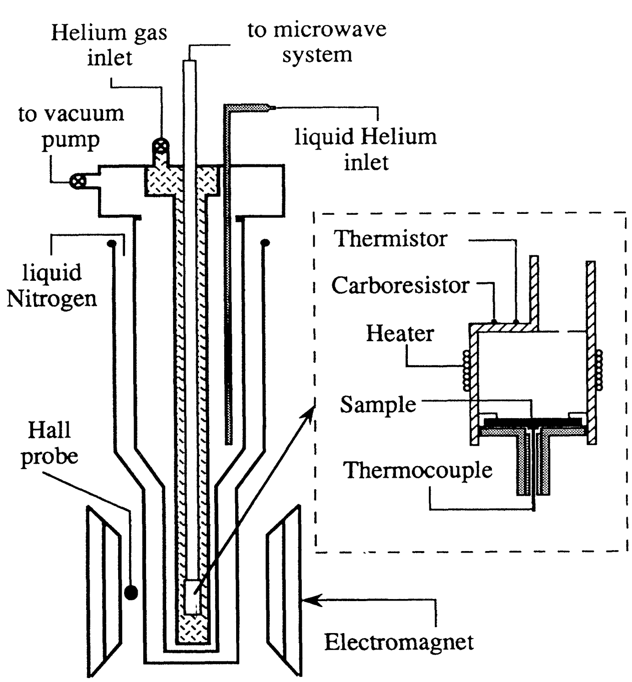

FIG. 1. Sketch of the cryogenic section of the experimental setup. In the inset, a detail of the microwave cavity and the sample arrangement is shown.

A. Expression of the viscosity coefficient $\eta(t)$

The first measurements based on the surface-resistance method, which have shown the existence of a dissipation due to the vortex motion in high-T_c superconductors, have been carried out by some of the authors in Y-Ba-Cu-O at 23 GHz.[14] Successively, measurements have been repeated in Bi-Sr-Ca-Cu-O, again at 23 GHz.[22] Due to the large value of H_{c2}, the behavior of Eq. (8) is nearly linear for T not strictly close to T_c. Therefore, at low fields, the best description of flux motion with the temperature is obtained through the slope $\beta(t)$ appearing in Eq. (9). Using Eq. (7), from Eq. (9) one obtains

$$\beta(t)=\frac{\Delta R_f}{\Delta B}=\frac{\phi_0}{2c^2}\frac{f(t)}{\eta(t)\lambda_L(t)} . \tag{12}$$

The experimental results for the slope, obtained in Y-Ba-Cu-O and Bi-Sr-Ca-Cu-O at 23 GHz, were fitted in a previous paper[22] by means of Eq. (12) under the assumption that the viscosity coefficient η was the sum of two terms

$$\eta(t)=\eta_c+\eta_{\rm BS}(t)=\eta_c+\frac{\phi_0}{c^2\rho_n}H_{c2}(t) ,$$

where $\eta_{\rm BS}(t)$ is the ordinary Bardeen-Stephen viscosity coefficient which is obtained by Eq. (7) with $\gamma=1$. The coefficient η_c is an additional viscosity term, empirically given by the expression $\eta_c=\eta_0(1-t)^{1/p}$ (with $p>3$) and arbitrarily introduced in order to avoid the divergence of Eq. (12) caused by $\eta_{\rm BS}(t)$ when $T\to T_c$. As will be seen, the divergence is not present in the measurement of $\beta(t)$ which, on the contrary, vanishes when $T\to T_c$. This experimental behavior of $\beta(t)$ can find a physical explanation if the temperature dependence of $\eta(t)$ close to T_c is reexamined. In the following we will give two expressions for $\beta(t)$, one valid for T not close to T_c $(t<1)$, the other for T close to T_c $(t\to 1)$.

Let us first examine the case $t<1$. The fraction $f(t)$ of weakly pinned fluxons is given by

$$f(t)=\exp\left[-\frac{U(t)}{k_BT}\right] , \tag{13}$$

$U(t)=(H_c^2/8\pi)\xi^3$ being the Anderson-Kim energy activation.[23] Assuming for the Ginzburg-Landau coherence length the expression

$$\xi=0.74\xi_0(1-t)^{-1/2} ,$$

and for the thermodynamical critical field the expression

$H_c=H_{c0}(1-t^2)$, Eq. (13) becomes

$$f(t)=\exp\left[-\frac{U_0}{T_c}\frac{(1-t^2)^2(1-t)^{-3/2}}{t}\right], \tag{14}$$

where U_0, the energy activation at $t=0$, is expressed in K. For the London penetration length, the strong-coupling expression is used,[24]

$$\lambda_L(t)=\lambda_0(1-t^4)^{-1/2} . \tag{15}$$

Finally, for the viscosity coefficient $\eta(t)$ we assume the Bardeen-Stephen expression which is obtained by Eq. (7) with $\gamma=1$; that is,

$$\eta(t)=\frac{\phi_0}{c^2\rho_n}H_{c2}(t) . \tag{16}$$

By means of Eqs. (14–(16) and by using, for $H_{c2}(t)$, the ordinary expression

$$H_{c2}(t)\sim(1-t)(1+t)^2 ,$$

the slope $\beta(t)$ given by Eq. (12) assumes the temperature dependence

$$\beta(t)=\beta_1\frac{(1-t^4)^{1/2}}{[(1-t)(1+t)^2]}\times\exp\left[-\frac{U_0}{T_c}\frac{(1-t^2)^2(1-t)^{-3/2}}{t}\right], \tag{17}$$

where U_0 is a parameter of the fit and β_1 is a factor depending on the electronic amplification chain.

In the second case ($t\rightarrow 1$), the influence of fluctuations on the temperature dependence of $\eta(t)$ must be taken into account. As it was shown by Gor'kov and Kopnin,[4] owing to the enlargement of the vortices when $T\rightarrow T_c$, the relaxation of the order parameter is slower than that proposed by Tinkham[25] and the Bardeen-Stephen viscosity coefficient has to be substituted by the expression[4]

$$\eta(t)=1.1(1-t)^{-1/2}\frac{\phi_0}{c^2\rho_n}H_{c2}(t) , \tag{18}$$

which is obtained from Eq. (7) with $\gamma=1.1(1-t)^{-1/2}$. Besides, in view of recent experimental results, the temperature dependence of $H_{c2}(t)$ in Eq. (18) has to be reexamined. In high-T_c superconductors, because of the small coherence length, the fluctuations play a much greater role than in traditional superconductors and one can expect an exotic power law for $H_{c2}(t)$ when $T\rightarrow T_c$. The accurate measurements carried out by Fang *et al.*[26] on the reversible onset of the magnetization in Y-Ba-Cu-O show that, from a few K near the transition, $H_{c2}(t)$ has the behavior

$$H_{c2}(t)\sim(1-t)^{1/2} .$$

Also, the data by Athreya *et al.*[27] seem to show the same trend. Then the right-hand side (rhs) of Eq. (18), with

$$H_{c2}(t)\sim(1-t)^{1/2} ,$$

becomes a constant near T_c. Thus, for $t\rightarrow 1$, assuming $\eta(t)=$const and using Eqs. (14) and (15), the slope $\beta(t)$ given by Eq. (12) becomes

$$\beta(t)=\beta_2(1-t^4)^{1/2}\exp\left[-\frac{U_0}{T_c}\frac{(1-t^2)^2(1-t)^{-3/2}}{t}\right] . \tag{19}$$

On the basis of the preceding arguments we will use two expressions in order to fit the experimental results: Eq. (17) for $t<1$ and Eq. (19) for $t\rightarrow 1$.

B. Experimental results

To test the consistence of the experimental data with theoretical expectations, a numerical check of the range of validity of the theory is in order. Equation (3) is obtained from Eq. (2) neglecting both the vortex mass m_v and the pinning coefficient p. Let us write Eq. (2) in the harmonic regime [exp $(-i\omega t)$]

$$\mathbf{X}[(p-\omega^2 m_v)-i\omega\eta]=\frac{\mathbf{J}\times\boldsymbol{\phi}_0}{c} . \tag{20}$$

The vortex mass m_v seems to show an appreciable dynamic effect only when the ac frequency $v=\omega/2\pi$ is larger than the reciprocal $1/\tau$ of the normal-state electron collision time.[16,17] For high-T_c superconductors, $1/\tau$ is estimated to be $\sim 10^3$ GHz or larger, well above our frequencies of measurement (23 and 48 GHz). Then Eq. (20) can be written

$$(\omega_p-i\omega)\mathbf{X}=\frac{\mathbf{J}\times\boldsymbol{\phi}_0}{c\eta} , \tag{21}$$

$\omega_p=p/\eta$ being the pinning frequency. A crude estimate of ω_p can be made following Gilchrist.[11] If the pinning extends to a distance r, the maximum restraining force is pr and it has to be equated to $J_c\phi_0/c$, J_c being the critical current density. Thus, we have $pr=J_c\phi_0/c$ or $\omega_p\eta r=J_c\phi_0/c$, that is $\omega_p=c^2J_c\rho_n/rH_{c2}$. At low temperature, assuming $J_c=10^6\,\mathrm{A/cm^2}$, $\rho_n=0.1$ mΩ cm, $H_{c2}=1$ MG, and $r=\xi=10^{-7}$ cm, one obtains $\omega_p\cong 10^{11}$ rad/s, which is of order of the measurement frequencies. Therefore, we should not be in a pure flux-flow regime for which the condition $\omega>\omega_p$ has to be satisfied. However, one has to take into account that there is a distribution of the pinning distance r (the value $r=\xi$, i.e., the vortex core dimension, is the minimum value that r can assume) and that only a fraction $f(t)$ of fluxons, dependent on the temperature, will be free or weakly pinned at the work frequencies ω. Hence, in Eq. (21), for the fluxon fraction $f(t)$, one can assume $\omega>\omega_p$ and Eq. (20) reduces to Eq. (3), used to fit of experimental results. On the other hand, doubling the microwave frequency, we do not observe a different behavior of the surface resistance slope $\beta(t)$ as a function of the temperature and this implies that $\omega>\omega_p$.

Let us now examine the magnetic response of the specimen. The influence of the sample shape may be described by a demagnetizing coefficient n. In our arrangement the specimen is equivalent to a thin slab with the external magnetic field H parallel to the surface and hence $n=0$

and $H_i = H$, H_i being the magnetic field in the material. Besides, since the Ginzburg-Landau parameter κ is very large, the susceptibility of the material is small if $H > H_{c1}$. Therefore, in the following we will assume $B = H$ in discussing the experimental results.

In Fig. 2, typical behaviors of the microwave surface resistance R_f as a function of the external magnetic field H at various temperatures T are shown.

In Fig. 3 the entire set of measurements of the slope $\beta(t) = \Delta R_f / \Delta H$ of the surface resistance R_f at low external magnetic fields as a function of the reduced temperature t is shown. Data in the figure refer to samples of $YBa_2Cu_3O_7$ ($T_c = 90.5$ K) and $Bi_2Sr_2Ca_2Cu_3O_{10}$ (axis $c \sim 37$Å, $T_c = 109.5$ K). For the specimen of Y-Ba-Cu-O, the measurements are performed at two frequencies, 23 and 48 GHz. It is clearly seen that, within the scattering of measurements, the slope $\beta(t)$, as it was predicted above, has univeral behavior which seems to be a truly remarkable result.[28]

In the same figure, the best fit of the experimental results (solid lines) carried out by means of Eq. (17) for $t < 0.97$ and of Eq. (19) for $t > 0.97$, is reported together with the experimental data. A fairly good agreement between theory and experiment is found with $U_0 = 0.8 T_c$, assuming U_0 to be the same for both Eqs. (17) and (19). This low value of U_0, as compared to those reported in literature under conditions of flux creep,[29] shows that we

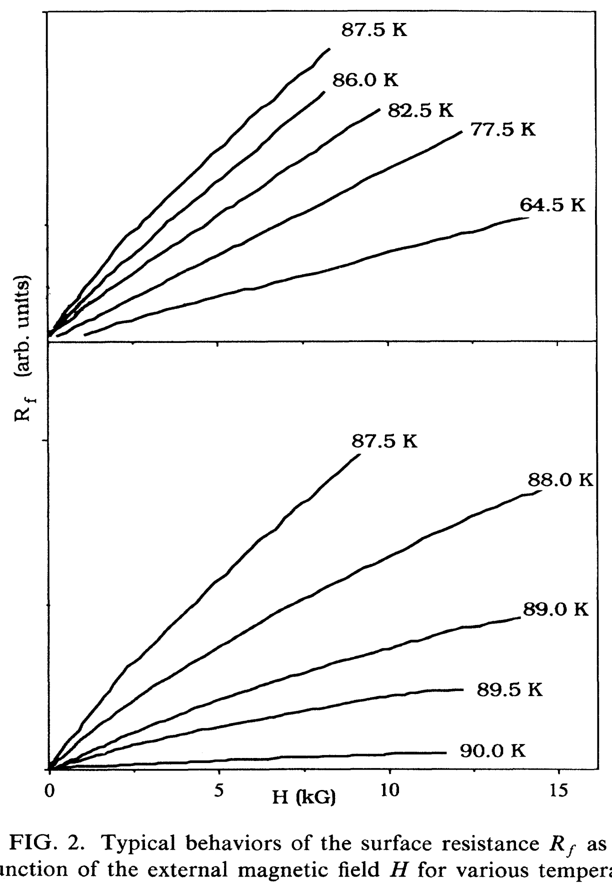

FIG. 2. Typical behaviors of the surface resistance R_f as a function of the external magnetic field H for various temperatures. The measurements refer to a sample of Y-Ba-Cu-O at 48 GHz.

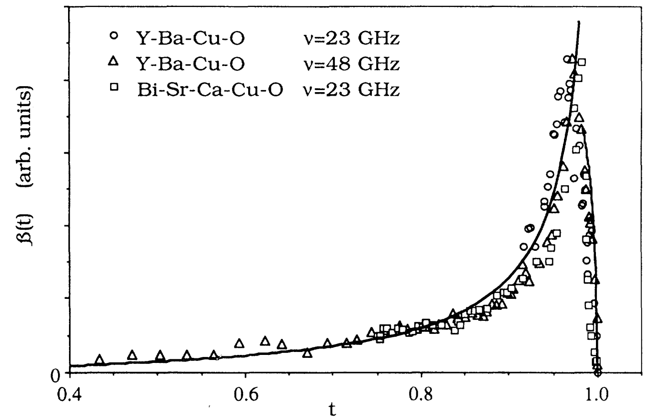

FIG. 3. Behavior of the microwave surface-resistance slope $\beta(t) = \Delta R_f / \Delta H$ at low magnetic fields as a function of the reduced $t = T/T_c$ temperature. Symbols refer to experimental results, solid lines to the best fit carried out by means of Eq. (17) for $t < 0.97$ and of Eq. (19) for $t > 0.97$ with $U_0 = 0.8 T_c$.

are at the edges of the pure flux-flow regime, in the sense previously explained.

As it can be checked, a temperature $T_0 = (0.97 \pm 0.01) T_c$ separates two regimes of dissipation. Above T_0, fluctuations seem to play a fundamental role in reducing dissipation. Should it be experimentally confirmed, an interpretation could be found in the Nelson model of a transition from a disentangled to an entangled vortex fluid.[30]

IV. CONCLUSIONS

(1) In high-T_c superconductors (Y-Ba-Cu-O and Bi-Sr-Ca-Cu-O), measurements of the microwave surface resistance R_f have shown the presence of a dissipation due to vortex motion, analogous to that of conventional type-II superconducting alloys.

(2) The dissipation can be described by a simple phenomenological model based on the London electrodynamcis and the vortex-motion equation in the flow regime.

(3) The dissipation is a universal function of the reduced temperature $t = T/T_c$ and of the reduced magnetic field $h = B/H_{c2}$. In particular, it is magnetically reversible and independent (at least, within certain limits) from the material and the microwave frequency.

(4) The behavior of R_f with H scales as H_{c2}. Thus, owing to the large value of H_{c2}, R_f can be considered a linear function of the magnetic induction B [see Eq. (9)] for external magnetic fields up to a few kG (provided the temperature is not close to T_c). In this frame the presence in Eq. (4) of the imaginary term due to the superconducting electrons is expedient for the agreement between theory and experimental results.

(5) Since the pinning frequency ω_p seems to be not much lower than the employed microwave frequencies, a thermally supported flux-flow dissipation must be considered [see Eq. (14)] to fit the data.

[1]See, for example, M. Tinkham, *Introduction to Superconductivity* (McGraw-Hill, New York, 1975); A. A. Abrikosov, *Fundamentals of the Theory of Metals* (North-Holland, Amsterdam, 1988).

[2]R. S. Thompson and C.-R. Hu, Phys. Rev. Lett. **27**, 1352 (1971); M. Yu Kupriyanov and K. K. Likharev, Pis'ma Zh. Eksp. Teor. Fiz. **15**, 349 (1972) [JETP Lett. **15**, 247 (1972)]; C.-R. Hu and R. S. Thompson, Phys. Rev. B **6**, 110 (1972); C.-R. Hu, *ibid*, **6**, 1756 (1972); C.-R. Hu and R. S. Thompson, Phys. Rev. Lett. **31**, 217 (1973).

[3]L. P. Gor'kov and N. B. Kopnin, Zh. Eksp. Teor. Fiz. **60**, 2331 (1971) [Sov. Phys.—JETP **33**, 1251 (1971)].

[4]L. P. Gor'kov and N. B. Kopnin, Usp. Fiz. Nauk **116**, 413 (1975) [Sov. Phys.—Usp. **18**, 496 (1976)].

[5]R. P. Huebener, *Magnetic Flux Structures in Superconductors,* Vol. 6 of *Springer Series in Solid State Sciences* (Springer-Verlag, Berlin, 1979), pp. 170–176.

[6]A. R. Strnad, C. F. Hempstead, and Y. B. Kim, Phys. Rev. Lett. **13**, 794 (1964); Y. B. Kim, C. F. Hempstead, and A. R. Strnad, Phys. Rev. **139**, 1163 (1965).

[7]Y. B. Kim, *Proceedings of the XII International Conference on Low Temperature Physics, Kyoto, 1971* (Academic Press of Japan, Japan, 1971), p. 231, and references therein.

[8]B. Rosenblum and M. Cardona, Phys. Rev. Lett. **12**, 657 (1964); B. Rosenblum, M. Cardona, and G. Fischer, RCA Rev. **25**, 491 (1964).

[9]J. I. Gittleman and B. Rosenblum, Phys. Rev. Lett. **16**, 734 (1966).

[10]A. G. Van Vijfeijken and A. K. Niessen, Phys. Lett. **16**, 23 (1965).

[11]J. le G. Gilchrist, Proc. R. Soc. London, Ser. A **295**, 399 (1966).

[12]P. Nozieres and W. F. Vinen, Philos. Mag. **14**, 667 (1966); W. F. Vinen and A. C. Warren, Proc. Phys. Soc. London **91**, 409 (1967); M. Rabinowitz, J. Appl. Phys. **42**, 88 (1971).

[13]J. le G. Gilchrist and P. Monceau, J. Phys. C **3**, 1399 (1970).

[14]R. Fastampa, M. Giura, R. Marcon, and C. Matacotta, Europhys. Lett. **9**, 719 (1989).

[15]A. M. Portis, K. W. Blazey, K. A. Müller, and J. G. Bednorz, Europhys. Lett. **5**, 467 (1988); E. B. Sonin and A. K. Tagantev, Zh. Eksp. Teor. Fiz. **95**, 994 (1989) [Sov. Phys.—JETP **68**, 572 (1989)].

[16]H. Suhl, Phys. Rev. Lett. **14**, 226 (1965).

[17]J. Bardeen and M. J. Stephen, Phys. Rev. **140**, A1197 (1965).

[18]J. le G. Gilchrist and P. Monceau, Philos. Mag. **18**, 237 (1968).

[19]R. Marcon, R. Fastampa, M. Giura, and C. Matacotta, Phys. Rev. B **39**, 2796 (1989).

[20]M. Giura, R. Marcon, and R. Fastampa, Phys. Rev. B **40**, 4437 (1989).

[21]M. Giura, R. Fastampa, R. Marcon, and E. Silva, Phys. Rev. B **42**, 6228 (1990).

[22]R. Marcon, R. Fastampa, and M. Giura, Europhys. Lett. **11**, 561 (1990).

[23]P. W. Anderson, Phys. Rev. Lett. **9**, 309 (1962); P. W. Anderson and Y. B. Kim, Rev. Mod. Phys. **36**, 39 (1964).

[24]The empirical function of temperature $F(t)=\lambda_0/\lambda_L(t)=F_1$, previously used by us to fit the microwave absorption connected to the decoupling of the junctions (Ref. 19), is very close to the strong-coupling behavior $F(t)=(1-t^4)^{1/2}$.

[25]M. Tinkham, Phys. Rev. Lett. **13**, 804 (1964).

[26]M. M. Fang, V. G. Kogan, D. K. Finnemore, J. R. Clem, L. S. Chumbley, and D. E. Farrell, Phys. Rev. B **37**, 2334 (1988).

[27]K. S. Athreya, O. B. Hyun, J. E. Ostenson, J. R. Clem, and D. K. Finnemore, Phys. Rev. B **38**, 11 846 (1988).

[28]We note that the steep descent of the dissipation near T_c, with increasing T, cannot be attributed to the amplification reduction of the microwave cavity caused by the superconducting transition, this effect having already been taken into account in the elaboration of experimental data.

[29]T. T. M. Palstra, B. Batlogg, L. F. Schneemeyer, and J. V. Wasczak, Phys. Rev. Lett. **61**, 1662 (1988).

[30]D. R. Nelson, Phys. Rev. Lett. **60**, 1973 (1988).