Fluctuation induced diamagnetism versus paraconductivity at high-reduced-temperatures in layered superconductors

2001 J. Phys.: Condens. Matter 13 2573

(http://iopscience.iop.org/0953-8984/13/11/313)

INSTITUTE OF PHYSICS PUBLISHING JOURNAL OF PHYSICS: CONDENSED MATTER

J. Phys.: Condens. Matter **13** (2001) 2573–2581 www.iop.org/Journals/cm PII: S0953-8984(01)17983-X

Fluctuation induced diamagnetism versus paraconductivity at high-reduced-temperatures in layered superconductors

C Carballeira[1], M V Ramallo[1] and F Vidal[1,2]

[1] Laboratorio de Baixas Temperaturas e Superconductividade[1], Departamento de Física da Materia Condensada, Universidade de Santiago de Compostela, Santiago de Compostela E-15706, Spain
[2] IRC in Superconductivity, Cavendish Laboratory, University of Cambridge, Cambridge CB3 0HE, UK

E-mail: fmvidal@usc.es (F Vidal)

Received 12 October 2000

Abstract
By introducing a total-energy or a momentum cut-off in the Gaussian–Ginzburg–Landau (GGL) approach for layered superconductors, the previous calculations of the in-plane paraconductivity, the fluctuation induced diamagnetism and the ratio between both observables are extended to the high-reduced-temperature region (typically for reduced temperatures, $\epsilon \equiv \ln(T/T_{c0})$, above 0.1), where the short-wavelength fluctuation effects may be important. As a first check of their interest, these theoretical results are used to briefly analyse the experimental data recently obtained in optimally doped $YBa_2Cu_3O_{7-\delta}$ samples. This comparison strongly suggests the adequacy of a total-energy cut-off to explain the observed high-reduced-temperature behaviour of the thermal fluctuations in copper oxide superconductors.

1. Introduction

It is now well established that many of the properties of the so-called high-temperature cuprate superconductors (HTSC) are appreciably affected above their superconducting transition temperature, T_{c0}, by the presence of Cooper pairs created by thermal fluctuations [1]. In turn, these thermal fluctuation effects may be used as a useful tool to access various central parameters of any microscopic or phenomenological description of these superconductors, including their superconducting coherence length amplitudes in all directions [1, 2]. Two of the observables best adapted to analyse these thermal fluctuations in HTSC are the in-plane paraconductivity (which affects the electrical conductivity parallel to the superconducting CuO_2 layers), $\Delta\sigma_{ab}$, and the fluctuation induced diamagnetism (for the magnetic field applied perpendicularly to the superconducting layers), $\Delta\chi_{ab}$. In HTSC, these observables are among

[1] Unidad Asociada al ICMM, Consejo Superior de Investigaciones Científicas, Madrid, Spain.

the ones with the highest amplitude relative to their normal or background behaviours [1, 2]. This is indeed a very important experimental advantage. However, in multilayered HTSC, with N superconducting CuO_2 layers per periodicity length, s, the analyses of both magnitudes are made difficult by their dependence on various parameters which are not always well settled, including the relative Josephson coupling strength between adjacent superconducting layers [1, 2]. Such a dependence may be easily calculated on the grounds of the so-called Gaussian–Ginzburg–Landau approach, the resulting expressions being [2, 3]

$$\Delta\sigma_{ab}(\epsilon) = N_e \frac{A_{AL}}{\epsilon}\left(1 + \frac{B_{LD}}{\epsilon}\right)^{-1/2} \tag{1}$$

and

$$\frac{\Delta\chi_{ab}(\epsilon)}{T} = N_e \frac{A_S}{\epsilon}\left(1 + \frac{B_{LD}}{\epsilon}\right)^{-1/2} \tag{2}$$

where N_e is an effective number of independent fluctuating superconducting CuO_2 layers per periodicity length, $\epsilon \equiv \ln(T/T_{c0})$ is the reduced temperature, $A_{AL} \equiv e^2/16\hbar s$ is the Aslamazov–Larkin paraconductivity amplitude, e is the electron charge, $\hbar$ is the reduced Planck constant, $B_{LD} \equiv [2\xi_c(0)/s]^2$ is the Lawrence–Doniach (LD) parameter which controls the dimensionality of the thermal fluctuations in single-layered materials, $\xi_c(0)$ is the superconducting coherence length amplitude in the c-direction (perpendicular to the superconducting layers), $A_S \equiv \mu_0\pi k_B\xi_{ab}^2(0)/3\phi_0^2 s$ is the Schmidt diamagnetism, μ_0 is the vacuum magnetic permeability, k_B is the Boltzmann constant, $\xi_{ab}(0)$ is the in-plane coherence length amplitude and $\phi_0 = \hbar\pi/e$ is the magnetic flux quantum. Equation (2) applies in the weak magnetic field limit (also called Schmidt-limit). This last limit corresponds to the condition $h \equiv H/H_{c2}(0) \ll \epsilon$, where $H_{c2}(0)$ is the upper critical magnetic field amplitude with H applied perpendicularly to the superconducting CuO_2 layers.

In spite of the existence of these explicit expressions of $\Delta\sigma_{ab}(\epsilon)$ and $\Delta\chi_{ab}(\epsilon)$, the presence in equations (1) and (2) of B_{LD} and, mainly, of N_e (which, in particular, depends on the relative Josephson coupling strength between adjacent layers [2, 3]) introduces, as noted before, important ambiguities when analysing separately the experimental results on the in-plane paraconductivity and on the fluctuation induced diamagnetism in multilayered HTSC [1–4]. However, as was recognized earlier [4], some of these complications may be easily overcome by using the relationship between both effects, which may be written (in MKSA units) as [2–5]

$$\frac{\Delta\chi_{ab}(\epsilon)}{T\Delta\sigma_{ab}(\epsilon)} = \frac{16\mu_0 k_B\xi_{ab}^2(0)}{3\pi\hbar} = 2.79\times 10^5\xi_{ab}^2(0). \tag{3}$$

This relationship does not depend on N_e (or, equivalently, on a 'not well defined'[2] effective periodicity length), but it is also ϵ-independent. Actually, it depends only on the in-plane coherence length amplitude. As a consequence, equation (3) has already provided one of the most direct and reliable checks of the applicability to the HTSC of the GGL-like approaches in the reduced temperature region bounded by $10^{-2} \lesssim \epsilon \lesssim 10^{-1}$ [2–5]. Moreover, equation (3) also provides one of the most reliable ways to determine the superconducting coherence length amplitude in the ab-plane in HTSC, a central parameter for any description of the superconductivity in these materials and that, due to the presence of

[2] Some of these difficulties are stressed in [1] p 325: 'Quantitative comparisons between the LD theory and the real materials are problematic because the crystallographic unit cell typically contains two or more inequivalent CuO_2 planes which, in principle, would require a generalization of the LD model; thus, even the appropriate choice of interplane spacing s in the model is not well defined. Moreover, it has proved difficult to obtain a consensus on accurate experimental values of ξ_i because they are inferred from $H_{c2}(T)$ data, and H_{c2} is poorly defined because of fluctuation rounding of the transition in high-temperature superconductors.'

thermal fluctuations, is difficult to measure using other types of experiment[2]. Complementarily, the agreement at a quantitative level of equation (3) with the experimental results obtained in different HTSC compounds in this ϵ-region was the first experimental demonstration of the absence of appreciable indirect (Maki–Thompson and DOS) fluctuation effects on the in-plane paraconductivity, a result in turn related to the unconventional (non $^1s_0$) pairing state in these superconductors [2–5]. It was observed earlier, however, that equation (3) applies only in the reduced temperature range bounded by $10^{-2} \lesssim \epsilon \lesssim 10^{-1}$ [2–5]. In fact, these two limits of validity in reduced temperature are very general: the failure for $\epsilon \lesssim 10^{-2}$ may be attributed to the penetration in the so-called full-critical region. This reduced temperature agrees quite well with the so-called Levanyuk–Ginzburg reduced temperature, ϵ_{LG}, which provides an estimation of the frontier between the mean-field and the full-critical regions [6]. In fact, the ϵ-dependence of $\Delta\chi_{ab}(\epsilon)/(T\,\Delta\sigma_{ab}(\epsilon))$ observed experimentally for reduced temperatures below 10^{-2} agrees, at least at a qualitative level, with the behaviour of $\Delta\sigma_{ab}(\epsilon)$ and $\Delta\chi_{ab}(\epsilon)$ predicted by the scaling approaches in the full-critical region [2–5].

The failure of equation (3) in the so-called high-reduced-temperature region, for $\epsilon \gtrsim 10^{-1}$, may be attributed to the fact that the conventional GGL approach strongly overestimates the statistical weight of the thermal fluctuations with wavelengths of the order of $\xi_{ab}(0)$. These short-wavelength effects on the thermal fluctuations were already observed in low-temperature superconductors (LTSC) through measurements of both the fluctuation induced diamagnetism [1, 7] and the paraconductivity [8]. In the case of the HTSC, the breakdown of the GGL theory at high-reduced-temperatures was mainly studied until now through the paraconductivity [9, 10]. These studies clearly show that these failures cannot be appreciably mitigated by introducing a momentum cut-off in the fluctuation spectrum, i.e. by imposing in the GGL approach the condition

$$k^2 < c\,\xi_{ab}^{-2}(0) \tag{4}$$

where k is the momentum (in units of $\hbar$) of each fluctuating mode, Ψ_{k}, and c is a constant (temperature independent) cut-off amplitude close to 1. However, it was recently proposed [10] that these difficulties in the paraconductivity may be eliminated if, instead of a momentum cut-off, a *total-energy cut-off* is imposed in the 2D-limit by (in units of $\hbar^2/2m^*$, where m^* is the effective mass of the Cooper pairs)

$$[k^2 + \xi_{ab}^{-2}(\epsilon)] < c\,\xi_{ab}^{-2}(0). \tag{5}$$

As already stressed in [10], this total-energy cut-off may be easily justified by taking into account that the probability of each fluctuating mode is controlled by its total energy $[k^2 + \xi_{ab}^{-2}(\epsilon)]$, and not only by its momentum [11][3]. By using the mean-field reduced-temperature dependence of the coherence length, $\xi_{ab}(\epsilon) = \xi_{ab}(0)\epsilon^{-1/2}$, equation (5) may be rewritten as $k^2 < (c-\epsilon)\xi_{ab}^{-2}(0)$. We see, therefore, that for $\epsilon \ll c$, the total-energy cut-off reduces to the conventional momentum cut-off. Nevertheless, the differences between both cut-off conditions may be very important at high-reduced-temperatures, when ϵ becomes of the order of c.

The central aim of this paper is to extend these results on the cut-off effects at high-reduced-temperatures to the relationship between the in-plane paraconductivity and the fluctuation induced diamagnetism in layered HTSC, i.e. to calculate the modifications that the momentum cut-off and the total-energy cut-off conditions may introduce in equation (3). Then, we will briefly compare these theoretical results with recent measurements in the high-reduced-temperature region in optimally doped $YBa_2Cu_3O_{7-\delta}$ (Y-123) samples. For simplicity, our

[3] A total-energy cut-off was already suggested by Patton and coworkers and by Nam when analysing in terms of a microscopic approach the fluctuation induced diamagnetism at high applied magnetic fields in LTSC.

calculations are going to be centred in single-layered superconductors. However, they may be easily extended to multilayered superconductors in various important limit situations as, for instance, when the Josephson coupling strength between adjacent superconducting layers are similar.

2. Theory

2.1. Paraconductivity under different cut-off conditions

Our starting point to calculate the paraconductivity under different cut-off conditions on the grounds of the phenomenological GGL approach is the relationship between $\Delta\sigma_{ab}(\epsilon)$ and the momentum of the fluctuating modes. Such a relationship was calculated for the first time by Schmidt for bulk superconductors [13], and it may be straightforwardly extended to the layered case by changing the fluctuation energy in the z-direction, $\xi_c^2(0)k_z^2$ (where k_z is the z-momentum of the fluctuation modes), by $B_{LD}[1-\cos(k_z s)]/2$ with $|k_z| \leqslant \pi/s$ [14]. The resulting expression is

$$\Delta\sigma_{ab}(\epsilon) = \frac{e^2\xi_{ab}^4(0)}{8\pi\hbar}\int \mathrm{d}k_z \int \mathrm{d}k_{xy}\frac{k_{xy}^3}{\left\{\epsilon + B_{LD}[1-\cos(k_z s)]/2 + \xi_{ab}^2(0)k_{xy}^2\right\}^3} \tag{6}$$

where k_{xy} is the modulus of the in-plane momentum of the fluctuation modes. Equation (6) is general and the in-plane paraconductivity for any cut-off criterion may be obtained by simply imposing the corresponding upper limits on the k-integrals. To calculate $\Delta\sigma_{ab}(\epsilon)$ under a momentum cut-off, we note first that since the z-spectrum of the fluctuations is already modulated through $-\pi/s \leqslant k_z \leqslant \pi/s$, the inclusion of a momentum cut-off in this direction is not necessary. For Y-123, this is a correct approach because the effective periodicity is $s = 5.9$ Å, whereas $\xi_c(0) \simeq 1.1$ Å and, therefore, the condition $|k_z| \leqslant \pi/s$ is stronger than $k_z^2 < c\xi_c^{-2}(0)$ (if $c \simeq 1$, see below). So, equation (4) becomes

$$k_{xy}^2 < c\,\xi_{ab}^{-2}(0) \tag{7}$$

and the in-plane paraconductivity under the momentum cut-off criterion, $\Delta\sigma_{ab}(\epsilon, c)_M$, is then found to be

$$\Delta\sigma_{ab}(\epsilon, c)_M = \frac{e^2}{16\hbar s}\left\{\frac{1}{\epsilon}\left(1+\frac{B_{LD}}{\epsilon}\right)^{-1/2} - \frac{c(c+\epsilon+B_{LD}/2)}{[(c+\epsilon+B_{LD})(c+\epsilon)]^{3/2}} - \frac{1}{\epsilon+c}\left(1+\frac{B_{LD}}{\epsilon+c}\right)^{-1/2}\right\}. \tag{8}$$

This expression has two interesting asymptotic limits: the conventional LD in-plane paraconductivity, equation (1) with $N_e = 1$, which is recovered by imposing ϵ, $B_{LD} \ll c$, and the 2D-limit of the paraconductivity under a momentum cut-off which may be obtained by imposing $B_{LD} \ll \epsilon$ (we find again the $\Delta\sigma_{ab}(\epsilon, c)_M$ expression first obtained for this 2D-limit by Gauzzi and Pavuna [9]).

To obtain from equation (6) the in-plane paraconductivity under a total-energy cut-off, $\Delta\sigma_{ab}(\epsilon, c)_E$, we must first note that the total-energy of the fluctuation modes is given by [2, 3]

$$E(\Psi_{\boldsymbol{k}}) = k_{xy}^2 + \xi_{ab}^{-2}(0)[\epsilon + B_{LD}(1-\cos(k_z s))/2]. \tag{9}$$

Therefore, the total-energy cut-off limits the in-plane momentum of the fluctuations through,

$$k_{xy}^2 < [c - \epsilon - B_{LD}(1-\cos(k_z s))/2]\;\;\xi_{ab}^{-2}(0) \tag{10}$$

whereas k_z is restricted, as in the case of the momentum cut-off, to the interval $-\pi/s \leqslant k_z \leqslant \pi/s$. By introducing these limits of integration in equation (6) we obtain,

$$\Delta\sigma_{ab}(\epsilon, c)_E = \frac{e^2}{16\hbar s}\left[\frac{1}{\epsilon}\left(1+\frac{B_{LD}}{\epsilon}\right)^{-1/2} - \frac{2}{c} + \frac{\epsilon + B_{LD}/2}{c^2}\right]. \tag{11}$$

Here again, the LD-limit (equation (1) with $N_e = 1$) can be recovered by simply imposing ϵ, $B_{LD} \ll c$, whereas the 2D-limit corresponds to $B_{LD} \ll \epsilon$. Other details of these calculations are going to be published elsewhere [10].

2.2. Fluctuation induced diamagnetism under different cut-off conditions

The simplest way to obtain $\Delta\chi_{ab}$ under both cut-off conditions in the LD scenario is to start by doing the calculations in the Prange-regime (i.e. under a finite magnetic field [1, 7]) and then going to the Schmidt-limit. In doing such a calculation, our starting point will be the well known divergent expression for the fluctuational part of the free energy per unit of volume [15, 16],

$$\langle\Delta F(\epsilon, h)\rangle = \frac{k_B T}{8\pi^2\xi_{ab}(0)^2}2h\sum_{n=0}^{\infty}\int_{-\pi/s}^{\pi/s}\mathrm{d}k_z\left[\ln\left(n+\frac{\epsilon+h+w_{k_z}}{2h}\right)+\ln(2h) - \ln\left(\frac{k_B T}{2a_0}\right)\right] \tag{12}$$

where a_0 is the so-called Ginzburg–Landau normalization constant and $n = 0, 1\ldots$ is the Landau-level index. This arises because in the presence of H the in-plane momentum components of the GGL-spectrum, k_x and k_y, are no longer good quantum numbers simultaneously and so k_{xy}^2 must be replaced by

$$k_{xy}^2 \rightarrow \frac{4e\mu_0 H}{\hbar}\left(n+\frac{1}{2}\right). \tag{13}$$

The momentum and total-energy cut-off lead to an upper limit, n_c, in the sum over the Landau levels in equation (12) which depends on the particular cut-off condition. In the case of the momentum cut-off, by combining equations (7) and (13) we obtain

$$n_c = \frac{c}{2h} - 1. \tag{14}$$

The fluctuation induced diamagnetism for finite fields under a momentum cut-off may be then obtained by imposing this upper limit in the sum of equation (12) and by using $\Delta\chi_{ab}(\epsilon, h) \equiv -(H\mu_0 H_{c2}(0))^{-1}\partial\langle\Delta F(\epsilon, h)\rangle/\partial h$. This gives

$$\Delta\chi_{ab}(\epsilon, h, c)_M = \frac{k_B T}{2\pi\phi_0 H}\int_{-\pi/s}^{\pi/s}\mathrm{d}k_z\left\{\ln\Gamma\left(\frac{\epsilon+h+w_{k_z}+c}{2h}\right) - \ln\Gamma\left(\frac{\epsilon+h+w_{k_z}}{2h}\right) - \frac{\epsilon+w_{k_z}+c}{2h}\psi\left(\frac{\epsilon+h+w_{k_z}+c}{2h}\right) + \frac{\epsilon+w_{k_z}}{2h}\,\psi\left(\frac{\epsilon+h+w_{k_z}}{2h}\right) + \frac{c}{2h}\right\} \tag{15}$$

where Γ and ψ are, respectively, the Gamma and Digamma functions. The fluctuation induced diamagnetism in the Schmidt-limit and under a momentum cut-off may be now obtained from equation (15) by imposing $h \ll \epsilon, c, B_{LD}$:

$$\frac{\Delta\chi_{ab}(\epsilon, c)_M}{T} = \frac{\pi\mu_0 k_B\xi_{ab}^2(0)}{3\phi_0^2 s}\left[\frac{1}{\epsilon}\left(1+\frac{B_{LD}}{\epsilon}\right)^{-1/2} - \frac{1}{\epsilon+c}\left(1+\frac{B_{LD}}{\epsilon+c}\right)^{-1/2}\right]. \tag{16}$$

Note that the Schmidt-limit without cut-off and for single-layered superconductors (equation (2) with $N_e = 1$) may be recovered by applying ϵ, $B_{LD} \ll c$ in the above equation.

Also note that for $B_{LD} = 0$ (i.e. in the 2D-limit) equation (15) reduces to the expression proposed to analyse the fluctuation induced diamagnetism well inside the finite magnetic field regime in a $La_{1.9}Sr_{0.1}CuO_4$ superconductor [17].

In the case of the total-energy cut-off, by combining equations (10) and (13) we found

$$n_c = \frac{1}{2h}\left[c - \epsilon - \frac{B_{LD}}{2}\left[1 - \cos(k_z s)\right]\right] - 1. \tag{17}$$

So, following the procedure described above, the fluctuation induced diamagnetism at finite applied magnetic fields under a total-energy cut-off is found to be,

$$\Delta\chi_{ab}(\epsilon, h, c)_E = \frac{k_B T}{2\pi\phi_0 H}\int_{-\pi/s}^{\pi/s} \mathrm{d}k_z\left\{-\frac{c}{2h}\psi\left(\frac{h + c_E}{2h}\right) - \ln\Gamma\left(\frac{\epsilon + h + w_{k_z}}{2h}\right)\right.$$
$$\left. + \ln\Gamma\left(\frac{h+c}{2h}\right) + \frac{\epsilon + w_{k_z}}{2h}\,\psi\left(\frac{\epsilon + h + w_{k_z}}{2h}\right) + \frac{c - \epsilon - w_{k_z}}{2h}\right\}. \tag{18}$$

Finally, by applying $h \ll \epsilon, c, B_{LD}$ the above equation simplifies to

$$\frac{\Delta\chi_{ab}(\epsilon, c)_E}{T} = \frac{\pi\mu_0 k_B \xi_{ab}^2(0)}{3\phi_0^2 s}\left[\frac{1}{\epsilon}\left(1 + \frac{B_{LD}}{\epsilon}\right)^{-1/2} - \frac{1}{c}\right] \tag{19}$$

which corresponds to $\Delta\chi_{ab}(\epsilon)$ in the Schmidt-limit and under a total-energy cut-off. Here again, equation (2) (with $N_e = 1$) is recovered by simply imposing $\epsilon, B_{LD} \ll c$.

2.3. The ratio $\Delta\chi_{ab}(\epsilon)/(T\Delta\sigma_{ab}(\epsilon))$ under different cut-off conditions

The ratio $\Delta\chi_{ab}(\epsilon)/(T\Delta\sigma_{ab}(\epsilon))$ in layered superconductors (and in the low-magnetic-field limit) under different cut-off conditions may now be directly obtained by using the expressions calculated above. For instance, in the case of a total-energy cut-off, we obtain

$$\frac{\Delta\chi_{ab}(\epsilon, c)_E}{T\Delta\sigma_{ab}(\epsilon, c)_E} = 2.79 \times 10^5 \xi_{ab}^2(0)\left\{1 + (c - \epsilon - B_{LD}/2)\right.$$
$$\left. \times\left(\frac{c^2}{\epsilon}\left(1 + \frac{B_{LD}}{\epsilon}\right)^{-1/2} - 2c + \epsilon + B_{LD}/2\right)^{-1}\right\}. \tag{20}$$

Note that by imposing $\epsilon, B_{LD} \ll c$ in the above equation we recover equation (3). Moreover, as we are interested on the cut-off effects in the high-reduced-temperature region, this type of relationship may be strongly simplified by assuming $B_{LD} \ll \epsilon, c$. This leads (in MKSA units) to

$$\frac{\Delta\chi_{ab}(\epsilon, c)_M}{T\Delta\sigma_{ab}(\epsilon, c)_M} = 2.79 \times 10^5 \xi_{ab}^2(0)\left(1 + \frac{\epsilon}{c}\right) \tag{21}$$

and

$$\frac{\Delta\chi_{ab}(\epsilon, c)_E}{T\Delta\sigma_{ab}(\epsilon, c)_E} = 2.79 \times 10^5 \xi_{ab}^2(0)\left(1 + \frac{\epsilon}{c - \epsilon}\right) \tag{22}$$

for, respectively, a momentum cut-off and a total-energy cut-off. As in the case of equation (3), these two expressions do not depend on any effective periodicity length. Also note that in the absence of any cut-off condition (i.e. with $c \to \infty$) or in the low-reduced-temperature region (i.e. $\epsilon \ll c$), both expressions reduce to equation (3). In contrast, strong differences with equation (3) appear in the high-reduced-temperature region, i.e. when ϵ becomes of the order of c, mainly under a total-energy cut-off. In this last case, the ratio between $\Delta\chi_{ab}(\epsilon)$ and $\Delta\sigma_{ab}(\epsilon)$ diverges when ϵ becomes of the order of c. Let us finally stress that these two expressions apply with all generality to multilayered strongly anisotropic HTSC, whose thermal fluctuations are in the 2D-limit in the accessible ϵ-region.

3. Application to Y-123

As an example of the usefulness of the above results, in figure 1 we summarize the comparison between the theoretical $\Delta\chi_{ab}(\epsilon)/(T\Delta\sigma_{ab}(\epsilon))$ predictions under different cut-off conditions and the experimental results obtained in almost optimally doped Y-123 superconductors. In these bilayered superconductors, which are moderately anisotropic, the thermal fluctuations behave within the 2D- and the 3D-limit. However, the Josephson coupling between the different neighbour layers are expected to be similar to each other [2]. Therefore, we may use the theoretical results for single-layered superconductors, but with $s/2$ instead of s as an effective periodicity length. A more detailed account of these experiments is going to be published elsewhere. Let us note here that the error bars in figure 1 are mainly due to the uncertainties in the estimation of the normal (or background) contributions to the measured $\chi_{ab}(T)$ and $\sigma_{ab}(T)$. To correctly analyse the high-ϵ region, such background contributions must be estimated by extrapolating through the transition the $\chi_{ab}(T)$ and the $\sigma_{ab}(T)$ data measured as far as possible from T_{c0}. So, these data have been obtained by using backgrounds estimated above $T \gtrsim 225$ K (which corresponds to $\epsilon \gtrsim 0.9$).

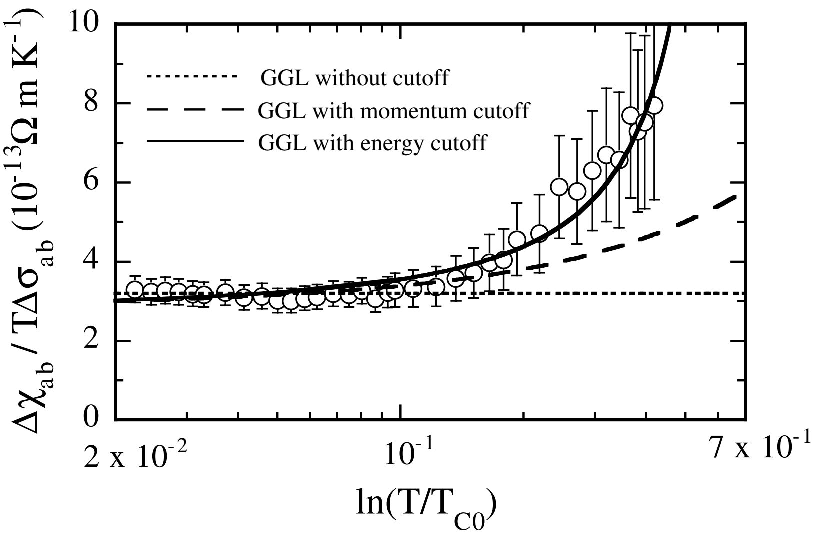

Figure 1. Comparison in the accessible ϵ-region above 2×10^{-2} between the $\Delta\chi_{ab}(\epsilon)/(T\Delta\sigma_{ab}(\epsilon))$ ratio measured in $YBa_2Cu_3O_{7-\delta}$ samples and the theoretical ratio obtained by imposing different cut-off conditions to the conventional Gaussian–Ginzburg–Landau approach for layered superconductors . These results illustrate how the GGL approach extends its validity to the high-ϵ (for ϵ above 10^{-1}) region when a total-energy cut-off is included in the fluctuation spectrum.

The dotted line in figure 1 corresponds to the GGL theory without any cut-off (equation (3)). As expected [2–5], the agreement between equation (3) and the data is excellent in the region $2\times10^{-1} \lesssim \epsilon \lesssim 10^{-1}$ and it leads to $\xi_{ab}(0) = 11$ Å. However, it can be clearly seen that the experimental data do not follow such a constant behaviour at higher temperatures. The solid line in this figure corresponds to the best fit of equation (20) to the experimental data in the ϵ-region $2\times10^{-2} \leqslant \epsilon \leqslant 5\times10^{-1}$ with $\xi_{ab}(0)$, $\xi_c(0)$, and c as free parameters. As can be seen, the agreement is excellent in almost the entire ϵ-region and it leads to $\xi_{ab}(0) \simeq 10$ Å, $\xi_c(0) \simeq 1.1$ Å (which are well within the accepted values [2]) and $c = 0.7$, which is comparable with the values of c that we have found recently when analysing the fluctuation induced diamagnetism in the high magnetic field regime in another cuprate superconductor [17]. The dashed line corresponds to the GGL theory with a momentum cut-off, i.e. to equation (16) divided by equation (8), with again $\xi_{ab}(0) \simeq 10$ Å, $\xi_c(0) \simeq 1.1$ Å and $c = 0.7$. The disagreement for $\epsilon \gtrsim 0.2$ is well beyond the experimental uncertainties, and these differences at high-reduced-temperatures cannot be mitigated by using other values of $\xi_{ab}(0)$, $\xi_c(0)$ and c without breaking the good agreement for $2 \times 10^{-2} \lesssim \epsilon \lesssim 10^{-1}$. These results strongly suggest, therefore, that to overcome the difficulties associated with the short-wavelength fluctuations, which mainly manifest in the high-reduced-temperature region ($\epsilon \gtrsim 0.1$), one must introduce in the GGL approach a total-energy cut-off, instead of the conventional momentum cut-off used until now to analyse the thermal fluctuations in LTSC [8] and HTSC [9].

4. Conclusions

In conclusion, to take into account the presence of the short-wavelength effects that appear at high-reduced-temperatures, the in-plane paraconductivity, the fluctuation induced diamagnetism, and the relationship between both observables have been calculated on the grounds of the Gaussian–Ginzburg–Landau approach for layered superconductors by imposing, for the first time, a momentum and a total-energy cut-off condition. Although the central motivation of this paper was the presentation of these theoretical results, we have also summarized here a preliminary comparison with the experimental $\Delta\chi_{ab}(\epsilon)/(T\,\Delta\sigma_{ab}(\epsilon))$ data recently obtained in $YBa_2Cu_3O_{7-\delta}$ samples. This comparison strongly suggests the adequacy of the total-energy cut-off condition to explain the behaviour of these fluctuation effects in the high-reduced-temperature region ($\epsilon \gtrsim 0.1$). This total energy is obtained by adding the kinetic energy to the confinement energy of the Cooper pairs formation. A more thorough comparison of these theoretical results with different experimental data in $YBa_2Cu_3O_{7-\delta}$ superconductors, and those we are now measuring in other HTSC and LTSC, will be presented elsewhere. However, the present results already clearly suggest that the calculations presented here open a promising way for the understanding, at a phenomenological level, of the thermal fluctuations of Cooper pairs in superconductors in the ϵ-region not too close to the superconducting transition, including the high-reduced-temperature region.

Acknowledgments

This work has been supported by the CICYT, Spain, under grant MAT98-0371, by the CICYT-FEDER Funds, Spain, under grant 1FD97-0146 and by Unión Fenosa under grant 0666-98. F V wishes to acknowledge the hospitality of Professor Yao Liang during his stay at IRC at the University of Cambridge as BBV Foundation Visiting Professor.

References

[1] See, for instance,
Tinkham M 1996 *Introduction to Superconductivity* (New York: McGraw-Hill) ch 9
[2] See, for instance,
Vidal F and Ramallo M V 1998 *The Gap Symmetry and Fluctuations in High-T_c Superconductors* ed J Bok *et al* (New York: Plenum) p 443
Vidal F, Torrón C, Viña J and Mosqueira J 2000 *Physica* C **332** 166 and references therein
[3] Ramallo M V, Torrón C and Vidal F 1994 *Physica* C **230** 97
Ramallo M V, Pomar A and Vidal F 1996 *Phys. Rev.* B **54** 4341
[4] Vidal F, Torrón C, Veira J A, Miguélez F and Maza J 1991 *J. Phys.: Condens. Matter* **3** L5219
[5] Ramallo M V and Vidal F 1999 *Phys. Rev.* B **59** 4475 and references therein

[6] Ramallo M V and Vidal F 1997 *Europhys. Lett.* **39** 177
[7] See, for instance,
Skocpol W J and Tinkham M 1975 *Rep. Prog. Phys.* **38** 1049
[8] Johnson W L and Tsuei C C 1976 *Phys. Rev.* B **13** 4827
Johnson W L, Tsuei C C and Chaudhari P 1978 *Phys. Rev.* B **17** 2884
[9] Freitas P P, Tsuei C C and Plaskett T S 1987 *Phys. Rev.* B **36** 833
Hopfengärtner R, Hensel B and Saemann-Ischenko G 1991 *Phys. Rev.* B **44** 741
Gauzzi A and Pavuna D 1995 *Phys. Rev.* B **51** 15420
[10] Carballeira C, Currás S R, Viña J, Veira J A, Ramallo M V and Vidal F 2001 *Phys. Rev.* B to be published
[11] For some considerations about these aspects of the GGL approach see Tinkham M in section 8.4 of [1]. See also [7] §3.1.
[12] See, respectively,
Patton B R, Ambegaoker V and Wilkins J W 1969 *Solid State Comm.* **7** 1287
Nam S B 1971 *Phys. Rev. Lett.* **26** 1369
[13] Schmidt H 1968 *Z. Phys.* **216** 336
See also
Abrikosov A A 1988 *Fundamentals of the Theory of Metals* (Amsterdam: North-Holland) p 478
Cyrot M 1973 *Rep. Prog. Phys.* **36** 103
[14] Lawrence W E and Doniach S 1970 *Proc. 12th Int. Conf. Low-Temp. Phys. (Kyoto)* ed E Kanda (Tokyo: Keigatu) p 361
See also [3] and
Klemm R A 1990 *Phys. Rev.* B **41** 2073
[15] For a detailed derivation of equation (12) in the 3D-case see, for instance,
Landau L and Lifchitz E 1978 *Statistical Physics* vol 2 (Oxford: Pergamon) section 49.
[16] An expression similar to equation (12) was already used by Baraduc and coworkers to calculate $\Delta\chi_{ab}(\epsilon)$ in layered superconductors without any cut-off. The 2D-case was studied, also without cut-off, by Koshelev. See, respectively,
Baraduc C, Buzdin A, Henry J-Y, Brison J-P and Puech L 1995 *Physica* C **248** 138
Koshelev A E 1994 *Phys. Rev.* B **50** 506
[17] Carballeira C, Mosqueira J, Revcolevschi A and Vidal F 2000 *Phys. Rev. Lett.* **84** 3157