Existence of two dissipation peaks in a superconducting glass model

1993 J. Phys. A: Math. Gen. 26 L535

(http://iopscience.iop.org/0305-4470/26/12/010)

J. Phys. A: Math. Gen. 26 (1993) L535–L541. Printed in the UK

LETTER TO THE EDITOR

Existence of two dissipation peaks in a superconducting glass model

Jorge V José† and Guillermo Ramirez-Santiago‡

† Department of Physics, Northeastern University, Boston, Massachusetts 02115, USA
‡ Instituto de Física, Universidad Nacional Autónoma de México, Postal 20-364, 01000 México DF, Mexico

Received 9 March 1993

Abstract. A superconducting glass model with unidirectional correlated disorder is studied via extensive Monte Carlo simulations in two and three dimensions. A calculation of the magnetic fluctuations shows two dissipation peaks at temperatures T_1 and T_2, with $T_1 < T_2 < T_c$. Results for the 3D phase correlation functions indicate that the $T^- < T_1$ phase is less ordered than the $T_1 < T^+$ phase. The T_1 transition is found to be due to an enhancement of short-range phase correlations, while the one at T_2 is longer ranged. The possible connection of these results to recent experiments in high-T_c oxide superconductors is discussed.

One of the most important goals in the theoretical study of high-temperature superconductors has been to understand the magnetic field against temperature phase diagram (H against T diagram). The existence of an irreversibility line (IL) in the H against T diagram was first reported by Müller *et al* [1]. Their results were later experimentally confirmed in virtually all other high-T_c superconductors. In a recent series of experiments, however, evidence for more than one thermal instability in the H against T diagram has been reported [2–4]. These results are in striking contrast with previous studies that, for a given experiment, have reported seeing only one thermal instability in the H against T diagram. In YBaCuO compounds (YBCO) a second thermal instability was seen as a 'knee' in resistivity measurements [2], while in more recent experiments evidence for having a melting transition *above* the IL has been presented [3]. Two dissipation peaks were also seen in experiments on BiSrCaCuO (BSCCO) single crystals when an external magnetic field *was rotated* with respect to the $\hat{c}$ axis [4]. A recent paper by Brandt [5] provides a possible explanation for these BSCCO experiments in terms of *geometrical* effects that depend directly on the physical dimensions of the samples used in the experiments. Most theoretical studies based on the superconducting glass model suggested by Müller *et al* have so far concentrated in calculating its thermodynamic properties. However, in the torsional oscillator and transport experiments what is measured is related to fluctuations, for example correlation functions.

Motivated by these experimental results we have carried out a detailed Monte Carlo (MC) analysis of the *fluctuations* in the magnetic and spatial correlation functions of a 2D and 3D *superconducting glass model* [6, 7] (SGM), that entails including correlated disorder (CSGM) [8] with a magnetic field *perpendicular* to the layers. Previous studies of the remanent magnetic properties of these types of models were successful in yielding qualitative agreement with experiments in the oxide superconductors [7, 8], including well-separated branches of zero-field-cooled (M_{ZFC}) and field-cooled (M_{FC}) magnetizations [8].

0305-4470/93/120535+07$07.50

Variations of the SGM have also been used as the basis of quantitative studies that consider the possibility of an equilibrium second-order vortex glass transition [9]. There are also recent studies in *non-disordered* models that show two *melting* transitions [10]. The present study is fundamentally different in that disorder plays an essential role.

In the calculations described here we found two thermal instabilities in the ZFC branch at finite temperatures T_1 and T_2 with $T_1 < T_2 < T_c$, where T_c is the critical temperature above which M is essentially zero. In contrast to Brandt's explanation of the BSCCO experiments, the two thermal instabilities found here are intrinsic to the model and appear when the magnetic field is perpendicular to the layers. The two thermal instabilities, or two dissipation peaks, are an inherent property of the CSGM, not present in the uncorrelated case, and are clearly seen only in the fluctuations of the magnetization and phase correlation functions in our 2D and 3D calculations. This fact shows that the two thermal instabilities are directly related to the vortex properties of the model. Clear differences between the 2D and 3D results emerge from the calculation of the phase correlation functions. We present strong evidence that in three dimensions the T_1 instability entails an increase in the local coherence of the phases, while the one at T_2 involves longer-ranged correlations. Furthermore, we find that the phase correlations are smaller in the $T^- < T_1$ phase than in the $T_1 < T^+$ phase, indicating that the $T^- \to T^+$ entails a disorder→order transition. We shall discuss the similarities and differences between our results and those found in recent experiments at the end of this letter.

The CSCG model studied here is defined by the Hamiltonian

$$\mathcal{H} = \sum_{ij,z} \{[E^{\parallel}[1 - \cos(\phi_i(z) - \phi_j(z) + 2\pi f_{ij}(z))] + E^{\perp}[1 - \cos(\phi_i(z) - \phi_i(z+1))]\} \quad (1)$$

with i, j denoting two-dimensional vectors and z the distance perpendicular to the layers; $\phi_i(z)$ stands for the phase of the Ginzburg–Landau order parameter of the ith (z) 'grain'; $E^{\parallel}$ and $E^{\perp}$ are the intralayer and interlayer Josephson coupling constants, respectively; the link variable $f_{ij}(z) = (1/\Phi_0)\int_{i(z)}^{j(z)} \boldsymbol{A}\cdot \mathrm{d}\boldsymbol{l}$, with $\boldsymbol{A}$ the magnetic vector potential, and Φ_0 the quantum of flux. In our calculations we use the Landau gauge $\boldsymbol{A} = (0, Hx, 0)$. The frustration parameter, F, is then defined as $F = \sum_P f_{ij}$, with $\sum_P$ the sum over plaquettes. We will measure energies normalized by $E^{\parallel}$ and fields by Φ_0/a_0^2, where a_0 is the lattice spacing, so that we write $F = H$. The disorder in the CSGM is introduced by independently displacing the y columns in each plane, with initially square lattice geometry, by $x_i = ia_0 + ra_0$ with r a random number uniformly distributed in the interval $[-\frac{1}{2}\delta, \frac{1}{2}\delta]$. In the 3D calculations the disorder along the y columns is uncorrelated between planes. Since the amount of disorder in the model can be changed by varying δ or H, we fix $\delta = 0.1$ in our calculations, as in [8], and vary only H. This type of disorder has the effect of preferentially enhancing the trapping of vortices in the system along the y axis. Another type of correlated disorder, of a columnar type, has been considered in recent experiments and theory [11]. The similarity between the columnar disorder and the one considered here is that the correlation enhances the pinning along preferential directions.

Near thermodynamic equilibrium, the linear response functions are described in terms of susceptibilities which in turn are related to the thermodynamic fluctuations in the system. Thus, to make contact with the quantities measured in, say, the oscillator experiments, we present results for the fluctuations of the magnetization as a function of field and temperature. We also present results of calculations of the fluctuations of the quenched gauge-invariant phase correlation functions, related to the current–current correlation functions of the model, which allow us to gain a microscopic understanding of the nature of the thermal instabilities found in our calculations.

The quenched averaged magnetization is obtained from calculating

$$M = \frac{1}{N}\left[\left\langle \sum_{\langle i,j,z\rangle} x^2_{i(z),j(z)} \sin\left(\phi_i(z) - \phi_j(z) + 2\pi f_{ij}(z)\right) \frac{E^{\parallel}}{\phi_0} X_{i(z),j(z)} \right\rangle\right]_{\rm c} \tag{2}$$

with $x_{i(z),j(z)} = \frac{1}{2}(x_{i(z)} + x_{j(z)})$ and $X_{i(z),j(z)} = \frac{1}{2}(x_{i(z)} + x_{j(z)})(y_{i(z)} - y_{j(z)})$. The bracket $[,]_{\rm c}$ denotes the configurational average over an ensemble of different realizations of the disordered lattice; $\langle\ ,\rangle$ stands for the thermodynamic average and $\langle i, j, z\rangle$ for a nearest-neighbour sum in the planes (2D) and along the z axis (3D). The quenched averaged gauge-invariant phase correlation function is defined by

$$G(i(z), j(z')) = \left[\left|\left\langle \exp({\rm i}\phi_i(z)) \prod_{\Gamma} \exp({\rm i}f_{i,i+1}(z)) \exp(-{\rm i}\phi_j(z'))\right\rangle\right|^2\right]_{\rm c}. \tag{3}$$

Here Γ denotes a path connecting the points $i(z)$ to $j(z')$. It is convenient to calculate the zero-momentum correlation function, which is known to have only one asymptotic correlation length, or Lyapunov exponent [12]. Note that, since the disorder is along the y columns, the $G(r)$ along the x and y directions are expected to show different behaviours, as found in our calculations. This is typical of gauge-invariant correlation functions which are path dependent [12].

We discuss first our 2D results. We started by calculating the ZFC and FC branches of the magnetization following a similar procedure as in the experiments, e.g. [1]. As found from a finite-size analysis in [8], we take a lattice of size 16×16 as representative of the properties of the model. Since we are interested in calculating fluctuations here, which are notoriously harder to evaluate than thermodynamic properties, we need to have long runs and vary the temperature very slowly. Specifically, the ZFC branch is obtained by first equilibrating at $T = 0.10$ in zero field; the field is then switched on and the system is warmed up first in steps of $\Delta T = 2\times 10^{-2}$ from $T = 0.10$ up to $T = 0.90$, and then in steps of $\Delta T = 0.10$ from $T = 0.90$ to $T = 1.5$. The $T_{\rm c}$ for this lattice size is $T_{\rm c} \sim 1.2$. The FC branch is obtained by cooling at the same variable rate down to $T = 0.10$, for the same H. Typically, the system was allowed to equilibrate for 10K MCS/site at each temperature and the averages were calculated over 60K MCS/site. A complete scan of ZFC plus FC branches consisted of 92 temperatures, entailing about 6.4×10^6 MCS/site. This process was repeated for five different configurations of disorder to obtain the quenched averages. We found that with this many temperatures and members of the ensemble we got statistically significant results for the 2D calculations.

In figure 1 we show representative results for the thermal fluctuations of $M_{\rm ZFC}$, σ_M/T, as a function of the renormalized temperature $T/T_{\rm c}$ and H. The fluctuations were calculated from the standard expression,

$$\sigma_M = \frac{1}{\sqrt{Nm-1}}\left|\frac{1}{m}\frac{1}{N}\sum_{j=1}^{m}\sum_{k=1}^{N} M_k^2 - \frac{1}{m}\sum_{j=1}^{m}\left(\frac{1}{N}\sum_{k=1}^{N} M_k\right)^2\right|^{1/2}. \tag{4}$$

Here m is the number of disordered samples in the ensemble and N stands for the number of subsets in the thermodynamic averages. The curves shown in figure 1 were obtained from smoothing the raw data obtained from equation (4) using a cubic-spline fit. This process gives different weights to the data points and, for example, figures 1(a) and (b) show a

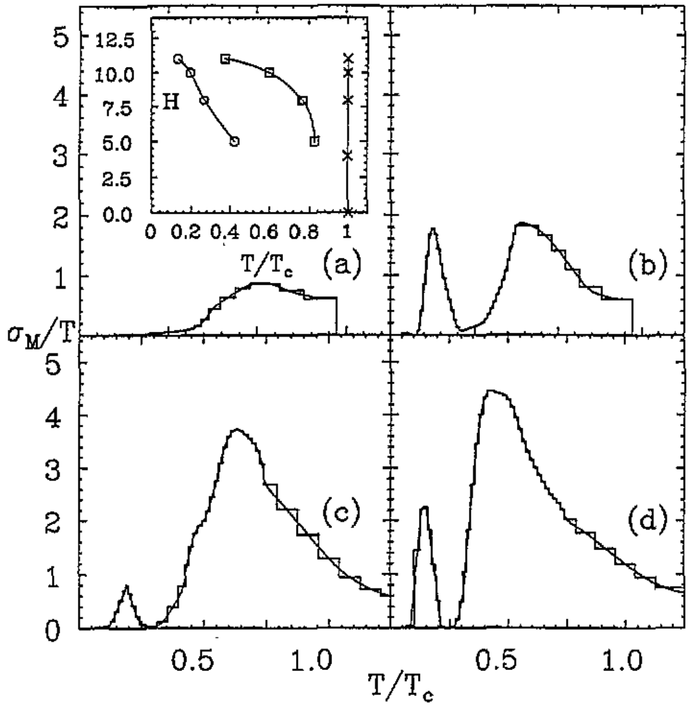

Figure 1. 2D ZFC magnetization fluctuations σ_M/T plotted against T/T_c, for different fields H. Results obtained from averages over five configurations of 16×16 lattices. The fields are: (*a*) $H = 0$, (*b*) $H = 0.008$, (*c*) $H = 0.010$ and (*d*) $H = 0.011$. The (*a*)–(*d*) vertical axis should be multiplied by 10^{-2} and the vertical axis in the inset by 10^{-3}. The inset shows $T_1(H)$ (○), $T_2(H)$ (□) and $T_c(H)$ (×) determined from an approximate estimate of the maxima in σ_M. The curves are a guide to the eye. More details are found in the text.

jump for $T/T_c \sim 1$ since the spline fit gives a very small weight for the σ_M values for $T/T_c \leqslant 1$. We clearly see that, for $H \neq 0$, there are two thermal instabilities in σ_M as a function of T. The lower peak corresponds to the temperature $T_1(H)$ where M_{ZFC} and M_{FC} join. This temperature has been used as the defining boundary between ergodic and non-ergodic behaviour [1]. The second increase in fluctuations occurs about $T_2(H)(< T_c)$, the width of the $T_1(H)$ peak being narrower than the one at $T_2(H)$. For values of $H \leqslant 0.005$ (below the lowest points for T_1 and T_2 shown in the inset) the two branches of M coincide corresponding to a purely diamagnetic Meissner phase. Note that, in the case when $H = 0$, there is only one maximum for the fluctuations, thus showing that disorder is essential to the existence of the two thermal instabilities. The $H = 0$ peak is not negligible since the 2D model has larger fluctuations than in three dimensions. This difference is clear from looking at the $H = 0$ result in figure 2(*a*). In contrast to the peaks in the M_{ZFC}, the results for the fluctuations of M_{FC}, not shown here, exhibit a monotonically increasing behaviour as T decreases, as would be expected from an equilibrium susceptibility with perhaps a $T = 0$ critical point. This important difference in the dissipation structure between the ZFC and FC magnetization branches also appears to agree with recent magnetization data obtained in BSCCO single crystals [13]. In the inset of figure 1 we show an H against T/T_c phase diagram obtained from an approximate determination of the maxima of the two thermal instabilities. A fit to our limited numerical results for $T_1(H)$ is not possible since it appears to have inflection points. A reasonable fit to the power law $H = A(T - T_c)^\alpha$ is obtained, however, for $T_2(H)$ with $A = 0.12$, $T_c = 1.199$ and the exponent $\alpha = 0.35$. In doing this fit we have assumed that $T_2(H = 0) = T_c(H = 0)$. We should emphasize that the existence of the second peak is clear only from the analysis of the fluctuations of M_{ZFC} and not from M_{ZFC} itself.

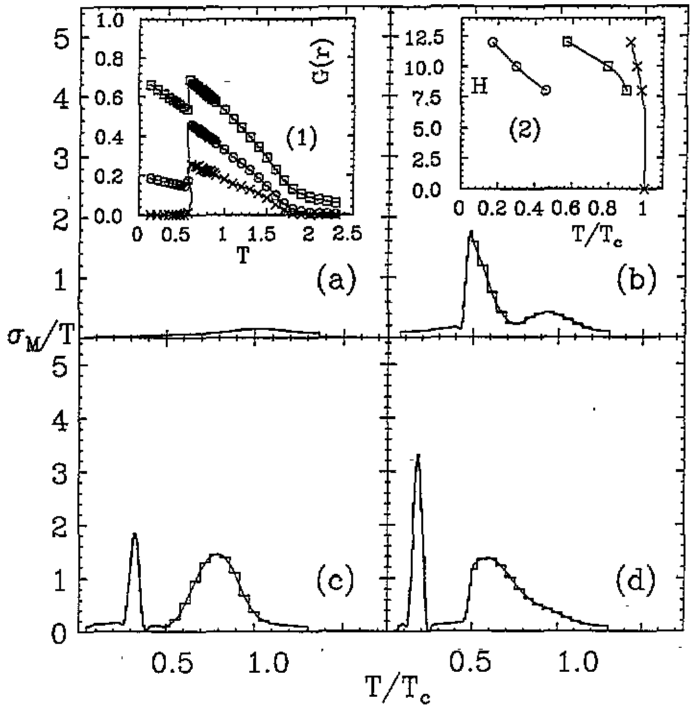

Figure 2. The same as in figure 1 for three dimensions including inset (2), for a lattice of dimension $16 \times 16 \times 5$ with $E^{\perp}/E^{\parallel} = 0.5$. Here (*a*) $H = 0$, (*b*) $H = 0.008$, (*c*) $H = 0.01$ and (*d*) $H = 0.012$. The scales are the same as in figure 1. Inset (1) shows results for the ZFC $G_{\parallel}(r)$ as a function of r, for $r = 1$ (□), $r = 2$ (○) and $r = 3$ (×) for $H = 0.01$. Note that $G_{\parallel}(r)$ is larger above T_1.

We now discuss our quasi-3D results. In figure 2 we show the corresponding results for σ_M/T as a function of T/T_c ($T_c \sim 1.8$) for different values of H, with $E^{\perp}/E^{\parallel} = 0.5$. Again, since we wanted to calculate fluctuations in the 3D model we had to find a compromise between many temperatures and ensemble averages. We opted for doing simulations in one system but with many closely spaced temperatures. We found good statistics for a lattice of size $16 \times 16 \times 5$ with 116 temperatures, with about the same number of MCS/site as in two dimensions, for a total of 8.1×10^6 MCS/site for a complete cycle, which entails about 10^{10} updates for each H. We followed the same procedure of warming and cooling as in the 2D calculations. The results of figure 2 show a two-peaked structure in the fluctuations of M_{ZCF}, with the peak at $T_1(H)$ more pronounced than the one at $T_2(H)$, in contrast to the 2D result. The rate of change of the peak widths as a function of field and temperature is qualitatively similar to that found in two dimensions, including the fact that the two-peak structure is not seen in the FC branch of the magnetization. In this case fits to the critical temperature of the form $A(T - T_c(H))^\alpha$ are more reasonable than in two dimensions and give the following results: $H = 0.8(T_1(H) - 1.83)^{0.931}$, $H = 0.13(T_2(H) - 1.83)^{0.3}$ and $H = 0.24(T_c(H) - 1.83)^{0.34}$.

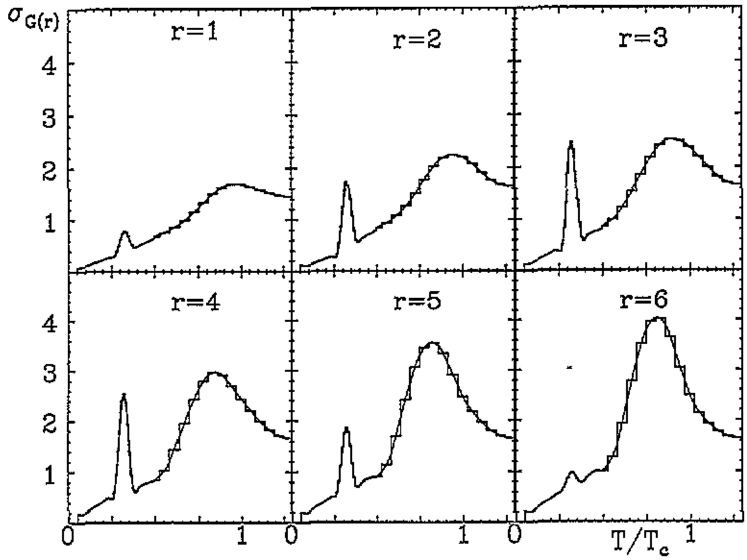

Figure 3. 3D $\sigma_{G(r)}$ plotted against T/T_c for different distances r. The parameters are the same as in inset (1) in figure 2.

We now discuss the results from the zero-momentum correlation function calculations, and their fluctuations, in two and three dimensions along the y and x directions, denoted as $G_\parallel(\boldsymbol{r})$ and $G_\perp(\boldsymbol{r})$, respectively. Recall that $\boldsymbol{A} \parallel y$. In the $T < T_1(H)$ region, we find that the $G_\perp$ as a function of r have a monotonic exponential decay, with correlation length decreasing as T increases. The ZFC and FC branches of $G_\perp(\boldsymbol{r})$ are essentially identical to each other in this temperature range. In contrast, the ZFC and FC branches of $G_\parallel(\boldsymbol{r})$ are different, with the ZFC branch decaying faster than the FC branch. This difference disappears for $T_1 < T$, both in two and three dimensions. The ZFC branch of $G_\parallel(\boldsymbol{r})$ as a function of T for fixed r shows a discontinuity about $T_1(H)$, that tends to zero for $r \geqslant 4$. This behaviour is shown for the 3D case in inset (1) of figure 2. We note the important fact that the correlations are smaller just below T_1 and larger above. This means that the phases are less correlated for $T < T_1$ than above. As in the magnetization case, there is no evidence in the ZFC $G_\parallel(\boldsymbol{r})$ for a T_2 instability. The T_2 instability becomes evident only by calculating the fluctuations of the ZFC $G_\parallel(\boldsymbol{r})$, $\sigma_{G(r)}$. In figure 3 we show the results for the 3D $\sigma_{G(r)}$, as a function of T for different values of r. A two thermal instabilities structure is seen in this figure, as in the σ_M case. However, we note that the $T_1(H)$ peak has a maximum for $r \sim (3, 4)$, while the peak at $T_2(H)$ increases monotonically and becomes narrower as r grows. The results for $G_\parallel(\boldsymbol{r})$ and $\sigma_{G(r)}$ indicate that the $T_1(H)$ instability separates a disordered from an ordered phase and is connected to an increase in the short-range phase correlation while the one at $T_2(H)$ involves an increase of longer-ranged correlations. By contrast, the corresponding 2D calculation of $\sigma_{G(r)}$ shows a small peak at $T_1(H)$ superimposed on a broad background, both of which remain almost of the same magnitude as r increases.

In conclusion, we have presented evidence for the existence of two thermal instabilities in a superconducting glass model with correlated disorder. Our results suggest a depinning transition about T_1 with no free vortex motion until there is enough thermal energy to induce longer-ranged motions about T_2. The second instability would appear to better fit the melting transition scenario. Although the model considered does not have an explicit representation to the experimental systems studied, the results are qualitatively analogous to the experimental results found in YBCCO [1], with a depinning→melting transition. More work is needed to ascertain the explicit relevance of the results found here to the experimental findings. An extensive discussion of these and other results will appear elsewhere [14].

We want to thank F de la Cruz and collaborators as well as R Markiewicz for very valuable discussions. This work has been supported in part by NSF grants DMR-9211339, INT-NSF-89-01622, and by the NSF Pittsburgh Supercomputing Center (JVJ). GRS was supported by a CONACYT fellowship, NSF-CONACYT grant no G001-1720/001328 and by the Supercomputing Center, DGSCA, of UNAM, Mexico.

References

[1] Muller K A, Takashige and Bednorz J G 1987 *Phys. Rev. Lett.* **58** 1143
[2] Worthington T K *et al* 1990 *Cryogenics* **30** 417
[3] Beck R G *et al* 1992 *Phys. Rev. Lett.* **68** 1594
Worthington T K *et al* 1992 *Bull. Am. Phys. Soc.* **37** 592; 1992 unpublished
Farrell D E *et al* 1991 *Phys. Rev. Lett.* **67** 1165
Safar H *et al* 1991 *Phys. Rev. Lett.* **68** 2672
[4] Durán C J *et al* 1991 *Phys. Rev.* B **44** 7737
Yasyi J *et al* 1991 *Physica* C **184** 254
Tuominen M *et al* 1990 *Phys. Rev.* B **42** 8740
Testelin C and Markiewicz R S 1993 to be published
[5] Brandt E H 1992 *Phys. Rev. Lett.* **68** 3769
[6] Ebner C and Stroud D 1985 *Phys. Rev.* B **31** 165
Shi W Y *et al* 1984 *Phys. Rev.* B **30** 1341
Chakrabarty A and Dasgupta C 1988 *Phys. Rev.* B **37** 7557
[7] Morgenstern I, Muller K A and Bednorz J G 1987 *Z. Phys* B **69** 33
[8] Choi J and José J V 1989 *Phys. Rev. Lett.* **62** 320
José J V 1991 *Proc. Int. Conf. on Transport Properties of Superconductors* ed R Nicolsky (Singapore: World Scientific)
[9] Fisher M P A 1989 *Phys. Rev. Lett.* **62** 1415
Huse D A and Seung H 1990 *Phys. Rev.* B **42** 1059
Cieplak M *et al* 1991 *J. Phys. A: Math. Gen.* **24** L145
Reger J D *et al* 1991 *Phys. Rev.* B **44** 7147
Gringas M J P 1991 *Phys. Rev.* B **44** 7139
[10] Ryu S *et al* 1992 *Phys. Rev. Lett.* **68** 710
[11] Civale L *et al* 1991 *Phys. Rev. Lett.* **67** 648
Nelson D R and Vinokur V M 1992 *Phys. Rev. Lett.* **68** 2398
[12] Ramirez-Santiago G and José J V 1991 *Phys. Rev. Lett.* **68** 1224
[13] Pastoriza H and de La Cruz F Private communication
[14] José J V and Ramirez-Santiago G 1993 in preparation