Fluctuation-induced diamagnetism below the critical temperature in high-T_c superconductors

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(Received 5 May 1998)

We have calculated the fluctuation-induced diamagnetism ($\Delta\chi$) of the high-T_c superconductors below the critical temperature T_c and in the weak-field region. The theoretical formulation estimates the relation between $\Delta\chi$ and the reduced temperature. [S0163-1829(98)02242-5]

I. INTRODUCTION

The thermodynamic fluctuations are important in high-T_c superconductors (HTSC) because of their short coherence length. There are numerous studies both theoretical as well as experimental of the fluctuation-induced conductivity, magnetoconductivity, specific heat, and diamagnetism above the critical temperature.[1,2] The fluctuation in the diamagnetic susceptibility was first theoretically studied by Schmidt above the critical temperature.[3] According to this formulation the diamagnetic susceptibility above T_c caused by the thermal fluctuation is expressed as follows:

$$\chi_D \propto (T-T_c)^{-1}. \tag{1}$$

Above the critical temperature, the fluctuation in diamagnetism is calculated by Schmid starting from the Ginzburg-Landau free-energy functional.[4] According to this formulation the fluctuation in diamagnetism is expressed as follows:

$$\chi = -\frac{\pi T \xi_{\parallel}^2(0)}{3 d \phi_0 \epsilon}. \tag{2}$$

Here ϵ is the reduced temperature defined as

$$\epsilon = \frac{T-T_c}{T_c}$$

and T_c is the critical temperature where a HTSC starts to show negative susceptibility during the cooling. d and $\xi_{\parallel}(0)$ are the interlayer seperation and the ab-plane coherence length at absolute zero, respectively, and ϕ_0 is the flux quanta. Koshelev has estimated the magnetization as a function of the applied magnetic field and reduced temperature of high-T_c superconductors above the critical temperature using the Ginzburg-Landau free-energy functional.[5] The magnetization is estimated both for the high and low magnetic fields. Baraduc *et al.* have reported the fluctuation in magnetization in the high-T_c superconductors above the critical temperature with the help of the Lawrence-Doniach free-energy functional.[6] Klemm has described fluctuation in diamagnetism in multiperiodic layered superconductors.[7] Lee, Klemm, and Johnston have reported the fluctuation in diamagnetism of biperiodic superconductors above T_c. The result shows that fluctuation is observed from around $2T_c$ in the weak-field regime.[8] The magnetization measurement of Bi-2212 up to 5 T below T_c by Kes *et al.* has revealed that quasi-two-dimensional fluctuations are responsible for the field-induced suppression of phase transition.[9] Tesanovic *et al.* have studied the thermodynamic quantities in HTSC in the critical region by taking into account the high magnetic field.[10] Ramallo, Torron, and Vidal have studied the fluctuation-induced diamagnetism above T_c in biperiodic HTSC systems.[11] Dorin *et al.* have studied fluctuation conductivity of layered superconductors in a magnetic field parallel to the c direction. The contributions from the fluctuation in the single quasiparticle density of states (DOS) and the Maki-Thompson contribution are compared to the Aslamazov-Larkin contribution in terms of an effective quasiparticle nearest-neighbor interlayer hopping energy.[12]

Skocpol and Tinkham have discussed the fluctuation below the critical temperature.[1] Bulaevskii, Ginzburg, and Sobyanin have suggested that the fluctuation can be studied below the critical temperature in the layered superconductors.[13] Later Bulaevskii, Ledvij, and Kogan have reported the fluctuation of vortices in layered superconductors.[14] The entropy contribution to the free energy is not too small in a broad temperature region below the critical temperature and they have calculated magnetization for a magnetic field $B \gg B_{cr}$ below the critical temperature. According to this formulation magnetization is field and temperature dependent.[14] Varlamov and Livanov have investigated the effect of superconducting fluctuations on the electronic part of the thermoelctric force and thermal conductivity at a temperature above the critical temperature.[15] They have reported that thermoelectric force possesses a maxima at the critical temperature and then rapidly decreases to zero in the superconducting state. In a recent paper Houssa *et al.* have described the fluctuation contribution in thermal conductivity below the critical temperature.[16] Below the critical temperature thermal conductivity is enhanced both from the electronic and phonon contribution. A dimensional crossover from three to two dimensions is observed for normalized fluctuation contribution below the critical temperature. The Ginzburg-Landau parameter is extracted in this paper from the crossover temperature and the critical temperature.[16] To our knowledge, there is no explicit study of fluctuation-induced diamagnetism below the critical temperature in the weak-field region in HTSC. We have estimated the fluctuation-induced diamagnetism $\Delta\chi$ below the critical temperature as a function of the reduced temperature in the high-T_c system.

II. THEORETICAL FORMULATIONS

A. Magnetic moment of high-T_c superconductors below T_c

The Lawrence-Doniach functional describes the fluctuation phenomena in layered high-T_c superconductors with the weak coupling between the CuO_2 planes. If a magnetic field is applied, the free-energy functional takes the form as[6]

$$F[\psi]=\sum_n \int d^2r\left[a|\psi_n|^2+\frac{b}{2}|\psi_n|^4+\frac{\hbar^2}{4m_\parallel}\left|\left(\nabla_\parallel-\frac{2ie}{\hbar c}A_\parallel\right)\right|^2 + t|\psi_{n+1}-\psi_n|^2\right], \quad (3)$$

where ψ_n is the order parameter of the nth superconducting layer and $a=\alpha(T-T_c)/T_c=\alpha\epsilon$.

As the magnetic field is applied to a superconductor, the order parameter will be perturbed. The thermal energy causes another perturbation in ψ_n. So the free-energy functional will be perturbed by ΔF_{fl}, which is termed as the fluctuation contribution to the free energy. This is written as[7]

$$\Delta F_{fl}(h,T)=-\frac{Vk_BTh}{4\pi^2\xi_\parallel^2(0)d}\int_{-\pi}^{+\pi}dz \times\sum_{n=0}^{\infty}\ln\frac{\pi k_BT}{\alpha[(2n+1)h+r(1-\cos z)+\epsilon]}, \quad (4)$$

with $h=B/H'_{c2}$ and the magnetic field is applied along the z direction or the c axis. The parameter r is given by[6]

$$r=2\frac{\xi_c^2(0)}{d^2},$$

where $\xi_c(0)$ is the coherence length along the c direction.

Baraduc *et al.* have calculated the magnetic moment using the fluctuation in the free energy below and above the critical temperature as a function of the magnetic field and reduced temperature.[6]

The magnetic moment below the critical temperature is given as follows:[6]

$$M(-|\epsilon|)=\Delta M+2M(\epsilon=0)-M(|\epsilon|). \quad (5)$$

B. Calculation of ΔM and its variation with the reduced temperature

ΔM is expressed as

$$\Delta M=-\frac{Vk_BT}{2\pi\xi_\parallel^2(0)dH'_{c2}}\frac{\partial}{\partial h}h\sum_{n=0}^{\infty}\ln\frac{\phi(P)^2}{\phi(P+|\theta|)\phi(P-|\theta|)}, \quad (6)$$

where

$$\phi(N)=N+\sqrt{N^2-\rho^2}, \quad (7)$$

with $N=n+\epsilon/2h+\rho$ and $\rho=r/2h$. The n has the cutoff with a value $n_0\simeq 1/h$, because the Ginzburg-Landau approximation is valid only for this choice.[4]

We have derived ΔM presented in Appendix A that can be expressed as

$$\Delta M=-\frac{Vk_BT}{\phi_0 d}\left[\sum_{n=0}^{n_0}\ln\frac{\phi(P)^2}{\phi(P+|\theta|)\phi(P-|\theta|)}+hS_T\right], \quad (8)$$

with

$$S_T=\sum_{n=0}^{n_0}T_1+T_2+T_3, \quad (9)$$

where T_1, T_2, and T_3 can be expressed as follows:

$$T_1=-\frac{r}{h^2[P+(P^2-\rho^2)^{1/2}]}\left[1+P(P^2-\rho^2)^{-1/2}\right], \quad (10)$$

$$T_2=-\frac{1+\{(P+|\theta|)^2-\rho^2\}^{-1/2}(P+|\theta|)}{(P+|\theta|)+\{(P+|\theta|)^2-\rho^2\}^{1/2}}\left(\frac{r+\epsilon}{2h^2}\right), \quad (11)$$

$$T_3=-\frac{1+\{(P-|\theta|)^2-\rho^2\}^{-1/2}(P-|\theta|)}{(P-|\theta|)+\{(P-|\theta|)^2-\rho^2\}^{1/2}}\left(\frac{r-\epsilon}{2h^2}\right). \quad (12)$$

In expressions for T_1, T_2, and T_3, we get

$$P+|\theta|=n+\frac{1}{2}+\frac{r+\epsilon}{2h} \quad (13)$$

and

$$P-|\theta|=n+\frac{1}{2}+\frac{r-\epsilon}{2h}. \quad (14)$$

In the limit $r\gg\epsilon$, P becomes equal to $P+|\theta|$ and $P-|\theta|$, which can be expressed as follows:

$$P=P\pm|\theta|=n+\frac{1}{2}+\frac{r}{2h}. \quad (15)$$

Therefore, in the above limit, ΔM becomes zero and is independent of ϵ otherwise so that

$$\frac{d\Delta M}{d\epsilon}=0. \quad (16)$$

C. Calculation of $M(|\epsilon|)$

The third part of the magnetic moment deduced by Koshelev below the critical temperature $M(|\epsilon|)$ is written as[4]

$$M(B,\epsilon)=-\frac{T}{d\phi_0}m_{sc}(b), \quad (17)$$

where

$$m_{sc}(b)=\frac{d}{db}\left[b\sum_{n=0}^{n_0}\int_{-1/2}^{+1/2}dx\,\ln\frac{1+b[2(n+x)+1]}{1+b(2n+1)}\right], \quad (18)$$

with $b=h/\epsilon$ and $\phi_0=2\pi\xi_\parallel^2(0)H'_{c2}$. After the integration followed by the differentiation in Eq. (18), given in Appendix B, we get $m_{sc}(b)$ as given below,

$$m_{sc}(b)=\sum_{n=0}^{n_0}\left[(n+1)\ln(1+2nb+2b)-n\ \ln(1+2nb)\right.$$
$$\left.-\ln(1+2nb+b)-b\frac{2n+1}{1+(2n+1)b}\right]. \qquad (19)$$

So, the term for magnetization $M(|\epsilon|)$ is obtained in the form

$$M(|\epsilon|)=-\frac{T}{d\phi_0}\sum_{n=0}^{n_0}\left[(n+1)\ln(1+2nb+2b)\right.$$
$$-n\ \ln(1+2nb)-\ln(1+2nb+b)$$
$$\left.-b\frac{2n+1}{1+(2n+1)b}\right]. \qquad (20)$$

D. Calculation of $M(\epsilon=0)$

Now to calculate $M(\epsilon=0)$, we reform $m_{sc}(b)$ as (see Appendix C) follows:

$$m_{sc}(\epsilon)=\sum_{n=0}^{n_0}\left[(n+1)\{\epsilon+2(n+1)h\}-n\ \ln(\epsilon+2nh)\right.$$
$$\left.-\ln\{\epsilon+(2n+1)h\}-\frac{(2n+1)h}{\epsilon+(2n+1)h}\right]. \qquad (21)$$

So, from Eq. (21), we get

$$m_{sc}(\epsilon=0)=\sum_{n=0}^{n_0}[(n+1)\ln(2n+1)h-n\ \ln 2nh$$
$$-\ln(2n+1)h-1]. \qquad (22)$$

On simplification, Eq. (22) is transformed to

$$m_{sc}(\epsilon=0)=\sum_{n=0}^{n_0}\left[n\ \ln\frac{n+1}{n}+\ln\frac{n+1}{n+\frac{1}{2}}-1\right]. \qquad (23)$$

This is same as that for $m_{sc}(b\gg 1)$ and is equal to -0.346. At $\epsilon=0$, $T=T_c$, so that $M(\epsilon=0)$ can be written as follows:

$$M(\epsilon=0)=\frac{0.346T_c}{d\phi_0}. \qquad (24)$$

Therefore, we have

$$\frac{dM(\epsilon=0)}{d\epsilon}=0. \qquad (25)$$

E. The diamagnetic susceptibility χ and $d\chi/d\epsilon$

We know that the susceptibility χ is obtained by the relation

$$\chi=\frac{M(-|\epsilon|)}{B}. \qquad (26)$$

So, the variation of χ with the reduced temperature is given by the relation

$$\frac{d\chi}{d\epsilon}=\frac{1}{B}\frac{dM(-|\epsilon|)}{d\epsilon}. \qquad (27)$$

It follows from Eq. (3),

$$\frac{dM(-|\epsilon|)}{d\epsilon}=\frac{d\Delta M}{d\epsilon}+2\frac{dM(\epsilon=0)}{d\epsilon}-\frac{dM(|\epsilon|)}{d\epsilon}. \qquad (28)$$

Using Eqs. (16), (25), and (28) we get

$$\frac{dM(-|\epsilon|)}{d\epsilon}=-\frac{dM(|\epsilon|)}{d\epsilon}. \qquad (29)$$

We know that $b=B/H'_{c2}\epsilon$. Therefore it is found that

$$d\epsilon=-\frac{B}{H'_{c2}b^2}db. \qquad (30)$$

Therefore, the variation of χ with the reduced temperature is found from Eqs. (27), (29), and (30) as follows:

$$\frac{d\chi}{d\epsilon}=-\frac{1}{H'_{c2}\epsilon^2}\frac{dM(|\epsilon|)}{db}. \qquad (31)$$

F. Calculation of $dM(|\epsilon|)/db$

Differentiating Eq. (20), as shown in Appendix D, we get

$$\frac{dM(|\epsilon|)}{db}=-\frac{T_c(1+\epsilon)}{d\phi_0}\sum_{n=0}^{n_0}\left[\frac{2(n+1)^2}{1+2nb+2b}-\frac{4n}{1+2nb}\right.$$
$$\left.+\frac{(2n+1)(2+2nb+b)}{(1+2nb+b)^2}\right]. \qquad (32)$$

G. Calculation of diamagnetic fluctuation below the critical temperature T_c

If the applied magnetic field is very small or zero, then using $b\simeq 0$ we have

$$\frac{dM(|\epsilon|)}{db}=-\frac{2T_c(1+\epsilon)}{d\phi_0}\sum_{n=0}^{n_0}(n^2+2n+2). \qquad (33)$$

Using Eqs. (31) and (33) we get

$$\frac{d\chi}{d\epsilon}=\frac{4\pi T_c\xi_{\parallel}^2(0)}{d\phi_0^2}\frac{1+\epsilon}{\epsilon^2}\Omega(n), \qquad (34)$$

with

$$\Omega(n)=\sum_{n=0}^{n_0}(n^2+2n+2). \qquad (35)$$

Integrating Eq. (34) we have the susceptibility χ below the critical temperature T_c represented by

$$\chi=\frac{4\pi T_c\xi_{\parallel}^2(0)}{d\phi_0^2}\Omega(n)\int\left(\frac{1}{\epsilon^2}+\frac{1}{\epsilon}\right)d\epsilon. \qquad (36)$$

The above expression for χ can be rewritten as

$$\chi=\eta\left[\ln \epsilon-\frac{1}{\epsilon}\right]+C, \tag{37}$$

with

$$\eta=\frac{4\pi T_c\xi_{\parallel}^2(0)}{d\phi_0^2}\Omega(n). \tag{38}$$

The diamagnetic susceptibility saturates with a value χ_0 at some lower temperature $T<T_c$, where $|\epsilon|\simeq 1$. With this boundary condition we can write the diamagnetic susceptibility as follows.

$$\chi=\eta\left[\ln \epsilon-\frac{1}{\epsilon}\right]+\chi_0+\eta. \tag{39}$$

At the critical temperature the superconductors should show perfect diamagnetism with the susceptibility χ_0. But the perfect diamagnetism is found at a temperature lower than T_c. The difference $\chi_0-\chi$ is the fluctuation in diamagnetism $(\Delta\chi)$ caused by the thermal fluctuations. The temperature dependence of $\Delta\chi$ is obtained from the above equation as expressed below,

$$\Delta\chi=\chi_0-\chi=-\eta\left[\ln \epsilon-\frac{1}{\epsilon}+1\right]. \tag{40}$$

Now we have $\ln \epsilon\simeq\epsilon-1$ because it is known that $2\geqslant\epsilon\geqslant 0$. Using this we can write $\Delta\chi$ as follows:

$$\Delta\chi=\chi_0-\chi=-\eta\left[\epsilon-\frac{1}{\epsilon}\right]. \tag{41}$$

Since $\epsilon^2\ll 1$ for the discussed temperature region, the fluctuation-induced diamagnetism is given by the following equation:

$$\Delta\chi=-\frac{\eta}{\epsilon}=-\frac{4\pi T_c\xi_{\parallel}^2(0)}{d\phi_0^2}\Omega(n)\epsilon^{-1}. \tag{42}$$

III. SUMMARY

The fluctuation-induced diamagnetism of HTSC is estimated as a function of the reduced temperature below the critical temperature. Below the critical temperature, $\Delta\chi$ varies inversely proportional to the reduced temperature ϵ. Also, the critical temperature and the number of superconducting layers influence the absolute magnitude of $\Delta\chi$. There may be a dimensional crossover below the critical temperature obtained between two regimes with different ϵ as is reported in the case of thermal conductivity.[16] Depending upon the nature of the pair-breaking (weak or strong) Maki-Thompson contribution below T_c may be an important study. Though the contribution from the density of states is larger above the crossover temperature above T_c, the contribution to $\Delta\chi$ may be affected to a certain extent in weak magnetic field below the critical temperature. These two contributions are controlled by the applied magnetic field. In addition, the ab-plane coherence length is an important factor to alter the fluctuation effect in diamagnetism in HTSC. This is unlike the fluctuation above the critical temperature.

ACKNOWLEDGMENT

One of the authors (A.K.G.) acknowledges the University Grants Commission, Government of India, for financial support.

APPENDIX A

We have the expression for ΔM as follows:

$$\Delta M=\text{const}\times\frac{\partial}{\partial h}h\sum_{n=0}^{n_0}\frac{\phi(P)^2}{\phi(P+|\theta|)\phi(P-|\theta|)}. \tag{A1}$$

We get from the above equation,

$$\begin{aligned}\frac{\partial}{\partial h}\sum_{n=0}^{n_0}&\frac{\phi(P)^2}{\phi(P+|\theta|)\phi(P-|\theta|)}\\
&=\frac{\partial}{\partial h}\sum_{n=0}^{n_0}[2\ln \phi(P)-\ln \phi(P+|\theta|)-\ln \phi(P-|\theta|)]\\
&=\sum_{n=0}^{n_0}\frac{2\phi'(P)}{\phi(P)}\frac{\partial P}{\partial h}-\frac{\phi'(P+|\theta|)}{\phi(P+|\theta|)}\frac{\partial(P+|\theta|)}{\partial h}\\
&\quad-\frac{\phi'(P-|\theta|)}{\phi(P-|\theta|)}\frac{\partial(P-|\theta|)}{\partial h}.\end{aligned} \tag{A2}$$

Now, T_1 can be expressed as

$$T_1=\frac{2\phi'(P)}{\phi(P)}\frac{\partial P}{\partial h}=\frac{2\phi'(P)}{\phi(P)}\frac{\partial}{\partial h}[n+\rho+\tfrac{1}{2}],$$

so that

$$T_1=\frac{2}{P+\sqrt{P^2-\rho^2}}\phi'(P)\frac{\partial}{\partial h}\left(\frac{r}{2h}\right).$$

Again putting $\phi'(P)$ in the above expression we can rewrite T_1 as follows:

$$T_1=\frac{r}{P+\sqrt{P^2-\rho^2}}\left(-\frac{1}{h^2}\right)[1+P(P^2-\rho^2)^{-1/2}]. \tag{A3}$$

The second term T_2 can be evaluated as follows:

$$T_2=\frac{\phi'(P+|\theta|)}{\phi(P+|\theta|)}\frac{\partial}{\partial h}(P+|\theta|). \tag{A4}$$

Finally, after the differentiation we have

$$T_2=\frac{1+\{(P+|\theta|)^2-\rho^2\}^{-1/2}(P+|\theta|)}{P+|\theta|+\sqrt{(P+|\theta|)^2-\rho^2}}\left(-\frac{r+\epsilon}{2h^2}\right). \tag{A5}$$

Similarly, the third term T_3 can be expressed as follows:

$$T_3=\frac{1+\{(P-|\theta|)^2-\rho^2\}^{-1/2}(P-|\theta|)}{P-|\theta|+\sqrt{(P+|\theta|)^2-\rho^2}}\left(\frac{-r+\epsilon}{2h^2}\right). \tag{A6}$$

APPENDIX B

$$m_{sc}=\frac{d}{db}\left[b\sum_{n=0}^{n_0} I_1\right], \tag{B1}$$

with

$$I_1=\int_{-1/2}^{+1/2}\ln\frac{1+b[2(n+x)+1]}{1+b(2n+1)}dx. \tag{B2}$$

Let us assume $1+b[2(n+x)+1]=y$ so that I_1 can be written as

$$I_1=\frac{1}{2b}\int_{1+2nb}^{1+b(2n+2)}\ln\frac{y}{1+b(2n+1)}dy. \tag{B3}$$

On integration of the above we have I_1 as given below,

$$I_1=\frac{1}{2b}[(1+2nb+2b)\ln(1+2nb+2b)-2b -(1+2nb)\ln(1+2nb)-2b\ \ln(1+2nb+b)]. \tag{B4}$$

Now we can rewrite $m_{sc}(b)$ as follows:

$$m_{sc}(b)=\frac{d}{db}\left[b\sum_{n=0}^{n_0}\frac{1}{2b}\{(1+2nb+2b)\times\ln(1+2nb+2b)-2b-(1+2nb)\times\ln(1+2nb)-2b\ \ln(1+2nb+b)\}\right]. \tag{B5}$$

After the differentiation we have

$$m_{sc}(b)=\sum_{n=0}^{n_0}\left[(n+1)\ln(1+2nb+2b)-n\ \ln(1+2nb)-\ln(1+2nb+b)-b\frac{2n+1}{1+(2n+1)b}\right]. \tag{B6}$$

APPENDIX C

The b is defined with the help of the reduced temperature as follows:

$$b=\frac{h}{\epsilon}. \tag{C1}$$

Putting b in the expression for $m_{sc}(b)$ we get after a simplification,

$$m_{sc}(b)=\sum_{n=0}^{n_0}\left[(n+1)\ln\{\epsilon+2(n+1)h\}-n\ \ln\{\epsilon+2nh\}-\ln\{\epsilon+(2n+1)h\}-\frac{(2n+1)h}{\epsilon+(2n+1)h}\right]. \tag{C2}$$

Now putting $\epsilon=0$ in the above expression we rewrite $m_{sc}(\epsilon=0)$ as follows:

$$m_{sc}(\epsilon=0)=\sum_{n=0}^{n_0}[(n+1)\ln\{2h(n+1)\}-n\ \ln 2nh-\ln\{(2n+1)h\}-1]. \tag{C3}$$

After rearrangement of the above equation we get

$$m_{sc}(\epsilon=0)=\sum_{n=0}^{n_0}\left[n\ \ln\frac{n+1}{n}+\ln\frac{n+1}{n+\frac{1}{2}}-1\right]. \tag{C4}$$

APPENDIX D

Differentiating both sides of Eq. (18) with respect to b we obtain

$$\frac{dM(|\epsilon|)}{db}=-\frac{T}{d\phi_0}\sum_{n=0}^{n_0}\left[\frac{2(n+1)^2}{1+2nb+2b}-2\frac{2n}{1+2nb}-\frac{2n+1}{1+2nb+b}-\frac{2n+1}{1+(2n+1)b}+b\frac{(2n+1)^2}{(1+2nb+b)^2}\right]. \tag{D1}$$

After simplification we get

$$\frac{dM(|\epsilon|)}{db}=-\frac{T}{d\phi_0}\sum_{n=0}^{n_0}\left[\frac{2(n+1)^2}{1+2nb+2b}-\frac{4n}{1+2nb}+\frac{(2n+1)(2+2nb+b)}{(1+2nb+b)^2}\right]. \tag{D2}$$

Putting $T=T_c(1+\epsilon)$ in the above equation we get

$$\frac{dM(|\epsilon|)}{db}=-\frac{T_c(1+\epsilon)}{d\phi_0}\sum_{n=0}^{n_0}\left[\frac{2(n+1)^2}{1+2nb+2b}-\frac{4n}{1+2nb}+\frac{(2n+1)(2+2nb+b)}{(1+2nb+b)^2}\right]. \tag{D3}$$

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