PHYSICAL REVIEW B VOLUME 60, NUMBER 14 1 OCTOBER 1999-II

Second magnetization peak in flux lattices: The decoupling scenario

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(Received 10 March 1999)

The second peak phenomenon of flux lattices in layered superconductors is described in terms of a disorder induced layer decoupling transition. For weak disorder the tilt modulus undergoes an apparent discontinuity which leads to an enhanced critical current and reduced domain size in the decoupled phase. The Josephson plasma frequency is reduced by decoupling and by Josephson glass pinning; in the liquid phase it varies as $1/[BT(T+T_0)]$, where T is temperature, B is field, and T_0 is the disorder dependent temperature of the multicritical point. [S0163-1829(99)51738-4]

Vortex matter in the presence of disorder has emerged as a fundamental problem of elastic manifolds in a random media.[1] Impurity disorder does not allow long-range translational order of the flux lattice and finite domains are expected.[2] At low temperatures and fields the system is a Bragg glass,[3,4] i.e., the lattice is dislocation free, and at long scales the displacement correlations decay as a power law and Bragg peaks are expected. The impurity-induced domains are essential for the description of both equilibrium, e.g., thermodynamic phase transitions and nonequilibrium, e.g., critical current phenomena.

The critical current j_c measures the pinning force in the domains.[1,2] Increasing the magnetic field or temperature reduces the pinning force and j_c is decreased. However, in many type II superconductors, a sharp enhancement of j_c is observed at a ''second peak'' field B_0. This peak phenomena is most pronounced in layered superconductors such as $Bi_2Sr_2CaCu_2O_8$ (BSSCO),[5–7] $YBa_2Cu_3O_7$ (YBCO),[8] in $NbSe_2$[9,10] and in Pb/Ge multilayers[11] for fields perpendicular to the layers. The second peak phenomena signals that pinning becomes more effective, e.g., due to softening of the flux lattice.[2] The reason for softening could be the approach to melting;[12] however, neutron-scattering data on BSCCO[13] show that Bragg spots of the flux lattice persist well above B_0.

Disorder plays an essential role also in the equilibrium phase diagram of layered superconductors. This has been most extensively studied in BSSCO.[5–7,14,15] The second peak corresponds to a phase transition[7] in the range 500–900 G (decreasing with disorder) and is weakly temperature dependent up to a temperature $T_0 \approx 40$ K. The point B_0, T_0 is a multicritical point where the second peak transition meets a first order transition as well as two depinning lines. Thus the second peak manifests both equilibrium and nonequilibrium phenomena of disorder in flux lattices and its understanding presents a fundamental challenge.

For a flux lattice with point impurities, by using renormalization group (RG) and replica symmetry breaking (RSB) methods we have derived[16] a phase diagram with four phases, which all meet at a multicritical point B_0, T_0, in remarkable correspondence with data on BSCCO. The present work focuses on the layer decoupling transition at a temperature independent field B_0 for $T<T_0$. As shown here, the fusion of Bragg glass concepts with decoupling accounts for the peculiar second peak phenomenon, i.e., the enhanced j_c. The Josephson plasma resonance is also considered as a probe of the Josephson coupling,[17,18] being reduced by decoupling and by a Josephson glass parameter. Very recent data on BSCCO have indeed shown a significant reduction in the resonance frequency at the second peak transition.[19,20]

It has been recently shown that decoupling coalesces with a defect unbinding transition[21] which has analogs in isotropic systems.[22] The resulting vacancies and interstitials lead to a reduction in the elastic tilt mudulus,[23] consistent with the decoupling scenario as described below. It is possible then that a decoupling-defect transition accounts for the peak phenomenon in all type II superconductors. The analysis below is, however, presented for layered anisotropic systems where quantitative predictions can be made.

In a layered superconductor each flux line is composed of one point singularity, or a pancake vortex, in each layer. When the pancake vortices fluctuate they can generate a divergence in the Josephson phase, leading to a renormalized Josephson coupling E_J^R which vanishes in the decoupled $B>B_0$ phase.[16,24,25] The three-dimensional flux lattice is still present in the decoupled phase (in the Bragg glass sense), with interlayer coupling mediated by the magnetic field. Before presenting a microscopic model, I started with a rather simple description of elasticity within domains, which shows the second peak transition, i.e., j_c enhancement at decoupling.

The transverse tilt modulus of a flux lattice in a layered superconductor for fields perpendicular to the layers is given by[26–28]

$$c_{44}(\mathbf{q},k)=\frac{\tau}{32\pi\lambda_{ab}^2 d}+\frac{B^2}{4\pi}\frac{1}{1+\lambda_c^2q^2+\lambda_{ab}^2k^2}+\frac{2B\phi_0}{(8\pi\lambda_c)^2}\ln(a^2/4\pi\xi_0^2), \quad (1)$$

where $\mathbf{q}$ and k are momenta parallel and perpendicular to the layers, respectively, λ_{ab} and λ_c are the London penetration lengths parallel and perpendicular to the layers, respectively, ϕ_0 is the flux quantum, $a^2=\phi_0/B$ is the unit-cell area, d is the interlayer spacing, ξ_0^2 is the in-layer coherence length, and $\tau=\phi_0^2d/(4\pi^2\lambda_{ab}^2)$ sets the energy scale. The first term of Eq. (1) is due to the magnetic coupling, while the second and third terms originate from the Josephson coupling energy per unit area E_J, i.e., $\lambda_c^2=\tau\lambda_{ab}^2/(4\pi E_J d^2)$. The second term is peculiar: at $q\neq 0$ it vanishes when E_J vanishes and $\lambda_c\to\infty$, as it should. However, at $q=0$ this term seems to survive even if $\lambda_c\to\infty$. The origin of this peculiarity is that the harmonic expansion of the Josephson cosine term which identifies c_{44} fails[28] when both q, $1/\lambda_c\to 0$. In fact, the nonlinear cosine term generates a renormalized λ_c^R which diverges at decoupling.

The Bragg glass domain size R_{BG} (parallel to the layers) sets a scale for the relevant q values. When $R_{BG}>\lambda_c^R$ the tilt mudulus is large, containing the $B^2/4\pi$ term of Eq. (1). However, as decoupling at the field B_0 is approached λ_c^R diverges, and when $R_{BG}<\lambda_c^R$ Eq. (1) fails to describe c_{44} on the scale of $q\approx 1/R_{BG}$. This defines an anharmonic crossover regime where usual elasticity cannot be used to derive Bragg glass properties. Finally, at $B>B_0$ elasticity is restored and c_{44} is reduced to the first term in Eq. (1). The main interest is in the regime of strong fields, i.e., $a\lesssim 2\lambda_{ab}$ where $T_0<\tau$ is below melting.[16] Thus at $B<B_0$ and for sufficiently large domains the second term in Eq. (1) dominates and c_{44} has an apparent discontinuity,

$$c_{44}=\pi\lambda_{ab}^2\tau/da^4 \qquad \lambda_c^R<R_{BG}, \tag{2a}$$

$$c_{44}=\tau/(32\pi\lambda_{ab}^2 d) \qquad \lambda_c^R=\infty. \tag{2b}$$

Hence c_{44} is reduced within the anharmonic regime by the small factor $\epsilon=a^4/(32\pi^2\lambda_{ab}^4)$.

The apparent discontinuity in c_{44} affects also the domain sizes which can be estimated by a dimensional argument.[2,3] Consider the tilt c_{44} and shear c_{66} terms of the elasticity Hamiltonian for the displacement $\mathbf{u}(\mathbf{r})$ and its transverse component $\mathbf{u}_T(\mathbf{r})$. Rescaling parallel and perpendicular lengths yields an isotropic form[1,4]

$$\mathcal{H}=\int d^3r\Big\{\tfrac{1}{2}c_{44}^{1/3}c_{66}^{2/3}[\nabla u_T(\mathbf{r})]^2 - (\xi_0^2/a^2d)U_{pin}(\mathbf{r})\sum_{\mathbf{Q}}\cos\mathbf{Q}\cdot[\boldsymbol{\rho}-\mathbf{u}(\mathbf{r})]\Big\}, \tag{3}$$

where $U_{pin}(\mathbf{r})$ is a random potential in three-dimensional $\mathbf{r}=(\boldsymbol{\rho},z)$ which couples to the flux density modulations with wave vectors $\mathbf{Q}$; its disorder average is $\langle U_{pin}(\mathbf{r})U_{pin}(\mathbf{r}')\rangle=\frac{1}{2}d\overline{U}\delta^3(\mathbf{r}-\mathbf{r}')$. Disorder average over configurations $\mathbf{u}(\mathbf{r})$ and $\mathbf{u}'(\mathbf{r})$ yields $\Sigma_{\mathbf{Q}}\cos\mathbf{Q}\cdot[\mathbf{u}(\mathbf{r})-\mathbf{u}'(\mathbf{r})]$; the sum is cut off by $Q\lesssim\langle u_T^2\rangle^{-1/2}$, where $\langle u^2\rangle\approx\langle u_T^2\rangle$ are the fluctuations in a domain of size R'. Thus averaging Eq. (3) yields

$$\langle H\rangle/R'^3=\tfrac{1}{2}c_{44}^{1/3}c_{66}^{2/3}\langle u_T^2\rangle R'^{-2}-\overline{U}^{1/2}\xi_0^2/[a^2d\langle u_T^2\rangle R'^3]^{1/2}. \tag{4}$$

Minimizing with respect to R' yields $R'\sim\langle u_T^2\rangle^3$, i.e., the Flory exponent.[3] The domain size parallel to the layers is [up to $\ln(a/d)$ and a numerical prefactor]

$$R\approx(\lambda_{ab}/a)^5\langle u_T^2\rangle^3/(s\xi_0^4 d) \qquad \lambda_c^R<R,$$

$$R\approx(\lambda_{ab}/a)^3\langle u_T^2\rangle^3/(4\pi s\xi_0^4 d) \qquad \lambda_c^R=\infty, \tag{5}$$

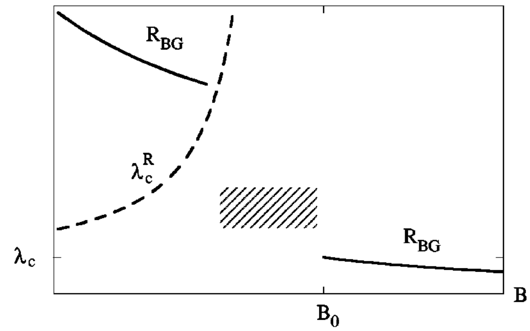

FIG. 1. Bragg glass domain size R_{BG} parallel to the layers and the renormalized London length perpendicular to the layers λ_c^R; the latter diverges at the decoupling field B_0. The R_{BG} can be found from elasticity for $B<B_0$ only if $R_{BG}>\lambda_c^R$; otherwise, as in the hatched region, the elastic tilt modulus is ill defined.

where $c_{66}=\tau/16da^2$,[26–28] $s=4\pi\overline{pU}\lambda_{ab}^4/[\tau^2a^2s\ln^2(a/d)]$ defines the decoupling transition at $s=\frac{1}{2}$, and $\overline{p}\approx 1$ is defined below.

The pinning length $R=R_p$ is given by Eq. (5) with $\langle u_T^2\rangle\approx\xi_0^2$. To allow for large pinning domains one needs either $a<\lambda_{ab}$ or to allow for domains with a somewhat larger fluctuations in $\langle u_T^2\rangle$; the latter increases R_p very rapidly since it increases with the 6th power of u_T. The critical current can now be estimated[1,2] by balancing the Lorenz force j_cBR^3/c with the pinning force $\langle H\rangle/\xi_0$ [evaluated at the minimum of Eq. (4)], leading to $j_c\sim 1/c_{44}$. Increasing the field within the anharmonic regime decreases c_{44} by the factor ϵ so that j_c is significantly enhanced when $a\lesssim\lambda_{ab}$. Note that the measured magnetization changes (and inferred j_c) at B_0 decrease with temperature due to the strongly temperature-dependent relaxation rates,[29] approaching the much smaller equilibrium magnetizations.

A second length scale $R=R_{BG}$ is identified by Eq. (5) with the fluctuations $\langle u_T^2\rangle\approx a^2$. The proper definition of R_{BG} is the scale for the onset of the $\ln r$ form for the displacement correlation function. While the derivation from Eq. (4) cannot capture this $\ln r$, it does give the right expression for R_{BG}.[3] Thus, R_{BG} depends on c_{44} and is reduced by $\epsilon^{1/2}$ through the anharmonic regime. The latter depends also on λ_c^R for which $\ln\lambda_c^R\sim(B-B_0)^{-1}$ in the RSB or first-order RG solutions,[16] though $\ln\lambda_c^R\sim(B-B_0)^{-1/2}$ in second-order RG;[30] decoupling may also be of first order,[25] leading to a narrower anharmonic regime. Figure 1 illustrates the lengths R_{BG} and λ_c^R, demonstrating the anharmonic regime within which R_{BG} has a significant drop and correspondingly j_c has an apparent jump. Note that even in the decoupled phase ($B>B_0$) R_{BG} is large for typical type II superconductors, $R_{BG}\approx\lambda_{ab}^3a^3/(4\pi s\xi_0^4 d)\gg a$, consistent with a decoupling transition within the Bragg glass phase, i.e., below a melting transition.

I proceed now to derive the lattice displacement correlation allowing for a renormalized Josephson coupling and for a Josephson glass order parameter. This derivation avoids the harmonic expansion for the elastic modulii and shows how the Bragg glass domain sizes are directly affected by the renormalized λ_c^R. The Josephson phase between the layers n and $n+1$ at position $\mathbf{r}$ in the layer involves contributions from a nonsingular component $\theta_n(\mathbf{r})$ and from singular vortex terms.[30] Consider a flux lattice with an equilibrium position of the lth flux line at $\mathbf{R}_l$. The singular phase around a pancake vortex at position $\mathbf{R}_l+\mathbf{u}_l^n$ is $\alpha(\mathbf{r}-\mathbf{R}_l-\mathbf{u}_l^n)$, where $\alpha(\mathbf{r})=\arctan(y/x)$ with $\mathbf{r}=(x,y)$. Expansion of the interlayer phase difference $\alpha(\mathbf{r}-\mathbf{R}_l-\mathbf{u}_l^n)-\alpha(\mathbf{r}-\mathbf{R}_l-\mathbf{u}_l^{n+1})$ yields for the singular part of the Josephson phase $b_n(\mathbf{r})=\Sigma_l(\mathbf{u}_l^{n+1}-\mathbf{u}_l^n)\nabla\alpha(\mathbf{r}-\mathbf{R}_l)$. The Hamiltonian for the transverse displacements involves also the magnetic contributions to the shear modulus $c_{66}=\tau/(16da^2)$ and the tilt mudulus

$$c_{44}^0(k)=[\tau/(8da^2\lambda_{ab}^2k_z^2)]\ln(1+a^2k_z^2/4\pi),$$

where $k_z=(2/d)\sin(kd/2)$; its $k\to0$ form is the first term in Eq. (1). This leads to the Hamiltonian of the pure system

$$\mathcal{H}_{pure}/T=\tfrac{1}{2}\sum_{q,k} G_f^{-1}(\mathbf{q},k)|\theta(\mathbf{q},k)|^2 +\tfrac{1}{2}\sum_{\mathbf{q},k} c(\mathbf{q},k)q^2|b(\mathbf{q},k)|^2 -\frac{E_J}{T}\sum_n\int d^2r\cos[\theta_n(\mathbf{r})+b_n(\mathbf{r})]. \qquad (6)$$

Here $c(\mathbf{q},k)=(a^2/2\pi d)^2[k_z^2c_{44}^0(k)+q^2c_{66}]/Tk_z^2$, E_J is the Josephson coupling, and the coefficient of the nonsingular phase is[30] $G_f(q,k)=4\pi d^3T(\lambda_{ab}^{-2}+k_z^2)/(\tau q^2)$. The conventional c_{44} is obtained by expanding the cosine term in Eq. (6) and shifting $\theta(\mathbf{q},k)$ to eliminate the cross term. The latter shift leads to an expansion parameter[28] with terms $\sim q^2k_z^2|u_T(\mathbf{q},k)|^2/[q^2+\lambda_c^{-2}(1+\lambda_{ab}^2k_z^2)]^2$, i.e., these diverge when both $q,1/\lambda_c\to0$ and the expansion becomes invalid.

Consider now a pinning potential $U_{pin}^n(\mathbf{r})$ which couples to the vortex shape function $p(\mathbf{r})$ leading to a pinning energy $\int d^2r\Sigma_{n,l}U_{pin}^n(\mathbf{r})p(\mathbf{r}-\mathbf{R}_l-\mathbf{u}_l^n)$. The aim is to identify domain sizes R_p (and infer R_{BG}), hence the pinning energy is expanded in $\mathbf{u}_l^n$ and a replica average with the weight $\exp\{-\int d^2r\Sigma_n[U_{pin}^n(\mathbf{r})]^2/\overline{U}\}$ then leads to $\exp[(\overline{Up}/4T^2)\Sigma_{n,l}\Sigma_{\alpha,\beta}\mathbf{u}_l^{n,\alpha}\cdot\mathbf{u}_l^{n,\beta}]$, where $\int\partial_ip(\mathbf{r})\partial_jp(\mathbf{r})d^2r=\overline{p}\,\delta_{i,j}$ and $\alpha,\beta=1,2,\ldots,n$ are replica indices.

The $b^\alpha(\mathbf{q},k)$ variables can be decoupled from the total Josephson phase $\tilde{b}_n(\mathbf{r})=b_n(\mathbf{r})+\theta_n(\mathbf{r})$ by shifting to $d^\alpha(\mathbf{q},k)=b^\alpha(\mathbf{q},k)-B_{\gamma,\alpha}(\mathbf{q},k)G_f^{-1}(\mathbf{q},k)\tilde{b}^\alpha(\mathbf{q},k)$, where

$$B_{\alpha,\beta}^{-1}(\mathbf{q},k)=G_f^{-1}(\mathbf{q},k)\alpha(\mathbf{q},k)\delta_{\alpha,\beta}-s_0q^2/k_z^2,$$

$\alpha(\mathbf{q},k)=1+G_f(\mathbf{q},k)c(\mathbf{q},k)q^2$ and $s_0=\overline{Up}a^2d/(4\pi d^2T)^2$. The resulting replicated Hamiltonian is

$$\mathcal{H}_r=\tfrac{1}{2}\sum_{q,k;\alpha,\beta}B_{\alpha,\beta}^{-1}d^\alpha(\mathbf{q},k)d^{\beta*}(\mathbf{q},k) +\tfrac{1}{2}[c(\mathbf{q},k)\alpha^{-1}(\mathbf{q},k)q^2\delta_{\alpha,\beta} -s_0\alpha^{-2}(\mathbf{q},k)q^2/k_z^2]\tilde{b}^\alpha(\mathbf{q},k)\tilde{b}^{\beta*}(\mathbf{q},k) -\frac{E_J}{T}\sum_{n;\alpha}\int d^2r\cos\tilde{b}_n^\alpha(\mathbf{r}) -\frac{E_v}{T}\sum_{n;\alpha\neq\beta}\int d^2r\cos[\tilde{b}_n^\alpha(\mathbf{r})-\tilde{b}_n^\beta(\mathbf{r})]. \qquad (7)$$

The inter-replica E_v term is generated from the Josephson coupling in second order RG. It is essential to keep it from the start since it generates a Josephson glass parameter and affects the value of the decoupling field.[16]

The $\alpha(\mathbf{q},k)$ factor, which results from the nonsingular phase, is for $\epsilon\ll1$ very close to 1 for all $\mathbf{q},k$ values except when both $k<1/\lambda_{ab}$ and $q>ka/\lambda_{ab}$. The phase transitions are dominated by $k>1/a$ modes so that our previous phase diagram is recovered (Ref. 16 and Fig. 1). In particular there is a multicritical point at a field B_0 where $s=\frac{1}{2}$ and temperature $T_0=\tau a^2\ln(a/d)/8\pi\lambda_{ab}^2$. At $B=B_0$ and $T<T_0$ there is a decoupling transition at which the renormalized Josephson coupling z (with bare value $z_{bare}=E_J/Td$) vanishes. Note that the higher B_0 of YBCO as compared to BSCCO is consistent with a shorter λ_{ab} and a somewhat weaker disorder.

The fluctuations in $u_T(\mathbf{q},k)$ in terms of the shifted variables, using the RSB solution[16] are given by

$$\langle|u_T(\mathbf{q},k)|^2\rangle=(2\pi d^2)^{-2}\frac{q^2}{k_z^4}\Bigg[s_0q^2G_f(\mathbf{q},k)\alpha^{-1}(\mathbf{q},k)\times\left(c(\mathbf{q},k)q^2+\frac{G_f^{-1}(\mathbf{q},k)z}{G_f^{-1}(\mathbf{q},k)+z}\right)^{-1} +\frac{s_0}{c(\mathbf{q},k)\alpha^2(\mathbf{q},k)}\left(\frac{c(\mathbf{q},k)}{\alpha(\mathbf{q},k)}q^2+z\right)^{-1}\Bigg]+\cdots, \qquad (8)$$

where $\cdots$ stands for terms which converge in $(\mathbf{q},k)$ integration. Note the term $G_f^{-1}(\mathbf{q},k)z/[G_f^{-1}(\mathbf{q},k)+z]$, which depends on the order of $q\to0$ and $z\to0$ limits; this limit dependence leads to the apparent discontinuity in c_{44} as discussed above. For $z\neq0$ and small q, i.e., $G_f^{-1}(\mathbf{q},k)\ll z$ the first term in Eq. (8) dominates, leading to

$$\langle|u_T(\mathbf{q},k)|^2\rangle\approx\frac{4\pi^2s_0T^2}{a^8[c_{44}k^2+c_{66}q^2]^2}\quad q<1/\lambda_c^R, \qquad (9)$$

where c_{44} is from Eq. (2a) and the condition $G_f^{-1}(\mathbf{q},k)\ll z$ is written in terms of a renormalized London length $\lambda_c^R=[\lambda_{ab}^2\tau/(4\pi Td^3z)]^{1/2}$. The correlations at distance r parallel to the layers are then

$$\langle[u_T(r)-u_T(0)]^2\rangle\approx\frac{4d^2s_0T^2}{a^4c_{44}^{1/2}c_{66}^{3/2}}r\equiv\xi_0^2\frac{r}{R_p}. \qquad (10)$$

The last equality defines the pinning length R_p where the fluctuations become of order ξ_0^2. This result for R_p (up to a numerical prefactor) is the same as the one obtained from Eq. (5) with $\langle u_T^2\rangle\approx a^2$. The Bragg glass domain size is enhanced by $R_{BG}\approx R_p(a/\xi_0)^6$, as discussed above.

In the decoupled phase with $z=0$ the second term in Eq. (8) dominates. To leading order in ϵ the result is identical to Eq. (10) except that c_{44} is replaced by its $z=0$ value Eq. (2b), i.e., the pinning and Bragg glass lengths are reduced. The main result is then that the fluctuations in $u_T(r)$ behave with an effective c_{44} which is large when $q<1/\lambda_c^R$ [Eq. (2a)], i.e., for domain sizes $R_{BG}>\lambda_c^R$, while for $z=0$ c_{44} is reduced [Eq. (2b)]. In the anharmonic region below decoupling (see Fig. 1), where $R_{BG}<\lambda_c^R$, the full form of Eq. (8) is required to interpolate between these limits; this form avoids the ill-defined harmonic expansion in this regime.

Consider next the Josephson plasma frequency, given by $\omega_{pl}^2=(c^2/\epsilon_0\lambda_c^2)\langle\cos\tilde{b}_n(r)\rangle$, where ϵ_0 is the dielectric constant.[17,18] The average in $\langle\cos\tilde{b}_n(r)\rangle$ is on both thermal fluctuations and disorder and can yield significant information on the phase diagram. As shown by Koshelev[18] the local $\langle\cos\tilde{b}_n(r)\rangle$ is finite even at high temperatures, e.g., above the decoupling transition. A high-temperature expansion yields[18] $\langle\cos\tilde{b}_n(r)\rangle=(E_J/2T)\int d^2r \exp[-A(r)]$, where $A(r)=\Sigma_{\mathbf{q},k}(1-\cos\mathbf{q}\cdot\mathbf{r})\langle|\tilde{b}^\alpha(\mathbf{q},k)|^2\rangle$. The solution with disorder[16] yields (up to a $\ln B$ dependence) $A(r)=B(T+T_0)q_u^2r^2/(2B_0T_0)$ for $r<1/q_u$, where $q_u=2\ln^{1/2}(a/d)/\lambda_{ab}$ while $A(r)\sim\ln q_u r$ or $\sim r$ for larger r. The r integration is dominated by the short r correlation which yields

$$\langle\cos\tilde{b}_n(r)\rangle\approx\frac{\pi E_J\lambda_{ab}^2}{2\ln(a/d)}\times\frac{B_0T_0}{BT(T+T_0)}. \qquad (11)$$

A $1/BT$ dependence has been obtained by Koshelev[18] with a weakly temperature-dependent prefactor for an *XY* model, i.e., infinite λ_{ab}. Data on BSCCO[17] have shown that $\langle\cos\tilde{b}_n(r)\rangle\sim B^{-0.8}T^{-1}$ is in reasonable agreement with the $1/BT$ form. The present result shows that in fact the $1/BT$ form is valid in the disorder dominated regime, i.e., $T<T_0$, though in general the fluctuation term yields $\omega_{pl}^2\sim1/[BT(T+T_0)]$.

Using the RSB solution, it can be shown that the Josephson glass parameter contributes a negative term to $\langle\cos\tilde{b}_n(r)\rangle$ so that ω_{pl} is reduced, while the Josephson coupling contributes a positive $\approx z/z_{bare}$ term which vanishes at decoupling. These are mean-field results to which fluctuation terms, as Eq. (11), should be added. The recent data on BSCCO[19,20] are consistent with these results, i.e., a drop at the second peak transition followed by a field dependent fluctuation term at higher fields.

In conclusion, it is shown that a decoupling transition leads to an apparent reduction in c_{44} within an anharmonic region where the harmonic expansion fails. The proper interpolation across the anharmonic region is achieved by Eq. (8). The reduction in c_{44}, the resulting reduction in domain sizes, and the enhanced j_c account for the hallmark feature of the second peak transition. Furthermore, B_0 being weakly T dependent and decreasing with disorder,[6] as well as the Josephson plasma resonance data,[19,20] lend substantial support for the identification of the second peak transition as a disorder induced decoupling.

I thank P. Le Doussal, T. Giamarchi, T. Natterman, G. Blatter, and V. B. Geshkenbein for most valuable discussions. This research was supported by The Israel Science Foundation founded by the Israel Academy of Sciences and Humanities.

[1] For a review, see G. Blatter *et al.*, Rev. Mod. Phys. **66**, 1125 (1995).
[2] A. I. Larkin and Y. N. Ovchinnikov, J. Low Temp. Phys. **34**, 409 (1979).
[3] T. Giamarchi and P. Le Doussal, Phys. Rev. B **52**, 1242 (1995).
[4] T. Nattermann, Phys. Rev. Lett. **64**, 2454 (1990); J. Kierfeld, T. Nattermann and T. Hwa, Phys. Rev. B **55**, 626 (1997).
[5] For a review, see P. H. Kes, J. Phys. I **6**, 2327 (1996).
[6] B. Khaykovich *et al.*, Phys. Rev. B **56**, R517 (1997); Czech. J. Phys. **46-S6**, 3218 (1996).
[7] B. Khaykovich *et al.*, Phys. Rev. Lett. **76**, 2555 (1996).
[8] K. Deligiannis *et al.*, Phys. Rev. Lett. **79**, 2121 (1997).
[9] M. J. Higgins and S. Bhattacharya, Physica C **232**, 232 (1996).
[10] A. Marley *et al.*, Phys. Rev. Lett. **74**, 3029 (1995).
[11] Y. Bruynseraede *et al.*, Phys. Scr. **T42**, 37 (1992).
[12] A. van Otterlo, R. T. Scalettar, and G. T. Zimanyi, Phys. Rev. Lett. **81** 1497 (1998).
[13] E. M. Forgan *et al.*, Czech. J. Phys. **46**, 1571 (1996); S. L. Lloyd *et al.* (unpublished).
[14] D. T. Fuchs *et al.*, Nature (London) **391**, 373 (1998).
[15] D. T. Fuchs *et al.*, Phys. Rev. Lett. **80**, 4971 (1998).
[16] B. Horovitz and T. R. Goldin, Phys. Rev. Lett. **80**, 1734 (1998).
[17] Y. Matsuda *et al.*, Phys. Rev. Lett. **78**, 1972 (1997).
[18] A. E. Koshelev, Phys. Rev. Lett. **77**, 3901 (1996).
[19] T. Shibauchi *et al.*, Phys. Rev. Lett. **83**, 1010 (1999).
[20] Y. Matsuda (private communication).
[21] M. J. W. Dodgson, V. B. Geshkenbein, and G. Blatter, cond-mat/9902244 (unpublished).
[22] E. Frey, D. R. Nelson, and D. S. Fisher, Phys. Rev. B **49**, 9723 (1994).
[23] M. C. Marchetti and L. Radzihovsky, Phys. Rev. B **59**, 12 001 (1999).
[24] L. I. Glazman and A. E. Koshelev, Physica C **173**, 180 (1991).
[25] L. L. Daemen, L. N. Bulaevskii, M. P. Maley, and J. Y. Coulter, Phys. Rev. Lett. **70**, 1167 (1993).
[26] A. Sudbø and E. H. Brandt, Phys. Rev. Lett. **66**, 1781 (1991).
[27] L. I. Glazman and A. E. Koshelev, Phys. Rev. B **43**, 2835 (1991).
[28] T. R. Goldin and B. Horovitz, Phys. Rev. B **58**, 9524 (1998).
[29] Y. Yeshurun, N. Bontemps, L. Burlachkov, and A. Kapitulnik, Phys. Rev. B **49**, 1548 (1994).
[30] B. Horovitz, Phys. Rev. B **47**, 5947 (1993).