Decoupling and Decommensuration in Layered Superconductors with Columnar Defects

A. Morozov, B. Horovitz
Department of Physics and Ilze Katz center for nanotechnology, Ben Gurion university, Beer Sheva 84105 Israel

P. Le Doussal
CNRS-Laboratoire de Physique Théorique de l’Ecole Normale Supérieure, 24 rue Lhomond,75231 Cedex 05, Paris France.

We consider layered superconductors with a flux lattice perpendicular to the layers and random columnar defects parallel to the magnetic field \(B\). We show that the decoupling transition temperature \(T_d\), at which the Josephson coupling vanishes, is enhanced by columnar defects by an amount \(\delta T_d/B^2\). Decoupling by increasing field can be followed by a reentrant recoupling transition for strong disorder. We also consider a commensurate component of the columnar density and show that its pinning potential is renormalized to zero above a critical long wavelength disorder. This decommensuration transition may account for a recently observed kink in the melting line.

The phase diagram of layered superconductors in a magnetic field \(B\) perpendicular to the layers is of considerable interest in view of recent experiments on high temperature superconductors [1–4]. Columnar defects (CD) induced by heavy ion bombardment provide an additional interesting probe [1]. In particular the irreversibility line at low temperatures is enhanced [5,6] while within the liquid phase an onset of enhanced 6 axis correlation [7–9] was observed. Recent data [10] indicates that CD produce a porous vortex matter in which ordered vortex crystallites are embedded in the ‘pores’ of a rigid matrix of vortices pinned on the CD’s. A sharp kink in the melting curve signals an abrupt change from melting enhanced by the matrix at high fields to a more weakly enhanced melting at lower fields. Theoretical studies on CD’s have shown a “localization” transition within a Bose glass phase [11,12]. Recent simulations [13] have been interpreted in terms of a Bragg-Bose glass with positional order which sets in as field increases. Also a “recoupling” crossover transition [14] was studied in the vortex liquid phase.

In the absence of CD theoretical studies have shown layer decoupling due to thermal fluctuations [15–17] or due to disorder [17,18]. At this phase transition the Josephson coupling between layers vanishes at long scales, i.e. the critical current perpendicular to the layers vanishes and superconducting correlations in the \(z\) direction (perpendicular to the layers) become short range. Decoupling involves in principle also proliferation of point defects - vacancies and interstitials (VI) [19]. The flux lattice is present even in the decoupled phase with the \(z\) axis positional correlations maintained by magnetic couplings. In the case of point disorder this phase would thus still exhibit Bragg glass type order without dislocations.

An increase in the critical current at a “second peak” transition has been interpreted as due to an apparent discontinuity in the tilt modulus at decoupling [20]. Plasma resonance data [21,22] has shown a significant jump at this transition, consistent with the decoupling scenario. Whether this transition is driven by decoupling alone rather than by a sudden dislocation proliferation [23] remains to be investigated.

In the present work we consider the effects of CD within the flux lattice phase, neglecting VI (whose role is discussed below). We find that the decoupling transition temperature \(T_d(B)\) is enhanced by CD. In particular for strong disorder the low field form \(T_d(B) \sim 1/B\) becomes \(T_d(B) \sim B\) at strong fields, hence decoupling followed by a reentrant transition into a coupled state, i.e. recoupling, is possible with increasing field. These predictions can test wether the “second peak” transition is of a decoupling type. We also allow for a finite component of the CD density which is commensurate with the flux lattice, a component which usually needs to be specifically prepared [24]. We find that long wavelength disorder renormalizes the commensurate coupling to zero, i.e. decommensuration, above a critical value of disorder. We propose that the matrix component of the porous vortex matter provides a commensurability potential for the embedded crystallites. At high fields this enhances the crystallite melting temperature, while below the decommensuration transition the crystallites decouple from the matrix, leading to a weaker enhancement of melting.

We study the classical partition function of \(L/d\) Josephson coupled layers where \(L \rightarrow \infty\) is the total length in the \(z\) direction perpendicular to the layers and \(d\) is the interlayer spacing. The elastic energy of the transverse displacement fields \(u(q,k)\) in the absence of Josephson coupling can be written as [25,26]

\[
H_{el} = \frac{1}{2L} \sum_{k,a} \int_{q} (c_{66} q^2 + c_{44}^0(k) k_z^2) |u^a(q,k)|^2 \quad (1)
\]

where a replica index \(a = 1,2,...,n\) is needed below for the disorder average. The elastic constants are [25,26]

\[
c_{44}^0(k) = \tau/(8d a_0^2 \lambda_{ab} k_z^2 \ln(1+a_0^2 k_z^2/4\pi)), c_{66} = \tau/(16d a_0^2),
\]

where \(k_z = (2/d) \sin(kd/2)\), \(\tau = \Phi_0^2 d/(4\pi^2 \lambda_{ab})\), \(\lambda_{ab}\) is the magnetic penetration length parallel to the layers, \(a_0^2\) is the area of the flux lattice unit cell and \(\Phi_0\) is the flux quantum.
i.e. \( \Phi_0 = Ba_0^2 \). Note that the Josephson coupling induces an additional term in \( c_{44} \) [26], as also shown below. The decoupling transition of the pure system (for weak Josephson coupling) is given by [16,17]:

\[
T_d = \frac{4a_0^2}{d^2} \left( \int_k \sum_{\kappa} |\tilde{c}_{44}(k)|^{-1} \right)
\]

and our principal aim is to obtain the corresponding \( T_d \) in presence of correlated disorder.

Consider a distribution of CD whose positions within a layer are random and uncorrelated. Each of the CD has a radius \( b_0 \) and their average areal density \( n_{CD} \) is low, \( n_{CD}b_0^2 \ll 1 \). A flux line has a core of radius \( \xi_0 \) which usually satisfies \( \xi_0 < b_0 \). Once a flux line is partially inside a CD it gains its core energy \( E_c \) per layer. The pinning potential per unit area is then \( U_{pin}(r) = (E_c/\xi_0^2) \sum p(r - r_i) \) with the sum on the CD positions and \( p(r) \) is a shape function, e.g. \( p(r) = 1 \) for \( r < b_0 \) and vanishes for \( r > b_0 \). The variance, neglecting CD overlaps, is therefore

\[
U_{pin}(r)U_{pin}(r') \approx E_c^2 n_{CD}(b_0/\xi_0)^4 \delta^2(r - r').
\]

The average with respect to a flux density involves an additional factor \( (\xi_0/b_0)^4 \) due to the decomposition of a sharply peaked flux into harmonics with reciprocal vectors \( \mathbf{Q} \). The replica average at temperature \( T \) is then [27]

\[
\mathcal{H}_{\text{dis}} = -\frac{1}{12L} \sum_{k} \int \frac{d^2q}{(2\pi)^2} s^2 q^2 L\delta_{k,0} u^a(q, k) u^b(-q, -k)
\]

\[
+ \frac{W}{2} \sum_{n,n'} \int d^2r \cos[\mathbf{Q} \cdot (\mathbf{u}_n^a(r) - \mathbf{u}_n^b(r))] / 2T
\]

where \( W = E_c^2 n_{CD}(b_0/\xi_0)^4 \) and only the shortest most relevant [27] \( \mathbf{Q} \) is retained. The cos term above involves vectors \( \mathbf{Q} \) and \( \mathbf{u}_n \) which in the averages below yield \( \langle \mathbf{Q} \cdot \mathbf{u}^2 \rangle = \frac{Q^2 (u_1^2 + u_2^2)}{2} \); \( u_1 \) is the longitudinal displacement which is reconsidered below, but is neglected for now as it has no effect on the decoupling transition. The parameter \( s \) is related to a long wavelength random torque coupled to a local bond angle [28] \( \gamma = (\partial_x u_y - \partial_y u_x) /2 \) since for transverse modes \( (\nabla u)^2 \) is the usual long wavelength disorder couples to \( \nabla \cdot \mathbf{u} \), hence it involves only longitudinal motions.

The long range Bragg glass properties depend on the nonlinear cos term in Eq. (4). If this cos is expanded, it yields \( \sum_{a,b} \int d^2r u^a(r, k = 0) u^b(r, k = 0) = a(k = 0) \) quadratic term which has no effect on the decoupling transition. It is therefore essential to treat the Bragg glass nonlinearities properly.

We also allow for a commensurate term of the CD density of the form

\[
\mathcal{H}_{\text{com}} = -y_c (2d/Q^2) \sum_{n,a} \int d^2r \cos[\mathbf{Q} \cdot \mathbf{u}_n^a(r)].
\]

This term assumes a predesigned component of the CD density, in addition to the random one.

Consider next the Josephson phase, i.e. the relative superconducting phase of two neighboring layers. Each flux line can be viewed as a collection of point singularities, or pancake vortices, positioned one on top of the other in consecutive layers. Around each pancake vortex the superconducting phase follows the angle \( \alpha(r) \) which changes by \( 2\pi \) in a complete rotation. The Josephson phase involves then a nonsingular component \( \theta_n(r) \) and a singular contribution from pancake vortices. The latter are positioned at \( \mathbf{R}_l + \mathbf{u}_l^a \) in the \( n \)-th layer and at \( \mathbf{R}_l + \mathbf{u}_l^{n+1} \) in the \( (n+1) \)-th layer, where \( \mathbf{R}_l \) is the undistorted position of the \( l \)-th flux line. The total Josephson phase is then

\[
\theta_n(r) = \sum_i [\alpha(\mathbf{r} - \mathbf{R}_i - \mathbf{u}_l^a) - \alpha(\mathbf{r} - \mathbf{R}_i - \mathbf{u}_l^{n+1})]
\]

\[
\approx \theta_n(r) + \sum_i (\mathbf{u}_l^n - \mathbf{u}_l^{n+1}) \mathbf{\nabla} \alpha(\mathbf{r} - \mathbf{R}_l)
\]

where the expansion is justified in the Bragg glass since the correlation length in the \( z \) direction is \( \gg d \). We define (including now the replica index) \( b^a(q, k) = -2\pi d e^{ikd/2} u^a(q, k) k_z / (qa^2) \) so that the Josephson phase is \( \theta_n^a(r) + b_n^a(r) \). Fluctuations of the \( \theta_n(r) \) field involve the Josephson energy as well as magnetic field terms, \( H_J \)

\[
\mathcal{H}_J = \frac{1}{2L} \sum_{k,a} \int \frac{d^2q}{(2\pi)^2} G_f^{-1}(q, k) |\theta^a(q, k)|^2
\]

\[
- y_J \sum_{n,a} \int d^2r \cos[\theta_n^a(r) + b_n^a(r)]
\]

where [29] \( G_f(q, k) = 4\pi d^3(\delta_{ab}^2 + k_z^2) / (\tau q^2) \). The full Hamiltonian is then \( H = H_{el} + H_{dis} + H_{com} + H_J \).

We proceed to solve this system by the variational method allowing for replica symmetry breaking (RSB) [27,30]. The form of the variational Hamiltonian \( H_0 \) is obtained by expanding the cos terms and replacing \( y \), \( y_c \), and \( W \) by variational parameters \( z_j, z_c \) and \( \sigma_{ab}(k) \), respectively. The Josephson term involves a \( \theta \) cross term which is eliminated by a shift \( \delta \theta^a(q, k) = \theta^a(q, k) - u^a(q, k) z_j (2\pi k_z / a_0^2) \exp(ikd/2) / (G_f^{-1} + z_j / d) \). Hence (repeated indices are summed)

\[
\mathcal{H}_0 = \frac{1}{2L} \sum_k \int \frac{d^2q}{(2\pi)^2} [G_a^{-1}(q, k) u^a(q, k) u^{a*}(q, k)]
\]

\[
+ [G_f^{-1}(q, k) + z_j / d] |\delta \theta^a(q, k)|^2 \quad \text{,}
\]

\[
G_a^{-1}(q, k) = \frac{[c_{66} q^2 + c_{44}(q, k) k_z^2 + z_c] \delta_{ab}}{-sL^2 \delta_{k,0} - \sigma_{ab}(k)}.
\]

The effect of the nonsingular \( \theta^a \) to shift \( c_{44}^a(q, k) \) of Eq. (1) into \( c_{44}(q, k) \),
\[ c_{44}(q, k) = c^0_{44}(k) + \frac{B^2}{4\pi(1 + \lambda^2_{\text{c}}q^2 + \lambda^2_{\text{cd}}k^2)} \]  
(10)

where \( \lambda^2_{\text{c}} = \Phi^2_0/(16\pi^3 z_d d) \). Note that the limit \( \lambda_{\text{c}} \to \infty \).

\[ \sigma_{ab}(k) = (WQ^2/2d^2T) \int_0^L dz [\cos k z \exp(-(B_{ab}(z)/2) - \delta_{ab} \sum_c \exp(-B_{ac}(z)/2)] \]
(11)

\[ z_J = y_J \exp\{-(T/2) \int_{q,k} \frac{k^2}{q^2} \left( \frac{2\pi d}{a^2} \right)^2 G^{-1}_f(q,k) \} \]
(12)

\[ z_c = y_c \exp\{-(TQ^2/4) \int_{q,k} G_{aa}(q,k) \} \]
(13)

where \( \int_{q,k} = \int d^2qd\tau/(2\pi)^3 \) and \( B_{ab}(z) \) is given by

\[ B_{ab}(z) = TQ^2 \int \frac{d^2qd\tau}{(2\pi)^3} (G_{aa}(q,k) - \cos(kz)G_{ab}(q,k)). \]

Since the disorder is \( z \) independent the off diagonal terms \( B_{ab} \) are zero independent so that \( \sigma_{ab} \) has only a \( k = 0 \) component; hence RSB is present only at \( k = 0 \). It is convenient to define \( G^{-1}_c(q,k) = \sum_a G^{-1}_a(q,k) \) so that for \( k \neq 0 \) \( G_c(q,k) = G_{aa}(q,k) \). The RSB solution reduces here to a one step form [27], hence \( G_c(q,k) \) can be written with self energies in the form

\[ G^{-1}_c(q,k) = c_{66}q^2 + c_{44}(q,k)k^2_z + z_c + \Sigma_1(1 - \delta_{k,0}) + I(k) \]
(14)

\[ \Sigma_1 + z_c = (WQ^4/16\pi d^2 c_{66}) \exp(-B_{+}/2). \]
(15)

\( B_{+} = TQ^2 \int q,k G_c(q,k) \) is a Debye Waller factor which is dominated by large \( q,k \) so that \( c_{44}^0(k) \) can be used to obtain

\[ |B_{+}| \lesssim \frac{16\pi T}{\tau} \ln \frac{c_{66}Q^2}{(\tau/8d a^2 \lambda^2_{\text{ab}} + \Sigma_1 + I(\tau/d) + z_c} \]
(16)

while the function \( I(k) \) satisfies for \( T \ll \tau \)

\[ I(k) = 4\pi c_{66}(\Sigma_1 + z_c) \int \frac{d^2q}{(2\pi)^2} \left[ (c_{66}q^2 + z_c + \Sigma_1)^{-1} - G_c(q,k) \right]. \]
(17)

Note that for \( k \to 0 \) this yields \( I(k) \sim |k| \) while \( I(k) \approx \Sigma_1 + z_c \) for large \( k \), up to logarithmic terms. A condition for melting can be estimated by a Lindemann number \( c_L \approx 0.15 \) [11] so that \( (a^2_0/\tau) = B_{+}/4\pi^2 \approx c^2_L \). Hence \( \tau \) is a measure of the melting temperature and for \( T \ll \tau \) \( B_{+} \) is small. We note that it is essential to keep the nonsingular phase \( \theta \) to obtain the correct structure factor \( G_c(q,k) \) in Eq. (14).

The decoupling transition is determined by the vanishing of \( z_J \). Eq. (12) can be written in the form

\[ z_J = y_J \exp \left\{ -\frac{1}{4} \int_{q,k} G_f(q,k) + (2\pi d/a^2)^2 G_c(q,k) + (G^{-1}_f(q,k) + z_J/d)^{-1} \right\} \]
(18)

at decoupling must be taken before \( q \to 0 \).

The variational method minimizes the free energy \( F_0 + (\mathcal{H} - \mathcal{H}_0) \) where the free energy \( F_0 \) and the average \( \langle ... \rangle_0 \) correspond to \( \mathcal{H}_0 \). This yields

\[ T_d \approx T^0_d \left[ 1 + (\Sigma_1 + z_c) \frac{8d a^2 \lambda^2_{\text{cd}}}{\tau \ln(a_0/d)} \right] \]
(19)

where an effective elastic modulus is \( c_{44}^{eff}(k) = c_{44}^0(k) + (\Sigma_1 + z_c + I(k))/k^2_z \). The \( z_c \) term is obvious here by an expansion of \( \mathcal{H}_{\text{com, Eq. (5). However the}} \Sigma \)

\[ \delta T_d \approx 2(4\pi^3 E^2_c n_{CD} b d^2 \lambda^2_{\text{cd}}) \approx 10^2 (b_0/a_0)^4 n_{CD} \lambda^2_{\text{cd}} \]
(20)

where \( E_c \approx 0.2 \tau \). For strong disorder and strong fields the CD can dominate and then \( T_d \sim 1/B \) and then a recoupling would occur at a higher field, assuming this field is still below melting.
Next we address the commensurability term, which, unlike the Josephson coupling, depends also on the longitudinal \(u(q,k)\) component. We therefore add longitudinal energy terms: first an elastic energy of the form Eq. (1) with \(c_{66}, c_{44}^{11}\) replaced by \(c_{11}\) and \(c_{44}^{11}\), respectively [26], and secondly the usual long wavelength disorder coupled to \(\nabla \cdot \mathbf{u}\) [27] which yields the form of the first term of Eq. (4) with \(s\) replaced by \(s'\). Since \(c_{22}\) originates from the \(W\) term in Eq. (4) \(\Sigma_1\) and \(I(k)\) are common to both longitudinal and transverse parts while the location of disorder parameters \(H\) hence at some critical \(s\) the effect on \(\Sigma_1\) is small, yet the structure factor for the longitudinal modes (analog of Eq. (14)) is significantly modified by the same \(\Sigma_1\) and \(I(k)\).

The equation for \(z_c\) depends also on the \(k = 0\) component of \(G_{ab}(q,k)\) which involves the long wavelength disorder parameters \(s, s'\). Using inversion methods for \(G_{ab} [27,30]\) we obtain our second principal result,

\[
z_c \sim y_c(e^{z_c/\Sigma_1 + z_c})^{1/2} (z_c/16\pi c_{66}^2 + s'Q^2/(16\pi c_{64}^2)). \tag{21}
\]

Hence at some critical \(s, s'\) (where the powers of \(z_c\) on both sides of (21) equal) the commensurability potential is renormalized to zero. We note that this renormalization is driven by \(k = 0\) terms, i.e. the same derivation is valid for a 2D system with point disorder [31].

Long wavelength disorder can generate dislocations [28] at \(s > c_{66}^2/16\pi\), i.e. below the critical value for the vanishing of \(z_c\). Furthermore, on very long scales dislocations will be induced by short wavelength CD disorder as the system is effectively two-dimensional [28,32]. We limit our discussion to a Bragg glass domain which ignores these very long scale effects.

The decoupling description neglects point defects, i.e. the nucleation of VI. The latter were studied in the absence of Josephson coupling and were shown to be generated by point disorder [19], leading to logarithmically correlated disorder for VI. Disordered CD, however, induce only a \(k = 0\) component of disorder which has exponentially decreasing correlations \(\sim (q^2 + \lambda_{ab})^{-2}(q^2 + \Sigma_1)^{-1}\), i.e. a straight flux line has exponentially decaying interactions even though each of its pancake components has a logarithmic interaction. This type of disorder cannot overcome the logarithmic interaction of pancake vortices; hence, although for any finite (Gaussian) CD disorder the ground state contains a finite density of VI perfectly aligned along \(z\), the defect transition temperature at which VI uncorrelated between layers unbind is not affected by the CD. The true decoupling, which allows for both Josephson phase fluctuations and for point defects, lies in between the above decoupling and the defect transition. For not too small Josephson coupling the transition is near the decoupling one [29,19] and therefore the results above should apply.

Finally, we address the data [10] on the melting curve showing a kink at fields \(B_k \gg B_\phi\). Within the proposed porous vortex model [10] we suggest that the ”vortex matrix”, pinned by the random CD, forms a commensurate potential. The lowest harmonic of this potential which couples to the flux periodicity has wavevector \(Q\) [harmonics with \(Q' > Q\) have a \((z_c/(z_c + \Sigma_1))^Q^2/2\pi\) factor in Eq. (21), forcing a \(z_c = 0\) solution]. Since \(s\) is a second order effect, we consider \(s' = W/d^2 \sim B_k B^2\), hence decommensation occurs at \(B \sim B_k\), the bare proportionality constant being, however, too small to account for the data. The parameters \(s, s'\) are relevant parameters within RG [28] so that their renormalized values can be large. The main result is then that elasticity dominates at large \(B\) while disorder dominates at low \(B\), driving \(z_c \to 0\), in qualitative agreement with the data. We propose then to search for an additional phase transition line within the solid phase, corresponding to decommensation, which meets the melting curve at \(B_k\).

In conclusion we have shown that columnar defects enhance the decoupling transition so that \(\delta T_\phi/T_d \sim B^2\). In contrast, the melting temperature involves the same ratio, however within a logarithm [see Eq. (16)]; hence at weak disorder the enhancement is also \(\sim B^2\) while at strong disorder only a weak \(\ln B\) effect. The \(B^2\) enhancement at strong disorder can therefore be useful in identifying a decoupling transition. Furthermore, for strong CD disorder a possibility of a reentrant transition has been found, i.e. with increasing field decoupling is followed by recoupling. We have also studied effects of a commensurate CD density and shown that its potential vanishes above a critical value of long wavelength disorder. This decommensation transition may account for the unusual kink in the melting curve data.

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