Boson Localization and Pinning by Correlated Disorder in High-Temperature Superconductors

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The physics of flux lines in the cuprate superconductors pinned by columnar defects is mapped onto boson localization in two dimensions. The theory predicts a Bose glass phase with an infinite tilt modulus and zero linear resistivity, as well as an entangled flux liquid. We describe correlations and transport in these phases, and propose a scaling theory for the irreversibility line which separates them.

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Understanding assemblies of flux lines in the cuprate high-temperature superconductors requires new ideas and phases, including entangled flux liquids [1] with a linear resistivity, and, if pointlike disorder is important, a possible vortex glass phase [2] whose nonlinear resistivity may be estimated via the collective pinning theory [3]. These theories have been stimulated by a number of striking experiments, including transport measurements [4,5] which suggest an underlying phase transition associated with the "irreversibility line" [6] in these materials.

Recently, Civale *et al.* [7] have reported greatly enhanced pinning in $YBa_2Cu_3O_7$ crystals with long aligned columns (diameter $c_0 \approx 50$ Å) of damaged material, produced by energetic Sn-ion radiation. These columnar pins produce, as well, a remarkable upward shift of several tesla in the irreversibility line. Below this line, but above the irreversibility line of the *undamaged* material, transport is presumably dominated by the correlated disorder embodied in the columnar pins, whose linear dimensions (15 μm or more) are comparable to the sample thickness. It was recently suggested [8] that flux lines under these circumstances could be understood via a mapping onto two-dimensional boson localization [9], similar to an analogy proposed earlier for vortices in pure systems [1].

In this paper, we show that this mapping predicts a low-temperature "Bose glass" phase, with flux lines localized on columnar pins, separated by a sharp phase transition from an entangled liquid of delocalized lines. Pinning in the Bose glass leads to an infinite tilt modulus, while disorder in the liquid shows up as a ridge of scattering observable via neutron diffraction [10]. We describe transport in these phases and propose a scaling theory for the current-voltage characteristics near the transition. A possible "Mott insulator" phase, in which both the tilt and compressional modulus are infinite, is also discussed.

We start with a simple model free energy F_N for N flux lines in a sample of thickness L, defined by their trajectories $\{\mathbf{r}_j(z)\}$ as they traverse a sample with columnar pins and magnetic field aligned with the z axis perpendicular to the CuO_2 planes,

$$F_N = \frac{1}{2}\tilde{\varepsilon}_1 \sum_{j=1}^{N} \int_0^L \left| \frac{d\mathbf{r}_j(z)}{dz} \right|^2 dz + \frac{1}{2} \sum_{i \neq j} \int_0^L V(|\mathbf{r}_i(z) - \mathbf{r}_j(z)|) dz + \sum_i \int_0^L U_D(\mathbf{r}_i(z)) dz . \tag{1}$$

Here, $\tilde{\varepsilon}_1$ is the local tilt modulus ($\tilde{\varepsilon}_1 < \varepsilon_1$, the line tension, due to anisotropy), $V(|\mathbf{r}_i - \mathbf{r}_j|)$ is the interaction potential between lines, and $U_D(\mathbf{r}_j)$ represents a z-independent random pinning potential. We model $U_D(\mathbf{r})$ for simplicity by a random array of identical cylindrical traps of average spacing d passing completely through the sample with well depth per unit length U_0 and effective diameter b_0 ($b_0 \approx \min\{c_0, \xi\}$), with $\overline{U_D(\mathbf{r}) U_D(\mathbf{r}')} = 2\Delta_1 \delta^{(2)}(\mathbf{r} - \mathbf{r}')$, where $\Delta_1 \approx U_0^2 b_0^4 / d^2$ and the overbar represents an average over disorder. Equation (1) is the simplest model which captures the physics of boson localization—interactions, in particular, are essential to obtain a sharp phase transition [9]. In our estimates of transport in the Bose glass phase we will suppress all factors of order unity. *Quantitative* accuracy would require $U_0 \ll \tilde{\varepsilon}_1$, which allows a simple description of vortex elasticity via the term $\frac{1}{2}\tilde{\varepsilon}_1 (d\mathbf{r}/dz)^2$. Nonlocal interactions will be discussed below.

The classical statistical mechanics associated with (1) is equivalent to the quantum mechanics of interacting bosons in two dimensions with a random static potential $U_D(\mathbf{r})$. The partition function, in particular, is determined by the ground-state energy of a fictitious quantum Hamiltonian [1,8]. We assume that each flux line spends most of its time near one of the attractive columnar pins, so that this boson Hamiltonian can be replaced by a tight-binding model defined on a lattice of sites determined by the pin positions $\{\mathbf{R}_j\}$ in a plane perpendicular to $\hat{\mathbf{z}}$. The grand canonical partition function associated with (1) is then $Z_{gr} \propto \exp(-E_0 L / k_B T)$, where E_0 is the ground-state energy of the tight-binding Hamiltonian

$$H = -\mu \sum_j a_j^\dagger a_j - \sum_{i \neq j} t_{ij} (a_i^\dagger a_j + a_j^\dagger a_i) + V_0 \sum_j a_j^\dagger a_j a_j^\dagger a_j . \tag{2}$$

Here, $\mu \propto H - H_{c1}$ is a chemical potential which fixes the flux-line density, $a_j^\dagger$ and a_j are boson creation and annihilation operators at site $\mathbf{R}_j$, $V_0 > 0$ is a typical interaction energy for two lines on the same site, and [11]

$$t_{ij} \propto \frac{U_{\rm eff}}{(E_{ij}/T)^{1/2}} \exp\left[\frac{-c_1 E_{ij}}{T}\right] \tag{3}$$

with $E_{ij} = (U_{\rm eff}\tilde{\varepsilon}_1)^{1/2} d_{ij}$ the hopping matrix element connecting sites i and j, separated by a distance d_{ij}. The parameter c_1 is a dimensionless constant of order unity and $U_{\rm eff}$ is an effective binding energy renormalized by thermal fluctuations [8]: $U_{\rm eff} = U_0 f(T/T^*)$, with $T^* = (2\tilde{\varepsilon}_1 b_0^2 U_0)^{1/2}$, $f(0)=1$, and $f(x) \approx x^2 e^{-x^2}$ for large x.

The physics of Eq. (2) was discussed by Fisher *et al.* in the context of boson localization [9]. The three phases found by these authors and their meaning for flux lines are illustrated in Fig. 1. In the entangled liquid or "superfluid" phase, flux lines are delocalized and hop freely from one columnar pin to another as they traverse the sample. All vortex trajectories "diffuse" as they wander across the sample, $\lim_{z\to\infty}\overline{\langle|\mathbf{r}_j(z)-\mathbf{r}_j(0)|^2\rangle} \approx 2D_R z$, where the brackets denote a thermal average. Both the tilt modulus c_{44} and bulk modulus c_{11} are finite. As discussed further below, the "superfluid density" in this phase is proportional to c_{44}^{-1}. The flux lines are localized to the vicinity of a few columnar pins in the Bose glass phase, which defines a localization length via $\lim_{z\to\infty}\overline{\langle|\mathbf{r}_j(z)-\mathbf{r}_j(0)|^2\rangle} = l_\perp^2$. Since the superfluid density vanishes in a Bose glass, the tilt modulus is infinite—the external field must be tipped a finite critical angle $\theta_c \approx \mathcal{O}((U_0 b_0^2 n_0/d^2\tilde{c}_{44})^{1/2})$ away from alignment with the pins before the average flux-line direction changes significantly [11]. Here, n_0 is the flux-line density and $\tilde{c}_{44}$ is the tilt modulus in the absence of pins evaluated at the relevant wave vector. The *shear* modulus c_{66} vanishes in both of these phases. The last phase occurs at low temperatures when there is one fluxon localized on every pin—the Mott insulator. The fluxon density will remain locked at the pin density for a finite range of external field strengths. Hence, *both* c_{44} and c_{11} are infinite. We defer further discussion of the Mott insulator [11], and concentrate now on the flux liquid and Bose glass phases, whose transition line $T_{\rm BG}(H)$ we identify with the irreversibility line of Civale *et al.* [7].

As $T \to T_{\rm BG}$ from below, the localization length diverges, $l_\perp \sim 1/(T_{\rm BG}-T)^{\nu_\perp}$, with $\nu_\perp \geq 1$ [9]. There is also a correlation length along the z axis, $l_\parallel \sim l_\perp^2/D_0$, which is the distance along z it takes a flux line to "diffuse" across a tube of diameter the localization length. Note that D_0 is a short-distance "diffusion" constant—the macroscopic large-scale parameter D_R *vanishes* in the vortex glass. The simplest scaling hypothesis is that D_0 remains finite at $T_{\rm BG}$, so that [9] $l_\parallel \sim 1/(T_{\rm BG}-T)^{\nu_\parallel}$ with $\nu_\parallel = 2\nu_\perp$.

To determine how disorder affects the flux liquid phase, we use a hydrodynamic approach [12], and find that the Fourier-transformed structure function, $S(\mathbf{q}_\perp, q_z) = \overline{\langle|\delta\hat{n}(\mathbf{q}_\perp,q_z)|^2\rangle}$, is [11]

$$S(\mathbf{q}_\perp,q_z) = \frac{Tn_0^2 q_\perp^2}{\hat{c}_{44}(\mathbf{q}_\perp,q_z)q_z^2 + \hat{c}_{11}(\mathbf{q}_\perp,q_z)q_\perp^2} + \Delta_1 \frac{n_0^4}{\hat{c}_{11}^2(\mathbf{q}_\perp,0)}\delta(q_z)\,, \tag{4}$$

where $\hat{c}_{44}(\vec{q})$ and $c_{11}(\vec{q})$ are nonlocal elastic moduli. In addition to the contours of constant scattering expected for neutron diffraction off flux liquids in clean systems [1], there is now a sharp ridge of intensity running down the $q_\perp$ axis due to the columnar pins. This spike measures how well the vortex lines track columnar pin trajectories, and would show up as a "central peak" in a q_z scan.

We now consider long-wavelength distortions and ask for the singularities in $c_{11} \equiv \hat{c}_{11}(\vec{0})$ and $c_{44} \equiv \hat{c}_{44}(\vec{0})$ as $T \to T_{\rm BG}$ from above. Using the boson representation for Eq. (1), and integrating out density fluctuations, leads to a free energy expressed only in terms of the boson phase $\theta(\mathbf{r}_\perp, z)$ [8],

$$F = \tfrac{1}{2} T^2 n_0^2 \int d^2r_\perp dz [c_{44}^{-1}|\nabla_\perp\theta|^2 + c_{11}^{-1}(\partial_z\theta)^2]\,,$$

which shows that c_{44}^{-1} is proportional to the boson superfluid density. Taking over the Josephson scaling analysis of Ref. [9], we find that $c_{11} = Tn_0^2 l_\perp^2/l_\parallel$ and $c_{44} = Tn_0^2 l_\parallel$, where $l_\perp$ and $l_\parallel$ are liquid-phase analogs of the Bose glass correlation lengths discussed above. The assumption that $\nu_\parallel = 2\nu_\perp$ means that c_{11} remains finite at the transition, while c_{44} diverges, $c_{44} \sim 1/(T-T_{\rm BG})^{2\nu_\perp}$. The wave-vector-dependent tilt modulus takes the scaling form $\hat{c}_{44}(\mathbf{q}_\perp,q_z) = l_\parallel \Phi(q_\perp l_\perp, q_z l_\parallel)$ with $\Phi(0,y) \sim 1/y$ for small y. Using this result, we find that the contours of constant scattering near the origin in Eq. (4) (given by $q_z \propto q_\perp$) are "pinched" down so that $q_z \propto q_\perp^2$ as $T \to T_{\rm BG}^+$.

The dynamics of the Bose glass is determined by the competition between the 2D array of columnar pinning centers and 3D thermal fluctuations of vortex lines. A detailed picture can be obtained in the region $T < T^*$, $B < B^* = \Phi_0/d^2$, where each vortex is localized on *one* columnar defect and the pins outnumber the vortices. For Y-Ba-Cu-O we estimate $T^* = 60$–80 K and B^* may be of order several tesla depending on the radiation dose, so this is a sizable regime. The boson mapping then reduces the single vortex dynamics to a problem of hopping conductivity for 2D localized bosons with hopping

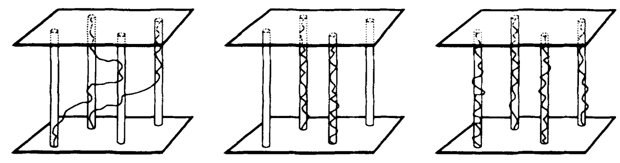

FIG. 1. Schematic of flux lines attracted by columnar pins in the flux liquid, Bose glass, and Mott insulator phases.

probabilities t_{ij}. The electric field in the usual hopping conductivity problem corresponds to the in-plane current density J in vortex dynamics, whereas current density and conductivity from hopping map, respectively, onto vortex velocity (i.e., voltage) and resistivity in vortex Bose glass dynamics. Upon noting that the inverse sample thickness plays the role of temperature in the boson representation [1], one can reproduce for Bose glass dynamics the rich variety of hopping conductivity phenomena in semiconductors [13] by transcribing the proper quantities. It can be shown that a current J_z applied perpendicular to the CuO_2 planes is equivalent to a fictitious "magnetic field" acting on the bosons. We discuss below only the most characteristic cases leaving detailed consideration for a subsequent paper [11].

The critical current in this strongly pinned Bose glass is just $J_c = cU_0/\Phi_0 b_0$. When $J_1 < J < J_c$, where $J_1 = cU_0/\Phi_0 d$, the flux motion occurs via a thermally activated "half loop" configuration with energy $E^* \approx (\tilde{\varepsilon}_1)^{1/2} U_0^{3/2}/f_L$, where $f_L = c^{-1} J\Phi_0$. The length of the critical unbound line segment is $l^* \approx (U_0 \tilde{\varepsilon}_1)^{1/2}/f_L$. This implies a nonlinear current-voltage characteristic, $V \sim \exp[-(E_k/T)J_1/J]$, with $E_k = d(\tilde{\varepsilon}_1 U_0)^{1/2}$, as found recently by Konczykowski *et al.* [14]. Dispersion in vortex binding energies, due to interactions, nearby pins, etc., will change this result only in the limit $B \ll B^*$ [11].

For $J < J_1$ the transverse displacement of the liberated vortex segment exceeds the mean distance between defects and the transition of the vortex line from one rod to another takes place via a thermally activated double kink configuration which throws a vortex segment onto an adjacent columnar defect in strong analogy to dislocation motion over a (random) Peierls potential. Samples thin enough that the dispersion of energy levels γ associated with (2) is negligible $(L\gamma < E_k)$ will now exhibit a nearest-neighbor percolative hopping conductivity with $V/J \sim \exp(-c_2 E_k/T)$, where the constant c_2 can be found explicitly from percolation theory [13]. Vortex motion when $B \ll B^*$ in this regime is determined by dead ends in the percolation network [15], and nonmonotonic current-voltage curves will result [11].

The dispersion of flux binding energies is essential for $J < J_1$ and $L > L_1 = E_k/\gamma$. We now have variable-range hopping conductivity of the most weakly found flux lines, with all lower energies filled [13]. We assume that the repulsive intervortex interaction effectively excludes multiple occupancy of pinning sites. Vortex motion then requires hopping (via double *superkink* formation) to a distant defect with nearly the same energy, since there is a large $(\sim L)$ barrier to hopping to nearby higher-energy states in a sufficiently thick sample. At very small currents for $L_1 < L < \infty$ Ohmic Mott variable-range hopping (VRH) conduction takes place: $V/J \sim \exp[-(L/L_0)^{1/3}]$, where $L_0 \simeq T^3 g(\mu)/\tilde{\varepsilon}_1 U_0$, and $g(\mu)$ is the density of states obtained from (2) evaluated at the energy of the highest occupied state. Stronger currents reduce the typical jump distance $r_\perp$, which is now determined self-consistently from $r_\perp = [g(\mu) f_L r_\perp]^{-1/2}$ ($f_L r_\perp$ is the fall in vortex potential energy over a typical jump $r_\perp$). This gives $r_\perp = [f_L g(\mu)]^{-1/3}$ and determines a current-dependent hopping probability $\sim \exp(-r_\perp E_k/dT)$. The transition to this non-Ohmic VRH behavior with current-voltage characteristic $V/J \sim \exp[-(E_k/T)(J_0/J)^{1/3}]$, where $J_0 = c/\Phi_0 g(\mu) d^3$, occurs at $J_L = cE_k/\Phi_0 L d \ll J_0$. When $J > J_2 = J_1 [U_0 g(\mu) d^2]^{1/2}$, transport is dominated by thermal activation from a single rod as discussed above. Note that interactions will lower $g(\mu)$ (as in the "Coulomb gap" of semiconductors), but not to zero provided the interaction *range* (here, of order λ) remains finite [13].

Inhomogeneity in the z direction can change the Bose glass dynamics in thick samples, although probably not above the irreversibility line in *undamaged* materials. The kink propagation along the rods may slow down significantly at large scales with kink velocity decaying with time as $v \sim t^{-(1-\mu)}$, $\mu < 1$, possibly leading to dynamically entangled vortex configurations [11].

For $B \gg B^*$, every pin is occupied and the excess vortices go into the interstitial spaces between the random columnar pins. These vortices will still be localized at low temperatures by interactions with vortices trapped on the pinning sites. The critical current is now $J_c \approx c\tilde{\varepsilon}_1/\Phi_0 d^2 n_0^{1/2}$, provided the Larkin-Ovchinnikov length r_c estimated from 2D collective pinning theory [3] is less than d. When $r_c \gg d$, standard 2D collective pinning formulas for J_c apply. For $T^* < T < T_{BG}(H)$, the localization length $l_\perp \gtrsim d$, the vortex line wanders between different columnar defects and can be pinned effectively only by the ensemble of the rods. Flux creep is now determined by hopping conductivity in a continuous random potential [13] with characteristic spatial scale $l_\perp$. In *all* regions below $T_{BG}(H)$ (but above the irreversibility line in undamaged materials) we expect the asymptotic laws $V \sim \exp[-(J_0/J)^{1/3}]$ in the thermodynamic limit and $V/J \sim \exp[-(L/L_0)^{1/3}]$ in finite thickness samples, as found explicitly for $T < T^*$, $B < B^*$.

By looking for an instability in the hydrodynamic theory for $T > T_{BG}$ when $B < B^*(U_0/\tilde{\varepsilon}_1)$ we find, following [12],

$$T_{BG}(B) = \text{const} \times T^* (\Phi_0/d^2 B)^{1/4} . \qquad (5)$$

Using the dynamic collective pinning approach of [16] we estimate the critical current in this regime as $J_c \approx (c\Delta_1^{1/2}/\Phi_0 d^2)(T^*/T)^4$ and find a localization length $l_\perp(T) \approx d(T/T^*)^2$. Note that Eq. (5) follows by setting $l_\perp(T) \approx n_0^{-1/2}$. Of course, the true localization length actually *diverges* as $T \to T_{BG}^-$. When $B > B^*(U_0/\tilde{\varepsilon}_1)$, we expect that $T_{BG}(B)$ will be close to the pure system melting temperature, with $J_c \sim 1/T^{5/2}$. For $T > T_{BG}(B)$, one expects the vortex liquid to exhibit the usual linear resistivity [12,16].

A scaling theory of the dynamics near the irreversibility line $T_{BG}(B)$ is easily constructed. We use the diverging lengths $l_\perp$ and $l_\parallel$, and assume that the time scale τ to relax a fluctuation of this size is given by a new critical exponent z', $\tau \sim l_\perp^{z'}$. Following [2], the natural scaling assumption is that the electric field and current are related by $El_\perp^{1+z'} = F_\pm(l_\perp l_\parallel J\Phi_0/cT)$, where $F_\pm(x)$ is a universal scaling function. Henceforth, we set $\nu_\parallel = 2\nu_\perp \equiv 2\nu'$. Above T_{BG}, we expect at small x $F_+(x) \sim x$, which leads to a linear resistivity $\rho \sim (T - T_{BG})^{\nu'(z'-2)}$. Below T_{BG}, we expect $F_-(x) \sim \exp(-\text{const}/x^{1/3})$ for small x, and a consistent scaling theory *at* T_{BG} requires that $E \sim J^{(1+z')/3}$. The analogous predictions for the vortex glass model are [2] $\rho \sim (T - T_{VG})^{\nu(z-1)}$ and $E \sim J^{(1+z)/2}$, where ν and z are the vortex glass correlation length and dynamic exponents.

Despite its formal resemblance to the vortex glass scaling theory, the Bose glass model differs in a number of important ways. In contrast to point disorder, which *promotes* flux-line wandering and entanglement [1], correlated disorder *inhibits* wandering and promotes localization. This difference shows up clearly upon considering the response to a perpendicular field $H_\perp$ at the transition: The generalized scaling ansatz $El_\perp^{1+z'} = F_\pm(l_\perp l_\parallel J\Phi_0/cT, H_\perp l_\perp l_\parallel/\Phi_0)$ leads to the prediction that the critical angle $\theta_c \sim (T_{BG} - T)^{3\nu'}$ as $T \to T_{BG}^-$. This upward cusp in the apparent irreversibility temperature as a function of angle is clearly evident in the data of Worthington *et al.* [17], suggesting that correlated disorder (in the form of twin boundaries) plays a key role in these experiments. The vortex glass hypothesis would predict a *smooth* variation of the irreversibility line with angle. Such differences arise because the correlated volume near the vortex glass transition is assumed to diverge isotropically [2], while the Bose glass correlations diverge with two anisotropic correlation lengths. Note also that the very existence of a vortex glass is open to question (its lower critical dimension appears to be close to 3 [2]), while the Bose glass rests on a firmer foundation.

We note in conclusion that the experiments of Refs. [4] and [5] may well be affected by correlated disorder, in the form of twin boundaries [4,5], and possibly forests of screw dislocations [4]. Recent low-temperature neutron diffraction experiments [10] and twin boundary pinning measurements [18] support this view, at least for twinned single crystal Y-Ba-Cu-O samples. Simple estimates of the relative importance of point and correlated disorder suggest that correlated disorder may dominate at long wavelengths [11]. A fit of the Bose glass scaling laws to the data of [4] and [5] yields exponents $\nu' \approx 1.1-1.6$ and $z' \approx 6.5-8.0$. Precise transport experiments on untwinned single-crystal samples (to test the vortex glass hypothesis) and on the columnar pin samples of Civale *et al.* [7] (to test the Bose glass hypothesis) would be highly desirable.

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