Theory of Plastic Vortex Creep

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We develop a theory for plastic vortex creep in a topologically disordered (dislocated) vortex solid phase in type-II superconductors in terms of driven thermally activated *dislocation* dynamics. *Plastic barriers* for dislocations show a power-law divergence at small driving currents j, $U_{\mathrm{pl}}(j) \propto j^{-\mu}$, with $\mu = 1$ for a single dislocation and $\mu = 2/5$ for creep of dislocation bundles. This implies a suppression of the creep rate at the transition from the ordered vortex phase ($\mu = 2/11$) to the dislocated glass and can manifest itself as an observed increase of the apparent critical current (*second peak*). Our approach applies to general dynamics of disordered elastic media on a random substrate.

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One of the most fascinating dynamic phenomena of complex systems with internal degrees of freedom is the thermally activated motion of elastic media in a random environment (creep) characterized by a highly nonlinear response to a dc driving force, F [1]: $v \sim \exp(-\mathrm{const}/TF^{\mu})$, where v is the velocity, T is the temperature, and μ is the exponent depending on the geometry and the dimensionality of the driven medium. The concept of thermally activated creep is ubiquitous in disordered systems and describes a wealth of low temperature transport phenomena including the dynamics of dislocation and/or domain walls in inhomogeneous environments [1,2], driven vortex lattices and charge density waves [3–6]. The derivation of the fundamental creep feature, energy barriers $U(j) \sim j^{-\mu}$ diverging at small driving forces (currents j), was based on the *elastic* behavior of the pinned structures; thus in the common view creep behavior is implicitly attributed to the elastic medium free of topological defects.

The description of thermally activated dynamics of amorphous structures containing a large amount of topological defects is a long-standing problem that appeared first in the theory of work hardening and related relaxation processes in dislocated solids. In the context of vortex physics the quest for the description of creep in a topologically disordered medium was motivated by the observation of the disorder-induced transition between a low-field quasilattice or Bragg glass (BrG) [5,7], the phase free of topological defects, and a high-field phase, characterized by an enhanced apparent critical current [8]. The latter phase was suggested to be a *topologically disordered,* dislocated vortex state or amorphous vortex glass (AVG) [9–11]. In a recent series of experiments [12,13] the phase coexistence characteristic for a first-order transition was established, and creep barriers in the high-field vortex state were shown to diverge *faster* than creep barriers in the low-field elastic phase.

On the theoretical side the problem is to find an appropriate quantity enabling parametrization of the amorphous phase, since the starting point for the usual creep description–elastic manifold and/or perfect elastic lattice do not seem to apply. A first step towards such a quantitative description was made in Ref. [14], where all phase transitions between vortex lattice phases, including both disorder-induced transition and thermal melting, were described in terms of *dislocation-mediated* behavior, and a free energy functional $\mathcal{F}[n_D]$ for an ensemble of directed dislocations in the presence of thermal fluctuations *and* quenched disorder was derived (n_D is the areal dislocation density). The BrG-AVG transition was found to be of weak first order in accordance with the experimental results of Ref. [13]: at the BrG-AVG transition, dislocations enter with a density $n_D \simeq R_a^{-2}$ given by the positional correlation length R_a on which typical vortex displacements are of the order of the lattice spacing a [6]. Upon increasing the magnetic field up to the critical point the dislocation density of the AVG increases to vortex liquidlike values $n_D \simeq a^{-2}$ such that the AVG and vortex liquid phases become thermodynamically indistinguishable at the critical point.

In this Letter, building on the aforementioned ideas, we propose a quantitative description of plastic creep in terms of the dislocation degrees of freedom. We find a critical plastic current j_{pl} below which dislocations are *collectively pinned* and plastic creep occurs via the *activated* motion of collectively pinned dislocation lines. The critical plastic current is lower than the critical current for vortex depinning $j_{\mathrm{pl}} < j_c$, hence plastic motion of depinned dislocations sets in before viscous flow of the entire vortex lattice can occur. We derive the associated *plastic* creep energy barriers $U_{\mathrm{pl}}(j) \sim j^{-\mu_{\mathrm{pl}}}$ diverging *infinitely* at $j \to 0$. We calculate the pinning force acting on dislocations from the Peach-Köhler force exerted on vortices by the pinning centers. We show that an external current sent through a dislocated vortex lattice generates a Peach-Köhler force with a component causing dislocation glide. The interplay of these two forces determines the glassy dislocation dynamics, in particular the depinning threshold for dislocation glide and the energy barriers for plastic creep below the depinning threshold.

The energy of a single *straight* vortex lattice dislocation of length L and with Burger's vector $\mathbf{b}$ consists of the core energy and of the logarithmically diverging contribution

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from the long-range elastic strains [15]: $E_0 = LE_D[c_D + \ln(L_\perp/a)]$, where $L_\perp$ is the lateral system size, $E_D = Kb^2/4\pi$, $K = \sqrt{c_{44}c_{66}}$ is the *isotropized elastic constant* in the rescaled coordinate $z = \tilde{z}\frac{1}{2}\sqrt{c_{44}/c_{66}}$, and $c_D \approx 1$ is found numerically (c_{44} and c_{66} are the tilt and shear moduli of the vortex lattice, respectively). Bending of the dislocation line costs an elastic energy associated with its stiffness ϵ_D. Hence, the single directed dislocation line—parametrized by its displacement field $\mathbf{u}_D(z)$—is described by the Hamiltonian

$$\mathcal{H}_D[\mathbf{u}_D] = E_0 + \int dz\, \frac{1}{2}\, \epsilon_D(\partial_z \mathbf{u}_D)^2, \tag{1}$$

where the stiffness $\epsilon_D \approx E_D \ln(1/k_z a)$ has a logarithmic dispersion due to the long-range strain field.

We find the driving force acting on an edge dislocation with $b \parallel x$ when a transport current $j \parallel y$ is sent through the sample. The driving current creates a magnetization gradient, determined by Maxwell's equation $\nabla \times B = \frac{4\pi}{c}\mathbf{j}$. This gradient, in turn, induces shear strains in the vortex lattice: $\partial_x u_y = \partial_x a = a\frac{2\pi}{c}\frac{j}{B}$. The resulting shear stresses give rise to a glide component of the driving Peach-Köhler force [15] (per dislocation length):

$$F_x^{\mathrm{drive}} = \sigma_{yx} b = bKa\, \frac{2\pi}{c}\, \frac{j}{B}\,. \tag{2}$$

Note that compression stress leads only to dislocation climb, which can be neglected as a slow process requiring diffusion of interstitials [16].

The displacements induced by the magnetization gradient can be accommodated only via the creation of a *stationary* superstructure of regularly spaced bands of dislocations with Burger's vectors having a y component [17]. Our representative "test" dislocation moves through this superstructure, which is similar to grain boundaries appearing in bent atomic crystals [15]. Since such dislocation bands are essentially free of shear stresses [15], they do not contribute to the bulk driving shear experienced by the test dislocation (2) everywhere between the bands, and, therefore, the superstructure does not affect the glide motion.

The random pinning potential $V_{\mathrm{pin}}(\mathbf{r})$ "seen" by the vortex array also produces Peach-Köhler–type forces acting on dislocations. To find these pinning-induced Peach-Köhler forces, we first have to determine the random stress exerted by the pinning potential on a frozen-in elastic displacement configuration $u_{\mathrm{el}}(\mathbf{R}, z)$ of the vortex lattice: $V_{\mathrm{pin}}(\mathbf{R} + \mathbf{u}_{\mathrm{el}}, z) = \sigma_{ij}^{\mathrm{pin}}(\mathbf{R} + \mathbf{u}_{\mathrm{el}}, z)\nabla_i u_{\mathrm{el},j}$. The spatial distribution of the pinning stresses is thus governed by the quenched distribution of the elastic displacements u_{el} of the *dislocation-free* collectively pinned vortex array. The latter shows different scaling behaviors depending on the spatial regime in question: (i) Small distances where vortex displacements u are smaller than the coherence length ξ, and perturbation theory applies [18]. (ii) The intermediate scales where $\xi \lesssim u \lesssim a$, and disorder potentials seen by different vortices are effectively *uncorrelated.* This regime is captured in so-called *random manifold* (RM) models [6,7], leading to a roughness $\tilde{G}(\mathbf{r}) = \overline{\langle[\mathbf{u}_{\mathrm{el}}(\mathbf{r}) - \mathbf{u}_{\mathrm{el}}(0)]^2\rangle} \simeq a^2(r/R_a)^{2\zeta_{\mathrm{RM}}}$, where $\zeta_{\mathrm{RM}} \approx 1/5$ for the $d = 3$-dimensional RM with two displacement components. The crossover scale to the asymptotic behavior is the *positional correlation length* R_a, where the average displacement is of the order of the vortex spacing: $u \approx a$. (iii) The asymptotic *Bragg glass* regime where the a periodicity of the vortex array becomes important for the coupling to the disorder, and the array is effectively subject to a *periodic* pinning potential with period a [5]. Here the *logarithmic* roughness $\tilde{G}(\mathbf{r}) \approx (a/\pi)^2 \ln(er/R_a)$, i.e., $\zeta_{\mathrm{BrG}} = \mathcal{O}(\log)$ [5,7] takes over.

For the physics of dislocations on scales $>a$, only the RM and BrG regimes are relevant. We obtain approximately Gaussian distributed quenched stresses with $\overline{\sigma_{ij}^{\mathrm{pin}}} = 0$ and $\overline{\sigma_{ij}^{\mathrm{pin}}(\mathbf{k})\sigma_{ij}^{\mathrm{pin}}(\mathbf{k}')} = \Sigma^{\mathrm{pin}}(k)\,(2\pi)^3 \times \delta(\mathbf{k} + \mathbf{k}')$ with $\Sigma^{\mathrm{pin}}(k) = K^2 k^2 G(k)$, i.e.,

$$\Sigma^{\mathrm{pin}}(k) = K^2 a^2 k^{-1} \begin{cases} \mathrm{BrG:} & 1 \\ \mathrm{RM:} & B_{\mathrm{RM}}(kR_a)^{-2\zeta_{\mathrm{RM}}}\,, \end{cases} \tag{3}$$

determined by the elastic correlations $G(k)$ with a numerical constant B_{RM}. The RM result holds for $kR_a > 1$, while the BrG behavior occurs for $kR_a < 1$.

To derive the correct Peach-Köhler pinning force it is crucial to take into account not only the "direct" quenched pinning stresses $\sigma_{ij}(\mathbf{r})$ but also the elastic stresses $\sigma_{ij}^{\mathrm{el}}$ themselves which are responding to the same pinning potential and hence tend to *relax* the (longitudinal) components of the stress. A simple but lengthy calculation shows that the pinning Peach-Köhler force on a dislocation element $d\mathbf{R}$,

$$dF_\alpha^{\mathrm{pin}} = \epsilon_{\alpha\beta l}(\sigma_{\beta k}^{\mathrm{pin}} + \sigma_{\beta k}^{\mathrm{el}})b_k dR_l\,, \tag{4}$$

is rotation-free $(\nabla \times d\mathbf{F}^{\mathrm{pin}})_\gamma = 0$; the corresponding potential plays the role of the *pinning Hamiltonian*:

$$\mathcal{H}_D^{\mathrm{pin}}[\mathbf{u}_D] = \int dz\, \frac{d\mathbf{u}_{D,l}}{dz}\, b_k g_{kl}^{\mathrm{pin}}(\mathbf{u}_D, z)\,,$$
$$\overline{g_{kl}^{\mathrm{pin}}(\mathbf{k})g_{k'l'}^{\mathrm{pin}}(-\mathbf{k})} \simeq k^{-2}\Sigma^{\mathrm{pin}}(k)\delta_{kk'}\delta_{ll'} = K^2 G(k)\delta_{kk'}\delta_{ll'}\,. \tag{5}$$

By combining Eqs. (1), (2), and (5), one arrives at the free energy $\mathcal{H}_D[\mathbf{u}_D] + \mathcal{H}_D^{\mathrm{pin}}[\mathbf{u}_D] - \int dz\, \mathbf{F}^{\mathrm{drive}} \cdot \mathbf{u}_D$ giving an adequate description of an ensemble of pinned dislocations.

Starting with statics we discuss roughening of the dislocation in the presence of disorder. The typical squared pinning energy fluctuations upon displacing a dislocation segment L over a distance u_D (in the xz glide plane) can be calculated from (5),

$$\begin{aligned} E_{\mathrm{pin}}^2(L, u_D) &\simeq b^2K^2 L u_D \int_0^L dz \int_0^{u_D} dx \\ &\quad\times \int \frac{d^3k}{(2\pi)^3}\, k^2 G(k) e^{ik_z z + ik_x x} \\ &\simeq E_D^2 L u_D \begin{cases} \mathrm{RM:} & (u_D/R_a)^{2\zeta_{\mathrm{RM}}} \\ \mathrm{BrG:} & 1 \end{cases}, \end{aligned} \tag{6}$$

whereas the corresponding elastic bending energy of the dislocation is $E_{el}(L, u_D) \simeq \frac{1}{2} E_D \ln(L/a) u_D^2/L$ (1). Optimization gives a dislocation roughness as

$$u_D(L) \simeq L \begin{cases} \text{RM:} & (L/R_a)^{2\zeta_{RM}/(3-2\zeta_{RM})} \mathcal{O}(\log) \\ \text{BrG:} & \ln^{-2/3}(L/a) \end{cases}, \quad (7)$$

i.e., exponents $\zeta_D \approx \frac{15}{13}$ for RM scaling ($L < R_a$) and $\zeta_D \approx 1 - \log^{2/3}$ for BrG scaling. The instability with respect to dislocation proliferation is signaled by anomalous energy gains if $\zeta_D > 1$, i.e., in the RM regime. In the BrG regime the energy balance is more subtle, and, to conclude on the stability at $\zeta_D < 1$, one has to convert the result (7) into an approximate renormalization group (RG) scheme: the energy gain due to roughening is $\Delta E \sim E_D L \ln^{-1/3}(L/a) \sim E_D L \tilde{\epsilon}_D(L)^{-1/3}$, where the logarithmic correction is identical to the dimensionless line tension $\tilde{\epsilon}_D = \epsilon_D/E_D$ on the scale L. Interpreting $\Delta E/(E_D L)$ as disorder correction to the line tension $\tilde{\epsilon}_D(L)$ on the scale L and summing these corrections up successively on each scale, together with the bare tension $\tilde{\epsilon}_D^0(L) = \ln(L/a)$, one obtains an integral RG equation,

$$\tilde{\epsilon}_D(\ln L) = \int_0^{\ln L} d\ell [1 \pm \tilde{\epsilon}_D(\ell)^{-1/3}], \quad (8)$$

equivalent to the result of Ref. [10]. Integration shows that corrections to $\tilde{\epsilon}_D^0(L)$ are irrelevant and hence the BrG regime is stable with respect to dislocation formation [10]. The detailed stability analysis for both regimes was given in Ref. [14]: the BrG-AVG transition is weakly first order, and dislocations proliferate with the density $n_d \simeq R_a^2$ defined by the scale R_a of crossover between the unstable RM and the stable BrG regimes.

Now we extend our scaling analysis to the dynamic behavior of the driven dislocation. A dislocation segment of length L ($<R_a$) and laterally displaced over u_D gains not only the energy (6) due to pinning potential but also an energy $L u_D F^{\text{drive}}$ by the driving force (2) while it loses bending energy, and has thus a free energy:

$$\frac{F(u_D, L)}{E_D} \simeq \frac{u_D^2}{L} - (L u_D)^{1/2} \left(\frac{u_D}{R_a}\right)^{\zeta_{RM}} - L u_D \frac{8\pi^2}{c} \frac{j}{B}. \quad (9)$$

Minima of the pinning potential are typically separated by distances $u_D \sim a$, determined by the core size of the dislocation. A dislocation confined within *one* such minimum is collectively pinned on a *plastic pinning length,*

$$L_{pl} \simeq a \left(\frac{R_a}{a}\right)^{2\zeta_{RM}/3} < R_a, \quad (10)$$

which is obtained by minimizing $F(u_D = a, L)$ with respect to L at zero current. This is the analog of the Larkin pinning length $L_c \simeq \xi \delta^{-1/3}$ of the single vortex, where δ is the dimensionless pinning strength of Ref. [6]. To depin the dislocation, the driving force in (9) has to exceed the pinning force $F^{\text{pin}}(L_{pl}) \simeq \frac{E_D}{a} L_{pl}(a/R_a)^{4\zeta_{RM}/3}$ on a segment of length L_{pl}. This determines a *critical plastic current* j_{pl},

$$j_{pl} \simeq \frac{c}{8\pi^2} \frac{B}{a} \left(\frac{a}{R_a}\right)^{4\zeta_{RM}/3} \simeq j_0 \left(\frac{2\pi H_{c2}}{B}\right)^{-7/4} \delta^{5/18}, \quad (11)$$

where $j_0 \simeq \frac{c}{6\sqrt{3}\pi} H_{c2}/\xi$ is the depairing current. Comparing this result to the depinning current $j_c \simeq j_0 \delta^{2/3}$ for a single vortex, one finds $j_{pl} < j_c$ for typical disorder strengths $\delta \sim 10^{-3}$, and the plastic motion of dislocations may occur even if the vortices themselves are still pinned. One concludes that at sufficient currents the plastic motion dominates transport in the dislocated AVG phase.

However, the existence of a current $j_{pl} > 0$ implies that dislocations are *pinned* at low currents and, therefore, the plastic motion for $j \ll j_{pl}$ occurs only via activation over *diverging* plastic energy barriers $U_{pl}(j) \sim j^{-\mu_{pl}}$, i.e., as a *plastic creep.* The typical segment size $L(j)$ for activated motion at $j \ll j_{pl}$ is determined from (9) by balancing the energy gain due to the driving force $E_{\text{drive}} \simeq \frac{8\pi^2}{c} \frac{j}{B} L^{1+\zeta_D}$ against the pinning energy $E_{\text{pin}} \simeq E_D L^{2\zeta_D - 1}$ of a dislocation line roughened according to the above result (7). This yields

$$U_{pl}(j) \simeq E_D a \left(\frac{a}{R_a}\right)^{2\zeta_{RM}/3} \left(\frac{j_{pl}}{j}\right)^{(2\zeta_D - 1)/(2-\zeta_D)} \quad (12)$$

and we obtain the scaling law $\mu_{pl} = (2\zeta_D - 1)/(2 - \zeta_D)$ (as for single vortex creep [6]) relating the plastic creep exponent μ_{pl} to the dislocation roughness. We find $\mu_{pl} = \frac{17}{11}$ in the RM regime $[L(j) < R_a]$ and $\mu_{pl} = 1$ in the BrG regime $[L(j) > R_a]$. Both exponents are considerably larger than their counterpart $\mu = 2/11$ for elastic single vortex creep, showing that plastic creep rates are much smaller than elastic creep rates.

So far we have focused on a *single* dislocation. Now we turn to an ensemble of interacting dislocations. On large scales exceeding the dislocation spacing R_D, which varies from $R_D \simeq R_a$ at the AVG-BrG transition to $R_D \simeq a$ at the critical point [14], interactions become essential and plastic creep is governed by the motion of dislocation *bundles* in a glide plane (xz plane). Deformations $u_D(x, z)$ of such a 2D bundle can be described by an elastic Hamiltonian with tilt modulus $K_z \simeq E_D/R_D$ and the compression modulus $K_x \simeq R_D \partial_{R_D}^2 [R_D^2 f(R_D)]$ which can be calculated from the dislocation-free energy $f(R_D)$, see Ref. [14] (in the absence of disorder, one finds $K_x \simeq E_D/R_D$). By including the pinning energies, we obtain the Hamiltonian,

$$\mathcal{H}[u_D(x, z)] = \int dx\, dz \frac{1}{2} \{K_x(\partial_x u_D)^2 + K_z(\partial_z u_D)^2\} + \sum_i \int dz\, \mathcal{H}_D^{\text{pin}}[\mathbf{b}_i, \mathbf{u}_D(iR_D, z)]. \quad (13)$$

The dislocation bundle contains dislocations of *opposite* signs with the same density to avoid the accumulation of stress. Hence the sum over the dislocation index i in (13) goes over *alternating* Burger's vectors $\mathbf{b}_i \parallel x$. On scales

$L_x \gg R_D$, dislocations couple effectively as dipoles to disorder and we obtain for the bundle disorder energy fluctuations $E^b_{\rm dis}(L_z, L_x, u_D) \simeq E_{\rm pin}(L_z, u_D)$ with $E_{\rm pin}(L_z, u_D)$ from (6). This has to be balanced against the elastic energy $E_{\rm el}(L_z, L_x, u_D) \simeq \sqrt{K_x K_y}\, u_D^2$ with $L_z \simeq \sqrt{K_z/K_x}\, L_x$ resulting in a roughness

$$u_D(L_z) \sim L_z^{1/3} R_D^{2/3} \begin{cases} \text{RM:} & (\frac{L_z^{1/3} R_D^{2/3}}{R_a})^{2\zeta_{\rm RM}/3 - 2\zeta_{\rm RM}} \\ \text{BrG:} & 1 \end{cases}. \tag{14}$$

Note that the bundle roughness is *reduced* as compared to that of single dislocations: $\zeta_D \approx \frac{5}{13}$ for RM scaling ($L < R_a$) and $\zeta_D \approx \frac{1}{3}$ for BrG scaling. Analogous to the case of the single dislocation, one easily establishes the plastic creep exponents for bundle creep, in particular one finds the same scaling relation $\mu_{\rm pl} = (2\zeta_D + d - 2)/(2 - \zeta_D)$ as for d-dimensional vortex bundles [6]. This gives $\mu_{\rm pl} = \frac{10}{21}$ in the RM regime and $\mu_{\rm pl} = \frac{2}{5}$ in the BrG regime. A crossover from the single dislocation to the bundle scaling occurs at currents $j < j^b$, where $L(j^b) \simeq \sqrt{K_z/K_x}\, R_D$. For $R_D \simeq a$, one finds $j^b \simeq j_{\rm pl}$, meaning that only plastic bundle creep can be measured above the critical point defined by $R_D \simeq a$.

By the spirit of the derivation, our results seem to apply to superconductors with pronounced vortex lines (like YBCO) rather than to the layered BSCCO. Yet the creep exponent measured within the AVG phase in Ref. [12] is strikingly close to $\mu_{\rm pl} = \frac{2}{5}$. In layered compounds such as BSCCO the vortex lattice consists of pancakes only weakly coupled across different layers by their magnetic interaction. Whereas the Bragg glass phase can persist at low magnetic fields due to the small interlayer coupling, the layers essentially decouple at higher fields, and the resulting 2-dimensional pancake lattices are unstable with respect to dislocation formation in the presence of pinning by point defects [19]. Also the 2D dislocations exhibit plastic creep, as can be seen from the following argument. Let us consider a pair of opposite edge dislocations a distance u_D apart with an interaction energy $E_{\rm int}(u_D) = E_D \ln(u/a)$, where $E_D = c_{66} = b^2/2\pi$ in 2D. The typical energy gain from the disorder has been calculated in Ref. [19] to be $E_{\rm pin}(u_D) \simeq E_D \ln^{3/2}(u_D/R_a)$. On the one hand, the 2D Bragg glass is unstable to dislocation formation because the disorder-induced *valleys* exceed the interaction energy: $E_{\rm pin}(u_D) \gg E_{\rm int}(u_D)$ for $u_D \gg R_a$. Furthermore, $E_{\rm pin}(u_D)$ also gives the typical size of the energy *barriers* between optimized dislocation positions, i.e., the barriers for plastic creep: $U_{\rm pl} \simeq E_{\rm pin}(u_D)$. In the presence of the driving current the dislocations will gain an energy $E_{\rm drive} \simeq \frac{8\pi^2}{c}\frac{j}{B} u_D$ from the force (2) which can pull the dislocation pair apart over these energy barriers. Balancing both terms, we find *logarithmically* diverging barriers for plastic creep in 2D:

$$U_{\rm pl}(j) \sim E_D \ln^{3/2}(1/j). \tag{15}$$

In conclusion, we have developed a theory of plastic creep in terms of the dislocation dynamics in the pinned vortex lattice. We have found diverging barriers for plastic vortex transport, in agreement with the experimentally observed low creep rates or high apparent critical currents. The obtained results are relevant for other systems where glassy dynamics is controlled by topological defects, for example, charge density waves in disordered crystals and/or work-hardened solids.

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