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Simple model for critical currents in anisotropic type-II superconductors

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Recently, a continuum theory of the mixed state in type-II superconductors was developed by two of us (Mathieu and Simon). An unorthodox model for the critical currents in soft materials proceeded from this theory. We show that this surface pinning model can be easily extended to soft high-T_c superconductors (HTSC's). It successfully predicts the order of magnitude of critical currents such as observed in monocrystals or epitaxial films, and accounts for the existence and the position in the (H,T) plane of the irreversibility line. In the framework of our model, the vanishing of critical currents appears as a simple consequence of the strong anisotropy of HTSC's; there is no need for any bulk phase transition of the vortex array.

The model for the critical current in soft conventional samples, that follows from the Mathieu-Simon (MS) theory,[1,2] provided a comprehensive view of the pinning and motion of vortices in the mixed state,[2–5] including some intricate, indeed unexplained, features of the flux-flow regime, in particular the flux-flow noise.[5] We believe that the MS theory and pinning model are quite as relevant to HTSC's. For the present purpose, we only need to recall the MS equations for vortex equilibrium, and how to use them to define a critical state in a soft material. By soft materials, we mean a wide class of samples of standard bulk homogeneity in which critical currents are governed essentially by surface defects. As discussed in Ref. 2, such samples are not exceptional. It has long been observed that the critical current I_c of foils of constant width W and varying thickness t, otherwise prepared in the same way, is nearly constant.[6] The MS model explains how the unavoidable roughness of the sample surface enables a supercurrent to flow close to the surface, over a small depth $\lambda_V(\xi \lesssim \lambda_V \lesssim \lambda)$, up to a superficial critical-current density i_c (in A/m). Granting that $I_c \simeq 2Wi_c$, we account for the order of magnitude of observed I_c, as well as for their H and T dependence.[3] By reporting I_c data in terms of a critical-current density $J_c = I_c/Wt$ (instead of $i_c = I_c/2W$) one obscures this striking and simple result. The same remark can be made about HTSC's. As an example, we shall refer below to critical currents data in $YBa_2Cu_3O_7$ monocrystalline platelets[7] and epitaxial films,[8] when the applied field $\mathbf{H}_0 = \mathbf{B}_0/\mu_0$ is aligned with the c axis, which itself is normal to the sample plane. Typically, at 77 K ($T_c - T \simeq 15$ K) and low fields, $J_c \sim 10^6$ A/cm^2 in films ($t \sim 2000$ Å),[8,9] and $J_c \sim 2\times 10^4$ A/cm^2 in monocrystals ($t \sim 10$ μm).[7,10] By merely inspecting experimental data, it is seen that $J_c \propto 1/t$ in a large range of thickness (2000 Å$\lesssim t \lesssim 200$ μm),[8,10] clearly meaning that, in this case, $i_c \simeq$ const~ 10 A/cm. So we claim that Y-Ba-Cu-O monocrystals are soft in the above restricted sense, and our first concern will be to infer this order of magnitude of i_c *at low fields*. Consistently, the occurrence of the irreversibility line [$I_c = 0$ at some field $H^*(T) < H_{c2}$] will appear as a simple consequence of the strong anisotropy.

In the MS theory, the macroscopic free-energy density $\bar{F}$ of the mixed state, regarded as a continuum, is expressed in terms of a few macroscopic variables. Among them are the magnetic field $\mathbf{B}$, the supercurrent density $\mathbf{J}_s$ (or the superfluid velocity $\mathbf{V}_s$), and the vortex field $\boldsymbol{\omega} = n\varphi_0 \boldsymbol{\nu}$, which describes both the local vortex density n and direction ν (φ_0 is the flux quantum). $\mathbf{B} = \bar{\mathbf{b}}$, $\mathbf{J}_s = \bar{\mathbf{J}}_s$ and $\mathbf{V}_s = \bar{\mathbf{v}}_s$ are the coarse-grained mean values of corresponding microscopic quantities, such as defined in the phenomenological Ginzburg-Landau theory (GL). Microscopic or macroscopic here means on a small or large scale compared with the vortex spacing a. Letting $\Psi = \rho e^{i\theta}$ the order parameter, $\mathbf{a}$ the vector potential, m the electron mass, and $q = -2e$ the effective charge, $2m\mathbf{v}_s = \hbar \nabla \theta - q\mathbf{a}$, and

$$j_{si} = \mu_{ik} \rho^2 q v_{sk}, \tag{1}$$

where μ_{ik} is a symmetric dimensionless tensor. On taking the z axis along the c axis, in an uniaxial crystal [Fig. 1(a)], μ_{ik} is diagonal, and its principal values can be written as[11] $\mu_1 = \mu_2 = \gamma^{2/3}$ and $\mu_3 = \gamma^{-4/3}$, where $\gamma^2 = \mu_1/\mu_3$ is the *anisotropic factor* ($\mu_1\mu_2\mu_3 = 1$). In isotropic materials, $\mu_{ik} = \delta_{ik}$. Averaging curl $\mathbf{v}_s$, and taking account of singularities, one finds the macroscopic London equation,[1] $\mathbf{B} - m/e$ curl $\mathbf{V}_s = \boldsymbol{\omega}$, stating that $\boldsymbol{\omega} \neq \mathbf{B}$ in regions where currents are not curlfree. London and Maxwell equations are constraints limiting the possible spatial variations of currents and fields throughout a vortex lattice. Nevertheless, $\mathbf{B}$, $\mathbf{J}_s$, and $\boldsymbol{\omega}$ must be considered, locally, as independent thermodynamic variables.[2] Whether vortices move or not, MS assume that $\Psi(\mathbf{r},t)$ rigidly satisfies its equilibrium conditions (namely, the first GL equations, including the boundary condition) arguing from the smallness of its relaxation time. In anisotropic crystals, the boundary condition reads:

$$\mu_{ik} \frac{\partial \rho}{\partial x_i} \mathbf{n}_k = 0, \tag{2}$$

which reduces to $\partial \rho / \partial n = 0$ in the isotropic case; $\mathbf{n}$ is the outward normal unit. Under these conditions, the macroscopic thermodynamic identity, obtained by averaging the usual GL expression for the free energy F (or dF), takes the simple form:[12]

$$d\bar{F} = -\bar{\sigma} dT + \frac{1}{\mu_0} \mathbf{B} \cdot d\mathbf{B} - \frac{m}{e} \mathbf{J}_s \cdot d\mathbf{V}_s + \boldsymbol{\varepsilon} \cdot d\boldsymbol{\omega}. \tag{3}$$

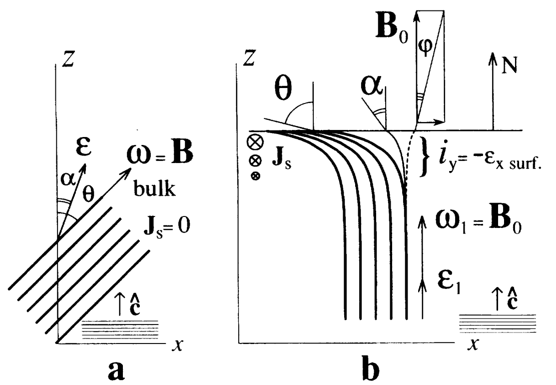

FIG. 1. (a) A uniform vortex array ($\boldsymbol{\omega}=\mathbf{B},\mathbf{J}_s=0$) in an uniaxial crystal, whose c axis has been taken as the z axis. The vortex potential $\boldsymbol{\varepsilon}$, such as defined by Eq. (3) and given by Eq. (4), lags the vortex field $\boldsymbol{\omega}$, when vortices are tilted with respect to the c axis. The lag is governed by the anisotropy factor γ; in the figure $\gamma\simeq1.6$. For $\gamma=7$ (YBaCuO) comparable angles $\alpha\sim20°$ would result from angles θ larger than 85°. (b) As explained in the text, $\boldsymbol{\varepsilon}$ lines (thin line) may end while making a variable angle α with the normal $\mathbf{N}$ to the mean smoothed surface. This entails a strong distortion of the vortex lattice over a small depth from the surface. A supercurrent flows so as to balance the restoring force associated with the vortex bending. Note that magnetic field lines (dashed line) and vortex lines (bold lines) incurve in opposite directions. Typical values (not followed in the figure) for YBaCuO, at $T=77$ K and $B_0=1$ T, are $i_y=10^3$ A/m, $\alpha=10°$, $\theta=83°$, $\varphi\lesssim0.1°$.

The local vortex potential $\boldsymbol{\varepsilon}(\boldsymbol{\omega},T,\mathbf{V}_s)$, here defined as $\partial\bar{F}/\partial\boldsymbol{\omega}$ (A/m, in rationalized units), plays a leading role in the MS model of pinning. Although it is related to the *total reversible* magnetization of an *ideal* sample (see below), by no means should it be regarded as a local magnetization.[2] In many practical situations, the $\mathbf{V}_s$ dependence of $\boldsymbol{\varepsilon}$ may be ignored;[2,12] in a London model, for instance, where the superfluid density ρ^2 is uniform outside the vortex core, $\mathbf{J}_s\propto\mathbf{V}_s$, so that strictly, $\partial\boldsymbol{\varepsilon}/\partial\mathbf{V}_s=0$. Therefore, we can adopt, as a first good approximation, the value $\boldsymbol{\varepsilon}(\boldsymbol{\omega},T)$ relating to a bulk uniform lattice. In this case, where $\mathbf{V}_s\equiv0$, $\mathbf{J}_s\equiv0$, $\boldsymbol{\omega}\equiv\mathbf{B}$, explicit calculations of $\bar{F}(T,\mathbf{B}=\boldsymbol{\omega},\mathbf{V}_s=0,\boldsymbol{\omega})=f(T,\boldsymbol{\omega})$ are available in various approximations, from which we derive $\boldsymbol{\varepsilon}=\partial f/\partial\boldsymbol{\omega}-\boldsymbol{\omega}/\mu_0$ (usually denoted as $-\mathbf{M}$ in the literature).

For vortices making an angle θ with the z or c axis of an uniaxial crystal, so that $\boldsymbol{\omega}(\omega\sin\theta,0,\omega\cos\theta)$, $\boldsymbol{\varepsilon}$ lies in the xz plane [Fig. 1(a)]. Over the whole range of vortex fields ($0<\omega<B_{c2}$), the $\boldsymbol{\varepsilon}$ components can be conveniently written in the compact form:

$$\mu_0\varepsilon_x=\frac{(B_{c2}-\omega)}{2\kappa_x^2}\,g_x\sin\theta,$$
$$\mu_0\varepsilon_z=\frac{(B_{c2}-\omega)}{2\kappa_z^2}\,g_z\cos\theta, \tag{4}$$

where $g_x(\omega)$ and $g_z(\omega)$ are numerical factors. $B_{c2}=B_c\sqrt{2}\kappa(\theta)$ is the θ-dependent upper critical field, with $\kappa(\theta)=\kappa_x[\gamma^2\cos^2\theta+\sin^2\theta]^{-1/2}$, $\kappa_x=\kappa(\pi/2)=\gamma\kappa_z$ and $\kappa_z=\kappa(0)$. In the low-field limit ($\omega\to0$), $\bar{F}=f=\omega H_{c1}(\theta)$,[13] where $H_{c1}=H_{c2}\ln\kappa^*/2\kappa^2$, and $g_x=g_z\simeq\ln\kappa^*=\ln\kappa+0.5$. When the vortex density is increased, ε_x and ε_z fall off and rapidly approach the Abrikosov line. In a London model ($B_{c1}\ll\omega\ll B_{c2}$),[14] $g_x=g_z=\frac{1}{2}\ln\eta B_{c2}/\omega(\eta\sim1)$. Near B_{c2}, but actually in a quite broad domain, $g_x=g_z\simeq$const, and $g=\beta^{-1}(\beta=1.16)$ if $\kappa_z\gg1$.[11] For numerical calculations we used an interpolating formula for $g(\omega)$, which is very close to that deduced from free-energy calculations by Hao *et al.*[15]

It is worth noting that, in a large range of vortex densities (including the limit $\omega=0$, but not a small range $\omega\lesssim B_{c1}$), $g_x=g_z$, so that θ and the angle α between ε and the c axis are simply related:[14]

$$\tan\theta=\gamma^2\tan\alpha, \tag{5}$$

meaning that ε tends to keep aligned with the c axis, or at least ε lags considerably, when θ is increased: if $\gamma=7$ (YBaCuO), $\alpha=10°$ when θ has exceeded 80°.

In equilibrium states, the second GL equation merely states that $\mathbf{j}=\mathbf{j}_s$ as defined by Eq. (1), so that $\mathbf{J}=\mathbf{J}_s=\mathrm{curl}\mathbf{B}/\mu_0$. Furthermore, by minimizing the magnetic free-enthalpy (T and H_0const) against small changes in $\mathbf{J},\mathbf{B}$, and ω distributions (subject to constraints mentioned above), MS derived the local condition for vortex equilibrium,[2] $\mathbf{J}_s=-\mathrm{curl}\boldsymbol{\varepsilon}$, and the boundary condition $\boldsymbol{\varepsilon}\times\mathbf{n}=\mathbf{0}$. On combining these equations with the fundamental equation of state $\varepsilon(\omega,T)$ and Maxwell equations, *one* equilibrium structure of the vortex lattice (at H_0,T) is determined. In simple shaped samples (ellipsoid, cylinder, slab in inclined field, . . .), vortex lines are uniformly distributed in the bulk with constant values $\boldsymbol{\omega}_1=\mathbf{B}_1$ and ε_1; but, in order to satisfy the boundary condition ($\boldsymbol{\varepsilon}\times\mathbf{n}=\mathbf{0}$), $\boldsymbol{\varepsilon}$ lines must bend to terminate normal to the surface. A correlative deformation of the vortex lattice takes place, and associated currents $\mathbf{J}_s=-\mathrm{curl}\boldsymbol{\varepsilon}$ flow, over a small depth λ_V from the surface.[1,2] These currents are nothing but Meissner-like diamagnetic currents, which are responsible for the whole magnetic moment $\mathcal{M}$ of the sample (reversible when changing H_0 or T); $\mathcal{M}_{\mathrm{rev}}\simeq-\boldsymbol{\varepsilon}_1V$, where V is the sample volume.[2] The continuum theory applies, and then the above results hold, on condition that the sample has no defects on the scale of $a\lesssim1000$ Å. But the best soft samples hardly are free of surface roughness on a scale comparable to or smaller than a. Nevertheless, MS suggested that the continuum description could be maintained, while still assuming idealized surfaces, provided that the macroscopic boundary condition was released and replaced by some inequality (like a friction angle condition in mechanics). We must go back to this step in connection with HTSC's.

The condition $\hat{\boldsymbol{\varepsilon}}=\mathbf{n}$ can be interpreted as a requirement about the direction of vortices at the surface. In tensorial form, Eq. (5) reads $\mu_{ik}\varepsilon_k=k\omega_i$, where k is a scalar function. So, $\hat{\boldsymbol{\varepsilon}}=\mathbf{n}$ at the surface implies that vortices end parallel to the vector $\mu_{ik}n_k$. When stated in this form, the boundary condition holds on a microscopic scale: as required by the GL boundary condition (2), a vortex line, i.e., a singular line $\rho=0$, must enter the sample along the direction $\mu_{ik}n_k$. It should also be noted that microscopic currents $\mathbf{j}_s$, round a vortex core, flow in planes normal to ε, whereas $v_s\cdot\boldsymbol{\omega}=0$, in accordance with Eq. (1). In isotropic media, vortices end normal to the boundary. We predicted and observed this effect for superfluid vortices in He II.[16]

In the presence of surface irregularities, there should be many ways for the vortices to terminate on the actual surface, so as to satisfy the microscopic boundary condition (2), implying that currents $\mathbf{j}_s$ flow tangential to the surface. As happens frequently in disordered systems such a circumstance allows for a large number of metastable solutions. According to the MS model, the surface roughness (of a slab in normal field) may induce a strain of the vortex array over a small depth λ_V inside the sample. In view of their strong interaction, vortex lines here cannot be considered individually, and the deformation is collective; but, on a large scale, both the amplitude and orientation of the vortex bending is randomly distributed, with some correlation length c: from recent noise measurements in conventional alloys we may infer that typically $c \gtrsim 1$ μm. More precisely, the vector ε makes with the normal $\langle \mathbf{n} \rangle = \mathbf{N}(0,0,1)$ to the mean smoothed surface a random angle α (brackets mean an average, in the xy plane, on a scale larger than c); the standard deviation α^* (say $\sim 1°$, or $\sim 10°$) should increase with the surface roughness.

According to $\mathbf{J}_s = -\text{curl}\boldsymbol{\varepsilon}$, and seeing that $\hat{\boldsymbol{\varepsilon}}_1 = \mathbf{N}$ in the bulk, the vortex curvature requires a macroscopic surface current $\mathbf{i} = \boldsymbol{\varepsilon}_{\text{surf}} \times \mathbf{N}$ $[i = \varepsilon(\alpha)\sin\alpha]$. Now, the statistical distribution of $\mathbf{i}$ can be more or less "polarized" so that a net current flows. For example, by systematically bending vortex and $\boldsymbol{\varepsilon}$ lines in the negative x direction [Fig. 1(b)], a net current will be transported in the y direction: $0 \leqslant \langle i_y \rangle \leqslant i_c$. The critical current density i_c will be estimated as the value of i for some mean representative values $\alpha_c \sim \alpha^*$:

$$i_c = \varepsilon(\alpha_c)\sin\alpha_c = |\varepsilon_x(\alpha_c)|. \qquad (6)$$

By constraining $\boldsymbol{\varepsilon}$ lines to make an angle $\alpha \sim \alpha_c$ with the z axis, an irregular surface should also give rise to a correlative bending for the vortices near the surface, up to a large angle $\theta \sim \theta_c$ [see Eq. (5)]. Then, vortex bending entails an increasing vortex density: $\omega = \omega_1/\cos\theta$, as required by the conservation of the number of vortex lines ($\text{div}\boldsymbol{\omega} = 0$). Here, $\omega_1 = B_0$ is the bulk vortex field. Given a value of α_c, $\theta_c = \arctan(\gamma^2 \tan\alpha_c)$, and substituting $\omega(\theta_c) = \omega_1/\cos\theta_c$ in the equation of state $\boldsymbol{\varepsilon}(\boldsymbol{\omega},T)$, we obtain i_c through Eq. (6). Conversely, from experimental data, we can estimate the critical angle α_c; for example, from the irreversible magnetization curve of a platelet (radius R, thickness t) in a normal field B_0: $\mathcal{M} = \mathcal{M}_{\text{rev}} \pm \frac{2}{3}\pi R^3 i_c$, where $\mathcal{M}_{\text{rev}} = \pi R^2 t \varepsilon_1$. Note that the reversible magnetization and its hysteretic deviations are both governed by the vortex potential $\boldsymbol{\varepsilon}$.

In conventional isotropic samples (lead alloys, pure Nb)[3,12] we found values of α_c ranging from a few degrees to 20°, depending on the surface treatment. Since there $\theta_c = \alpha_c$, ω and ε keep close to their bulk values, so that $i_c \simeq \varepsilon_1 \sin\theta_c$. It is obvious that any pinning model involves at least one adjustable parameter. However, we had to find values of $\sin\alpha_c$ reasonably smaller than unity. In this sense, the MS model is able to predict orders of magnitude of critical currents. In lead alloys, we observed that, at different temperatures (1.5 to 4.2 K) and fields ($0.8H_{c2} \lesssim H_0$), α_c was nearly constant.[3] This means that, in this range, α_c was a property of the rough surface, and that the (H_0,T) dependence of I_c was that of ε_1.

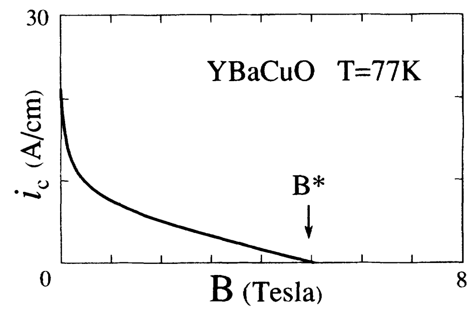

FIG. 2. Theoretical estimate of the critical current density i_c of an YBaCuO surface ($\gamma = 7$, $\kappa_z = 20$), at $T \simeq 77$ K ($B_{c2z} = 8$ T), in normal applied field B, as shown in Fig. 1(b). The critical angle α_c in Eq. (6), which characterizes the surface roughness, has been taken as 10°, a typical mean value suggested by previous experimental results in conventional soft materials ($1 < \alpha < 20°$).

Let us return to HTSC's. As an example, let us consider YBaCuO at $T = T_c - 15$ K$\simeq 77$ K. Reported critical field slopes $dH_{c2}/dT|_{T_c}$ (0.5 − 0.55 T/K$\|c$) yield typical anisotropy ratios of 6–8.[9,17] Reported values of κ_z are much more scattered, from $\kappa_z \lesssim 10$,[17] to $\kappa_z \lesssim 60$.[15] From $B_c(0) \sim 1$ T, one finds $\kappa_z \lesssim 20$ in the dirty limit. For definiteness we have picked out $\gamma = 7$, $\kappa_z = 20$, $B_{c2}(\theta = 0, 77\text{ K}) = 8$ T. Then taking $\alpha_c = 10°$ as a representative value over the whole field range, $i_c(B_0)$ has been calculated from Eqs. (4)–(6) and $\omega = \omega_1/\cos\theta$. Results are shown in Fig. 2 and could speak for themselves; the low-field limit, $i_c \sim 10$ A/cm, as also the field at which critical currents vanish, $B^* \simeq 5$ T$\simeq 0.6B_{c2}$, are in fair agreement with standard data from Y-Ba-Cu-O single crystals.[8,9,18]

In calculating i_c as $\varepsilon_x(\alpha_c)$ in Eq. (6), we make use of the first Eq. (4), where $\theta = \theta_c$, $\sin\theta_c \simeq 1$, and $\omega = \omega_1/\cos\theta_c \gg \omega_1$. As a practical estimate of maximum critical current densities at low fields, we may take the extrapolated value at $\omega = 0$ of the Abrikosov line:

$$i_c = \varepsilon_x(\alpha_c) \lesssim \frac{H_{c2}(\pi/2)}{2\kappa_x^2} < H_{c1}(\pi/2) < H_{c1}(0), \qquad (7)$$

For instance, with data of Fig. 2, Eq. (7) reads $i_c \lesssim 11$ A/cm, $H_{c1}(\pi/2) = 62$ A/cm (78 Oe), $H_{c1}(0) = 280$ A/cm (350 Oe). In isotropic materials, $\omega_{\text{surf}} \simeq \omega_1$, and we find similarly $i_c \lesssim H_{c1}\sin\alpha_c < H_{c1}$. Thus, in either case, H_{c1} (A/m) represents a wide upper bound for maximum surface critical currents. This explains why conventional and HTSC's materials finally have comparable performances, as far as critical currents are concerned. The occurrence of large J_c in hard samples should not alter this general remark; the introduction of many bulk inhomogeneities (cavities, precipitates, sintered powders) more or less amounts to increasing artificially surface effects by multiplying interfaces.

The vanishing of critical currents over a rather large field range below H_{c2z}, appears here as a simple consequence of the strong increase in vortex density near the surface. On stating that the surface vortex density, $\omega = B_0/\cos\theta_c$, has reached the upper critical value $B_{c2}(\theta=\theta_c)$, so that $\varepsilon\to 0$, we obtain a simple expression for B^*:

$$B^*=\frac{B_{c2z}}{(1+\gamma^2\tan^2\alpha_c)^{1/2}} . \qquad (8)$$

Note that, in any case, B^* is strictly smaller than B_{c2z}. In isotropic materials, however, this effect is hardly significant: for $\gamma=1$ and $\alpha_c\sim 10°$, $B^*\simeq B_{c2}$ within better than 2%. Taking $\gamma=7$ and $\alpha_c=10°$, we find $B^*/B_{c2z}=0.63$ (Fig. 2), in agreement with a number of reported results in YBaCuO.[8,9,18] Otherwise stated, B^* data yield, through Eq. (8), a reasonable value of the critical angle α_c; their fair reproducibility should not be surprising, insofar as crystals are made by using few standard methods. The ratio B^*/B_{c2z} decreases rapidly with increasing anisotropy, as observed in BiSrCaCuO;[18] taking, for example, $\gamma=50$ and $\alpha_c=10°$, $B^*/B_{c2z}=0{,}11$. According to the naive mean picture of Fig. 1(b), with $\alpha=\alpha_c=\text{const}$, a normal sheath develops at the surface above B^*; actually, an intricate distribution of growing normal spots should populate the surface at high fields. As far as nondissipative surface currents $\langle \mathbf{i}\rangle$, macroscopic on the sample scale, are concerned, a rough surface above B^* behaves as an ideal surface: $\langle \mathbf{i}\rangle=-\boldsymbol{\varepsilon}_1\times\mathbf{N}$, that is zero in normal field, is uniquely determined, so that magnetization becomes reversible. However, the response of a rough surface, in the reversible region, should be very different in general from that of an ideal surface, when the vortex array is subject to various probes or disturbances (surface impedance, flux flow noise).

In conclusion, we wish to point out some marked differences between the MS theory and the classical elastic continuum theory (EC). Besides the fact that it only deals with small deformations of the vortex lattice, the EC theory is unable to account for equilibrium configurations such as that shown in Fig. 1(b), where vortex lines and magnetic field lines separate ($\boldsymbol{\omega}\neq\mathbf{B}$). In fact, this question arises independently of any pinning model or anisotropy effect; near the surface of an ideal isotropic sphere at equilibrium, vortices and fields lines must bend in opposite directions.[2] Taking the vector product of $\mathbf{J}_s=-\text{curl}\boldsymbol{\varepsilon}$ by $\varphi_o\nu$, this can be interpreted as the equilibrium between the Lorentz force $\mathbf{J}_s\times\varphi_0$ and the restoring force $\text{curl}\boldsymbol{\varepsilon}\times\varphi_0$ associated with the vortex bending. Deformations such as described by the EC theory wrongly presuppose that $\boldsymbol{\omega}\equiv\mathbf{B}$, so that both forces, acting in the same direction, never can counterbalance one another. Thus, in order to account for critical currents, the EC theory must have recourse to the small shear modulus c_{66}; if $c_{66}=0$, pinning should be ineffective. Whence, the opinion widely held, that the vortex lattice melting along a phase transition line $B^*(T)$ would explain the vanishing of critical currents. Instead, by using only one parameter $\boldsymbol{\omega}$ to describe the vortex state, we have deliberately left out small differences in energy between, for example, triangular and square lattices, and thereby, we ignore shear stresses.

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