Fragmentation of Magnetic Flux in Anisotropic Superconductors

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It is shown that in anisotropic superconductors uniform current flow can be unstable for current densities $j_{f1} < j < j_{f2}$, where j_{f1} and j_{f2} essentially depend on the orientation of **j**. The instability leads to a stratification of the magnetic flux into macroscopic domains (macrovortices). Densities of closed shielding currents circulating inside the macrovortices can be well above the mean density of the transport current. The relation of such a macrovortex structure to the magnetic granularity of high-T_c oxides is discussed. The anisotropy is shown to result in the instability of both the Bean critical-state model and the uniform vortex glass with current-voltage characteristic of the form $V \propto \exp(-\text{const}/j^{\mu})$.

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Properties of high-T_c superconductors are highly anisotropic because of their layered crystalline structure and the existence of specific planar defects (twins, stacking faults, grain boundaries, etc.). The anisotropy results in a noncollinearity of the magnetic induction **B** and the magnetic field **H**,[1] and low values and anisotropy of tilt and shear elastic moduli of the vortex lattice,[2,3] which considerably reduces the stability of the mixed state with respect to the entanglement and reconnection of vortices and the creation of dislocations and other defects.[4-6] This may lead to novel vortex structures such as liquid phases,[3-5] vortex glass,[4,6] two-dimensional (2D) vortices localized at the CuO_2 planes,[7] etc.

On the other hand, the effect of the anisotropy on the absolute stability of both equilibrium and dissipative vortex structures has not been studied yet. This effect can prove to be of great importance especially for nonlinear resistive states, where vortex structures are always metastable and the anisotropy essentially influences the electrodynamics of a superconductor. In this paper I consider the macroscopic electromagnetic stability of the resistive states in anisotropic superconductors, provided that the details of the vortex structure are inessential. It is shown that the uniform current flow becomes unstable in the case of strong nonlinearity of the current-voltage (I-V) characteristic typical for the flux-creep regime, the Bean model with anisotropic critical current density $j_c(\Phi)$,[8] and the vortex-glass state.[6] The coupling of anisotropy and nonlinearity of $I(V)$ leads to a collective electromagnetic instability of the macroscopically uniform mixed state and its transition into a cellular structure which can be considered as an array of magnetic macrovortices in the system of the Abrikosov vortices.

The stability criterion of the uniform resistive state with respect to small perturbations $\delta\mathbf{B}(\mathbf{r},t)$ and $\delta\mathbf{E}(\mathbf{r},t)$ can be obtained using the Maxwell equations $\partial\delta\mathbf{B}/\partial t = -\text{curl}\,\delta\mathbf{E}$, $\text{curl}\,\delta\mathbf{B} = \mu_0\delta\mathbf{j}$, where $\delta E_\alpha = R_{\alpha\beta}\delta j_\beta$, $R_{\alpha\beta}(\mathbf{j}) = \partial E_\alpha/\partial j_\beta$ is the tensor of the differential resistivity, E_α and j_β are the components of the electric field and current density, respectively, $\alpha,\beta = (x,y,z)$, the orthogonal axes x,y,z are assumed to coincide with the symmetry axes, and $\mu_0\mathbf{H} = \mathbf{B}$, which corresponds to $H \gg H_{c1}$, with H_{c1} the lower critical field. Let, for simplicity, **B** be parallel to the z axis and let the current flow along the x-y plane, then the elements R_{xz}, R_{zx}, R_{yz}, and R_{zy} vanish due to symmetry. Assuming $\delta\mathbf{E}, \delta\mathbf{H} \propto \exp(\lambda t + i\mathbf{k}\cdot\mathbf{r})$ and neglecting the self-field effects, one gets a set of linear equations for the Fourier components: $\lambda\delta\mathbf{B}(\mathbf{k}) = -i[\mathbf{k}\times\delta\mathbf{E}(\mathbf{k})]$, $i[\mathbf{k}\times\delta\mathbf{B}(\mathbf{k})] = \mu_0\delta\mathbf{j}(\mathbf{k})$, $\delta E_\alpha(\mathbf{k}) = R_{\alpha\beta} \times \delta j_\beta(\mathbf{k})$; the condition of their solvability yields the spectrum of the electromagnetic perturbations $\lambda(\mathbf{k})$. After some algebra, one finds that $\lambda(\mathbf{k}) = -\mu_0^{-1} f(\mathbf{n}) k^2$, where $\mathbf{n} = \mathbf{k}/k$, $k = |\mathbf{k}|$, and the function f obeys the following quadratic equation:

$$f^2 - [(n_y^2+n_z^2)R_{xx} + (n_x^2+n_z^2)R_{yy} + (n_x^2+n_y^2)R_{zz} - n_x n_y(R_{xy}+R_{yx})]f + n_z^2(R_{xx}R_{yy} - R_{xy}R_{yx}) + [n_y^2 R_{xx} + n_x^2 R_{yy} - n_x n_y(R_{xy}+R_{yx})]R_{zz} = 0. \quad (1)$$

The solutions of Eq. (1) describe two modes with the increments $\lambda(\mathbf{k})$ depending on the orientation of **k**. Since the instability arises if $\text{Re}\lambda(\mathbf{k}) > 0$, let us consider the most "dangerous" mode with the minimum value of $\text{Re}f(\mathbf{n})$. As follows from Eq. (1), at $R_{\alpha\alpha} > 0$ such a mode corresponds to $n_z = 0$ and to the minimum of the function

$$Y = n_y^2 R_{xx} + n_x^2 R_{yy} - n_x n_y(R_{xy}+R_{yx}).$$

The mode becomes unstable if $Y(\Psi) < 0$ and $\partial Y/\partial\Psi = 0$, where $n_y = \cos\Psi$, $n_x = \sin\Psi$. These conditions give the instability criterion in the form

$$4R_{xx}R_{yy} < (R_{xy}+R_{yx})^2, \quad (2)$$

$$\tan 2\Psi = (R_{xy}+R_{yx})/(R_{yy}-R_{xx}). \quad (3)$$

Here Eq. (3) fixes the orthogonal principal axes **n** and **m** of $R_{\alpha\beta}$ which generally do not coincide with the symmetry axes in anisotropic nonlinear media, where **n** depends on the orientation of **j** as well. Equation (2) implies that the principal value of $R_{\alpha\beta}$ for the axis **m** is negative, which indicates a 2D instability in the x-y plane (Fig. 1).

In anisotropic superconductors Eq. (2) can hold in a wide region of j. In order to show that, let us present the I-V characteristic as follows:

$$E_\alpha = G(j,\Phi,\theta)\rho_{\alpha\beta}j_\beta , \quad (4)$$

where Φ and θ are the angles between **j** and the principal axes x and z of the resistivity tensor $\rho_{\alpha\beta}$, $j=|\mathbf{j}|$, and $G(j,\Phi,\theta)$ is a nonlinear function of j, Φ, and θ (see below). Equation (4) takes into account the noncollinearity of **E** and **j** and the angular dependence of the mean density of unpinned vortices $n_f(\mathbf{j},\mathbf{B}) \propto G$ contributing to **E**.[9] Such an approach enables one to conclude about the stability, regardless of the origin of the resitivity and anisotropy. From Eq. (4), one finds

$$\begin{aligned} R_{xx} &= (G + j\cos^2\Phi\,\partial G/\partial j - \cos\Phi\sin\Phi\,\partial G/\partial\Phi)\rho_x , \\ R_{xy} &= (j\cos\Phi\sin\Phi\,\partial G/\partial j + \cos^2\Phi\,\partial G/\partial\Phi)\rho_x , \\ R_{yx} &= (j\cos\Phi\sin\Phi\,\partial G/\partial j - \sin^2\Phi\,\partial G/\partial\Phi)\rho_y , \\ R_{yy} &= (G + j\sin^2\Phi\,\partial G/\partial j + \cos\Phi\sin\Phi\,\partial G/\partial\Phi)\rho_y , \end{aligned} \quad (5)$$

where ρ_x and ρ_y are the principal values of $\rho_{\alpha\beta}$. The explicit criterion of the instability can be obtained by means of Eqs. (2) and (5), if one assumes the scaling $G(j,\Phi) = G(j/j_k(\Phi))$. Then Eq. (2) reduces to the quadratic inequality for $\partial G/\partial j$, which yields the instability criterion in the form

$$s(j) > s_c = g/2 + (g^2/4 + g)^{1/2} , \quad (6)$$

$$g(\Phi) = \frac{16\rho_x\rho_y}{[b(\rho_x+\rho_y) + (\rho_x-\rho_y)(b\cos 2\Phi - \sin 2\Phi)]^2} , \quad (7)$$

where $s(j) = \partial \ln G/\partial \ln j > 0$ and $b(\Phi) = \partial \ln j_k(\Phi)/\partial\Phi$. In the isotropic case ($\rho_x = \rho_y$, $b=0$, $g=\infty$) the uniform current is stable; otherwise Eq. (6) can hold if the anisotropy and the nonlinearity of $I(V)$ are high enough.

Let us now consider the stability of the characteristic resistive states, making use of Eq. (6). For instance, in the conventional flux-creep model[10] the quantity $\rho_{\alpha\beta}$ is the resistivity tensor in the thermally assisted flux-flow (TAFF) regime,[11] $G = (j_1/j)\sinh(j/j_1)$, and $j_k = j_1$, where $j_1(\Phi) = k_B T j_c(\Phi)/U(\Phi)$ and $U(\Phi)$ is an activation energy. Then $s = (j/j_1)\coth(j/j_1) - 1$ and the instability arises at $j > j_{f1}$ with

$$j_{f1} = \begin{cases} g j_1, & g \gg 1 , \quad (8) \\ \sqrt{3} g^{1/4} j_1, & g \ll 1 . \quad (9) \end{cases}$$

In the critical-state model one has $G = 1 - j_c(\Phi)/j$, $j_k = j_c$, and $\rho_{\alpha\beta}$ is the flux-flow resistivity tensor, whence

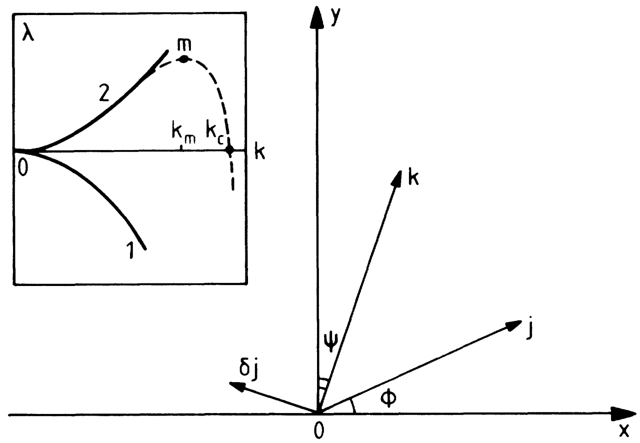

FIG. 1. The geometry of the instability. Inset: The dependence of λ on k: curve 1, $s < s_c$; curve 2, $s > s_c$. The dashed curve shows the change of $\lambda(k)$ due to the suppression of the short-wave instability (see text).

$s = j_c/(j - j_c)$ and Eq. (6) reduces to $j < j_{f2}$, where

$$j_{f2} = [(\tfrac{1}{4} + 1/g)^{1/2} + \tfrac{1}{2}] j_c . \quad (10)$$

A similar situation occurs for the I-V characteristic $V \propto \exp[-(j_0/j)^\mu]$, $\mu > 0$, predicted for the vortex-glass model.[6,12] Then $s(j) = \mu(j_0/j)^\mu$, so at $j \to 0$ the instability criterion $j < j_g = j_0(\mu/s_c)^{1/\mu}$ holds for any anisotropy, which indicates a transformation of the glassy state and the change of $V(I)$ at $j < j_g$. For $G \propto j^m$, $m > 0$, the stability criterion $m < s_c$ is weaker; nevertheless, for extremely high anisotropy ($s_c \to 0$) only the TAFF regime ($m=0$) may prove to be stable. Thus, the instability arises at $j_{f1}(\Phi) < j < j_{f2}(\Phi)$, if $s_c(\Phi) < s_m = \max[s(j)]$. At $s_c(\Phi) \to s_m$, one gets $j_{f1} \to j_{f2}$, but for $g \sim 1$, Eqs. (8)-(10) yield $j_{f2} - j_{f1} \sim j_c$. In the latter case the values j_{f2} and j_{f1} can be found, regardless of the crossover region between the flux-creep and flux-flow regimes, which was assumed above when deriving Eqs. (8)-(10).[13]

Now we examine the mechanisms of the anisotropy in more detail. The angular dependence of $j_k(\Phi)$ is determined by the crystalline symmetry in the x-y plane, which yields $b(\Phi) \propto \partial j_k/\partial\Phi = 0$ at $\Phi = \pi n/2$, $n = 0, 1, 2, \ldots$. This implies $g(\Phi) = \infty$ for $\mathbf{j} \parallel \hat{\mathbf{x}}$ or $\mathbf{j} \parallel \hat{\mathbf{y}}$, that is, the uniform current flow is stable. However, for other orientations of **j** the parameter $g(\Phi)$ is finite and the instability arises if $j_{f1}(\Phi) \lesssim j_c$, i.e., $s_c(\Phi) \lesssim j_c/j_1 \sim U/k_B T$. For instance, at $\rho_x \gg \rho_y$, one finds from Eqs. (7) and (9)

$$j_{f1} = \frac{\sqrt{6}(\rho_y/\rho_x)^{1/4} j_1(\Phi)}{|\cos\Phi[\sin\Phi - b(\Phi)\cos\Phi]|^{1/2}} , \quad (11)$$

where the function $b(\Phi)$ can be expanded into a Fourier series as $b(\Phi) = \sum_n b_n \sin(2n\Phi)$. We examine the instability region $j_{f1}(\Phi) \lesssim j_c$, taking into account, for simplicity, only the first harmonic b_1. If $b_1 < \frac{1}{2}$, there exist four stability sectors with the angular width $\delta\Phi \sim (\rho_y/\rho_x)^{1/2}(j_1/j_c)^2 \ll 1$, and for $b_1 > \frac{1}{2}$, additional sectors appear in the vicinity of $\Phi = \pm \arccos(2b_1)^{-1/2}$ (Fig. 2). The case $\rho_x \gg \rho_y$ is typical for high-T_c oxides at $\mathbf{B} \perp \mathbf{c}$ with $\rho_a \sim 0.1\rho_c$ for $YBa_2Cu_3O_x$, $\rho_a \sim (10^{-4}\text{-}10^{-5})\rho_c$ for Bi- and Tl-based compounds, and $j_c(0)/j_c(\pi/2) \sim b \sim 10\text{-}10^3$ (Refs. 14–17) ($x = c$, $y = a$). Then Eq. (11) yields $j_{f1} \simeq (0.1\text{-}0.01)j_c \ll j_c$, which implies that at $\rho_x \gg \rho_y$ the instability can be caused by any weak nonlinearity of $\mathbf{E}(j)$ with $s > 0$. The anisotropy in the a-b plane can be due to the Cu-O chains, twins, and grain boundaries as well. Consider, for example, a stack of planar crystalline defects in an isotropic superconductor, provided that they act as additional pinning centers and give a small contribution to the macroscopic resistivity. Then $\rho_x = \rho_y$, but the macroscopic values of $j_1(\Phi)$ and $j_c(\Phi)$ are anisotropic, which can lead to the instability, since the parameter $g(\Phi) = 4b^{-2}(\Phi)$ is finite. The stability diagram is similar to that shown in Fig. 2(a), with $\delta\Phi \sim \Phi_c$ and $b(\Phi_c) \simeq 6(j_1/j_c)^2 \ll 1$, i.e., $\Phi_c \sim (j_1/j_c)^2 \sim (k_B T/U)^2$. An analogous situation could also arise in superconducting multilayered structures.[18]

Within the framework of the macroscopic approach the instability has an explosive character, since at $s > s_c$ [$\mathrm{Re} f(\mathbf{n}) < 0$] the maximum increment $\lambda(\mathbf{k}) = -\mu_0^{-1} k^2 \times f(\mathbf{n})$ corresponds to $k \to \infty$. This implies that at large k one should take account of nonlocal effects in Eq. (4) caused by additional mechanisms suppressing the shortwave instability at some $k > k_c$. For example, such a mechanism might be the diffusion of heat or nonequilibrium quasiparticles accompanying the vortex motion. The wavelength k_c^{-1} could also be limited by the sizes of crystallographic grains. Indeed, for the above example of planar defects, the macroscopic instability arises only at $kd \ll 1$, where d is a spacing between the defects. At $kd \gg 1$ the electromagnetic perturbations are mainly localized between the defects, where the superconducting state is stable, whence $k_c \sim d^{-1}$. In any case, the length k_c^{-1} cannot be less than the vortex bundle size[10] (pinning correlation length[19]) which is a physical scale on which the vortex structure is always stable. All these

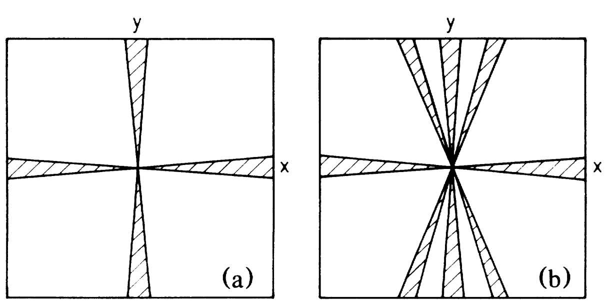

FIG. 2. The diagram of stability at (a) $b_1 < \frac{1}{2}$ and (b) $b_1 > \frac{1}{2}$. The hatched sectors correspond to the directions for which the uniform current flow is stable.

mechanisms result in a maximum of $\lambda(\mathbf{k})$ at some $k = k_m \sim k_c$ (Fig. 1).

Let us now outline the structure of the resistive state above the instability threshold $s > s_c$, where one of the principal values of the magnetic diffusivity tensor $D_{\alpha\beta} \propto R_{\alpha\beta}^{-1}$ becomes negative. This leads to the growth of inhomogeneous magnetic structure and a stratification of the current into domains with different values of $\mathbf{j}$ by analogy to the Gunn effect in semiconductors[20] or the spinodal decomposition of solids.[21] The instability induces transverse components to the initial $\mathbf{j}$ (see Fig. 1), resulting in the appearance of a 2D static perturbation $\delta\mathbf{j}(\mathbf{r})$ of the form $\sum_{\mathbf{k}} \delta\mathbf{j}_{\mathbf{k}} \cos(\mathbf{k}\cdot\mathbf{r} + \gamma_{\mathbf{k}})$, $0 < k \lesssim k_c$, with the amplitudes $\delta\mathbf{j}_{\mathbf{k}}$ and phases $\gamma_{\mathbf{k}}$ determined by the nonlinearity of $I(V)$ and the form of $\lambda(\mathbf{k})$.[22] Such a structure is generally nonperiodic, and its characteristic spatial scale is of the order of k_m^{-1}, since the main contribution to $\delta\mathbf{j}(\mathbf{r})$ comes from the harmonics $\delta\mathbf{j}_{\mathbf{k}}$ with $k \simeq k_m$ corresponding to the maximum increment of $\lambda(\mathbf{k})$ (see, e.g., Ref. 21). When a characteristic amplitude j_s of $\delta\mathbf{j}(\mathbf{r})$ becomes of the order of the mean transport current density j_t, partial closure of the current lines arises, which results in the creation of anisotropic current loops of size $\sim k_m^{-1}$. Such a state can be considered as a macrovortex structure,[23] where j_s plays the role of a characteristic current density circulating inside the macrovortices. The value j_s may be estimated, assuming the sample is separated into domains with $j \lesssim j_{f1}$ and $j \gtrsim j_{f2}$, their concentrations depending on j_t. This results in a suppression of the instability inside the domains, where $s < s_c$; hence $j_s \sim j_{f2} - j_{f1}$.

The amplitude j_s essentially depends on Φ. Indeed, if $\mathbf{j}_t$ is nearly parallel to one of the boundaries of the hatched sectors in Fig. 2, the instability leads to weak modulations of $\mathbf{j}(\mathbf{r})$ only, since $s \to s_c$ and $j_s \sim j_{f2}(\Phi) - j_{f1}(\Phi) \ll j_t$. However, an increase of the misalignment between $\mathbf{j}_t$ and one of the principal axes x or y results in the growth of $j_s(\Phi)$ and the appearance of the macrovortex structure in the case $j_s \gtrsim j_t$, which can easily hold at $g \sim 1$, where $j_t \sim j_{f1}(\Phi) \ll j_c$ and $j_s \sim j_{f2}(\Phi) - j_{f1}(\Phi) \sim j_c$.

Let us discuss possible manifestations of this instability. As mentioned above, in spite of the macroscopic origin of the instability, the structure of the nonuniform state at $s > s_c$ is determined by both nonlinear electromagnetic effects and details of pinning or crystalline lattice. For instance, the angular dependence of the intergrain $j_c(\Phi)$ may turn the crystallographic grains into magnetic macrovortices with high shielding current densities of the order of an intragrain j_c. This transition occurs at $j_t \sim j_{f1}$, where $j_{f1} \ll j_c$ in the case of high anisotropy. The instability may also be caused by a high crystalline anisotropy itself without invoking defects, especially in Bi- and Tl-based compounds, where the sizes of macrovortices could be determined by intrinsic pinning mechanisms. A similar magnetic granularity and a considerable difference between the transport (j_{tc}) and the shielding (j_{sc}) critical current densities in nonceramic high-T_c oxides were discussed before.[24,25] Notice that the instability can be suppressed upon decreasing the anisotropy by material inhomogeneities or, say, by the irradiation of the sample resulting in a partial amorphization of the crystalline lattice. This might clarify the fact that the difference between j_{tc} and j_{sc} was observed to be more pronounced for single crystals than for more inhomogeneous oriented-grain materials.[24] The appearance of the macrovortex structure can be considered as a dissipative phase transition from the laminar current flow into a peculiar "turbulent" regime, which could manifest itself in peaks of the magnetic susceptibility, the electric noise power, and the mechanical damping in vibrating-reed experiments,[26] or as a nonuniform distribution of the magnetic flux in magneto-optical experiments.[27] Such a transition may also lead to anomalous dependences of kinetic coefficients on T and B and novel mechanisms of nonlinearity of $V(I)$.

Finally, the anisotropy can restrict possible forms of the I-V characteristics for which uniform current flow is stable. If the anisotropy and/or the nonlinearity of $I(V)$ is high enough, macroscopic modulations of the magnetization arise, regardless of the structure and dynamics of the mixed state. Similar instabilities could also arise in other nonlinear anisotropic systems.

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