Low-angle resistivity anomaly in layered superconductors

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The pinning effect of vortex lines by the layered structure (intrinsic pinning) on the resistivity of high-T_c superconductors in the mixed state is investigated by means of perturbation theory. A sharp drop in the resistivity at small angles for which vortex lines are almost aligned with the *ab* planes is shown to occur even in a high-temperature region where the pinning potential is reduced by thermal fluctuations.

The phenomenon of intrinsic pinning in high-T_c superconductors (HTSC's), i.e., the confinement of vortex lines between Cu-O planes, has recently attracted a great deal of attention. The possibility of intrinsic pinning, which is one of the most appealing manifestations of the layered structure of the copper oxide superconductors, was suggested by Tachiki and Takahashi[1] and Barone, Larkin, and Ovchinnikov.[2] They pointed out that since the coherence length along the *c* direction, ξ_c, is comparable to the interplane distance *s*, it is more favorable for vortex lines to locate between the Cu-O planes in order to minimize their core energy. This implies that vortices parallel to the *ab* planes experience a periodic potential with maxima corresponding to positions of the cores of the vortex lines at the Cu-O planes. The amplitude of the periodic pinning potential is determined by the ratio ξ_c/s. For $\xi_c/s \ll 1$ the magnitude of the pinning potential is nearly as large as the condensation energy $(H_c^2/8\pi)s\xi_{ab}$, whereas when $\xi_c/s \gg 1$, the amplitude of the periodic pinning potential drops very rapidly with ξ_c/s [for example,[2] as $\exp(-\Gamma\xi_c/s)$]. The behavior of the intrinsic periodic potential for a vortex lattice near the upper critical field H_{c2} has been investigated in detail in Ref. 3 within the framework of mean-field theory.

The most direct way to investigate intrinsic pinning is to perform resistive measurements at small misalignments of magnetic field with respect to the Cu-O planes. Strong evidence of the intrinsic pinning was given in experiments by Refs. 4 and 5, where sharp increases in the critical current were observed in Y-Ba-Cu-O thin films for $H \parallel ab$ planes. Very precise angular measurements of the linear resistivity have been done for Y-Ba-Cu-O single crystals[6] and for thin films of Bi-Sr-Ca-Cu-O.[7] In both compounds below some typical temperature a very sharp drop in resistivity appears at very small angles θ between *H* and the *ab* planes, where $\theta < 0.5°$. This anomaly becomes more pronounced with decreasing temperature.

Two different temperature regimes of vortex motion through a periodic pinning potential can be distinguished. At low temperatures (the exact criterion will be specified below) the motion of the vortex lines along the *c* axis is completely suppressed and the onset of resistivity at finite misalignment is expected to be due to a lock-in transition, i.e., the multiple creation of kinks in the vortex lines at angles above some critical angle θ_0.[8] These kinks can move parallel to the *ab* plane and dissipate energy, giving rise to finite resistivity. At high temperatures the periodic potential is suppressed by thermal fluctuations of the vortex lines, and is to be considered rather as a small perturbation for the vortex motion. Note that the resistive measurements by Kwok *et al.*[6] were done mainly in this high-temperature region which is the subject of this paper. We show that even a small periodic potential, being treated within perturbation theory, can produce a noticeable anomaly in the angular dependence of the resistivity.

We consider the correction $\delta\rho$ to the resistivity ρ arising from the intrinsic pinning in the regime of ohmic resistivity: $\rho = E/J$, where $\mathbf{E}$ is the electric field induced by the vortex motion with velocity $\mathbf{v}$ and $\mathbf{E} = (1/c)[\mathbf{B}\times\mathbf{v}]$ in the presence of the transport current $\mathbf{J}$. The Ohmic resistivity reflects the viscous motion of the vortex lines. In the absence of pinning, the vortex velocity is determined by the viscous friction η:$v = BJ/c\eta$. In order to find $\delta\rho$ one has to calculate the correction to vortex velocity $\mathbf{v}$ due to the intrinsic pinning potential. We assume the pinning potential to be small; therefore the most adequate method is the dynamic approach introduced in Refs. 9 and 10 and explored successfully in investigations of vortex dynamics in HTSC's. An additional pinning force per unit volume due to a pinning potential $V(\mathbf{r})$ is of the form[11,12]

$$\mathbf{F}_{\text{pin}} = -\langle n(\mathbf{r},t)\nabla V(\mathbf{r}-\mathbf{v}t-\mathbf{u})\rangle , \quad (1)$$

where $n(\mathbf{r},t) = \sum_i \delta(\mathbf{r}_\perp - \mathbf{R}_{0i}(z,t))$ is the vortex density, and brackets denote averaging, both spatial and thermodynamic.

$\mathbf{R}_{0i}(z,t)$ are the positions of the vortex lines (which are assumed to be parallel to *z* axis) in the absence of pinning, but accounting for their thermal fluctuations, $\mathbf{v}$ is the mean velocity of vortex configuration, and $\mathbf{u}(\mathbf{r},t)$ is an additional distortion field due to the pinning potential. Since the displacement field $\mathbf{u}$ varies slowly from vortex

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to vortex, it can be treated as a continuous function. Within the linear approximation one has

$$u_\alpha(\mathbf{r},t)=\int d^3\mathbf{r}'dt'G_{\alpha\beta}(\mathbf{r}-\mathbf{r}';t-t')F_{\text{pin}\beta}(\mathbf{r}',t')\ , \qquad (2)$$

where $\alpha=x,y$ and $G_{\alpha\beta}(\mathbf{r},t)$ is the response function for the vortex configuration. Expanding the expression for the friction force (1) with respect to small distortions $\mathbf{u}(\mathbf{r},t)$ and using (2) one obtains

$$F_{\text{pin}\gamma}(\mathbf{v})=-\int d^3\mathbf{r}'dt'G_{\alpha\beta}(\mathbf{r}-\mathbf{r}',t-t')S(\mathbf{r}-\mathbf{r}',t-t') \times\langle\nabla_\alpha\nabla_\gamma V(\mathbf{r}-\mathbf{v}t)\nabla'_\beta V(\mathbf{r}'-\mathbf{v}t')\rangle\ . \qquad (3)$$

The structure factor is given by

$$S(\mathbf{r},t)=\langle n(\mathbf{r},t)n(0,0)\rangle\ . \qquad (4)$$

We choose the z axis to lie along the vortex lines and to constitute the angle θ with the ab planes, the x axis to enclose the angle θ with the c axis, and the y axis being in the ab plane and perpendicular to the magnetic field (see Fig. 1). We take the pinning potential to be of the form $V(\mathbf{r})=V_0[1-\cos(2\pi\mathbf{r}\cdot\mathbf{n}_c/s)]$, where $\mathbf{n}_c$ is the unit vector along the c direction. The chosen form of the pinning potential is justified for the temperature range where $\xi_c>s$ (e.g., Ref. 2). Substituting this potential in Eq. (3) and expanding with respect to the velocity we find

$$\frac{\delta\rho}{\rho}\equiv\frac{F_{\text{pin}}}{\eta v} =-\frac{V_0^2}{2\eta}\left[\frac{2\pi}{s}\right]^4\int dz\,dt\,t\dot{S}(\mathbf{q}_\perp;z,t) \times G_{xx}(\mathbf{r}_\perp=0;z,t)\exp(iq_zz)\ . \qquad (5)$$

The vector $\mathbf{q}$ describes the periodicity along the c axis and to first order in θ is given by $\mathbf{q}=(2\pi/s,0,2\pi\theta/s)$. The Fourier transform of the structure factor is given by

$$S(\mathbf{k},t)=n_v\int dz\sum_j\exp\{-ik_zz+ik[\mathbf{R}_{0j}(z,t) -\mathbf{R}_{00}(0,0)]\}\ , \qquad (6)$$

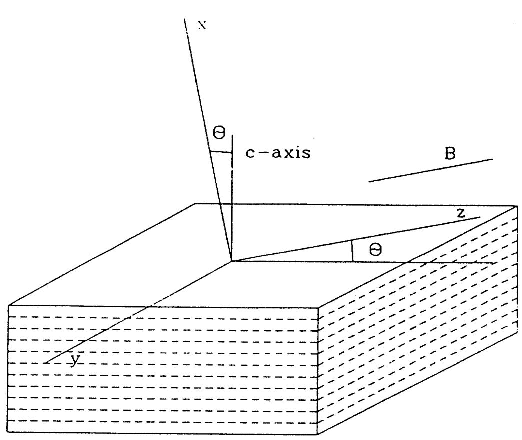

FIG. 1. Layered superconductor in tilted magnetic field. Choice of axes.

where $n_v=B/\Phi_0$ is an average density of vortex lines. In deriving Eq. (5) we used the condition $s/2\pi\ll a_x, a_x$ being the vortex lattice parameter in the x direction.

One can distinguish two different contributions to the integral in Eq. (5). The first one, $\delta\rho_1(\theta)$, comes from the small spatial and temporal scales at which mean fluctuational displacements $u_f(z,t)=\langle[u(z,t)-u(0,0)]^2\rangle^{1/2}$ become of the order of the interplane separation s. On these scales, fluctuations of vortex lines are uncorrelated, and the corresponding contribution can be viewed as from the individual vortices. Note that this single-vortex contribution should exist also in the vortex lattice phase. The second contribution, $\delta\rho_2(\theta)$, stems from the correlated long-range deformations of vortex lines with characteristic time τ_{pl} and corresponding spatial scale L_{pl} along the z axis. This correction diverges as temperature approaches the melting point.

Let us consider first the single-vortex contribution. At times scales shorter than the typical phonon time of the vortex lattice the interaction between vortex lines can be neglected and the structure factor transforms into

$$S(k_x,k_y=0;z,t)=\exp\left[-\frac{k_x^2}{2}u_f^2(z,t)\right]\ . \qquad (7)$$

The fluctuational displacement of a single-vortex line $u(z,t)$ is given by

$$u_f^2(z,t)=T\left[\frac{z}{\epsilon_1}+\frac{1}{\sqrt{\pi\eta_v\epsilon_1}}\int_0^t\frac{dt'}{\sqrt{t'}}\exp\left[-\frac{\eta_vz^2}{4\epsilon_1t'}\right]\right]\ , \qquad (8)$$

$\eta_v=\eta/n_v$ is the friction associated with the motion of a single vortex. The response function acquires the form

$$G(\mathbf{k}_\perp;z,t)=\frac{\exp(-\eta_vz^2/4\epsilon_1t)}{n_v\sqrt{4\pi\eta_v\epsilon_1t}}\ . \qquad (9)$$

ϵ_1 is the linear tension of the single vortex with respect to tilting deformations along the x direction which with logarithmic accuracy can be estimated as

$$\epsilon_1=\frac{\gamma\Phi_0^2}{(4\pi\lambda_{ab})^2}\ln\left[\frac{1}{k_z\gamma s}\right]\ , \qquad (10)$$

where λ_{ab} is the London penetration depth, γ is the anisotropy factor, and k_z^{-1} is the characteristic scale of vortex deformations. Substituting expressions (7) and (9) into Eq. (5) one finds [recall that the angular dependence is given by the factor $\exp(iq_zz)$ in the integrand in Eq. (5)]

$$\frac{\delta\rho_1}{\rho}=-rF(\theta/\theta_1)\ , \qquad (11)$$

with the amplitude of the anomaly,

$$r=\frac{6}{\sqrt{\pi}}\left[\frac{s}{2\pi}\right]^4\frac{V_0^2\epsilon_1^2}{T^4}\ . \qquad (12)$$

The universal function $F(\alpha)$ is given by

$$F(\alpha)=\int\frac{du\,\exp(-u^2)}{[(1/\sqrt{\pi})\exp(-u^2)+u\,\text{erf}(u)+i\alpha u]^4}\ , \qquad (13)$$

with the error function $\mathrm{erf}(u)=(2/\sqrt{\pi})\int_0^u ds\,\exp(-s^2)$. The angle $\theta_1=(2\pi/s)(T/2\epsilon_1)$ defines the characteristic tilt above which the pinning-induced single-vortex friction drops rapidly. A plot of the function $F(\alpha)$ is given in Fig. 2. The condition for the applicability of the perturbation theory is given by $rF(0)\ll 1$ which reduces to $T\gg 0.4s\sqrt{V_0\epsilon_1}$. Taking temperatures slightly above the melting temperature $T_m\approx\epsilon_1\sqrt{\Phi_0/B}\,c_L^2$, where c_L is the Lindemann number, one easily finds that this condition is well satisfied in the experimental range of fields.

Now we turn to the singular contribution from larger scales and times. This contribution is important for a strongly correlated vortex liquid in the vicinity of the melting temperature, where the remnance of crystalline order in the vortex configuration comes into play. In the solid vortex phase this contribution would give rise to a divergent correction to resistivity, but in the liquid phase the fluctuations at very large scales cut off this divergence. Slightly above the melting temperature the residual of the crystalline order gives rise to a considerable enhancement of $\delta\rho$. Additional complications arise from the effective mass anisotropy that leads to a complex form of the elastic energy in the vortex crystalline state. It is useful to perform scaling transformations analogous to those used in Refs. 13 and 14, $\tilde{x}_c=\gamma^{2/3}x_c$, $\tilde{x}_{a,b}=\gamma^{-1/3}x_{a,b}$, to avoid these complications. In the following we note all quantities related to the new coordinate system by the sign $\sim$. In this new set of coordinates, the system under consideration becomes isotropic and the elastic properties of the vortex crystal are described by the shear and nonlocal tilt module $\tilde{C}_{66}=\Phi_0\tilde{B}/(8\pi\tilde{\lambda})^2$, $\tilde{C}_{44}=\tilde{B}^2/4\pi/[1+(\tilde{\lambda}\tilde{k})^2]+\tilde{B}\tilde{\epsilon}_1/\Phi_0$, with $\tilde{\lambda}=\gamma^{1/3}\lambda_{ab}$ and $\tilde{B}=\gamma^{2/3}B$.

We assume that the melting occurs via the weak first-order transition. This means that the vortex liquid in the vicinity of the transition is expected to be strongly correlated, i.e., to behave as a solid at times not exceeding the characteristic plastic deformation times $\tau_{\rm pl}$ and for distances along the z direction less than the corresponding spatial scale $\tilde{L}_{\rm pl}$. The relation between these quantities can be estimated from the equation of dissipative motion of vortex configuration: The dissipative force $\tilde{\eta}u/\tau_{\rm pl}$ is of the same order of magnitude as the elastic restoring force $\tilde{C}_{44}u/\tilde{L}_{\rm pl}^2$, where u is the characteristic displacement corresponding to the time $\tau_{\rm pl}$. As a result we find $1/\tau_{\rm pl}\sim\tilde{C}_{44}/\tilde{\eta}\tilde{L}_{\rm pl}^2$. Such "viscoelastic" behavior can be quantitatively described by means of the $\tilde{k}_z$- and ω-dependent shear modulus $\tilde{C}_{66}(\omega,\tilde{k}_z)$ analogous to that introduced in the Maxwell description of very viscous liquid[15]:

$$\tilde{C}_{66}(\omega,k_z)=\frac{\tilde{C}_{66}}{1-1/(i\omega\tau_{\rm pl}-\tilde{k}_z^2\tilde{L}_{\rm pl}^2)} . \qquad (14)$$

Within this approximation the structure factor in the region $\tilde{k}_\perp\gg(\tilde{B}/\Phi_0)^{1/2}$ is given by Eq. (7). The behavior of $\tilde{u}_f^2(\tilde{z},t)$ at large spatial and time scales $t>\tau_{\rm ph}$, $\tilde{z}>\tilde{L}_{\rm ph}$, where $\tau_{\rm ph}^{-1}=4\pi\tilde{C}_{66}\tilde{B}/\Phi_0\tilde{\eta}$, $\tilde{L}_{\rm ph}=(\tilde{\epsilon}_1/\tilde{C}_{66})^{1/2}$, is affected by the intervortex interaction. At $t\sim\tau_{\rm ph}$ ($\tilde{z}\sim\tilde{L}_{\rm ph}$) the fluctuational displacement achieves its "Lindemann" value[17,18] (i.e., the fluctuation displacement in the crystal state in the vicinity of the melting transition) $\tilde{u}_f^2(t)\simeq\tilde{u}_L^2=T/(2\sqrt{\pi}\sqrt{\tilde{\epsilon}_1\tilde{C}_{66}})$, but in contrast to the behavior of the distortion in the crystalline phase, in the liquid phase $\tilde{u}_f(z,t)$ keeps growing slowly, the growth rate being determined by parameters $\tau_{\rm pl}$ and $\tilde{L}_{\rm pl}$. Time dependences of the fluctuation displacements in the crystalline and the liquid states are illustrated in Fig. 3.

The strength of correlation in the liquid is determined by the dimensionless parameter $\tau_{\rm pl}/\tau_{\rm ph}$. At $t>\tau_{\rm pl}$, $\tilde{z}>\tilde{L}_{\rm pl}$ the fluctuational displacement $\tilde{u}_f(z,t)$ and the response function $\tilde{G}(0,\tilde{z},t)$ behave like those for a single vortex but with renormalized expressions for the linear tension $\tilde{\epsilon}_{\rm eff}=4\pi\tilde{C}_{66}\tilde{L}_{\rm pl}^2/\ln(4\pi(\tilde{B}\tilde{R}_{\rm pl}^2/\Phi_0)$ and the viscous friction $\tilde{\eta}_{\rm eff}=2\pi\tilde{C}_{66}\tau_{\rm pl}/\ln(4\pi\tilde{B}\tilde{R}_{\rm pl}^2/\Phi_0)$, where $\tilde{R}_{\rm pl}=\sqrt{\tilde{C}_{66}/\tilde{C}_{44}}\tilde{L}_{\rm pl}$ (if nonlocality in $\tilde{C}_{44}$ is relevant, this gives $\tilde{R}_{\rm pl}=\sqrt{\tilde{a}\tilde{L}_{\rm pl}}$). In other words the effect of interactions gives rise to the "dressing" and the effective stiffening of the vortex line on large scales. This stiffening results, in particular, in the renormalization of the diffusion coefficient of the single-vortex line in a liquid state discussed in Ref. 16. Therefore, the large-scale contribution is once again given by Eq. (11) with renormalized ϵ_1 and reduced by the Debye-Waller factor

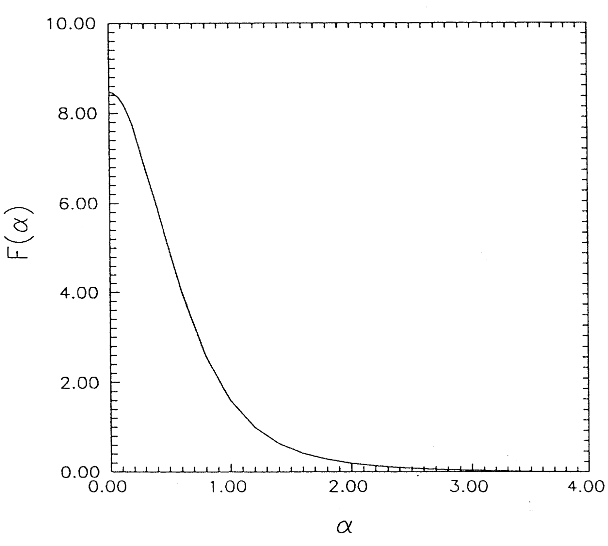

FIG. 2. Universal function $F(\alpha)$ (see text).

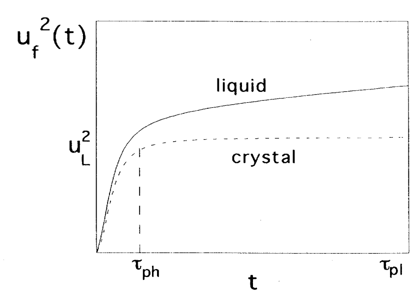

FIG. 3. Schematic time dependences of the fluctuation displacements in the crystal and strongly correlated liquid.

amplitude of the periodic potential:

$$\frac{\delta\rho_2}{\rho}=r\left[\frac{\epsilon_{\rm eff}}{\epsilon_1}\right]^2\exp\left[-\frac{1}{2}\left[\frac{2\pi u_L}{s}\right]^2\right]F(\theta/\theta_2)\ , \qquad (15)$$

where $\theta_2=2\pi/s\,(T/2\epsilon_{\rm eff})\ll\theta_1$. The latter equation is already written for the real coordinate system. For comparison with the single-vortex contribution it is convenient to rewrite the formula for $\epsilon_{\rm eff}$ as

$$\epsilon_{\rm eff}=\frac{\gamma\Phi_0^2}{(4\pi\lambda)^2}\left[\frac{L_{\rm pl}}{L_{\rm ph}}\right]^2\Big/\ln\left[\frac{L_{\rm pl}}{L_{\rm ph}}\right]\ . \qquad (16)$$

The value of the Debye-Waller factor is determined by the Lindemann fluctuation displacement near the melting point which can be estimated as $u_L\approx c_L a_x$, where $c_L\sim0.1$ is the Lindemann constant.

The result (15) holds as long as $\delta\rho_2/\rho\ll1$. One can see that near the melting temperature this condition may be violated if $L_{\rm pl}(T_m)$ grows too large.

Equations (11) and (15) represent the main result of our paper. The resulting angular dependence $\rho(\theta)$ is the superposition of the conventional angular resistivity dependence due to anisotropy and corrections (11) and (15) (Fig. 4). The correction from (11) varies over a wider angular interval than the correction from (15). The correction $\delta\rho_2$ typically has a smaller amplitude due to the Debye-Waller factor; however, it diverges quickly as the temperature approaches the melting point due to the divergent behavior of $\epsilon_{\rm eff}$.

The upper characteristic angle θ_1 may be connected to the experimentally accessible quantities as follows:

$$\theta_1=(2\pi)^2\frac{T_c}{\Phi_0\gamma s}\left[\frac{dH_{c1}}{dT}\right]^{-1}\frac{1}{T_c-T}\ . \qquad (17)$$

For typical values of these parameters for Y-Ba-Cu-O $dH_{c1}/dT=10$ Oe/K, $\gamma=8$, $s=13$ Å, $T_c-T=1$ K, and $\theta_1\simeq16°$. The latter estimate means that the range of the angular variation of $\delta\rho_1$ is the same as for the angular dependence of the flux flow resistivity $\rho_0(\theta)$ governed by anisotropy:

$$\rho_0(\theta)=\rho(\theta^\circ)\{\cos^2(\theta)+\gamma^2(\sin\theta)^2\}^{1/2}\ . \qquad (18)$$

This means that the resulting experimental angular dependence at large angle scales can be observed as a distorted and enhanced anisotropic resistivity behavior with a temperature dependent ratio $\rho(0°)/\rho(90°)$.

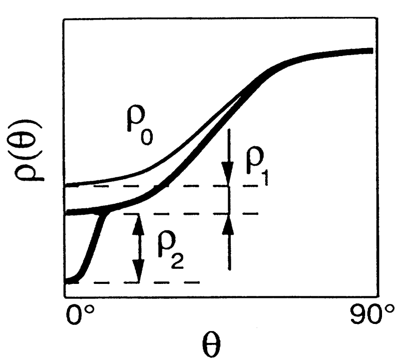

FIG. 4. Schematic angular dependence of the resistivity. The angular width of the anomalous correction $\delta\rho_2(\theta)$ is exaggerated for clarity.

The behavior of the second correction $\delta\rho_2(\theta)$ is determined by the critical behavior near the melting point T_m. The growth of correlations near the transition leads to a decrease of the typical angle θ_2 and to an increase of the amplitude of the anomaly. In fact, the correction $\delta\rho_2$ is noticeable only if the regime of strongly correlated vortex liquids takes place. Therefore, the observation of this correction would mean that the transition from a vortex solid to a vortex liquid is almost continuous. For the first-order phase transition both the typical angle θ_2 and the resistivity at $\theta=0$ should jump to zero at $T=T_m$ from finite values.

Real samples always have some misorientations of the c axis which tend to smear out the angular anomaly in resistivity. The anomaly can be resolved only in the temperature region where the typical angle θ_2 is larger than the typical misorientation angle θ_m. In the opposite case $\theta_2<\theta_m$ the intrinsic angular behavior is averaged out by misorientation. In this case the angular dependence is determined by the angular distribution function $P(\theta)[P(\theta)$ gives the probability to find a crystallite which constitutes the angle θ with average orientation]. In this region the angular dependence of resistivity simply follows that of $P(\theta)$:

$$\frac{\delta\rho_2}{\rho}=Cr\theta_2\left[\frac{\epsilon_{\rm eff}}{\epsilon_1}\right]^2\exp\left[-\frac{1}{2}\left[\frac{2\pi u_L}{s}\right]^2\right]P(\theta)\ , \qquad (19)$$

with $C=\int d\alpha F\alpha=2\pi^{5/2}/3=11.66$. The width of the angular anomaly is given by the temperature-independent misorientation angle θ_m, whereas the amplitude of the correction still grows as the temperature approaches the melting point from above. Such a behavior has been observed experimentally in the resistive measurement of Kwok *et al.*[6]

In conclusion, we have calculated the corrections to Ohmic resistivity due to the intrinsic periodic potential associated with the layered structure. We have found the complex angular dependence of the resistivity for small angles between the magnetic field and the *ab* planes. The sharp dip in $\rho(\theta)$, $\delta\rho_2(\theta)$ originates from the correlated motion of large vortex bundles in a vortex liquid, and can be found on the bottom of the wider dip in the resistivity $\delta\rho_1(\theta)$ resulting from the pinning of single vortices (see Fig. 4). The obtained results describe qualitatively the observe shape of the angular dependence of the resistivity and can be used to infer quantitative information about the magnitude of the intrinsic pinning potential and the correlations in the vortex liquid state near the melting transition.

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