1-1-2000

Interlayer quasiparticle transport in the vortex state of josephson coupled superconductors

I. Vekhter
University of Guelph

L. N. Bulaevskii
Los Alamos National Laboratory

A. E. Koshelev
Argonne National Laboratory

M. P. Maley
Los Alamos National Laboratory

Interlayer Quasiparticle Transport in the Vortex State of Josephson Coupled Superconductors

I. Vekhter[1], L. N. Bulaevskii[2], A.E. Koshelev[3], and M.P. Maley[2]
[1]*Department of Physics, University of Guelph, Guelph, Ontario N1G 2W1, Canada*
[2]*Los Alamos National Laboratory, Los Alamos, NM 87545*
[3]*Materials Science Division, Argonne National Laboratory, Argonne, Illinois 60439*
(May 19, 2018)

We calculate the dependence of the interlayer quasiparticle conductivity, σ_q, in a Josephson coupled d-wave superconductor on the magnetic field $\mathbf{B} \parallel c$ and the temperature T. We consider a clean superconductor with resonant impurity scattering and a dominant coherent interlayer tunneling. When pancake vortices in adjacent layers are weakly correlated at low T the conductivity increases sharply with B before reaching an extended region of slow linear growth, while at high T it initially decreases and then reaches the same linear regime. For correlated pancakes σ_q increases much more strongly with the applied field.

Experimental study of quasiparticle properties in high-temperature superconductors so far has focussed almost exclusively on thermodynamic and in-plane transport properties. Two remarkable theoretical results connect these properties with the symmetry of the superconducting gap: a) a massless Dirac spectrum of nodal quasiparticles in a d-wave superconductor leads to a finite density of states (DOS) at the Fermi level in presence of impurity scattering [1] and to a universal (impurity independent) low-temperature limit of the in-plane thermal conductivity, $\kappa_{00} \equiv \lim_{T\to 0} \kappa_{ab}(0,T)/T$, in the absence of a magnetic field. [2,3]; b) in the vortex state the DOS is locally enhanced due to the effect of the supercurrents around the vortex cores on the near-nodal quasiparticles (Volovik effect) [4]. Consequently, the effect of the magnetic field on the electrical and the thermal conductivities is twofold: the net enhancement of the DOS tends to increase the conductivity, while local changes in the phase and amplitude of the order parameter lead to its decrease. [5,6]

The enhancement of the DOS with the magnetic field, B, has been invoked to explain the increase of the specific heat with B observed in $YBa_2Cu_3O_{7-\delta}$ (YBCO). [7] The in-plane thermal conductivity, κ_{ab}, has been found to have a highly non-trivial dependence on the applied field and the temperature in both YBCO and $Bi_2Sr_2CaCu_2O_{8+\delta}$ (Bi-2212) compounds. [8,9] At low temperatures κ_{ab} increases with B, while at $T \geq 5-10$ K it decreases and then often saturates in fields above a few Tesla. In some zero-field cooled Bi-2212 samples a hysteretic field dependence with two plateaus at different values of $\kappa_{ab}(B)$ has been observed. [8] This indicates that the behavior of $\kappa_{ab}(B)$ is sensitive to the vortex arrangement in the sample.

Franz [6] has appealed to the effects of disorder in the vortex positions to explain the observed high-field plateau in $\kappa_{ab}(B)$. He has argued that randomly positioned vortices act similarly to impurities leading to the enhancement of both the DOS and the electron scattering. These two effects compensate each other at high fields (in analogy to the zero-field result being insensitive to impurity concentration) leading to the universal $\kappa_{ab}(B,T)/T \to \kappa_{00}$. However, to obtain this result Franz has made questionable approximations: he has replaced the spatial average of the product of two Green's functions by the product of the averages, and has assumed a Lorentzian shape of the distribution function of the in-plane supervelocity, which differs from a realistic distribution in the vortex state, see below. Clearly, this approach cannot explain the different (non-universal) plateau values. [8]

Recently it has been demonstrated [10,11] that the *interlayer* quasiparticle electrical transport can be studied directly in the resistive state of small area samples (mesas) fabricated from the Josephson coupled superconductors such as Bi-2212. In the energy and temperature range below ≈ 3 meV, the experimental data are well understood in the framework of a Fermi-liquid model for the near-nodal quasiparticles in a d-wave superconductor assuming (i) clean limit, (ii) resonant (unitarity) impurity scattering, and (iii) dominant contribution to the c-axis conductivity from *coherent* (conserving the in-plane momentum) *interlayer tunneling.* [11] The authors of Ref. [11] introduced the universal quasiparticle c-axis conductivity $\sigma_q(T=0, B=0)$. This quantity is insensitive to the impurity vertex corrections which modify the in-plane conductivity as discussed by Durst and Lee. [3]

So far the experimental data on the magnetic field dependence of interlayer quasiparticle conductivity have been quite scarce. Measurements of $\sigma_c(B)$ in high fields, where $\sigma_c \approx \sigma_q$, reveal a linear increase of σ_c with B on the scale of 40 T. [12] For mesas Yurgens *et al.* [10] have reported a weak field dependence for $B < 7$ T without quantitative results for its functional form.

In this Letter we address the effect of vortices on the c-axis quasiparticle conductivity, $\sigma_q(B,T)$, as a function of the magnetic field $\mathbf{B} \parallel c$, in Josephson coupled superconductors in the framework of the approach used in Ref. [11]. We find that at low (high) temperatures for c-axis uncorrelated vortices $\sigma_q(B,T)$ increases (decreases) with B, before reaching an extended region of slow linear-in-B increase. In this high field regime the increase in the DOS is largely compensated by the enhancement of the electron scattering at tunneling, due to disorder in the positions of vortices *in neighboring layers*.

To calculate the effect of vortices on the interlayer transport in the framework of the d-wave model we employ a semiclassical approach. [4] The supercurrents around vortex cores lead to a Doppler shift, $\epsilon_n(\mathbf{k},\mathbf{r}) = \mathbf{k}\cdot\mathbf{v}_{sn}(\mathbf{r})$, in the quasiparticle spectrum, $\mathcal{E}(\mathbf{k},\mathbf{v}_s) = E_{\mathbf{k}} + \epsilon_n$. Here $\mathbf{k}$ is the quasiparticle momentum and $\mathbf{v}_{sn}(\mathbf{r})$ is the supervelocity at point $\mathbf{r}$ inside layer n, $E_{\mathbf{k}} \approx [\xi_{\mathbf{k}}^2 + \Delta^2(\mathbf{k})]^{1/2}$ is the quasiparticle energy at $B=0$, $\xi_{\mathbf{k}} = \mathbf{v}_F\cdot(\mathbf{k}-\mathbf{k}_F)$ and $\Delta(\mathbf{k}) = \Delta_0(k_x^2 - k_y^2)/k^2$. As c-axis currents are parallel to the field (in contrast to the in-plane currents), [5] the net interlayer transport is determined by the spatial average of local transport coefficients. In the following we neglect correlations between the positions of impurities and those of vortices, and average over the distribution of each independently, see below. We also neglect the temperature dependence of Δ_0 and small effects due to Zeeman splitting. Then the interlayer quasiparticle conductivity is given by [11]

$$\frac{\sigma_q(B,T)}{\sigma_q(0,0)} = \frac{\Delta_0}{8TN(0)}\int_{-\infty}^{+\infty}\frac{d\omega}{\cosh^2(\omega/2T)}\int d\mathbf{k} \times (t^2(\mathbf{k})/t_0^2)\,\langle A(\mathbf{k},\omega+\epsilon_n)A(\mathbf{k},\omega+\epsilon_{n+1})\rangle\,, \tag{1}$$

where $\sigma_q(0,0) = 2e^2t_0^2N(0)s\eta/\pi\hbar\Delta_0$ is the universal c-axis conductivity, $N(0)$ is the 2D DOS, $t(\mathbf{k})$ is the interlayer transfer integral, $t_0 = t(\mathbf{k}_g)$, $\mathbf{k}_g$ are the positions of the gap nodes on the Fermi surface, and $0<\eta<1$ is the weight for the coherent tunneling. Further, $\langle ...\rangle$ denotes the spatial average, and

$$A(\mathbf{k},\omega) \approx \left(1+\frac{\xi_{\mathbf{k}}}{E_{\mathbf{k}}}\right)\frac{L(\omega,E_{\mathbf{k}})}{2} + \left(1-\frac{\xi_{\mathbf{k}}}{E_{\mathbf{k}}}\right)\frac{L(\omega,-E_{\mathbf{k}})}{2}$$

is the spectral density averaged over impurities with $L(\omega,E) = \gamma(\omega)/\pi[(E-\Omega(\omega))^2+\gamma^2(\omega)]$. The functions $\omega-\Omega(\omega)$ and $-\gamma(\omega)$ are the real and the imaginary part of the self-energy respectively. [2] In the unitarity limit the effective scattering rate of quasiparticles is $\gamma(\omega) \approx \gamma_0 - \omega^2/8\gamma_0$ when $\omega \ll \gamma_0$ and $\gamma(\omega) \approx \pi\gamma_0^2/2|\omega|$ when $\omega \gg \gamma_0$, where $\gamma_0 \approx (\hbar\nu_0\Delta_0)^{1/2}$ and ν_0 is the bare scattering rate. The renormalized frequency is $\Omega(\omega) \approx \omega/2$ when $\omega \ll \gamma_0$ and $\Omega(\omega) \approx \omega$ if $\omega \gg \gamma_0$.

At low temperatures $T \ll \Delta_0$ the quasiparticle current comes mainly from the regions near the gap nodes. In the vicinity of a node we can linearize the quasiparticle spectrum and obtain from Eq. (1)

$$\frac{\sigma_q(B,T)}{\sigma_q(0,0)} = \int_{-\infty}^{+\infty} d\omega \int_0^\infty dE \frac{3\pi^2 E}{8T\cosh^2(\omega/2T)} \times \langle L(\omega+\epsilon_{n+1},E)[L(\omega+\epsilon_n,E) + (1/3)L(\omega+\epsilon_n,-E)]\rangle, \tag{2}$$

where $\epsilon_n(\mathbf{r}) = \mathbf{k}_g\cdot\mathbf{v}_{sn}(\mathbf{r})$ and we have set $t(\mathbf{k}) \simeq t_0$. The characteristic energy scale for the Doppler shift, ϵ_n, is $\epsilon_B = \hbar v_F/a$, where $a = (\Phi_0/B)^{1/2}$ is typical intervortex distance. The semiclassical approach is valid for $\epsilon_B \ll \Delta_0$, i.e. for fields $B \ll B_\Delta$, where $B_\Delta = \Phi_0\Delta_0^2/\hbar^2v_F^2$. This approach also does not account correctly for the quasiparticles in the regions near the vortex cores, which leads to corrections of the order of $\delta\sigma_q/\sigma_q \approx \epsilon_B^2/\Delta_0^2 = B/B_\Delta$. Another important field scale can be obtained by comparing ϵ_B with the scattering rate γ_0; the field $B_\gamma = \Phi_0\gamma_0^2/\hbar^2v_F^2$ separates the regimes of impurity dominated and field dominated behavior. Using the parameters $\Delta_0 \approx 25$ meV, $\gamma_0 \approx 2-3$ meV [11] and $v_F \approx 1.5 - 2.5\cdot 10^7$ cm/s [13] we obtain $B_\gamma \approx 0.2-0.6$ T and $B_\Delta \approx 40-80$ T.

We now rewrite the spatial average in Eq.(2) as the average over the probability distribution of the Doppler shift, ϵ, which is fully determined by the vortex arrangement. In Bi-2212 crystals the 3D vortex lattice is destroyed by pinning at the "second peak field" $B \approx 0.02$-0.05 T (see, e.g., [14]), and at higher fields pancakes in neighboring layers are only weakly correlated. Consequently, we consider the limits of c-axis-correlated and uncorrelated pancakes. In both cases the average depends only on the probability distribution function of ϵ in a single layer, $\mathcal{P}(\epsilon)$, which is related to the distribution $P(p_x)$ of a single component of supermomentum $p_x = 2mv_{sx}$, where m is the effective mass, by $\mathcal{P}(\epsilon) = (2/v_F)P(p_x = 2\epsilon/v_F)$. The function $P(p_x)$ is determined by the pancake configuration $\{\mathbf{R}_i = (X_i,Y_i)\}$, and, when a is smaller than the London penetration depth λ_{ab}, is given by $P(p_x) = \langle\delta(p_x - \hbar\sum_i Y_i/R_i^2)\rangle_{\{R_i\}}$. For an isolated vortex $p \propto \hbar/R$ where $\xi_{ab} \ll R \ll a$ is the distance from the pancake center, and ξ_{ab} is the coherence length. Consequently, $\mathcal{P}(\epsilon)$ has a universal tail $\mathcal{P}(\epsilon) = \pi\epsilon_B^2/(8\epsilon^3)$ at $\epsilon_B \ll \epsilon \ll \Delta_0$, while its behavior at $\epsilon \lesssim \epsilon_B$ depends on the actual positions of vortices. However, in absence of correlation between the positions of vortices and impurities $\mathcal{P}(\epsilon)$ involves a single energy scale ϵ_B for *any* ϵ, and depends on ϵ only via ϵ/ϵ_B, i.e. $\mathcal{P}(\epsilon)d\epsilon = \mathcal{P}_B(\epsilon/\epsilon_B)d\epsilon/\epsilon_B$. In the clean limit we can extend the asymptotic behavior $\mathcal{P}(\epsilon) \propto 1/\epsilon^3$ to infinity when the integral over ϵ in Eq. (2) converges. Therefore, importantly, up to terms of the order of γ_0/Δ_0, the conductivity depends only on the dimensionless parameters $\epsilon_B/\gamma_0 \equiv \sqrt{B/B_\gamma}$ and T/γ_0.

For c-axis-correlated vortices $\epsilon_n \approx \epsilon_{n+1}$, and the average in Eq. (2) means $\langle\mathcal{F}(\epsilon_n,\epsilon_{n+1})\rangle \approx \int d\epsilon\mathcal{P}(\epsilon)\mathcal{F}(\epsilon,\epsilon)$. At low temperatures, $T \ll \gamma_0$, both in weak ($B \ll B_\gamma$) and in strong ($B \gg B_\gamma$) fields the leading field-dependent part of the conductivity varies as $(\epsilon/\gamma_0)^2$. We therefore separate σ_q into the contributions from the regions $\epsilon < \gamma_0$ and $\epsilon > \gamma_0$, and use the asymptotic behavior of $\Omega(\omega)$ and $\gamma(\omega)$ in each case to estimate

$$\frac{\sigma_q(B,0)-\sigma_q(0,0)}{\sigma_q(0,0)} \approx \frac{1}{6}\frac{\langle\epsilon^2\rangle_{\epsilon<\gamma_0}}{\gamma_0^2} + \frac{3\pi^2}{8}\frac{\langle\epsilon^2\rangle_{\epsilon>\gamma_0}}{\gamma_0^2}. \quad (3)$$

As the average is dominated by the tail of $\mathcal{P}(\epsilon)$ we can evaluate σ_q in weak and strong fields with logarithmic accuracy. Moreover, we interpolate between the two limits to obtain also a qualitative description at $B \sim B_\gamma$,

$$\sigma_q(B,T)/\sigma_q(0,0) \approx 1 + \pi^2T^2/18\gamma_0^2 + \frac{3\pi^3 B}{64B_\gamma}\left[\ln\frac{B_\Delta}{(B^2+B_\gamma^2)^{1/2}} + \frac{4}{9\pi^2}\ln\frac{(B^2+B_\gamma^2)^{1/2}}{B}\right]. \quad (4)$$

Therefore at low T for the 3D-ordered case σ_q increases quasi-linearly with B in the entire field range up to B_Δ.

For the c-axis uncorrelated vortices the average in Eq. (2) has to be taken independently in each layer, $\langle\mathcal{F}(\epsilon_n,\epsilon_{n+1})\rangle = \int d\epsilon_1\mathcal{P}(\epsilon_1)\int d\epsilon_2\mathcal{P}(\epsilon_2)\mathcal{F}(\epsilon_1,\epsilon_2)$. Again, for $T \ll \gamma_0$ and $B \ll B_\gamma$ we expand in ϵ_n/γ_0, and obtain

$$\frac{\sigma_q(B,T)}{\sigma_q(0,0)} \approx 1 + \frac{\pi^2T^2}{18\gamma_0^2} + \frac{\pi B}{96B_\gamma}\ln\frac{B_\gamma}{B}. \quad (5)$$

At high fields the variation of ϵ_n is of the order of $\epsilon_B \gg \gamma_0$ and consequently $\langle L(\omega+\epsilon_n,E)\rangle \approx \mathcal{P}(E-\omega)$. Then for $B_\gamma \ll B \ll B_\Delta$ the conductivity is given by

$$\frac{\sigma_q(B,T)}{\sigma_q(0,0)} = C_1 + \frac{B}{B_0}, \quad C_1 = 2\pi^2\int_0^\infty d\epsilon\epsilon\mathcal{P}^2(\epsilon), \quad (6)$$

The most important contribution to the linear, in B, term is due to the increase in the tunneling away from the nodes, $t(\mathbf{k}) \simeq t_0 + t_1\phi^2$, where ϕ is the angle between $\mathbf{k}$ and the nodal direction, and $t_1 \gg t_0$, [15] which yields $B_0 \sim (t_0/t_1)B_\Delta$. Other contributions, due to deviations of the quasiparticle spectrum from the massless Dirac form of the linearized dispersion and due to corrections to the semiclassical approximation in the vicinity of the vortex cores, only enhance σ_q on the scale of $B_\Delta \gg B_0$. Very importantly, due to scaling of $\mathcal{P}(\epsilon) = \mathcal{P}_B(\epsilon/\epsilon_B)/\epsilon_B$ (see above), C_1 in Eq. (6) is a constant which depends solely on the *shape* of the distribution $\mathcal{P}(\epsilon)$, but not on the magnetic field. If $\mathcal{P}_B(x)$ is monotonous C_1 depends on its asymptotic decay. For the Lorentzian (when vortices act *exactly* as the impurity scatterers do at $B=0$) $\mathcal{P}_B(x) \propto x^{-2}$ at $x \gg 1$, and $C_1 = C_L = 1$, while for the Gaussian distribution $C_1 = C_G = \pi/2$. In the vortex state $\mathcal{P}_B(x) \propto x^{-3}$ at large x, and we expect $1 < C_1 < \pi/2$. Therefore to find σ_q we need accurate information on pancake arrangement to determine the supervelocity distribution $\mathcal{P}(\epsilon)$.

Below the irreversibility line $T_{irr}(B)$ at high fields and in the vortex liquid state pancakes do not possess long range order inside layers due to pinning and thermal fluctuations. To calculate $\sigma_q(B,T)$ in these regimes we use the distribution function $P(p_x)$ obtained by numerical simulation of the 2D pancake liquid at different values of the dimensionless temperature $t = T_{eff}/2\pi E_0$, which characterizes pancake disorder inside the layers. Here $E_0 = \Phi_0^2 s/16\pi^3\lambda_{ab}^2$ is the characteristic energy of pancake interaction. In the liquid state $T_{eff} = T$, while in the glass state it is reasonable to take $T_{eff} \approx T_{irr}(B)$. The calculated function $P(p_x)$ is shown in Fig. 1 for several values of t; the inset shows that the parameter C_1 grows with t from 1.08 to 1.22. The full field dependence of $\sigma_q(B,T)/\sigma_q(0,0)$ at $T \ll \gamma_0$ obtained from these distribution functions is shown in Fig. 2. In computing σ_q we have determined $\gamma(\omega)$ and $\Omega(\omega)$ self-consistently [16] in the unitarity limit, choosing the scattering rate so that $\gamma_0 = 0.1\Delta_0$, typical of Bi-2212 samples. [11]

Due to short range correlations inherent to strongly interacting vortices at $B > \Phi_0/\lambda_{ab}^2$ the calculated function $P(p_x)$ differs significantly from the Gaussian (obtained in Ref. [17] assuming uncorrelated and random vortex positions). The distribution function depends on a single variable ϵ/ϵ_B only if the positions of vortices and those of impurities are uncorrelated; this is the case for a vortex solid and for the liquid phase in presence of weak pinning. In general, below the irreversibility line impurities act as pinning centers, and the distribution function depends on two variables, ϵ/ϵ_B and $B/\Phi_0 n_i$, where n_i is the impurity concentration in a layer. We believe nevertheless that in that regime our approach gives at least the correct qualitative behavior.

At $T \gg \gamma_0$, we also consider the intermediate fields, $B_\gamma \ll B \ll B_T = \Phi_0 T^2/\hbar^2 v_F^2$, when the quasiparticle concentration is determined by the temperature and the primary effect of vortices is to increase the scattering at tunneling. Consequently σ_q decreases with field

$$\frac{\sigma_q(B,T)}{\sigma_q(0,0)} \approx C_2\sqrt{\frac{B_T}{B}}, \quad C_2 = \frac{3}{2}\pi^2\ln 2\int_0^\infty dx\mathcal{P}_B^2(x). \quad (7)$$

At $B \gg B_T$ Eq. (6) holds. Therefore, for c-axis uncorrelated pancakes at $T \ll \gamma_0$, $\sigma_q(B,T)$ increases quasi-linearly with B with the slope $1/B_\gamma$ when $B \ll B_\gamma$, and with a smaller slope $\sim 1/B_0$, for $B_\gamma \ll B \ll B_\Delta$. At higher temperatures, $T \geq \gamma_0$, $\sigma_q(B,T)$ *decreases* with field for $B \lesssim B_T$, before crossing over to the slow linear growth with the slope $\sim 1/B_0$, see Fig. 3. Here the conductivity has been computed with a temperature-independent Δ_0. In fact Δ_0 is reduced with increasing T, and the plateau values increase. Note that linear interpolation of $\sigma(B,T)$ from high fields does *not* yield $\sigma_q(0,T)$.

To conclude, we find that: a) the field dependence of the quasiparticle interlayer conductivity is sensitive to the structure of vortex state; b) for the c-axis uncorrelated vortices the field dependence of the interlayer conductivity at low and high temperatures is quite different at low fields but becomes similar in high fields $B \gg \max[B_T, B_\gamma]$.

We thank M. J. Graf for helpful discussions. This work was supported by the Los Alamos National Laboratory under the auspices of the U.S. Department of Energy.

Work in Argonne was supported by the NSF Office of the Science and Technology Center under contract No. DMR-91-20000 and by the U.S. DOE, BES-Materials Sciences, under contract No. W-31-109-ENG-38. I.V. acknowledges the hospitality of Centre Émile Borel and Aspen Center for Physics, where part of this work was done.

[1] L.P. Gor'kov and P.A. Kalugin, Pis'ma ZhETF, **41**, 208 (1985) [JETP Lett. **41**, 253 (1985)].

[2] P.A. Lee, Phys. Rev. Lett. **71**, 1887 (1993); A.V. Balatsky, A. Rosengren, and B.L. Altshuler, Phys. Rev. Lett. **73**, 720 (1994); P.J. Hirschfeld, W.O. Putikka, and D.J. Scalapino, Phys. Rev. B **50**, 10250 (1994); M. J. Graf *et al.*, Phys. Rev. B **53**, 15147 (1996).

[3] Ref. [2] has also predicted the universal limit for the in-plane electrical conductivity. However, recent work has shown that the universality is violated if vertex and Fermi-liquid corrections are accounted for, see A. C. Durst and P. A. Lee, cond-mat/9908182.

[4] G.E. Volovik, Pis'ma Zh. Eksp. Teor. Fiz. **64**, 641 (1993) [JETP Lett. **58**, 6 (1993)]; see also S. K. Yip and J. A. Sauls, Phys. Rev. Lett. **69**, 2264 (1992).

[5] C. Kübert and P.J. Hirschfeld, Phys. Rev. Lett. **80**, 4963 (1998); Yu. S. Barash and A.A. Swidzinsky, Phys. Rev. B **58**, 6476 (1998); I. Vekhter and A. Houghton, cond-mat/9904077.

[6] M. Franz, Phys. Rev. Lett. **82**, 1760 (1999).

[7] K.A. Moler *et al.*, Phys. Rev. Lett. **73**, 2744 (1994); B. Revaz *et al.*, Phys. Rev. Lett. **80**, 3364 (1998); D. A. Wright *et al.*, Phys. Rev. Lett. **82**, 1550 (1999).

[8] K. Krishana *et al.*, Science **277**, 83 (1997); H. Aubin *et al.*, Science **280**, 9a (1998); Y. Ando *at al.*, cond-mat/9812265 v2.

[9] M. Chiao, Phys. Rev. Lett. **82**, 2943 (1999); H. Aubin *et al.*, Phys. Rev. Lett. **82**, 624 (1999); N. P. Ong *et al.*, cond-mat/9904160.

[10] A. Yurgens *et al.*, Phys. Rev. B **53**, R8887 (1996), Phys. Rev. Lett. **79**, 5122 (1997); **59**, 7196 (1999); K. Tanabe *et al.*, Phys. Rev. B **53**, 9348 (1996); K. Schlenga *et al.*, Phys. Rev. B **57**, 14518 (1998).

[11] Yu.I. Latyshev *et al.*, Phys. Rev. Lett. **82**, 5345 (1999).

[12] V.N. Zavaritskii and M. Springford, JETP Lett. **5**, 448 (1998), N. Morozov *et al.*, unpublished.

[13] H. Ding *et al.*, Phys. Rev. Lett. **74**, 2784 (1995).

[14] R. Cubitt *et al.*, Nature **365**, 407 (1993); C.M. Aegerter *et al.*, Phys. Rev. B **54** , R15661 (1996); B. Khaykovich *et al.*, Phys. Rev. B **56**, R517 (1997); A.E. Koshelev and V.M. Vinokur, Phys. Rev. B **57** 8026(1998)

[15] L.B.Ioffe and A.J.Millis, Phys. Rev. B **58**, 11631 (1998).

[16] We ignore here the change of the internode scattering with the Doppler shift. [5] If included, it only affects the intermediate field regime, and does not change the low and high field limits.

[17] F. Yu *et al.*, Phys. Rev. Lett. **74**, 5136 (1995).

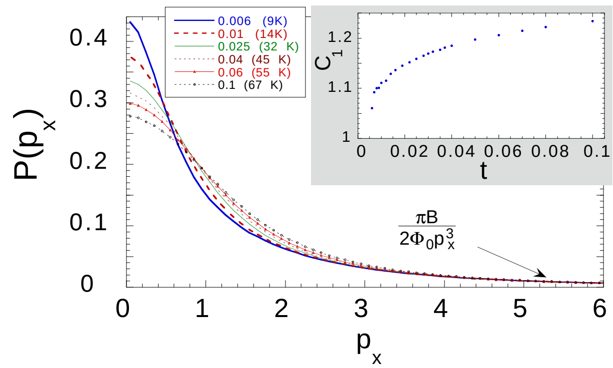

FIG. 1. The distribution functions $P(p_x)$ of the x-component of supermomentum. Here p_x is measured in units $\hbar/a_0$, where $a_0 = (2\Phi_0/\sqrt{3}B)^{1/2}$ is the lattice constant of a triangular vortex lattice. $P(p_x)$ has been calculated using Langevin dynamics simulations of a 2D liquid at different reduced temperatures $t = T/2\pi E_0$. The corresponding absolute temperatures obtained assuming $\lambda_{ab} = 200\text{nm}(1 - T^2/T_c^2)^{-1/2}$ are given in brackets. The distribution function of the Doppler shifts $\mathcal{P}(\epsilon)$ is related to $P(p_x)$ by $\mathcal{P}(\epsilon) = (2a_0/\hbar v_F)P(p_x = 2a_0\epsilon/\hbar v_F)$. Inset: temperature dependence of the parameter C_1 in Eq. (6).

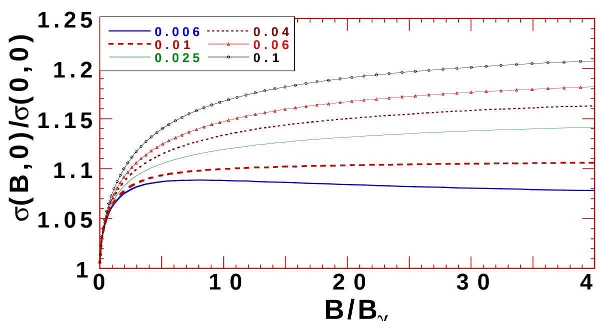

FIG. 2. Magnetic field dependence of the normalized quasiparticle conductivity $\sigma_q(B,0)/\sigma_q(0,0)$ at low temperatures $T \ll \gamma_0$ for c-axis uncorrelated vortex state calculated using functions $P(p_x)$ shown Fig. 1. Here we do not show the weak linear field dependence at $B/B_\gamma \gg 1$, see Eq. (6).

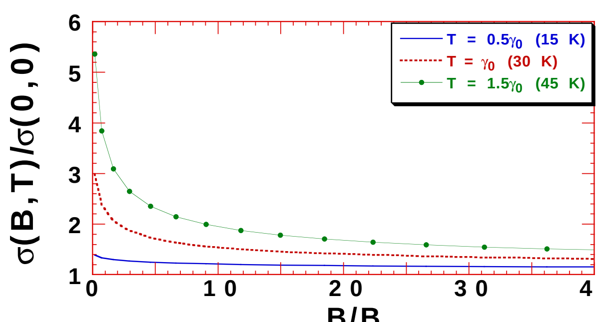

FIG. 3. Field dependence of the quasiparticle conductivity at temperatures $T \geq 0.5\gamma_0$, computed for the pancake liquid state with $P(p_x)$ obtained by simulations (Fig. 1) and $\gamma_0 = 0.16E_0 \approx 30$ K.