Absence of Cyclotron Resonance in the Anomalous Metallic Phase in InO_x

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It is observed that many thin superconducting films with not too high disorder level (generally $R_N/\square < 2000\ \Omega$) placed in magnetic field show an anomalous metallic phase where the resistance is low but still finite as temperature goes to zero. Here we report in weakly disordered amorphous InO_x thin films that this anomalous metal phase possesses no cyclotron resonance and hence non-Drude electrodynamics. The absence of a finite frequency resonant mode can be associated with a vanishing downstream component of the vortex current parallel to the supercurrent and an emergent particle-hole symmetry of this metal, which establishes its non-Fermi-liquid character.

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Conventional wisdom dictates that electrons confined to two dimensions and cooled to the limit of zero temperature can only be superconducting or insulating. Yet a number of examples of zero temperature 2D metallic states exist. A particularly interesting example is the anomalous metal where thin superconducting films in magnetic field may exhibit a saturated, but nonzero resistance much smaller than the normal state at low temperature. This is a widely observed phenomenon that manifests in a diverse number of physical systems: disordered thin films [1–8], Josephson junctions arrays [9,10], artificially patterned superconducting islands [11], and interfacial superconductivity [12]. It is underappreciated how anomalous this otherwise simple behavior is. The natural behavior of bosons in the limit of zero temperature is either condensed or localized. But here, due to the low resistance of the sample and the obvious manifestation of the superconducting transition, one requires the presence of Cooper pairs that move diffusively even in the apparent limit of zero temperature. In our below work we assume that this anomalous metal phase has been established by these previous experiments.

Although observed ubiquitously, the nature of this anomalous metal is still unclear [4,13–19]. It was observed in previous experiments [8] that its optical response is characterized by a very narrow conductance peak with a width that is orders of magnitude less than the normal state scattering rate (which is typically ∼100 THz). It was shown that over short timescales the system possesses a frequency-dependent phase stiffness [8], indicating that superconducting correlations are retained on small timescales and length scales. But how this is possible in a zero temperature dissipative phase, and to what extent this state's phenomenology is different from conventional metals, is unclear.

In this work, we have performed broadband microwave measurements of frequency, temperature, and field dependence of the complex microwave conductance on a low-disorder 2D superconducting InO_x film. These measurements have been extended to the very low field regime when samples are still highly conductive, giving unprecedented insight into how this anomalous metallic state develops from superconductivity. We show that even at small magnetic fields, although the superconducting delta function is retained, the low frequency dissipative modes almost immediately form. With further applied field the delta function is suppressed, leaving only the dissipative peak. Among other aspects we demonstrate that although on some level this anomalous metal can be characterized as a high mobility metal, the system possesses *no* cyclotron resonance, which is a ubiquitous feature of high mobility metals with conventional electrons. Our observation has much in common with recent observations that the anomalous metal state has—over a range of fields—no Hall effect [20]. Taken together, these works show that the anomalous metal state has an emergent particle-hole symmetry, that although was previously considered to be a property of the superconductor to insulator transition (SIT) [21] quantum critical point [5,27] is exhibited here over an entire intervening phase.

In Fig. 1(a), we show two terminal dc sheet resistance $R/\square$ measured in the Corbino geometry as a function of temperature. The resistance as a function of field at the low temperature of ≈0.4 K is presented in Fig. S3 of Supplemental Material (SM) [28]. One can see the characteristic phenomenon of the anomalous metallic phase [28] with apparently finite resistances persisting to the low temperature limit. The "mean field" superconducting transition temperature T_{c0} scale at zero field is estimated to be about 2.5 K. As the magnetic field increases, the low temperature ($\ll T_{c0}$) resistance departs from the zero-field curve and at base temperature (0.38 K) first becomes

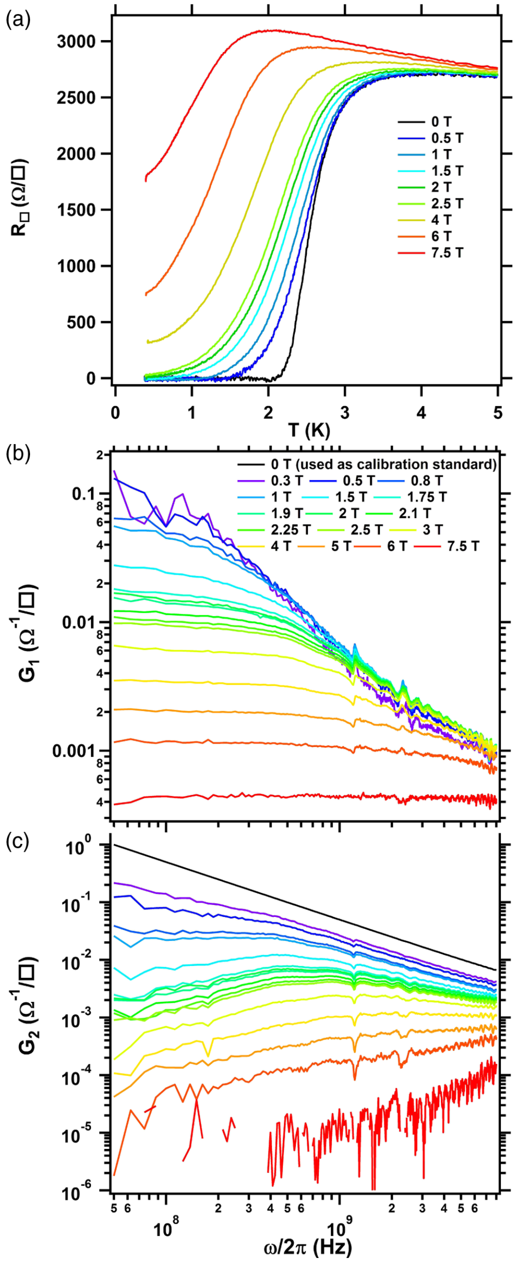

FIG. 1. (a) Temperature dependence of the dc sheet resistance under different magnetic fields. (b),(c) Complex sheet conductance at base temperature on a log-log scale. (b) Real part G_1. G_1 ($H = 0$ T) is not plotted, but can be found in Fig. 2. (c) The corresponding imaginary part.

distinguishable from zero above ≈2 T (see Fig. S3 in SM [28]), which gives us an estimate for the critical field for the superconductor-to-metal transition H_{sm}. For fields above 2 T, the temperature dependence is weak as $T \to 0$, and the resistance appears to saturate at finite values much smaller than the normal state. For 4, 6, and 7.5 T, a weak insulating dependence develops around T_{c0}, but the resistance still decreases and then levels out at low T. It is important to note that this is likely only a true phase transition in the limit of zero temperature and that any 2D superconductor in magnetic field is expected to have a small (but possibly undetectable) resistance at finite temperature.

In Figs. 1(b) and 1(c), we show the real (G_1) and imaginary (G_2) parts of the complex conductance at the base temperature at different fields on log-log plots. These data were measured in a novel broadband Corbino spectrometer [31] down to 0.38 K. Zero-field data were used as a superconducting short calibration standard by assuming the imaginary part is ideal [31]. At zero field, the real part has presumably the response of a Dirac delta function given by $G_1(\omega) = [(\pi n_s e^2)/(2m)]\delta(\omega)$, the spectral weight (e.g., the integrated area) of which corresponds to superfluid density ρ_s. By Kramers-Kronig consistency, the imaginary part $G_2(\omega) = [(n_s e^2)/(m\omega)] = [(2e^2)/(\pi\hbar)][(k_B T_\theta)/(\hbar\omega)]$.

Here one can rewrite the spectral response in terms of a phase stiffness T_θ, which is the energy scale to put a twist in the phase of the superconducting order parameter. With the presence of small external field, the real part immediately obtains a finite value at finite frequency alongside the delta function, which is retained over some range in field (the latter inferred from the vanishing dc resistance). Consistent with this, the imaginary conductance deviates from a $1/\omega$ dependence at low frequencies, which can be interpreted as suppressed long-range phase coherence. The high frequency response near 8 GHz retains its $1/\omega$ dependence but with reduced magnitude. For fields just below 1 T, G_2 is flat over almost a decade of frequencies, which further suggests that multiple conductance channels exist and a single Drude term or a $1/\omega$ dependence is insufficient to account for the low frequency response. Above H_{sm}, the imaginary part develops a broad maximum that was previously interpreted to indicate the fluctuation frequency Ω [8] on the approach to the superconductor-metal transition.

Notably, as the narrow low frequency peak in G_1 decreases in magnitude and broadens, to within experimental uncertainty, it does not shift from zero frequency. If the dissipative state above H_{sm} were a conventional high mobility metal, one would expect to observe classical cyclotron resonance whereby a finite frequency peak will be exhibited in G_1. Cyclotron resonance in which charges undergo periodic orbits in magnetic field is an essential property of conventional metals. Classically, this resonance frequency is expected to be $\omega_c = eB/m$. In a Galilean invariant system, via Kohn's theorem [40] this frequency is expected to be unchanged by interactions with an inferred mass m equal to the bare electron mass (m_e) itself. In low density 2DEGs with Fermi wavelengths much larger than the lattice constant, an approximate Galilean invariance is obtained and m is found to be the band mass [41,42]. From a minimal fitting of the complex conductance using the expression for semiclassical transport [43], we can set an upper bound on the cyclotron resonance to be less than 65 MHz at 6 T. This is an exceedingly small number and in the context of semiclassical transport implies an effective mass greater than $2500 m_e$ or internal effective field that is $1/2500$ of the applied field. Therefore, we ascertain that to

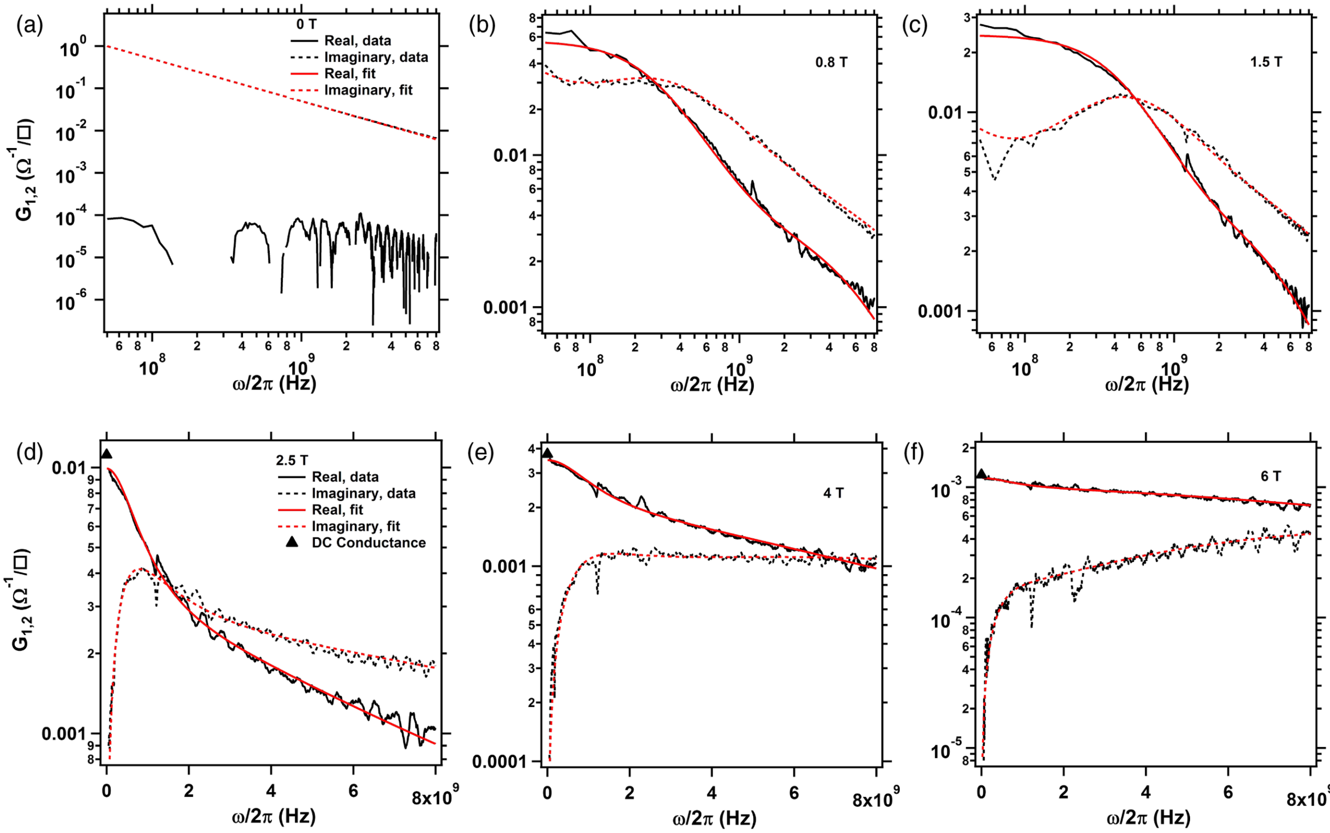

FIG. 2. Complex conductance fits for (a) 0 T, (b) 0.8 T, (c) 1.5 T, (d) 2.5 T, (e) 4 T, and (f) 6 T. Panels (a)–(c) are on log-log scales while panels (d),(e) only have the vertical axis on log scale to show dc data. The black triangle represents the comeasured dc data. Models used in the fitting are described in the text.

within our experimental uncertainty the anomalous metal state has no cyclotron resonance.

The spectra can in principle be fit to a combination of three phenomenologically assigned terms: a zero frequency delta function, a finite width peak centered at zero frequency that is characteristic of the anomalous metal, and a broader background (see SM [43] for a decomposition). In the fits these contributions are modeled as classical Lorentzians [43]. Figure 2 shows the fitting at 6 characteristic fields. In the superconducting state [Figs. 2(a)–2(c)], the imaginary conductance exhibits a diverging trend as frequency goes to zero. This observation can be related to the existence of the delta function in the real conductance. For small, but finite field (say 0.8 T), the contribution of multiple features in the spectra is apparent. In the superconducting regime, we fit the spectra to the three terms, and in the anomalous metal, only two. In the superconducting state, the presence of very small, but still finite width to the "delta" function contribution (say due to finite temperature) does not affect these fits, as any width is below the low frequency end of the spectrometer and indistinguishable from a delta function. In the anomalous metal, the delta function is obviously absent and the imaginary part extrapolates to zero at zero frequency. However, the conductance has non-Drude line shape, as demonstrated in Fig. 2(e) for 4 T where the imaginary part is almost flat from 1 to 8 GHz. These data fit well with a Drude term of the width of ∼8 GHz and a much narrower term (∼1 GHz wide) which compensates the imaginary part at higher frequencies. In the metallic regime, it is not clear if one should interpret the spectra as truly the sum of two channels, or if the transport is single channel and the fits should be considered only phenomenological.

However, given these fits it is interesting to see how the spectral weight of the various features evolves as a function of field. In phase models, the spectral weight can be identified with a phase stiffness (T_θ), which one finds by multiplying the plasma frequency [43] squared by $[\hbar/(2k_BG_Q)]$ (with the quantum of conductance $G_Q \equiv 4e^2/h$) to get the stiffness in the units of degrees kelvin. In Fig. 3, we plot the spectral weight of the delta function, the low frequency dissipative conductance, and the total as a function of field. The spectral weight of the low frequency dissipative conductance was obtained by integrating it up to 10 GHz.

With small applied field the spectral weight of the delta function falls extremely quickly with field. Although part of the spectral weight is transferred to the low frequency peak, the rest is not and presumably goes to energy scales of order the normal state scattering. By the time the metallic state is

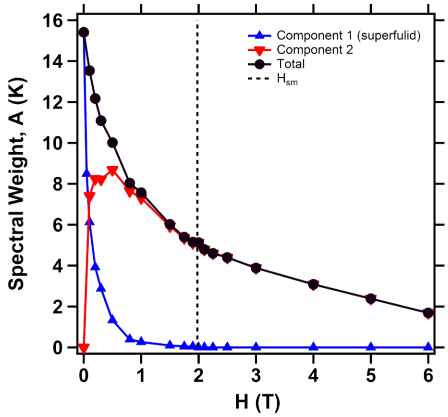

FIG. 3. Spectral weight from fits in the units of kelvin for the delta function and Drude terms. Blue and red correspond to the superfluid (delta function) and the slowly decaying component, respectively. The sum of the total spectral weight is in black. The spectral weight for component 2 at 0.05 T is not plotted because the real part of conductance falls below sensitivity. See Fig. S6 in SM [43] for details.

being approached near 2 T, the delta function spectral weight has been suppressed to very small values. The transition to anomalous metal state occurs by suppressing the delta function altogether and leaving behind the anomalous low frequency peak as the dominant conducting channel at low ω.

Our observation of no cyclotron resonance is consistent with recent transport experiments by Breznay and Kapitulnik on InO_x and TaN_x thin films [20]. They report that the Hall resistance of InO_x is indistinguishable from zero upon entering the anomalous metal regime at ≈2.0 T and remains indistinguishable from zero up to a higher field scale that is still within the anomalous metal region. We speculate that in the two component decomposition of our spectra, the narrow low frequency component exhibits no Hall response and the crossover to a regime with Hall response occurs when a broader feature (which we associate with the normal state) starts to dominate the spectrum. These results suggest that the anomalous metal state itself has a robust particle-hole symmetry. Previously it was suggested in Ref. [27] that the SIT quantum critical point had an emergent particle-hole symmetry that ensured that the pseudo-Lorentz force was zero at the transition giving zero Hall effect and a vanishing cyclotron resonance. Working from the perspective that here the critical point of the SIT broadens into a *phase*, it appears that the anomalous metal exhibits a similar symmetry. We should point out that a particle-hole symmetry as such *is not* conventionally a property of superconductors themselves. In the flux-flow regime, superconductors show a Hall response [46–48] that depends on the details of the vortex damping. Even in the fluctuation regime above T_c particle-hole symmetry appears to be broken by terms that depend on the derivative of the density of states with energy [49,50], giving the Aslamazov-Larkin superconducting fluctuation contribution a finite Hall effect. However, note that although a corresponding cyclotron resonancelike energy scale can be defined in such calculations, it is expected to be *larger* than the normal metal cyclotron resonance by a factor of $k_F l$ [51] and, moreover, does not necessarily result in a resonance [52]. Therefore, the particle-hole symmetry encountered here requires further perspective. It would be interesting to consider older theories [13,15,16,19] in the context of these newer experiments.

Our observations are in part reminiscent of the phenomenology of composite fermions of half filled Landau levels in 2D electron gasses. In such systems a cyclotron resonance mode is found that depends not on the applied magnetic field but instead an effective magnetic field B_{eff}, which is the physical applied field minus a Chern-Simons field [53]. The cyclotron resonance disappears at fields that correspond to even-denominator Landau level filling fractions [54]; however, the actual physical Hall effect remains finite. In contrast, in the present case the cyclotron resonance is absent over a range of fields and the Hall effect vanishes. Connection of the theory of composite fermions to the SIT quantum critical point and intervening metals have been recently made [18]. We point out in this regard that the density of charges at the transition is of order the number of vortices (see Fig. S8 in SM [43]).

Recently, Davison *et al.* employed a hydrodynamic approach and memory matrix formalism to calculate the dynamical conductivity of a phase fluctuating superconductor in the incoherent limit [17]. With strong time-reversal and parity symmetry breaking, the dynamical conductivity is allowed to have a *supercyclotron* resonance mode at $\omega^\star = \Omega^H - i\Omega$. This mode is allowed if the vortex current $\boldsymbol{J_v} \equiv n_v q_v \boldsymbol{v_v}$ has a downstream component with the supercurrent $\boldsymbol{J_s} \equiv \rho_s \boldsymbol{v_s}$, other than a transverse component due to the superfluid Magnus force [47]; i.e., $\boldsymbol{F_M} = -q_v \rho_s (\boldsymbol{v_s} - \boldsymbol{v_v}) \times \hat{z}\Phi_0$ for a single vortex with vorticity q_v in S.I. unit. The downstream vortex current then generates an emergent transverse electric field that bends the path of the supercurrent. Our data constrain Ω^H to be less than ~65 MHz, which implies that the downstream component is much smaller than the transverse component that dominates vortex dynamics.

In this work, we have investigated the low frequency dynamical conductivity of a thin superconducting film in the vicinity of the superconductor-to-metal transition. Remarkably, the system possesses *no* cyclotron resonance, which is a ubiquitous feature of high mobility metals composed of conventional electrons. Our observation taken with that the anomalous metal state has—over a range of fields—no Hall effect [20] shows an emergent particle-hole symmetry. This was previously considered to be a property of the SIT quantum critical point, but here is exhibited over an entire intervening phase. The anomalous metal should be considered a unique state of matter, which cannot be understood in terms of the conventional theory of metals.

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[1] H. M. Jaeger, D. B. Haviland, B. G. Orr, and A. M. Goldman, Phys. Rev. B **40**, 182 (1989).
[2] A. Yazdani and A. Kapitulnik, Phys. Rev. Lett. **74**, 3037 (1995).
[3] D. Ephron, A. Yazdani, A. Kapitulnik, and M. R. Beasley, Phys. Rev. Lett. **76**, 1529 (1996).
[4] N. Mason and A. Kapitulnik, Phys. Rev. Lett. **82**, 5341 (1999).
[5] M. A. Paalanen, A. F. Hebard, and R. R. Ruel, Phys. Rev. Lett. **69**, 1604 (1992).
[6] J. A. Chervenak and J. M. Valles, Phys. Rev. B **61**, R9245 (2000).
[7] Y.-H. Lin, J. J. Nelson, and A. M. Goldman, Phys. Rev. Lett. **109**, 017002 (2012).
[8] W. Liu, L. D. Pan, J. Wen, M. Kim, G. Sambandamurthy, and N. P. Armitage, Phys. Rev. Lett. **111**, 067003 (2013).
[9] H. S. J. van der Zant, F. C. Fritschy, W. J. Elion, L. J. Geerligs, and J. E. Mooij, Phys. Rev. Lett. **69**, 2971 (1992).
[10] Z. Han, A. Allain, H. Arjmandi-Tash, K. Tikhonov, M. Feigel'Man, B. Sacépé, and V. Bouchiat, Nat. Phys. **10**, 380 (2014).
[11] S. Eley, S. Gopalakrishnan, P. M. Goldbart, and N. Mason, Nat. Phys. **8**, 59 (2012).
[12] T. Schneider, A. D. Caviglia, S. Gariglio, N. Reyren, and J.-M. Triscone, Phys. Rev. B **79**, 184502 (2009).
[13] D. Das and S. Doniach, Phys. Rev. B **60**, 1261 (1999).
[14] A. Kapitulnik, N. Mason, S. A. Kivelson, and S. Chakravarty, Phys. Rev. B **63**, 125322 (2001).
[15] D. Dalidovich and P. Phillips, Phys. Rev. Lett. **89**, 027001 (2002).
[16] V. M. Galitski, G. Refael, M. P. A. Fisher, and T. Senthil, Phys. Rev. Lett. **95**, 077002 (2005).
[17] R. A. Davison, L. V. Delacrétaz, B. Goutéraux, and S. A. Hartnoll, Phys. Rev. B **94**, 054502 (2016).
[18] M. Mulligan and S. Raghu, Phys. Rev. B **93**, 205116 (2016).
[19] B. Spivak, P. Oreto, and S. A. Kivelson, Phys. Rev. B **77**, 214523 (2008).
[20] N. P. Breznay and A. Kapitulnik, Sci. Adv. **3**, e1700612 (2017).
[21] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.120.167002 for recent investiagtions of SIT, which includes Refs. [22–26].
[22] G. Sambandamurthy, L. W. Engel, A. Johansson, E. Peled, and D. Shahar, Phys. Rev. Lett. **94**, 017003 (2005).
[23] G. Sambandamurthy, A. Johansson, E. Peled, D. Shahar, P. Björnsson, and K. Moler, Europhys. Lett. **75**, 611 (2006).
[24] M. Steiner and A. Kapitulnik, Physica (Amsterdam) **422C**, 16 (2005).
[25] R. W. Crane, N. P. Armitage, A. Johansson, G. Sambandamurthy, D. Shahar, and G. Grüner, Phys. Rev. B **75**, 094506 (2007).
[26] V. Gantmakher and M. Golubkov, JETP Lett. **73**, 131 (2001).
[27] M. P. Fisher, Physica (Amsterdam) **177A**, 553 (1991).
[28] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.120.167002 for magneto resistance, and the lack of reentrance behavior which indicates the film is not granular. This includes Refs. [29,30].
[29] G. Sambandamurthy, L. W. Engel, A. Johansson, and D. Shahar, Phys. Rev. Lett. **92**, 107005 (2004).
[30] R. Crane, N. P. Armitage, A. Johansson, G. Sambandamurthy, D. Shahar, and G. Grüner, Phys. Rev. B **75**, 184530 (2007).
[31] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.120.167002 for details about this technique and calibration procedures, which includes Refs. [8,32–39].
[32] W. Liu, M. Kim, G. Sambandamurthy, and N. P. Armitage, Phys. Rev. B **84**, 024511 (2011).
[33] M. Scheffler and M. Dressel, Rev. Sci. Instrum. **76**, 074702 (2005).
[34] W. Liu, L. Pan, and N. Armitage, Rev. Sci. Instrum. **85**, 093108 (2014).
[35] J. C. Booth, D. H. Wu, S. B. Qadri, E. F. Skelton, M. S. Osofsky, A. Piqué, and S. M. Anlage, Phys. Rev. Lett. **77**, 4438 (1996).
[36] M. Lee and M. L. Stutzmann, Phys. Rev. Lett. **87**, 056402 (2001).
[37] T. Ohashi, H. Kitano, I. Tsukada, and A. Maeda, Phys. Rev. B **79**, 184507 (2009).
[38] M. Scheffler, Ph.D. thesis, University of Stuttgart, 2004.
[39] W. Liu, Ph.D. thesis, The Johns Hopkins University, 2013.
[40] W. Kohn, Phys. Rev. **123**, 1242 (1961).
[41] G. Abstreiter, P. Kneschaurek, J. Kotthaus, and J. Koch, Phys. Rev. Lett. **32**, 104 (1974).
[42] T. Kennedy, R. Wagner, B. McCombe, and D. Tsui, Phys. Rev. Lett. **35**, 1031 (1975).
[43] See Supplemental Material at http://link.aps.org/supplemental/10.1103/PhysRevLett.120.167002 for fitting models and techniques, which includes Refs. [44,45].
[44] G. Dresselhaus, A. Kip, and C. Kittel, Phys. Rev. **98**, 368 (1955).
[45] A. Kuzmenko, Rev. Sci. Instrum. **76**, 083108 (2005).
[46] S. J. Hagen, C. J. Lobb, R. L. Greene, M. G. Forrester, and J. H. Kang, Phys. Rev. B **41**, 11630 (1990).
[47] P. Nozières and W. Vinen, Philos. Mag. **14**, 667 (1966).
[48] T. W. Jing and N. P. Ong, Phys. Rev. B **42**, 10781 (1990).
[49] H. Fukuyama, H. Ebisawa, and T. Tsuzuki, Prog. Theor. Phys. **46**, 1028 (1971).
[50] N. P. Breznay, K. Michaeli, K. S. Tikhonov, A. M. Finkel'stein, M. Tendulkar, and A. Kapitulnik, Phys. Rev. B **86**, 014514 (2012).
[51] K. Michaeli, K. S. Tikhonov, and A. M. Finkel'stein, Phys. Rev. B **86**, 014515 (2012).
[52] A. Finkelstein (private communication).
[53] B. Halperin, *Perspectives in Quantum Hall Effects: Novel Quantum Liquids in Low-Dimensional Semiconductor Structures* edited by S. D. Sarma and A. Pinczuk (Wiley VCH, Weinheim, 2004), pp. 225–263.
[54] I. Kukushkin, J. Smet, K. Von Klitzing, and W. Wegscheider, Nature (London) **415**, 409 (2002).