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Weak antilocalization of composite fermions in graphene

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We demonstrate experimentally that composite fermions in monolayer graphene display weak antilocalization. Our experiments deal with fractional quantum Hall (FQH) states in high-mobility, suspended graphene Corbino disks in the vicinity of \( v = 1/2 \). We find a strong temperature dependence of conductivity \( \sigma \) away from half filling, which is consistent with the expected electron-electron interaction-induced gaps in the FQH state. At half filling, however, the temperature dependence of conductivity \( \sigma(T) \) becomes quite weak, as anticipated for a Fermi sea of composite fermions, and we find a logarithmic dependence of \( \sigma \) on \( T \). The sign of this quantum correction coincides with the weak antilocalization of graphene composite fermions, indigenous to chiral Dirac particles.

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I. INTRODUCTION

The fractional quantum Hall (FQH) state is a many-body phenomenon where fractionally charged elementary excitations lead to quantization of the Hall conductance at a fractional filling factor \( v = nh/(eB) \) at a carrier density \( n \) and magnetic field \( B \) [1]. The generation of these incompressible liquid states requires a large Coulomb interaction energy compared with the disorder potential, putting strict requirements on temperature, the quality of the two-dimensional electron gas (2DEG), and the strength of the magnetic field. Owing to reduced screening in atomically thin graphene, the electrons in graphene interact with a larger Coulomb interaction energy than electrons in semiconductor heterostructures, providing an excellent setting for studies of FQH states and their description in terms of composite fermions [2–4].

The composite fermion theory [2] and composite fermion Chern-Simon (CFCS) theory [5–8] have been very successful in outlining a unified picture of the fractional quantum Hall effect. Lopez and Fradkin [5] showed that the problem of interacting electrons moving in 2D in the presence of an external magnetic field is equivalent to a fermion system, described by a Chern-Simon gauge field, where electrons are bound to an even number of vortex lines. Fluctuations in the gauge field were soon realized to have a strong influence on the quantum correction of the composite fermion conductivity [6]. Subsequently, a Fermi-liquid type of theory was proposed for the half-filled Landau level [7] where various observables in the low-temperature limit are described in terms of Fermi-liquid parameters [8], involving most notably the effective mass \( m^* \) for composite fermions, which is expected to have a strong enhancement near half filling.

Extensive experimental evidence exists in favor of a weakly interacting Fermi sea of composite fermions effectively in a zero magnetic field at \( v = 1/2 \). Transport anomalies in the lowest Landau level of two-dimensional electrons at half filling were observed by Jiang et al. [9]. Distinct features related with compressibility in surface-acoustic-wave propagation on high-quality AlGaAs/GaAs heterostructures were observed by Willett et al. [10–12]. Furthermore, resonances at fields where the classical cyclotron orbit becomes commensurate with a superlattice have been found [13,14]. A strong enhancement of the composite fermion mass near \( v = 1/2 \) has been observed in experiments on similar 2DEG heterostructures [15,16]. A logarithmic temperature dependence of conductivity at half-integer filling factors has been observed, which has been interpreted to point towards residual interactions between composite fermions [17–19].

Graphene, owing to the pseudospin originating from the two-atom basis, may display reduced backscattering that will result in weak antilocalization (WAL). This has been observed for Dirac particles in graphene at small magnetic fields [20], and similar WAL behavior might carry over to graphene composite fermions instead of interaction-induced quantum corrections to conductivity [21]. In our present work, we demonstrate experimentally the presence of chiral, Dirac-particle-like composite fermions in monolayer graphene by measuring quantum corrections to conductivity for a half-filled quantum Hall state in magnetic fields of \( B = 5–9 \) T. By using a micron-wide suspended Corbino disk with high mobility, we find evidence of weak antilocalization behavior at half filling, in spite of the inherent field fluctuations in the composite fermion system. Our results display typical logarithmic temperature corrections for two-dimensional weak localization (WL), but with an opposite sign in comparison to previous observations in 2DEG experiments [18,19].

Interaction and weak localization corrections to conductivity in graphene are influenced by the pseudospin introduced by the two-atom basis of the hexagonal lattice. The pseudospin leads to an additional Berry phase of \( \pi \) in the backscattering wave interference, resulting in reduced backscattering, or in other words, antilocalization [22–24]. True time reversal symmetry (TRS) of graphene requires intervalley scattering, which makes its weak localization phenomena intriguing. In fact, low-field magnetoresistance in graphene may be either positive or negative, depending on the strength of the intervalley and dephasing scattering, as well as on the rate for TRS breaking in a single valley [20,24,25].

In the composite fermion picture, a charge carrier is combined with two flux lines, vortices. The underlying valley
structure will be reflected on the graphene composite fermions, and we expect scattering processes with similar relations among the graphene composite fermions as for graphene Dirac particles at small fields. Consequently, quantum corrections to conductivity at \( v = 1/2 \) can be either positive (WAL) or negative (WL) depending on the rates of intervalley, dephasing, and single-valley TRS breaking scattering for composite fermions denoted as \( \tau_{\text{d}}^{-1} \), \( \tau_{\text{w}}^{-1} \), and \( \tau_{\text{c}}^{-1} \), respectively. Even though fluctuations due to Chern-Simon fields will enhance the dephasing rate of composite carriers, we anticipate the tendency towards weak antilocalization to remain, which reflects the presence of a Berry phase due to the Dirac nature of these composite carriers.

Recently, Corbino geometry was used in graphene experiments which focused on studies of magnetoconductance in the quantum Hall regime [26,27]. These previous measurements on graphene Corbino disks have failed to show any FQH states, perhaps due to the strong charge inhomogeneity induced by the substrate. Our experiments at mK temperatures, however, display a multitude of FQH states in the lowest Landau level which are clearly visible in magneto- and transconductance measurements [28]. In our devices, conductance is governed by bulk properties instead of the chiral edge states of regular quantum Hall bars. Consequently, a current-annealed, suspended graphene Corbino disk forms an excellent system to study carrier dynamics in the free Fermi sea of composite fermions at half filling.

II. METHODS

The samples considered here are denoted by EV2 (\( r_2 = 1900 \) nm, \( r_1 = 750 \) nm), EV3 (\( r_3 = 1600 \) nm, \( r_1 = 400 \) nm), and C2 (\( r_2 = 1400 \) nm, \( r_1 = 450 \) nm), where \( r_0 \) and \( r_i \) denote the outer and inner radii of the Corbino disks. Our sample fabrication is explained in detail in Ref. [28]. First, we exfoliated graphene on LOR (lift-off resist) using a heat-assisted exfoliation technique to maximize the size of the exfoliated flakes [29]. The contacts were deposited by employing poly(methyl methacrylate) (PMMA)-based lithography in a scheme developed further from the ideas presented in Ref. [30]. A strongly doped silicon Si\( ^{++} \) substrate with 285 nm of thermally grown SiO\( _2 \) was used as a global back gate. Annealing of samples on LOR was typically performed at low temperatures using a bias voltage of 1.6 \( \pm \) 0.1 V, which is comparable with our HF etched, rectangular two-lead samples [31,32].

In our suspended devices (see Fig. 1), the Corbino disk is supported only by the inner and outer leads made of Au/Cr [28]. The conductivity was calculated from the measured conductance \( G \) using \( \sigma = \frac{G}{e^2} \ln(\frac{r_o}{r_i}) \). The field-effect mobility was determined using the equation \( \mu = \frac{\sigma - \sigma_0}{e n_0} \), after subtracting out the contact resistance from \( \sigma \) and the measured minimum conductivity \( \sigma_0 \) at the Dirac point (\( V_g = V_g^d \approx -2 \) V). Our samples had \( \mu \approx 1.3-2 \times 10^6 \) cm\(^2\)/V/s at low \( n \). The residual charge density \( n_0 \approx 1-2 \times 10^9 \) cm\(^{-2}\) for the best regions of the samples was identified by looking for a crossover between constant and power-law behavior in log \( G \) vs log \( n \) traces in the unipolar transport regime. Contact doping [32] is estimated to correspond to a charge density of \( n_c \approx 5 \times 10^{11} \) cm\(^{-2}\) under the contact metal. Details concerning the determination of these basic parameters are given in Ref. [28].

![FIG. 1. False-color scanning electron micrograph of a typical sample. The teal circular ring is the graphene sheet, yellow parts denote the gold contacts, dark green marks the LOR support, and the underlying substrate appears as gray.](image)

![FIG. 2. The magnetocconductance of sample EV2 displayed as a Landau fan diagram on the B vs gate voltage \( V_g \) plane; data measured at \( f = 3.333 \) Hz using ac peak-to-peak excitation current \( I_{ac} \) = 0.1 nA at 20 mK. The region of half filling is given by the slope 0.64 T/V starting from the zero-field Dirac point at \( V_g = -2.3 \) V.](image)
waiting times up to 5 s (for \(G > 2\times10^{-8} \text{ S}\)). Since the symmetry of the Corbino geometry enforces the azimuthal electric field to zero, our experiment is sensitive only to the longitudinal conductivity \(\sigma_{xx}\) of the sample.

### III. MAGNETOCONDUCTANCE AROUND \(1/2\) FILLING

Our data on the electronic conductivity \(\sigma_{xx}\) versus the gate-swept filling factor \(\nu = \frac{\hbar n}{eB}\) in sample EV3 are displayed in Fig. 3(a) for several temperatures in the range \(T = 0.01–4.2\ \text{K}\) measured at \(B = 5\ \text{T}\). Similar data from the sample EV2 in the temperature range \(T = 0.03–6\ \text{K}\) at \(B = 9\ \text{T}\) are displayed in Fig. 3(b). In both cases the half-filling regime is characterized by the weakest temperature dependence of \(\sigma_{xx}\), which is in accordance with the absence of any energy gap in this regime. Only a linear background variation of \(\sigma_{xx}(v)_{\text{bg}}\) around \(v = 1/2 = \frac{\hbar n}{eB_{1/2}}\) is observed at the highest measured temperatures, which indicates washing out of quantum corrections to conductance by thermal fluctuations.

The effective field seen by the composite fermions is given by \(\vec{B} = B - B_{1/2}\), which amounts to \(\vec{B} = 0\) at \(v = 1/2\). From the data in Fig. 3 we see that the magnetoconductance around \(\vec{B} = 0\) is negative at the lowest mK temperatures. This negative magnetoconductance with respect to \(\vec{B}\) applies to all our samples at the magnetic fields \(5\ \text{T} < B < 9\ \text{T}\). Foremost, however, this work deals with the transport behavior at \(B = 9\ \text{T}\), at which the composite fermion effects and the FQH states are the clearest. The main characteristics of the data around half filling at \(9\ \text{T}\) can be summarized by \(d\sigma_{xx}/dT > 0\) and \(d\sigma_{xx}/d|\vec{B}| < 0\). The regime of \(d\sigma_{xx}/d|\vec{B}| < 0\) covers only the range \(v \approx 0.49–0.51\) which corresponds to effective fields \(|\vec{B}| < 200\ \text{mT}\).

The temperature dependence of \(\sigma_{xx}\) can also be employed to determine the energy gaps \(E_{\nu}/k_B\) around \(v = 1/2\), for which we obtain, at \(5\ \text{T}\), \((B - 2h\nu/e)\times3\ \text{K}/T - \Gamma_0\) and \((B - 2h\nu/e)\times6\ \text{K}/T - \Gamma_0\) with \(\Gamma_0 \approx 1–2\ \text{K}\) for particles and holes, respectively. Using the slopes of \(E_{\nu}/dv\) and the relation \(R_{\text{eff}} = h/(e\nu)(1 - 2\nu)\) at a fixed magnetic field, we obtain for the effective masses 0.4\(m_e\) and 0.8\(m_e\), at fields at \(B_{\text{eff}} = 0\) and at \(B_{\text{eff}} < 0\), respectively, where \(m_e\) denotes the electron mass [34]. In graphene, the cyclotron mass increases with charge density as \(n^{1/2}\). We find an even faster, approximately linear, increase in the composite fermion mass \(\nu^*\) with \(n\) at \(B_{\text{eff}} < 0\) while this is absent at \(B_{\text{eff}} > 0\).

For a comparison of the magnetoconductivity in Fig. 3 at small \(\vec{B}\) with phenomena observed at small \(B\), we note that there may be either positive or negative magnetoresistance in graphene around \(B = 0\) [24]. The sign of the magnetoresistance has been demonstrated to depend fundamentally on the parameters \(\tau_e/\tau_i\) and \(\tau_e/\tau_i^*\); when both of these values are \(\lesssim 1\), weak antilocalization is preferred and \(d\sigma_{xx}/dT < 0\) [20]. In Ref. [20], small ratios were achieved at elevated temperatures \(T \gg 4.2\ \text{K}\). In our suspended sample, we were able to observe WAL even at 20 mK, and the range of weak antilocalization magnetoconductance was found to be within \(\delta B = 10–20\ \text{mT}\) [see the inset of Fig. 3(b)]. At high \(B\), the magnetoconductance measurements as a function of \(|\vec{B}|\) are complicated by the gauge field variation \(\sim 100\ \text{mT}\) originating from charge variation via the composite fermion flux attachment.

### IV. DISCUSSION

Figure 4 displays the temperature dependence of conductivity \(\sigma_{xx}\) for samples EV2 and EV3 at \(v = 1/2\). The data indicate \(\sigma_{xx} = \lambda T / \log T + \text{const}\) with an increase in \(\sigma_{xx}\) with temperature where \(\lambda = 0.005\). For our third sample C2, we found similar behavior with \(\lambda = 0.008\). Altogether, the observed quantum correction is quite large and it amounts to \(\sim 20\%\) of \(\sigma_{xx}\) over the measured range.

To relate the measured conductivity to the properties of composite fermions, we follow the treatment given in Ref. [2] for 2DEG composite fermions (for alternate views, see also Refs. [35–37]). Then, the connection between electronic conductivity \(\sigma_{xx}\) in units \(e^2/h\) and the (scaled) composite fermion
The combined effect of the diffusive motion of composite fermions and the gauge field interaction has been studied by Khveshenko [38,39] and by Mirlin and Wölfle [40]. In addition to the weak localization/antilocalization contributions, the magnetoresistance may be influenced by the localization of charge carriers due to the magnetic field fluctuations that lead to cyclotron radii which allow closed loops for orbiting particles to be formed. Reference [41] considers magnetic field fluctuations of magnitude of \( B_0 \) and correlation length \( d \) with a scaled correlation function specified as \( \langle B(r)B(0) \rangle = B_0^2 \langle F(r/d) \rangle \). Depending on whether \( d/R_0 = \alpha \ll 1 \), i.e., the ratio of the correlation radius to the cyclotron radius \( R_0 \) in the field \( B_0 \), their analysis deals with weak or strong random magnetic field (RMF) fluctuations, respectively. We estimate the correlation radius \( d \sim 100 \) nm as it is related to the size of charge puddles which in high-quality graphene may be on the order of 100 nm [42]. In the text below we will argue that \( B_0 \sim 100 \) mT is supported by our results. Consequently, \( \alpha \sim 2 \) and we are in the intermediate-\( \alpha \) regime, but closer to strong RMF fluctuations. Based on an analysis of snake states and localized trajectories [43] at \( \alpha \gg 1 \), Ref. [41] obtains for conductance \( \sim k_F d/\alpha^{1/2} \). Thus, \( \sigma_{xx} \) should enhance approximately as \( n^{1/2} \) as \( B \) is increased at \( \nu = 1/2 \), provided that the correlation radius \( d \) remains unchanged. Our data for samples EV2 and EV3 at fixed \( \nu = 1/2 \) follow a sum of a constant-linear increase in \( \sigma_{xx} \) with increasing \( n \) above \( \sim 3 \) and \( \sim 4 \) T, respectively, which is in agreement with a linearized form of the theoretical \( n^{1/2} \) dependence.

At \( \nu \neq 1/2 \), there will be a nonzero average field \( \overline{B} \) seen by the composite fermions. If \( \overline{B} \) is so large that the cyclotron radius is clearly smaller than \( d \), the majority of the trajectories become localized. The localization of composite fermions leads to a reduction in the number of the charge carriers, which diminishes conduction across the Corbino sample. This will lower the conductivity \( \sigma_{xx} \) of the sample as seen in Fig. 4 around \( \nu = 1/2 \). The amount of localized trajectories for \( \alpha \sim 1 \) increases quickly when the field \( \overline{B} \) seen by the composite fermions becomes \( \sim B_0 \). Thus, a decrease in conductivity is obtained at effective fields \( \overline{B} > B_0 \) [41]. By looking at the decay range of \( \sigma_{xx} \) around \( \nu = 1/2 \) at 9 T in Fig. 3(b), we obtain an estimate of \( B_0 \sim 100 \) mT. This is quite close to the RMF variation magnitude given by \( n_0^{2/3} \sim 75–150 \) mT. Therefore, we conclude that the negative, classical magnetococonductance around \( \nu = 1/2 \) dominates over the weak antilocalization corrections seen in \( \sigma_{xx}(T) \) at \( \nu = 1/2 \).

Our suspended graphene samples always contain built-in nonuniform strain, which can result in visible frozen ripples at room temperature. This strain is modified by the applied gate voltage, which can induce additional rippling around the perimeter if the graphene sheet is able to slide against the metallic contact. A variation in strain in suspended graphene will lead to locally varying pseudomagnetic fields that will lead to dephasing and TRS-breaking scattering within one cone (i.e., shorten \( \tau_\rho \) and \( \tau_\sigma \)). According to Ref. [25], nonuniform strain is the main contributor to both elastic scattering time \( \tau \) and \( \tau_\sigma \) and \( \tau \sim \tau_\sigma \) within a factor of 2–3. This scenario explains the observation of WAL also in our sample at \( B \sim 0 \).

Using similar arguments as above, Chern-Simon fluctuations governing \( \rho_{xx}^{\mathrm{CF}} \) will strongly enhance scattering rates
τ−1, τ−1, and τ−1 at ν = 1/2 at high fields. Typically, it is argued that weak localization effects cannot be observed in the presence of Chern-Simon field fluctuations that have an infrared divergence in the density of states of low-frequency excitations at small wave vectors q [8]. In our sample, however, there is a rather high infrared cutoff due to a small geometrical size and low carrier density. By taking a cutoff qmin ≃ 1/(τi − τr) and kF ≃ \sqrt{2πn}, we obtain for the ratio qmin/kF ≃ 0.01 at n ≃ 10^{11}\ cm^{-2}. Consequently, the Chern-Simon fields do not kill the weak localization effects of graphene composite fermions, but rather just reduce the dephasing time (and the ratios τϕ/τi and τϕ/τi) which favors weak antilocalization behavior among graphene Dirac particles [20].

V. CONCLUSIONS

In conclusion, we have investigated fractional quantum Hall states in a suspended graphene Corbino disk around half filling. We find a weak logarithmic temperature dependence of conductivity, the sign of which indicates a weak antilocalization behavior of graphene composite fermions. These observations with nearly zero effective field acting on composite fermions can be understood by drawing a parallel with the weak antilocalization theories of graphene Dirac particles at small magnetic fields and using classical arguments to analyze the role of gauge field fluctuations in magnetoconductance.

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