Signatures of a Novel Fermi Liquid in a Two-Dimensional Composite Particle Metal

H. C. Manoharan,[1,2] M. Shayegan,[1,2] and S. J. Klepper[2]

[1]*Department of Electrical Engineering, Princeton University, Princeton, New Jersey 08544*

[2]*Francis Bitter National Magnet Laboratory, Massachusetts Institute of Technology, Cambridge, Massachusetts 02139*

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Quasiparticle excitation gaps of the $\nu = p/(2p + 1)$ fractional quantum Hall states in a low-disorder two-dimensional hole system are observed to scale linearly with $(e^2/\epsilon\ell_B)/(2p + 1)$, predicted by the Chern-Simons gauge field theory if a hole-flux composite particle Fermi liquid forms around filling $\nu = \frac{1}{2}$ (ℓ_B is the magnetic length). The effective mass and quantum lifetime of these new particles reveal a remarkable renormalization structure, including a strongly diverging mass and enhanced scattering close to the Fermi surface, suggesting a novel class of *marginal* Fermi liquid behavior.

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The transmutability of particle statistics, unique to two-dimensional (2D) systems, underlies several recent theoretical treatments of the complex phenomena observed in these systems at high magnetic field. Jain [1] and Halperin, Lee, and Read [2] have exploited a singular gauge transformation that preserves Fermi statistics and binds $2m$ flux quanta ($2m\Phi_0$, where m is an integer and $\Phi_0 \equiv h/e$ is the quantum of flux) to each electron. Doing so provides an elegant mapping of the fractional quantum Hall effect (FQHE) at Landau level (LL) filling factor $\nu = p/(2mp + 1)$ onto the integral quantum Hall effect (IQHE) at $\nu^* = |p|$ (p integer), where ν^* is the quasi-LL filling factor of the electron-flux "composite fermions" (CFs). These new particles experience an effective magnetic field $B^* = B - 2mn_e\Phi_0 = B - B_{\nu=1/2m}$ (n_e is the electron density) and therefore experience zero net magnetic field at exactly half filling. Here they are expected to form a Fermi surface and have a renormalized mass dependent only on electron-electron interaction.

Recent experiments on 2D electron systems (2DESs) in GaAs heterostructures [3–8] seem consistent with this picture. In this Letter, we present the intriguing behavior of a new system comprised of *holes,* where the larger hole band mass $m_h^* \simeq 0.38m_e$ (m_e is the free electron mass) leads to significant LL mixing into the fundamental many-body electronic states of the system [9,10]. The data demonstrate the influence of LL mixing on the renormalization of the bare current-carrying particles into particle-flux composites, and furthermore reveal several surprises that point to an unconventional Fermi liquidlike state of these CFs around the half-filled LL.

We concentrate on the temperature (T) dependence of R_{xx} extrema around $\nu = \frac{1}{2}$ in a low-disorder 2D hole system (2DHS) of density $n_h = 1.6 \times 10^{11}\ \mathrm{cm}^{-2}$ and low-T mobility $\mu \simeq 1 \times 10^6\ \mathrm{cm^2/V\,s}$, in the range $0.04 \lesssim T \lesssim 1.1$ K. The sample, grown by molecular beam epitaxy on an undoped (311)A substrate, consists of a 200 Å GaAs quantum well surrounded by $Al_{0.3}Ga_{0.7}As$ spacer layers and Si-doped regions. Ohmic contacts to the confined hole layer were made by alloying eutectic In:Zn in a van der Pauw geometry. Together with a calibrated RuO_2 resistor with known magnetoresistance corrections used for thermometry, the sample was immersed in the dilute phase of a $^3He/^4He$ dilution refrigerator inside a superconducting/Bitter hybrid magnet system.

Figure 1(a) shows R_{xx} vs B around $\nu = \frac{1}{2}$ at $T \simeq 40$ mK. The density of the sample places $B_{\nu=1/2}$ at 13.2 T, the deviation from which is labeled by the B^* axis. Minima (maxima) are marked by integer (half-integer) p. The exceptional quality of this 2DHS is manifested in the high-order FQH states up to $p \simeq \pm 5$ around $\nu = \frac{1}{2}$ as well as states at $\nu = \frac{2}{7}$ and $\frac{3}{11}$ at higher field. Our measurements focus on the simplest form of CF renormalization, $m = 1$, corresponding to two flux quanta per hole. As alluded to above, the behavior around $\nu = \frac{1}{2}$ ($B^* = 0$) can be linked to that around $B = 0$ by canceling the average fictitious gauge flux with the externally applied magnetic flux; this results in a new average flux per particle, $1/p = 1/\nu - 2$, mapping the resistivity oscillations (i.e., the FQH effect) at high field around $\nu = \frac{1}{2}$ onto the low-field Shubnikov–de Haas (SdH) oscillations around $B = 0$. As such, the minima at $\nu = p/(2p + 1)$ correspond to filled quasi-LLs of CFs with effective filling $\nu^* = |p|$, where $p = n_h\Phi_0/B^*$ (in analogy to $\nu = n_h\Phi_0/B$ at low field). Maxima should then occur when p is half integer.

We now examine the T dependence of the R_{xx} minima corresponding to FQH states at $\nu = p/(2p + 1)$, with integer p up to $+4$ ($\nu = \frac{4}{9}$) and -5 ($\nu = \frac{5}{9}$). By plotting $\log_{10}R_{xx}$ vs $1/T$ [see Fig. 1(b)], we extract the quasiparticle excitation energy gaps Δ_ν of these incompressible quantum liquid states from fits to the activated regions, using $R_{xx} \propto \exp(-\Delta_\nu/2T)$. As shown in Fig. 2, these energy gaps obey a remarkable scaling law

$$\Delta_\nu = \frac{\kappa}{|2p + 1|}\frac{e^2}{\epsilon\ell_B} - \Gamma, \qquad (1)$$

where ϵ is the dielectric constant, $\ell_B \equiv \sqrt{\hbar/eB}$ is the magnetic length, and κ and Γ are constants. In the plot of Δ_ν vs $(e^2/\epsilon\ell_B)/(2p + 1)$, linear fits from both positive and negative B^* yield nearly identical slopes and intercepts, from which we obtain $\kappa \simeq 0.07$ and $\Gamma \simeq 1.4$ K.

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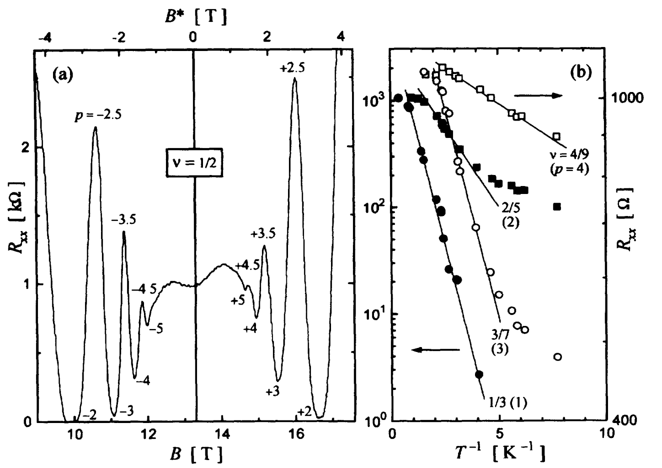

FIG. 1. (a) Magnetoresistance R_{xx} of the 2DHS vs applied magnetic field B and CF effective field $B^* = B - B_{\nu=1/2}$. (b) FQH energy gap analysis for $1 \le p \le 4$. Filled (open) symbols correspond to the left (right) axis.

The observation of the energy gap scaling (1) is an impressive finding for the following reasons. In the simplest interpretation of the CF theory, the incompressibility gaps of the FQH liquids map precisely onto the "single-particle" cyclotron energy gaps of the corresponding IQH states of CFs. One would then expect $\Delta_\nu \propto \hbar\omega_c^*$, where $\omega_c^* \equiv eB^*/m^*$ and m^* is the effective mass of the CFs. This interpretation was used in Ref. [4] to extract m^* for CFs in 2DESs, assuming a constant m^* for each sign of B^*; the fits there yielded differing values of mass depending on both B and the sign of B^*. The Chern-Simons gauge field theory [2], however, explicitly postulates modifications to this simple interpretation. First, since m^* is a result of electron-electron interaction, it cannot remain rigorously constant as B varies. In fact, the theory predicts from dimensionality arguments that

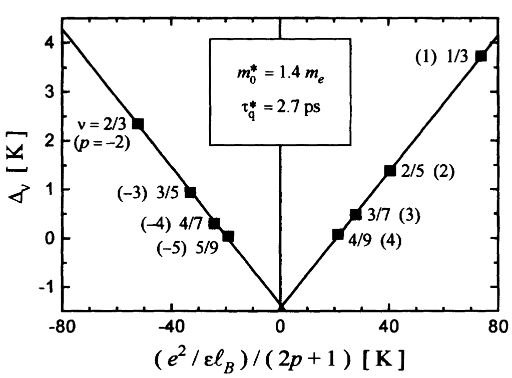

FIG. 2. Observed scaling of the quasiparticle excitation gaps of the FQH states around $\nu = \frac{1}{2}$. The estimated accuracy of the measured gaps is $\pm 10\%$ at $p = 1$, ± 2, -3; $\pm 20\%$ at $p = 3$, -4; and $\pm 50\%$ at $p = 4$, -5. Solid lines are linear fits to the data.

$$m^* = \left(\frac{\hbar}{e}\right)^2 \frac{\epsilon}{\kappa \ell_B} = \left(\frac{\Phi_0}{2\pi}\right)^{3/2} \frac{\epsilon}{\kappa}\sqrt{B}. \qquad (2)$$

Furthermore, the requirements of particle-hole symmetry in the FQHE posit the equivalency of states at p and $-(p+1)$. These two considerations predict directly the observed scaling (1). We therefore interpret this observation as strong evidence that a hole-flux composite particle Fermi liquid forms around $\nu = \frac{1}{2}$.

We can make several inferences about these renormalized fermions. Our value of κ, using Eq. (2) and $\epsilon = 13.1$, gives a CF effective mass at $B^* = 0$ of $m_0^* \simeq 1.4 m_e$ for this density. The observation of the gap scaling behavior then provides direct experimental evidence for a field-dependent mass $m^* = m_0^*(2\nu)^{-1/2} \propto \sqrt{B}$, *unlike any conventional Fermi liquid.* One may still think of the energy gaps (1) as CF cyclotron gaps $\hbar\omega_c^*$ corrected for disorder, $\Delta_\nu = \hbar e B^*/m^* - \Gamma$, as long as m^* is recognized to be field dependent in the form of Eq. (2). Interpreted in this way, Γ is the collision broadening of the CF LLs [4]. We may then estimate the quantum lifetime τ_q^* of CFs at $B^* = 0$ as $\tau_q^* = \hbar/2\Gamma \simeq 2.7$ ps, somewhat smaller than the transport lifetime $\tau_{\rm tr}^* = m^*/n_h e^2 \rho \simeq 7.5$ ps computed from the resistivity ρ at $\nu = \frac{1}{2}$, but similar to $\tau_q \simeq 2.7$ ps for bare holes near $B = 0$ deduced from the vanishing of low-field SdH oscillations. Since $\tau_q^*/\tau_{\rm tr}^* \gg \tau_q/\tau_{\rm tr}$ ($\tau_{\rm tr} \simeq 200$ ps from ρ at $B = 0$), we conclude that short-range (large-angle) scattering dominates for the hole-flux CFs around $B^* = 0$ in contrast to small-angle scattering for bare holes around $B = 0$, while the two types of particles scatter at the same rate (similar results for electrons were found in Refs. [4] and [6]). This can be explained by the nature of the impurities "seen" by the relevant particles: the bare holes feel a long-range perturbation from the setback ionized dopants, while the CFs are presumably more strongly scattered by the effects of these impurities on the renormalization dynamics, i.e., "magnetic" impurities in the form of gauge field fluctuations.

The data of Fig. 2 allow us to make an important conclusion regarding the effects of LL mixing, finite layer thickness, and disorder on the FQH energy gap scaling and CF mass renormalization. Consider the most fundamental FQH state at $\nu = \frac{1}{3}$ ($p = 1$). Since Γ, and thus disorder effects, may be independently determined from the observed scaling, we can use the extracted value of $\kappa \simeq 0.07$ to predict $\Delta_{1/3} \sim 0.023 e^2/\epsilon\ell_B$ in this 2DHS, about 4 times weaker than $\Delta_{1/3} \sim 0.1 e^2/\epsilon\ell_B$ theoretically expected [11] for ideal 2D systems, and approximately half as strong as $\sim 0.05 e^2/\epsilon\ell_B$ observed in high-mobility 2DESs [4,12]. The difference between the experimental and ideal values is most likely due to the finite thickness of the 2D layer and to LL mixing (particularly important for the 2DHS due to the large hole mass), both of which are known to reduce the strength of the FQH states [9]. We emphasize that in this picture, increased disorder is primarily manifested in a smaller τ_q^* and therefore larger Γ. Hence, *κ is a quantitative measure of LL mixing and finite layer thickness* *effects;* the further κ deviates from $\simeq 0.3$, the more dominant these effects are. As additional evidence for this interpretation, we have replotted the measured gaps of Ref. [4] vs $(e^2/\epsilon \ell_B)/(2p+1)$ and find the same scaling (1) with $\kappa \simeq 0.15$ and $\Gamma \simeq 1.5$ K. The higher value of κ compared to holes is consistent with smaller LL mixing for electrons ($m_e^* = 0.067 m_e$), while the similar value for Γ indicates comparable disorder effects. We have also confirmed that a higher density 2DHS obeys the same gap scaling behavior with a slightly higher value of κ, as expected. Finally we note that the CF mass is closer to the band mass in this 2DHS compared to 2DESs of similar density, an additional effect of LL mixing [2].

Assuming the presence of a Fermi surface, we have analyzed the R_{xx} oscillations around $\nu = \frac{1}{2}$ by treating them as the SdH oscillations originating from the B-field-fractured density of states of the CF metal [6,7]. The analysis is based on an accepted theoretical model which has been successfully employed for 2D carriers around zero B. We utilize the first Fourier component of the 2D Adams-Holstein-type formula for the oscillatory amplitude ΔR_{xx} normalized to the zero-field resistance R_0,

$$\frac{\Delta R_{xx}}{R_0} = 4 \exp\left(-\frac{\pi}{\omega_c \tau_q}\right) \frac{\xi}{\sinh \xi} \cos\left[\pi(2\nu - 1)\right], \quad (3)$$

where $\xi \equiv 2\pi^2 k_B T/\hbar\omega_c$ appears in the Dingle factor $D(\xi) \equiv \xi/\sinh \xi$, describing temperature-induced damping [13]. This formula, including the constant prefactor, describes extremely well the observed oscillations around zero field in 2D systems with varying kinds of scattering mechanisms [14]. We apply it to the metallic state of hole-flux composite particles around $\nu = \frac{1}{2}$ by evaluating R_0 at $B^* = 0$ and making the transformations $\omega_c \rightarrow \omega_c^* \equiv eB^*/m^*$, $\tau_q \rightarrow \tau_q^*$, and $\nu \rightarrow \nu^* \equiv |p|$.

From the T dependence of the R_{xx} extrema, where the cosine term of (3) is ± 1, we can extract m^* of the CFs from the Dingle factor [15]. This was done for each extremum by plotting $\log_{10}(\Delta R_{xx}/R_0 T)$ vs T as shown in Fig. 3, estimating the mass from the slope ($\propto -m^*/B^*$) at high T, and then refining the one-parameter fit by a nonlinear least-squares fit to the entire data set. Error bars are an upper estimate of the systematic uncertainty based on the χ^2-derived standard deviation; the absolute error in the mass is $\simeq \pm 15\%$, mainly from the uncertainty in the thermometry. As manifest in Fig. 3, we find outstanding agreement to the hypothesis of metallic behavior over more than 2 orders of magnitude in the fitted amplitude variable within our entire temperature range, giving further credence to FL formation around $\nu = \frac{1}{2}$. The renormalized mass, plotted vs $p^{-1} \propto B^*$ in Fig. 4(a), exhibits extraordinary features. We note that m^* very roughly follows the $\sqrt{B}$ dependence (dashed line) expected from the energy gap scaling of Fig. 2, but demonstrates strong symmetrical features on both sides of $\nu = \frac{1}{2}$, as well as an apparent divergence as $B^* \rightarrow 0$. This is incongruous with the experimental electron-flux CF behavior propounded by Leadley *et al.* [6]. The overall behavior also differs significantly

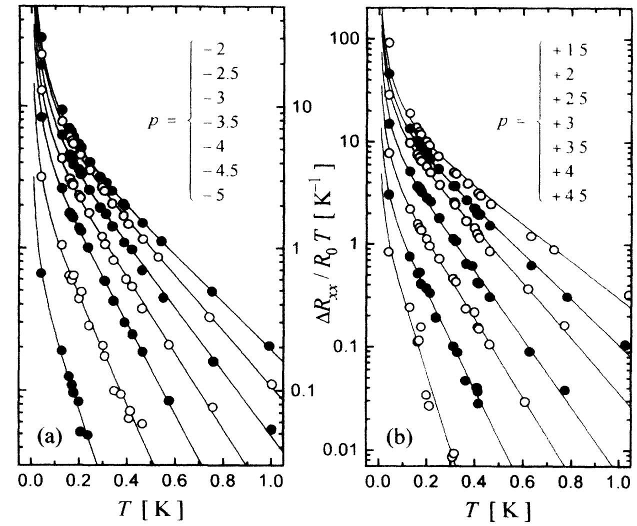

FIG. 3. SdH analysis for (a) $p < 0$ and (b) $p > 0$ ($p = -1.5$ and $+1$ omitted for clarity). Solid lines are fits by Eq. (3) for minima (closed circles) and maxima (open circles).

from the observations of Du *et al.* [7], who observe a constant SdH mass with a small enhancement near $\nu = \frac{1}{2}$ that is stated to be within their experimental error bars. In contrast, we observe significantly more pronounced m^* variations both near $\nu = \frac{1}{2}$ as well as at large B^* [16].

It is also enlightening to examine the quantum lifetime of the hole-flux composites. Typical Dingle plots of $\log_{10}[\Delta R_{xx}/R_0 D(\xi)]$ vs $1/B^*$ or m^*/B^*, however, turn out to be nonlinear, betraying a B dependence of the CF lifetimes. With confidence that Eq. (3) provides an accurate description of the CF SdH oscillations, we use it to extract τ_q^* *as a function of* B^* from the m^* fits. We plot this renormalized lifetime in Fig. 4(b) and note its striking behavior: in the low-B^* limit, scattering appears to be strongly enhanced, while away from $\nu = \frac{1}{2}$, τ_q^* recovers and then drops again. The enhanced values for τ_q^* measured at large positive B^* may be an artifact due to the proximity of the insulating phase near $\nu = \frac{2}{7}$, but the overall trends of the lifetime appear to be similar on both sides of $B^* = 0$.

We have measured values of m^* and τ_q^* for CFs using two temperature-variation techniques that differ in fundamental concept. Nevertheless, the results are surprisingly self-consistent. The gap scaling predicts that m^* scales $\sim\sqrt{B}$ at least far away from $\nu = \frac{1}{2}$ where we can measure the gaps, and the SdH results confirm an overall increasing m^* with comparable magnitude and functional dependence on B. The τ_q^* derived from the quasi-LL broadening Γ and the SdH analysis are in similar quantitative agreement at least close to $B^* = 0$. Moreover, the measured values confirm that in the regime where we have applied the SdH analysis, $k_B T < \hbar\omega_c^*$ and $\omega_c^* \tau_q^* \gtrsim 1$. Comparing to the gap scaling again, we find that the SdH oscillations of CFs around $\nu = \frac{1}{2}$ vanish when $\hbar\omega_c^* \lesssim \Gamma$.

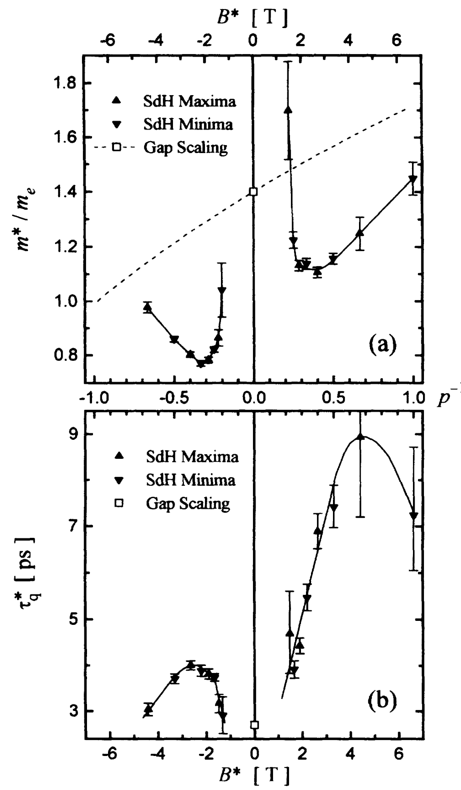

FIG. 4. The renormalized mass (a) and quantum lifetime (b) of the CFs. Open squares and the dashed line are expected values from the scaling in Fig. 2. Lines are guides to the eye.

Our observations have several implications relating to the fundamental nature of the CF metallic state. The energy gap scaling alone indicates that m^* is *not* constant; hence, conventional Fermi liquid (FL) behavior is ruled out. The SdH-derived m^* sheds more light on the proper characterization. Recently, in the context of normal-state high-T_c superconductors, there have been proposals of certain *marginal* 2D FLs [17,18], characterized by a diverging m^* close to the Fermi surface. Metallic properties nevertheless persist, and theory [19] predicts magneto-oscillatory behavior in the same form as Eq. (3). In this interpretation, our measured m^* is strongly suggestive of marginal FL behavior, and the structure observed in the renormalized particle lifetime is qualitatively consistent with the enhanced scattering expected near the Fermi surface of a marginal FL.

The gauge field approach [2], in fact, entails a singularity in the self-energy, leading to a diverging m^* close to the Fermi surface, but this behavior is reduced from an algebraic to a weaker logarithmic correction when Coulomb interactions are assumed. Within our experimental error, however, the observed divergence is clearly not logarithmic and appears to be faster than any power law in $(B^*)^{-1} \propto p$. Our best estimate, based on fitting the data of Fig. 4(a) for $p \leq -3$ and $p \geq 2.5$, is an exponential divergence $m^* \sim \exp(|p|^\gamma)$, where $\gamma \simeq 1.4 \pm 0.3$. The nature of this divergence suggests a novel universality class of marginal FL behavior. In this sense, the data have a striking parallel with the 2D Luttinger liquid class, involving strongly diverging mass enhancement factors [18,20]. The bulk of the observed renormalization structure of this exotic metal, however, awaits a comprehensive theoretical explanation.

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Note added.—After the submission of this manuscript we received a paper by R. Du *et al.* [21] reporting a CF mass enhancement near $\nu = \frac{1}{2}$ in a 2D *electron* system.

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[15] The data reduction method we employ involves drawing curve fits through the minima and maxima to obtain the oscillation envelope, averaging these curves to obtain the oscillation background, and measuring the resistivity difference between the various extrema and this background curve; this difference is the amplitude ΔR_{xx}.
[16] Density fluctuations are primarily responsible for the vanishing of oscillations around $\nu = \frac{1}{2}$. Numerical simulations which we have performed reveal that such sample inhomogeneities produce an artificially *suppressed* extracted mass close to $\nu = \frac{1}{2}$; the observed mass enhancement, therefore, cannot be explained by this type of disorder.
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