Temperature Dependent Scattering of Composite Fermions

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We have studied the temperature dependent resistivity at exactly half Landau level filling of a high-quality two-dimensional electron system in high magnetic field. The experimental data can be well described in terms of a standard degenerate system of fermions at *zero* magnetic field being scattered by impurities and acoustic phonons. The gauge-field-mediated phonon scattering contributes a T^{-3} dependence to the composite fermion mobility. The effective mass of composite fermions is obtained from the temperature dependence of composite fermion–impurity scattering.

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The composite fermion model has developed into a very attractive framework in which to interpret electron correlation phenomena in two-dimensional (2D) electron systems in high magnetic fields [1–3]. It postulates the existence of new particles called composite fermions (CF's), which consist of electrons to which an even number of magnetic flux quanta has been attached by virtue of the electron-electron interaction. Such flux attachment to the carriers turns out to be a very effective theoretical means to maintain optimum electron separation and, hence, achieve minimum Coulomb energy in a many particle system.

The CF model allows one to draw analogies between the behavior of electrons and CF's. The process of flux quanta attachment incorporates part of the external magnetic field into the new particle and the resulting CF experiences a reduced, *effective* magnetic field. In this way, the higher order states of the fractional quantum Hall effect (FQHE) at filling factor $\nu = p/(2p \pm 1)$ around $\nu = \frac{1}{2}$ are being related to the integral quantum Hall effect (IQHE) at $\nu = p$ [2]. Just as the IQHE is a consequence of the quantization of regular electrons into Landau levels, the FQHE can be regarded as resulting from the Landau quantization of these bizarre, new objects. An interpretation of the magneto-oscillatory data on the higher order FQHE states in terms of a simple, single-particle Shubnikov–de Haas (SdH) formalism has been very successful [4,5]. It supports the existence of new particles with a definite effective mass moving under the influence of an *effective* magnetic field.

One of the most startling implications of the CF model is the existence of a Fermi liquid [1] with a well-defined Fermi wave vector $k_F = \sqrt{4\pi n}$ at $\nu = \frac{1}{2}$, where n is the electron density. At this filling factor, the effective magnetic field for CF's is exactly zero and an analogy is drawn to a regular electron system in the *absence* of a magnetic field. In fact, a sequence of elegant experiments on the state at $\nu = \frac{1}{2}$ [6–8] confirms the existence of such a k_F of the CF liquid at $\nu = \frac{1}{2}$. So far, this is the only available experimental parameter of the proposed Fermi liquid state at exactly $\nu = \frac{1}{2}$. Quite clearly, the measurement of other parameters of this Fermi liquid would strengthen the case for the applicability of the CF model.

Even when accepting the CF picture, it has become quite clear in the recent past that the state at $\nu = \frac{1}{2}$ is not likely an ordinary Fermi liquid [1,9]. This theoretical expectation is supported by SdH measurements which detect the beginnings of such a divergence away from $\nu = \frac{1}{2}$, although its strength seems to exceed by far the theoretical estimates [10,11]. All these experimental findings and theoretical models beg the question: What is the mass of a CF in the Fermi liquid state at $\nu = \frac{1}{2}$?

Our experimental data on the T dependence of the resistivity at $\nu = \frac{1}{2}$ can be fit with a remarkable accuracy and over the whole temperature range from 60 mK to 10 K to the T dependent resistivity of a degenerate system of CF's being scattered by a combination of impurities and acoustic phonons. From the low T behavior we can deduce a density of states mass for CF's at exactly $\nu = \frac{1}{2}$.

The experiments were performed on a 2D electron gas at a GaAs/AlGaAs heterojunction interface, having an electron density $n = 1.5 \times 10^{11}$ cm^{-2} and a peak mobility $\mu = 8 \times 10^6$ cm^2/V sec. These conditions were achieved after standard, short-term illumination of the sample by light from an LED. The setback distance between the silicon dopant and electron gas was 800 Å. The sample consisted of a 5×5 mm square with eight indium contacts diffused symmetrically around the edges. All measurements were performed in a 14 T superconducting magnet system with the sample on a heating stage in thermal contact with a dilution refrigerator. We employed a low-frequency (17 Hz) lock-in technique. The resistivity ρ of the specimen at $B = 0$ and at $\nu = \frac{1}{2}$ was determined by the standard van der Pauw method, performing four separate measurements in four different I-V configurations at each temperature on a subset of four of the eight indium contacts. The inset to Fig. 1 shows a standard magnetoresistance (R_{xx}) trace at 50 mK. The bulk of Fig. 1 shows the result of the T dependent van der Pauw measurements at $B = 0$ and at $\nu = \frac{1}{2}$, translated into mobility $\mu = (ne\rho)^{-1}$. Traces $\mu^e(T)$ and $\mu^{\mathrm{CF}}(T)$ identify the data

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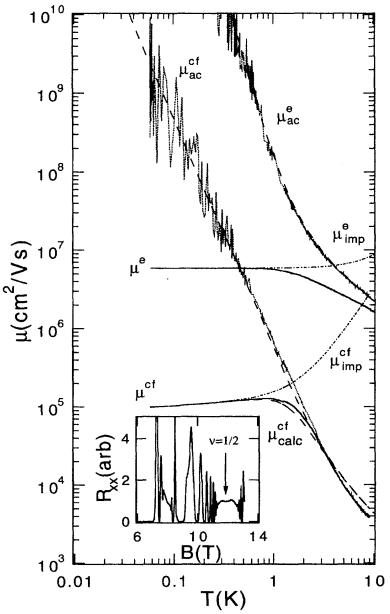

FIG. 1. Temperature dependence of the electron mobility μ^e and CF mobility μ^{CF} (solid curves). The traces μ^e_{imp} and $\mu^{\mathrm{CF}}_{\mathrm{imp}}$ (dash-dotted curves) represent the impurity scattering part for electrons and CF's, respectively. The dotted traces μ^e_{ac} and $\mu^{\mathrm{CF}}_{\mathrm{ac}}$ are the residues, $\mu_{\mathrm{ac}} = (\mu^{-1} - \mu^{-1}_{\mathrm{imp}})^{-1}$, representing the phonon contribution to μ^e and μ^{CF}. Since μ_{ac} is derived from the inverse of the difference of two very similar large numbers, the noise level increases significantly as $T \to 0$. The dashed curves underneath the noisy traces are the theoretical curves for phonons. In the case of $\mu^{\mathrm{CF}}_{\mathrm{ac}}$ the prefactor of the T^{-3} dependence is a fitting parameter (vertical shift). The dashed curve $\mu^{\mathrm{CF}}_{\mathrm{ac}}$ under μ^{CF} represents the total theoretical mobility based on the numerical calculation discussed in the text. The inset shows the magnetoresistance R_{xx} of the sample at base temperature of 50 mK.

for electrons and CF's, respectively. The other traces of Fig. 1 are a result of our data reduction.

In trying to understand the behavior of $\mu^{\mathrm{CF}}(T)$, we draw on the analogy between electrons and CF's at $\nu = \frac{1}{2}$. The electron mobility trace is well understood from earlier studies [12–15]. The mobility of 2D electrons in this T range is described in terms of scattering by acoustic phonons and by fixed crystal imperfections (impurities, defects, interface roughness, etc.). Scattering of a *degenerate* Fermi system at imperfections contributes a term $\mu_{\mathrm{imp}}(T) \propto c_1 + c_2T^2$. The ratio $c_1/c_2 \propto E_f^2 \equiv (\hbar^2k_f^2/2m^*)^2$ provides a measure for the density of state mass m^* at E_f. For an electron system with density of 1.5×10^{11} cm^{-2} and $m^* \sim 0.07m_0$, the Fermi energy $E_f \sim 5$ meV is large and, therefore, the temperature dependence of μ_{imp}, shown as a dash-dotted curve, is negligible over the temperature range in Fig. 1. Acoustic phonons contribute a T^{-1} term at high temperatures which changes to a T^{-5} dependence at low temperatures. The characteristic temperature for the transition is related to the sound velocity s by $kT_c = 2k_f\hbar s$ and amounts only to a few degrees in our system. The $\mu^e(T)$ data in Fig. 1 are exceedingly well reproduced by a combination of these two scattering mechanisms. The noisy part of μ^e_{ac} shows the residue $\mu^e_{\mathrm{ac}}(T) = [\mu^e(T)^{-1} - \mu^e_{\mathrm{imp}}(T)^{-1}]^{-1}$ which follows very well the underlying, smooth theoretical curve for acoustical phonon scattering.

Having succeeded in reproducing theoretically the temperature dependence of the electron scattering in Fig. 1, we now approach the temperature dependence of the CF scattering in a manner analogous to the electron case. We parametrized the CF's at half filling as a *degenerate* system of carriers moving in *zero* magnetic field with a Fermi energy E_f and a mass m^* [16]. Then the CF-impurity mobility is given by $\mu^{\mathrm{CF}}_{\mathrm{imp}} = c_1 + c_2T^2 + \mu^{\mathrm{CF}}_{\mathrm{int}}$, with an added term $\mu^{\mathrm{CF}}_{\mathrm{int}} = c_3 \log T$. This standard term, which results from interaction effects, is negligible for high-mobility electrons but becomes appreciable in the case of CF's [1]. Trace $\mu^{\mathrm{CF}}_{\mathrm{imp}} = c_1 + c_2T^2 + c_3 \log T$ of Fig. 1 shows the low temperature mobility which best fits the experimental data. The parameters are $c_1 = 1.25 \times 10^5$ cm^2/V sec, $c_2 = 2.50 \times 10^4$ cm^2/V sec K^2, $c_3 = (2.1 \times 10^4$ cm^2/V sec$)/\log K$.

Concentrating on the temperature range between 60 mK and 4 K, the ratio $c_1/c_2 = (8/\pi^2)(E_f)^2 = 5.0K^2$ leads us to a mass $m^* = (3.0 \pm 0.3)m_0$. While the quoted mass (and its uncertainty) is the only value of this parameter which fits the data of Fig. 1, the data themselves are illumination dependent and so is the derived mass. Masses increase from $m^* \sim 2m_0$ to $m^* \sim 6m_0$ with increasing level of illumination. This illumination *dependence* is not to be confused with an *uncertainty* in the mass determination which remains at the 10% level. It rather seems to indicate that different illumination levels create different conditions leading to different masses as discussed later in the text. The derived CF masses of $2m_0 \lesssim m^* \lesssim 6m_0$ at exactly $\nu = \frac{1}{2}$ are about a factor of 2 to 6 larger than the masses measured in the SdH effect, slightly away from $\nu = \frac{1}{2}$ [4,5,10,11,17]. It is tempting to speculate that our enhanced values indicate a continuation of the mass divergence seen in these experiments, smeared out by electron density fluctuations or lifetime effect. This may render the value at exactly $\nu = \frac{1}{2}$ large but finite. We clearly need to develop a better understanding of the CF liquid before pursuing such speculations. As to the CF-phonon interaction, the noisy trace marked $\mu^{\mathrm{CF}}_{\mathrm{ac}}$ represents the residue to the data determined as $\mu^{\mathrm{CF}}_{\mathrm{ac}}(T) = [\mu^{\mathrm{CF}}(T)^{-1} - \mu^{\mathrm{CF}}_{\mathrm{imp}}(T)^{-1}]^{-1}$. Over more than 4 orders of magnitude $\mu^{\mathrm{CF}}_{\mathrm{ac}}$ follows a T^{-3} dependence, which differs markedly from the T^{-5} dependence of the electrons.

At this point it is important to describe in some more detail our fitting procedure and its limitations. We have taken several precautions. (1) The $\mu(T)$ dependence is wholly based on a physical picture of the scattering processes involved. $\mu^{\mathrm{CF}}_{\mathrm{imp}}(T)$ is chosen in analogy to the

ordinary Fermi liquid case adding a $\log T$ term, theoretically predicted to be appreciable [1]. $\mu_{\rm ac}^{\rm CF}(T \to 0)$, again in analogy to a Fermi liquid, is expected to be a power law $c_4 T^{\alpha}$ with an integral exponent α. (2) We have performed numerous fits to $\mu^{\rm CF}$, on many experimental data sets (see below) constraining c_1, c_2, c_3, c_4, or α. We purposely imposed certain values on some of the constants and determined how it affected the others and whether there is more than one set of parameters that fit the experimental data. In all cases we find $\alpha = -3$ to be the only acceptable integer exponent and the constants c_1, c_2, c_3, c_4 to be tightly bounded. A full five-parameter optimization results in $\alpha = -2.83$. (3) The exponent $\alpha = -3$ is supported by theoretical calculations (see below). (4) The values of the resulting fitting parameters are physically sensible. We, therefore, feel very confident that the fit to $\mu^{\rm CF}(T)$ is valid and that we can use the so-determined parameters to deduce properties of the CF's.

The T^{-3} dependence of $\mu_{\rm ac}^{\rm CF}$ is in contrast to the T^{-5} dependence for $\mu_{\rm ac}^{e}$. Such a modified power law for CF acoustic phonon scattering has been theoretically derived by He [18]. As in the case of scattering with imperfection, the origin of this weaker temperature dependence lies in the enhanced CF-phonon scattering caused by gauge field fluctuations of the system.

The enhancement can be understood by the following simple argument. A phonon generates a distortion in the GaAs lattice which induced a local charge imbalance. Electrons will move to screen the charge imbalance giving rise to density fluctuations δn_q at wave vector q. At zero magnetic field electron scattering proceeds via the interaction $H_I^e \sim \delta n_q^\dagger \delta n_q$ which for $T \ll E_f$ and for the piezoelectric channel leads to the characteristic T^{-5} dependence [12]. In the case of CF's the density fluctuations δn_q give rise to fluctuations in the Chern-Simons gauge field [1,19] given by $\delta a_q \sim \delta n_q/q$ which scatters the *current* through the interaction $H_I^{\rm CF} \sim j_q^\dagger \delta a_q \sim j_q^\dagger \delta n_q/q \sim \delta n_q^\dagger \delta n_q/q$. Hence, the scattering rate which is proportional to the matrix element squared is enhanced by a factor $1/q^2 \propto 1/\omega^2$ (for phonons) $\propto 1/T^2$ for CF's as compared to regular electrons at $B = 0$.

We have chosen to use the memory function technique [20] to calculate the electron-phonon scattering rates [18]. When $T/E_f \ll 1$, the rates for the overriding scattering with piezoelectric phonons of longitudinal (ℓ) and transverse (t) polarization are given by [18]

$$\frac{1}{\tau_{p\ell}} = \frac{135\zeta(3)}{2^9\pi}\Re_\ell, \quad \frac{2}{\tau_{pt}} = \frac{135\zeta(3)}{2^9\pi}\left(3 + \frac{35}{8}\right)\Re_t, \tag{1}$$

$$\Re_{\ell,t} = \tilde{\phi}^2\left(\frac{m^*}{m}\right)^2 \frac{1}{\hbar k_F}\frac{(\hbar k_F)^2}{m_b}\frac{m_b(eh_{14})^2}{\kappa_{\ell,t}^2\hbar^2} \times \frac{[(\hbar k_F)^2/m_b](k_B T)^3}{(\hbar s_{\ell,t}k_F)^2\,(e^2k_F/\varepsilon)^2}, \tag{2}$$

where $\tilde{\phi} = 2$ for $\nu = \frac{1}{2}$, m^* is the density of states mass at E_f, m_b is the GaAs band mass, $\varepsilon = 12.8$ its dielectric constant, $s_{\ell,t}$ are sound velocities of longitudinal and transverse piezoelectric phonons, h_{14} is the relevant piezoelectric tensor component, and $\kappa_{\ell,t} = \eta s_{\ell,t}$ with the mass density η of GaAs. There remains a theoretical ambiguity as to whether the effective mass of CF's or the electronic band mass should be used in the CF-phonon scattering rate. Therefore $\overline{m}$ is either m^* or m_b.

It is very satisfying to observe that the experimental T^{-3} dependence is, indeed, reproduced by theory. It provides further validation for our interpretation of the resistivity data. Using *ad hoc* $\overline{m} = m_b$ and $1/\tau_p = 1/\tau_{p\ell} + 2/\tau_{pt}$, we obtain an effective mass $m^* = 12.7 m_b = 0.89 m_0$ from the experiment. This mass compares favorably with the reported mass of CF's [1,4,5,10,11,17] and not too dissimilar from the mass we obtain from CF-impurity scattering. Presently we cannot rely on such a procedure to deduce the CF mass from the prefactor of the T^{-3} term.

At higher temperatures the $\mu^{\rm CF}$ data of Fig. 1 show a deviation from the T^{-3} behavior. We have not made a detailed attempt to fit this section of $\mu^{\rm CF}$. Several factors may contribute to this deviation: (1) The theoretical model does not include deformation potential scattering [12] which may contribute to the scattering at these elevated temperatures. (2) At these high temperatures magnetic flux quanta may become unbound and the Fermi liquid picture may no longer be appropriate. (3) The Fermi liquid is no longer degenerate and the softening of the Fermi edge needs to be taken into account. The latter of these effects can be incorporated integrating numerically and in analogy to the electron case, across the Fermi distribution. The dashed curve, continuing into the T^{-3} dependence at low T, is the result of this calculation. It shows that a large part of this "bending" can be explained in these terms.

Summarizing our findings, we may state that the low temperature transport behavior at $\nu = \frac{1}{2}$ Landau-level filling is extraordinarily well reproduced by the temperature dependent impurity and phonon scattering of a Fermi liquid of particles with a mass m^*. From the ratio of the impurity scattering parameters of the data of Fig. 1 a mass $m^* = 3.0 m_0$ is deduced which is somewhat larger but not too different from the masses derived by SdH measurements away from $\nu = \frac{1}{2}$. Phonon scattering shows the theoretically deduced T^{-3} behavior, although there remains theoretical uncertainty in its prefactor. At higher temperature, the dependence seems to be quite well described by a softening of the Fermi edge. We regard our measurement as a determination of the CF density of states mass at exactly $\nu = \frac{1}{2}$.

The success of such an extraordinarily simple interpretation of the transport behavior is quite remarkable. Yet there is one important feature of our data for which we

have no firm explanation. It may point to an interesting characteristic of the CF mass: The $\mu(T)$ dependence for CF's in Fig. 1 depends on the illumination level of the specimen and subsequent data reduction yields a *set* of *different* masses. Under increasing illumination the derived masses vary from $m^* = 2.1m_0$ to $m^* = 6.2m_0$. The requisite data are shown in Fig. 2. To be sure, in all cases, the above model produces fits comparable in quality to the one shown in Fig. 1. The T^{-3} behavior is maintained in all cases. Furthermore, each data set is internally consistent: As the mass deduced from the impurity scattering rises, so does the mass deduced from the prefactor of the phonon scattering and their ratio, at least in the $\overline{m} = m_b$ case, varies by less than a factor of 1.5. This is not due to a variation in electron density: At the illumination range covered in Fig. 2, the electron density is constant. Hence, $\nu = \frac{1}{2}$ always remains at fixed $B_{1/2}$.

We can provide two enticing interpretations of this observation although we do not have the means to prove either. The electron density saturation with increasing light level is well understood. At saturation, although more donors in the δ layer are ionized, their chemical potential equals that of the 2D electron system and electron transfer from the donors into the 2D channel ceases. However, the carriers remaining at the site of the δ layer, 800 Å away from the 2D electron system, form a highly polarizable reservoir. This polarizable sheet increases the screening of the carriers in the 2D system, thereby moving the *e-e* interaction from Coulombic towards short range. Theory suggests [1,9] that the mass of CF's depends on the length scale of the *e-e* interaction. Short-range interaction is expected to create heavier CF's than long-range interaction. This is in accord with the trend in our measurements. Alternatively, increasing electron concentration in the δ layer, which is known to improve mobility, may ameliorate the density fluctuations

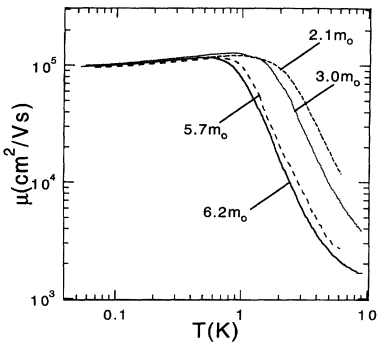

FIG. 2. Temperature dependent mobility of CF's for different levels of illumination. The numbers indicate the derived masses increasing with increasing light level.

at $\nu = \frac{1}{2}$. This may reduce the smearing of the mass divergence observed in SdH experiments and lead to a larger effective mass value at exactly $\nu = \frac{1}{2}$. Whether as much as a factor of 3 increase in the CF mass can be attributed to either of the mechanisms needs to await a better understanding of the CF characteristics.

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