Phonon Excitations of Composite Fermion Landau Levels


1 Institut für Festkörperphysik, Universität Hannover, Appelstraße 2, 30167 Hannover, Germany
2 High Field Magnet Laboratory, NSRIM, University of Nijmegen, Toernooiveld 1, 6525 ED Nijmegen, The Netherlands
3 Lehrstuhl für Angewandte Festkörperphysik, Ruhr-Universität Bochum, Universitätsstraße 150, 44780 Bochum, Germany

(Dated: February 26, 2004)

Phonon excitations of fractional quantum Hall states at filling factors $\nu = \frac{1}{3}, \frac{2}{5}, \frac{3}{7}, \frac{4}{9}, \frac{5}{11}, \frac{6}{13}$ are experimentally shown to be based on Landau level transitions of Composite Fermions. At filling factor $\nu = \frac{3}{5}$, however, a linear field dependence of the excitation energy in the high-field regime rather hints towards a spin transition excited by the phonons. We propose to explain this surprising observation by an only partially polarized $\frac{\sqrt{3}}{2}$-ground-state making the energetically lower lying spin transition also allowed for phonon excitations.

PACS numbers: 73.43.-f, 73.20.Mf, 72.10.Di

The Coulomb interaction in a two-dimensional electron system (2DES) subjected to a quantizing magnetic field leads to the formation of new, fractionally charged quasiparticles at Landau level filling factors $\nu = \frac{p}{q}$ ($q$ is an odd-integer) [1, 2]. In the last decade this fractional quantum Hall (FQH) effect was very effectively described in the framework of Composite Fermions (CFs) [3]. At a fractional filling factor with even denominator, $\nu = \frac{1}{2m}$, these quasiparticles are formed by attaching an even number $2m$ of flux quanta $\phi_0$ to each electron, i.e. two flux quanta at $\nu = 1/2$. Their effective mass $m^*$ is originating from the Coulomb interaction [4].

The ground state of a 2DES at a fractional filling factor $\nu = \frac{p}{q}$ is a collective wave function [2] with finite wave-vector collective excitations [5, 6] directly accessible by, e.g. Raman techniques [7, 8], photoluminescence [9] or phonon absorption experiments [10, 11]. In a simple picture these excitations originate from the level structure of CFs in an effective magnetic field, $B_{\text{eff}} = B - B_{\nu}$ ($\nu = \frac{1}{2}$), with an effective (integer) filling factor, $p = \frac{\nu}{2\nu - \nu}$ [3]. The levels can then be described as spin-split Landau levels of CFs and, therefore, excitations can be interpreted as either Landau level transitions or spin-excitations.

In this Letter we use phonons to probe the FQH excitation spectrum at filling factors $\nu = \frac{1}{3}, \frac{2}{5}, \frac{4}{7}, \frac{3}{5}, \frac{2}{3}, \frac{3}{5}, \frac{4}{9}, \frac{5}{11}, \frac{6}{13}$, and $\frac{5}{11}$ are 2DES are placed at the edges of the sample far away from the meander to avoid any phonon interactions with the contacts. We took great care that the contact resistances (< 10 $\Omega$) only play a negligible role in the measured two-terminal resistance. Using the persistent photoconductivity we varied the electron concentration in several steps from $n = 0.89 \times 10^{15}$ m$^{-2}$ (mobility $\mu = 193$ m$^2$/Vs) to $n = 1.50 \times 10^{15}$ m$^{-2}$ (mobility $\mu = 102$ m$^2$/Vs).

Our sample consists of a high-mobility AlGaAs/GaAs heterostructure grown on a 2 mm thick GaAs wafer. On the front side containing the 2DES we patterned a 60 $\mu$m wide meander extending over a total length $l = 10$ mm on an area $A = 1 \times 1$ mm$^2$. The huge aspect ratio $l/w = 111$ maximizes the $\rho_{xx}$-contribution to the two-terminal resistance and thereby allows to measure smallest changes in $\rho_{xx}$. The Ohmic contacts to the 2DES are placed at the edges of the sample far away from the meander to avoid any phonon interactions with the contacts. We took great care that the contact resistances ($< 10 \Omega$) only play a negligible role in the measured two-terminal resistance. Using the persistent photoconductivity we varied the electron concentration in several steps from $n = 0.89 \times 10^{15}$ m$^{-2}$ (mobility $\mu = 102$ m$^2$/Vs) to $n = 1.50 \times 10^{15}$ m$^{-2}$ (mobility $\mu = 193$ m$^2$/Vs).

The sample is mounted on the tail of a dilution refrigerator in a superconducting magnet with maximum fields up to 13 T and connected to high frequency coaxial cables. Great care was taken to assure a proper thermal anchoring of the cables. We achieve a 2DES temperature $T_{\text{DES}} \lesssim 100$ mK for a cryostat base temperature of 75 mK.

The experimental setup is shown schematically in Fig. 1a: A thin constantan film acting as a phonon emitter is placed on the polished back side of the sample. By passing a short current pulse during a time $t_H$ through this heater, non-equilibrium phonon pulses are created at the heater-GaAs interface. They are characterized by a black-body spectrum at a temperature $T_H = (P_H/\sigma A_H + T_0^4)^{1/4}$, where $A_H$ is the heater area, $P_H$ is the power dissipated in the heater, $\sigma = 524$ W/m$^2$ K$^4$ is the acoustic mismatch constant between constantan and GaAs, and $T_0$ is the GaAs lattice temperature [12]. When entering the GaAs, the non-equilibrium phonons travel ballistically through the 2 mm-thick substrate. Af-
we have plotted a characteristic signal at the measured relative specific heats, \( C(T) \), is measured as a function of time and the experiment is repeated a few million times. When the total power dissipated inside the GaAs substrate, the sample has cooled down back to its base temperature and the experiment is repeated a few million times. Depending on the total power dissipated inside the GaAs substrate, the sample has cooled down back to its base temperature and the experiment is repeated a few million times.

In Fig. 1 we have plotted a characteristic signal at filling factor \( \nu = 2/3 \). Due to their strong focusing [11] mainly TA phonons are visible, with a first peak starting around \( \tau_{TA} = 0.6 \) \( \mu s \) and a second due to multiply reflected phonons after \( 3\tau_{TA} \). After typically 1 ms (depending on the total power dissipated inside the GaAs substrate), the sample has cooled down back to its base temperature and the experiment is repeated a few million times.

By using the temperature dependent resistance measured under equilibrium conditions (as plotted in Fig. 1b) the raw phonon signal curve can be translated into a 2DES temperature versus time. The reliability of this procedure is checked as follows: After a certain time (\( \approx 10 \) \( \mu s \)), all non-equilibrium phonons induced with the heater are thermalized in the GaAs substrate and the 2DES and the substrate are in thermal equilibrium. This is experimentally measured by a merely changing 2DES temperature. This measured temperature agrees well with the theoretically expected one as deduced from the total energy dissipated in the heater, \( P_H \tau_H \), and the specific heat of the GaAs substrate.

In order to extract quantitative data from our experiments we use a simple model to describe the phonon absorption in the 2DES. In the most general case, the differential temperature gain \( dT \) of the 2DES within a time interval \( dt \) is given by

\[
C(T) \, dT = r(T, T_H) P_H \, dt - P_e(T, T_0) \, dt
\]

where \( C(T) \) is the 2DES’s specific heat, \( r(T, T_H) P_H \) is the phonon energy absorbed by the 2DES with an absorption coefficient \( r \) depending on the heater temperature \( T_H \) and the 2DES temperature \( T \), and \( P_e(T, T_0) \) is the energy emitted by the 2DES, depending on \( T \) and the equilibrium substrate temperature \( T_0 \). In our experiments we use very short (10 ns) heater pulses with a moderate heater power \( P_H \). Consequently, the peak height of the first ballistic phonon signal peak is dominated by absorption and the emission term can be ignored on these short time scales.

In a first set of experiments, we calibrate the relative specific heat of the 2DES at given fractional filling factor. The maximum 2DES temperature on the first ballistic phonon peak, \( T_1 \) (see Fig. 1d), is measured as a function of the 2DES base temperature \( T_0 \), with a fixed duration and a constant amplitude of the heater pulse. Since all the \( T_0 \) used are distinctively lower than the energy gaps at these filling factors we always deal with a situation where the quasiparticle ground states are almost full and their excitations are almost empty. As a consequence, the relative proportion of phonons absorbed is independent from \( T_0 \), and we can approximate \( r(T, T_H) \to r_0 \). Integrating Eq. (1) over the pulse length with these assumptions, we get:

\[
\int_{T_0}^{T_1} C(T) \, dT = r_0 P_H \tau_H
\]

Using the mean value theorem for this integral equation we can determine the relative specific heat \( C(T)/r_0 \) of the 2DES from our set of experiments where we measured \( T_1 \) for fixed heater power and varying \( T_0 \). As a consistency check we performed the same set of experiments with different heater powers and find a very comparable temperature dependence of \( C(T) \).

In Fig. 2a the measured relative specific heats, \( C(T)/r_0 \), are shown exemplarily for filling factors \( \nu = 2/3 \) and \( 2/5 \). We should note that we cannot determine the absolute value of \( C(T) \) due to the unknown absorption
coefficient \(r_0\) (0 < \(r_0\) < 1). The lines shown in Fig. 2a are fits to the theoretical predictions for the specific heat of a 2DES \(C_{2DES} \propto \frac{1}{T^2} e^{-\Delta/T} \) [14] plus a small empirical constant, taking into account an additional contribution to \(C(T)\) possibly resulting from a finite (thermodynamic) density of states inside the excitation gap [11].

In a second set of phonon absorption experiments we can now determine the energy gaps at fractional filling factors. This time, the heater temperature \(T_H\) is varied for a fixed base temperature \(T_0\). By increasing \(T_H\) the number of phonons for every wavelength is increased. Since the major contribution to the phonon absorption is predominantly due to excitations around a gap \(\Delta\), the total energy absorbed by the 2DES increases as

\[
E_{abs} \propto \frac{1}{\exp(\Delta/T_H)} - 1
\]

The absorbed phonon energy, \(E_{abs} = r(T, T_H) P_{hT_H}\), as a function of \(T_H\) is deduced from the measured \(T_1\) and \(T_0\) by integrating Eq. (1) using the previously determined specific heat. Again, the emission term is neglected for the short time scales considered. In Fig. 2b we show results for filling factors \(\nu = 2/3\) and 2/5. The gap values are obtained by averaging over several experiments using different heater powers and base temperatures. All the individual gaps measured are within ±10% of the average value. The solid lines in Fig. 2b show fits using Eq. (3) and indeed, the experimental data match a model with phonon excitation across a single energy gap \(\Delta\).

Using such an elaborated series of calibrations we can now investigate in detail the excitation gaps. As listed in Table I, we have measured the gaps for eight different electron concentrations at filling factors \(\nu = 2/3, 2/5,\) and 1/3 up to the highest magnetic fields accessible in our magnet. These filling factors all showed well developed minima for all the concentrations used. Additional filling factors \(\nu = 4/3, 5/3, 3/5\) and 2/7 are only clearly pronounced for the highest electron concentration and, accordingly, phonon gaps were only measured in sample #8 for these filling fractions.

All gaps measured at various filling factors and electron concentrations are compiled in Fig. 3a. For comparison, we also show the energy gaps deduced from activated transport measurements [17] in Fig. 3b. In order to discuss these results for different filling factors \(\nu\) we describe them within the CF picture [3], where the FQH filling factors \(\nu = \frac{p}{p+1}\) are mapped to integer CF filling factors \(p\). The energy levels for \(|p|\) filled CF Landau levels can then be written [15] as:

\[
E_{ns}(p) = (n + \frac{1}{2}) \hbar \omega_c^*(p) + sg^* \mu_B B
\]
TABLE I: Electron concentrations, mobilities and energy gaps measured by phonon absorption at filling factors 2/3, 2/5, and 1/3. When no values are given, the field required was too large to be accessed in our magnet.

<table>
<thead>
<tr>
<th>sample #</th>
<th>n [10^{15} \text{ m}^{-2}]</th>
<th>\mu [\text{m}^2/\text{Vs}]</th>
<th>\Delta [K] for \nu = \frac{2}{3}</th>
<th>\frac{2}{5}</th>
<th>\frac{1}{3}</th>
</tr>
</thead>
<tbody>
<tr>
<td>1</td>
<td>0.89</td>
<td>102</td>
<td>3.2</td>
<td>4.9</td>
<td>9.2</td>
</tr>
<tr>
<td>2</td>
<td>0.95</td>
<td>109</td>
<td>4.1</td>
<td>5.3</td>
<td>8.2</td>
</tr>
<tr>
<td>3</td>
<td>1.03</td>
<td>119</td>
<td>4.5</td>
<td>5.9</td>
<td>8.9</td>
</tr>
<tr>
<td>4</td>
<td>1.13</td>
<td>131</td>
<td>4.8</td>
<td>5.8</td>
<td>—</td>
</tr>
<tr>
<td>5</td>
<td>1.21</td>
<td>144</td>
<td>4.9</td>
<td>5.7</td>
<td>—</td>
</tr>
<tr>
<td>6</td>
<td>1.36</td>
<td>168</td>
<td>6.0</td>
<td>—</td>
<td>—</td>
</tr>
<tr>
<td>7</td>
<td>1.46</td>
<td>187</td>
<td>6.7</td>
<td>—</td>
<td>—</td>
</tr>
<tr>
<td>8</td>
<td>1.50</td>
<td>193</td>
<td>6.6</td>
<td>—</td>
<td>—</td>
</tr>
</tbody>
</table>

Here the CF cyclotron energy is given purely by the Coulomb interaction and thus follows the form \( h\omega_{\text{c}}(p) = \frac{eB}{m^* (2p \pm 1)} \) with a CF mass \( m^* = m_e \alpha \sqrt{B/T} \) [16].

Since phonons carry no spin we expect that, in a phonon absorption experiment, the lowest lying excitations of a CF state \( p \) are Landau level transitions from \( n \) to \( n+1 \) with the same spin \( s \). The corresponding energy gap \( h\omega_{\text{c}}(p) \) can now be adjusted to our data by one single dimensionless mass parameter \( \alpha \). In Fig. 3 the fits of such Landau level mass transitions to all the data at \( \nu = 1/3, 2/5, 4/7, 3/5, 4/3, \) and \( 5/3 \) are shown yielding \( \alpha \approx 0.158 \) (±0.006) (dotted lines). Here, \( \nu = 4/3 = 1 + 1/3 \) and \( 5/3 = 1 + 2/3 \) are treated as \( 1/3 \) respectively \( 2/3 \) plus one inert fully occupied Landau level. The experimentally determined CF mass parameter \( \alpha \) is in astonishing agreement with the theoretical predictions in Eq. (1) of Ref. 16.

Compared to all phonon gaps measured at the above mentioned filling factors, the phonon absorption data at filling factor 2/3, are distinctly different: The measured excitation gaps can in no way be described with the square-root dependence of CF Landau level excitations. They rather show a linear dependence, \( \Delta_{2/3} \propto (B - B_p) \), strongly suggesting that they are related to a spin gap, which is normally not directly accessible by phonon excitations. Here \( B_p \approx 2.8 \text{ T} \) is the field where the 2/3 state changes from a spin-unpolarized to a spin-polarized state. Indeed, when we quantitatively compare the expected field dependence of the spin gap with our data we find a remarkable agreement (dashed line in Fig. 3a). This interpretation is also supported by the transport gaps measured at \( \nu = 2/3 \) which have the same linear behavior, reduced by a constant due to disorder as seen in Fig. 3b. All other transport gaps are also systematically lower than the phonon excitation gaps due to disorder effects and because temperature can also couple to spin flip excitations.

The fact that we observe spin-related excitation gaps with phonons indicates that the \( \nu = 2/3 \) state can not be described with independent spin and Landau level indexes. This supposition is also supported by recent theoretical [18] and experimental [19] evidence suggesting that \( \nu = 2/3 \) state is not fully polarized, even in high magnetic fields. As a result, we may speculate that the complexity of the 2/3 state is responsible for the appearance of a spin-forbidden transition in the phonon absorption.

In conclusion, we have measured phonon excitation gaps in the FQH regime for filling factors \( \nu = 1/3, 2/5, 4/7, 3/5, 2/3, 4/3, \) and \( 5/3 \) for eight different electron densities. For all filling factors besides \( \nu = 2/3 \) the measured gaps can be well described in the framework of Landau level transitions of CF involving no spin flip. The gaps measured at \( \nu = 2/3 \), however, correspond to a normally forbidden spin transition, pointing towards a complex not fully polarized ground state.

We acknowledge financial support by BMBF and DFG priority program “quantum Hall systems”.