Phonon Excitations of Composite-Fermion Landau Levels

F. Schulze-Wischeler,[1] F. Hohls,[1] U. Zeitler,[1,2] D. Reuter,[3] A. D. Wieck,[3] and R. J. Haug[1]

[1]*Institut für Festkörperphysik, Universität Hannover, Appelstraße 2, 30167 Hannover, Germany*

[2]*High Field Magnet Laboratory, NSRIM, University of Nijmegen, Toernooiveld 7, 6525 ED Nijmegen, The Netherlands*

[3]*Lehrstuhl für Angewandte Festkörperphysik, Ruhr-Universität Bochum, Universitätsstraße 150, 44780 Bochum, Germany*

(Received 2 March 2004; published 7 July 2004)

Phonon excitations of fractional quantum Hall states at filling factors $\nu = 1/3$, $2/5$, $4/7$, $3/5$, $4/3$, and $5/3$ are experimentally shown to be based on Landau-level transitions of composite fermions. At filling factor $\nu = 2/3$, 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 2/3 ground state, making the energetically lower lying spin transition also allowed for phonon excitations.

DOI: 10.1103/PhysRevLett.93.026801 PACS numbers: 73.43.–f, 72.10.Di, 73.20.Mf

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 $\nu = p/q$ (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 an *even* denominator, $\nu = 1/2m$, these quasiparticles are formed by attaching an even number $2m$ of flux quanta ϕ_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 $\nu = p/q$ 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, $B_{\rm eff} = B - B(\nu = \frac{1}{2})$, with an effective (integer) filling factor, $p = \frac{\nu}{1-2\nu}$ [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 $\nu = 1/3$, $2/5$, $4/7$, $3/5$, $4/3$, and $5/3$ when varying the electron densities in the same sample over a wide range. At a given filling factor nearly all gaps measured show a square-root dependence on the magnetic field, strongly suggesting that we probe Landau-level transitions of CFs. All these data are described by one single fit parameter related to the CF effective mass. Surprisingly, the gap at $\nu = 2/3$ displays a *linear* field dependence, rather related to a spin transition forbidden for phonon excitations. This observation is a clear hint that the 2/3 ground state is not fully spin polarized in high B fields and that a strict separation between spin transitions and CF Landau-level transitions is no more possible.

Our sample consists of a high-mobility AlGaAs/GaAs heterojunction grown on a 2 mm thick GaAs wafer. On the front side containing the 2DES we patterned a $w = 90$ μm 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 ρ_{xx} contribution to the two-terminal resistance and thereby allows us to measure the smallest changes in ρ_{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 Ω) play only 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/V s) to $n = 1.50 \times 10^{15}$ m^{-2} (mobility $\mu = 193$ m^2/V s).

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 $T_{\rm 2DES} \lesssim 100$ mK for a cryostat base temperature of 75 mK.

The experimental setup is shown schematically in Fig. 1(a): 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 $\tau_{\rm H}$ through this heater, nonequilibrium phonon pulses are created at the heater-GaAs interface. They are characterized by a blackbody spectrum at a temperature $T_{\rm H} = (P_{\rm H}/\sigma A_{\rm H} + T_0^4)^{1/4}$, where $A_{\rm H}$ is the heater area, $P_{\rm H}$ 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 nonequilibrium phonons travel ballistically through the 2-mm-thick substrate. After a time of flight $\tau_{\rm LA} = 0.42$ μs and $\tau_{\rm TA} = 0.6$ μs for longitudinal acoustic (LA) phonons and for transverse acoustic (TA) phonons, respectively, they hit the 2DES and a small part of their energy is absorbed [13]. As a consequence, the 2DES temperature increases, directly measured by a

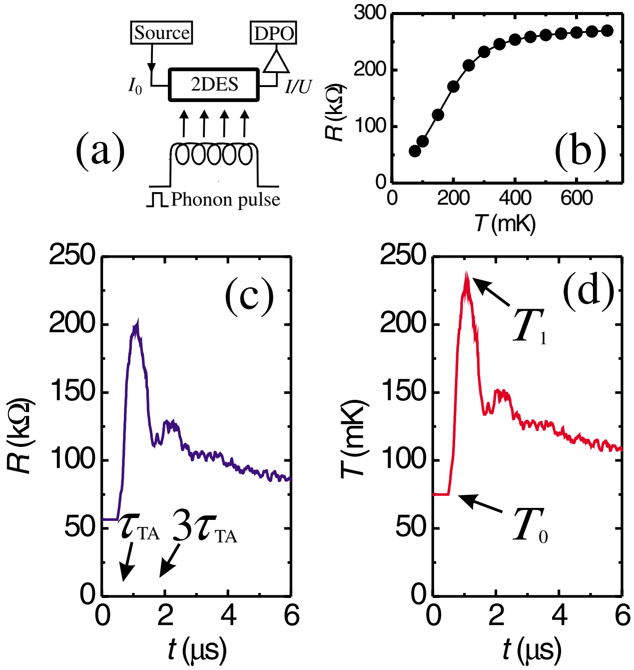

FIG. 1 (color online). (a) Schematic experimental setup with the phonon emitter on the back side and the 2DES at the front. Time resolved readout is achieved by an ultrafast digital phosphore oscilloscope (DPO). (b) Steady-state calibration of the 2DES resistance at $\nu = 2/3$ for $n = 0.89 \times 10^{15}$ m^{-2}. (c) Phonon signal for the same ν and n after emitting $T_{\mathrm{H}} = 2.09$ K phonons during $\tau_{\mathrm{H}} = 10$ ns at $t = 0$. The 2DES was excited from a base temperature $T_0 = 75$ mK. (d) 2DES temperature versus time as deduced from the raw phonon signal curve (c) using the calibration (b).

change of its resistance. This resistance change is detected by the current change at constant voltage using a 10-MHz current amplifier and a digital phosphor oscilloscope to average over up to a few 10^6 traces.

In Fig. 1(c) we have plotted a characteristic signal at filling factor $\nu = 2/3$. Because of their strong focusing [11] mainly TA phonons are visible, with a first peak starting around $\tau_{\mathrm{TA}} = 0.6$ μs and a second due to multiply reflected phonons after $3\tau_{\mathrm{TA}}$. 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. 1(b)] 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 (≈ 10 μs), all nonequilibrium 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, $P_{\mathrm{H}}\tau_{\mathrm{H}}$, 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

$$C(T)dT = r(T, T_{\mathrm{H}})P_{\mathrm{H}}dt - P_{\mathrm{e}}(T, T_0)dt, \tag{1}$$

where $C(T)$ is the 2DES's specific heat, $r(T, T_{\mathrm{H}})P_{\mathrm{H}}$ 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 $P_{\mathrm{e}}(T, T_0)$ 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 a given fractional filling factor. The maximum 2DES temperature on the first ballistic phonon peak, T_1 [see Fig. 1(d)], 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 $r(T, T_{\mathrm{H}}) \rightarrow r_0$. Integrating Eq. (1) over the pulse length with these assumptions, we get

$$\int_{T_0}^{T1} C(T)dT = r_0 P_{\mathrm{H}}\tau_{\mathrm{H}}. \tag{2}$$

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. 2(a) 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. 2(a) are fits to the theoretical predictions for the specific heat of a 2DES $C_{\mathrm{2DES}} \propto \frac{1}{T^2} e^{-\Delta_C/T}$ [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].

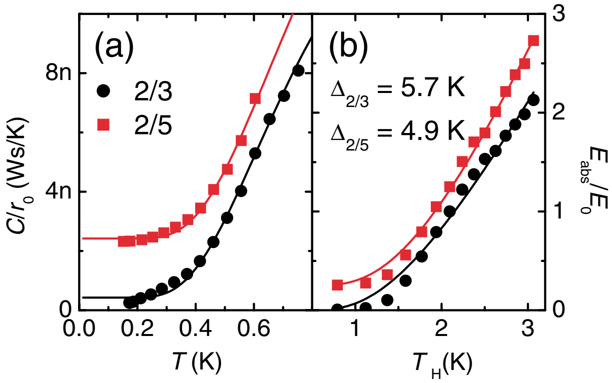

FIG. 2 (color online). (a) Examples of the relative specific heat at filling factors $\nu = 2/3$ and $2/5$ measured by phonon absorption for sample #5 ($n = 1.21 \times 10^{15}$ m^{-2}). Data for 2/5 are shifted for clarity by $2n$. The lines are fits according to Ref. [14] plus a constant. (b) Relative phonon energy absorbed by the 2DES, normalized to its value E_0 for $T_H = 2.09$ K, for a 10-ns pulse at filling factors $\nu = 2/3$ and $2/5$ (data for 2/5 are shifted by 0.25 for clarity). The curves are fits for an excitation across a gap $\Delta_{2/3} = 5.7$ K and $\Delta_{2/5} = 4.9$ K.

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 Δ, the total energy absorbed by the 2DES increases as

$$E_{abs} \propto \frac{1}{\exp(\Delta/T_H) - 1}. \quad (3)$$

The absorbed phonon energy, $E_{abs} = r(T, T_H)P_H\tau_H$, 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. 2(b) 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 $\pm 10\%$ of the average value. The solid lines in Fig. 2(b) show fits using Eq. (3) and indeed, the experimental data match a model with phonon excitation across a single energy gap Δ.

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 $4/7$ are clearly pronounced only for the highest electron concentration

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.

Sample #	n (10^{15} m^{-2})	μ (m^2/V s)	Δ (K) for $\nu =$ 2/3	2/5	1/3
1	0.89	102	3.2	4.9	9.2
2	0.95	109	4.1	5.3	8.2
3	1.03	119	4.5	5.9	8.9
4	1.13	131	4.8	5.8	⋯
5	1.21	144	4.9	5.7	⋯
6	1.36	168	6.0	⋯	⋯
7	1.46	187	6.7	⋯	⋯
8	1.50	193	6.6	⋯	⋯

and, accordingly, phonon gaps were measured only in sample #8 for these filling fractions.

All gaps measured at various filling factors and electron concentrations are compiled in Fig. 3(a). For comparison, we also show the energy gaps deduced from

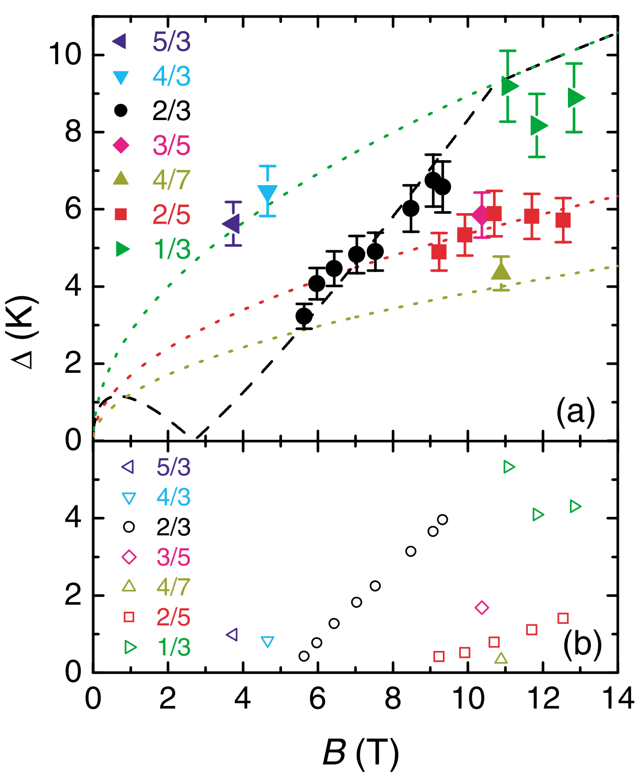

FIG. 3 (color online). (a) Compilation of all phonon excitation gaps Δ measured at various filling factors as a function of the magnetic field (i.e., the electron concentration for each filling factor). The dotted lines are CF Landau level transitions adjusted to our phonon data with one single mass parameter $\alpha = 0.158$ (± 0.006). The dashed line represents the expected form of a normally forbidden spin transition at $\nu = 2/3$. (b) The transport gaps are systematically lower than the phonon excitation gaps in (a).

activated transport measurements [15] in Fig. 3(b). In order to discuss these results for different filling factors ν we describe them within the CF picture [3], where the FQH filling factors $\nu = \frac{p}{2p+1}$ are mapped to integer CF filling factors p. The energy levels for $|p|$ filled CF Landau levels can then be written [16] as

$$E_{ns}(p) = (n + \tfrac{1}{2})\hbar\omega_c^*(p) + sg^*\mu_B B. \quad (4)$$

Here the CF cyclotron energy is given purely by the Coulomb interaction and thus follows the form $\hbar\omega_c^* = \hbar eB/(m^*(2p \pm 1))$ with a CF mass $m^* = m_e\alpha\sqrt{B[\mathrm{T}]}$ [17].

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 $\hbar\omega_c^*(p)$ can now be adjusted to our data by one single dimensionless mass parameter α. In Fig. 3 the fits of such Landau-level transitions to all the data at $\nu = 1/3$, $2/5$, $4/7$, $3/5$, $4/3$, and $5/3$ are shown yielding $\alpha = 0.158\ (\pm 0.006)$ (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 α is in astonishing agreement with the theoretical predictions in Eq. (1) of Ref. [17].

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 $B_p \approx 2.8$ 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 remarkable agreement [dashed line in Fig. 3(a)]. 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. 3(b). 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 cannot 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."

[1] D. C. Tsui, H. L. Stormer, and A. C. Gossard, Phys. Rev. Lett. **48**, 1559 (1982); H. L. Stormer, D. C. Tsui, and A. C. Gossard, Rev. Mod. Phys. **71**, S298 (1999).
[2] R. B. Laughlin, Phys. Rev. Lett. **50**, 1395 (1983); Rev. Mod. Phys. **71**, 863 (1999).
[3] J. K. Jain, Phys. Rev. Lett. **63**, 199 (1989).
[4] B. I. Halperin, P. A. Lee, and N. Read, Phys. Rev. B **47**, 7312 (1993).
[5] S. M. Girvin, A. H. MacDonald, and P. M. Platzman, Phys. Rev. Lett. **54**, 581 (1985); Phys. Rev. B **33**, 2481 (1986).
[6] R. K. Kamilla, X. G. Wu, and J. K. Jain, Phys. Rev. Lett. **76**, 1332 (1996); Phys. Rev. B **54**, 4873 (1996).
[7] A. Pinczuk, J. P. Valladares, D. Heiman, A. C. Gossard, J. H. English, C. W. Tu, L. Pfeiffer, and K. West, Phys. Rev. Lett. **61**, 2701 (1988).
[8] Moonsoo Kang, A. Pinczuk, B. S. Dennis, M. A. Eriksson, L. N. Pfeiffer, and K. W. West, Phys. Rev. Lett. **84**, 546 (2000).
[9] I. V. Kukushkin, K. v. Klitzing, and K. Eberl, Phys. Rev. Lett. **82**, 3665 (1999).
[10] C. J. Mellor, R. H. Eyles, J. E. Digby, A. J. Kent, K. A. Benedict, L. J. Challis, M. Henini, C. T. Foxon, and J. J. Harris, Phys. Rev. Lett. **74**, 2339 (1995).
[11] U. Zeitler, A. M. Devitt, J. E. Digby, C. J. Mellor, A. J. Kent, K. A. Benedict, and T. Cheng, Phys. Rev. Lett. **82**, 5333 (1999).
[12] A. J. Kent, R. E. Strickland, K. R. Strickland, and M. Henini, Phys. Rev. B **54**, 2019 (1996).
[13] F. Schulze-Wischeler, U. Zeitler, F. Hohls, R. J. Haug, D. Reuter, and A. D. Wieck, Physica (Amsterdam) **12E**, 474 (2002).
[14] T. Chakraborty and P. Pietiläinen, Phys. Rev. B **55**, R1954 (1997).
[15] F. Schulze-Wischeler, E. Mariani, F. Hohls, and R. J. Haug, Phys. Rev. Lett. **92**, 156401 (2004).
[16] E. Mariani, R. Mazzarello, M. Sassetti, and B. Kramer, Ann. Phys. (Berlin) **11**, 926 (2002).
[17] K. Park and J. K. Jain, Phys. Rev. Lett. **80**, 4237 (1998).
[18] V. M. Apalkov, T. Chakraborty, P. Pietiläinen, and K. Niemelä, Phys. Rev. Lett. **86**, 1311 (2001).
[19] N. Freytag, Y. Tokunaga, M. Horvatić, C. Berthier, M. Shayegan, and L. P. Lévy, Phys. Rev. Lett. **87**, 136801 (2001).