RAPID COMMUNICATIONS

PHYSICAL REVIEW B VOLUME 41, NUMBER 11 15 APRIL 1990-I

Evidence for a spin transition in the $\nu = \frac{2}{3}$ fractional quantum Hall effect

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(Received 29 December 1989)

Magnetotransport measurements on a low-density two-dimensional electron system have revealed a reentrant dependence of the activation energy on magnetic field for the fractional quantum Hall state at $\frac{2}{3}$ filling of the lowest Landau level. The data are consistent with a change in the spin structure of the ground state at $\frac{2}{3}$ filling, but do not provide a simple picture of the quasiparticle excitation process. As yet we find no similar effect for the $\nu = \frac{2}{5}$ state.

The spin degree of freedom in two-dimensional electron systems (2D ES) at high magnetic fields can play an essential role in forming the many-body ground state and its quasiparticle excitations. While the original theoretical description[1] of the fractional quantum Hall effect (FQHE) assumed that the large Zeeman energy completely polarized the spins, it was soon suggested[2] that at lower fields spin reversal should be considered. The first evidence for an *unpolarized* ground state in the FQHE came from tilted-field experiments[3] on the even-denominator state at filling fraction $\nu = \frac{5}{2}$. The collapse of the $\frac{5}{2}$ state upon tilting suggests an unpolarized ground state since the primary effect of the added in-plane field is an enhancement of the spin-flip energy while the Coulomb correlation energies remain roughly fixed. Among the more conventional odd-denominator FQHE states, recent experiments[4,5] have revealed apparent phase transitions in both the $\frac{8}{5}$ and $\frac{4}{3}$ ground states. These transitions, driven by tilting the magnetic field, suggest a change from an unpolarized ground state at small total magnetic fields (small angle) to a polarized state at higher total fields.

The FQHE at $\frac{2}{5}$ filling was recognized early on to present a good possibility for observation of an unpolarized ground state. Halperin[2] proposed a Laughlin-like wave function for such a state and subsequent small-system calculations[6] confirmed that the ground state was unpolarized, at least in the absence of the Zeeman energy. As later noted by Haldane,[7] just as the Laughlin $\nu = \frac{1}{3}$ state can be viewed as correlations applied to the fully polarized $\nu = 1$ state, the *same* correlations applied to the unpolarized $\nu = 2$ fully-filled lowest Landau level yield an unpolarized $\nu = \frac{2}{5}$ ground state. In the absence of the Zeeman energy, Halperin's unpolarized $\frac{2}{5}$ state is a primitive FQHE state exactly analogous to the polarized Laughlin $\frac{1}{3}$ state. With finite Zeeman energy, however, some different, polarized $\frac{2}{5}$ state may lie lower in energy.

In contrast to this relatively clear theoretical situation at $\nu = \frac{2}{5}$, the $\nu = \frac{2}{3}$ FQHE state is not well understood and is often assumed to be a polarized particle-hole conjugate of the $\frac{1}{3}$ state. The first tilted-field studies[8,9] of the $\frac{2}{3}$ state, at magnetic fields above 6 T, while showing an asymmetry between the $\frac{2}{3}$ and $\frac{1}{3}$ states, did not provide evidence for an unpolarized ground state. Only recently have numerical studies[10-12] suggested that again, in the absence of Zeeman energy, the ground state at $\nu = \frac{2}{3}$ might be unpolarized.

The results of Eisenstein *et al.*[4] on the $\nu = \frac{8}{5}$ state were taken as representative of the expected physics at $\frac{2}{5}$ since, neglecting Landau-level mixing, the correct particle-hole conjugate of $\frac{2}{5}$ is $2 - \frac{2}{5} = \frac{8}{5}$. Only if spin mixing can be neglected is the symmetry between ν and $1-\nu$. The reentrant dependence of the $\nu = \frac{8}{5}$ activation energy on total magnetic field was cited[4] as strong evidence for the expected phase transition from an unpolarized to a polarized ground state. Somewhat similar results were obtained by Clark *et al.*[5] for the $\nu = \frac{4}{3}$ state (conjugate to the $\frac{2}{3}$ state).

In this paper we present the first detailed activation energy study of the $\nu = \frac{2}{3}$ state in which a clear reentrant transition is observed. In addition to driving the transition via tilted magnetic fields, we present the first data in which the transition is swept at *zero tilt* using a back-gate bias to change the 2D carrier concentration. We present a simple model relating the transitional total magnetic field as measured in these two types of experiments. While the reentrant behavior seen at $\frac{2}{3}$ is crudely similar to that seen at $\frac{8}{5}$ or $\frac{4}{3}$, it differs in detail. Clark *et al.*[13] have also observed reentrant behavior at $\nu = \frac{2}{3}$, although only the fixed-temperature tilt dependence of the resistivity minimum has been reported. In common with earlier work,[8] we find no evidence of any reentrance in the primitive $\nu = \frac{1}{3}$ state. Finally, at filling fraction $\frac{2}{5}$ our data so far shows no reentrant structure and the observed Arrhenius plots for the resistivity minimum are anomalous. Thus the expected phase transition at $\frac{2}{5}$ filling remains unobserved.

The sample used in this work is a conventional $GaAs/Al_xGa_{1-x}As$ heterostructure grown by molecular-beam epitaxy. Low-temperature magnetotransport measurements are performed in the dark, i.e., in the absence of additional light-induced carriers. A back-gate bias voltage V_g is employed to allow continuous variation of the carrier concentration from about $N = 2.6 \times 10^{10}$ cm^{-2} at $V_g = -200$ V to 7×10^{10} cm^{-2} at $V_g = +130$ V. Over this same range the measured mobility increases, roughly linearly with density, from about 2.2 to 6×10^6 cm^2/V s. Employing 10 nA, 5 Hz excitation currents avoids any observable electron heating throughout the temperature range of the measurements. For the tilted-field studies we employ an *in situ* rotation device and measure the angle θ

between the normal to the 2D ES plane and the applied magnetic field by observing the $\cos\theta$ shift of strong features in the diagonal resistivity ρ_{xx}.

Figure 1 shows representative ρ_{xx} traces obtained under different back-gate bias and tilt-angle conditions. The two cases correspond to densities $N=4.0$ and 6.5×10^{10} cm^{-2} with mobilities $\mu\sim3.5$ and 5.5×10^{6} cm^2/Vs. In both cases prominent FQHE features are indicated. The enhanced FQHE structure apparent in the higher density trace ($V_g=+100$ V) is due in part to the higher mobility. Also important is the increased magnitude of the Coulomb energy $e^2/\varepsilon l_0 \propto B^{1/2}$ (l_0 being the magnetic length) for any given FQHE state.

The conclusions of this paper are based on measurements of the temperature dependence of ρ_{xx} for various FQHE states. Assuming the activated form $\rho_{xx}=\text{const}\times\exp(-\Delta/2T)$ allows determination of the energy gap Δ taken to be the energy required for creating a quasielectron-quasihole pair out of the condensate. At $\nu=\frac{2}{3}$ we find ρ_{xx} obeys this activated form over almost two decades, except under special conditions which will be described below. For the FQHE at $\nu=\frac{2}{5}$, however, we find strong deviations from this dependence at almost all densities and angles examined.

Arrhenius plots for the $\frac{2}{3}$ state are exhibited in Fig. 2. These are representative of the many obtained at this filling for various values of V_g and θ. In Fig. 2(a) the data shown were obtained at zero tilt ($\theta=0$) at three different values of V_g. In Fig. 2(b), however, the data were obtained at constant V_g (i.e., constant density N) for three different tilt angles θ. In both cases the intermediate data set has the smallest activation energy. As the figure shows, while the Arrhenius plots are typically linear for well over a decade in resistivity, those obtained near the transition are not. For these, we have estimated Δ from the linear portion of the plots. This nonlinearity is found only in a narrow window about the bottom of the transition (either in angle or density) and is reminiscent of the anomalous behavior observed at the $\nu=\frac{8}{5}$ transition.[4]

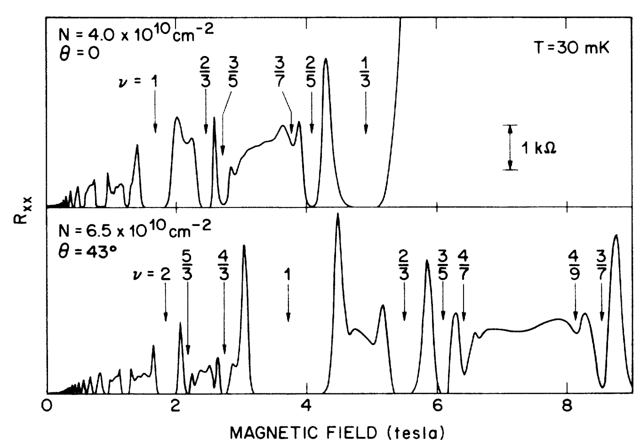

FIG. 1. Longitudinal magnetoresistance at $T=30$ mK for same sample at different densities and angles. Top trace obtained with $\theta=0$ and $V_g=-100$ V. Lower trace at $\theta=43°$ and $V_g=+100$ V. Major fractional features are indicated.

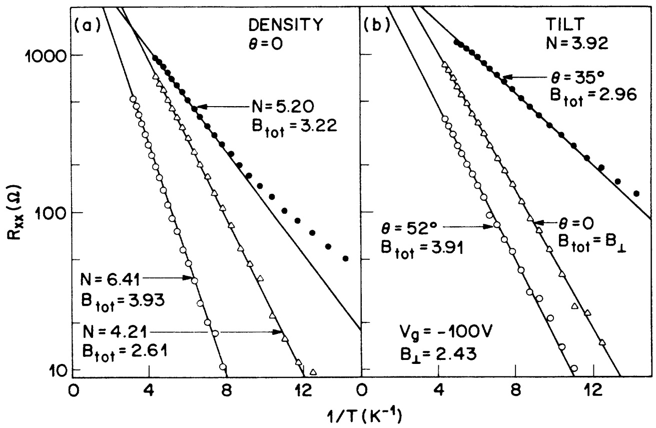

FIG. 2. Arrhenius plots showing reentrant behavior of the $\nu=\frac{2}{3}$ FQHE state. (a) Three densities at $\theta=0$. (b) Three angles at constant density. Densities in units of 10^{10} cm^{-2}, magnetic fields in teslas.

Unlike the $\frac{8}{5}$ case, however, we do not resolve a splitting of the resistivity minimum at the $\frac{2}{3}$ transition.

The overall picture of this reentrant behavior of the $\frac{2}{3}$-state energy gap is displayed in Fig. 3 where three plots of energy gap versus total magnetic field, B_{tot}, are shown. The solid circles were obtained at $\theta=0$, shifting the density, and thus the magnetic field for the $\frac{2}{3}$ state, via the back-gate bias V_g. The gap Δ, at first rising slowly from $B_{\text{tot}}=2.5$ T ($V_g=-100$ V, $N=4.0\times10^{10}$ cm^{-2}), begins to fall at around 2.9 T, reaching a minimum at about 3.3 T. On increasing B_{tot} further, Δ rises steadily to beyond 4.3 T ($V_g=130$ V, $N=7\times10^{10}$ cm^{-2}). A similar reentrant behavior of the $\frac{2}{3}$-state gap is observed on tilting the

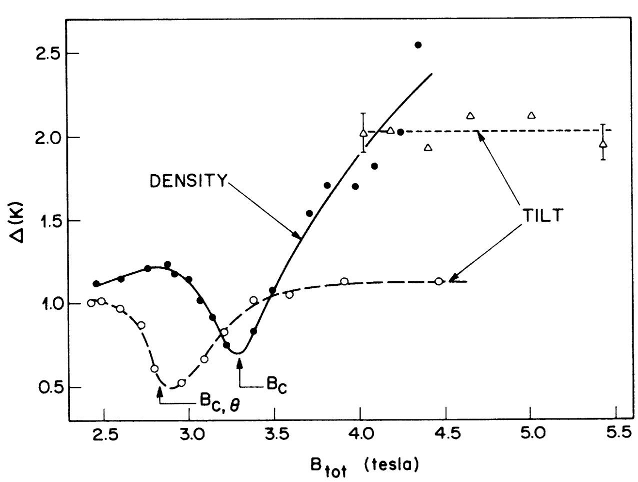

FIG. 3. Energy gap vs total magnetic field. Solid circles are obtained at $\theta=0$, varying density via back-gate bias. Open circles and triangles show tilt dependence at two different fixed densities. Open circles: $N=3.9\times10^{10}$ cm^{-2}, $B_\perp=2.43$ T. Triangles: $N=6.5\times10^{10}$ cm^{-2}, $B_\perp=4.02$ T. Error bars represent typical scatter in gap determinations made at high back-gate bias.

sample at fixed density (fixed V_g). The open circles in Fig. 3 were obtained at $V_g = -100$ V. For these, and other tilt data taken at constant density, the perpendicular magnetic field $B_\perp = B_{\rm tot}/\cos\theta$ is constant, fixed by the filling fraction $\nu = Nh/eB_\perp = \frac{2}{3}$. As shown in the figure, the energy gap Δ at $V_g = -100$ V exhibits a minimum at $B_{\rm tot} \sim 2.9$ T, well below that seen in the $\theta = 0$ data set. Beyond the minimum, Δ rises and then levels off, in contrast to the continuing increase seen in the zero-tilt data. Although omitted from the figure for clarity, an analogous tilt dependence is also observed at $V_g = -40$ V; the minimum occurring at $B_{\rm tot} \sim 3.1$ T and the high-field plateau value of Δ being somewhat higher than that obtained with $V_g = -100$ V. Last, by increasing the density to well beyond that showing the minimum gap in the $\theta = 0$ data, the reentrant tilt dependence of the gap can be suppressed. Although the gap determinations were found to exhibit greater scatter at high back-gate bias, the triangles shown in Fig. 3 (obtained at $V_g = +100$ V) reveal a gap essentially independent of $B_{\rm tot}$ at this density.

These observations suggest a transition in the spin configuration of the FQHE state at $\frac{2}{3}$ filling. Since the transition may be driven either by changing the density or the sample tilt angle, the possibility that it arises from an artifact of the in-plane magnetic field component unrelated to the spin of the system is highly unlikely. A plausible model for our results assumes that the $\frac{2}{3}$ ground state in this sample is spin unpolarized for 2D densities below a critical value $N_c \sim 5.3 \times 10^{10}$ cm^{-2}. This value corresponds to the magnetic field $B_c = 3hN_c/2e$ at which the zero-tilt energy gap (solid circles, Fig. 3) is a minimum, about 3.3 T. Above this density a polarized ground state takes over. The condensation energy per particle for each phase contains a Coulomb term proportional to $e^2/\epsilon l_0$. While carrying a different prefactor in each phase, this term depends only on the perpendicular magnetic field (equivalently, the density) and scales as $B_\perp^{1/2}$. In the polarized phase there is an additional stabilizing Zeeman term of the form $-\frac{1}{2} g\mu_B B_{\rm tot}$. Setting the condensation energies of the two phases equal results in an equation of the form

$$\alpha B_\perp^{1/2} = \tfrac{1}{2} g\mu_B B_{\rm tot}\,, \tag{1}$$

where α is the *difference* between the prefactors of the Coulomb terms for the unpolarized and polarized phases, respectively, g is an appropriate g factor, and μ_B is the Bohr magneton. For the zero-tilt case Eq. (1) gives $B_c^{1/2} = 2\alpha/g\mu_B$. Employing recent numerical estimates[11] of α and the bare g factor for electrons in GaAs ($g \sim 0.4$), we obtain $B_c \sim 11$ T, well in excess of the 3.3 T value mentioned above. This discrepancy may not be significant owing to several approximations made in the theoretical estimate. Important among these is the neglect of the finite thickness of the 2D electron wave functions. This effect reduces all Coulomb energies significantly, reducing B_c and making it sample dependent.

Within this model, if the 2D density N is less than N_c, at zero tilt the $\frac{2}{3}$ state is unpolarized. Tilting will drive the system into the polarized phase at a total magnetic field $B_{\rm tot} = B_\perp/\cos\theta$ that is less than B_c. This follows because tilting holds the Coulomb terms constant as the Zeeman energy increases. The critical field observed via tilt is found from Eq. (1) to be $B_{c,\theta} = (B_\perp B_c)^{1/2}$. The arrow in Fig. 3 gives this result, $B_{c,\theta} = 2.83$ T, for the $V_g = -100$ V data (open circles), assuming $B_c = 3.3$ T. Although lying slightly below the minimum at 2.9 T for this data set, the rough correspondence supports the overall model. A similar level of agreement is observed for analogous data with $V_g = -40$ V. Last, for $N > N_c$ the $\frac{2}{3}$ state begins in the polarized phase and tilting cannot drive a ground-state spin transition; this is the situation for the triangles in Fig. 3.

The data do not, however, evoke a simple picture for the spin states of the quasiparticle excitations, except at the highest magnetic fields. This is in contrast to that observed by Eisenstein *et al.*[4] in tilted-field studies of the $\nu = \frac{8}{5}$ state. In that case the energy gap Δ exhibited a simple linear dependence on $B_{\rm tot}$ in both phases. From the slope $d\Delta/dB_{\rm tot}$ in each phase a g factor quite close to 0.4, the bare GaAs value, was obtained if simple $\Delta S = \pm 1$ transitions were assumed. As Fig. 3 shows, such a concise description does not fit our new data at $\nu = \frac{2}{3}$.

For the zero-tilt data, a simple description is unlikely in any case since changing the density alters the Coulomb and Zeeman energies simultaneously. It also changes the sample mobility and this is well known[14] to influence FQHE gaps. Tilting, however, allows (approximately) independent control of the spin Zeeman term. As the open circles in Fig. 3 show, the $\frac{2}{3}$ state energy gap under tilt does not exhibit simple linear dependences on $B_{\rm tot}$. The plateau observed at high $B_{\rm tot}$ suggests that spin preserving, $\Delta S = 0$, quasiparticle excitation is operative. This is obviously expected for a polarized system at sufficiently high fields. On either side of the minimum gap, around 2.9 T for the $V_g = -100$ V data, the derivative $d\Delta/dB_{\rm tot}$ reaches values several times larger than that expected if $\Delta S = \pm 1$ transitions with $g \sim 0.4$ dominate. In this region we have no clear picture of the quasiparticle excitation process. Based on finite-system calculations, Chakraborty[12] has claimed that the $\frac{2}{3}$ state exhibits zero gap over a magnetic-field window surrounding the spin transition in the ground state. Our data do not support this insofar as a strong $\frac{2}{3}$ state is observable at all fields, but perhaps the large variation of Δ observed in this region is related to an anomaly like that suggested by Chakraborty.[12]

According to particle-hole symmetry, an analogous transition should be present at $\nu = \frac{4}{3}$. Recent work by Clark *et al.*[5] shows reentrant behavior at $\frac{4}{3}$ in a sample of considerably higher density than ours. In that experiment the $\frac{4}{3}$ state *does* appear to be absent over a range of magnetic fields. It is not clear whether this is an intrinsic effect or is due to the disorder in the sample. In our sample we find the $\frac{4}{3}$ state to steadily weaken under tilt. Apparently the reentrant phase lies at higher field than we can obtain. Clark *et al.*[5] further claim the $\nu = \frac{7}{5}$ daughter state disappears just as the $\frac{4}{3}$ state reappears at high angles. In the present case, the relevant daughter to the $\frac{2}{3}$ state is $\nu = \frac{3}{5}$. We find this state fully formed both below and above the $\frac{2}{3}$ transition, as Fig. 1 reveals. The relation, if any, between the $\frac{2}{3}$ and $\frac{4}{3}$ transitions requires further study.

Finally, we turn to the $\nu=\frac{2}{5}$ FQHE state. We find the Arrhenius plots for this state to be highly nonlinear at all angles and densities studied. Unambiguous energy gap determinations are thus not yet possible. This is in sharp contrast to the situation we find at $\nu=\frac{2}{3}$ and that observed earlier[4] at $\nu=\frac{8}{5}$. In both of those cases, Arrhenius plots linear for well over one decade in resistivity were found, except in narrow ranges surrounding the transition. We have, however, examined the tilt dependence of the $\nu=\frac{2}{5}$ Arrhenius plots for several densities N, the lowest producing $\nu=\frac{2}{5}$ at 4.1 T at $\theta=0$. In no case has a clear reentrant behavior been found. At first sight, this is puzzling since a transition has been observed[4] in the conjugate state $\nu=\frac{8}{5}$. That observation, however, was made in a sample of much higher density ($N=2.3\times10^{11}$ cm^{-2}) than the present one, for which the $\frac{8}{5}$ state is too weak to study (occurring around 1.3 T). As already mentioned, there are good reasons to expect the transition field B_c to be sample dependent. Even in the same sample B_c may not be the same for the conjugate states ν and $2-\nu$. While we find a $\frac{2}{3}$ transition around 3 T, the $\frac{4}{3}$ state shows no reentrance out to 9 T. Recent numerical work[11] suggests that the energetic advantage of the unpolarized ground state over the polarized state (in the absence of the Zeeman energy) for the $\frac{2}{3}$ state exceeds that for the $\frac{2}{5}$ state. On this basis one expects a lower B_c for $\frac{2}{5}$ than $\frac{2}{3}$, which puts it out of reach with this sample.

To summarize, we have observed a reentrant dependence of the $\nu=\frac{2}{3}$ FQHE energy gap on density and tilt angle. Our data are suggestive of a spin Zeeman origin for the effect. We have proposed a model in which the ground state at this filling fraction makes a transition from being spin unpolarized at low total magnetic fields to being polarized at higher fields. The relation of between the critical field observed at zero tilt by changing the density and that seen by tilting the sample is consistent with the model. The nature of the quasiparticle excitation in the two phases remains obscure, raising the possibility that the transition has some other origin. We have also searched for a transition in the $\nu=\frac{2}{5}$ FQHE; our present results suggest that, if it occurs, it lies below 4 T.

It is a pleasure to thank A. H. MacDonald for discussions and M. A. Chin and K. W. Baldwin for their technical support.

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