Microscopic manifestation of the spin phase transition at filling factor $2/3$

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First-order phase transitions are known to be accompanied by hysteresis due to the formation of domains with different order parameters. As the transition is approached, the domain sizes increase to macroscopic dimensions. It is therefore natural to expect that hysteresis will be absent from samples of microscopic size. Here, we explore from a microscopic standpoint the hysteretic behaviour across the spin phase transition that occurs at filling $\nu = 2/3$ in the fractional-quantum-Hall regime^{1–5}. Using a single-electron transistor, we follow the evolution of localized states across the spin transition by measuring the local compressibility. Localized-state spectra clearly reveal the hysteretic behaviour accompanying the transition. Using electrostatic gating we continuously vary the size of the sample undergoing the phase transition. For submicrometre dimensions the hysteresis disappears, indicating domain sizes in excess of 500 nm.

The weak Zeeman energy in GaAs enables us to investigate various fractional quantum-Hall phases with different spin polarizations^{1–4}. Recent photoluminescence measurements confirmed directly that indeed many of the fractions such as $2/3$ and $2/5$ are not necessarily spin polarized⁵. In the composite-fermion (CF) picture^{6,7}, $\nu = 2/3$ corresponds to two filled Landau levels (LLs) of the CF. Each level is characterized by its orbital ($n = 0, 1, 2, \dots$) and spin quantum number ($\uparrow$ or $\downarrow$). When $\nu = 2/3$ occurs at low enough magnetic field, the two occupied levels correspond to the two possible spin orientations: levels $(0, \uparrow)$ and $(0, \downarrow)$. As the magnetic field B increases, the energy spacing between these different spin states increases linearly with B whereas the energy separation to the next orbital level increases only as $\sqrt{B}$. Therefore, a degeneracy between the spin-down state of the lowest LL $(0, \downarrow)$ and the spin-up state of the second LL $(1, \uparrow)$ may occur, as illustrated in Fig. 1a, and the system will undergo a first-order quantum phase transition to a fully polarized state. The disappearance of dissipationless current flow due to the closure of the energy gap signals this transition in transport^{3,4}. Figure 1b shows an example at $\nu = 2/3$. We can form the analogy with easy-axis ferromagnetism by assigning an isospin to each of the degenerate levels^{8–10}. As in ferromagnetism, the degeneracy gives rise to domains of different isospin that are separated by domain walls. So far, the ferromagnetic character of this transition has been revealed through hysteretic transport^{11–14} and NMR

Figure 1 Composite-fermion Landau level diagram and transport data. **a**, Energy dependence of the CF LLs as a function of B field (see the Methods section). The dashed red curve indicates the Fermi energy at filling $2/3$. For $B < B_c$ the lowest two occupied levels are $(0, \uparrow)$ and $(0, \downarrow)$ and the system is unpolarized. For $B > B_c$ the $(1, \uparrow)$ level becomes lower in energy than $(0, \downarrow)$ and the system becomes fully polarized. **b**, Colour plot of the longitudinal resistance in the B -field–density plane (obtained from sample A, described in the Methods section). The density is measured relative to $n_{2/3} = (2/3)(B/\phi_0)$. At B_c the quantum Hall effect disappears and the resistance becomes finite as the system undergoes the spin transition.

spectroscopy¹⁵ on macroscopic samples. Here, using a single-electron transistor (SET), we explore the $\nu = 2/3$ transition from a microscopic standpoint.

Figure 2 Compressibility measurement across the spin transition. **a**, Schematic drawing of the measured inverse compressibility using a SET as a function of both density and B field. Each black line corresponds to a discrete charging event of a dot or antidot that appears near integer and fractional fillings. The charging spectrum near filling 1 corresponds to single-electron additions. Near fillings $1/3$ and $2/3$ the line spacing is three times denser and manifests the fractional quasi-particle charge $1/3e$. The solid blue line marks the degeneracy condition between $(1, \uparrow)$ and $(0, \downarrow)$. **b**, Schematic drawing of the expected behaviour for the localized states as a function of B and density (measured relative to $n_{2/3} = (2/3)(B/\phi_0)$) across the spin phase transition. The same spectrum of localized states is expected to appear on both sides of the transition. **c**, The local inverse compressibility acquired as a function of B field and carrier density (measured relative to $n_{2/3} = (2/3)(B/\phi_0)$) across the spin phase transition. The data were taken at a base temperature of 10 mK on sample B. Data points making up this colour graph were recorded by stepping B from 4.5 to 7 T in 5 mT steps and sweeping the density from high values to low values for each value of the field. Each dark line corresponds to the charging of a localized state. The spectra of localized states below the transition and above it are clearly different, indicating a change in the underlying screened potential across the spin phase transition. Conversely, below or above the transition the spectra are nearly invariant to changes in magnetic field. The white dotted lines serve as guides to the eye. Roughly within the 0.2 T field interval separating these white lines, the charging lines fade, bend or disappear and the spin transition takes place.

Localization in the integer- and fractional-quantum-Hall effect is governed by the competition between disorder and Coulomb interaction^{16,17}. Despite its complexity, a straightforward nonlinear screening model has been able to accommodate all previous experimental manifestations of localization in local compressibility measurements. For simplicity we review here the model near $\nu = 1$ and generalize it later to other integer and fractional fillings. At large B field far away from $\nu = 1$, the density of states within an LL is large enough for electrons to rearrange so as to completely screen

the external disorder potential. The resulting spatially varying electron density, $n(\mathbf{r}) = n_{av} + \Delta n(\mathbf{r})$, is invariant to changes in the B field or the average density, which can be tuned using a back gate. Here n_{av} is the average density and $\Delta n(\mathbf{r})$ the spatially varying part. As the filling increases towards unity, the required density for perfect screening in certain regions of the sample exceeds the maximal available density within the LL, $n_{max} = B/\phi_0$ (ϕ_0 is the flux quantum), and hence the local density saturates. As the density is increased further only small pockets with $n < n_{max}$ remain (see the

Figure 3 Hysteretic behaviour in the spectrum of localized states. The left and centre panels show the spectra of localized states measured on sample B as a function of magnetic field and density (measured relative to $n_{2/3} = (2/3)(B/\phi_0)$). In the left panel the spectra are measured while scanning the density upwards, whereas in the centre panel the density is scanned downwards. The hysteretic behaviour is emphasized in the right panel, where the difference between the two spectra is shown. The hysteretic behaviour is only present in the vicinity of the spin phase transition.

inset of Fig. 2b) and the addition and removal of charge into these pockets happens one particle at a time. These pockets form antidots as they have a lower density compared with their surroundings. Discrete charge additions can be readily observed by measuring the local compressibility with a SET (see the Methods section). Each charge addition appears as a compressibility spike in the n_{av} - B plane (Fig. 2). As B increases, n_{max} is raised as well and we need to increase the average density according to $\nu = 1$ to maintain exactly the same microscopic disorder potential. Therefore, spikes that correspond to antidots of $\nu = 1$ follow a single slope given by $\nu = 1$. As the density increases further, regions corresponding to the peaks in the density profile start occupying the next LL and dots form. Such dots, being surrounded by incompressible regions of $\nu = 1$, still follow the same slope in the n_{av} - B plane. The resulting spectrum of localized states, schematically shown in Fig. 2a, depends only on the bare disorder potential and the presence of a gap. Therefore, identical charging spectra, following a slope of the underlying filling, are seen for different integer fillings¹⁶. The same holds for fractional fillings except that the quasi-particle charge is fractional^{18–23}, giving rise, for example, to three times more charging lines for fillings $1/3$ and $2/3$ (ref. 17). As we approach the spin transition near filling $2/3$, the gap is expected to close. This in turn will lead to improved screening and the disappearance of the localized states as illustrated in Fig. 2a,b.

A typical measurement with a local SET near $\nu = 2/3$ is depicted in Fig. 2c. The sample and fabrication details are described in the Methods section. The plot shows a colour map of the local compressibility across the $(\Delta n_{2/3}, B)$ plane, where $\Delta n_{2/3}$ is the

deviation from the electron density at which the average filling is $2/3$, $\Delta n_{2/3} = n_{av} - (2/3)B/\phi_0$. A large set of dark lines departs from the left as well as the right. These compressibility spikes run parallel to the horizontal $\Delta n_{2/3} = 0$, marking constant average filling $2/3$. This is consistent with the localization scenario outlined previously. Each spike represents a single charging event filling a localized state near the SET in one of the dots or antidots, which form where the degeneracy of the Landau level being filled has been exhausted. Because this potential is rigid as we tune B , the configuration of dots and antidots remains identical at different B and spikes recur at the same $\Delta n_{2/3}$ irrespective of B . The resulting horizontal charging lines continue until the spin transition is approached. There, the gap closes and the prerequisite for observing localization, the existence of a gap, is no longer fulfilled. Hence, lines fade or vanish altogether rather abruptly. The transition seems slanted in the local compressibility just as in transport (Fig. 1). This can readily be understood by determining the zero-gap contour in the $(\Delta n_{2/3}, B)$ plane where the $(1, \uparrow)$ and $(0, \downarrow)$ levels coincide (see the Methods section). The local compressibility shares another important signature of the spin transition with transport data, namely hysteresis^{11–14}. Figure 3 plots the local compressibility in a narrow B -field region around the spin transition and compares data acquired during up and down sweeps of $\Delta n_{2/3}$. The right panel has been generated by subtracting the up and down sweeps to bring out the hysteresis explicitly. Further away from the transition, the hysteresis is nonexistent. Hysteresis is indicative of glassy behaviour in this disordered electron system when it undergoes the ferromagnetic first order quantum phase transition from the unpolarized to the completely polarized fractional quantum Hall liquid. Different domain configurations evolve as the path to approach the transition is varied. These microscopic differences obviously affect the local compressibility detected by the SET as it is sensitive to the spatial location of charging events and, hence, they are responsible for hysteresis. However, they also influence the macroscopic resistance measurement, for instance as a result of alterations in the accumulated domain-wall resistance^{9,24}. Sufficiently far away from the transition where domains with opposite spin configurations are no longer relevant, the behaviour is independent of the path taken. For completeness, we note that in transport studies it has been argued that changes in the nuclear spin polarization of the GaAs host can also contribute to hysteretic behaviour. Dynamic nuclear polarization by the externally imposed current, for instance, may produce such changes in the nuclear spin polarization near the transition^{11,25,26}. The local-compressibility experiments here, though, have been carried out in the absence of an imposed current.

Unlike transport, which inevitably senses a macroscopic area, local probe studies can focus on only a microscopic region. If the hysteresis is intimately related to domain formation, we anticipate its disappearance as the area undergoing the transition is scaled down. By applying a voltage difference, $V_{SET} - V_{2D}$, between the SET and the two-dimensional system, the electron density directly underneath the SET ($50 \times 500 \text{ nm}^2$) can either be enhanced or depleted compared with the surrounding bulk. By tuning this imposed density difference, it is possible to ensure that the bulk is still compressible and far away from filling $2/3$, whereas underneath the SET the density is closer to filling $2/3$ and screening is no longer capable of entirely compensating the bare disorder potential. Under these circumstances, the bulk will not contribute any $2/3$ -compressibility spikes and the observed features can be assigned solely to the artificial electron pocket or antidot defined by $V_{SET} - V_{2D}$. An example of compressibility data in this regime for up and down density sweeps is illustrated in Fig. 4. Hysteretic behaviour has vanished. This observation conforms to our expectations for a glassy system studied on a length scale

Figure 4 Absence of hysteresis when downscaling the active region. Using the SET as a local gate, we deform the density profile of the two-dimensional electron gas (2DEG) such that only the region underneath the SET (50 nm × 500 nm) undergoes the spin phase transition. The left and centre panels show the spectra of localized states taken during up and down sweeps of the density respectively. Unlike the data in Fig. 3, the two spectra are identical, indicating the absence of hysteresis in this effectively small specimen. The right panel shows the difference of the two spectra, emphasizing once more the absence of hysteresis. The arrowhead marks the local filling factor 2/3 (underneath the biased SET).

smaller than or comparable to the typical domain size. Because the monitored region is comparable to the SET size it is reasonable to conclude that domain structure is larger than 500 nm. By applying intermediate $V_{\text{SET}} - V_{\text{2D}}$ voltages it is possible to gradually tune the ‘active’ sample area. In the Supplementary Information, a graph is shown at a bias voltage in between the voltages at which the data of Figs 2 and 4 were taken to illustrate the gradual disappearance of hysteresis with shrinking ‘active’ sample area.

So far the intuitive localization picture governed by Coulomb interaction^{16,17} has performed remarkably well to describe the experiment. However, a notable inconsistency is exposed when carefully comparing the localized-state spectrum at low B fields deep in the unpolarized regime, with the spectrum well above the spin transition, where the system is fully spin polarized. Contrary to the expectation highlighted in Fig. 2a,b, the spectra on the two sides of the transition do not match. Charging lines, which vanish as the spin transition comes closer, do not re-emerge on the other side. Instead, new ones at different $\Delta n_{2/3}$ appear with no apparent correlation between the new and old values. In addition, the expected gradual disappearance of localized states when the gap closes is not observed. Rather, the spectrum disappears abruptly and is replaced by a new one. This observation cannot be reconciled with the above localization picture. Our experiment implies that the configuration of dots and antidots is modified, yet a change in the bare disorder potential on crossing the spin transition is not plausible. A more likely scenario is that the nature of the Coulomb interaction and screening alters when crossing the transition, because there LLs with different spin and orbital quantum numbers become occupied.

It is worth emphasizing that a comparison of spectra with different spin and orbital quantum numbers can also be made for integer fillings. However, there the spectra for different integer fillings would be collected at substantially different densities or B fields. Despite this constraint our earlier work¹⁶ shows qualitatively similar spectra for different integer fillings with only a few modifications to them. These modifications may be a result of the change in quantum numbers. The spin transition at filling 2/3 offers the unique opportunity of comparing charging spectra of different spin and orbital quantum numbers at the same B field and density. The screening-induced density profile, which compensates for the bare disorder potential when the system is compressible, determines the size and shape of the dots and antidots, which form when screening breaks down. Apparently, this density profile is significantly different below and above the transition. Because spikes follow horizontals, we conclude however that it remains largely unaltered when staying on the same side of the transition. We conjecture that the difference in the density profile below and above the transition stems from modified screening conditions. Differences in orbital wavefunction spreading as well as the gap due to exchange enhancement may produce differences in the polarizability of the electronic system²⁷. This discrepancy between the spectra on either side of the transition survives when zooming in on a small enough area, as in Fig. 4, for which the absence of hysteresis signals the absence of domains. Both observations substantiate that the mismatch between the localized-state spectra on the two sides of the spin transition is unrelated to domain physics near the transition itself.

METHODS

SAMPLE PREPARATION

Our studies were carried out on two separate samples: a single GaAs/Al_{0.32}Ga_{0.68}As heterostructure denoted as sample A or a 22 nm wide single GaAs quantum well, surrounded on either side by an Al_{0.32}Ga_{0.68}As barrier denoted as sample B. The two-dimensional electron systems were located at a distance of 240 nm (sample A) or 250 nm (sample B) from the top crystal surface. A single Si delta-doped layer, displaced 80 nm (sample A) or 90 nm (sample B) from the heterojunction towards the top surface, provided the charge carriers, with a density that varied in the dark between 8×10^{10} and 10×10^{10} cm⁻². A 200-nm-thick n⁺-GaAs layer underneath the two-dimensional electron system at a distance of approximately 1.8 μm served as an *in situ* grown backgate in both samples. It enables us to tune the carrier density. Square van der Pauw geometries with an ohmic contact in each corner were prepared with optical lithography. After electron-beam lithography, a two-angle shadow evaporation technique with an oxidation step in between was used to fabricate a pair of aluminium-based SETs on each mesa.

DETERMINATION OF THE ZERO-GAP CONTOUR

The spin transition takes place when the composite-fermion LLs ($1, \uparrow$) and ($0, \downarrow$) coincide (top panel Fig. 1). The energy of these levels is composed of a kinetic-energy term associated with the cyclotron motion and a Zeeman term describing the spin splitting: $E_{(0,\downarrow)} = \hbar\omega_c/2 + |g|\mu_B B/2$ and $E_{(1,\uparrow)} = 3\hbar\omega_c/2 - |g|\mu_B B/2$. Here, the cyclotron energy $\hbar\omega_c$ is equal to $\hbar e|B - 2n\Phi_0|/m^*$, with m^* the effective mass of composite fermions. Bearing in mind that composite fermions originate from Coulomb interaction, a dimensional analysis stipulates that m^* scales with the average inter-particle distance, $m^* = \beta\sqrt{n}$ (ref. 28). The proportionality constant β depends on the details of the system, such as the spatial extent of the wavefunction in the confinement direction. The Zeeman term contains the effective g -factor, which will be assumed constant, and the Bohr magneton μ_B . From the condition $E_{(0,\downarrow)} = E_{(1,\uparrow)}$, it is straightforward to derive an analytical expression for the critical magnetic field B_n^* at which the transition occurs as a function of a single variable, the density n : $B_n^* = 2n\Phi_0/(\alpha\sqrt{n} + 1)$. The constant α is a fit parameter, which, apart from fundamental constants, absorbs the unknown g -factor and proportionality constant β . This functional dependence can be fitted well to both the transport and compressibility data for the investigated

heterostructures. At very low temperatures, the thermal spin polarization of the nuclei of the crystal hosting the two-dimensional electron system is no longer negligible and the electronic Zeeman energy is affected through the hyperfine interaction. It adds a constant contribution to the Zeeman term, which produces moderate shifts of the critical fields. Here, it has been ignored for simplicity as its inclusion does not alter in any essential way the overall behaviour of the spin transition with increasing density.

MEASUREMENT OF LOCALIZED STATES

Unlike extended electrons, whose charge is spread over the entire volume of the system, a localized electron is confined to a small region in space, determined by its localization length. On population, a single electronic charge enters this confined region. Unlike filling an extended state, this results in a discrete jump in the local chemical potential μ and a spike in its derivative with respect to the back-gate voltage $\delta\mu/\delta V_{\text{BG}}$. Hence, in contrast to transport and macroscopic compressibility measurements, the local derivative $\delta\mu/\delta V_{\text{BG}}$, which is inversely proportional to the local electronic compressibility, provides direct information about the localized states.

To measure this derivative we use a SET electrometer. The current through the SET is a sensitive measure of the local electrostatic potential. Changes in the SET current directly reflect the changes in the local electrostatic potential underneath. When the 2DEG is in electrochemical equilibrium, the local chemical potential, μ , and the local electrostatic potential, φ , add up to a constant value determined by the electrochemical potential, $eV = \mu + e\varphi$, where e is the charge of the electron. Therefore, the measured changes in the local electrostatic potential are exactly equal and opposite in sign to the changes of μ , $d\mu + ed\varphi = 0$.

The inverse local compressibility can be directly measured by modulating the back-gate voltage with respect to the 2DEG. A uniform density modulation is thereby induced in the sample, $e\Delta n_{2\text{D}} = C\Delta V_{\text{BG}}$, where C is the capacitance between the back gate and the 2DEG. A local change in density causes a local change in the chemical potential, $\Delta\mu$. Because the 2DEG is under electrochemical equilibrium, this local change in chemical potential is obtained through a local change in the electrostatic potential, $\Delta\varphi$ and, hence, $\Delta n = (\partial n/\partial\mu)\Delta\mu = e(\partial n/\partial\mu)\Delta\varphi = e\kappa\Delta\varphi$, where κ is the compressibility. Therefore, by measuring the local change in potential induced by the back gate, we can infer directly the local inverse compressibility, $\partial\varphi/\partial V_{\text{BG}} = C/e^2\kappa$.

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References

- Halperin, B. I. Theory of the quantized Hall conductance. *Helv. Phys. Acta* **56**, 75–102 (1983).
- Xie, X. C., Guo, Y. & Zhang, F. C. Fractional quantum Hall effect with spin reversal. *Phys. Rev. B* **40**, R3487–R3490 (1989).
- Eisenstein, J. P., Stormer, H. L., Pfeiffer, L. N. & West, K. W. Evidence for a spin transition in the $\nu = 2/3$ fractional quantum Hall effect. *Phys. Rev. B* **41**, R7910–R7913 (1990).
- Engel, L. W., Hwang, S. W., Sajoto, T., Tsui, D. C. & Sayegan, M. Fractional quantum Hall effect at $\nu = 2/3$ and $3/5$ in tilted magnetic fields. *Phys. Rev. B* **45**, 3418–3425 (1992).
- Kukushkin, I. V., von Klitzing, K. & Eberl, K. Spin polarization of composite fermions: Measurements of the Fermi energy. *Phys. Rev. Lett.* **82**, 3665–3668 (1999).
- Heinonen, O. (ed.) *Composite Fermions: A Unified View of the Quantum Hall Regime* (World Scientific, Singapore, 1998).
- Das Sarma, S. & Pinczuk, A. (eds) *Perspectives on Quantum Hall Effects* (Wiley, New York, 1996).
- Jungwirth, T., Shukla, S. P., Smrcka, L., Shayegan, M. & MacDonald, A. H. Magnetic anisotropy in quantum Hall ferromagnets. *Phys. Rev. Lett.* **81**, 2328–2331 (1998).
- Jungwirth, T. & MacDonald, A. H. Resistance spikes and domain wall loops in Ising quantum Hall ferromagnets. *Phys. Rev. Lett.* **87**, 216801 (2001).
- MacDonald, A. H., Rajarman, R. & Jungwirth, T. Broken-symmetry ground states in $\nu = 2$ bilayer quantum Hall systems. *Phys. Rev. B* **60**, 8817–8826 (1999).
- Kronmüller, S. *et al.* New resistance maxima in the fractional quantum Hall regime. *Phys. Rev. Lett.* **81**, 2526–2529 (1998).
- Cho, H. *et al.* Hysteresis and spin transitions in the fractional quantum Hall effect. *Phys. Rev. Lett.* **81**, 2522–2525 (1998).
- Eom, J. *et al.* Quantum Hall ferromagnetism in a two-dimensional electron system. *Science* **289**, 2320–2323 (2000).
- Smet, J. H., Deutschmann, R. A., Wegscheider, W., Abstreiter, G. & von Klitzing, K. Ising ferromagnetism and domain morphology in the fractional quantum Hall regime. *Phys. Rev. Lett.* **86**, 2412–2415 (2001).
- Stern, O. *et al.* NMR study of the electron spin polarization in the fractional quantum Hall effect of a single quantum well: Spectroscopic evidence for domain formation. *Phys. Rev. B* **70**, 075318 (2004).
- Ilani, S. *et al.* The microscopic nature of localization in the quantum Hall effect. *Nature* **427**, 328–332 (2004).
- Martin, J. *et al.* Localization of fractionally charged quasi-particles. *Science* **305**, 980–983 (2004).
- Laughlin, R. B. Anomalous quantum Hall effect: An incompressible quantum fluid with fractionally charged excitations. *Phys. Rev. Lett.* **50**, 1395–1398 (1983).
- Goldman, V. J. & Su, B. Resonant tunneling in the quantum Hall regime: Measurement of fractional charge. *Science* **267**, 1010–1012 (1995).
- Franklin, J. D. F. *et al.* The Aharonov-Bohm effect in the fractional quantum Hall regime. *Surf. Sci.* **361–362**, 17–21 (1996).
- Saminadayar, L., Glatli, D. C., Jin, Y. & Etienne, B. Observation of the $e/3$ fractionally charged Laughlin quasiparticle. *Phys. Rev. Lett.* **79**, 2526–2529 (1997).
- Reznikov, M., de Picciotto, R., Griffiths, T. G., Heiblum, M. & Umansky, V. Observation of quasiparticles with one-fifth of an electron's charge. *Nature* **399**, 238–241 (1999).
- de Picciotto, R. *et al.* Direct observation of fractional charge. *Nature* **389**, 162–164 (1997).
- Brey, L. & Tejedor, C. Spins, charges, and currents at domain walls in a quantum Hall Ising ferromagnet. *Phys. Rev. B* **66**, R041308 (2002).
- Hashimoto, K., Muraki, K., Saku, T. & Hirayama, Y. Longitudinal resistance anomaly around the $2/3$ filling factor observed in a GaAs/AlGaAs single heterostructure. *Physica B* **298**, 191–194 (2001).
- Kraus, S. *et al.* From quantum Hall ferromagnetism to huge longitudinal resistance at the $2/3$ fractional quantum Hall state. *Phys. Rev. Lett.* **89**, 266801 (2002).
- Aleiner, I. L. & Glazman, L. I. Two-dimensional electron liquid in a weak magnetic field. *Phys. Rev. B* **52**, 11296–11312 (1995).
- Halperin, B. I., Lee, P. A. & Read, N. Theory of the half-filled Landau level. *Phys. Rev. B* **47**, 7312–7343 (1993).

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Competing financial interests

The authors declare no competing financial interests.

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