Anomalous-Filling-Factor-Dependent Nuclear-Spin Polarization in a 2D Electron System

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Spin-related electronic phase transitions in the fractional quantum Hall regime are accompanied by a large change in resistance. Combined with their sensitivity to spin orientation of nuclei residing in the same plane as the 2D electrons, they offer a convenient electrical probe to carry out nuclear magnetometry. Despite conditions which should allow both electronic and nuclear-spin subsystems to approach thermodynamic equilibrium, we uncover for the nuclei a remarkable and strongly electronic filling-factor-dependent deviation from the anticipated thermal nuclear-spin polarization.

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The spin of 2D electrons with density n_{2D} has long been recognized as a relevant degree of freedom in fractional quantum Hall (FQH) studies [1]. Both theory and experiment have established that numerous FQH states come in several distinct spin arrangements [2–4]. The competition between the Zeeman (E_Z) and Coulomb energy (E_C) settles a particular spin configuration. When maintaining fixed filling, E_C grows with $\sqrt{n_{2D}}$, whereas E_Z rises linearly. This disparate functional dependence offers the possibility to alter the relative strength of both energy scales and force transitions between different spin configurations by simply adjusting n_{2D}.

This spin transition physics can be lucidly phrased in single particle terms when invoking the composite fermion (CF) description of the FQH effect [5]. The FQH state at filling $\nu = 2/3$ serves as a prototypical example [3]. At this filling, two composite fermion levels are completely filled. The evolution of the energy spectrum with n_{2D} is schematically illustrated in the bottom right inset of Fig. 1. Each level is characterized by its orbital radius ($n = 0, 1, \dots$) and spin quantum number ($\uparrow, \downarrow$), and the chemical potential is marked by a black dashed line. At small n_{2D}, the ground state is unpolarized. At higher n_{2D}, E_Z increases more rapidly than the CF Landau quantization energy $\hbar\omega_c^* \propto E_C$. The fully spin-polarized ground state ensues when E_Z exceeds $\hbar\omega_c^*$. This phase transition leaves unambiguous fingerprints in the transport properties [6]. Because the gap diminishes as level $(0, \downarrow)$ overtakes the $(1, \uparrow)$ level, it is accompanied by a loss of quantization in the Hall resistance and the longitudinal resistance no longer vanishes.

The strong resistance change when the electronic system is poised between these two spin configurations can be exploited to detect minute changes in the spin polarization of nuclei residing in the same plane as the 2D electrons [7,8]. In the presence of nonzero average nuclear-spin orientation $\langle I \rangle$, the hyperfine interaction [9] causes an effective nuclear field B_N, which alters the electronic Zeeman energy $E_Z \propto (B + B_N)$, but leaves $\hbar\omega_c^*$ unaffected [see the shifted dashed blue $(0, \downarrow)$ level for $B_N \neq 0$]. Consequently, a gradual polarization of the nuclear spins will progressively relocate the 2/3-spin transition to a different carrier density. This shift is reflected in transport measurements when choosing a fixed working point near the spin transition, where a small change in E_Z/E_C induces a large variation in the resistance. Here, this technique is generalized to perform nuclear magnetometry over a wide range of filling factors in order to detect the nuclear-spin polarization "background" experienced by the 2D electrons as a function of the filling factor at which the sample is left to relax and equilibrate. Contrary to expectations, the polarization of the nuclear-spin subsystem $\langle I \rangle$ exhibits an astonishingly large variation with the electronic filling. As a matter of fact, $\langle I \rangle$ was anticipated to be equal to the thermal polarization—non-negligible at the temperature of the experiment—without any significant dependence on filling at all, except for the tiny changes produced by the Knight shift.

These studies were carried out on a field effect transistor with a 3 μm long and 250 μm wide channel, obtained with cleaved edge overgrowth. A cross section is depicted in Fig. 1. Fabrication details were described elsewhere [7,10]. Doped GaAs layers act as source (S) and drain (D). The gate voltage V_g electrostatically induces the two-dimensional electron system (2DES) at the GaAs/AlAs interface. The density can be tuned widely from depleted to beyond $4 \times 10^{11}/\mathrm{cm}^2$ by virtue of the contact scheme and the absence of modulation doping. The sample is cooled in a dilution unit with 20 mK base temperature. The source-drain resistance R_{SD} is measured with a lock-in technique and a 1 nA sinusoidal drive current (5.36 Hz). The excellent density tunability in this geometry goes at the expense of the ability to separately record the different resistivity tensor components. Condensation in a quantum Hall state is announced by the appearance of a plateau in R_{SD} or a maximum in case the state is not fully developed. In between quantum Hall

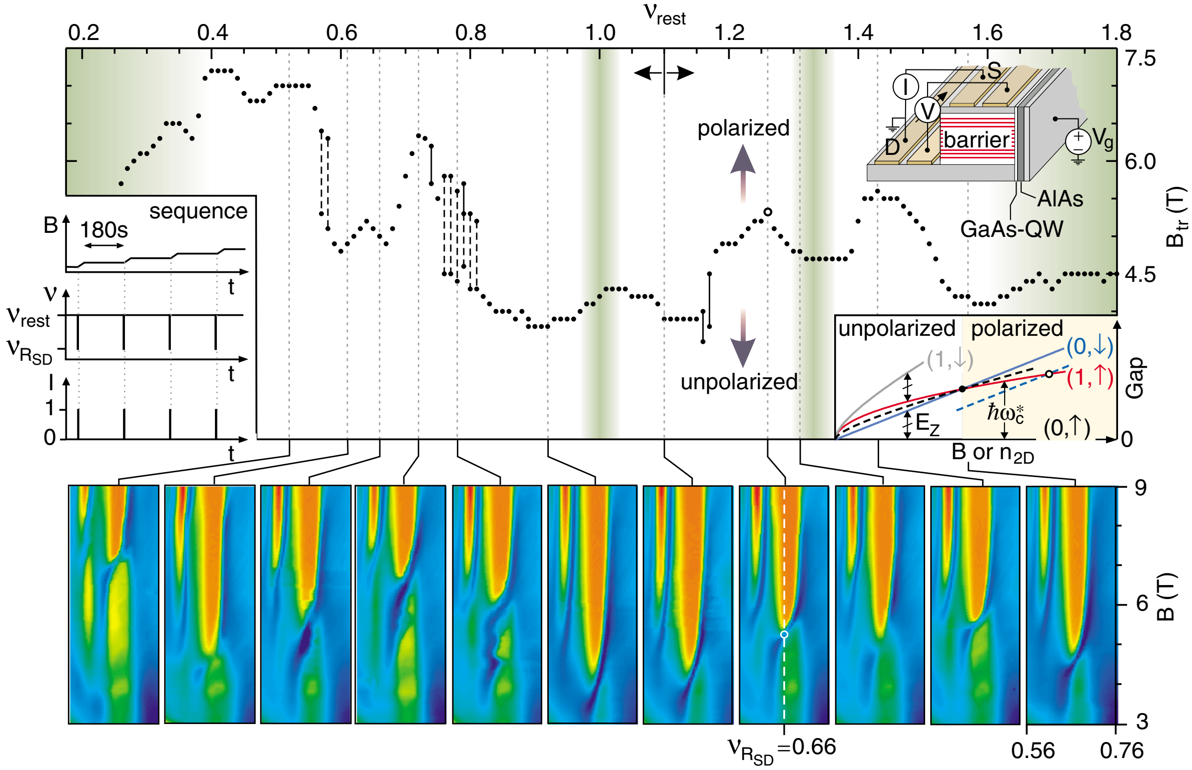

FIG. 1 (color). Location of the 2/3-spin transition as a function of filling ν_{rest} at which the sample is allowed to relax. Right insets: Sample cross section and evolution of the composite fermion Landau levels with density or B field at fixed $\nu = 2/3$. Left inset: Time sequence for current (in nA), B field, and filling factor (gate voltage). For each value of ν_{rest}, B is swept discontinuously from 3 to 9 T in 0.1 T steps. The gate voltage tracks B during sweep operations to maintain fixed filling ν_{rest} at all times. After a new field value has been set, the 2DES is left to relax for 180 s. Subsequently, the gate voltage is adjusted for a short excursion to $\nu_{R_{SD}} = 0.66$ and simultaneously the current is turned on. R_{SD} is recorded after 2 s to account for time constants of the signal recovery scheme. The current is turned off and the original filling ν_{rest} restored. B is run up to the next value and the entire procedure is repeated up to 9 T. The field value at which $R_{SD}(\nu_{R_{SD}} = 0.66)$ reaches a minimum (spin transition) is plotted at abscissa ν_{rest}. To acquire the next data point, the sample is kept at the upcoming value of ν_{rest} while ramping down the field back to 3 T. ν_{rest} is scanned downwards in 0.01 steps from 1.1 to 0.25 and upwards from 1.1 to 1.8. The bottom color panels are the outcome of a prolonged measurement sequence. They visualize the spin transition for selected values of ν_{rest} by plotting R_{SD} for the full B-field range and the filling factor interval covering the 2/3 and 3/5 FQH states. The color scale for R_{SD} spans values from $0.38h/e^2$ (blue) to $1.55h/e^2$ (red). Here also, a 180 s pause at ν_{rest} precedes the collection of every single data point making up these graphs. For some values of ν_{rest}, either a broad minimum or two separate minima are observed in $R_{SD}(\nu_{R_{SD}} = 0.66)$. This is pointed out by placing multiple dots at the same abscissa, connected either by a solid black line or a dashed black line, respectively. The color panel for $\nu_{rest} = 0.78$ illustrates this issue.

states, R_{SD} drops, because the nonzero longitudinal conductance short circuits the Hall voltage. When scanning B while maintaining constant filling $\nu = 2/3$, the transition from unpolarized to fully spin-polarized manifests itself in R_{SD} as a minimum [10].

The crux of this Letter is highlighted in Fig. 1. The location of the 2/3-spin transition is recorded depending on the resting filling factor ν_{rest} at which the sample is allowed to equilibrate: The 2DES is placed at the chosen ν_{rest}. B is swept from 3 to 9 T while maintaining this filling. For consecutive values of B, 0.1 T apart, the sweep is interrupted and the system is left to relax for 180 s. Subsequently, *only a brief excursion* from ν_{rest} to $\nu_{R_{SD}} = 0.66$ is carried out to probe R_{SD} *by rapidly adjusting* V_g. A minimum in the R_{SD} values acquired in this fashion as a function of B signals the transition [11]. Figure 2 shows an example for $\nu_{rest} = 0.52$. The spin transition occurs near $B_{tr} = 7$ T. This field is plotted on the ordinate of Fig. 1 at abscissa ν_{rest}. The detailed sequence of operations to collect the data is described in Fig. 1 and illustrated in a (B, ν) plane in Fig. 2. In accordance with the previous exposition, we conclude that a filling dependent nuclear-spin polarization develops. The modified electronic Zeeman splitting via the static Overhauser shift [9] triggers the transition at different B_{tr}. From the matching condition $\hbar\omega_c^* = E_Z$, we derive $\alpha\sqrt{B_{tr}} = \beta(B_{tr} + B_N)$ at the transition, where α relates $\hbar\omega_c^*$ to E_C and β contains the electronic g factor. The ordinate can then be converted to obtain an estimate of B_N as a function of ν_{rest}, if the ratio α/β is determined from an appropriate reference

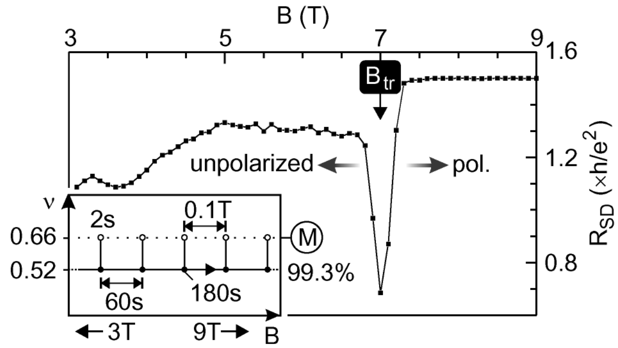

FIG. 2. Example of determining the location of the spin transition for $\nu_{rest} = 0.52$. The inset depicts the path traced in the (B, ν) plane in the course of collecting the R_{SD} values. The corresponding time sequence is plotted in Fig. 1. For the entire experiment, 99.3% of the time is spend at ν_{rest}.

point. However, to first order α/β is irrelevant, and the change in B_N when comparing two different ν_{rest} approximately equals half the shift on the B_{tr} ordinate. A ΔB_N of more than 1.5 T is extracted for the extreme cases $\nu_{rest} = 0.5$ and 0.92 (largely insensitive to the precise value of α/β). For GaAs, it implies a variation in the nuclear-spin polarization $\langle I \rangle$ of order 0.2 and more [12].

Its origin is puzzling. Most of all, what is the impetus for such a large change in $\langle I \rangle$ upon switching to a different ν_{rest}? What determines the final nuclear-spin orientation for a given ν_{rest} and how is the drastic redistribution among the nuclear levels achieved? Large changes in the electronic spin configuration and polarization $\langle S \rangle$ are frequently met when varying the filling in a 2DES. It is natural to surmise that the electronic system adapts instantaneously on the time scale of the experiment, because we are dealing with an open system with a source and sink for electronic spin and abundant relaxation channels for electronic spin exist [13]. This change in $\langle S \rangle$ upon altering the filling comes along with a different Knight shift [9]. Although this evokes a nonequilibrium with respect to the thermal nuclear-spin population, even for the maximum possible Knight shift [14] this mechanism falls 3 orders of magnitude short of accounting for the observed size of ΔB_N.

Because nuclei immensely outnumber electrons and possess gyromagnetic ratios 3 orders of magnitude smaller, large ΔB_N were reported only when saturating electronic transitions with suitable electromagnetic excitation such as in optical orientation or under electron spin resonance conditions [9,15,16]. The sustained nonequilibrium electron spin distribution dynamically polarizes nuclei through hyperfine interaction mediated flip-flop scattering events as though their magnetic moment were comparable to that of electrons. Such a scheme would settle the magnitude issue, but it is far from obvious that the conditions for this dynamic Overhauser effect are fulfilled here. No electromagnetic radiation to establish a nonequilibrium is incident on the sample. Current may in principle serve a similar role [17]. Note, however, that the externally imposed ac current is turned off (Fig. 1) during relaxation at ν_{rest}. Moreover, the experiment has been repeated for a limited set ν_{rest} with either (i) a 5.36 Hz current ranging from 0.1–10–nA or (ii) a 1 nA ac current on top of a larger dc-current component (up to 20 nA) permanently left on. In most instances, no influence is observed. For these amplitudes, only dc currents modify the location of the transition and only for certain values of ν_{rest} such as those near 2/3 itself [7], which might be ascribed to the intricate current flow through a domain pattern of distinct phases. Nevertheless, the shift is insignificant on the scale of Fig. 1.

Hence, the incentive and mechanism to produce a substantial change in the nuclear polarization upon adjusting ν_{rest} remains unidentified. The detailed shape of the curve in Fig. 1 does suggest in some ways that the final state of the nuclear-spin subsystem for filling ν_{rest} may be dictated by the electronic spin polarization at this filling. The symmetry around $\nu_{rest} = 1$ for filling interval $[0.8, 1.2]$ is conspicuous. Away from exact filling 1, Skyrmion physics makes $\langle S \rangle$ drop to small values [14]. The transition is observed at low B_{tr}. High B_{tr} values are obtained near $\nu_{rest} = 0.5$, where a nearly fully spin-polarized CF metallic state is anticipated [18]. Unfortunately, at this filling $\langle S \rangle$ gradually rises as B is swept. Bearing in mind our hypothesis, the final state of the nuclear-spin subsystem would alter accordingly, and therefore follows from self-consistent considerations. In principle, fillings 1 and 2, for which $\langle S \rangle$ is B independent, well known, and equal to 1 and 0, would make ideal reference points to corroborate our premise that the final degree of nuclear-spin orientation is correlated with the electronic spin polarization at ν_{rest}, if it were not for excessively long nuclear lattice relaxation times governing at these fillings. The nuclear-spin subsystem typically faces the flip-flop scattering bottleneck when attempting to adjust $\langle I \rangle$. In the absence of appropriate electronic excitations, which energetically match the nuclear Zeeman splitting, energy conservation cannot be fulfilled and nuclear-spin lattice relaxation via flip-flop scattering is quenched. It prevents the system from reaching equilibrium or a stationary state on the time scale of the experiment in Fig. 1. We have ascertained that the final state is essentially achieved [19] for the chosen parameters for the larger part of this diagram except for, for instance, fillings (marked by green regions) near the diagram borders, near 4/3 and also exactly at 1 and its close vicinity, where very long relaxation times have been previously reported based on NMR [14]. Accordingly, at these fillings the observed location of the spin transition depends on the previous history and provides no valuable information to construct an explanation. For this reason, firming up our conjecture has proven problematic.

When all nuclear species are fully polarized, the nuclear field in GaAs adds up to as much as 5.3 T [12]. Under thermal equilibrium at $T \leq 20$ mK and B up to 9 T, the

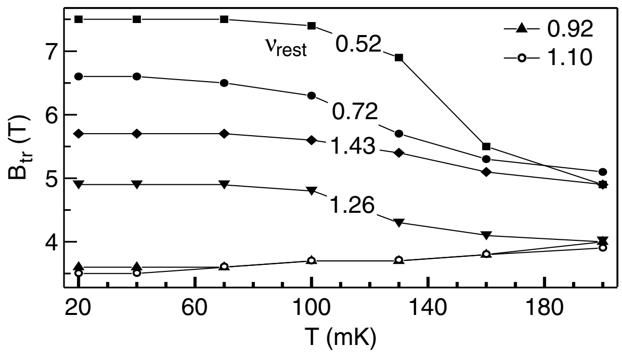

FIG. 3. Temperature dependence of the transition field B_{tr} for selected ν_{rest}. Sample and temperature sensor are in close proximity and both submerged in the He^3/He^4 mixture. The data were taken during a separate cooldown. This accounts for the small discrepancy with the B_{tr} fields plotted in Fig. 1.

thermal nuclear polarization is oddly comparable to the observed maximum value of $|\Delta B_N| \approx 1.5$ T. Does the thermal nuclear-spin polarization put an upper bound on $|\Delta B_N|_{max}$ or is this agreement downright fortuitous? A temperature dependent study in Fig. 3 of B_{tr} for selected values of ν_{rest} discloses the answer. The large difference in B_N is held up to beyond 100 mK for which the thermal nuclear-spin polarization has dropped nearly 1 order of magnitude. It precludes a close connection between the thermal polarization and $|\Delta B_N|_{max}$. Note that the convergence of the B_{tr} curves at higher T may be a measurement artifact. The short time interval at filling 0.66 needed for lock-in acquisition of R_{SD} can no longer be ignored. The nuclear lattice relaxation at $\nu_{R_{SD}} = 0.66$ is accelerated by the temperature increase; the R_{SD}-probe sequence may modify the nuclear polarization and falsify the extracted value of B_{tr} for filling ν_{rest}. Large $|\Delta B_N|$ may therefore persist even up to higher temperatures, although we cannot prove this in our arrangement. The temperature insensitivity of B_{tr} does not contradict the suggestion that B_{tr} tracks the electron spin polarization at ν_{rest}. For instance, NMR studies suggest no or very weak changes in the electron spin polarization near filling 0.5 and 0.8–1.2 for the temperature range of Fig. 3 where the interpretation is unambiguous ($<$ 130 mK) [14,20].

In summary, an anomalous, large filling factor dependence of the spin polarization of nuclei residing in the same plane as the 2D electrons has been revealed. A mechanism to accomplish the implied numerous nuclear-spin flips remains enigmatic. Common schemes of dynamic nuclear polarization based on suitable electromagnetic radiation or an external current where either incident photons or metallic reservoirs supply spin angular momentum seem irrelevant. A role for spin orbit coupling [13] in this problem is difficult to reconcile with the intricate filling dependence. Even if not well understood, the experimental facts can be exploited to imprint a particular value of net nuclear spin at the touch of a gate voltage. This write capability for nuclear-spin polarization should be particularly instrumental for devising RF-free pendants of NMR techniques to investigate electron spin-nuclear spin interactions in a single layer system.

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