Spin Reversal of a Quantum Hall Ferromagnet at a Landau Level Crossing

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When Landau levels (LLs) become degenerate near the Fermi energy in the quantum Hall regime, interaction effects can drastically modify the electronic ground state. We study the quantum Hall ferromagnet formed in a two-dimensional hole gas around the LL filling factor $\nu = 1$ in the vicinity of a LL crossing in the heavy-hole valence band. Cavity spectroscopy in the strong-coupling regime allows us to optically extract the two-dimensional hole gas’ spin polarization. By analyzing this polarization as a function of hole density and magnetic field, we observe a spin flip of the ferromagnet. Furthermore, the depolarization away from $\nu = 1$ accelerates close to the LL crossing. This is indicative of an increase in the size of Skyrmion excitations as the effective Zeeman energy vanishes at the LL crossing.

The quantum Hall state at the Landau level (LL) filling factor $\nu = 1$ constitutes an exciting platform for studies of many-body physics. It has a rich phase diagram depending on the Coulomb, Zeeman and disorder potentials. At low disorder potential, the ground state at $\nu = 1$ is an itinerant quantum Hall ferromagnet (QHF) which can persist even in the absence of a single-particle energy gap due to strong many-body interactions [1–5]. Furthermore, in the absence of Zeeman coupling, a zero-temperature phase transition from a ferromagnet to a quantum Hall Skyrmion glass (QHSG) occurs for a large but finite Coulomb to disorder ratio, in the presence of density fluctuations [6].

Previous investigations showed a high degree of spin polarization in a two-dimensional electron gas (2DEG) at $\nu = 1$ measured by different techniques such as NMR [7, 8] and optical spectroscopy [9, 10]. In these studies, the relevance of Skyrmionic excitations in the vicinity of $\nu = 1$ was discussed in order to explain the fast depolarization of the system away from $\nu = 1$. It was also pointed out that, both in GaAs 2DEGs and two-dimensional hole gases (2DHGs), Skyrmions are the lowest-lying charged excitations in the vicinity of $\nu = 1$ [1, 7, 9, 11–13]. Skyrmionic excitations rely on the interplay between Zeeman energy $E_Z$ and the Coulomb energy $E_C$. When the ratio between these two quantities $\tilde{g} = E_Z/E_C$ is small, large Skyrmions are favoured. Eventually, in an ideal system with vanishing $E_Z$, the Skyrmion size would become macroscopic [1, 14, 15]. Vanishing $E_Z$ was experimentally achieved by hydrostatic pressure on the sample. In this way, the g-factor was reduced to zero leading to the observation of large Skyrmions [14]. A very small g-factor was also achieved in an Al0.13Ga0.87As/GaAs quantum well (QW) structure leading to a Skyrmionic excitation composed of approximately 50 spins [15]. Another possibility to obtain a vanishing $E_Z$ is to exploit LL crossings in a 2DHG, which occur at finite magnetic field due to heavy- and light-hole mixing in the valence band [16–23]. In this scenario, an interesting so far unresolved question, is the robustness of the QHF and the role of Skyrmions at the LL crossing.

In this Letter, we study the $\nu = 1$ QHF in a density-tunable 2DHG sample in the vicinity of a LL crossing. We extract the absolute value of the spin polarization using cavity spectroscopy from which we obtain a phase diagram of the $\nu = 1$ state, clearly showing a spin reversal of the ferromagnetic ground state. A narrow transition region with vanishing average polarization separates the two highly-spin-polarized phases, while the transition region broadens away from $\nu = 1$. By approaching the LL crossing, thereby tuning $E_Z$ to zero, we observe a faster depolarization of the QHF as a function of the filling factor, suggesting a diverging Skyrmion size.

The sample is grown by molecular beam epitaxy and its structure is shown in Fig. 1(a). The 2DHG in a 15-nm-thick, single-side modulation-doped GaAs QW is embedded at the centre of a single-wavelength long Al0.21Ga0.79As optical microcavity and the carbon δ-doping is located 30 nm above the QW. The top and bottom distributed Bragg reflectors (DBRs) are composed of 21 and 25 pairs of AlAs/Al0.21Ga0.79As layers, respectively. The microcavity has a quality factor of $Q \approx 5 \times 10^3$ which was measured by white light reflection. A silicon-doped n-layer is grown below the bottom DBR acting as a gate which allows tuning the hole density $p$ of the 2DHG. At the border of the sample 16 DBRs pairs were etched with selective wet etching in order to contact the 2DHG which was achieved with In0.96Zn0.04. A second etching step was performed in order to etch under the 2DHG and contact the n-type gate with In. The measured mobility, from a sample of the same wafer with a patterned 100-μm-wide Hall bar at 250 mK, is about $1.14 \times 10^6$ cm$^2$/Vs at a density of $1.92 \times 10^{11}$ cm$^{-2}$ obtained by Hall measurement. Applying a negative bias to the diode-like structure, the 2DHG is linearly depleted. The hole mobility decreases linearly with decreasing the hole density, down to $0.33 \times 10^6$ cm$^2$/Vs at a density of $0.57 \times 10^{11}$ cm$^{-2}$. 

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In a dilution refrigerator at about 70 mK, the optical resonances of the 2DHG are probed by white light reflection as a function of the gate voltage in the region where the top mirror was partially removed. At large negative bias (\(-8 \text{ V}\)) the 2DHG is depleted, only the heavy-hole and light-hole exciton modes at about 1528.5 meV and 1536 meV are observed (Fig. 1(b)). The shift to lower energies from \(-8\) to \(-10\) V is attributed to the quantum confined Stark effect [24]. At about \(-7.9 \text{ V}\), a low density 2DHG is formed which interacts with the excitons and splits the resonance into two polaronic modes. These two modes arise from the dressing of the optically excited excitons by excitations of the heavy-hole Fermi sea, thereby forming an attractive and a repulsive exciton-polaron resonance [25–31]. The heavy-hole attractive exciton-polaron, previously referred to as the charged exciton or trion [28], appears at an energy of approximately \(2\) meV lower than the exciton mode. The heavy-hole repulsive exciton-polaron is also well visible from \(-5.5 \text{ V}\) up to \(-8 \text{ V}\). Due to phase-space filling, as well as due to the quantum confined Stark effect, the modes shift to higher energies with increasing gate voltage, or equivalently the hole density. We also observe the light-hole exciton dressed by the heavy-hole Fermi sea forming an attractive polaron at an energy \(\sim 0.5 \text{ meV}\) lower than the corresponding (light-hole) exciton. In contrast to the heavy-hole exciton-polaron, there appears to be a continuous transition from the bare light hole exciton mode distinguishable from the polaron branch.

In the following, we will focus on the region of the sample with the microcavity where the exciton-polarons hybridize with the cavity mode to form polaron-polaritons [25, 26, 31, 32]. In particular, we select the cavity energy to be close to resonance with the heavy-hole attractive exciton-polaron at low densities \(p < 6 \times 10^{10} \text{ cm}^{-2}\). Given the complex behaviour of the heavy-hole LLs, we performed numerical calculations based on the 8 \times 8 Kane model [19]. Many-body interactions were taken into account within the self-consistent Hartree approximation. Due to mixing between the heavy- and light-hole states, the LLs are non-linearly dependent on the magnetic field \(B\) (Fig. 1(c)), resulting in multiple LL crossings. In this work we are interested in studying the QHF in the vicinity of the crossing between the two lowest LLs occurring at a magnetic field of approximately \(B_{\text{cr}} \approx 1.6 \text{ T}\) (see green arrow in Fig. 1(c)). The LL labeled “a” is a pure heavy-hole state (spin \(z\) component \(S_z = +3/2\)) for all values of \(B\), whereas the LL labeled “b” is pure heavy-hole \((S_z = -3/2)\) only in the limit \(B = 0\), but it acquires some light-hole character \((S_z = +1/2)\) at finite \(B\) (about 10% at \(B = 5 \text{ T}\)). Note that the spin \(z\) component is expressed according to the “hole notation” [33]. The effective Zeeman gap \(\tilde{E}_Z\) between the LLs a and b is very small at magnetic fields \(B < B_{\text{cr}}\). For magnetic fields above \(B_{\text{cr}}\), \(\tilde{E}_Z\) quickly increases to hundreds of \(\mu\)eV. In Fig. 1(c), the optical transitions are represented according to the selection rules for dipole transitions. At \(B < B_{\text{cr}}\) the lower energy optical transition is expected to be \(\sigma_+\) while at \(B > B_{\text{cr}}\), it is \(\sigma_-\).

To study the itinerant QHF around \(\nu = 1\), we perform polarization-resolved optical spectroscopy in perpendicular magnetic fields to obtain the spin polarization of the 2DHG. The white light reflection is probed independently for \(\sigma_-\) and \(\sigma_+\) polarized light. In the top panel of Fig. 1(d) the difference between the \(\sigma_-\) (red) and \(\sigma_+\) (blue) reflectivities is shown for a hole density of...


\[ p = 4.9 \times 10^{10} \text{ cm}^{-2} \] as a function of \( \nu \). The polariton normal-mode splitting \( \Omega \) corresponds to half the energy splitting between the upper and lower polariton at zero cavity detuning. For the displayed range of filling factors, a large polariton normal-mode splitting is observed for \( \sigma_+ \) polarization. However, the normal-mode splitting for \( \sigma_- \) vanishes at the \( \nu = 1 \) state suggesting a complete spin polarization of holes in the lowest LL spin subband. Indeed, the normal-mode splitting squared \( \Omega_{\sigma_+}^2 \) reflects the polaron oscillator strength and therefore the available density of states of each spin subband of the first LL. This allows for a quantitative determination of the spin polarization

\[ P = (\Omega_{\sigma_+}^2 - \Omega_{\sigma_-}^2)/(\Omega_{\sigma_+}^2 + \Omega_{\sigma_-}^2) \quad (1) \]

of the 2DHG [31, 34].

For each filling factor, the white light reflection spectra for both circular polarizations were fitted with Lorentzian lineshapes. This allows extracting the normal-mode splitting and therefore the spin polarization according to Eq. 1, as shown in the bottom panel of Fig. 1(d). Almost complete spin polarization (98 \%) was measured at \( \nu = 1 \) consistent with the fact that the system is in a robust ferromagnetic state. Depolarization of the system is observed on both sides of \( \nu = 1 \). We present experimental evidence for the spin flip of the ferromagnetic ground state caused by the LL crossing in Fig. 2, where white light reflection measurements as a function of the magnetic field are presented for both circular polarizations. At a density of \( 2.9 \times 10^{10} \text{ cm}^{-2} \) (left column), the \( \nu = 1 \) state is at about 1.20 T, where the polariton normal-mode splitting vanishes for \( \sigma_+ \) optical excitation. This is expected for a ferromagnetic state, where all the spins occupy the lower spin-subband of the first LL \( |\uparrow\rangle \) and therefore reduces the oscillator strength of the \( \sigma_+ \) exciton resonance. In contrast, the normal-mode splitting is enhanced for \( \sigma_- \) due to an increase of oscillator strength, due to availability of all holes for dynamic screening, as well as due to complete absence of phase-space filling for \( \sigma_- \) exciton formation [31]. At the density of \( 3.6 \times 10^{10} \text{ cm}^{-2} \), a remarkably sharp transition is observed at a magnetic field of approximately 1.45 T. At this density and magnetic field, the \( \nu = 1 \) state is exactly at the LL crossing and the polariton modes experience a sharp jump in \( \Omega \) as a function of \( B \). At a density of \( 4.2 \times 10^{10} \text{ cm}^{-2} \) the normal-mode splitting vanishes for \( \sigma_- \) polarized light indicating that the ferromagnet is in the \( |\downarrow\rangle \) state. The magnetic field and density at which the LL crossing occurs are in good agreement with the numerical simulations (Fig. 1(c)). Moreover, the numerical simulations show that for \( B < B_{cr} \), the LL gap between the first and second levels is much smaller than for \( B > B_{cr} \). This explains the small difference between \( \sigma_- \) and \( \sigma_+ \) optical excitations at \( B < B_{cr} \) in Fig. 2 (bottom row) where a clear spin splitting is observed only at \( \nu = 1 \).

In Fig. 3(a), additional evidence for the LL crossing is presented in a phase diagram of the 2DHG spin polarization as a function of density and filling factor. The LL crossing appears as a transition region with low polarization (shown in white) between two highly-spin-polarized regions. The upper (red) region is the one where the itinerant holes are mainly in the \( |\downarrow\rangle \) state, while in the bottom (blue) region, they are mainly in the \( |\uparrow\rangle \) state. We show, in Figs. 3(b-d), three line cuts at constant filling
factors of the absolute value $|P|$ to illustrate the evolution of the spin polarization away from $\nu = 1$. At exactly $\nu = 1$ and $p > p_{cr} \approx 3.6 \times 10^{10}$ cm$^{-2}$, the spin polarization reaches a maximum value of 98% as expected for a robust QHF state at this integer quantum Hall state [7, 8, 35]. After a sharp transition at the LL crossing, occurring at a density of approximately $p_{cr} \approx 3.6 \times 10^{10}$ cm$^{-2}$, the system repolarizes in the opposite spin state ($|\uparrow \rangle$), up to an absolute value of spin polarization of 80%. Such a low polarization at the $\nu = 1$ state is expected for samples with a low Coulomb energy to disorder ratio [2, 6]. In this scenario, we would expect a gradual decrease of the spin polarization while decreasing the 2DHG density. In contrast, we observe a step-like feature with two plateaus at $P \approx 80\%$ for $p < p_{cr}$ and at 98% for $p > p_{cr}$. Therefore, a more likely possibility is a reduction of $\tilde{E}_Z$: indeed, at $p < p_{cr}$, a substantially smaller $\tilde{E}_Z$ is expected from the numerical simulations (Fig. 1(c)) and the magnetoreflection measurements (lowest panels of Fig. 2). As a consequence of the small $\tilde{E}_Z$ and the presence of disorder and density fluctuations, charges are introduced into the system in the form of Skyrmions and anti-Skyrmions. In this situation, a QHSG is formed at a sufficient Skyrmion density, at which they begin to overlap [6, 36]. For $0 < |\nu - 1| < 0.15$, Skyrmionic excitations reduce the many-body energy gap, leading to a further lowering of the spin polarization [1, 7–10]. The line cuts in Figs. 3(b, d) at $\nu = 0.92$ and 1.08 show a broadening of the transition region (white area), which indicates a lower energy gap compared to the line cut at $\nu = 1$, as expected. The asymmetry of the polarization at $\nu > 1$ and $\nu < 1$ is a consequence of the asymmetry of $\tilde{E}_Z$ around the LL crossing.

In order to evaluate the importance of Skyrmions near the crossing, we studied the depolarization of the system when the filling factor is tuned away from $\nu = 1$, i.e. when a magnetic flux quantum is added or removed from the system. The depolarization at $\nu = 1$ can be described by the Skyrmion model proposed in [7]:

$$P = \begin{cases} P_S^1 \left[ S \frac{2}{\nu} - (S - 1) \right], & \nu > 1 \\ P_A^1 \left[ \frac{1}{\nu} - (2A - 1) \frac{1-\nu}{\nu^2} \right], & \nu < 1 \end{cases}$$

(2)

where $S$ and $A$ describe the sizes of the Skyrmions and anti-Skyrmions in terms of number of holes whose spin flips upon removing or adding a flux quantum, respectively. $P_S^1$ or $P_A^1$ are two fitting parameters representing the amplitude of the spin polarization in the limits $\nu \to 1^+$ and $\nu \to 1^-$, respectively. This allows for a better fitting of the depolarization in the vicinity of $\nu = 1$ even when the system is not fully polarized. In a single-particle picture, the parameters $A$ and $S$ are simply equal to 1 [7]. However it has been shown in previous studies that in GaAs 2DEGs and 2DHGs, $S$ and $A$ are larger than one in the vicinity of $\nu = 1$ [7–10, 12, 13]. The values of $S$ and $A$ depend on the ratio $\tilde{\eta}$ of $\tilde{E}_Z$ to the Coulomb energy. $\tilde{E}_Z$ promotes single-spin excitations reducing the sizes of Skyrmions and anti-Skyrmions and therefore the $S$ and $A$ parameters. In contrast, a large Coulomb energy promotes large Skyrmions and anti-Skyrmions.

Figure 4(a) shows the magnitude of the parameters $A$ and $S$ obtained by fitting the data (shown in Fig. 3(a)) using Eq. 2 as a function of the 2DHGs density. As the density gets closer to $p_{cr}$, the parameters $A$ and $S$ diverge. According to particle-hole symmetry, $A$ and $S$ should be equal at any density [1, 9]. However, this is not observed in the vicinity of $p_{cr}$. At $p > p_{cr}$, $A$ is only weakly affected by the crossing compared to $S$ as can also be seen in the line cuts at $p = 3.69, 3.96$, and $4.32 \times 10^{10}$ cm$^{-2}$ in Figs. 4(b–d). For $p < p_{cr}$, the situation is the opposite, with $A$ being more influenced by the LL crossing compared to $S$. In Fig. 4(b) the crossing is clearly visible and occurs at approximately $\nu = 1.07$. The polarization quickly drops from about 0.75 down to 0 in a very short range of filling factor. According to the Skyrmion model, the Skyrmion size is $S = 11.7$, while the anti-Skyrmion size $A = 2.8$ remains relatively small. We believe that the difference between the $A$ and $S$ is a direct consequence of the non-linear dependence of $\tilde{E}_Z$ on $B$. Indeed, a small deviation away from $B_{cr}$, especially in the case where $B > B_{cr}$, leads to a large increase of $\tilde{E}_Z$ as can be seen in the simulations of Fig. 1(c). Furthermore, this leads to a reduction of the Skyrmion and anti-Skyrmion size.

To conclude, we studied the spin polarization of the $\nu = 1$ state in the vicinity of a LL crossing occurring in 2DHGs by means of cavity spectroscopy in the strong coupling regime. The LL crossing leads to a rich ferromagnetic phase diagram with a peculiar transition of the ground state from the $|\uparrow \rangle$ to the $|\downarrow \rangle$ state which is re-

![FIG. 4](image-url)
markably sharp at \( \nu = 1 \) and broadens upon tuning away from it. Moreover, this transition exhibits a vanishing effective Zeeman energy, which promotes large Skyrmions and therefore accelerates the depolarization of the ferromagnet upon tuning the filling factor away from \( \nu = 1 \). Further studies could focus on fractional quantum Hall states in the proximity of LL crossings, enabled in 2DHGs by the richness of the valence band structure. We envision that this exciting platform will lead to new insights on quantum Hall ferromagnetism and Skyrmion physics in the limit of a vanishing Zeeman energy.

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