

7/3 fractional quantum Hall state in GaAs double quantum wells

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We observed the fractional quantum Hall effect at a total filling factor of $7/3$ in the Hall resistance, accompanied by a single minimum in longitudinal magnetoresistance in uncoupled, bilayer electron gases formed in GaAs double quantum well systems. When the density of the top electron gas is varied, the total $7/3$ ($1 + 4/3$) state can be tuned such that while the Hall plateau value essentially remains unchanged, the associated longitudinal magnetoresistance develops into a double minimum. When Zeeman energy is introduced by tilting the applied magnetic field, the two minima respond differently; one remains essentially unchanged, while the other deepens, consistent with the differing spin dependences of the integer and fractional gaps. Our results demonstrate that the detailed evolution of the $7/3$ plateau probes layer-resolved incompressibility, interlayer charge redistribution, and spin-dependent gap behavior in a nontunneling bilayer system.

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I. INTRODUCTION

Two-dimensional electron gases (2DEGs) in high magnetic fields provide a controlled setting in which electron-electron interactions dominate over kinetic energy and give rise to a variety of incompressible quantum Hall states. In the lowest Landau level (LLL), the hierarchy of Haldane pseudopotentials stabilizes robust integer and fractional plateaus at filling factors $\nu = 1, 1/3, 2/3, 4/3$, etc. [1–4], whose gap structures and spin configurations have been extensively characterized through transport, thermodynamic compressibility, and capacitance measurements [5,6]. These studies establish the fundamental energy scales and density dependence of LLL incompressibility in high-mobility GaAs heterostructures.

Bilayer electron systems (BESs) provide an expanded platform for exploring interaction effects, as the layer degree of freedom introduces an additional axis of tunability. A BES is characterized by the fact that if the barrier separating the two 2DEGs is narrow enough to allow electron tunneling between them, the exotic fractional quantum Hall effect (FQHE) may be realized collectively by electrons from both layers [5,7–15]. Additionally, BESs host interlayer coherent phases, including the excitonic condensate at $\nu = 1$, characterized by Josephson-like interlayer transport and quantized Hall drag [11,16–18]. In contrast, when the barrier is sufficiently wide, interlayer tunneling is strongly suppressed, and the two layers form independent sets of Landau levels [19]. In this limit, the layers are often described as electronically decoupled, with

the total filling factor of the BES given by the sum of the individual layer fillings [19,20]. However, even in the absence of tunneling, the layers remain electrostatically coupled and share a common electrochemical potential through the Ohmic contacts. As a result, their densities are not strictly independent: When one layer becomes incompressible, its chemical potential cannot adjust continuously, and charge redistribution between the layers can occur to maintain equilibrium. Thus, gating one layer can indirectly influence the other, particularly in the quantum Hall regime where compressibility plays a central role [7,19–21].

Experiments utilizing capacitance spectroscopy, local compressibility, and transport have demonstrated that such systems exhibit negative electronic compressibility [22,23], oscillatory chemical potentials associated with integer and fractional quantum Hall transitions [24–26], and magnetic-field-driven interlayer charge redistribution [27–30]. More recent studies further highlighted that transport depends sensitively on layer compressibility, coupling to reservoirs, and the requirement of a common electrochemical potential [19–21,31].

In this work, we study two spatially separated 2DEGs confined in a GaAs double quantum well. The large interwell barrier suppresses interlayer tunneling, so that each layer retains its own Landau level (LL) spectrum. At the same time, the two layers remain electrostatically coupled and are constrained to share a common electrochemical potential via the Ohmic contacts. A top gate primarily tunes the carrier density of the upper 2DEG; however, when one of the layers enters an incompressible quantum Hall state, this constraint leads to a redistribution of charge between the layers, effectively modifying both densities.

By electrostatically tuning the upper layer while nominally keeping the lower layer fixed, we access a regime in which a quantized Hall plateau at a total filling factor of $7/3$ emerges from the coexistence of integer and fractional states at fillings 1 and $4/3$ in the two layers, respectively. We investigate the

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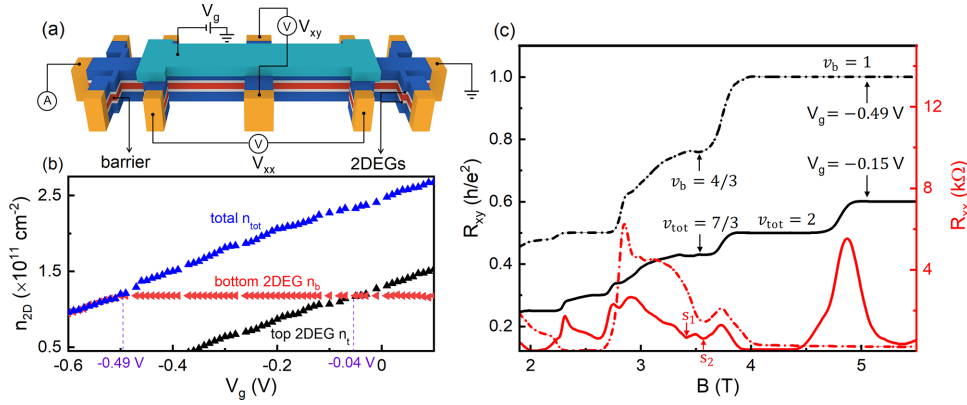


FIG. 1. (a) Layout of the Hall bar and the measurement setup. The Hall bar, Ohmic contacts, and top gate are shown in dark blue, yellow, and light blue, respectively. A red barrier separates the top and bottom 2DEGs (shown in white). (b) The gate voltage dependences of electron densities for the top (n_t , black) and bottom (n_b , red) layers, as well as the total for the bilayer system (n_{tot} , blue) for sample 1. A balanced bilayer condition where $n_t = n_b$ is achieved at $V_g \approx -0.04 \text{ V}$, and a complete depletion of the top layer is achieved at $V_g \approx -0.49 \text{ V}$. (c) Transverse and longitudinal magnetoresistances, R_{xy} (black) and R_{xx} (red), as a function of B for $V_g \approx -0.15 \text{ V}$ (solid traces) and -0.49 V (dash-dotted traces). The $V_g \approx -0.49 \text{ V}$ trace corresponds to the effective single-layer regime where the top layer has been depleted, and only the bottom layer contribution is measured. At $V_g \approx -0.15 \text{ V}$, the $\nu_{\text{tot}} = 7/3$ plateau in R_{xy} is accompanied by two R_{xx} minima, labeled s_1 and s_2 ; for comparison, the bottom layer $\nu_b = 4/3$ state at $V_g \approx -0.49 \text{ V}$ is indexed.

evolution of this composite $7/3$ state under controlled density imbalance and enhanced Zeeman coupling using tilted magnetic fields, allowing us to probe how layer-resolved incompressibility, interlayer electrostatics, and spin energetics collectively shape the resulting quantum Hall structure.

II. EXPERIMENT

Experimental samples were converted into Hall bars with a top gate using molecular beam epitaxy (MBE) grown GaAs/Al_{0.33}Ga_{0.67}As heterostructures containing a double quantum well (each well was 15 nm thick), separated by a 30 nm thick Al_{0.33}Ga_{0.67}As barrier. The Hall bar had a length of 1.8 mm and a width of 80 μm and was covered by an 860 μm long Ti/Au top gate. Several Ohmic contacts were formed at the edges of the Hall bar by evaporating and then rapidly annealing an alloy of Ni, Ge, and Au to make electrical contacts via both 2DEGs. The top quantum well (containing the top 2DEG) was located approximately 275 nm beneath the surface of the heterostructure [Fig. 1(a)]. The magnetoresistance measurements were performed based on the four-terminal lock-in technique using an ac excitation current of 10 nA at 77 Hz at a base temperature $T \approx 20 \text{ mK}$ in a cryofree dilution refrigerator (DR). The electron density (mobility) of both 2DEGs is $\sim 1.2 \times 10^{11} \text{ cm}^{-2}$ ($1.2 \times 10^6 \text{ cm}^2/\text{Vs}$). The magnetic field B was ramped at a rate of 6 T/h; all measurements were recorded during the downward sweep of the magnetic field [19]. Experiments were performed on two samples (samples 1 and 2) across multiple cooldown processes in a DR to ensure the reproducibility of reported experimental results.

III. RESULTS

A. Double minimum for $\nu_{\text{tot}} = 7/3$

The experimental sample (sample 1) was studied by measuring transverse (R_{xy}) and longitudinal (R_{xx}) magnetoresistance curves as a function of magnetic field for

a range of top gate voltages (V_g). With low magnetic field Shubnikov-de Haas oscillations in the R_{xx} data, the variation of 2D electron density n_{2D} as a function of V_g is shown in Fig. 1(b). While V_g is varied, only the density of the top 2DEG n_t (and the total density n_{tot}) is modulated; the density of the bottom 2DEG n_b remains nearly constant until the top 2DEG is depleted at $V_g \approx -0.49 \text{ V}$ [19]. When dn_t/dV_g is used in the capacitance formula, the distance between the top gate and the top 2DEG is estimated to be $\sim 278 \text{ nm}$, consistent with the MBE-grown layer thickness.

In uncoupled 2DEGs, the total filling factor is given by $\nu_{\text{tot}} = \nu_t + \nu_b$, where ν_t and ν_b are the filling factors for the top and bottom 2DEGs, respectively. Figure 1(c) shows selected magnetoresistance curves for two gate voltages: -0.15 and -0.49 V . The R_{xy} magnetoresistance curve (solid black line) for $V_g \approx -0.15 \text{ V}$ shows a Hall plateau at $0.429 (h/e^2)$, representing $\nu_{\text{tot}} = 7/3$, due to contributions from both 2DEGs. The deviation between the measured and expected R_{xy} values is less than 0.1%. The corresponding R_{xx} curve (solid red line) shows two sequential R_{xx} minima, denoted state 1 (s_1) and state 2 (s_2). A minor step at the center of the $7/3$ plateau aligns with the point where the double minimum merges, indicating a possible transition between s_1 and s_2 . To estimate contributions from each 2DEG, magnetoresistance curves for $V_g \approx -0.49 \text{ V}$, representing contributions from only the bottom 2DEG, are shown for reference. For instance, for $V_g \approx -0.15 \text{ V}$, at $B \approx 3.50 \text{ T}$, $\nu_{\text{tot}} = 7/3$, the corresponding $\nu_b = 4/3$, and $\nu_t = 1$.

A detailed magnetoresistance investigation was carried out over a range of top gate voltages from $\sim +0.15$ to -0.50 V ; the results are shown in the form of a Landau fan diagram of $R_{xx}(B, V_g)$ as a function of magnetic field and the applied top gate voltage in Fig. 2(a). The regions in black (red) represent minima (risers) in R_{xx} and plateaus (transitions between plateaus) in R_{xy} . As stated before, the bottom 2DEG density remains unchanged until $V_g \lesssim -0.49 \text{ V}$, which suggests the LL contributions from the bottom 2DEG will remain nearly

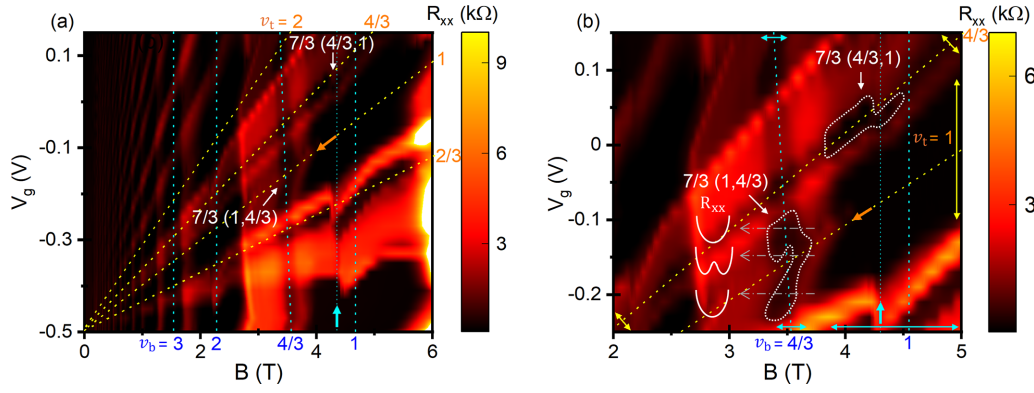


FIG. 2. (a) Fan diagram of R_{xx} as a function of B and V_g for sample 1, where dark (red) regions indicate minima (risers) in R_{xx} . Yellow and cyan dashed guide lines indicate the positions of filling factors of the top, ν_t , and bottom, ν_b , layers, respectively. White arrows mark the regions where $\nu_{tot} = 7/3$ is observed. The filling factors for the top and bottom layers are indicated at the top and bottom of the plot, respectively, in orange and blue. A vertical cyan arrow near the dashed cyan line shows $\nu_b = 1$, and the orange arrow (aligned along the yellow dashed line which represents $\nu_t = 1$) points to discontinuities between compressible and incompressible states. (b) Enlarged view of a section of (a), highlighting the gate voltage regions corresponding to $\nu_{tot} = 7/3$: (1) For negative V_g , $\nu_{tot} = 7/3 = 1(\nu_t) + 4/3(\nu_b)$; for positive V_g , $\nu_{tot} = 7/3 = 4/3(\nu_t) + 1(\nu_b)$. These regions where the $\nu_{tot} = 7/3$ state exhibits a double-minimum structure in R_{xx} are enclosed by two white-dotted contours ($R_1^{7/3}$, the lower one, and $R_2^{7/3}$, the upper one). Three horizontal gray guide arrows piercing through the lower contour ($R_1^{7/3}$) indicate representative V_g values at which $\nu_{tot} = 7/3$ displays single-, double-, and single-minimum behavior. The incompressible states 1 and $4/3$ corresponding to the bottom (top) layer are shown with cyan double-headed arrows (yellow double-headed arrows).

unchanged for the range of V_g under consideration. This is evident from the vertical dashed cyan lines, which correspond to LL fillings ν_b from the bottom layer. The top 2DEG's density changes linearly with gate voltage [Fig. 1(b)]; its LL fillings are shown as dashed yellow lines. The intersections of the yellow and cyan lines indicate that both layers are in incompressible states, and their individual filling factors add up to form the total filling of the bilayer system. For example, $\nu_{tot} = 7/3$ (shown by white arrows) appears for two ranges of gate voltage. We will focus on the evolution of $\nu_{tot} = 7/3$ in this work.

Figure 2(b) shows a magnified view of the region in Fig. 2(a), where $\nu_{tot} = 7/3$ is evolving. For $V_g \approx -0.15$ V, $\nu_{tot} = 7/3 = 1(\nu_t) + 4/3(\nu_b)$; for $V_g \approx 0.05$ V, $\nu_{tot} = 7/3 = 4/3(\nu_t) + 1(\nu_b)$. The total $7/3$ for negative V_g appears for a narrow range of $B \approx 3.5 \pm 0.2$ T, over which $\nu_b = 4/3$ (small cyan double-headed arrow) and $\nu_t = 1$ (large yellow double-headed arrow); the V_g range is ~ -0.1 to -0.23 V. It may be noted that there is a deformation in the compressible state (red regions) in the center of the total $7/3$ region, shown by the white dotted contour (we call it $R_1^{7/3}$), which is located at $B \approx 3.5$ T and $V_g \approx -0.15$ V; we call it the $V_g^- 7/3$ center. The $V_g^- 7/3$ center location, interestingly, appears to align with the center of the incompressible state, $\nu_t = 1$, as shown by the orange arrow aligned along the $\nu_t = 1$ dashed yellow line. The appearance of the R_{xx} minimum along different vertical sections of $R_1^{7/3}$ is shown in a cartoon, with three horizontal arrows (top, middle, and bottom) indicating their locations. In the top region, R_{xx} exhibits a single minimum, followed by a double minimum in the middle and a single minimum at the bottom section. This is a reflection of the experimentally measured response, which will be discussed later. Similarly, for the positive top gate voltage, the total $7/3$ appears for a range of V_g from ~ 0.01 to 0.06 V, over which R_{xx} takes a shape with a single minimum, double minima,

and a single minimum; we call this region $R_2^{7/3}$ and denote its center where the double minimum occurs at $V_g \approx 0.05$ V and $B \approx 4.30$ T as $V_g^+ 7/3$. The $R_1^{7/3}$ and $R_2^{7/3}$ regions, particularly their centers, are reminiscent of the transition between compressibilities. The V_g^+ central $7/3$ region ($R_2^{7/3}$) aligns with a discontinuity observed at the compressible-incompressible junction of $\nu_{tot} = 5/3$ [$2/3(\nu_t) + 1(\nu_b)$], which is clearly visible in Fig. 2(a) near $B \approx 4.30$ T. A similar discontinuity may be traced for $R_1^{7/3}$ at the junction of $\nu_b = 2$ and $\nu_t = 1$ near $B \approx 2.20$ T. The discontinuity at the transition between the compressible and incompressible states is a hallmark of uncoupled bilayer systems [7,9,19,20].

Essentially, based on Figs. 1 and 2, we see signatures of $\nu_{tot} = 7/3$ for two ranges of the top gate voltage, and the occupancy of the layers is reversed with gate voltage tuning. In addition, each $7/3$ region, as a function of small variations in gate voltage, exhibits switching between single- and double-minimum behavior in R_{xx} while maintaining the total Hall plateau at the $7/3$ filling.

B. Single and double minimum in R_{xx}

Figure 3 shows the magnetoresistance plots of R_{xy} and R_{xx} , measured as a function of magnetic field at selected V_g values, emphasizing the regions where $\nu_{tot} = 7/3$ was observed. Figure 3(a) shows the FQHE at $\nu_{tot} = 7/3$ accompanied by a double minimum in R_{xx} measured at $V_g \approx 0.05$ V. The magnetic field at the center of the $7/3$ Hall plateau $B_{V_g^+}^{7/3}$ is ~ 4.29 T. With a reduction of the top gate to 0.01 V, the plateau width corresponding to $\nu_{tot} = 7/3$ decreases alongside the development of a single minimum in R_{xx} , and $B_{V_g^+}^{7/3}$ decreases to ~ 3.95 T. Figure 3(c) shows $\nu_{tot} = 7/3$ exhibiting a double minimum at $V_g \approx -0.15$ V. With a reduction of the voltage to -0.17 V [Fig. 3(d)], the $7/3$ plateau width decreases like in

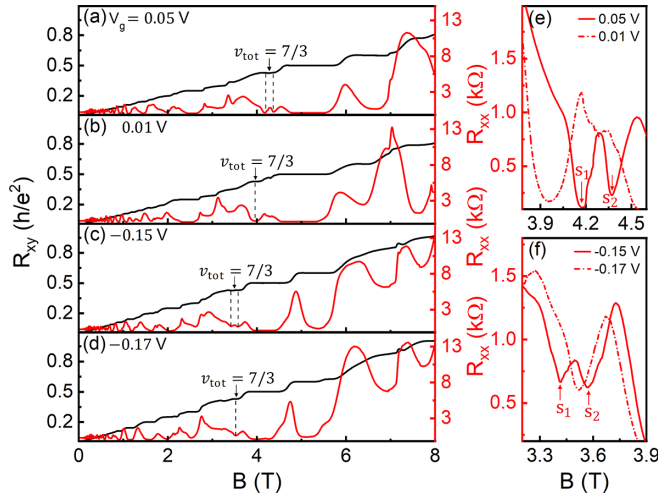


FIG. 3. R_{xy} and R_{xx} as a function of magnetic field for (a) $V_g \approx 0.05$ V, (b) 0.01 V, (c) -0.15 V, (d) and -0.17 V in sample 1. The $\nu_{\text{tot}} = 7/3$ state exhibits a double (single) R_{xx} minimum for the $(\nu_t, \nu_b) = (4/3, 1)$ configuration in (a) and (b) [for the $(\nu_t, \nu_b) = (1, 4/3)$ configuration in (c) and (d)]. (e) and (f) Enlarged views of the $\nu_{\text{tot}} = 7/3$ features in R_{xx} for (a), (b) and (c), (d), respectively, highlighting the two minima as s_1 and s_2 .

the case for the positive top gate voltage and is accompanied by a single minimum. Interestingly, $B_{V_g^-}^{7/3}$ remains almost the same, ~ 3.60 T, for these two gate voltages, unlike the case with positive gate voltages, where $B_{V_g^+}^{7/3}$ changes with V_g . The effect of gate voltages on the R_{xx} minimum is summarized in two magnified plots in Figs. 3(e) and 3(f) for positive and negative gate voltages, respectively. Here the two minima appearing in R_{xx} are labeled s_1 and s_2 . It should also be noted from Figs. 3(a)–3(d) that the width of the total $7/3$ plateau corresponding to the double minimum is almost twice that of the single-minimum case. Similar observations were reported for the $2/3$ FQHE as a function of increasing in-plane magnetic field [32], which suggests the variation in carrier density and the resultant effects are a reflection of a possible spin reconfiguration with the total $7/3$ FQHE.

As discussed, the total $7/3$ state in our bilayer system is composed of 1 and $4/3$ states. For $V_g > 0$ V, $\nu_t = 4/3$, and $\nu_b = 1$ [Figs. 3(a) and 3(b)]. For $V_g < 0$ V, $\nu_t = 1$, and $\nu_b = 4/3$ [Figs. 3(c) and 3(d)]. The application of negative gate voltage in the range of ~ -0.12 to -0.17 V will affect only the top 2DEG ($\nu_t = 1$); the bottom 2DEG ($\nu_b = 4/3$) should remain unaffected. Therefore, a change in carrier density in the top 2DEG is expected to be adjusted by the populated $\nu_t = 1$ state. We know that $\nu_t = 1$ is a highly degenerate, incompressible QH liquid which resists change. It has been established in bilayer electron systems that to maintain equilibrium in the chemical potential, a net transfer of charge carriers should occur between the layers, resulting in hysteresis and spikes in the magnetoresistance and magnetocapacitance [7, 19, 20]. Therefore, to maintain the global electrochemical potential, the change in carrier density migrates to the bottom 2DEG, resulting in a many-body-induced compressible-incompressible transition at the $4/3$ filling. A change in carrier density $\delta n_{2D} \approx 2.0 \times 10^9 \text{ cm}^{-2}$ estimated from B values corresponding to s_1

and s_2 suggests a change in the nature of compressibility leads to a transfer of charges between the layers. Additionally, note that the 2DEGs are 45 nm apart (the center-to-center distance between the two quantum wells), so tunneling or hybridization between them is highly unlikely.

Let us consider the different energy scales, Coulomb (E_C), cyclotron (E_{ω_c}), and Zeeman (E_Z):

$$E_C = e^2/4\pi\epsilon_0\epsilon\ell_B, \quad \ell_B = \sqrt{\hbar/eB},$$

$$E_{\omega_c} = \hbar\omega_c/m^*, \quad E_Z = |g^*|\mu_B B, \quad (1)$$

where ℓ_B is the magnetic length, ϵ is the dielectric constant of GaAs (~ 12.9), $\hbar$ is the reduced Planck constant, $|g^*| = 0.44$ is the effective Landé g factor, $m^* = 0.067m_e$, m^* (m_e) is the effective (rest) mass of the electron, and μ_B is the Bohr magneton. For $B \approx 3.50$ – 4.40 T, Eq. (1) yields $\ell_B \approx 12$ – 14 nm and $E_C \approx 8.14$ – 9.13 meV $> E_{\omega_c} \approx 6.05$ – 7.60 meV $\gg E_Z \approx 0.09$ – 0.11 meV.

Moreover, the observation of the double minimum when the occupancy of both 2DEGs is swapped, as discussed before, additionally confirms the role of nonequilibrium phenomena or compressibility changes in our bilayer systems [19, 20]. Our intuition is corroborated by the observations made in Figs. 2 and 3, where, for variations in positive V_g , the change in B corresponding to the double-minimum region (the center of $R_{xx}^{7/3}$) is ~ 0.30 T; in contrast, for negative V_g , it is ~ 0.05 T (the center of $R_{xx}^{1/3}$), which implies no appreciable change in carrier density in the upper layer. Our results suggest the layer compressibility and incompressibility, as well as the associated electrostatic potentials, play a role.

C. Temperature dependence

Figures 4(a)–4(c) show the temperature dependence of the $\nu_{\text{tot}} = 7/3$ plateau and the corresponding double minimum, s_1 and s_2 , for $V_g \approx 0.07$ V in sample 2. Figures 4(a) and 4(c) show the Hall plateau at $\nu_{\text{tot}} = 7/3$ and the corresponding R_{xx} , respectively, for several temperatures ranging from approximately 30 to 700 mK. We note that the $7/3$ plateau width decreases with temperature, and the point at which s_1 and s_2 combine shifts to higher B . Additionally, s_2 smears out more quickly than s_1 , indicating different mobility gaps for these states. This is shown in the color plot in Fig. 4(b). Due to the limited number of data points in the low-temperature regime (< 100 mK), we could not accurately estimate the activation gaps corresponding to s_1 and s_2 . However, the results qualitatively show that both states responsible for the total $7/3$ Hall plateau evolve asymmetrically with increasing temperature.

The temperature dependence of the total $7/3$ was also measured in sample 1 at $V_g \approx 0.03$ V at a different cooldown; the results are shown in Fig. 4(d). At $T \approx 22$ mK, a broad s_1 , accompanied by a weakly defined s_2 , is shown in R_{xx} . As the temperature increases, s_1 remains fairly stable until around 100–200 mK and subsequently smears out due to thermal broadening. In contrast, s_2 displays strengthening in that the R_{xx} valley deepens until 200 mK and then weakens and gradually disappears at 400–500 mK. The contrasting temperature dependence of s_1 and s_2 points to a possible role of disorder in shaping the observed behavior, particularly

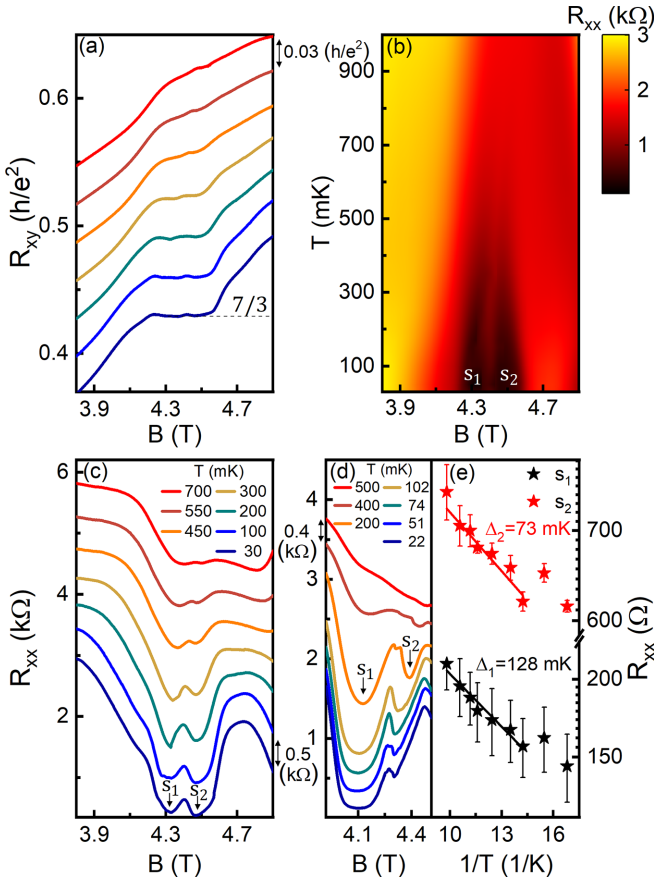


FIG. 4. Temperature dependence of the $\nu_{\text{tot}} = 7/3$ state exhibiting a double-minimum structure in R_{xx} . (a) R_{xy} and (c) R_{xx} versus magnetic field measured at different temperatures at $V_g \approx 0.07$ V in sample 2, with traces vertically offset (offset values are indicated in both plots) for clarity. (b) Color scale plot of R_{xx} results in (c) as a function of B and temperature T (dark: minima), highlighting the two minima, s_1 and s_2 . Temperature-dependent measurements performed on sample 1 at $V_g \approx 0.03$ V: (d) R_{xx} versus B at selected temperatures (vertically offset) and (e) Arrhenius plots of the R_{xx} minima used to estimate activation energies for s_1 and s_2 using data in (d). (a)–(c) and (d), (e) show, respectively, nearly symmetric and asymmetric double-minimum structures.

during the strengthening of s_2 . Figure 4(e) shows the Arrhenius plots using $R_{xx} \propto \exp(-\Delta/2k_B T)$, where k_B is the Boltzmann constant, and the corresponding linear fits for s_1 and s_2 . The estimated activation gaps Δ for s_1 and s_2 are ~ 128 and 73 mK, respectively. The gap of s_2 is consistent with a fractional state at $\nu = 4/3$, while the gap at s_1 ($\nu = 1$) is associated with spin splitting within the LLL, i.e., the Zeeman energy enhanced by exchange interactions.

Note that the dissipated power for a 10 nA ac excitation current lies in the femtowatt range, suggesting that electron heating is negligible and unlikely to affect the extracted activation gaps. Additionally, at sufficiently low temperatures, transport in quantum Hall systems is known to deviate from simple activated behavior due to disorder-induced localization and the onset of variable-range hopping conduction [33–36]. As a result, the extracted activation gap corresponds to a mobility gap determined from the linear regime of the Arrhenius

plot. We therefore restrict our fits to the temperature range where the data exhibit activated behavior.

Early transport measurements in GaAs reported activated behavior of the $\nu = 4/3$ FQH state with an energy gap on the order of 300 mK [37]. Subsequent experiments in high-mobility GaAs heterostructures consistently observed activation gaps in the few hundred millikelvin range, with the precise value depending sensitively on disorder, finite thickness, and Landau level mixing, among other factors [36,38,39]. In tilted magnetic fields, it was shown that the $\nu = 4/3$ activation gap is strongly sensitive to the Zeeman energy and undergoes a pronounced reduction from ~ 300 mK in the spin-unpolarized regime to ~ 100 mK in the spin-polarized regime, indicating that the lowest-energy charged excitations depend critically on the spin degree of freedom [35]. In contrast, the activation energy of the $\nu = 1$ integer QHE state is not determined solely by the cyclotron gap and reflects spin-related excitations and exchange-enhanced spin splitting, which increases the effective gap scale compared to the bare Zeeman energy [40]. This separation of energy scales becomes particularly relevant in the tunnel-decoupled bilayer system at $\nu_{\text{tot}} = 7/3$, realized as $\nu_t = 4/3$ and $\nu_b = 1$. In this limit, the bilayer transport response corresponds to two parallel quantum Hall subsystems with negligible interlayer tunneling, such that the total longitudinal conductivity is approximately additive, $\sigma_{xx}^{\text{tot}} \approx \sigma_{xx}^{(t)} + \sigma_{xx}^{(b)}$, with each contribution exhibiting activated behavior $\sigma_{xx}^{(i)} \propto \exp[-\Delta_i/(2k_B T)]$. Consequently, $\sigma_{xx}^{\text{tot}}(T) \propto \exp[-\Delta_{4/3}/(2k_B T)] + \exp[-\Delta_1/(2k_B T)]$, implying that at sufficiently low temperature, the measured dissipation is dominated by the smaller gap layer. Therefore, the apparent activation gap extracted from ρ_{xx} [or $R_{xx}(B)$] at $\nu_{\text{tot}} = 7/3$ is expected to asymptotically approach $\Delta_{7/3} \approx \Delta_{4/3}$, unless interlayer Coulomb correlations, drag-induced coupling, and/or disorder-induced current partitioning modify the effective transport response [10,29,41,42]. Disorder is expected to play an enhanced role in such a composite bilayer state. Since the $\nu = 4/3$ state has a relatively small intrinsic gap, even moderate LL broadening or density inhomogeneity in that layer can substantially suppress the observed $\Delta_{7/3}$, while the $\nu = 1$ layer remains comparatively robust. This picture is consistent with observations in bilayer systems that even integer quantum Hall gaps, such as the balanced $\nu = 2$ state, exhibit strong dependence on carrier density and disorder, with reported activation gaps ranging from ~ 1 to 3 K [43], emphasizing that in two parallel quantum Hall subsystems, the experimentally extracted activation energy reflects not only intrinsic many-body excitation energies but also disorder broadening and the interplay of parallel conducting layers.

D. Angle dependence

Figure 5 shows the angle-dependent magnetoresistance measurements performed for $V_g \approx 0.055$ (single minimum) and ≈ -0.145 V (asymmetric double minimum). The effective applied perpendicular and in-plane magnetic fields are given by $B_{\perp} = B^* = B_{\text{total}} \cos \theta$ and $B_{\parallel} = B_{\text{total}} \sin \theta$, where θ is the tilt angle, respectively. In Fig. 5(a), for $V_g \approx 0.055$ V, we notice that R_{xx} slightly increases in resistance with increasing tilt angle until $\theta = 45^\circ$. Apart from this change, there are hardly

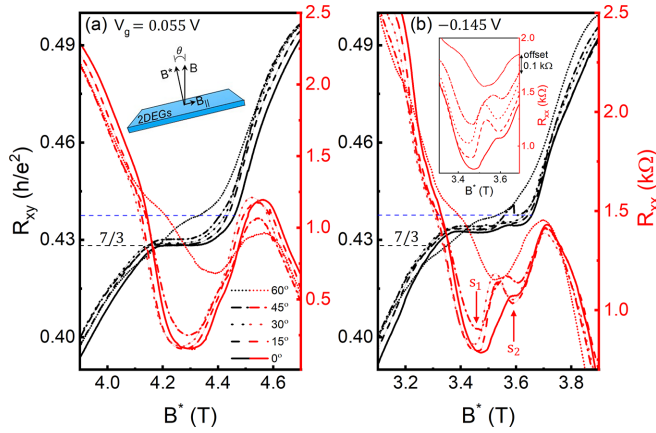


FIG. 5. The evolution of the s_1 and s_2 features as a function of tilting the magnetic field from 0° to 60° in steps of 15° for two gate voltages in sample 1: (a) $V_g \approx 0.055$ V and (b) $V_g \approx -0.145$ V. Both plots show R_{xy} and R_{xx} as a function of the effective magnetic field B^* . The inset in (a) shows a schematic representation of tilting the magnetic field relative to the sample. The inset in (b) shows a magnified view of s_1 and s_2 in the main plot, with traces vertically offset by 0.1 k Ω for clarity.

any significant variations. At $\theta = 60^\circ$, however, R_{xx} shifts to a higher magnetic field with a shallower minimum; R_{xy} also becomes shallower and shifts to $0.437(h/e^2)$. Figure 5(b), at $\theta = 0^\circ$, shows the double minimum s_1 and s_2 , which is asymmetric toward s_1 , with a much shallower s_2 ; the inset shows a magnified view of the evolution. With an increase in tilt to 30° , the valley of s_2 deepens, whereas it weakens substantially at 45° . Similar to the case for the single minimum [Fig. 5(a)], the 60° tilt in Fig. 5(b) results in a single, shallower minimum (and a shallower R_{xy} plateau), centered at the middle of the double minimum observed at other tilt angles in R_{xx} .

To quantitatively understand the tilt-dependent anomalous behavior observed at $\nu_{\text{tot}} = 7/3$, we should account for the specific bilayer geometry of our system. When the two layers reach their maximum incompressibility at slightly different magnetic fields, R_{xx} exhibits two distinct dips, s_1 toward low field and s_2 on the high-field side [Fig. 5(b)]. However, if their densities are (further) tuned with the gate voltage such that they still collectively form $\nu_{\text{tot}} = 7/3$ with their maximum incompressibilities perfectly aligned, they exhibit a single R_{xx} minimum, as in Fig. 5(a). Tilt-dependent magnetoresistance measurements of the former case reveal that the two minima (s_1 and s_2) behave entirely differently under an applied in-plane field. The s_1 state remains stable and predominantly unchanged under sample tilt. This confirms its origin tied to the $\nu_t = 1$ layer, which acts as a rigid, fully spin-polarized integer quantum Hall state. In contrast, the higher-field minimum s_2 (originating from the $\nu_b = 4/3$ layer) exhibits a dramatic dynamic evolution under sample tilt. It begins relatively shallow at $\theta = 0^\circ$, deepens significantly up to a 30° tilt, and subsequently shallows at 45° , clearly reflecting a spin polarization transition taking place in this layer. At $\theta = 60^\circ$ tilt, s_1 and s_2 eventually unite and form a shallow, single central minimum. The collapse of the $7/3$ R_{xx} minimum at large tilt angles is driven by strong $B_{\parallel}$ magneto-orbital coupling effects. This coupling softens intralayer Coulomb repulsion

by broadening electron wave functions and drives diamagnetic shifts that induce Landau level mixing and interlayer density imbalance, ultimately destabilizing the correlated many-body states [44–48]. In the other situation, when only a single minimum in R_{xx} is observed for the total $7/3$ filling [Fig. 5(a)], the tilt hardly influences the dip in R_{xx} minimum until $\theta = 45^\circ$. One possibility is that locked incompressibility and parallel conduction through the 2DEGs allow the robust $\nu = 1$ state to mask underlying $\nu = 4/3$ spin transitions, yielding moderate-tilt immunity. The abrupt extreme-tilt collapse suggests magneto-orbital coupling effects.

Halperin [49] initially proposed that the ground state trial wave function for the $2/5$ FQHE is unpolarized in GaAs due to the small Landé g factor, a view supported by multiple experiments. For filling factors less than 1, the unpolarized or partially polarized ground state is energetically preferred when the Coulomb energy E_C surpasses the Zeeman energy E_Z [50]. Such systems may undergo spin phase transitions as the Zeeman energy to Coulomb exchange energy ratio is adjusted either by increasing Zeeman energy (proportional to B) or altering electron density (proportional to $B^{1/2}$). Experimental results indicate that FQH states at $3/5$, $2/3$, $4/3$, $7/5$, and $8/5$ experience spin phase transitions as the Zeeman energy is varied through sample tilting or density changes [32,35,50–56]. These studies showed transitions from partially or unpolarized to fully spin-polarized states. Furthermore, phenomena like doubling of the R_{xx} minimum or subminima [32] and reentrant minima have been observed for $2/3$, $3/5$, and $8/3$ states [35,51,57], as well as for holes in GaAs at a filling factor of $4/3$ [55]. The characteristic tilt dependence of s_2 ($\nu = 4/3$ layer) is similar to reported results, suggesting a spin polarization transition occurring in the composite $7/3$ state formed in tunnel-decoupled bilayer electrons.

The total $\nu = 7/3$ state reported in this work presents a striking parallel to conventional high-index fractional states. While our state emerges in the LLL as a synthetic composite of two uncoupled layers ($\nu = 1$ and $\nu = 4/3$), its tilt-dependent characteristics resemble the Zeeman-driven spin features of the $\nu = 7/3$ state formed in the second Landau level (SLL; $N = 1$) of single 2DEGs [58–60]. The spin polarization of the SLL $\nu = 7/3$ state has long been a subject of intense scrutiny; foundational tilt-angle measurements by Eisenstein *et al.* [58] characterized it as a spin-polarized state, yet subsequent studies revealed a highly complex landscape, yielding contrasting reports on its exact spin nature under varying magneto-orbital conditions [45,59,61]. Therefore, our structurally distinct bilayer $7/3$ LLL state provides a unique platform to examine the interplay between spin and orbital dynamics. In the following section, we leverage the composite fermion framework to systematically untangle these competing many-body effects and elucidate the physical origin of this anomalous tilt dependence.

E. Composite fermion framework

The dynamic evolution of s_1 and s_2 in response to the variations in the tilt angle clearly suggests that a spin polarization transition occurs in the $\nu_b = 4/3$ layer [28]. This dynamic evolution is consistent with the standard composite fermion framework.

Within the composite fermion framework [4,49], the strongly interacting 2D electrons are mapped onto a system of weakly interacting quasiparticles known as composite fermions (CFs). Each CF is formed by the attachment of an even number of magnetic flux quanta ($2p'$, where p' is an integer, typically $2p' = 2$) to an electron. By incorporating the complex Coulomb interactions into the topological formation of these quasiparticles, the CFs effectively experience a reduced magnetic field $B_{\text{cf}}^* = B_{\perp} - 2p'\phi_0 n_{2D}$, where $\phi_0 = h/e$ is the magnetic flux quantum.

Under B_{cf}^* , CFs execute cyclotron motion and condense into discrete CF Λ levels. These levels are separated by a CF cyclotron gap, $E_{\omega_c}^{\text{cf}} = \hbar e |B_{\text{cf}}^*| / m_{\text{cf}}^*$, where m_{cf}^* is the dynamically generated many-body CF effective mass. Within this picture, the complex FQH states of electrons at filling factor ν map directly to integer quantum Hall states of CFs at an integer filling factor p . Specifically, for the $1 < \nu < 2$ regime studied here, the $\nu = 4/3$ state maps to a CF filling factor of $p = 2$, corresponding to the occupation of the lowest two Λ levels. This mapping provides a useful framework to analyze energy level crossings and spin polarization transitions driven by the crossing of the CF Λ levels. Specifically, the unpolarized ground state (occupying $\Lambda_0 \uparrow$ and $\Lambda_0 \downarrow$) transitions to the fully polarized ground state (occupying $\Lambda_0 \uparrow$ and $\Lambda_1 \uparrow$) driven by the competition between $E_{\omega_c}^{\text{cf}}$ and E_Z [4,28].

To provide a quantitative context for this crossover or spin polarization transition, we estimate the CF effective mass at $\theta = 0^\circ$, where the shallow R_{xx} minimum (s_2) indicates the system sits near the level crossing condition ($E_Z \approx E_{\omega_c}^{\text{cf}}$). In the $\nu_b = 4/3$ layer, the $\nu = 1$ condition occurs at $B_1 \approx 4.9$ T (the point of maximum incompressibility; Fig. 2). At the crossing field of $B_{\perp} \approx 3.6$ T, the CFs experience an effective magnetic field of $|B_{\text{cf}}^*| = |3B_{\perp} - 2B_1| \approx 1.0$ T. Equating the Zeeman energy to the CF cyclotron gap,

$$|g^*| \mu_B B_{\text{total}} = \frac{\hbar e B_{\text{cf}}^*}{m_{\text{cf}}^*}, \quad (2)$$

yields an estimated CF effective mass $m_{\text{cf}}^* \approx 1.26m_e$, where we used $|g^*| = 0.44$. This value falls within the canonical experimental range of $\sim 0.7m_e$ to $1.5m_e$ widely reported for CFs [34,62–64].

This estimation is corroborated by tracking the $E_Z/E_{\omega_c}^{\text{cf}}$ ratio through the experimental tilt angles. Since $E_{\omega_c}^{\text{cf}}$ depends only on $B_{\perp}$ (which is fixed), the ratio scales purely with the total field: $E_Z/E_{\omega_c}^{\text{cf}} \propto 1/\cos\theta$. Setting the baseline $E_Z/E_{\omega_c}^{\text{cf}} \approx 1.00$ at $\theta = 0^\circ$ (where the $\Lambda_0 \downarrow$ and $\Lambda_1 \uparrow$ levels overlap, causing the shallow reentrant gap), we can map the transition.

(1) At $\theta = 15^\circ$, $E_Z/E_{\omega_c}^{\text{cf}} \approx 1.04$. The Zeeman gap begins to exceed the CF cyclotron gap, initiating the separation of the Λ levels.

(2) At $\theta = 30^\circ$, $E_Z/E_{\omega_c}^{\text{cf}} \approx 1.15$. The condition $E_Z > E_{\omega_c}^{\text{cf}}$ is robustly satisfied, pulling the Λ levels entirely apart. This isolates the spin-polarized ground state, manifesting experimentally as deepening of the $\nu = 4/3$ minimum. The total field reaches $B_{\text{total}} \approx 4.16$ T.

(3) At $\theta = 45^\circ$ ($E_Z/E_{\omega_c}^{\text{cf}} \approx 1.41$) and $\theta = 60^\circ$ ($E_Z/E_{\omega_c}^{\text{cf}} \approx 2.00$), the extreme total magnetic field drives massive Zeeman dominance. The subsequent shallowing and collapse into a single minimum at 60° reflect the eventual tilt-induced

rebalancing of the layer densities and the suppression of the fractional gap due to magneto-orbital coupling effects at high total fields ($B_{\text{total}} \approx 7.20$ T).

We emphasize that the extracted $m_{\text{cf}}^* \approx 1.26m_e$ serves as an estimation to reflect the energy crossover, shaped by two competing many-body effects. In a low-density regime, electron-electron exchange interactions are known to enhance the effective g factor, which would lower the estimated mass to $\sim 0.8m_e$ to $1.0m_e$ [62]. However, this is counterbalanced by the capacitive coupling to the adjacent $\nu_t = 1$ layer, which screens and softens the effective intralayer Coulomb interaction within the active $\nu_b = 4/3$ layer. Given the CF effective mass is dynamically generated by Coulomb interactions ($m_{\text{cf}}^* \propto e^2/\epsilon\ell_B$), this interlayer screening inherently reduces $E_{\omega_c}^{\text{cf}}$, naturally increasing the CF effective mass relative to an isolated single layer [63]. The estimated mass of $1.26m_e$ and the tracked $E_Z/E_{\omega_c}^{\text{cf}}$ tilt dynamics are therefore highly consistent with a Λ level spin polarization transition occurring in a nontunneling bilayer system.

Finally, we discuss the absence of any noticeable tilt-dependent feature when the 7/3 state exhibits a single minimum, as shown in Fig. 5(a). Electrostatically tuning the charge distribution between the 2DEGs to a precise 4:3 density ratio via the top gate forces a magnetic field commensurability, superposing the thermodynamic gaps into a singular R_{xx} minimum. Strikingly, this commensurate state appears unresponsive to moderate $B_{\parallel}$ up to the 45° tilt, which may be attributed to the simultaneous transport taking place via two parallel 2DEGs. The energetically robust, fully spin-polarized $\nu = 1$ layer acts as a nondissipative electrical shunt, effectively masking the underlying CF spin dynamics of the $\nu = 4/3$ layer. At an extreme tilt of 60° , the R_{xx} minimum abruptly shallows, indicating that the system remains locked in resonance, resisting $B_{\parallel}$ -induced diamagnetic charge transfer [44–48]. Therefore, the quench is of a purely magneto-orbital origin, driving a sharp divergence in the CF effective mass that ultimately collapses the $\nu = 4/3$ fractional energy gap.

IV. PHENOMENOLOGICAL MODEL

Our system consists of a GaAs bilayer containing two nontunneling 2DEGs. The top gate is capacitively coupled to the bilayer system. For the model, we will focus on the positive gate voltage window, where we observe a quantized Hall plateau at $\nu_{\text{tot}} = 7/3$; the layer fillings satisfy $\nu_b \simeq 1$ and $\nu_t \simeq \frac{4}{3}$. The key experimental observation is that the 7/3 plateau is accompanied by (1) a single minimum in $R_{xx}(B)$ at some gate voltages, (2) a resolvable double-minimum structure in an intermediate gate voltage range, and then again (3) a single minimum before the 7/3 plateau disappears. Below, we develop a phenomenological model that connects this behavior to gate-dependent gap centers in the two layers, their disorder broadening, and the associated compressibility and electrostatic constraints.

A. Electrochemical equilibrium and capacitance

Let n_i and ϕ_i be the electron density and electrostatic potential of layer $i = t, b$, respectively. Each layer satisfies

$$\tilde{\mu}_i = \mu_i(n_i, B) + e\phi_i, \quad (3)$$

where $\mu_i(n_i, B)$ is the thermodynamic chemical potential of a disorder-broadened, Landau quantized 2DEG. Because both layers contact the same reservoir, the electrochemical potentials must be equal:

$$\tilde{\mu}_b = \tilde{\mu}_t = \tilde{\mu}_{\text{res}}. \quad (4)$$

The gated bilayer structure can be described by a parallel plate capacitance network. Denoting the gate-top layer and the interlayer (top-bottom) capacitances per unit area by C_{gt} and C_{tb} , respectively, we can write the electrostatic relations compactly as

$$\begin{pmatrix} -en_t \\ -en_b \end{pmatrix} = \begin{pmatrix} C_{gt} + C_{tb} & -C_{tb} \\ -C_{tb} & C_{tb} \end{pmatrix} \begin{pmatrix} \tilde{\phi}_t \\ \tilde{\phi}_b \end{pmatrix}, \quad (5)$$

where $\tilde{\phi}_t = \phi_t - V_g$ and $\tilde{\phi}_b = \phi_b - V_g$ are potentials relative to V_g . To include the effect of finite (quantum) compressibility κ_i [41], we define

$$\kappa_i = \frac{\partial n_i}{\partial \mu_i}, \quad C_{Q,i} = e^2 \kappa_i, \quad (6)$$

with $C_{Q,i}$ being the quantum capacitance of layer i . Incompressible quantum Hall states have $\kappa_i \rightarrow 0$ (and thus $C_{Q,i} \rightarrow 0$), whereas compressible regions have large κ_i . Linearizing the electrochemical equilibrium condition, Eq. (4), and the electrostatic relation, Eq. (5), for small changes δV_g , δn_i , and $\delta \phi_i$, we find

$$\begin{pmatrix} \delta n_t \\ \delta n_b \end{pmatrix} = \mathbf{M}(B, V_g) \delta V_g, \quad (7)$$

which defines a 2×2 matrix $\mathbf{M}(B, V_g)$ that depends on C_{gt} , C_{tb} , and the quantum capacitances $C_{Q,i}(B, V_g)$. Because C_{tb} is finite, when one layer becomes incompressible and $\partial \mu_i / \partial n_i$ becomes large, Eq. (4) forces small, but non-negligible, charge transfer into the other layer. Using $d \simeq 45$ nm, the distance between the middles of the quantum wells, we estimate

$$C_{tb} \simeq 2.5 \times 10^{-3} \text{ F/m}^2, \quad (8)$$

so a chemical potential jump $\Delta \mu$ [8,19,20,29,65–68] in the range of 1–3 meV produces a density shift

$$\Delta n \sim \frac{C_{tb} \Delta \mu}{e^2} \approx (1.6\text{--}4.8) \times 10^9 \text{ cm}^{-2}, \quad (9)$$

which is consistent with the difference in carrier density associated with the magnetic fields at the double minima (Figs. 2 and 3). These shifts are large enough to move the layer-specific incompressible points by $\sim 0.1\text{--}0.2$ T. When a given layer is incompressible, $C_{Q,i} \rightarrow 0$, and its density is effectively pinned, so that most gate-induced charge flows into the other layer or is accommodated by changes in the electrostatic potentials. When both layers are simultaneously incompressible, Eq. (7) admits only a very narrow range of densities consistent with Eq. (4), resulting in a correspondingly narrow gate voltage window in which the condition $(\nu_t, \nu_b) \simeq (4/3, 1)$ and hence $\nu_{\text{tot}} = 7/3$ can be satisfied.

B. Interaction-driven chemical potential structure

The internal chemical potential may be written as

$$\mu_i(n_i, B, \theta) = \frac{\partial E_{\text{int}}^{(i)}}{\partial n_i}, \quad (10)$$

where $E_{\text{int}}^{(i)}$ includes interaction effects strongly nonlinear in the LLL [22]. Near an incompressible filling, the chemical potential exhibits sharp cusps and jumps [41,65,69,70], defined as

$$\mu_i(n_i, B, \theta) = \mu_i^{(0)}(n_i, B) + \Delta_i(B, \theta) g_i(\nu_i), \quad (11)$$

where $\mu_i^{(0)}(n_i, B)$ is a smooth, slowly varying background part of the chemical potential (compressible behavior inside a broadened LL), $g_i(\nu_i)$ is sharply peaked at each layer's incompressible filling, and $\Delta_i(B, \theta)$ is the effective incompressible gap of layer i . The microscopic origin of these incompressible states differs: The $\nu = 1$ gap is expected to be dominated by exchange-enhanced spin splitting, while the $\nu = 4/3$ gap arises from many-body correlations and LLL pseudopotentials. This asymmetry explains the differing widths, gap magnitudes, and disorder sensitivity of the two minima. In most cases, it is the interaction-induced gap that controls how sharply the layer becomes incompressible [22,41,71,72]. These nonlinearities produce strong compressibility minima (i.e., a jump in chemical potential),

$$\kappa_i^{-1} = n_i^2 \frac{\partial \mu_i}{\partial n_i}, \quad (12)$$

and amplify the effects of the capacitive coupling in Eq. (4), which in turn strongly enhances the sensitivity of each layer's incompressible point to the small density shifts produced by interlayer capacitive coupling. Thus, interaction effects determine the widths, magnitudes, and stability of the two minima, while capacitive coupling determines their relative positions.

C. Gap centers and their gate dependence

For a fixed gate voltage, the optimal field $B_i^o(V_g)$ at which the i th incompressible state is most robust is defined by the condition that the filling factor in that layer equal the target value:

$$\nu_b(B_b^o, V_g) = 1, \quad \nu_t(B_t^o, V_g) = \frac{4}{3}. \quad (13)$$

Equivalently, in terms of the layer densities,

$$B_b^o(V_g) = \frac{h n_b(V_g)}{e}, \quad B_t^o(V_g) = \frac{3 h n_t(V_g)}{4 e}. \quad (14)$$

The functions $n_i(V_g)$ are implicit solutions of the coupled electrostatic and compressibility equations, (4) and (7). Because the bottom layer typically has a large integer gap and small quantum capacitance $C_{Q,b}$ in the relevant field range, its density is more strongly pinned, and $B_b^o(V_g)$ varies only weakly with gate voltage. By contrast, the top layer responds more strongly to the gating effect, so that $n_t(V_g)$ and hence $B_t^o(V_g)$ evolve more rapidly with V_g . This different lever arm is central to the emergence and gate evolution of the double-minimum structure in R_{xx} .

D. Spin-resolved gaps

Each layer experiences a set of spin-resolved Landau levels in the LLL. The orbital quantization is governed by the perpendicular magnetic field $B_{\perp} = B^* = B \cos \theta$, while the Zeeman splitting depends on the total magnetic field B (or

$B_{\text{total}})$,

$$E_Z(\theta) = g^* \mu_B B = g^* \mu_B \frac{B_{\perp}}{\cos \theta}. \quad (15)$$

Tilting the sample therefore increases E_Z at fixed $B_{\perp}$, allowing spin-dependent gaps to be selectively strengthened or weakened without altering the orbital Landau quantization. In the magnetic field range of interest, $\nu_b = 1$, and $\nu_t = 4/3$. Their ideal gaps can be written as

$$\Delta_1^0(B, \theta) \simeq E_Z(\theta) + \Delta_1^{\text{ex}}(B), \quad (16)$$

$$\Delta_{4/3}^0(B, \theta) \simeq \alpha \frac{e^2}{\varepsilon \ell_B(B)} f\left(\frac{E_Z(\theta)}{E_C(B)}\right). \quad (17)$$

The integer gap contains an exchange-enhanced Zeeman component, Δ_1^{ex} , and is typically only weakly angle dependent. By contrast, the fractional 4/3 state has a strong spin component and can undergo spin transition [35,55]. Its gap depends on the ratio E_Z/E_C , and the function f interpolates between two regimes: (1) a partially spin-polarized phase at small E_Z/E_C and (2) a fully polarized phase once E_Z/E_C exceeds a critical value.

In experimental samples the ideal gaps are reduced by the disorder broadening Γ_i , which represents the width of the localized states separating the extended states at the mobility edges [34,73,74]. We define an effective mobility gap

$$\Delta_i^{\text{eff}}(B, \theta) = \Delta_i^0(B, \theta) - \Gamma_i, \quad (18)$$

which is the energy scale governing longitudinal activation. In this formulation, tilting the sample modifies the fractional gap strongly through its enhanced Zeeman energy, while affecting the integer gap only weakly. This distinction explains why one of the observed minima in R_{xx} remains nearly unchanged under tilt while the other deepens and, eventually, both unite and weaken for $\theta = 60^\circ$ tilt.

E. Zeeman versus magneto-orbital coupling

Tilted field measurements separate the roles of Zeeman energy and orbital coupling. The Zeeman scale grows with the total magnetic field, Eq. (15), and can drive spin reconfigurations of fractional states. However, in GaAs the in-plane component, $B_{\parallel} = B_{\perp} \tan \theta$, can also couple orbitally to a finite thickness wave function, renormalizing interaction pseudopotentials and often suppressing fragile FQH gaps [45,48,75–77]. Distinct responses of the 7/3 and 5/2 states to tilt, with emphasis on the importance of magneto-orbital coupling in finite thickness 2DEGs, have been extensively reported [45,48,75,76].

A useful order of magnitude criterion for the onset of strong magneto-orbital coupling is that the in-plane magnetic field threads a significant fraction of a flux quantum through the effective area spanned by the magnetic length ℓ_B and the vertical extent w_{eff} of the electronic wave function [44,78]:

$$\zeta \equiv \frac{B_{\parallel} w_{\text{eff}} \ell_B}{\phi_0}, \quad \phi_0 = \frac{h}{e}. \quad (19)$$

For $B_{\perp} \simeq 4.3$ T and $\ell_B \simeq 12.4$ nm at $\theta = 60^\circ$, $B_{\parallel} \simeq 7.45$ T. If we take w_{eff} between a conservative rms value (~ 4 nm) and an upper bound comparable to the well width (15 nm),

TABLE I. Energy scales at the two R_{xx} minima (s_1 and s_2). Here $E_C^{\text{intra}} = e^2/(4\pi \varepsilon \varepsilon_0 \ell_B)$, $E_C^{\text{inter}} = e^2/(4\pi \varepsilon \varepsilon_0 d)$, and $\kappa_{\parallel} = E_C^{\text{intra}}/E_{\omega_c}$.

Quantity	s_1 (4.20 T)	s_2 (4.35 T)
ℓ_B (nm)	12.52	12.31
E_C^{intra} (meV)	8.91	9.07
E_{ω_c} (meV)	7.26	7.52
$\kappa_{\parallel}$	1.23	1.21
E_C^{inter} (meV)	2.48	2.48
E_Z (meV), 0°	0.107	0.111
E_Z (meV), 60°	0.214	0.222

Eq. (19) yields $\zeta \sim 0.1$ – 0.3 at 60° tilt, suggesting weak to moderate orbital coupling. Because $B_{\parallel}$ grows rapidly with $\tan \theta$, modest changes in θ around 60° can move the system across a regime where magneto-orbital renormalization becomes quantitatively important, producing an apparently abrupt change in the R_{xx} line shape and a degradation of quantization.

As $B_{\parallel}$ introduces an orbital coupling [45,48,75–77], the bilayer system with finite layer separation d produces a relative displacement of the Landau guiding centers between the two layers,

$$\Delta x \simeq d \tan \theta, \quad (20)$$

which strongly modifies the interlayer Coulomb form factor for $\Delta x \gtrsim \ell_B$. The guiding center shift has been analyzed in tilted field effects and tunneling in bilayer systems [76,77,79]. With $d \simeq 45$ nm, we estimate $\Delta x \simeq 78$ nm ($\sim 6.2 \ell_B$) at 60° tilt. Such large displacements strongly modify the Coulomb form factor and can disrupt the self-consistent balance implied by Eq. (4), particularly when one layer is incompressible and the other is near a spin-sensitive fractional gap. In this sense, the abrupt evolution near 60° tilt is plausibly a combined effect. The increase in Zeeman energy modifies the spin configuration and therefore the excitation gap of the 4/3 fractional state, while the rapidly growing in-plane magnetic field component and the associated interlayer guiding center displacement introduce strong orbital coupling. These tilt-induced orbital effects modify the effective Coulomb interaction within and between the two quantum wells, weaken the incompressibility of the fractional layer, and, ultimately, drive the collapse of the double-minimum structure into a single shallow minimum.

F. Energy scales of the double minimum

Table I summarizes the relevant microscopic scales at the two longitudinal minima ($B_{\perp} \approx 4.20$ and 4.35 T). Several conclusions follow.

(1) Coulomb and cyclotron energies are comparable, implying moderate LL mixing. At both minima the intralayer Coulomb scale is $E_C^{\text{intra}} \simeq 9$ meV, while the cyclotron energy is $E_{\omega_c} \simeq 7.3$ – 7.5 meV. The resulting LL mixing parameter $\kappa_{\parallel} = E_C^{\text{intra}}/E_{\omega_c} \simeq 1.2$ indicates *moderate* LL mixing, which means for the FQH state, the LLL description remains qualitatively appropriate, but quantitative gap values and line shapes can be renormalized by LL mixing and finite thickness form

factors [80,81]. In particular, small changes in effective interaction can produce noticeable changes in fragile fractional gaps, consistent with the observed sensitivity of the 4/3 related minimum.

(2) Interlayer Coulomb coupling is substantial but not dominant. For $d \simeq 45$ nm, the characteristic interlayer Coulomb scale is $E_C^{\text{inter}} \simeq 2.48$ meV, approximately 27% of the intralayer scale at these fields:

$$\frac{E_C^{\text{inter}}}{E_C^{\text{intra}}} \simeq \frac{\ell_B}{d} \approx 0.27\text{--}0.28. \quad (21)$$

Thus, although the layers do not tunnel, they are not independent: Interlayer Coulomb coupling is strong enough to enforce self-consistent electrostatic constraints and to mediate charge redistribution when one layer becomes incompressible. At the same time, because $E_C^{\text{inter}} \ll E_C^{\text{intra}}$, the dominant correlation physics that sets the fractional $\nu = 4/3$ gap remains primarily intralayer. This hierarchy supports our interpretation in which the total 7/3 plateau arises from parallel conduction of two incompressible fluids, while the detailed R_{xx} line shape is modified by interlayer electrostatics.

(3) Zeeman energy is small on the Coulomb scale but can still control the spin sector of the fractional state. At zero tilt ($\theta = 0^\circ$), $E_Z \simeq 0.11$ meV, i.e.,

$$\frac{E_Z}{E_C^{\text{intra}}} \simeq 1.2 \times 10^{-2}, \quad (22)$$

while at $\theta = 60^\circ$, E_Z increases to ~ 0.22 meV. Even at the largest tilt explored, $E_Z/E_C^{\text{intra}} \lesssim 0.02$, so Zeeman splitting remains a small correction relative to Coulomb correlations. Nevertheless, for spin-sensitive fractional states, changes in E_Z on the order of a few tenths of a meV can be comparable to differences between competing spin configurations and therefore can selectively modify the effective 4/3 gap. This accounts for the observed asymmetric tilt response of the two minima: The integer-related minimum is comparatively insensitive, while the fractional-related minimum evolves strongly with tilt.

(4) Alignment of incompressible layers merges the double minima. The small field separation between the two minima (~ 4.20 versus 4.35 T) corresponds to only a modest change in ℓ_B and hence only an $\sim 2\%$ change in E_C^{intra} across the two features, while E_C^{inter} remains fixed by geometry. This supports the picture that the double-minimum structure is controlled mainly by the relative alignment of the incompressible points in the two layers, governed by gate-controlled density partitioning and the electrochemical equilibrium constraint. When the gate tuning brings the incompressible points into coincidence, the dissipative minima overlap, and the total R_{xx} exhibits a single minimum.

(5) Connection to the abrupt evolution near large tilt. Although E_Z grows with the total magnetic field and can change the spin sector of the fractional state, the rapid evolution near $\theta \sim 60^\circ$ is also consistent with the onset of strong magneto-orbital effects (see Sec. IV E) from the rapidly increasing in-plane magnetic field component. The values in Table I highlight that such tilt-induced orbital renormalization acts in a system where the dominant energy scales are Coulomb

and cyclotron, while Zeeman energy provides a smaller, but symmetry-selective, perturbation.

Overall, the quantitative hierarchy $E_Z \ll E_{\omega_c} \sim E_C^{\text{intra}}$ with $E_C^{\text{inter}} \sim 0.27 E_C^{\text{intra}}$ supports a conservative physical interpretation: The bilayer 7/3 plateau reflects parallel conduction of two LLL incompressible fluids, while the double-minimum structure and its evolution reveal how the layer-resolved incompressibility is reshaped by finite inter- and intralayer Coulomb coupling, constrained charge redistribution, and spin-sensitive renormalization of the fractional gap under tilt.

G. Activated Gaussian profiles

The longitudinal conductivity of layer i near its incompressible state may be approximated by an activation law across the disorder-broadened mobility gap,

$$\sigma_{xx}^{(i)}(B, V_g) \propto \exp\left[-\frac{\Delta_i^{\text{eff}}(B, V_g)}{2k_B T}\right], \quad (23)$$

with Δ_i^{eff} given by Eq. (18). We now focus on the dependence of the gap on magnetic field near its maximum at $B = B_i^o(V_g)$. We expand the ideal gap as

$$\Delta_i^0(B, \theta, V_g) \simeq \Delta_{i,\text{max}}^0(\theta, V_g) - \frac{1}{2} \kappa_i(V_g) [B - B_i^o(V_g)]^2, \quad (24)$$

where $\Delta_{i,\text{max}}^0(V_g)$ is the maximum of the ideal gap and $\kappa_i(V_g) > 0$ is the curvature that characterizes how rapidly the gap decreases as we detune B from the optimal value. Substituting Eq. (24) into Eq. (18), we obtain

$$\Delta_i^{\text{eff}}(B, \theta, V_g) \simeq \Delta_{i,\text{max}}^{\text{eff}}(\theta, V_g) - \frac{1}{2} \kappa_i(V_g) [B - B_i^o(V_g)]^2, \quad (25)$$

with

$$\Delta_{i,\text{max}}^{\text{eff}}(\theta, V_g) \equiv \Delta_{i,\text{max}}^0(\theta, V_g) - \Gamma_i. \quad (26)$$

Inserting this into the activation law (23) gives

$$\begin{aligned} \sigma_{xx}^{(i)}(B, \theta, V_g) &\propto \exp\left[-\frac{\Delta_{i,\text{max}}^{\text{eff}}(V_g)}{2k_B T} + \frac{\kappa_i(V_g)}{4k_B T} [B - B_i^o(V_g)]^2\right] \\ &= A_i(V_g) \exp\left[-\frac{[B - B_i^o(V_g)]^2}{2 \delta B_i^2(V_g)}\right], \end{aligned} \quad (27)$$

where we have defined

$$A_i(V_g) \propto \exp\left[-\frac{\Delta_{i,\text{max}}^{\text{eff}}(V_g)}{2k_B T}\right], \quad \delta B_i^2(V_g) = \frac{2k_B T}{\kappa_i(V_g)}. \quad (28)$$

Equation (27) shows that the longitudinal conductivity of each layer exhibits a Gaussian dip in B centered at $B_i^o(V_g)$, with the width δB_i set by the curvature of the gap and the temperature. In practice, the experimental widths manifest both the true temperature and the disorder-induced broadening Γ_i ; this can be absorbed into an effective temperature in Eq. (28).

To include a smooth background $\sigma_{xx,\text{bg}}^{(i)}$, we write

$$\begin{aligned} \sigma_{xx}^{(i)}(B, \theta, V_g) &= \sigma_{xx,\text{bg}}^{(i)}(B, \theta, V_g) \\ &\quad - \tilde{A}_i(V_g) \exp\left[-\frac{[B - B_i^o(V_g)]^2}{2 \delta B_i^2(V_g)}\right], \end{aligned} \quad (29)$$

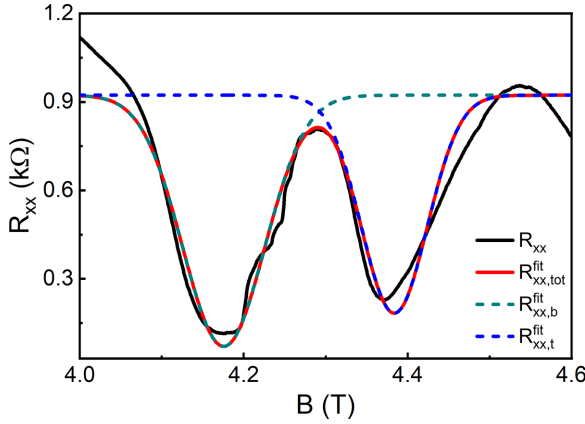


FIG. 6. $R_{xx}(B)$ near $\nu_{\text{tot}} = 7/3$ measured at $V_g \approx 0.05$ V. The black solid curve shows the experimental data. The red solid curve represents the total fit $R_{xx,\text{tot}}^{\text{fit}}$ obtained from a two-component model consisting of a contribution from the bottom layer $R_{xx,b}^{\text{fit}}$ at $\nu_b = 1$ (cyan dashed curve) and a gate-tunable fractional contribution from the top layer $R_{xx,t}^{\text{fit}}$ at $\nu_t = 4/3$ (blue dashed curve). The relative separation of the two fitted minima evolves with gate voltage, while the bottom layer minimum remains pinned, consistent with electrostatic tuning of the upper layer density at fixed lower layer density.

where $\tilde{A}_i$ absorbs proportionality factors from Eq. (27).

Writing the measured longitudinal resistance as $R_{xx}^{\text{tot}}(B, V_g) \propto \rho_{xx}^{\text{tot}}(B, V_g)$ and substituting Eq. (29) for both layers, we arrive at the two-dip form

$$R_{xx}^{\text{tot}}(B, \theta, V_g) = R_{bg}(B, V_g) - A_b(V_g, \theta) \exp \left[-\frac{[B - B_b^o(V_g)]^2}{2\delta B_b^2} \right] - A_t(V_g, \theta) \exp \left[-\frac{[B - B_t^o(V_g)]^2}{2\delta B_t^2} \right], \quad (30)$$

where $R_{bg}(B, V_g)$ is a slowly varying background and the amplitudes $A_i(V_g)$ inherit the gap dependence through Eq. (28). Equation (30) is the central phenomenological expression we use to interpret the evolution of R_{xx} with gate voltage and magnetic field.

H. Gaussian gap profile fit to $R_{xx}(B)$

The phenomenological model described above is not intended as a microscopic theory of the bilayer system, but rather as a minimal framework to interpret how two independent incompressible quantum Hall fluids, coupled only electrostatically, can generate a total $R_{xx}(B)$ response at a fixed total Hall plateau. In this picture, the measured $R_{xx}(B)$ reflects the sum of two activated dissipative channels (layers) whose gap magnitudes peak at their respective incompressible points, which we model phenomenologically using Gaussian field profiles. We show the internal consistency of this approach by fitting a representative trace in which the Hall plateau at a total filling factor $7/3$ is accompanied by a pronounced double-minimum structure in $R_{xx}(B)$. As shown in Fig. 6, the experimental line shape (R_{xx} , black curve) is well

reproduced (R_{xx}^{fit} , red curve) by the superposition of a fixed integer component associated with the bottom layer $\nu = 1$ gap (R_b^{fit} , cyan dashed curve) and a second, gate-tunable fractional component associated with the top layer $\nu = 4/3$ gap (R_t^{fit} , blue dashed curve) using fitting parameters A_b , A_t , R_{bg} , B_b , and B_t set to 0.85, 0.74, 0.92, 4.18, and 4.38, respectively. The resulting decomposition supports the qualitative interpretation that the two minima arise from a mismatch of the incompressible gap centers in the two layers. Within this view, the collapse of the double minimum into a single minimum occurs when the two gap centers coincide in the magnetic field, yielding a single combined activated minimum while preserving the total $7/3$ plateau. More generally, the phenomenological fit highlights that the detailed evolution of the longitudinal response in a nontunneling bilayer is governed not only by the existence of incompressible states in each layer but also by electrostatic compressibility constraints and interlayer charge redistribution, which can shift the effective filling of the fractional layer while leaving the integer layer comparatively pinned.

I. Phenomenological interpretation of $R_{xx}(B)$

Using Eq. (30), we discuss the observed sequence of R_{xx} line shapes as a function of V_g in the experimental results shown in Figs. 2 and 3.

(1) *Low gate voltage.* For $V_g \approx 0.01$ – 0.02 V near the lower bound of the $7/3$ window, the conditions $\nu_b \simeq 1$ and $\nu_t \simeq 4/3$ are already satisfied in the accessible field range, so both amplitudes A_b and A_t are nonzero, and the Hall plateau is quantized at $7/3$. However, the two gap centers $B_b^o(V_g)$ and $B_t^o(V_g)$ may be nearly coincident, or the fractional gap amplitude A_t may be significantly smaller than A_b , so the two Gaussian dips in Eq. (30) overlap and appear experimentally as a single minimum in $R_{xx}(B)$.

(2) *Intermediate gate voltage.* As V_g is increased to ~ 0.05 V, the coupled electrostatics and compressibility shown in Eq. (7) produce a larger change in $n_t(V_g)$ than in $n_b(V_g)$, so the centers $B_t^o(V_g)$ and $B_b^o(V_g)$ separate [Eq. (14)]. In a range of V_g where both gaps are still robust (both A_b and A_t are sizable), the difference $|B_b^o - B_t^o|$ becomes comparable to or larger than the widths δB_i , and the total R_{xx} clearly resolves two distinct minima: one associated with the integer $\nu_b = 1$ state in the bottom layer and one associated with the fractional $\nu_t = 4/3$ state in the top layer.

(3) *Higher gate voltage within the $7/3$ window.* With a further increase of V_g to ~ 0.06 V, the fractional gap in the top layer weakens (due to changes in either density, spin polarization, or the enhanced relative importance of disorder), reducing $A_t(V_g)$. Alternatively, the gap centers may get closer together again because of interlayer charge redistribution constraints. In either case, Eq. (30) predicts that the two Gaussian dips merge into a single apparent minimum, even though the Hall response remains quantized at $\nu_{\text{tot}} = 7/3$ over this gate voltage range.

(4) *Beyond the $7/3$ window.* At still higher gate voltages, the density configuration consistent with Eq. (4) and the capacitive coupling can no longer support a robust $\nu_t = 4/3$ state within the experimental field window. The effective gap $\Delta_{4/3}^{\text{eff}}$

closes, $A_t(V_g) \rightarrow 0$, and the top layer becomes compressible, leading to the collapse of the $7/3$ plateau.

V. DISCUSSION

Our measurements reveal a rich structure in the longitudinal magnetoresistance of a nontunneling GaAs bilayer at $\nu_{\text{tot}} = 7/3$. As the density of the top layer is tuned by a gate voltage over a narrow ~ 60 mV window, the $R_{xx}(B)$ line shape evolves from a single minimum, through an asymmetric then symmetric double-minimum structure, and back to a single minimum before the $7/3$ Hall plateau disappears altogether. Within the phenomenological framework, these features can be understood as direct signatures of (1) the gate-dependent positions of the integer ($\nu_b \simeq 1$) and fractional ($\nu_t \simeq 4/3$) gaps in the two layers, (2) their distinct disorder-broadened activation profiles, and (3) the constraints imposed by electrochemical equilibrium and interlayer capacitance in a low compressibility bilayer.

At a qualitative level, the emergence of a double minimum is a hallmark of two distinct incompressible transitions occurring at nearby but nonidentical magnetic fields. In our nontunneling bilayer geometry, these transitions can be identified with the formation of the bottom layer $\nu_b = 1$ state and the top layer $\nu_t = 4/3$ state. The electrostatic and compressibility analysis of Eqs. (4)–(7) shows that the density in the bottom layer is strongly pinned whenever the $\nu_b = 1$ gap is open, while the top layer density responds more strongly to the gate. In other words, when one layer becomes incompressible, its quantum capacitance collapses, and $\partial n/\partial \mu$ becomes small, causing its density to become effectively pinned. Additional induced charge must then be absorbed primarily by the other layer or by shifts in the electrostatic potentials. For the present device parameters ($C_{tb} > C_{gt}$, symmetric wells, and negligible tunneling), the top layer responds more strongly to the gate voltage, while the bottom layer changes only weakly once $\nu_b = 1$ is reached. Consequently, as V_g is increased, the effective field $B_t^o(V_g)$ at which the top layer reaches $\nu_t = 4/3$ moves more rapidly than the corresponding $B_b^o(V_g)$ of the bottom layer, which changes only weakly [Eq. (14)]. In the language of our two-dip expression [Eq. (30)], the relative motion of the centers $B_b^o(V_g)$ and $B_t^o(V_g)$ with gate voltage controls whether the two Gaussian dips in $R_{xx}^{\text{tot}}(B, V_g)$ are overlapped (single minimum), partially separated (asymmetric double minimum), or well resolved (symmetric double minimum).

A second crucial component is the role of disorder and its interplay with the intrinsic gaps. In our phenomenological model, disorder broadening is captured in the energy scale Γ_i appearing in Eq. (18), which reduces the ideal gaps to effective mobility gaps. Upon expanding the gap near its maximum in B , the longitudinal conductivity acquires a Gaussian profile in field, Eq. (27), with a width δB_i determined by both the curvature of the gap and the effective temperature; the latter is dominated by the disorder broadening at the lowest temperatures.

If we change V_g while remaining within the $7/3$ window, the detailed balance between electrostatics, spin, and

disorder begins to disfavor the fractional state. In the model this can appear as a reduction in A_t , a broadening of δB_t , or a partial reapproach of B_t^o toward B_b^o due to interlayer charge redistribution under the constraint of electrochemical equilibrium. Experimentally, this is observed as a merging of the two minima into a single feature, even though the Hall resistance remains quantized at $\nu_{\text{tot}} = 7/3$. Finally, at gate voltages beyond the upper edge of the $7/3$ window, the density configuration compatible with $(\nu_t, \nu_b) \simeq (4/3, 1)$ can no longer be realized at any field between 3 and 5 T; the top layer $4/3$ gap collapses ($A_t \rightarrow 0$), and the total $7/3$ plateau disappears.

Tilted field measurements provide an additional and more selective probe of the two gaps. In our experiments, one of the minima in R_{xx} is essentially insensitive to tilt, while the other deepens with increasing tilt and eventually weakens as the $7/3$ plateau itself becomes fragile. In the present framework, this behavior is interpreted as the distinct spin structure of the two gaps. The integer $\nu_b = 1$ state is expected to be fully spin polarized in the relevant field and tilt range, so that its effective gap is only weakly modified by changes in Zeeman energy. By contrast, the fractional $\nu_t = 4/3$ gap is known to be sensitive to spin polarization; increasing E_Z can initially stabilize a polarized state, enhancing $\Delta_{4/3, \text{max}}^{\text{eff}}$ and thus A_t , but at larger tilt it may destabilize the fractional order altogether due to magneto-orbital effects. This is the trend inferred from the tilt dependence of the phenomenological amplitudes and is consistent with the eventual collapse of the double-minimum structure and the weakening of the $7/3$ plateau at the highest tilt angles.

VI. CONCLUSIONS

Our results show that a nontunneling bilayer system can reveal detailed information about layer-resolved incompressibility and gap evolution in the line shape of longitudinal magnetoresistance as a function of carrier density. The non-monotonic appearance and disappearance of two minima in R_{xx} , together with the limited gate window over which the total $7/3$ plateau exists, directly reflect the interplay among capacitance, compressibility, Coulomb coupling, and spin-dependent fractional gaps in the LLL. The same framework should be applicable to other nontunneling bilayers, including electron-hole and electron-composite fermion bilayers, for which correlated states at nontrivial combinations of layer fillings have been proposed. In this sense, the gate-tunable $7/3$ bilayer studied here provides a platform for exploring how incompressible fluids in different layers compete, coexist, and, ultimately, reconstruct under the combined influence of disorder, spin, and Coulomb interactions. Additionally, our experimental results have important implications for understanding the spin states associated with the total $7/3$ FQHE in uncoupled bilayer electron systems.

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DATA AVAILABILITY

The data that support the findings of this study are available from the corresponding author upon reasonable request.

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