

Magnetoelectric subbands in the one-dimensional to quasi-one-dimensional confined regime

Yingshi Duo¹, Joshua Coop¹, Ian Farrer², David A. Ritchie³ and Sanjeev Kumar^{1,*}

¹*Department of Electronic and Electrical Engineering, University College London, Torrington Place, London WC1E 7JE, United Kingdom*

²*School of Electrical & Electronic Engineering, University of Sheffield, Mappin Street, Sheffield S1 3JD, United Kingdom*

³*Cavendish Laboratory, J.J. Thomson Avenue, Cambridge CB3 0HE, United Kingdom*



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We investigate the evolution of conductance quantization and subband structure in a weakly confined quasi-one-dimensional channel as a function of carrier density and perpendicular magnetic field. As the carrier density is reduced, the $2e^2/h$ plateau is progressively suppressed and an avoided crossing develops between the lowest two subbands in transconductance, indicating a strong intersubband coupling. This behavior is consistent with an interaction-driven restructuring of the electron system into a zigzag or two-row configuration in the regime of dominant Coulomb interactions. The application of a perpendicular magnetic field suppresses the anticrossing through enhanced magnetoelectric confinement, which reduces inter-row tunneling and hybridization via magnetic compression of the wave function. At intermediate fields (~ 0.5 T), further reduction of carrier density leads to the emergence of the $0.7(2e^2/h)$ feature coexisting with the restored $2e^2/h$ plateau. In addition, spin-split features, such as $1.5(2e^2/h)$, begin to resolve in higher subbands in the low-density, weakly confined regime. At higher fields (~ 0.7 T), where the anticrossing is fully quenched, the low-density regime exhibits the coexistence of conductance features near $2e^2/h$, $0.7(2e^2/h)$, and $0.5(2e^2/h)$. These observations indicate that perpendicular magnetic fields predominantly modify the orbital spectrum while leaving the strongly correlated spin sector intact. The resulting transport reflects the interplay between Coulomb interactions, exchange coupling, and Zeeman splitting, placing the system near the crossover between spin-coherent and spin-incoherent regimes.

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I. INTRODUCTION

Due to the control over the spin and charge of electrons in restricted dimensions, one-dimensional (1D) quantum wires based on GaAs/AlGaAs heterostructures have generated significant interest for potential applications in spintronics, charge sensing, emerging quantum technologies, and nanoelectronics [1–13]. A 1D system is created by defining a pair of split gates on the surface of a GaAs/AlGaAs heterostructure containing the two-dimensional electron gas (2DEG) [14]. This gating method creates a 1D constriction or quantum wire, which exhibits ballistic transport where conductance is quantized in units of $N(2e^2/h)$ (with the factor 2 accounting for the spin degeneracy of electrons, e is the charge of an electron, h is Planck's constant, and $N = 1, 2, 3, \dots$) for noninteracting electrons [15,16]. However, when Coulomb interactions between 1D electrons are considered, deviations from the standard single-particle model can occur [1,3,4,12,17–20].

It has long been predicted that a 3D system of electrons on a metal surface would form an electron lattice as the carrier density decreases, first proposed by Wigner in 1934

[21]. Later theoretical developments in the field led to the suggestion of a 2D version of the Wigner lattice on the surface of liquid helium, which was subsequently observed in experiments [22–24]. The concept of a 2D Wigner lattice was initially proposed for electrons in the inversion layer of semiconductors at low densities [25]. Subsequently, it was shown that a low-density 1D chain, driven by increased interaction energy and reduced kinetic energy, could transform into a zigzag string [26]. Further theoretical developments supported the emergence of multiple rows of electrons in the 1D Wigner regime in quantum wires with a reduction in the carrier density and the confinement potential [27–30]. In the low confinement regime, Coulomb interaction dominates over kinetic energy, causing 1D electrons to shift transversely and form a zigzag pattern; by further tuning the instability along the transverse displacement, the system settles into two or more rows of electrons. However, quantum and potential fluctuations and the absence of long-range order within 1D channels make it difficult to directly observe the 1D Wigner lattice [31].

Experimentally, the formation of the 1D Wigner lattice was reported in weakly confined quasi-1D quantum wires. In conductance measurements, at low carrier density, the first plateau at $2e^2/h$ melted, resulting in the new ground state at $2(2e^2/h)$, indicating that the 1D channel has localized into two rows, each contributing $2e^2/h$ to the conductance [3,19,32]. This was suggested as a reflection of a quantum phase transition in ballistic quantum wires in the Wigner regime [33]. This phenomenon was directly imaged through transverse electron

*Contact author: sanjeev.kumar@ucl.ac.uk

focusing experiments, which visualized the formation of the 1D Wigner zigzag phase [34]. Also, the 1D Wigner lattice has been investigated in carbon nanotubes using imaging techniques [35]. In one of the unexpected revelations, it was shown that the conductance of a low-density, weakly confined 1D quantum wire was quantized at a fractional filling of $2/5$ (in units of e^2/h) in the ground state [36]. Recently, a variety of fractional quantized states have been observed in high-quality, asymmetrically defined, low-density 1D regimes due to the highly correlated motion along the zigzag, which produces cyclic current as a means to lower the ground state energy [1,4,5,12,37,38]. Such developments highlight the importance of weakly confined regimes where the interplay between kinetic energy and electron-electron interactions can serve as a basis for exploring many-body physics, especially in the transition from 1D to quasi-1D or 2D regimes.

In this work, we investigated 1D quantum wires formed in high-quality GaAs/AlGaAs heterostructures via transport measurements. We show the emergence of the 1D Wigner state by fine-tuning the carrier density and confinement potential and investigate its evolution under magnetoelectric confinement.

II. EXPERIMENTAL METHODS

A. Sample details

Two samples, A and B, were fabricated from high-mobility 2DEGs formed in GaAs/AlGaAs heterostructures grown by molecular beam epitaxy. Sample A consists of a 30-nm-wide quantum well (QW) (hosting the 2DEG) located 90 nm beneath the surface, whereas sample B utilizes a 2DEG formed at the GaAs/AlGaAs interface, positioned 300 nm below the surface. At low temperature, sample A (B) exhibits an electron mobility $\mu = 2.38 \times 10^6 (3.5 \times 10^6)$ cm²/V s and a carrier density $n_{2D}^{\text{bulk}} = 2.24 \times 10^{11} (2.0 \times 10^{11})$ cm⁻² after illuminating with a red LED.

1D devices were defined using a multigate architecture comprising a pair of split gates and an overlapping top gate, electrically isolated by a cross-linked poly(methyl methacrylate) (PMMA) dielectric (upper inset, Fig. 1). This geometry enables independent control of confinement strength and carrier density, allowing access to the weak-confinement, low-density regime [3,5,6,12]. The lithographic channel length L and width W are $0.7 \mu\text{m}$ and $1 \mu\text{m}$ for sample A and $0.4 \mu\text{m}$ and $0.7 \mu\text{m}$ for sample B, respectively.

Transport measurements were performed in a cryogen-free dilution refrigerator with a base temperature of ~ 20 mK. The two-terminal differential conductance, $G = dI/dV_{SD}$, where V_{SD} is the applied dc source-drain bias, was measured using standard low-frequency lock-in techniques with an excitation voltage of $10 \mu\text{V}$ at 77 Hz. The observed effects were reproduced in several cooldowns. All measurements reported here correspond to the linear-response regime unless otherwise stated.

B. GaAs 2DEGs

The devices investigated here are realized in two different GaAs/AlGaAs heterostructure architectures, which lead to distinct vertical confinement potentials and electronic ground

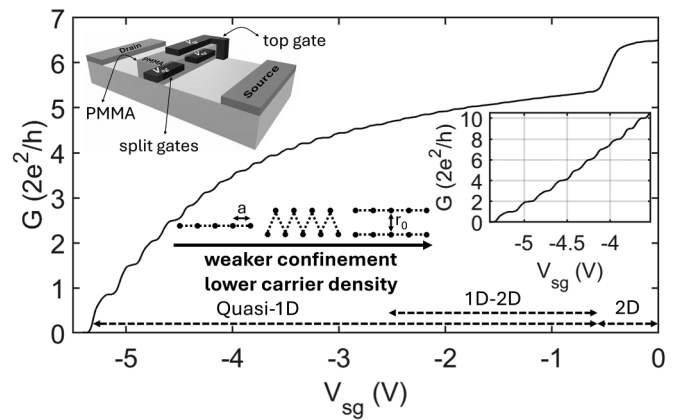


FIG. 1. Representative differential conductance G versus split-gate voltage V_{sg} plot for sample A at zero top-gate voltage V_{tg} in the absence of a magnetic field. Upper inset: schematic of the device showing a pair of split gates, top gate, source (S), and drain (D) for the two-terminal measurement. The main trace shows the as-measured data without removing the series resistance. Right inset: a magnified view of the main G trace, after correction for series resistance. Lower middle inset: illustration of the quasi-1D and the 1D-2D transition in the 1D channel, and the formation of the Wigner lattice, where a is the interparticle distance and r_0 is the transverse displacement due to Coulomb interactions exceeding the transverse confinement.

states. In the heterojunction case (sample B), the 2DEG forms at the GaAs/AlGaAs interface approximately 300 nm below the surface via modulation doping. The conduction-band offset together with the electrostatic field from remote donors produces a triangular quantum well at the interface, resulting in an asymmetric ground state envelope wave function that is localized on the GaAs side and decays into the bulk over a characteristic length scale of order 10–20 nm for typical densities [39,40]. In contrast, the QW structure consists of a 30 nm GaAs layer confined between AlGaAs barriers, with the 2DEG located approximately 105 nm below the surface. In this case the vertical confinement potential is closer to a finite square well, yielding a more symmetric ground state wave function centered within the GaAs layer and decaying into the surrounding barriers. Although the nominal well width exceeds the characteristic width of the triangular heterojunction state, the spatial extent of the electronic ground state in both systems remains comparable for similar carrier densities. The principal distinction therefore lies in the symmetry and location of the wave function, rather than its overall thickness.

The electrostatically defined quantum wires in the two heterostructures experience different confinement conditions arising from both the depth of the 2DEG and the lithographic gate geometry. In the QW device, the 2DEG resides relatively close to the surface (~ 105 nm), which allows strong electrostatic coupling to the gates. However, the split-gate opening is comparatively large ($\sim 700 \times 1000$ nm), producing a broad saddle-point potential and consequently a relatively small transverse confinement energy. The resulting 1D subbands are therefore closely spaced and the conductance plateaus $G = N(2e^2/h)$ may appear compressed in gate voltage and less sharply defined. In contrast, the heterojunction device has

a smaller lithographic constriction ($\sim 400 \times 700$ nm), which tends to enhance the transverse confinement and increase the subband spacing. Nevertheless, because the 2DEG lies deeper below the surface (~ 300 nm), the electrostatic potential from the gates spreads laterally before reaching the electron layer, resulting in smoother confinement than the lithographic dimensions alone would suggest. As a consequence, while well-defined 1D conductance plateaus can still form in both systems, their detailed evolution with gate voltage reflects the interplay between gate geometry and 2DEG depth. Importantly, the relatively soft confinement in these wide channels allows the electron system to expand laterally as the carrier density is reduced using the top gate. In this regime the kinetic energy decreases and electron-electron interactions become increasingly dominant, which can destabilize the 1D chain of electrons and promote the formation of zigzag or two-row charge arrangements. Such interaction-driven rearrangements could modify the conventional subband structure and are experimentally manifested as anomalies in the conductance plateaus, providing a platform for exploring many-body physics in low-density 1D systems [3,5,19,29,32,41].

III. RESULTS

A. Differential conductance

Differential conductance G of sample A as a function of split-gate voltage, V_{sg} , at zero top-gate voltage, V_{tg} , in the absence of a magnetic field is shown in Fig. 1, spanning the transition from the 2D regime to complete pinch-off of the channel. The split gates provide lateral confinement, transforming the 2DEG into a quasi-1D channel. Sweeping V_{sg} progressively negative narrows the quasi-1D channel, thereby controlling both the effective width and the number of occupied 1D subbands, resulting in the characteristic staircase of conductance quantized in units of $2e^2/h$. A negative V_{tg} further reduces the carrier density within the channel and modifies the overall confinement potential. To ensure accurate quantization, the series resistance arising from ohmic contacts and measurement leads has been subtracted in all plots, aligning the conductance plateaus with integer multiples of $2e^2/h$.

Figure 2 presents the results of G of sample A as a function of V_{sg} for selected top-gate voltages (≈ -2.40 to -3.12 V) and perpendicular magnetic field, $B_{\perp}$ (0.4–0.7 T). The data are shown both as conductance traces and as grayscale maps of the transconductance, dG/dV_{sg} . Figure 2(a) illustrates the evolution of G as the carrier density in the 1D channel is progressively reduced using V_{tg} . For V_{tg} between 0 and ~ -2.40 V, corresponding to strong and intermediate confinement regimes, respectively, well-defined integer plateaus at multiples of $2e^2/h$ are observed (not shown) [1,3,12,37]. As V_{tg} is made more negative from ≈ -2.40 V (intermediate confinement) to ≈ -3.12 V (weak confinement), the integer plateaus gradually weaken. In the regime of higher density (left side of the plot), the $2e^2/h$ plateau corresponds to the ground state, while a weaker plateau at $2(2e^2/h)$ corresponds to the first excited state. With decreasing V_{tg} (reducing confinement), the $2e^2/h$ plateau is progressively suppressed and eventually disappears, such that $2(2e^2/h)$ becomes the lowest observable

plateau. This transition occurs around $V_{tg} \approx -2.88$ V (V_{sg} near pinch-off ~ -2.40 V), as indicated by an arrow pointing to the dashed trace. Concurrently, the $2(2e^2/h)$ plateau becomes more pronounced as the $2e^2/h$ plateau weakens. Upon further reduction of V_{tg} (weak confinement), a shallow feature near $\sim 1.5(2e^2/h)$ develops alongside the $2(2e^2/h)$ plateau. At even lower density range, the $2(2e^2/h)$ plateau broadens and eventually weakens, accompanied by the reemergence of the $2e^2/h$ plateau and a strengthening of the $1.5(2e^2/h)$ -like feature.

It is observed that, for the most negative values of V_{tg} , the measured quantized conductance values are slightly lower than the expected values, even after removing the series resistance from ohmic contacts and measurement leads; however, this does not invalidate the identification of the underlying conducting modes. Such deviations are attributed to limitations of two-terminal measurements in the weak confinement regime, where strong electron-electron interactions can lead to nonuniversal conductance values [20,42].

The evolution of the 1D subbands is further visualized in the transconductance grayscale map shown in Fig. 2(f). Figure 2(k) reproduces the same data with dash-dotted red guide lines highlighting the evolution of the lowest two subbands. As seen in Fig. 2(f), the first excited state at $2(2e^2/h)$ remains relatively shallow compared to the ground and second excited states for $V_{tg} \gtrsim -2.80$ V, indicating weak coupling between them. As V_{tg} is made more negative and approaches ≈ -2.88 V, the $2e^2/h$ plateau is progressively suppressed and an anticrossing (AX) between the ground and first excited states becomes evident, as emphasized by the guide lines in Fig. 2(k) and shown with an arrow in red. The inset of Fig. 2(f) illustrates a typical avoided crossing between two energy levels, E_0 and E_1 .

The observed AX signals a reorganization of the 1D chain of electrons driven by interactions. In particular, it is consistent with the formation of a zigzag configuration, corresponding to a quasi-1D Wigner lattice, as expected in the regime of weak confinement and low density [1,3,12,18,32,37]. Although the two lowest modes approach closely, the finite residual gap reflects hybridization between them, characteristic of a finite system and indicative of a crossover rather than a phase transition [33]. Importantly, as V_{tg} is reduced further, the $2e^2/h$ plateau reemerges. This behavior indicates that the original first excited state evolves into the new ground state through the AX process, consistent with level repulsion and adiabatic continuity of the eigenstates [3,37,43–46].

Figures 2(b)–2(e) show the evolution of the conductance under $B_{\perp}$ in the range 0.4–0.7 T, in steps of 0.1 T. The immediate effect of applying $B_{\perp}$ is a systematic strengthening and widening of the quantized plateaus, as evident already at 0.4 T [Fig. 2(b)]. Concurrently, the AX observed at $V_{tg} \approx -2.88$ V in Fig. 2(a) is weakened, accompanied by the reemergence of the $2e^2/h$ plateau. At more negative V_{tg} (~ -3.12 V), an additional shallow structure develops near $0.5(2e^2/h)$.

The transconductance maps in Figs. 2(g) and 2(l) provide enhanced visibility of these features. In the range $V_{tg} \sim -2.40$ to -2.90 V, identified as the “interaction zone” before AX, the $2e^2/h$ plateau weakens and a sub- $2e^2/h$ structure emerges, evolving between $\sim 0.7(2e^2/h)$ and $0.5(2e^2/h)$. The interaction zone is highlighted by vertical dashed lines spanning $V_{sg} \sim -2.90$ to -2.40 V. As V_{tg} is reduced further

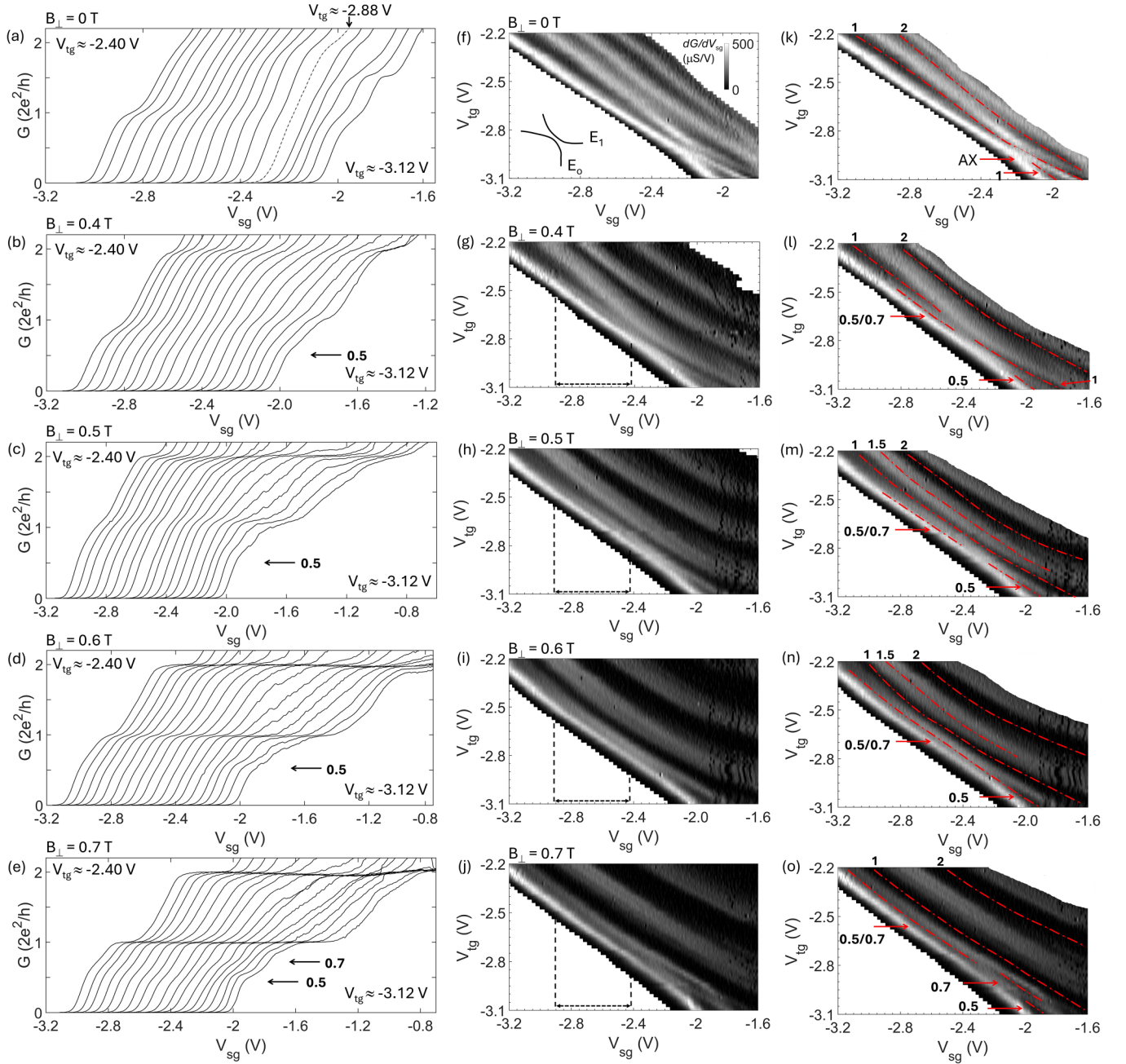


FIG. 2. G as a function of V_{sg} for sample A, measured at V_{tg} varying from approximately -2.40 V (left, intermediate confinement) to -3.12 V (right, weak confinement), at perpendicular magnetic field, $B_{\perp}$ (a) 0 T, (b) 0.4 T, (c) 0.5 T, (d) 0.6 T, and (e) 0.7 T. In (a), the dashed trace at $V_{tg} \approx -2.88$ V highlights the transition from pinch-off to the $2(2e^2/h)$ ground state, associated with the suppression of the $2e^2/h$ plateau. In (b)–(e), the progressive emergence of sub- $2e^2/h$ features is observed, including the development of $0.7(2e^2/h)$ and/or $0.5(2e^2/h)$ structures, as well as their evolution ($0.7 \rightarrow 0.5$) for $V_{tg} \lesssim -2.95$ V. Grayscale maps of the transconductance dG/dV_{sg} as a function of V_{sg} and V_{tg} are shown in (f)–(j), with corresponding indexed plots in (k)–(o) for (a)–(e), respectively. Bright (dark) regions correspond to conductance risers (plateaus). The inset in (f) schematically illustrates the anticrossing between the ground state E_0 and first excited state E_1 . In (f)–(j), black dashed lines and arrows delineate the “interaction zone” ($V_{sg} \sim -2.40$ to -2.90 V), where sub- $2e^2/h$ features $0.5(2e^2/h)$ or $0.7(2e^2/h)$ emerge. In (k)–(o), integer and fractional plateaus are highlighted by dash-dotted red lines. The integer and half plateaus, 1, 1.5, and $2(2e^2/h)$, are indexed on top of each plot. The sub- $2e^2/h$ structures are pointed out by red arrows and indexed next to them. They are also annotated as 0.5/0.7 to denote their continuous evolution between these values.

(< -2.80 V), the sub- $2e^2/h$ feature is suppressed near the center of the reduced AX, while the $2e^2/h$ plateau strengthens beyond it, together with the development of a more pronounced $0.5(2e^2/h)$ -like structure [Fig. 2(l)].

At $B_{\perp} = 0.5$ T [Figs. 2(c), 2(h), and 2(m)], the quantized plateaus become further stabilized and the anticrossing is strongly suppressed. The sub- $2e^2/h$ feature (0.7 or 0.5) in the “interaction zone” is enhanced, evolving toward a

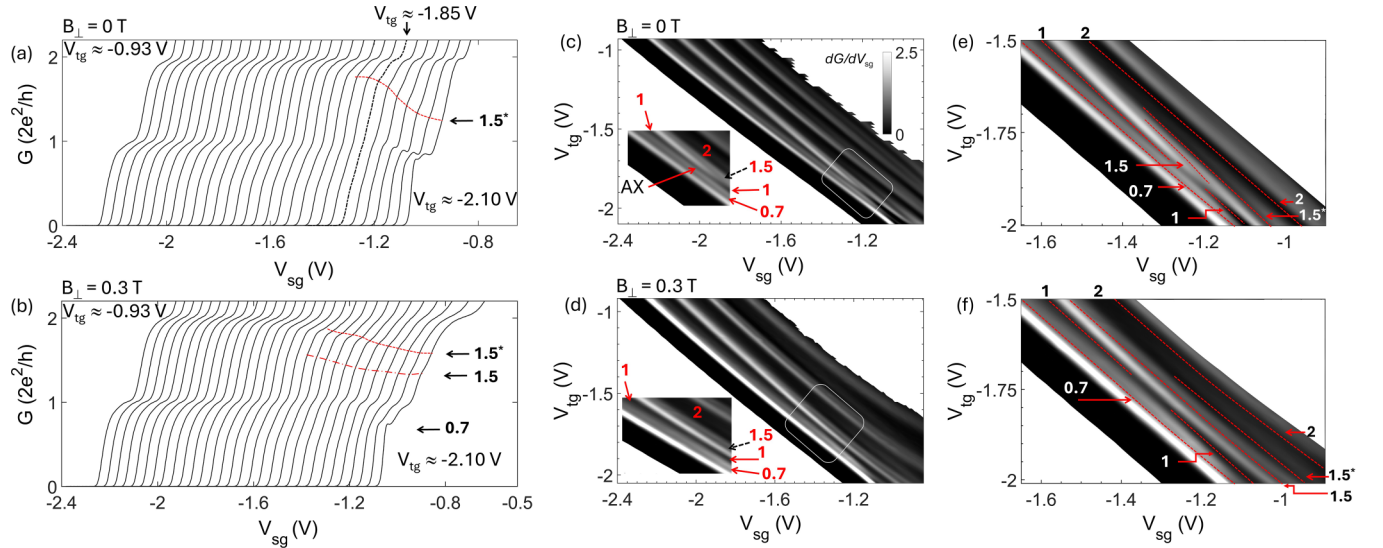


FIG. 3. G as a function of V_{sg} for sample B at V_{tg} ranging from approximately -0.93 V (left, intermediate confinement) to -2.10 V (right, weak confinement), measured (a) at $B_{\perp} = 0$ T and (b) at $B_{\perp} = 0.3$ T. In (a), the suppression of the $2e^2/h$ plateau at $V_{tg} \approx -1.85$ V is highlighted by the dash-dotted trace, marking the onset of an anticrossing regime. The red-dotted line guides the evolution of an additional feature emerging from the $2(2e^2/h)$ plateau, which evolves downward and stabilizes between $1.5(2e^2/h)$ and $2(2e^2/h)$; this feature is labeled 1.5^* indicated by a black arrow. In (b), the application of $B_{\perp} = 0.3$ T restores the $2e^2/h$ plateau and resolves a spin-split structure near $1.5(2e^2/h)$, while the 1.5^* feature persists between $1.5(2e^2/h)$ and $2(2e^2/h)$. The evolution of both features, 1.5^* and 1.5 , with decreasing V_{tg} is highlighted by dotted and dash-dotted red guide lines, respectively. Grayscale maps of the transconductance dG/dV_{sg} as a function of V_{sg} and V_{tg} are shown for (c) $B_{\perp} = 0$ T and (d) $B_{\perp} = 0.3$ T, with corresponding annotated plots in (e) and (f), respectively. Bright (dark) regions correspond to conductance risers (plateaus). In (c) and (d), the anticrossing between the $2e^2/h$ and $2(2e^2/h)$ states is highlighted by a white rectangle and magnified in the insets, where the plateaus (0.7, 1) are indicated by red arrows; the emergence of the 1.5 feature is shown with a dotted arrow in black. Panels (e) and (f) provide magnified views of the grayscale maps in (c) and (d), respectively, highlighting the detailed development of conductance structures with guided dotted-red lines. The grayscale bar represents the magnitude of dG/dV_{sg} .

well-defined $0.5(2e^2/h)$ structure for $V_{tg} < -2.95$ V. In addition, a weak feature near $1.5(2e^2/h)$ becomes discernible in the grayscale plots [Figs. 2(h) and 2(m)], despite remaining unresolved in the conductance traces, suggesting the emergence of exchange-correlation-induced spin splitting in the first excited state [47].

At $B_{\perp} = 0.6$ T [Figs. 2(d), 2(i), and 2(n)], AX is no longer observable and the conductance is dominated by well-defined plateaus. The sub- $2e^2/h$ structure now persists over the full range of V_{tg} , evolving continuously from a $0.7(2e^2/h)$ - or $0.5(2e^2/h)$ -like structure toward a robust $0.5(2e^2/h)$ plateau at a lower range of density (right of the plot). In addition, the spin-split feature near $1.5(2e^2/h)$ is resolved in the grayscale plots, reflecting the increasing role of Zeeman energy relative to the reduced exchange scale [41,47].

At the highest field, $B_{\perp} = 0.7$ T [Figs. 2(e), 2(j), and 2(o)], the evolution of the sub- $2e^2/h$ structure with V_{tg} becomes more distinct. In the higher-density regime (left of the plot), a $0.7(2e^2/h)$ -like feature is observed, which gradually evolves into a $0.5(2e^2/h)$ -like structure as the density is reduced within the “interaction zone”. Outside this zone ($V_{tg} \sim -2.90$ V, $V_{sg} \sim -2.50$ V), the $0.5(2e^2/h)$ feature weakens, while at further lower density range all three structures $0.7(2e^2/h)$, $0.5(2e^2/h)$, and $2e^2/h$ coexist within a single conductance trace, as highlighted in Figs. 2(j) and 2(o). This coexistence indicates the emergence of distinct, interaction-driven states as the system evolves from weak to strong correlation under combined electrostatic and magnetic confinement. Enhanced spin correlations and thermal

depopulation of spin subbands have been suggested to produce anomalous 0.5 and 0.7 conductance features [48,49].

It may, however, be noted that the higher subband feature $1.5(2e^2/h)$ is not as distinctly resolved as in Figs. 2(h) and 2(i). The weaker resolution of the $1.5(2e^2/h)$ feature at $B_{\perp} = 0.7$ T may reflect the fact that, in the low-density weak-confinement regime, a higher magnetic field not only increases Zeeman energy but also reorganizes the orbital spectrum. Once magnetoelectric confinement becomes dominant, the first excited subband can broaden, shift, or hybridize with nearby interaction-driven features, reducing the visibility of the spin-resolved 1.5 structure. At the same time, the suppressed exchange energy places the system close to the spin-correlated/spin-incoherent crossover [41], where spectral weight can transfer to the more robust 0.7- and 0.5-like features.

Figure 3(a) shows the conductance of sample B as a function of split-gate voltage for a range of top-gate voltages. Reducing the confinement by making V_{tg} more negative from approximately -0.93 V leads to a systematic weakening of the $2e^2/h$ plateau, which disappears at $V_{tg} \approx -1.85$ V, where the $2(2e^2/h)$ plateau becomes the new ground state (indicated by an arrow). Upon further reduction of V_{tg} , the $2e^2/h$ plateau reemerges together with the development of the $0.7(2e^2/h)$ feature in the weakest confinement regime ($V_{tg} \approx -2.10$ V). It may be noted that the $2(2e^2/h)$ plateau remains comparatively robust throughout the confinement tuning, from strong to weak, indicating that the ground state is governed by interaction effects rather than single-particle confinement.

In addition, a distinct structure emerges from the $2(2e^2/h)$ plateau [see the trace marked with $V_{tg} \approx -1.85$ V and its adjacent traces in Fig. 3(a)]. This feature evolves downward and stabilizes between $1.5(2e^2/h)$ and $2(2e^2/h)$, shown with a red-dotted line as a guide to the eye and labeled 1.5^* . The evolution of this and other features is more clearly resolved in the transconductance maps shown in Figs. 3(c) and 3(e). As the confinement is reduced, AX between the ground and first excited states develops near $V_{tg} \approx -1.85$ V, highlighted by a white rectangle and magnified in the inset of Fig. 3(c). Although only weakly visible in the conductance traces, a $0.7(2e^2/h)$ -like feature begins to develop for $V_{tg} \lesssim -1.2$ V and persists up to ~ -2.1 V. The inset further reveals a splitting of the riser between the first and second plateaus, shown by a black dotted arrow, suggestive of a $1.5(2e^2/h)$ -like feature due to interaction-driven spin splitting [47,49]. Meanwhile, the first excited state remains comparatively unaffected, while the 1.5^* feature becomes more pronounced for $V_{tg} < -1.85$ V.

Figure 3(b) shows the corresponding data in the presence of $B_{\perp} = 0.3$ T. In contrast to the zero-field case, the $2e^2/h$ plateau is restored, particularly in the weak-confinement regime. As V_{tg} is reduced, the first plateau weakens for the range of $V_{tg} \sim -1.70$ to -1.85 V before strengthening again at the lower-density regime ($V_{tg} \lesssim -1.85$ V), consistent with the behavior observed in sample A. The transconductance map in Fig. 3(d) shows that AX is strongly suppressed, while the $0.7(2e^2/h)$ feature remains visible within this regime.

Additionally, as observed in Figs. 3(c) and 3(e) at zero magnetic field, the two structures labeled 1.5^* and 1.5 for $V_{tg} \lesssim -1.5$ V become significantly enhanced in the presence of $B_{\perp}$ [Figs. 3(d) and 3(f)]. The 1.5 feature originates from the splitting of the riser between the first and second plateaus [inset of Fig. 3(d)] and evolves into a well-developed spin-resolved structure with decreasing V_{tg} , as highlighted by the red-dotted guide line in Fig. 3(f). The 1.5^* feature, however, remains located between $1.5(2e^2/h)$ and $2(2e^2/h)$. This structure is analogous to the so-called 0.7 analog observed in higher subbands [6,50,51] and is attributed to interaction-driven many-body effects. The simultaneous presence and distinct evolution of 1.5^* and 1.5 highlight the coexistence of interaction-driven and spin-resolved phenomena in this regime.

Comparing Figs. 2 and 3, several consistent features emerge between samples A and B: first, in both samples, the first conductance plateau weakens with decreasing V_{tg} , culminating in an AX between the ground and first excited states. Upon further reduction of V_{tg} , the $2e^2/h$ plateau reemerges, indicating a reordering of the lowest modes. Second, the application of a perpendicular magnetic field suppresses AX, consistent with enhanced transverse confinement and reduced hybridization. Third, within the AX regime, associated with the onset of a zigzag or quasi-1D Wigner configuration, both samples exhibit a sub- $2e^2/h$ structure, whose conductance evolves between $0.7(2e^2/h)$ and $0.5(2e^2/h)$. Fourth, in the weakest confinement regime at zero field, both samples display an additional $1.5(2e^2/h)$ -like structure (1.5^*) and the evolution of spin-split feature (1.5) in the first excited state under finite $B_{\perp}$. The reproducibility of these observations across two distinct device architectures indicates that the underlying

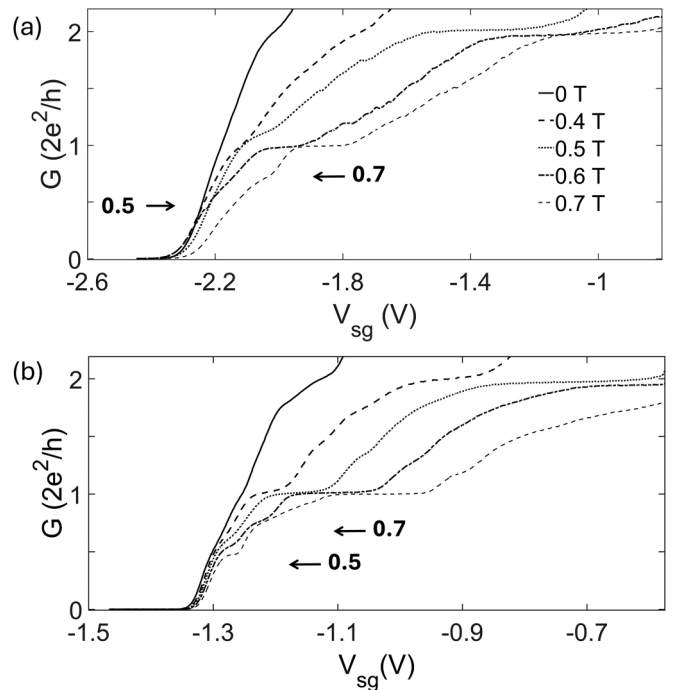


FIG. 4. Dependence of the anticrossing between the ground and first excited states on $B_{\perp}$, measured at $V_{tg} \approx -2.88$ V (sample A) and $V_{tg} \approx -1.85$ V (sample B). The plots show G as a function of V_{sg} for $B_{\perp}$ varying from 0 to 0.7 T for (a) sample A and (b) sample B. With increasing $B_{\perp}$, the anticrossing is progressively suppressed, consistent with reduced intersubband hybridization arising from enhanced magnetoelectric confinement. Concurrently, sub- $2e^2/h$ features evolve, with the emergence and strengthening of $0.7(2e^2/h)$ and $0.5(2e^2/h)$ structures in the low-density regime.

phenomena are robust and reflect universal interaction-driven physics in low-density quasi-1D systems, rather than sample-specific effects.

Figures 4(a) and 4(b) show the evolution of the conductance under $B_{\perp}$ in the range 0.4–0.7 T, measured at gate configurations corresponding to the AX regime identified in Figs. 2(a) and 3(a) for samples A and B, respectively. In the absence of a magnetic field, the $2e^2/h$ plateau is highly suppressed and the plateau at $2(2e^2/h)$ becomes the ground state. As $B_{\perp}$ increases, the previously suppressed $2e^2/h$ plateau progressively reemerges, while the $2(2e^2/h)$ plateau becomes more pronounced, consistent with the formation of magnetoelectric subbands and enhanced transverse confinement. In addition, sub- $2e^2/h$ features evolve systematically with increasing field. A $0.7(2e^2/h)$ -like structure develops in both samples, while a well-defined $0.5(2e^2/h)$ feature is more clearly resolved in sample B, indicating stronger interaction-driven spin effects at lower density regimes [4,5]. This behavior is consistent with Fig. 2, where in the ultra-low-density regime ($V_{tg} < -2.88$ V), pronounced $0.5(2e^2/h)$ and $0.7(2e^2/h)$ structures emerge in sample A for $0.5 \lesssim B_{\perp} \lesssim 0.7$ T. The qualitative agreement between the two samples demonstrates that the evolution of fractional conductance features [$0.7(2e^2/h)$ and $0.5(2e^2/h)$] with the perpendicular magnetic field is a robust characteristic of the low-density regime, despite differences in confinement and carrier den-

sity. A detailed discussion of these observations, including their connection to density estimates and interaction effects, is presented in Sec. III B (1D density estimation) and Sec. III C (influence of $B_{\perp}$).

B. Interaction-driven reconfiguration

The suppression of the $2e^2/h$ conductance plateau and the emergence of $2(2e^2/h)$ as the lowest plateau (new ground state) in both samples indicate that the conventional single-mode ground state is no longer realized. Instead, this behavior is consistent with an interaction-driven restructuring of the electron system into a quasi-1D configuration comprising two coupled rows, commonly described as a zigzag or two-row Wigner lattice, illustrated in the lower middle inset in Fig. 1 [27,29,52]. In this regime, it is essential to determine the local carrier density within the constriction, the corresponding 1D density, and the relative strength of Coulomb interactions and transverse confinement.

1. Local density within the 1D constriction

To account for the influence of both gates (split gates and the top gate), we introduce an effective gate voltage

$$V_{\text{eff}} = V_{\text{tg}} + \beta V_{\text{sg}}, \quad (1)$$

where $\beta \approx 0.6$ is the relative split-gate lever arm, extracted from the slope of the pinch-off line in the transconductance maps (Figs. 2 and 3), reflecting the relative capacitive coupling of the split gate and top gate to the 1D channel [53]. Accordingly, the pinch-off condition is defined by

$$V_{\text{eff}}^{(p)} = V_{\text{tg}} + \beta V_{\text{sg}}^{(p)}. \quad (2)$$

The local density in the constriction is then estimated as

$$n_{2\text{D}}^{\text{local}} \approx n_{2\text{D}}^{\text{bulk}} \left(1 - \frac{|V_{\text{eff}}|}{|V_{\text{eff}}^{(p)}|} \right), \quad (3)$$

which captures the leading order electrostatic depletion of the 2DEG [40,54]. We extract the gate voltages for both samples where the first plateau appeared at $2(2e^2/h)$. For sample A, $V_{\text{tg}} \approx -2.88$ V, $V_{\text{sg}}^{(p)} \approx -2.25$ V (pinch-off voltage), and $V_{\text{sg}} \approx -2.10$ V [voltage at the center of the $2(2e^2/h)$ plateau], we obtain $V_{\text{eff}}^{(p)} \approx -4.23$ V and $V_{\text{eff}} \approx -4.14$ V, yielding

$$n_{2\text{D}}^{\text{local(A)}} \sim 4.7 \times 10^9 \text{ cm}^{-2}. \quad (4)$$

For sample B ($V_{\text{tg}} \approx -1.85$ V, $V_{\text{sg}}^{(p)} \approx -1.35$ V, and $V_{\text{sg}} \approx -1.15$ V), we obtain $V_{\text{eff}}^{(p)} \approx -2.66$ V and $V_{\text{eff}} \approx -2.54$ V, yielding

$$n_{2\text{D}}^{\text{local(B)}} \sim 9.0 \times 10^9 \text{ cm}^{-2}. \quad (5)$$

The $n_{2\text{D}}^{\text{local}}$ estimation suggests that both samples may have been driven into an ultra-low-density regime within the constriction. We will now estimate the effective linear density within the 1D channel.

2. Effective width and 1D density

In the reconfigured regime where the $2e^2/h$ plateau is substantially weakened, the transverse extent is governed by the

hybridized two-row wave function. Therefore, the effective channel width may be written as

$$W_{\text{eff}} \sim \lambda_F, \quad \lambda_F = \sqrt{\frac{2\pi}{n_{2\text{D}}^{\text{local}}}}. \quad (6)$$

The 1D density may be defined as

$$n_{1\text{D}} \approx n_{2\text{D}}^{\text{local}} W_{\text{eff}}, \quad (7)$$

giving

$$n_{1\text{D}} \sim 1.7 \times 10^5 \text{ cm}^{-1} \text{ (sample A),}$$

$$n_{1\text{D}} \sim 2.4 \times 10^5 \text{ cm}^{-1} \text{ (sample B).} \quad (8)$$

3. Interaction strength

The mean interparticle spacing in a 1D chain is defined as

$$a = \frac{1}{2n_{1\text{D}}} \quad (9)$$

and using Eq. (8) we estimate

$$a \sim 29 \text{ nm (sample A),} \quad a \sim 21 \text{ nm (sample B).} \quad (10)$$

Along the 1D wire, the interaction strength may be characterized by the dimensionless Wigner-Seitz parameter [29,55–58],

$$r_s^{\text{1D}} = \frac{a}{a_B^*}, \quad (11)$$

where $a_B^* = 10$ nm is the effective Bohr radius for GaAs, yielding

$$r_s^{\text{1D(A)}} \sim 2.9 \text{ (sample A),} \quad r_s^{\text{1D(B)}} \sim 2.1 \text{ (sample B).} \quad (12)$$

Thus both samples satisfy the Wigner lattice condition, $r_s^{\text{1D}} \gtrsim 1$, in which Coulomb and exchange interactions qualitatively modify the low-energy state [25,27,29,59].

4. Zigzag or two rows formation

The transverse confinement introduces a characteristic length scale, which describes the spatial extent of the lowest transverse mode defined as

$$\ell_{\perp} = \sqrt{\frac{\hbar}{m^* \omega}}, \quad (13)$$

where $m^* = 0.067m$; $m^*(m)$ is the effective (rest) mass of the electron and $\hbar$ is the reduced Planck constant. Using $\hbar\omega \sim 0.5$ meV and $\omega \sim 0.76 \times 10^{12} \text{ s}^{-1}$ we obtain

$$\ell_{\perp} \sim 48 \text{ nm}, \quad (14)$$

while the Coulomb energy at the mean interparticle spacing, using Eqs. (9) and (10), is

$$E_C(a) = \frac{e^2}{4\pi\epsilon\epsilon_0 a} \quad (15)$$

and using $\epsilon = 12.9$, the dielectric constant for GaAs, we get

$$E_C(a) \sim 3.85 \text{ meV (sample A);} \quad 5.31 \text{ meV (sample B).} \quad (16)$$

To quantify the onset of transverse displacement of electrons, we consider the balance between Coulomb and

confinement energies [27–29],

$$E_C(r_0) = \frac{e^2}{4\pi\epsilon_0 r_0}, \quad V_{\text{conf}} = \frac{1}{2}m^*\omega^2 r_0^2, \quad (17)$$

which yields the characteristic transverse displacement

$$r_0 \approx \left(\frac{2e^2}{4\pi\epsilon_0 m^*\omega^2} \right)^{1/3} \sim 100 \text{ nm}, \quad (18)$$

$E_C(r_0) \sim 1.11 \text{ meV}$ [$\ll E_C(a)$], and $V_{\text{conf}} \sim 1.10 \text{ meV}$, demonstrating that Coulomb interactions dominate even at zigzag or two row formation. The hierarchy of length scales is therefore $a \lesssim \ell_\perp \lesssim r_0$, while $E_C \gg \hbar\omega$, which is the defining condition for transverse instability of a 1D chain. In this regime, electrons minimize their mutual repulsion by developing a finite transverse separation r_0 relative to both $\ell_\perp$ and a , indicating that electrons are no longer confined to a single 1D chain. This provides evidence that the observed $2(2e^2/h)$ ground state arises from an interaction-driven reconfiguration into zigzag or two rows of electrons [3,19,32].

Theoretically, it has been shown that a strictly 1D electron chain becomes unstable toward a zigzag configuration as the confinement is relaxed [27,29,41,52]. In finite gate-defined systems, this instability is rounded into a smooth crossover, producing an avoided crossing rather than a true level crossing. Weakening the confinement brings a competing low-energy configuration close to the lowest subband and finite hybridization between these configurations results in level repulsion and an adiabatic interchange of character, consistent with the generic behavior of interacting quantum systems [60,61]. Similar interaction-driven deviations from the conventional subband picture have been reported in weakly confined quantum wires [3,19,32].

Although the energy separation between the two lowest levels is substantially reduced near the avoided crossing, it remains finite throughout the accessible parameter range, indicating that the system does not undergo a true phase transition. Instead, the observed behavior is consistent with a finite-size manifestation of the interaction-driven zigzag transition, which in the thermodynamic limit has been modeled as a second-order quantum phase transition in the Ising universality class [33]. The evolution of the subbands is consistent with the adiabatic evolution of the instantaneous eigenstates of an interacting two-level system under quasistatic gate tuning. Consequently, the Landau-Zener formalism, which describes nonadiabatic transitions driven through an avoided crossing by time-dependent perturbations [43,44], does not appear to be directly applicable to the equilibrium spectra measured here. Nevertheless, the observed anticrossing gap provides the characteristic energy scale governing such dynamics.

C. Effect of perpendicular magnetic field: Orbital compression and spin correlations

1. Sample A

We investigate the evolution of conductance under the combined effects of electrostatic and magnetic confinement in a perpendicular magnetic field ($B_\perp \sim 0.4\text{--}0.7 \text{ T}$). At zero field, the suppressed $2e^2/h$ plateau and the resultant anticrossing between the lowest subbands indicate a restructured

zigzag or two-row ground state. With increasing $B_\perp$, the anticrossing is weakened and the $2e^2/h$ plateau is restored, while higher plateaus sharpen. Concurrently, sub- $2e^2/h$ features near $0.7(2e^2/h)$ and/or $0.5(2e^2/h)$ strengthen in the low-density regime. This evolution reflects magnetic confinement suppressing inter-row hybridization while enhancing interaction-driven spin effects.

The underlying orbital mechanism can be understood in terms of the magnetic length,

$$\ell_{B_\perp} = \sqrt{\frac{\hbar}{eB_\perp}}, \quad (19)$$

which sets the characteristic transverse extent of the orbital wave function in a perpendicular field. For $B_\perp \sim 0.4 \rightarrow 0.7 \text{ T}$,

$$\ell_{B_\perp} \sim 41 \rightarrow 31 \text{ nm}, \quad (20)$$

comparable to both the zero-field transverse confinement length $\ell_\perp \sim 48 \text{ nm}$ [Eq. (14)] and the mean interparticle spacing $a \sim 20\text{--}30 \text{ nm}$ [Eq. (10)]. When $\ell_{B_\perp} \lesssim \ell_\perp$, the perpendicular magnetic field compresses the transverse wave function and suppresses lateral fluctuations, thereby quenching the zigzag instability. Importantly, this orbital compression does not remove interaction effects. Rather, by restricting transverse motion and reducing the available orbital phase space, the field enhances the relative importance of Coulomb correlations within the quasi-1D channel. Thus the perpendicular field suppresses the transverse two-row reconfiguration while preserving, and in some respects amplifying, the correlated nature of the low-density electron system.

A complementary picture is provided by the cyclotron radius,

$$r_c = \frac{\hbar k_F}{eB_\perp}, \quad (21)$$

where the Fermi wave vector $k_F = n_{1D}\pi/2$. Although a quantum wire does not support complete cyclotron orbits, r_c provides an estimate of the spatial extent over which magnetic confinement influences electron motion. Using the experimentally relevant densities (e.g., for sample A) from Eq. (8), we obtain

$$r_c \sim 44 \rightarrow 25 \text{ nm}, \quad (22)$$

for $B_\perp \sim 0.4 \rightarrow 0.7 \text{ T}$, which is comparable to both $\ell_\perp$ and a . As $B_\perp$ increases, r_c decreases and the electron motion becomes confined on a spatial scale comparable to the mean electron separation. This has two direct consequences: first, the orbital trajectories can no longer sustain the laterally extended two-row arrangement required for the zigzag state, favoring restoration of a single mode and hence the reappearance of the $2e^2/h$ plateau. Second, the reduced spatial extent of motion suppresses wave function overlap between neighboring electrons, thereby lowering the exchange coupling and potentially enhancing spin-correlated effects [7,18,19,29,41,58].

The relevant energy scales further support this picture. The cyclotron energy

$$\hbar\omega_c = \frac{\hbar e B_\perp}{m^*} \sim 0.9 \rightarrow 1.2 \text{ meV}, \quad (23)$$

for $B_{\perp} \sim 0.5 \rightarrow 0.7$ T. The cyclotron energy $\hbar\omega_c > \hbar\omega \sim 0.5$ meV, which confirms that the field strongly reorganizes the orbital spectrum into magnetoelectric subbands [62,63]. At the same time, Zeeman energy

$$E_Z = g^* \mu_B B_{\perp} \sim 13 \rightarrow 18 \text{ } \mu\text{eV}, \quad (24)$$

for $B_{\perp} \sim 0.5 \rightarrow 0.7$ T, where $|g|^* \approx 0.44$ is the effective Landé g -factor for GaAs and μ_B is the Bohr magneton. E_Z is small compared with the subband spacing but comparable to the low-energy scale of the dilute ground state. This is sufficient to make spin splitting visible in transconductance, particularly in higher subbands where splitting of the riser is observed at $B_{\perp} \geq 0.5$ T [47,49].

A quantitative description of the spin sector is obtained from the exchange coupling J between neighboring electrons. In a low-density 1D system approaching the Wigner crystal regime, the exchange interaction is exponentially suppressed with increasing interparticle spacing [29,41,52,64],

$$J \sim E_F \exp(-\eta \sqrt{r_s^{\text{1D}}}), \quad (25)$$

where $\eta \sim 1$ –2 is a constant that depends weakly on the confinement potential. Using $r_s^{\text{1D}} \sim 2.1$ –2.9 [Eq. (12)] and $E_F \sim 0.05$ –0.25 meV (representing the occupation of the lowest subband), we estimate

$$J \sim 2\text{--}60 \text{ } \mu\text{eV}, \quad (26)$$

corresponding to a temperature scale $J/k_B \sim 20$ –700 mK. This energy scale is comparable to both the experimental temperature (~ 20 mK) and the Zeeman energy at $B_{\perp} \sim 0.5$ –0.7 T [Eq. (24)]. Thus the system lies close to the crossover regime $J \sim k_B T$, in which spin correlations are only partially developed. For $J \gtrsim k_B T$, the system retains significant spin correlations, which can give rise to conductance anomalies such as the 0.7 structure [7,17,19,41,58,65,66]. As J is further suppressed and approaches $J \lesssim k_B T$, the system crosses over to a spin-incoherent regime in which spin excitations are thermally randomized while charge modes remain coherent. In this limit, the conductance deviates from the noninteracting value and can be reduced toward $0.5(2e^2/h)$ [7,18,58].

The suppression of J originates from the increasing spatial separation of electrons driven by Coulomb interactions. This same mechanism underlies the transverse instability of the 1D chain and the formation of a zigzag or two-row Wigner configuration [27,29,52]. The evolution from weak $0.7(2e^2/h)$ - or $0.5(2e^2/h)$ -like features to a robust $0.5(2e^2/h)$ ground state can therefore be understood as a direct consequence of the continuous reduction of J with increasing r_s^{1D} , linking spin correlations, spin incoherence, and spatial reorganization within a unified interaction-driven framework. This connection between anomalous sub- $2e^2/h$ conductance features and incipient electron-lattice formation is consistent with earlier studies of weakly confined quantum wires, where strong interactions and suppressed exchange were primarily associated with the emergence of a $0.5(2e^2/h)$ feature in comparable low-density regimes [7,18,19]. In the present work, however, we demonstrate the magnetic-field-assisted

coexistence of 0.5, 0.7, and 1 structures within a similar density and confinement regime. The sub- $(2e^2/h)$ features have also been suggested to arise due to the localized charge density waves formed in the strongly interacting 1D Wigner lattice regime, thus serving as a macroscopic signature of a 1D Wigner phase [30].

2. Sample B

Sample B exhibits $2(2e^2/h)$ ground state at zero field, indicating a spatial reconfiguration of 1D electrons in a zigzag or two rows in a strongly interacting regime with $r_s^{\text{1D}} \sim 2.1$. In this regime, a shallow $0.7(2e^2/h)$ -like feature is observed together with a $1.5(2e^2/h)$ -like structure (labeled 1.5^*), consistent with a higher-subband analog of the $0.7(2e^2/h)$ anomaly [6,17,50,51,66].

Upon applying $B_{\perp} = 0.3$ T, a qualitative reorganization of the spectrum occurs. The previously missing $2e^2/h$ plateau reappears, while the $0.7(2e^2/h)$ feature strengthens at lower density. In the low-density regime, the zero-field 1.5^* structure shifts upward toward $\sim 1.7(2e^2/h)$, while a new magnetic-field-resolved structure emerges near $1.5(2e^2/h)$ (labeled 1.5) in the transconductance map. The persistence of the 1.5^* structure, already present at zero field, indicates its many-body origin as a higher-subband $0.7(2e^2/h)$ [51], while the emergent 1.5 feature reflects Zeeman-driven spin correlation splitting [49].

This behavior can be understood from the combined effects of orbital compression and exchange suppression. The magnetic length, $\ell_{B_{\perp}} \sim 47$ nm, becomes comparable to the confinement length, $\ell_{\perp} \sim 48$ nm, causing the magnetic confinement to compete directly with the electrostatic confinement [16,62]. As a result, transverse fluctuations are suppressed, destabilizing the zigzag configuration and promoting the recovery of a single-row channel. Simultaneously, the low-density regime is characterized by a strongly reduced exchange coupling, $J \sim 2$ –60 μeV [52,58]. Since this energy scale is comparable to both the experimental temperature ($J \sim k_B T$) and the Zeeman energy ($E_Z \sim 8 \text{ } \mu\text{eV}$), even a modest perpendicular magnetic field can significantly influence the spin sector, enhancing spin-related conductance anomalies such as the $0.7(2e^2/h)$ feature and its higher-subband analog [36,47,49,67].

The splitting of the first excited subband therefore reflects the coexistence of conventional spin splitting and a many-body-derived $0.7(2e^2/h)$ analog (1.5^*) inherited from the nearby zigzag regime [3,36]. These results demonstrate that a perpendicular magnetic field suppresses transverse charge reconstruction while preserving, and enhancing, interaction-driven spin correlations across both ground and excited subbands in a low-density quantum wire.

Results on both samples establish a unified picture in which magnetic confinement suppresses inter-row hybridization while driving the system deeper into an interaction-dominated regime. The resulting transport behavior is governed by the interplay between Coulomb interactions, exponentially suppressed exchange coupling, and weak Zeeman splitting, linking the disappearance of the anticrossing to the emergence of spin-related conductance anomalies.

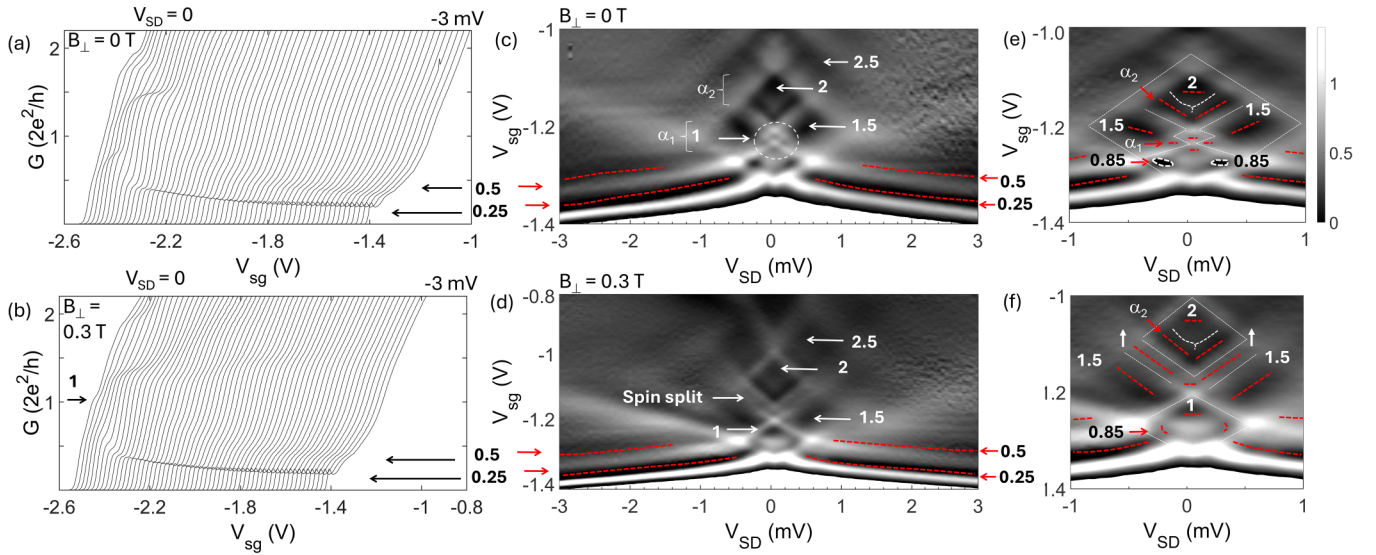


FIG. 5. Dc-bias spectroscopy of sample B measured at $V_{\text{tg}} \approx -1.85$ V. (a),(b) G as a function of V_{sg} for dc source-drain bias V_{SD} ranging from 0 to -3 mV at (a) $B_{\perp} = 0$ T and (b) $B_{\perp} = 0.3$ T. Successive traces are offset horizontally for clarity, with the leftmost (rightmost) trace corresponding to $V_{\text{SD}} = 0$ V (-3 mV) and increments of 0.5 mV. At large bias, the emergence of $0.5(2e^2/h)$ and $0.25(2e^2/h)$ features is observed and indicated. Grayscale maps of transconductance (dG/dV_{sg}) as a function of V_{sg} and V_{SD} are shown in (c) and (d), with corresponding annotated plots in (e) and (f) for (a) and (b), respectively. Bright (dark) regions correspond to conductance risers (plateaus). In (c), a broad diamond centered at zero bias spans the $2(2e^2/h)$ state and a weak structure near $2e^2/h$, indicating hybridization of the lowest modes into bonding and antibonding states. The weak $2e^2/h$ -like feature is highlighted by a white-dotted circle, while integer and half-integer plateaus are marked by white arrows. The development of $0.5(2e^2/h)$ - and $0.25(2e^2/h)$ -like features at higher bias is indicated by red-dotted guide lines. A magnified view in (e) shows the broad diamond, annotated with a white-dotted border; within it, the 2 and structures below it, α_2 , and additional low-energy structures such as α_1 are indexed and highlighted by red-dotted lines; both α_1 and α_2 are also highlighted in (c). Also, the appearance of $0.85(2e^2/h)$ features at finite bias is marked by white-dotted ovals. In (d), the application of $B_{\perp} = 0.3$ T leads to the formation of well-defined diamonds corresponding to both $2e^2/h$ and $2(2e^2/h)$, while the $0.5(2e^2/h)$ - and $0.25(2e^2/h)$ -like features persist. The arms of the V-shaped diamond associated with the second subband exhibit splitting, consistent with enhanced spin resolution. The magnified view in (f) shows that the $0.85(2e^2/h)$ features are significantly suppressed, while the split diamond arms are indicated by white arrows pointing up, highlighting their spin origin; adjacent α_2 features may also be seen. The grayscale bar represents the magnitude of dG/dV_{sg} .

D. Source-drain bias spectroscopy

The differential conductance of sample B as a function of applied dc source-drain bias V_{SD} , in the confinement regime where the anticrossing is observed corresponding to the formation of a zigzag or double-row configuration, is shown in Fig. 5(a) for $B_{\perp} = 0$ T and Fig. 5(b) for $B_{\perp} = 0.3$ T. The corresponding transconductance maps are shown in Figs. 5(c) and 5(d). In general, in a noninteracting 1D channel at sufficiently small bias, the conductance is quantized in integer multiples of $2e^2/h$. When V_{SD} becomes comparable to the subband spacing, the source and drain chemical potentials probe different regions of the subband dispersion, lifting the momentum degeneracy and giving rise to half-integer plateaus at $1.5(2e^2/h)$ and $2.5(2e^2/h)$ on either side of the zero bias [13,54,68,69].

However, the ground-state evolution deviates from this simple picture. Instead of a symmetric half-integer counterpart, a feature near $0.85(2e^2/h)$ first emerges, followed by the development of structures consecutively near $0.5(2e^2/h)$ and $0.25(2e^2/h)$ with increasing bias. These features are attributed to the progressive lifting of spin degeneracy, from partial to near-complete spin polarization in the lowest subband [3,6,32,69,70]. In the present case of interacting 1D electrons, at large bias ($V_{\text{SD}} \sim \pm 3$ mV), a well-defined plateau near

$0.25(2e^2/h)$ is accompanied by an additional feature close to $0.5(2e^2/h)$ [Figs. 5(a) and 5(c)]. This dual conductance structure is consistent with the presence of a zigzag or two-row configuration, in which each row contributes approximately $0.25(2e^2/h)$ under strong bias [1,3,6,19].

In the grayscale plot in Fig. 5(c) and a magnified version in Fig. 5(e), the discussed features and additional fine structures are better resolved. A broad diamond-shaped feature spans both the ground and first-excited states, indicative of hybridized (bonding and antibonding) states arising from the avoided crossing at zero bias. Within this diamond, a circled region (denoted “1”), corresponding to a weak remnant of the $2e^2/h$ plateau at $V_{\text{SD}} = 0$, together with a darker region (denoted “2”) near $2(2e^2/h)$, defines the hybridized states. The circled region exhibits two vertical features (vertical- α_1) at zero bias and corresponding horizontal features (horizontal- α_1), at finite bias ($V_{\text{SD}} \sim \pm 150$ μ V). As the magnitude of V_{SD} increases, these features broaden; a $1.5(2e^2/h)$ -like plateau at higher bias ($\sim \pm 500$ μ V) is then resolved, resulting in the estimation of the subband spacing of ≤ 500 μ eV. We also observe the evolution of the $0.85(2e^2/h)$ features at finite bias, shown in white-dotted ovals [Fig. 5(e)].

Similarly, the region associated with the $2(2e^2/h)$ state (“2”) exhibits an additional structure just below it at zero

bias, together with closely spaced features extending to finite bias ($\sim \pm 200 \mu\text{V}$), shown as α_2 . The presence of low-energy features (α_1 and α_2), with characteristic energy scales below $250 \mu\text{eV}$, is most likely evidence for bonding and antibonding states in the Wigner regime [19,32,45,46,71].

Figure 5(b) shows the counterpart of Fig. 5(a) in the presence of $B_\perp = 0.3 \text{ T}$, with the corresponding transconductance maps presented in Figs. 5(d) and 5(f). The application of $B_\perp$ restores the $2e^2/h$ plateau at zero bias, which is unresolved in the zero-field case, indicating suppression of the transverse zigzag or double-row configuration. At the same time, the dual conductance features near $0.25(2e^2/h)$ and $0.5(2e^2/h)$ at large source-drain bias persist, although the overall features become sharper and more well defined in the transconductance maps.

A comparison with the zero-field data reveals a significant reorganization of the bias-dependent low-energy structures. The horizontal- α_1 features at finite V_{SD} , previously observed in the circled region of Fig. 5(c), are largely suppressed, while the vertical- α_1 features become more pronounced. In particular, the lower vertical- α_1 feature evolves into a well-defined $2e^2/h$ plateau (“1”), whereas the upper- α_1 feature, together with the $2(2e^2/h)$ state (“2”), forms a distinct conductance diamond, giving rise to a characteristic V-shaped structure. The lower branches of this diamond exhibit a clear splitting as V_{SD} is varied from 0 to $\pm 500 \mu\text{V}$; the α_2 features may also be seen above the spin-split branches. This splitting is consistent with Zeeman-assisted spin resolution in the presence of interactions [Figs. 5(d) and 5(f)] [3,72].

In contrast, the overall structure of the $2(2e^2/h)$ state and its surrounding features, α_2 , remain largely unchanged compared to the zero-field case, indicating that the excited-state spectrum retains signatures of the underlying interaction-driven reconstruction. In the ground-state sector, the $2e^2/h$ plateau is significantly strengthened, while the intermediate $0.85(2e^2/h)$ feature is strongly suppressed and nearly absent. Importantly, the $2e^2/h$ (“1”) and $2(2e^2/h)$ (“2”) states now form separate, well-defined conductance diamonds. This contrasts with the zero-field case, where a single, broader diamond encompasses both states, reflecting hybridization between bonding and antibonding modes.

In low-density quasi-1D systems, disorder can in principle couple strongly to such correlated states by pinning the electron lattice and suppressing long-range correlations [29,31,52,58]. However, disorder-dominated transport is typically characterized by resonant features, irregular plateau structure, and nonreproducible behavior [6,73,74]. In contrast, the present measurements exhibit smooth and systematic evolution across conductance, transconductance, and finite-bias spectroscopy, indicating that the observed features arise from interaction-driven reorganization rather than disorder-induced effects.

In the interaction-driven framework mentioned earlier, the key control parameter $r_s^{\text{1D}} > 2$ indicates a regime where the 1D electron chain becomes unstable to transverse distortions, resulting in a zigzag (two-row) pattern and hybridization of the lowest energy modes. In cleaner systems, such as higher-mobility heterostructures [75,76], this regime is likely

to exhibit a broader range of phenomena, including well-defined collective transverse modes, larger hybridization gaps, and complex spin ground states influenced by multiparticle interactions, such as ring exchange [27,29,52]. Improved mobility mainly reduces disorder effects, sharpening these features. Ultraclean systems could thus offer insights into more delicate aspects of quasi-1D Wigner physics, such as detailed structural transitions, collective excitations, and the evolution of the spin phase diagram.

E. Discussion on dc-bias spectroscopy of a two-row state

The dc-bias spectroscopy provides direct insight into the internal structure of the reconstructed state beyond what can be inferred from linear-response conductance. Under a finite dc bias, the source and drain chemical potentials differ by eV_{SD} ,

$$\mu_S = E_F + \frac{eV_{SD}}{2}, \quad \mu_D = E_F - \frac{eV_{SD}}{2}. \quad (27)$$

In the linear-response limit ($V_{SD} \rightarrow 0$), conductance plateaus appear whenever a subband bottom crosses the Fermi level, yielding integer multiples of $2e^2/h$ for spin-degenerate modes [15,16]. At finite bias, however, the presence of two distinct chemical potentials allows subband edges to lie within the bias window, producing half-integer conductance features when $eV_{SD} \gtrsim \Delta E$ [3,5,54,68,69]. Here ΔE denotes the spacing between consecutive 1D subbands, set by the transverse confinement potential ($\hbar\omega \sim 0.5 \text{ meV}$).

In the weak-confinement, low-density regime, Coulomb interactions drive the electron system into a configuration comprising two laterally displaced rows. When these rows are sufficiently close, finite overlap of their wave functions allows tunneling between them, resulting in hybridization of the two row states into bonding and antibonding combinations [19,32,71]. The energy separation between these hybridized states may be written as [45,46]

$$\Delta = \sqrt{\varepsilon^2 + 4t^2}, \quad (28)$$

where ε represents the energy imbalance between the two rows (set by confinement asymmetry or electrostatics) and t is the inter-row tunneling amplitude. This relation has two important implications: (1) when tunneling dominates ($t \gg \varepsilon$), the two rows are strongly hybridized, producing symmetric and antisymmetric states with a splitting $\Delta \simeq 2t$; (2) in contrast, when the detuning dominates ($\varepsilon \gg t$), the rows behave as nearly independent modes and the hybridization is weak. Thus the strength of hybridization is controlled by the competition between tunneling and electrostatic imbalance, both of which depend sensitively on density and confinement.

In dc-bias spectroscopy, this hybridization manifests as a broad conductance diamond spanning both the ground (“1”) and first excited (“2”) states, as shown in Fig. 5(e), indicating coupled modes rather than independent subbands. In addition,

the finite splitting Δ introduces an internal energy scale that is much smaller than the single-particle subband spacing. This gives rise to additional low-energy features, typically below $|V_{SD}| \lesssim 250 \mu\text{eV}$, represented as α_1 and α_2 within the broader diamond structure. These features correspond to the onset of transport through the bonding and antibonding states as they enter the bias window [45,46]. Accordingly, the internal energy scale (Δ) may be estimated,

$$\Delta \approx \beta e|V_{SD}^*| \sim (75-100) \mu\text{eV} \\ \text{for } \beta \sim 0.5, |V_{SD}^*| \sim 150-200 \mu\text{V}. \quad (29)$$

For near-zero detuning ($\varepsilon \ll 2t$), Eq. (28) then gives

$$t \sim \Delta/2 \sim (30-50) \mu\text{eV}. \quad (30)$$

This magnitude is consistent with the experimental observation that coupled-row states form avoided crossings and hybridized spectra in the weak confinement.

The spin sector introduces a second, much smaller energy scale. In the zigzag or two-row regime, the exchange coupling is exponentially suppressed, $J \ll \hbar\omega < E_C$, with the experimentally estimated hierarchy $E_C \sim 5.3 \text{ meV}$, $\hbar\omega \sim 0.5 \text{ meV}$, and $J \sim 2-60 \mu\text{eV}$. This places the system in a regime where charge reconstruction and hybridization are controlled by Coulomb energy and tunneling, while the spin sector is soft and easily perturbed by dc bias or magnetic field [3,52]. Consequently, the small-bias evolution through $0.85(2e^2/h)$ and the appearance of $0.5(2e^2/h)$ - and $0.25(2e^2/h)$ -like structures at larger bias can be understood as spin-related nonequilibrium features developing on top of the hybridized two-row background, rather than as simple single-particle half plateaus. These structures are thought to be emergent features of bonding and antibonding states when subjected to a high source-drain bias [32]. Prior bias-spectroscopy studies have linked such $0.25(2e^2/h)$ -type structures to bias-driven spin polarization, while weak-confinement experiments have associated the coexistence of $0.25(2e^2/h)$, $0.5(2e^2/h)$, and row-coupling effects with strongly correlated quasi-1D states [3,6,32].

The $0.85(2e^2/h)$ feature is a well-established signature of partial spin polarization in conventional 1D quantum wires [69]; its appearance in the hybridized states reflects the distinct nature of the zero-field ground state. At zero magnetic field, the channel resides in a hybridized zigzag configuration, where bonding and antibonding states and additional low-energy modes ($<250 \mu\text{eV}$) contribute to a broad and continuous evolution of the nonlinear conductance. In this regime, the $0.85(2e^2/h)$ feature represents an intermediate state arising from the interplay between hybridization and a reduced exchange energy.

A perpendicular magnetic field reorganizes this spectrum in two complementary ways. First, orbital compression reduces the transverse extent of the wave function, thereby suppressing the two-row configuration and weakening the

effective tunnel-coupled reconstruction; this accounts for the recovery of the $2e^2/h$ plateau and the separation of the broad hybridized diamond into more conventional subband diamonds. Second, the Zeeman energy ($\sim 8 \mu\text{eV}$) becomes comparable to the small exchange scale J , making spin splitting more visible once the hybridization is reduced. The field therefore suppresses the transverse charge configuration while preserving, and often sharpening, the spin-related structures in the bias spectroscopy [3,6,19,32,69]. The system evolves more abruptly from the spin-degenerate $2e^2/h$ plateau toward spin-resolved structures at $0.5(2e^2/h)$ and $0.25(2e^2/h)$ under large bias, effectively bypassing the intermediate $0.85(2e^2/h)$ state. Thus the disappearance of the $0.85(2e^2/h)$ feature reflects a redistribution of spectral weight associated with the suppression of hybridization rather than the absence of spin-related physics. The ordering $J \sim E_Z \lesssim \Delta \lesssim \hbar\omega \ll E_C$ thus summarizes why dc-bias spectroscopy resolves (i) hybridization scales in the $100 \mu\text{eV}$ range and (ii) spin-related reorganizations under modest $B_\perp$.

IV. CONCLUSIONS

We have studied the interplay of confinement, carrier density, and perpendicular magnetic field in weakly confined quasi-1D quantum wires in the strongly interacting regime. At low density, the suppression of the $2e^2/h$ plateau and the emergence of an avoided crossing reveal an interaction-driven restructuring into a spatial zigzag or two-row configuration with hybridized subbands. A perpendicular magnetic field suppresses this anticrossing and restores the $2e^2/h$ plateau through magnetic compression, which reduces inter-row hybridization. Despite this orbital reorganization, interaction effects remain robust: conductance features near $0.7(2e^2/h)$ and $0.5(2e^2/h)$ are enhanced in the low-density regime, with a narrow window where $2e^2/h$, $0.7(2e^2/h)$, and $0.5(2e^2/h)$ coexist. This behavior reflects a crossover toward a spin-incoherent Wigner regime driven by the suppression of exchange. Our results demonstrate that magnetic confinement decouples orbital reconstruction from spin physics, establishing a unified framework in which hybridization, Coulomb interactions, and reduced exchange govern transport in low-density quantum wires.

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DATA AVAILABILITY

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

- [1] S. Kumar and M. Pepper, Interactions and non-magnetic fractional quantization in one-dimension, *Appl. Phys. Lett.* **119**, 110502 (2021).
- [2] A. G. Pogosov, A. A. Shevyrin, D. A. Pokhobov, E. Yu. Zhdanov, and S. Kumar, Suspended semiconductor nanostructures: Physics and technology, *J. Phys.: Condens. Matter* **34**, 263001 (2022).
- [3] S. Kumar, K. J. Thomas, L. W. Smith, M. Pepper, G. L. Creeth, I. Farrer, D. Ritchie, G. Jones, and J. Griffiths, Many-body effects in a quasi-one-dimensional electron gas, *Phys. Rev. B* **90**, 201304(R) (2014).
- [4] S. Kumar, M. Pepper, D. Ritchie, I. Farrer, and H. Montagu, Formation of a non-magnetic, odd-denominator fractional quantized conductance in a quasi-one-dimensional electron system, *Appl. Phys. Lett.* **115**, 123104 (2019).
- [5] S. Kumar, M. Pepper, S. N. Holmes, H. Montagu, Y. Gul, D. A. Ritchie, and I. Farrer, Zero-magnetic field fractional quantum states, *Phys. Rev. Lett.* **122**, 086803 (2019).
- [6] V. Kumar, Y. Duo, P. See, J. P. Griffiths, D. A. Ritchie, I. Farrer, and S. Kumar, Correlation effects mediated by disorders in one-dimensional channels, *Phys. Rev. B* **112**, 035404 (2025).
- [7] H. Montagu, I. Farrer, D. Ritchie, and S. Kumar, Spin polarised quantum conductance in 1D channels, *Appl. Phys. Express* **18**, 015002 (2025).
- [8] B. Antonio, A. Bayat, S. Kumar, M. Pepper, and S. Bose, Self-assembled Wigner crystals as mediators of spin currents and quantum information, *Phys. Rev. Lett.* **115**, 216804 (2015).
- [9] C. Yan, S. Kumar, K. Thomas, M. Pepper, P. See, I. Farrer, D. Ritchie, J. P. Griffiths, and G. A. C. Jones, Direct observation of exchange-driven spin interactions in one-dimensional system, *Appl. Phys. Lett.* **111**, 042107 (2017).
- [10] C. Yan, S. Kumar, K. Thomas, P. See, I. Farrer, D. Ritchie, J. Griffiths, G. Jones, and M. Pepper, Engineering the spin polarization of one-dimensional electrons, *J. Phys.: Condens. Matter* **30**, 08LT01 (2018).
- [11] C. Yan, S. Kumar, K. Thomas, P. See, I. Farrer, D. Ritchie, J. Griffiths, G. Jones, and M. Pepper, Coherent spin amplification using a beam splitter, *Phys. Rev. Lett.* **120**, 137701 (2018).
- [12] S. Kumar and M. Pepper, Advances in interaction effects in the quasi one-dimensional electron gas, *Frontiers of Nanoscience* (Elsevier, Amsterdam, 2021), Vol. 20, pp. 7–29.
- [13] K.-F. Berggren and M. Pepper, Electrons in one dimension, *Philos. Trans. R. Soc. A* **368**, 1141 (2010).
- [14] T. J. Thornton, M. Pepper, H. Ahmed, D. Andrews, and G. J. Davies, One-dimensional conduction in the 2D electron gas of a GaAs-AlGaAs heterojunction, *Phys. Rev. Lett.* **56**, 1198 (1986).
- [15] B. J. van Wees, H. van Houten, C. W. J. Beenakker, J. G. Williamson, L. P. Kouwenhoven, D. van der Marel, and C. T. Foxon, Quantized conductance of point contacts in a two-dimensional electron gas, *Phys. Rev. Lett.* **60**, 848 (1988).
- [16] D. A. Wharam, T. J. Thornton, R. Newbury, M. Pepper, H. Ahmed, J. E. F. Frost, D. G. Hasko, D. C. Peacock, D. A. Ritchie, and G. A. C. Jones, One-dimensional transport and the quantisation of the ballistic resistance, *J. Phys. C* **21**, L209 (1988).
- [17] K. J. Thomas, J. T. Nicholls, M. Y. Simmons, M. Pepper, D. R. Mace, and D. A. Ritchie, Possible spin polarization in a one-dimensional electron gas, *Phys. Rev. Lett.* **77**, 135 (1996).
- [18] W. K. Hew, K. J. Thomas, M. Pepper, I. Farrer, D. Anderson, G. A. C. Jones, and D. A. Ritchie, Spin-incoherent transport in quantum wires, *Phys. Rev. Lett.* **101**, 036801 (2008).
- [19] W. K. Hew, K. J. Thomas, M. Pepper, I. Farrer, D. Anderson, G. A. C. Jones, and D. A. Ritchie, Incipient formation of an electron lattice in a weakly confined quantum wire, *Phys. Rev. Lett.* **102**, 056804 (2009).
- [20] A. Yacoby, H. L. Stormer, N. S. Wingreen, L. N. Pfeiffer, K. W. Baldwin, and K. W. West, Nonuniversal conductance quantization in quantum wires, *Phys. Rev. Lett.* **77**, 4612 (1996).
- [21] E. Wigner, On the interaction of electrons in metals, *Phys. Rev.* **46**, 1002 (1934).
- [22] R. S. Crandall and R. Williams, Crystallization of electrons on the surface of liquid helium, *Phys. Lett. A* **34**, 404 (1971).
- [23] C. C. Grimes and G. Adams, Evidence for a liquid-to-crystal phase transition in a classical, two-dimensional sheet of electrons, *Phys. Rev. Lett.* **42**, 795 (1979).
- [24] Yu. P. Monarkha and V. E. Syvokon, A two-dimensional Wigner crystal, *Low Temp. Phys.* **38**, 1067 (2012).
- [25] A. V. Chaplik, Possible crystallization of charge carriers in low-density inversion layers, *Zh. Eksp. Teor. Fiz.* **62**, 746 (1972) [*Sov. Phys. JETP* **35**, 395 (1972)].
- [26] A. V. Chaplik, Instability of quasi-one-dimensional electron chain and the “string-zigzag” structural transition, *Pis'ma Zh. Eksp. Teor. Fiz.* **31**, 275 (1980) [*JETP Lett.* **31**, 252 (1980)].
- [27] G. Piacente, I. V. Schweigert, J. J. Betouras, and F. M. Peeters, Generic properties of a quasi-one-dimensional classical Wigner crystal, *Phys. Rev. B* **69**, 045324 (2004).
- [28] G. Piacente, G. Q. Hai, and F. M. Peeters, Continuous structural transitions in quasi-one-dimensional classical Wigner crystals, *Phys. Rev. B* **81**, 024108 (2010).
- [29] J. S. Meyer and K. A. Matveev, Wigner crystal physics in quantum wires, *J. Phys.: Condens. Matter* **21**, 023203 (2009).
- [30] O. P. Sushkov, Conductance anomalies in a one-dimensional quantum contact, *Phys. Rev. B* **64**, 155319 (2001).
- [31] L. I. Glazman, I. M. Ruzin, and B. I. Shklovskii, Quantum transport and pinning of a one-dimensional Wigner crystal, *Phys. Rev. B* **45**, 8454 (1992).
- [32] L. W. Smith, W. K. Hew, K. J. Thomas, M. Pepper, I. Farrer, D. Anderson, G. A. C. Jones, and D. A. Ritchie, Row coupling in an interacting quasi-one-dimensional quantum wire investigated using transport measurements, *Phys. Rev. B* **80**, 041306(R) (2009).
- [33] A. Bayat, S. Kumar, M. Pepper, and S. Bose, Quantum phase transition detected through one-dimensional ballistic conductance, *Phys. Rev. B* **96**, 041116(R) (2017).
- [34] S.-C. Ho, H.-J. Chang, C.-H. Chang, S.-T. Lo, G. Creeth, S. Kumar, I. Farrer, D. Ritchie, J. Griffiths, G. Jones, M. Pepper, and T.-M. Chen, Imaging the zigzag Wigner crystal in confinement-tunable quantum wires, *Phys. Rev. Lett.* **121**, 106801 (2018).
- [35] I. Shapir, A. Hamo, S. Pecker, C. P. Moca, Ö. Legeza, G. Zarand, and S. Ilani, Imaging the electronic Wigner crystal in one dimension, *Science* **364**, 870 (2019).
- [36] S. Kumar, K. J. Thomas, L. W. Smith, M. Pepper, I. Farrer, D. A. Ritchie, G. A. C. Jones, and J. Griffiths, Effect of low transverse magnetic field on the confinement strength in a quasi-1D wire, *AIP Conf. Proc.* **1566**, 245 (2013).

- [37] S. Kumar, M. Pepper, H. Montagu, D. Ritchie, I. Farrer, J. Griffiths, and G. Jones, Engineering electron wavefunctions in asymmetrically confined quasi one-dimensional structures, *Appl. Phys. Lett.* **118**, 124002 (2021).
- [38] R. Davies, I. V. Lerner, and I. V. Yurkevich, Resonant fractional conductance through a 1D Wigner chain, *Commun. Phys.* **7**, 67 (2024).
- [39] T. Ando, A. B. Fowler, and F. Stern, Electronic properties of two-dimensional systems, *Rev. Mod. Phys.* **54**, 437 (1982).
- [40] J. H. Davies, *The Physics of Low-dimensional Semiconductors: An Introduction* (Cambridge University Press, Cambridge, UK, 1998).
- [41] K. A. Matveev, Conductance of a quantum wire at low electron density, *Phys. Rev. B* **70**, 245319 (2004).
- [42] S. Tarucha, T. Honda, and T. Saku, Reduction of quantized conductance at low temperatures observed in 2 to 10 μm -long quantum wires, *Solid State Commun.* **94**, 413 (1995).
- [43] C. Zener, Non-adiabatic crossing of energy levels, *Proc. A* **137**, 696 (1932).
- [44] L. Landau, Zur theorie der energieübertragung. II, *Phys. Z. Sowjetunion* **2**, 46 (1932).
- [45] C. Cohen-Tannoudji, B. Diu, and F. Laloë, *Quantum Mechanics, Volume III: Fermions, Bosons, Photons, Correlations, and Entanglement* (John Wiley & Sons, New York, 2019).
- [46] W. G. Van der Wiel, S. De Franceschi, J. M. Elzerman, T. Fujisawa, S. Tarucha, and L. P. Kouwenhoven, Electron transport through double quantum dots, *Rev. Mod. Phys.* **75**, 1 (2002).
- [47] C.-K. Wang and K.-F. Berggren, Spin splitting of subbands in quasi-one-dimensional electron quantum channels, *Phys. Rev. B* **54**, R14257 (1996).
- [48] H. Bruus, V. V. Cheianov, and K. Flensberg, The anomalous 0.5 and 0.7 conductance plateaus in quantum point contacts, *Phys. E* **10**, 97 (2001).
- [49] K.-F. Berggren and I. I. Yakimenko, Effects of exchange and electron correlation on conductance and nanomagnetism in ballistic semiconductor quantum point contacts, *Phys. Rev. B* **66**, 085323 (2002).
- [50] H. Hou, Y. Kozuka, J.-W. Liao, L. W. Smith, D. Kos, J. P. Griffiths, J. Falson, A. Tsukazaki, M. Kawasaki, and C. J. B. Ford, Quantized conductance of one-dimensional strongly correlated electrons in an oxide heterostructure, *Phys. Rev. B* **99**, 121302(R) (2019).
- [51] A. C. Graham, K. J. Thomas, M. Pepper, N. R. Cooper, M. Y. Simmons, and D. A. Ritchie, Interaction effects at crossings of spin-polarized one-dimensional subbands, *Phys. Rev. Lett.* **91**, 136404 (2003).
- [52] A. D. Klironomos, J. S. Meyer, T. Hikihara, and K. A. Matveev, Spin coupling in zigzag Wigner crystals, *Phys. Rev. B* **76**, 075302 (2007).
- [53] L. L. Sohn, L. P. Kouwenhoven, and G. Schön, *Mesoscopic Electron Transport*, NATO Science Series E (Springer, Dordrecht, 1997).
- [54] L. I. Glazman and A. V. Khaetskii, Nonlinear quantum conductance of a lateral microconstraint in a heterostructure, *Europhys. Lett.* **9**, 263 (1989).
- [55] D. L. Maslov and M. Stone, Landauer conductance of Luttinger liquids with leads, *Phys. Rev. B* **52**, R5539 (1995).
- [56] L. I. Glazman, G. B. Lesovik, D. E. Khmel'nitskii, and R. I. Shekhter, Reflectionless quantum transport and fundamental ballistic-resistance steps in microscopic constrictions, *Pis'ma Zh. Eksp. Teor. Fiz.* **48**, 218 (1988) [*JETP Lett.* **48**, 238 (1988)].
- [57] T. Giamarchi, *Quantum Physics in One Dimension* (Oxford University Press, Oxford, UK, 2003).
- [58] K. A. Matveev, Conductance of a quantum wire in the Wigner-crystal regime, *Phys. Rev. Lett.* **92**, 106801 (2004).
- [59] A. C. Mehta, C. J. Umrigar, J. S. Meyer, and H. U. Baranger, Zigzag phase transition in quantum wires, *Phys. Rev. Lett.* **110**, 246802 (2013).
- [60] J. von Neumann and E. P. Wigner, Über merkwürdige diskrete eigenwerte, in *The Collected Works of Eugene Paul Wigner: Part A: The Scientific Papers*, edited by A. S. Wightman (Springer Berlin Heidelberg, Berlin, Heidelberg, 1993), pp. 291–293.
- [61] C. Cohen-Tannoudji, B. Diu, F. Laloë, S. R. Hemley, N. Ostrowsky, and D. B. Ostrowsky, in *Quantum Mechanics: Angular Momentum, Spin, and Approximation Methods*, edited by C. Cohen-Tannoudji, B. Diu, and F. Laloë (Wiley-VCH, Berlin, 2020), Vol. 2.
- [62] C. W. J. Beenakker and H. van Houten, Quantum transport in semiconductor nanostructures, *Solid State Phys.* **44**, 1 (1991).
- [63] K.-F. Berggren, T. J. Thornton, D. J. Newson, and M. Pepper, Magnetic depopulation of 1D subbands in a narrow 2D electron gas in a GaAs:AlGaAs heterojunction, *Phys. Rev. Lett.* **57**, 1769 (1986).
- [64] M. M. Fogler and E. Pivovarov, Exchange interaction in quantum rings and wires in the Wigner-crystal limit, *Phys. Rev. B* **72**, 195344 (2005).
- [65] S. M. Cronenwett, H. J. Lynch, D. Goldhaber-Gordon, L. P. Kouwenhoven, C. M. Marcus, K. Hirose, N. S. Wingreen, and V. Umansky, Low-temperature fate of the 0.7 structure in a point contact: A Kondo-like correlated state in an open system, *Phys. Rev. Lett.* **88**, 226805 (2002).
- [66] A. P. Micolich, What lurks below the last plateau: Experimental studies of the $0.7 \times 2e^2/h$ conductance anomaly in one-dimensional systems, *J. Phys.: Condens. Matter* **23**, 443201 (2011).
- [67] G. A. Fiete, *Colloquium: The spin-incoherent Luttinger liquid*, *Rev. Mod. Phys.* **79**, 801 (2007).
- [68] N. K. Patel, J. T. Nicholls, L. Martin-Moreno, M. Pepper, J. E. F. Frost, D. A. Ritchie, and G. A. C. Jones, Evolution of half plateaus as a function of electric field in a ballistic quasi-one-dimensional constriction, *Phys. Rev. B* **44**, 13549 (1991).
- [69] T.-M. Chen, A. C. Graham, M. Pepper, I. Farrer, and D. A. Ritchie, Bias-controlled spin polarization in quantum wires, *Appl. Phys. Lett.* **93**, 032102 (2008).
- [70] H. Lind, I. I. Yakimenko, and K.-F. Berggren, Electric-field control of magnetization in biased semiconductor quantum wires and point contacts, *Phys. Rev. B* **83**, 075308 (2011).
- [71] K. J. Thomas, J. T. Nicholls, W. R. Tribe, M. Y. Simmons, A. G. Davies, D. A. Ritchie, and M. Pepper, Bonding and antibonding states in strongly coupled ballistic one-dimensional wires, *Phys. E* **6**, 581 (2000).
- [72] C. Rössler, S. Baer, E. de Wiljes, P.-L. Ardel, T. Ihn, K. Ensslin, C. Reichl, and W. Wegscheider, Transport properties of clean quantum point contacts, *New J. Phys.* **13**, 113006 (2011).
- [73] K. Nikolić and A. MacKinnon, Conductance and conductance fluctuations of narrow disordered quantum wires, *Phys. Rev. B* **50**, 11008 (1994).

- [74] P. L. McEuen, B. W. Alphenaar, R. G. Wheeler, and R. N. Sacks, Resonant transport effects due to an impurity in a narrow constriction, *Surf. Sci.* **229**, 312 (1990).
- [75] Y. J. Chung, A. Gupta, K. W. Baldwin, K. W. West, M. Shayegan, and L. N. Pfeiffer, Understanding limits to mobility in ultrahigh-mobility GaAs two-dimensional electron systems: 100 million cm^2/Vs and beyond, *Phys. Rev. B* **106**, 075134 (2022).
- [76] V. Umansky, M. Heiblum, Y. Levinson, J. Smet, J. Nübler, and M. Dolev, MBE growth of ultra-low disorder 2DEG with mobility exceeding $35 \times 10^6 \text{cm}^2/\text{Vs}$, *J. Cryst. Growth* **311**, 1658 (2009).