Spin-orbit coupling and transport in strongly correlated two-dimensional systems
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Spin-orbit coupling and the transport of strongly correlated two-dimensional systems

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Measuring the magnetoresistance (MR) of ultraclean GaAs two-dimensional holes in a large $r_s$ range of 20-50, two striking behaviors in relation to the spin-orbit coupling (SOC) emerge in response to strongly electron-electron interaction. First, in exact correspondence to the zero-field metal-to-insulator transition (MIT), the sign of the MR switches from being positive in the metallic regime to being negative in the insulating regime when the carrier density crosses the critical density $p_c$ of MIT ($r_s \sim 39$). Second, as the SOC-driven correction $\Delta \rho$ to the MR decreases with reducing carrier density (or the in-plane wave vector), it exhibits an upturn in the close proximity just above $p_c$ where $r_s$ is beyond 30, indicating a substantially enhanced SOC effect. This peculiar behavior echoes with a trend of delocalization long suspected for the SOC-interaction interplay. Meanwhile, for $p < p_c$ or $r_s > 40$, in contrast to the common belief that a magnet field enhances Wigner crystallization, the negative MR is likely linked to enhanced interaction.

Electron systems governed by strong Coulomb interaction manifest unique charge transport behaviors characterized by the collective modes such as those in the Wigner crystal (WC)\(^1\), fractional quantum Hall states (FQHS)\(^2\), and the zero-field metal-to-insulator transition (MIT)\(^3\). There is an important interaction-driven effect associated with the spin-orbit coupling (SOC) which has attracted a lot of interests due to its promising spintronic applications. SOC studies often utilize systems that lack inversion symmetries, both bulk and structural inversion asymmetries (BIA and SIA), which give rise to Dresselhaus and Rashba SOCs\(^5\). As a result, the spin degeneracy of the energy bands is lifted even in a zero magnetic field ($B$). The band splitting in response to controllable quantities, especially electron-electron interaction, is a fundamental effect that is not well understood. SOC-driven effects in weakly interacting systems are usually perturbations, i.e. those tested $p$-type two-dimensional (2D) $InGaAs$, GaAs/AlGaAs heterostructures with carrier densities $\sim 2-5 \times 10^{11}$ cm$^{-2}$, where $r_s = E_{\text{ee}}/E_F = a/a_B \sim 1 - 5$ and $k_F l \gg 1$. $E_{\text{ee}} \sim e^2/er$-Coulomb energy, $E_F = \pi n \hbar^2/m^*\text{-Fermi energy, } a\text{-average charge spacing, } a_B\text{-Bohr radius, } k_F\text{-Fermi wave vector, and } l\text{-mean free path. However, in a strongly correlated system, the SOC-interaction interplay modifies the exchange interaction and leads to more prominent effects}\(^6-8\). For example, a diverging density of states (DOS) known as the Van Hove singularity is expected in the limit of small wavevectors $k_\parallel$ (or very low carrier density). However, such effects have not been studied previously due to lack of access to strongly interaction-driven systems. Large $r_s = E_{\text{ee}}/E_F = a/a_B \gg 1$ in the absence of a magnetic field requires very low carrier concentrations for which the usual level of disorder renders an Anderson localization (or percolation transition). This study utilizes ultra-high purity two-dimensional (2D) holes in GaAs heterostructures and demonstrates enhanced SOC-driven effects in a truly interaction-driven limit.

In $p$-type (accumulation-type) GaAs heterostructures, SOC results in splittings of both the light hole (LH) and heavy hole (HH) bands. While the BIA-induced effect remains constant, the SIA (Rashba) contribution can be tuned externally, i.e. via a metal gate. For high carrier densities, BIA-induced splitting of the HH band is dominated by a $k_\parallel^3$ contribution ($k_\parallel$ is the in-plane carrier momentum). However, the situation becomes complicated when the carriers are sufficiently dilute. Apart from weakening the external electric field, reducing the gate bias modifies the confinement potential which affects the HH-LH separation and thus influences the Rashba coefficient ($\alpha$). Meanwhile, the effect of the LH-HH mixing (neparabolic dispersion) is expected to rise. Recently MR measurements on the $p$-dependence shows a moderately increasing splitting with reducing $k_\parallel^{12}$, suggesting an enhancement of $\alpha$ possibly related to the LH-HH mixing. However, due to limited density range, such effects are not explored into the strongly correlated metal-to-insulator (MIT) transition regime\(^7\) where larger enhancement is anticipated for small enough $k_\parallel^{11}$.

We adopt undoped ultra-dilute two-dimensional (2D) holes in GaAs HIGFETs (hetero-junction-insulated-gate field-effect-transistors) in which the carrier density $p$ can be continuously tuned from $0.2$ to $2 \times 10^{10}$ cm$^{-2}$ (or $r_s$ from 60 to 25). This allows a first probe to the MR across MIT with the critical density $p_c$ being only $4 \times 10^9$ cm$^{-2}$ (or $r_s \sim 39$ if an effective mass $m^* = 0.3 m_0$ is assumed). The zero-field temperature ($T$) dependence of the conductance shows non-activated transport, excluding the domination of an Anderson-localization, which is

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a crucial indicator for an interaction-driven nature. For
a decreasing $p$ from well above $p_c$ down to the proximity
of $p_c$, the measured $p$-dependence of the MR captures a
substantial resistivity variation (or correction) $\Delta \rho$ that
Corresponds to a nearly three-fold rise around $6 \times 10^9$
cm$^{-2}$. Moreover, crossing into an insulator at $p_c$, the
positive sign of MR switches rapidly to negative, in con-
tact to the usual expectation of a WC stabilized by a $B$
field.

The devices adopted are 2D holes in undoped $p$-
channel (100) HIGFET$^{13-15}$ patterned into 1 mm $\times$ 6
mm Hallbar shapes. As illustrated in Fig. 1 (a), the 2D
holes are confined in the $\sim$10-30 nm triangular well at
GaAs/Al$_{0.3}$GaAs hetero-interface. The Ohmic contacts
are made with AuBe alloy annealed at 420°C and the
contacts to the top gate layer of n$^+$-GaAs are made with
Cr/Au alloy without annealing. The top gate, a 30 nm
heavily doped GaAs layer, and the 2D hole layer form a
capacitor with 600 nm of dielectric AlGaAs layer in be-
tween. The 2D holes are, as shown in Fig. 1 (b), only
capacitively induced by gating the $n^+$ layer beyond the
threshold voltage $V_T \approx -1.3$ V with respect to 2D channel.
$p$ is tuned from $7 \times 10^8$ cm$^{-2}$ to $1.5 \times 10^{10}$ cm$^{-2}$
varying $V_{gate}$ and is determined through quantum Hall
measurements with a field sweep rate of 100 Gauss/min
in a dilution refrigerator.

We first present the $T$-dependence of the conductivity ($\sigma$)
and resistivity ($\rho$) results for a range of $p$ from
2 to $18 \times 10^9$ cm$^{-2}$ in a zero-field. It is well known
that hopping conductance is often assumed for all di-
lute systems. This means that, with small $E_{cc} \sim 1$meV
or 10 K and $E_F \sim 300$mk, disordered systems usually
give in to unscreened (or poorly screened) disorders and
undergoes Anderson Localization$^{16}$. As a result, one
observes exponentially suppressed conductance by cool-
ing: $\sigma \sim e^{-(T/T)\nu}$. $\nu = 1$ is for Arrhenius case at
higher $T$ and $\nu = 1/3$ and 1/2 for variable-range hop-
ing (VRH)$^{17,18}$. Even with relatively clean systems,
Anderson scenario can still hold if the single particle
localization length $\xi$ is exceeded by the average charge
spacing $a = 1/\sqrt{\pi \rho}$ at low enough $p$. Thus, ultrahigh
purity systems are required so that interaction-driven na-
ture persists even at the onset of a WC$^{15}$, $r_s \sim 40$, which
Corresponds to $p \sim 4 \times 10^8$ cm$^{-2}$ (or $1 \times 10^9$ cm$^{-2}$ for
electrons).

As shown in Fig. 2(a) in log-log scales, the apparent
MIT occurs around $p_c \sim 4 \times 10^9$ cm$^{-2}$ below which a
striking non-activated power law behavior$^{15,20}$ $\sigma \sim
(T/T^*)^\gamma$, is observed in multiple tested samples. $p_c$
Corresponds to $r_s \sim 39^{19}$, the anticipated onset point of
a WC. Notice that the critical density obtained pre-
viously in more disordered systems varies substantially
from sample to sample, i.e. $8 \times 10^{10}$ to $3 \times 10^9$ cm$^{-2}$
(or $3 < r_s < 15$) for electron devices and $8 \times 10^{10}$
to $8 \times 10^9$ cm$^{-2}$ (or $5 < r_s < 25$ for holes. However, the $p_c$
found in HIGFETs is significantly lower and varies little
among many different samples: 3.8 to $4.4 \times 10^9$ cm$^{-2}$.
The nonactivated power-law $T$-dependence on the insu-
lating side contrasts the disorder-dominated hopping sce-
nario. The exponent $\gamma$ varies from 1 to 2 with decreasing
$p$. This power law $T$-dependence, which has been
reported for both GaAs$^{20}$ and SiGe$^{21}$ systems, likely be-
longs to a strongly correlated liquid (i.e. a melted WC
or glass) since the experimental temperature is above the
WC melting temperature $T_m$.

An important fact about the power-law $T$-dependence
is its vulnerability to even a slight increase of disorder.
Long-range disorders are introduced through LED illu-
mination to the sample within the same cooling cycle.
The photons energy is approximately the band gap in
GaAs and AlGaAs. There is at least 24 hours of wait
time after the illumination before measurement resumes
and the uniformity of the carriers are verified via the
Hall measurement. The amount of disorder introduced
is roughly controlled by the time duration of the light
exposure, which is usually 0.5-2s, at a constant 0.05 $\mu$A
current excitation. This causes $p_c$ to rise$^{22}$: i.e. $p_c$ is
almost doubled after 0.5s of illumination and tripled after
1.5 s illumination. Meanwhile, as shown in the $log – log$
plot in Fig. 2(b) for $p = 1.5 - 2.8 \times 10^9$ cm$^{-2}$, $\rho(T)$
undergoes a qualitative change from the power-law (prior
to the illumination) into an exponential law $\rho \sim e^{-(T/T)}$.
The dimensionless parameter \( \mu \) for low \( T \) MR results. The interaction-driven nature which is foundational to the following MR results.

We performed further work to confirm this by studying \( \rho_{xx} \) jumps by orders of magnitude. This qualitative change of behaviors indicates that the power-law represents a different state likely due to strong correlation. We performed further work to confirm this by studying the exponent \( \gamma \) of the power-law and found it scales with a dimensionless parameter \( a/d \) with \( a \) being the charge spacing, and \( d \) the distance to the gate. This is due to the onset of a dipolar screening and thus confirms the interaction-driven nature which is foundational to the following MR results.

The longitudinal MR, \( \rho_{xx}(B) \), shown in Fig. 3(a), is for \( p \approx 1.2 \times 10^{10} \text{ cm}^{-2} \) (or \( r_s \approx 30 \)) with a carrier mobility of \( \mu \approx 300,000 \text{ cm}^2/(\text{V} \cdot \text{s}) \) and \( \rho_{xx}(0) = 1.25 \text{ k}\Omega/\mu \approx 0.1 \text{h/}\varepsilon^2 \). \( h/\varepsilon^2 \) is the quantum resistance. Shubnikov-de Haas oscillations (SdH) [Fig. 3(b)] are observed between 0.05 to 0.25T before a substantial (re-entrant insulating) peak develops at \( B = 0.37 \text{T} \) proceeding the filling factor 1 (at \( <0.5 \text{T} \)). Fourier analysis of \( \rho_{xx}(1/B) \), shown in Fig. 3(c), resolves only a single frequency \( f \) peak at 0.25T. The width of the peak, \( \Delta f \approx 0.09 \text{T} \), is the level of uncertainty corresponding to a density difference \( \Delta \rho \) less than \( 2 \times 10^9 \text{ cm}^{-2} \). Therefore, the heavy hole (HH) band is approximately degenerate. The corresponding subband density \( p = (g_\ell e/h)f = 6 \times 10^9 \text{ cm}^{-2} \) with \( g_\ell = 2 \), leading to a total \( \rho \) of \( 1.2 \times 10^{10} \text{ cm}^{-2} \).

The correction \( \Delta \rho \) to the MR for the degenerate HH band differs from the scenario of amply split subbands that possess not only unequal charge carrier masses (due to the warped dispersion caused by SOC), but also the non-proportionally different densities. Significantly split HH subbands exhibit different mobility \( \mu_i \) (band index \( i = 1, 2 \)) and the interband scattering gives rise to a correction \( \Delta \rho(B)/\rho(0) \approx (\mu_1 - \mu_2)^2 \). However, for the (nearly) degenerate case, the two-band classical Drude term vanishes, so does the inter-band scattering since \( \mu_1 = \mu_2 \). In addition, even the single-band Drude term is out, if isotropic scattering assumed, because the Lorentz force is cancelled by the Hall field. Therefore, \( \Delta \rho(B)/\rho(B = 0) \) here is mainly due to the quantum corrections, especially by SOC, as long as a Fermi Liquid or, more generally, \( T < T_F \) is valid. As shown below, large \( \Delta \rho \) in the \( r_s \approx 1 \) limit grows progressively with decreasing \( p \) (or increasing interaction) in a non-monotonic fashion as the system undergoes MIT.

Fig. 4(a) shows the MR for a series of \( p \) from 0.2 to \( 1.5 \times 10^{10} \text{ cm}^{-2} \), approximately 1-5% of the carrier densities used in Refs.\(^{25,26} \). \( \rho \) reaches \( h/\varepsilon^2 \) (or 5.8 k\( \Omega \)) around \( p_c \approx 4 \times 10^9 \text{ cm}^{-2} \) where the sign of MR switches. For \( p > p_c \), the positive MR measures of the strength of SOC which produces corrections \( \Delta \rho(B) = \rho(B) - \rho(0) \) (in the unit of \( h/\varepsilon^2 \) dependent on \( p \) [Fig. 4(b)]; larger \( \Delta \rho(B) \) for lower \( p \), in the opposite trend of decreasing SOC with lower \( k_F \). The derivative \( dp/dB \) shown in Fig. 4(d), taken within 30 mT, captures this increasing WAL effect with decreasing \( p \) peaked around \( 6 \times 10^9 \text{ cm}^{-2} \). Negative MR develops rapidly below \( p_c \) and becomes stronger with decreasing \( p \). Fig. 4(c) shows the same \( \Delta \rho \) result in (b) normalized by \( \rho(0) \). Taking the derivatives \( dp/dB \) normalized by \( \rho(0) \), a local maximum is again captured as shown in Fig. 4(e).

As \( k_F \) is significantly reduced, the band splitting \( \Delta \) may exhibit a crossover from both the linear (BIA) and the cubic (SIA) dependence for larger \( k_F \) to the linear-\( k_F \) dependence for smaller \( k_F \). Unlike the depletion cases, decreasing the gate voltage \( V_g \) reduces the electric field according to \( E = (V_e - V_{fb})/d = -\nabla V \) (where \( V_{fb} \) is the flab-band voltage, \( d \) is the barrier thickness, and \( V \) is the approximately triangular confinement potential with a width \( w \)). The electric field is expressed as

\[
E = \frac{2E_0g_e\rho}{\epsilon}
\]
where $E_g$ is the triangular well height, $e$ the electron charge, and $\epsilon = 13$ is the dielectric constant. For a 50\% decrease in $p$ from 1.55 to $0.85 \times 10^{10}$ cm$^{-2}$, $E$ is reduced by approximately $\sim 15\%$. This agrees well with the gate voltage change from $V_g = -1.62$V to $V_g = -1.37$V corresponding to $p = 1.55$ and $0.85 \times 10^{10}$ cm$^{-2}$ respectively. The well width $w = (2eE_g/e\rho)^{1/2}$ is widened by 1.4 times, resulting in a 30\%-decrease in the x-direction (growth direction) wave-vector $k_z = \pi/\omega$ of the Airy-like wavefunctions. The Dresselhaus coupling parameter $\beta \propto k_z^2$ is halved. If we consider the $k_z^2$ splitting, it has a $p^{1.5}$ dependence since $p = 1.55$ and $0.85 \times 10^{10}$ cm$^{-2}$. Because the observed overall correction is linearly decreasing for this range, it suggests a lesser Rashba effect than the Dresselhaus for higher $p$.

Notice that for $p = 1.55, 1.22, and 0.85 \times 10^{10}$ cm$^{-2}$, the system is already strongly correlated with $\tau_s = 20 \sim 28$. It turns out the usual fitting of MR corrections based on the non-interacting model already becomes problematic. Our fitting attempts produces $\tau_s, 1.5$, and 0.6 ps respectively for $\tau_s$, however, with large errors ($R^2 = 0.94 \sim 0.89$). Note that $k_F$, which is between 1 and 10, is much smaller than the cases in ref.$^{26}$ and $\Delta p/p(0)$, amounting to $\sim 15\%$ at $B = 50m$T, is beyond perturbation. Clearly, a model appropriately incorporating interaction is needed. The theory on $\Delta \sigma(B)$ due to interaction by Althuizer et al$^{11}$ might be more relevant, even though it does not produce non-monotonic corrections for the large $r_s$ situations discussed below.

Remarkably, the decreasing trend in $[1/p(0)]d\rho/dB$ with lowering of $p$ is replaced by a striking 40\% increase (or 300\% in $d\rho/dB$) around $6 \times 10^9$ cm$^{-2}$. It signifies a substantially rising SOC coupling coefficient, particularly, the Rashba coefficient $\alpha$. In principle, as the $k_z^2$ dependence fades at low $p$, the linear-$k_z^2$ (or $p^{1/2}$) dependence eventually gives into to the rising $\alpha \propto p^{-4/3}$, resulting in the upturn of $[1/p(0)]d\rho/dB$. This is qualitatively in agreement with the data. There are other relevant effects that should also be recognized. Lowering $p$ to extremely dilute limits shrinks $k_z$ to the HH-LH anti-crossing point$^{30}$ where the consequent effect on $\alpha$ is not fully known. Meanwhile, growing effective interaction with reducing $p$ raises $r_s$ beyond 30 and it gives rise to effects via exchange interaction$^{11}$ through an enhancement factor $\lambda_{SO}$. Theory predicts an approximately linear relation $\lambda_{SO} \sim r_s$ for moderate $r_s$ values less than 20$^{11}$. Another interaction-driven effect derives from the enhancement of the effective mass $m^*$ which has been known to exist in the close vicinity of the $p_c$ of the MIT$^{31}$. Due to the non-parabolic dispersion relation at low $p$ and the HH-LH mixing, $m^*$ is a complicated quantity. Actual comparison would require a measurement of $m^*$ which is difficult to achieve with cyclotron resonance (due to the small energy) or the SdH (due to the large Coulomb energy). Nevertheless, the non-monotonic rising is consistent with an enhanced SOC which influences the MR opposite to the trend of localization as $p$ crosses below $p_c$. Therefore, SOC supports MIT$^2$. $\Delta p(B) \rightarrow 0$ at $p = p_c$ corresponds to the diminishing WAL eventually overcome by the localization. A Wigner crystal (WC) regime is arrived at $p < p_c$ or $r_s > 40$. The role of SOC depends on the actual carrier state which is likely a WC liquid because of the following: Due to the tiny energy scales, $E_F \sim 150 \sim 400$ mK, a WC is fragile to the influences of disorder, thermal$^{32}$ and quantum fluctuations (on a scale of $\sim 400$ nm/$\sqrt{T}$). Consequently, the melting temperature $T_m$ is usually reduced well below the classical estimate: $\sim E_{ee}/130 \sim 120$mK. A recent study demonstrates a dynamical pinned WC marked by enormous pinning strength and extremely sharp dc VI threshold$^{33}$. The $T$-dependence of pinning suggest a $T_m \sim 30$ mK which is lower than the temperature for the MR measurement. Meanwhile, since those data also support a second-order phase transition, an intermediate phase$^{34}$ is relevant for the MR results for $T > T_m$. The quantum scenario still holds because $T$ is less than $T_F = E_F/k_B$, plasma frequency $\Omega = \sqrt{\tau_s E_F}$, and $E_{ee}$. Thus, the hydrodynamic flow model for the semi-quantum case$^{35,36}$ does not apply.

The negative MR for $p < p_c$ is intuitively surprising since an external $B$-field should stabilize a WC. On the other hand, the $B$-field raises the Zeeman energy $g\mu_B B$ which increases the overall carrier energy, favoring delocalization. Though, the Zeeman effect is small for higher $p$ (or weakly interacting) cases, it becomes prominent as the $g$-factor is greatly enhanced with increasing interaction (or lower $p$)$^{37}$. For a reducing $p$ (below $p_c$), both the absolute Coulomb energy $E_{ee} = e^2/\sqrt{\pi \rho}$ and the $E_F = (\pi \hbar^2/m^*)p$ decrease. Moreover, for lower enough $k_F$, 2D DOS diverges (towards the van Hove singularity), resulting in lower kinetic energy states. Thus, interaction effect, including the $g$-factor, is further enhanced. As the Zeeman energy grows with increasing $B$, a larger negative MR for lower $p$ is a possibility. Relevant transport theory is needed.

In summary, interaction-driven effects are realized in high purity ultra-dilute systems possessing unique transport behaviors distinct from the activated transport in Anderson insulators (or percolation) governed by disorder. SOC-driven band splitting is small in the dilute limit. Yet, there exists a substantial non-monotonic rise in the WAL in the metallic side, consistent with a delocalization effect. The switching of the signs of MR is in excellent correspondence to the zero-field MIT. The rising SOC coupling parameters are due to the large interaction including the scrambled exchange energy. The
negative MR in the insulating side seems to be consistent with a rising carrier energy through the Zeeman term facilitated by the reducing absolute Coulomb energy.

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1. E. P. Wigner, Physical Review B 46, 1002 (1934).
2. D. C. Tsui, H. L. Stormer, and A. C. Gossard, Phys. Rev. Lett. 48, 1559 (1982).
3. S. V. Kravchenko, W. E. Mason, G. E. Bowker, J. E. Furneaux, V. M. Pudalov, and M. D’Iorio, Phys. Rev. B 51, 7038 (1995).
4. G. Dresselhaus, Phys. Rev. 100, 580 (1955).
5. E. Rashba, Sov. Phys. Solid State 2, 1109 (1960).
6. B. Al’Tshuler, A. Aronov, A. Larkin, and D. Khmel’Nitskii, Zh. Eksp. Teor. Fiz. 81, 768 (1981).
7. V. Pudalov, Journal of Experimental and Theoretical Physics Letters 66, 175 (1997).
8. G. Giuliani and G. Vignale, Quantum theory of the electron liquid (Cambridge university press, 2005).
9. Y. Yaish, O. Prus, E. Buchstab, S. Shapira, G. B. Yoseph, U. Sivan, and A. Stern, Phys. Rev. Lett. 84, 4954 (2000).
10. G. H. Chen and M. E. Raikh, Phys. Rev. B 60, 4826 (1999).
11. R. Winkler, H. Noh, E. Tutuc, and M. Shayegan, Phys. Rev. B 65, 155303 (2002).
12. B. Kane, L. Pfeiffer, and K. West, Applied physics letters 67, 1262 (1995).
13. H. Noh, M. P. Lilly, D. C. Tsui, J. A. Simmons, L. N. Pfeiffer, and K. W. West, Phys. Rev. B 68, 241308 (2003).
14. J. Huang, D. S. Novikov, D. C. Tsui, L. N. Pfeiffer, and K. W. West, Phys. Rev. B 74, 201302 (2006).
15. P. W. Anderson, Physical review 109, 1492 (1958).
16. N. Mott, Journal of Non-Crystalline Solids 1, 1 (1968).
17. B. I. Shklovskii and A. L. Efros, Electronic properties of doped semiconductors, Vol. 45 (Springer Science & Business Media, 2013).
18. B. Tanatar and D. M. Ceperley, Phys. Rev. B 39, 5005 (1989).
19. J. Huang, L. N. Pfeiffer, and K. W. West, Phys. Rev. B 85, 041304 (2012).
20. K. Lai, W. Pan, D. C. Tsui, S. Lyon, M. Mühberger, and F. Schäffler, Phys. Rev. B 75, 033314 (2007).
21. J. Huang, L. Pfeiffer, and K. West, Applied Physics Letters 98, 092105 (2011).
22. J. Huang, L. N. Pfeiffer, and K. W. West, Phys. Rev. B 85, 041304(R) (2012).
23. J. Huang, L. N. Pfeiffer, and K. W. West, Phys. Rev. B 85, 041304(R) (2012).
24. R. Winkler, S. Papadakis, E. De Poortere, and M. Shayegan, Spin-Orbit Coupling in Two-Dimensional Electron and Hole Systems, Vol. 41 (Springer, 2003).
25. B. Habib, M. Shayegan, and R. Winkler, Semiconductor Science and Technology 24, 064002 (2009).
26. B. Grbić, R. Leturcq, T. Ihn, K. Ensslin, D. Reuter, and A. D. Wieck, Phys. Rev. B 77, 125312 (2008).
27. M. Rahimi, S. Anissimova, M. R. Sakr, S. V. Kravchenko, and T. M. Klapwijk, Phys. Rev. Lett. 91, 116402 (2003).
28. F. Stern, Phys. Rev. Lett. 33, 960 (1974).
29. S. Hikami, A. I. Larkin, and Y. Nagaoka, Progress of Theoretical Physics 63, 707 (1980).
30. J. N. Schulman and Y.-C. Chang, Phys. Rev. B 31, 2056 (1985).
31. E. Abrahams, S. V. Kravchenko, and M. P. Sarachik, Rev. Mod. Phys. 73, 251 (2001).
32. N. D. Mermin and H. Wagner, Phys. Rev. Lett. 17, 1133 (1966).
33. T. Knighton, A. Serafin, Z. Wu, J. Xia, J. Huang, L. N. Pfeiffer, and K. West, to be published.
34. B. Spivak and S. A. Kivelson, Phys. Rev. B 70, 155114 (2004).
35. A. V. Andreev, S. A. Kivelson, and B. Spivak, JETP Lett. 28, 557 (1978).
36. A. V. Andreev, S. A. Kivelson, and B. Spivak, Phys. Rev. Lett. 106, 256804 (2011).
37. F. F. Fang and P. J. Stiles, Phys. Rev. 174, 823 (1968).