

Hierarchy of quantum corrections to resistivity of a γ -Al₂O₃/SrTiO₃ two-dimensional electron gas

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A resistance upturn that occurs in SrTiO₃-based two-dimensional electron gases at low temperatures is often attributed to Kondo scattering of carriers from localized spins. We show that the low-temperature transport behavior of the electron gas that forms at the γ -Al₂O₃/SrTiO₃ interface is governed by a whole hierarchy of quantum corrections in addition to the Kondo contribution. Upon cooling, the resistance shows a typical Kondo-like upturn, but then decreases due to weak antilocalization, and finally vanishes at the superconducting transition. When superconductivity is suppressed by applying a magnetic field or by reducing the carrier density, a logarithmic correction due to electron-electron interactions appears at millikelvin temperatures. These results show that the low-temperature resistance behavior cannot be accounted for by a Kondo-only interpretation, but requires consideration of several quantum corrections.

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I. INTRODUCTION

Low-temperature resistance upturns that are often seen in SrTiO₃-based two-dimensional electron gases (2DEGs) were initially discussed in terms of Kondo scattering of itinerant electrons from localized spins, despite the fact that SrTiO₃ is a nominally nonmagnetic oxide and even the heterostructure component layers, such as LaAlO₃, contain no magnetic ions [1].

Subsequent experiments demonstrated that while a resistivity upturn does not occur in lightly-doped bulk SrTiO₃ [2,3] above the superconducting transition temperature, a Kondo-type resistivity correction can be induced even at nondoped SrTiO₃ surfaces when carriers are accumulated by electrostatic ion-gel gating [4]. First-principles studies have shown that in *n*-type SrTiO₃, each oxygen vacancy acts effectively as a single donor, trapping one electron as a small polaron on a Ti site while donating the other electron to the conduction band [5]. Consequently, singly ionized oxygen vacancies in SrTiO₃ can provide localized electrons that hybridize with the Ti 3*d* states, explaining the coexistence of itinerant carriers and localized magnetic moments [5–8]. On this basis, the low-temperature resistance upturn in SrTiO₃ 2DEGs can be understood as an intrinsic electronic Kondo effect rather than being caused by extrinsic magnetic dopants.

However, enhanced electronic phase coherence at low temperature can give rise to additional quantum corrections to resistivity in the same temperature range with Kondo scattering [9,10]. For the diffusive transport regime in oxide heterostructure 2DEGs, possible candidates are the electron-electron interaction (EEI) and weak antilocalization (WAL), both of which can produce a logarithmic temperature dependence of conductivity [11,12]. Among these, WAL originates from the Rashba-type spin-orbit interaction (SOI) induced by the broken inversion symmetry at an interface [13] and is well established in SrTiO₃-based heterostructures through magnetoresistance measurements [14,15].

Although the WAL and EEI corrections may coexist with the Kondo scattering, the resistivity behavior has often been analyzed under the assumption that Kondo scattering dominates, ignoring other quantum corrections [4,16–19]. Indeed, at the LaAlO₃/SrTiO₃ interface, multiband transport and orbital hybridization associated with the Lifshitz transition [20] complicate the interpretation of low-temperature transport, making it difficult to disentangle the contributions of individual quantum corrections.

In this work, we show that the more broadly tunable 2DEG at the γ -Al₂O₃/SrTiO₃ heterointerface can be used to conveniently separate the various resistivity corrections. Unlike polar oxide interfaces such as LaAlO₃/SrTiO₃, where the electrostatic interface charge accumulation due to the polar discontinuity sets a lower bound on the carrier density [21], the carrier density at the nonpolar γ -Al₂O₃/SrTiO₃ interface is controlled by the density of interfacial oxygen vacancies [22–24], enabling continuous tuning of conductivity from a metallic to an insulating state without electrostatic gating [25]. Additionally, the carriers at the γ -Al₂O₃/SrTiO₃ interface have a heavy out-of-plane *d*_{yz}/*d*_{zx} orbital character, with the light in-plane *d*_{xy} orbital lying at a higher energy [26–28],

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which means that there is no Lifshitz transition within the experimentally accessible carrier density range. These features allow low-temperature normal-state transport to be studied without gating in a single heavy-electron subband.

We disentangle the WAL and EEI contributions from the Kondo contribution at the γ -Al₂O₃/SrTiO₃ interface by transport measurements down to the mK regime in different magnetic field orientations. While the Kondo contribution approaches saturation in the strong-coupling limit at low temperatures [29], quantum corrections become stronger at lower temperatures [12] and respond differently to magnetic fields [12,30]. This allows us to separate the low-temperature transport mechanisms.

II. EXPERIMENTAL DETAILS

Crystalline γ -Al₂O₃ thin films with a total thickness of six unit cells (4.75 nm) were grown on nondoped wet-etched TiO₂-terminated SrTiO₃ (001) substrates with a 0.2° miscut (Shinkosha) by pulsed laser deposition (PLD) using a KrF excimer laser ($\lambda = 248$ nm) operated at a repetition rate of 1 Hz. Prior to film deposition, the substrates were annealed at 950 °C for 10 min under an oxygen partial pressure of 1×10^{-5} Torr to obtain a well-defined surface termination [31]. Film growth was carried out at a substrate temperature of 600 °C.

The carrier density at the γ -Al₂O₃/SrTiO₃ interface was controlled by the background oxygen pressure during deposition and by the laser fluence during the initial stage of film growth. The oxygen pressure influences the oxygen vacancy density near the interface [32–34], while the laser fluence affects the oxygen vacancy creation rate during PLD growth [35,36].

Two γ -Al₂O₃ films were grown by a two-step deposition process consisting of an initial one-unit-cell layer followed by a five-unit-cell capping layer. Sample A was grown at an oxygen partial pressure of 5×10^{-6} Torr, whereas sample B was grown at 1×10^{-5} Torr to obtain superconducting and nonsuperconducting interfaces, respectively. The laser fluence was varied during the deposition of the first unit cell layer. For sample A, the first unit cell was deposited using a laser fluence of 4.4 J/cm², whereas for sample B the first unit cell was deposited using a lower laser fluence of 1.0 J/cm² to reduce the number of oxygen vacancies generated at the SrTiO₃ substrate surface. Following the deposition of the initial unit cell, the remaining five unit cells were deposited as a capping layer using the same laser fluence of 1.3 J/cm² for both samples. This procedure kept the total deposition time similar for the two samples, reducing the risk of having different oxygen vacancy depth distributions near the interface [37].

Hall bar devices were fabricated by mechanical diamond milling to avoid unintended chemical reactions or contamination that may arise from wet processing. The Hall bars had a channel width of $W = 166$ μ m and a longitudinal voltage probe separation of $L = 537$ μ m. For both samples, the electrical resistance between the Hall bar channel and the surrounding regions exceeded 500 M Ω , indicating effective electrical isolation of the transport channel.

Electrical four-point transport and Hall measurements were carried out down to 1.8 K in a Physical Property Measurement

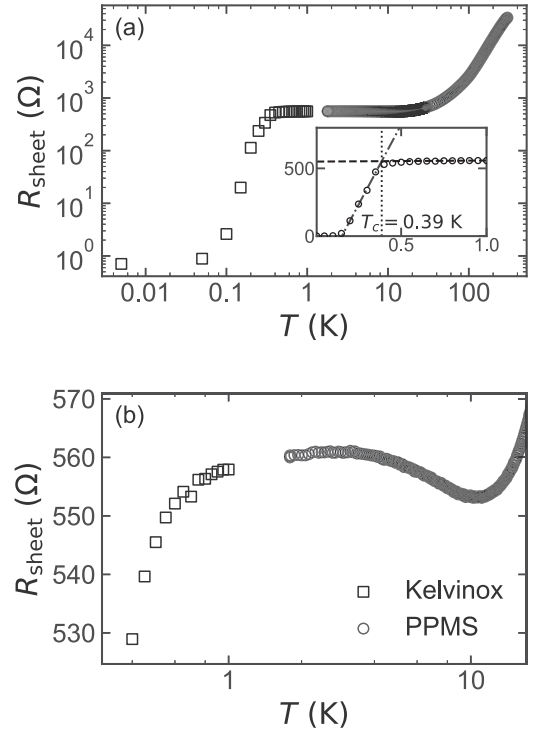


FIG. 1. (a) Temperature dependence of the sheet resistance of sample A. The inset shows an enlarged view of the low-temperature region below 1 K, highlighting the superconducting transition onset temperature $T_c = 0.39$ K. (b) Enlarged view of the low-temperature region, highlighting the logarithmic sheet resistance increase below 10 K and subsequent decrease due to WAL and superconductivity.

System (PPMS) using both ac and dc techniques and extended below 1 K using a top-loading dilution refrigerator (Kelvinnox), where only the ac lock-in technique was employed to minimize measurement-induced electron heating. The excitation current was 10 nA at 19 Hz for sample A and 23 Hz for sample B.

X-ray absorption spectroscopy (XAS) measurements were performed at BL-2A MUSASHI of the Photon Factory, KEK, using linearly polarized light via the measurement of the sample drain current. For Ti L -edge x-ray linear dichroism (XLD) measurements, XAS spectra were acquired at a fixed angle of 30° between the sample surface and the incident light under two polarization geometries [38]. In the in-plane (IP) geometry, the polarization vector $\mathbf{E}$ lay entirely within the sample plane, making the acquired spectrum I_{IP} primarily sensitive to the in-plane d_{xy} and $d_{x^2-y^2}$ orbitals. In the out-of-plane (OP) geometry, $\mathbf{E}$ had both in-plane and out-of-plane components but for the 30° incident angle, the I_{OP} spectrum mostly probes the out-of-plane d_{xz}/d_{yz} and d_{z^2} orbitals [26,27,39]. The XLD spectrum is defined as $I_{OP} - I_{IP}$.

III. RESULTS AND DISCUSSION

A. Transport characteristics of the γ -Al₂O₃/SrTiO₃ interface

The basic transport characteristics of the γ -Al₂O₃/SrTiO₃ interface are summarized in Fig. 1(a), which shows the temperature dependence of the sheet resistance of sample A from room temperature to 5 mK. Upon cooling, the resistance

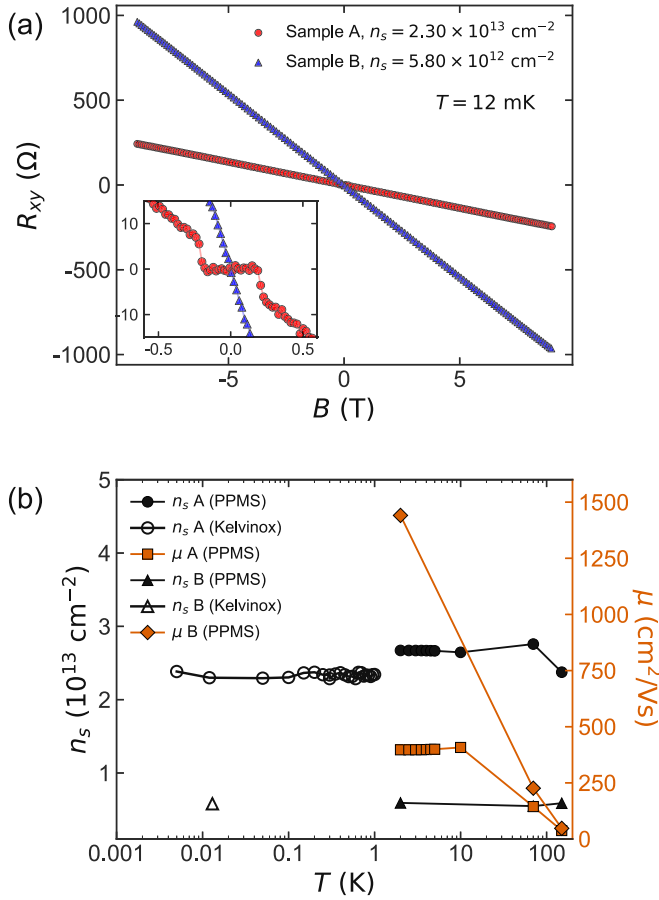


FIG. 2. (a) Hall resistance $R_{xy}(B)$ measured at $T = 12$ mK for samples A and B. The inset shows the low-field region on an expanded scale, highlighting the vanishing Hall resistance of sample A in the superconducting state below 200 mT. (b) Temperature dependence of the sheet carrier density n_s and mobility μ for samples A and B.

decreases overall, indicating metallic behavior, and drops to zero as the system enters the superconducting state. The superconducting transition onset temperature is $T_c = 0.39$ K, consistent with values reported for various SrTiO_3 -based systems [2,3,15,40–42]. Figure 1(b) shows an enlarged view of the low-temperature region. In this range, the resistance first decreases and then increases again below 10 K with a logarithmic temperature dependence. Below 3 K, this increase does not saturate and instead the resistance begins to decrease again.

Figure 2(a) shows the Hall resistance of samples A and B measured at 12 mK. Both samples exhibit a linear Hall response up to 9 T. The sheet carrier densities at 12 mK are $2.3 \times 10^{13} \text{ cm}^{-2}$ for sample A and $5.8 \times 10^{12} \text{ cm}^{-2}$ for sample B. In SrTiO_3 -based 2DEGs, a nonlinear low-temperature Hall response has often been attributed to multicarrier transport either due to a Lifshitz transition or carrier depth distribution [20,42–47]. No such nonlinearity is observed in either sample, showing that multicarrier transport is not present.

The low-field inset of Fig. 2(a) shows that the two samples are in different transport states at 12 mK. Below 200 mT,

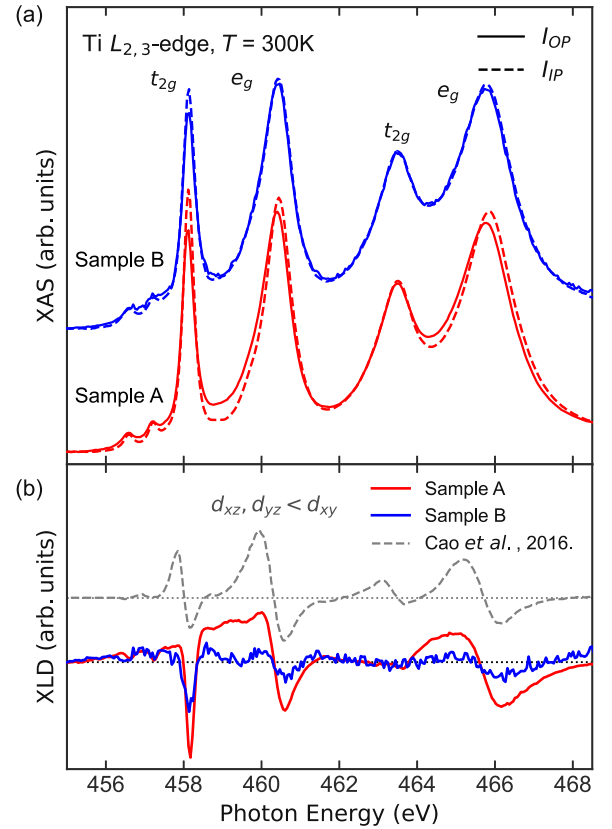


FIG. 3. (a) Ti $L_{2,3}$ -edge x-ray absorption spectra acquired at room temperature in the OP and IP geometries. (b) X-ray linear dichroism of samples A and B, defined as $I_{OP} - I_{IP}$ (solid lines). The gray dashed line denotes the reference data from Ref. [26], rescaled by the geometrical projection factor for comparison.

sample A shows zero Hall resistance due to entering the superconducting state. In contrast, sample B remains in the normal state and retains a linear Hall response down to zero field. This shows that the two samples represent superconducting and normal transport regimes within the same $\gamma\text{-Al}_2\text{O}_3/\text{SrTiO}_3$ interface system due to a slight carrier density difference. Sample B thus serves as a useful zero-field normal-state reference for the low-temperature transport.

Figure 2(b) shows the temperature dependences of the carrier density and mobility for samples A and B. The carrier density changes little upon cooling in either sample, ruling out carrier freeze-out as the origin of the anomalous low-temperature resistance behavior. In sample A, the mobility increases upon cooling and saturates below 10 K, following the temperature dependence of the dielectric permittivity of bulk SrTiO_3 [48]. SrTiO_3 is a well-known quantum paraelectric [49], where the permittivity increases sharply below 100 K, saturating around 10 K [50]. At 2 K, the mobility of sample A is about $400 \text{ cm}^2/\text{Vs}$, which is much lower than the highest value of about $5000 \text{ cm}^2/\text{Vs}$ reported for $\text{LaAlO}_3/\text{SrTiO}_3$ heterostructures [51]. This relatively low mobility is consistent with the heavy effective mass of carriers in the $\gamma\text{-Al}_2\text{O}_3/\text{SrTiO}_3$ 2DEGs.

Figure 3(a) shows the Ti $L_{2,3}$ -edge XAS spectra of samples A and B measured at room temperature in both out-of-plane

(I_{OP}) and in-plane (I_{IP}) polarizations. For both samples, the XAS spectra show the spin-orbit-split L_3 and L_2 structures, both of which further resolve into the t_{2g} ($d_{xy,xz,yz}$) and e_g ($d_{x^2-y^2,z^2}$) components [26,39,52].

Figure 3(b) shows the XLD spectra, defined as $I_{OP} - I_{IP}$. Both samples show XLD intensity dips at the high-energy edge of each absorption peak, indicating that the I_{OP} spectrum is shifted to a slightly lower energy. This XLD behavior is similar to previous reports for the γ - Al_2O_3 /SrTiO₃ interface [26] and indicates that the d_{xz}/d_{yz} orbitals lie at a lower energy than the d_{xy} orbitals. The γ - Al_2O_3 /SrTiO₃ 2DEG thus has the d_{xz}/d_{yz} heavy-electron character, consistent with the relatively low mobility [Fig. 2(b)].

B. Temperature-dependent resistance under different magnetic field configurations

We first examine to what extent the nonmonotonic low-temperature resistance of sample A in Fig. 1(a) can be explained by the Kondo effect alone. We express $R(T)$ at zero field as the sum of a metallic background term and a conventional Kondo term

$$R(T) = R_{\text{metal}}(T) + R_K(T), \quad (1)$$

with the metallic background defined as

$$R_{\text{metal}}(T) = R_0 + aT^2, \quad (2)$$

where R_0 is the residual resistance, and the aT^2 term represents the Fermi liquid electron-electron scattering [53]. In this temperature range, the electron-phonon scattering term can be neglected. The Kondo contribution to the resistance is described as

$$R_K(T) = R_{K0} \left[1 + (2^{1/s} - 1) \left(\frac{T}{T_K} \right)^2 \right]^{-s}, \quad (3)$$

where $s = 0.225$ is fixed for a spin-1/2 Kondo scatterer, and the Kondo temperature T_K is defined by the condition $R_K(T_K) = R_{K0}/2$ [4,54]. Fits were performed for the PPMS data of sample A above 1.8 K, and the fitted parameters were used to extrapolate the expected resistance behavior to the millikelvin range.

The dashed line in Fig. 4(a) shows the fit obtained above 1.8 K by considering only the Kondo effect on top of the metallic background. The Kondo-only fit appears to reproduce the overall temperature dependence of $R(T)$, but it shows a systematic deviation near 3 K and fails to reproduce the observed resistance decrease.

We test whether WAL can be responsible for this low-temperature resistance decrease. In a two-dimensional diffusive system, quantum interference effects induce a logarithmic temperature dependence of the conductance [12]. For WAL, this contribution leads to an increase in conductance upon cooling [10], which we can introduce phenomenologically at the resistance level as

$$R_{\text{WAL}}(T) = A_{\text{WAL}} \ln T, \quad (4)$$

where $A_{\text{WAL}} > 0$ sets the magnitude of the WAL-induced correction.

The total resistance, including both the Kondo contribution and the WAL correction, can thus be written as

$$R(T) = R_{\text{metal}}(T) + R_K(T) + R_{\text{WAL}}(T). \quad (5)$$

The solid line in Fig. 4(a) reproduces the resistance decrease below 3 K that is missed by the Kondo-only fit. Including a $\ln T$ -type WAL term gives a better description of the zero-field $R(T)$. The fitting parameters for all magnetic field configurations analyzed in this section are summarized in Table I.

Below 800 mK, a much more rapid resistance decrease is observed than predicted by the Kondo + WAL model. As already seen in Fig. 1(a), sample A shows a superconducting transition, so the most likely origin of this additional decrease is superconductivity. However, given that onset $T_c = 390$ mK, it is not reasonable to interpret the resistance decrease starting at 800 mK as the transition itself. Unlike bulk SrTiO₃ [2,3], where the superconducting transition is relatively sharp, fluctuation effects are enhanced as the dimensionality is reduced [55], and fluctuation-driven resistance reductions above T_c have been quantitatively modeled with Aslamazov-Larkin and Maki-Thompson contributions in SrTiO₃-based 2DEGs [56]. The resistance-decrease residual below 800 mK can thus be understood consistently in terms of superconducting fluctuation effects, although quantitative analysis is not possible due to the limited temperature range and overlap with the WAL and Kondo contributions.

Next, we examine the magnetic field response of $R(T)$. If the low-temperature transport is governed by Kondo scattering and two-dimensional WAL, the two contributions should show different responses to an applied magnetic field. The Kondo contribution is expected to be insensitive to the field orientation [30], whereas the WAL contribution is strongly affected by orbital dephasing [12]. The WAL contribution should thus be suppressed by a perpendicular magnetic field because of Aharonov-Bohm phase accumulation, but remains comparatively robust under an in-plane magnetic field. The analysis of $R(T)$ under different magnetic field orientations thus helps identify the WAL contribution and its dimensionality.

Figure 4(b) shows $R(T)$ of sample A at $B_{\perp} = 1$ T. Unlike the zero-field $R(T)$, the data obtained under a perpendicular magnetic field are well described by the Kondo effect alone down to about 200 mK, without any additional WAL contribution. Below about 200 mK, another $\ln T$ -type upturn appears when both the superconducting state and the WAL contribution are suppressed by the applied $B_{\perp} = 1$ T field. This behavior suggests that another low-temperature quantum correction appears that is robust in an applied field but has a different form of $\ln T$ dependence. The detailed origin of this correction is discussed in Sec. III D.

The Kondo + WAL (solid line) and Kondo-only (dashed line) fits for an in-plane field $B \parallel I = 1$ T are shown in Fig. 4(c). Even under an in-plane magnetic field, the Kondo-only fit fails to reproduce the resistance decrease below 3 K and a WAL term is still required to describe the data over the full temperature range, although the necessary correction is diminished, as shown in the inset. The low-temperature decrease is weaker than at zero field, as shown by the fitting

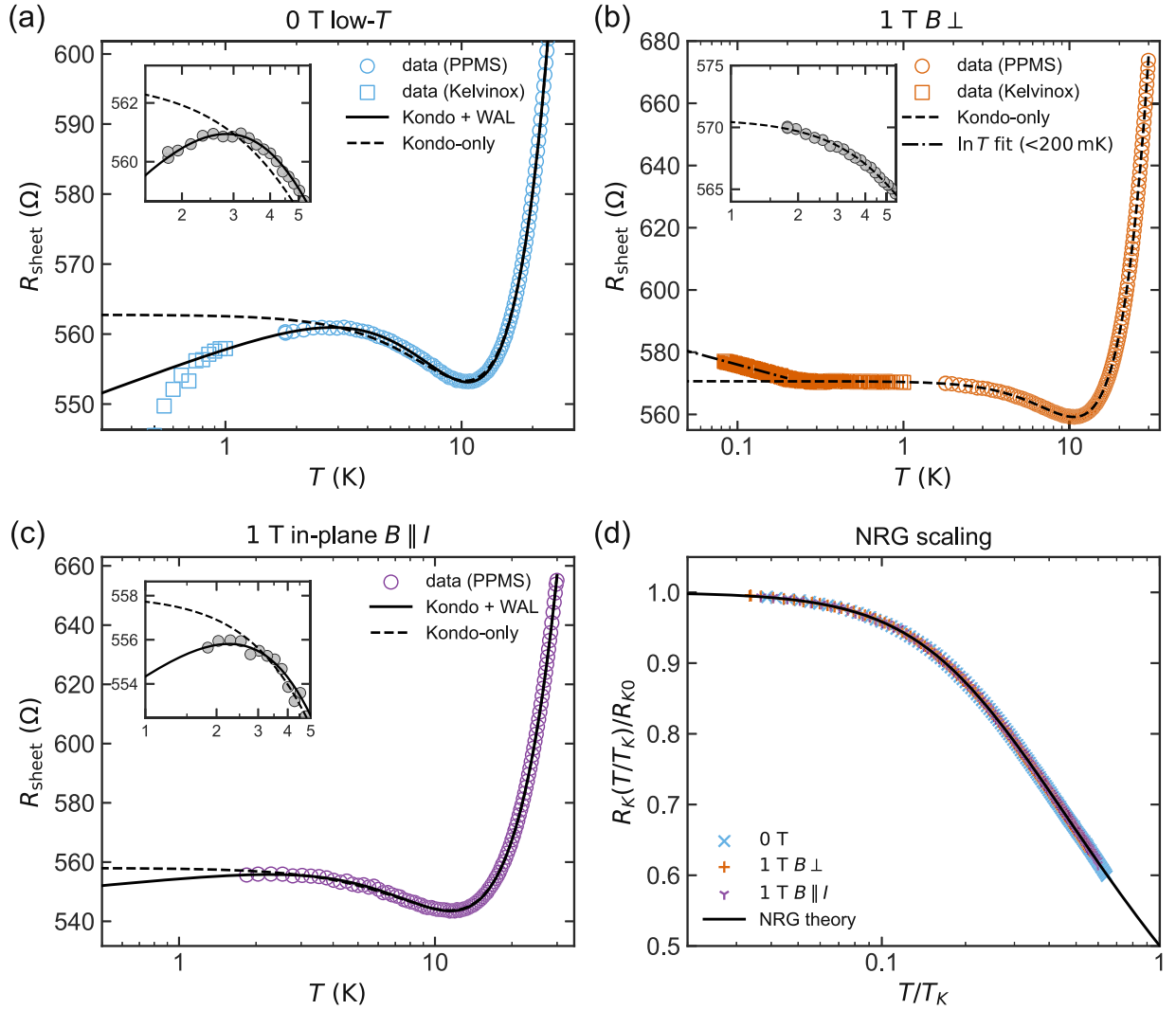


FIG. 4. Temperature dependence of the sheet resistance $R(T)$ of sample A measured under different magnetic field configurations together with the corresponding fits. Open circles denote data measured above 1.8 K using a PPMS, and open squares denote data measured below 1 K using a Kelvinox system. Solid and dashed lines represent fits using the Kondo+WAL and Kondo-only models, respectively. Insets show enlarged views of the data and fits. (a) $R(T)$ measured at $B = 0$. (b) $R(T)$ measured under a perpendicular field $B_{\perp} = 1$ T. The dash-dotted line shows the $\ln T$ fit below 200 mK, extrapolated to 0 K. (c) $R(T)$ measured under an in-plane field $B \parallel I = 1$ T. (d) Universal scaling of the Kondo contribution $R_K(T/T_K)/R_{K0}$ compared with the numerical renormalization group (NRG) result. Crosses, plus symbols, and downward triangles correspond to $B = 0$, $B_{\perp} = 1$ T, and $B \parallel I = 1$ T, respectively.

results for A_{WAL} in Table I, dropping from 5.5 at 0 T to 3.7 at $B \parallel I = 1$ T.

TABLE I. Fitting parameters obtained for Sample A under different magnetic field configurations. The Kondo+WAL model is applied for $B = 0$ and $B \parallel I = 1$ T, and the Kondo-only model for $B_{\perp} = 1$ T where WAL is suppressed.

Parameter	$B = 0$ T	$B_{\perp} = 1$ T	$B \parallel I = 1$ T
R_0 (Ω)	270.9 ± 1.3	265.1 ± 1.4	211.9 ± 2.8
R_{K0} (Ω)	287.3 ± 1.3	305.6 ± 1.4	342.8 ± 2.8
T_K (K)	46.5 ± 0.3	53.3 ± 0.2	50.0 ± 0.5
A_{WAL}	5.5 ± 0.2	—	3.7 ± 0.3
a ($\Omega \text{ K}^{-2}$)	0.2289 ± 0.0005	0.2405 ± 0.0003	0.2447 ± 0.0009

This result shows that the WAL $\ln T$ correction remains relatively robust under an in-plane magnetic field but is not completely unaffected. The $R(T)$ analysis alone is therefore not sufficient to identify the origin of this correction as being caused by WAL. We test the origin and dimensionality of the $\ln T$ term by magnetoresistance analysis in Sec. III C.

Finally, Fig. 4(d) shows that the renormalized Kondo components extracted from the $R(T)$ fitting results at zero field, perpendicular field, and in-plane field all collapse onto the numerical renormalization group (NRG) theoretical curve. The NRG theory predicts that the Kondo contribution follows a universal scaling curve as a function of the reduced temperature T/T_K [54,57,58]. The agreement here shows that the fitted term captures the Kondo contribution rather than a simple phenomenological $\ln T$ dependence.

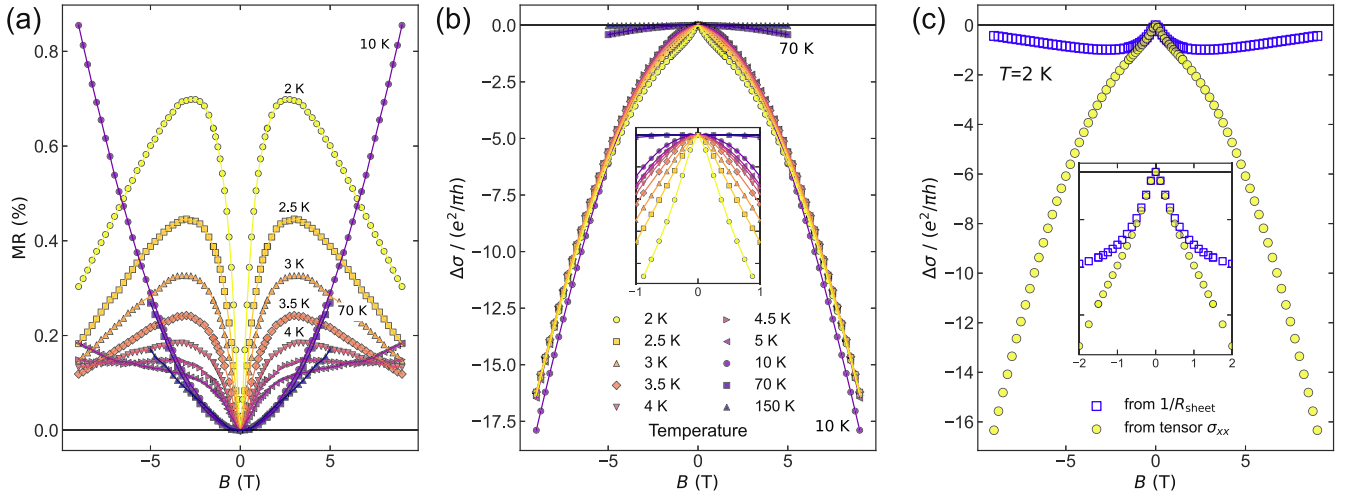


FIG. 5. (a) Magnetoresistance of sample A measured under an out-of-plane magnetic field at different temperatures. The temperature labels are given in panel (b). (b) Normalized magnetoconductance $\Delta\sigma(B)$ at different temperatures. The inset shows an enlarged view of the low-field region. (c) Comparison at 2 K of the normalized conductance obtained from $1/R_{\text{sheet}}$ and from the tensor-defined conductivity σ_{xx} . The inset shows the low-field region.

C. Identification of quantum corrections by magnetoresistance measurements

1. Perpendicular-field magnetotransport

A perpendicular magnetic field provides a direct transport probe of WAL because the magnetic field destroys the electron phase coherence and restores the resistance reduced by the WAL. Figure 5(a) shows the perpendicular magnetoresistance of sample A at different temperatures. The magnetoresistance is parabolic at high temperature but a characteristic cusp signature of WAL appears near zero magnetic field at lower temperatures due to the enhancement of quantum corrections as the electron phase coherence length increases. The magnetoresistance curves in which the WAL cusp appears can be used to extract the strength of the SOI and the degree of electron phase coherence. The γ - $\text{Al}_2\text{O}_3/\text{SrTiO}_3$ heterostructures contain no heavy elements, which means that the atomic SOI would be expected to be weak [59]. Due to the large internal electric field at the interface, however, a strong Rashba-type SOI is generally present in SrTiO_3 heterostructures [60] and its presence can be probed directly by photoelectron spectroscopy [61] or by perpendicular magnetoresistance analysis [14,15]. Various theoretical quantum correction models exist for the analysis of the magnetoconductance behavior of two-dimensional electron systems [62–64]. The normalized magnetoconductance was therefore calculated from the full resistivity tensor. The tensor definition, fitting model, parameter definitions, and fitting results are presented in Appendix A.

Figure 5(b) shows the normalized magnetoconductance $\Delta\sigma(B)$ obtained from this tensor definition. Unlike the magnetoresistance in Fig. 5(a), the magnetoconductance obtained from the tensor definition shows a more emphasized parabolic background, but the sharp cusp still remains, as shown in the inset. The reason for using the conductivity tensor analysis can be seen by comparing the magnetoresistance data for 10 K and 70 K in Fig. 5(a) and the corresponding normalized magnetoconductance in Fig. 5(b). The magnetoresistance plots in

Fig. 5(a) for 10 K and 70 K show almost overlapping parabolic forms. However, the normalized magnetoconductance plots in Fig. 5(b) do not overlap at all. This difference arises because, at low carrier densities, the Hall resistance R_{xy} is not negligibly small and strongly affects the longitudinal conductivity. Figure 5(c) shows a comparison of the magnetoconductance obtained from the tensor definition and the magnetoconductance obtained from the inverse of the resistance at 2 K, where the WAL analysis is carried out. A difference appears already in the low-field region below 1 T, and a completely different form is clearly seen at high field. We emphasize that in WAL analysis of samples with low carrier densities, the magnetoconductivity must be treated as a tensor and fitting simple inverse resistance data does not yield consistent results.

The normalized magnetoconductance was analyzed using the Maekawa-Fukuyama (MF) model [14,64] with an added Kohler term [15] for the orbital effect. Figure 6(a) shows the normalized magnetoconductance of sample A measured from 2 K to 4 K over the full magnetic field range, together with the MF model fits shown with solid lines. The model explains the data up to the high-field range of 9 T without structured residuals. The low-field fitting within 1 T in the inset is also well reproduced, which shows that the low-field region containing the WAL information is not missed.

The fitting parameters, the characteristic fields B_{s0} and B_{ϕ} , and the corresponding spin relaxation time τ_{s0} and phase coherence time τ_{ϕ} obtained from these characteristic fields are shown in Figs. 6(b) and 6(c). B_{s0} and B_{ϕ} are inversely proportional to τ_{s0} and τ_{ϕ} , respectively. In Fig. 6(b), B_{s0} shows no clear trend with temperature, whereas B_{ϕ} decreases slightly upon cooling. This trend appears more clearly in the corresponding scattering time plots in Fig. 6(c). The spin relaxation time τ_{s0} remains almost constant and independent of temperature, whereas the phase coherence time τ_{ϕ} increases clearly as temperature decreases. This is consistent with SOI being determined by the confining potential asymmetry of the heterostructure while the electron phase coherence improves as the thermal energy decreases upon cooling.

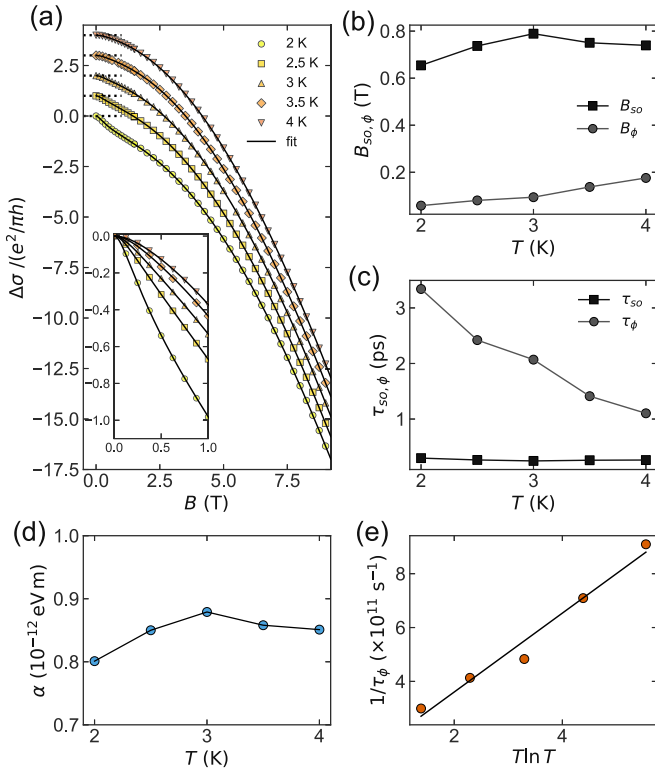


FIG. 6. (a) Out-of-plane magnetoconductance over the full magnetic field range with weak antilocalization fits. For clarity, the traces are vertically offset. The black dotted line indicates the zero-conductance reference, $\Delta\sigma = 0$, at zero magnetic field. The inset shows the low-field region, highlighting the WAL cusp. (b) Temperature dependence of the characteristic fields B_{so} and B_ϕ extracted from the fits. Error bars are smaller than the marker size. (c) Corresponding spin relaxation and phase coherence times, τ_{so} and τ_ϕ . (d) Temperature dependence of the Rashba coupling constant α . (e) Inverse phase coherence time $1/\tau_\phi$ as a function of $T \ln T$, together with a linear fit.

The Rashba coupling constant α in Fig. 6(d) also shows almost no temperature dependence and is about 8×10^{-13} eV m at 2 K. This is similar to previously reported values for the $\text{LaAlO}_3/\text{SrTiO}_3$ interfaces [14,65,66]. As shown in Fig. 6(e), the inverse phase coherence time $1/\tau_\phi$ shows a linear dependence on $T \ln T$. This behavior is consistent with the characteristic Nyquist electron-electron scattering in a diffusive two-dimensional electron system and shows that the WAL in sample A is a two-dimensional effect [67,68].

2. In-plane magnetoresistance and coexistence of WAL and Kondo contributions

Under an in-plane magnetic field, the Kohler-type orbital magnetoresistance is strongly suppressed. In this case, the in-plane magnetoresistance, where the magnetic field is parallel to the current, can contain a negative component due to a slight suppression of the Kondo effect by spin alignment [4] and a positive component caused by the suppression of WAL either by the Zeeman effect [69] or by Aharonov-Bohm phase accumulation due to the finite thickness of the 2DEG layer [70]. Because the two contributions have different signs, the

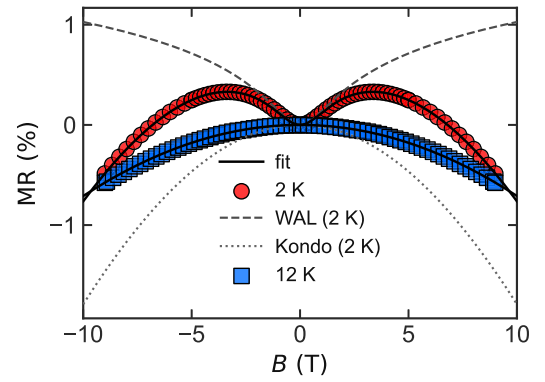


FIG. 7. In-plane magnetoresistance ($B \parallel I$) measured at 2 K and 12 K. Solid lines represent the total response, including both weak antilocalization (WAL) and Kondo contributions at 2 K and only the Kondo contribution at 12 K. Dashed and dotted lines show the WAL and Kondo components, respectively, and are presented only for the 2 K data.

form of the in-plane magnetoresistance can be used to confirm the presence of WAL and Kondo scattering independently of the $R(T)$ experiment. It also allows the two-dimensional quantum corrections to be separated and interpreted under conditions where the classical orbital background is strongly reduced. The in-plane magnetoresistance was fitted using a phenomenological model including WAL and Kondo contribution terms. The fitting model, parameter definitions, and fitting results are summarized in Appendix B.

Figure 7 shows the in-plane magnetoresistance measured at 2 K and 12 K. The magnetoresistance at 2 K increases at low field and then decreases at high field, whereas at 12 K only negative magnetoresistance is observed. Considering the possible resistance changes in the in-plane magnetoresistance, this can be interpreted as follows. At 2 K, the WAL component counteracts the Kondo scattering, producing a nonmonotonic magnetoresistance field dependence. At 12 K, no clear WAL component can be resolved, and the magnetoresistance is described mainly by the Kondo component. The solid fit lines in Fig. 7 show that the WAL and Kondo-suppression model can phenomenologically describe well the magnetoresistance at 2 K and 12 K.

This analysis shows that the temperature range where the Kondo effect appears is higher than the temperature range where WAL becomes pronounced. At 2 K, both effects affect the in-plane magnetoresistance, but at 12 K the system leaves the quantum-diffusive regime because of thermal decoherence, the WAL contribution becomes weak, and only the Kondo effect is measurably affected by the in-plane magnetic field.

The in-plane magnetoresistance of $\gamma\text{-Al}_2\text{O}_3/\text{SrTiO}_3$ is different from the behavior of the $\text{LaAlO}_3/\text{SrTiO}_3$ and $\text{LaTiO}_3/\text{SrTiO}_3$ interfaces, where the in-plane magnetoresistance shows giant negative magnetoresistance (GNMR) with resistance reductions often exceeding 50% [71–73]. In the $\text{LaAlO}_3/\text{SrTiO}_3$ and $\text{LaTiO}_3/\text{SrTiO}_3$ interfaces, this has been understood to originate from a Lifshitz transition when both the d_{xy} and d_{xz}/d_{yz} subbands are occupied [72,73]. The absence of GNMR at the present interface independently supports, from the transport side, the XLD result that no Lifshitz

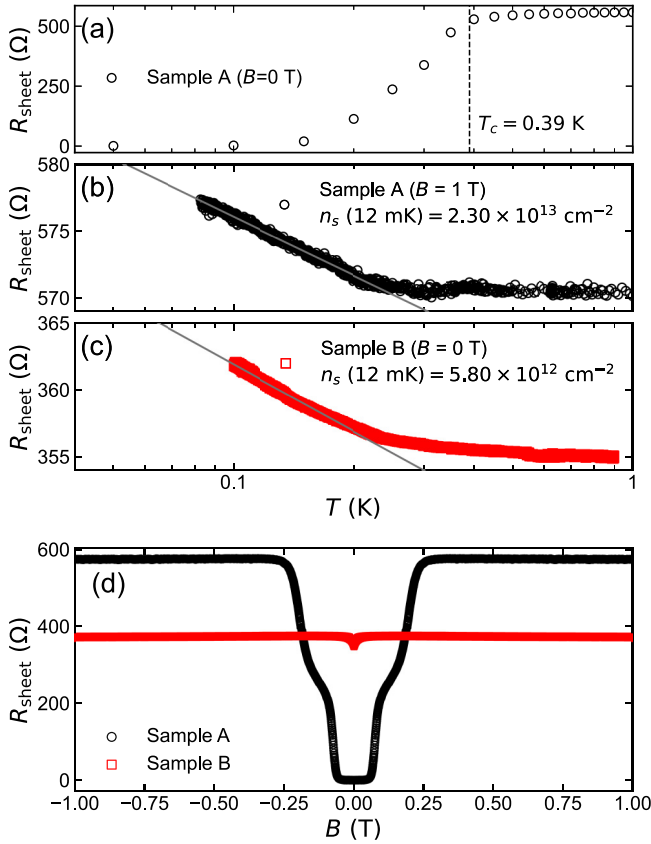


FIG. 8. Temperature dependence of the sheet resistance $R(T)$ of sample A at $B = 0$ (a), showing a superconducting transition with an onset $T_c = 0.39$ K, sample A at $B_{\perp} = 1$ T (b), and sample B at zero applied field (c). The gray solid lines in (b) and (c) are $\ln T$ fits below 200 mK. (d) Magnetoresistance of samples A and B measured at 12 mK.

transition occurs within the carrier density range accessed in this study.

D. Two-dimensional EEI-dominated transport below 200 mK

In Sec. III B, a $\ln T$ -type resistance upturn was observed below 200 mK in sample A under a perpendicular magnetic field. To examine the origin of this ultralow-temperature quantum correction under conditions where superconductivity and WAL are suppressed, we show in Fig. 8 the sub-1 K $R(T)$ response and the 12 mK $R(B)$ response of both samples A and B. For sample A, the normal state is obtained by applying a 1 T perpendicular magnetic field, while sample B does not show a superconducting transition even at ultralow temperature because of its lower carrier density.

Figure 8(a) shows that sample A undergoes a superconducting transition at $T_c = 0.39$ K in zero magnetic field, which makes direct access to the normal-state resistance at the lowest temperature impossible. When a perpendicular magnetic field of $B_{\perp} = 1$ T is applied to suppress the superconducting transition, a resistance upturn proportional to $\ln T$ appears below 200 mK after the Kondo resistance saturation, as shown in Fig. 8(b).

To confirm that this resistance upturn is not induced by the magnetic field, we show in Fig. 8(c) the $R(T)$ of sample

B, which remains in the normal state but shows the same resistance upturn below 200 mK, even in zero magnetic field.

A comparison of the 12 mK magnetoresistance data for the two samples in Fig. 8(d) shows that no anomalous field-dependent conductivity corrections remain once the resistance reductions from superconductivity and WAL are removed by a magnetic field. This indicates that the most likely cause of the $\ln T$ -type correction observed below 200 mK is the electron-electron interaction correction, which is relatively insensitive to magnetic field, and other magnetic-field-sensitive quantum-interference corrections can be ruled out.

The electron-electron interaction is expected to produce a $\ln T$ -type correction to conductivity in a disordered two-dimensional system [12]. Indeed, a $\ln T$ -type temperature dependence of resistance that is relatively robust against magnetic field in the liquid helium temperature range has been widely observed in various two-dimensional systems, including Si MOSFET inversion layers, where this behavior has been interpreted as a signature of EEI [74–77]. A similar low-temperature resistance upturn has recently been observed in Pb monolayers on Si(111), supporting the interpretation that the ultralow-temperature resistance upturn can be a normal-state transport correction in a two-dimensional electron system [78,79].

A possible explanation as to why this correction becomes pronounced only below 200 mK is a dimensional crossover related to the finite thickness of the conducting layer in γ -Al₂O₃/SrTiO₃. In the diffusive regime, the EEI correction is governed by the thermal diffusion length $L_T = \sqrt{\hbar D/k_B T}$ [80], where D is the diffusion constant (Appendix A). As discussed by Lee and Ramakrishnan [12], films of finite thickness are expected to show a dimensionality crossover on a temperature scale determined by the thickness. Experimentally, the emergence of a $\ln T$ resistance upturn with decreasing layer thickness has also been reported in thin metallic systems, where it was interpreted as a two-dimensional EEI correction associated with dimensionality [81].

The particularly low temperature scale in the present interface can be related to transport occurring in the relatively heavy d_{xz}/d_{yz} subbands at the γ -Al₂O₃/SrTiO₃ interface. Such heavy-electron subbands have a smaller diffusion constant, which means that L_T becomes shorter and the onset of observable two-dimensional EEI behavior is shifted to lower temperature.

For example, in LaTiO₃/SrTiO₃ heterostructures, where the 2DEG is derived from the light-electron d_{xy} orbital [82], a similar $\ln T$ -type correction to conductivity has been reported to be observable already at 4 K [83]. In contrast, at the γ -Al₂O₃/SrTiO₃ interface, where the heavy d_{xz}/d_{yz} subbands lie at lower energy, a clear $\ln T$ -type resistance upturn appears only below about 200 mK.

IV. CONCLUSION

We have investigated the ultralow-temperature transport of a nonmagnetic, light-element-based γ -Al₂O₃/SrTiO₃ two-dimensional electron gas where the carrier density was tuned by adjusting the heterostructure growth conditions. The strong interfacial Rashba spin-orbit interaction was found to

induce a pronounced two-dimensional weak antilocalization transport correction, which contributes not only to the magnetoresistance but also to the temperature-dependent resistance together with the usual Kondo scattering from localized spins.

In zero magnetic field, the transport evolves from a Kondo-dominated regime around 10 K through a WAL-induced resistance reduction near 3 K and finally enters the superconducting state below 390 mK. In contrast, under a perpendicular magnetic field, both WAL and superconductivity are suppressed, revealing a two-dimensional electron-electron interaction correction on top of the saturated low-temperature Kondo contribution. The separation of these low-temperature corrections was found to be consistent with the electronic structure of the γ -Al₂O₃/SrTiO₃ interface.

Specifically, transport is dominated by the relatively heavy d_{xz}/d_{yz} subbands, which give rise to low carrier mobility and a reduced diffusion scale, thereby separating the WAL and EEI regimes on the temperature axis. Finally, we emphasize that in the low-carrier-density regime that can be reached in the γ -Al₂O₃/SrTiO₃ 2DEG, the Hall contribution cannot be neglected in magnetoconductance analysis.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX A: PERPENDICULAR-FIELD MAGNETOCONDUCTANCE

To analyze the perpendicular-field magnetoconductance quantitatively, we use the Maekawa-Fukuyama model [14,64], which is adopted here as the fitting form for the WAL contribution in the presence of SOI and Zeeman splitting. We analyze the normalized magnetoconductance, defined as $\Delta\sigma(B)/G_0 = [\sigma(B) - \sigma(0)]/G_0$, where $G_0 = e^2/(\pi h) \approx 1.2 \times 10^{-5}$ S is the quantum of conductance used in the present analysis.

We obtain the conductivity tensor from the full resistivity tensor, defined in terms of the longitudinal and Hall resistivities, ρ_{xx} and ρ_{xy} , as follows:

$$\rho = \begin{pmatrix} \rho_{xx} & \rho_{xy} \\ -\rho_{xy} & \rho_{xx} \end{pmatrix}, \quad \sigma = \rho^{-1}. \quad (\text{A1})$$

This yields

$$\sigma_{xx} = \sigma_{yy} = \frac{\rho_{xx}}{\rho_{xx}^2 + \rho_{xy}^2}, \quad \sigma_{xy} = -\sigma_{yx} = -\frac{\rho_{xy}}{\rho_{xx}^2 + \rho_{xy}^2}. \quad (\text{A2})$$

The magnetoconductance is written as the sum of the WAL contribution described by the Maekawa-Fukuyama model and a classical background term,

$$\Delta\sigma(B) = \Delta\sigma_{\text{WAL}}(B) + \Delta\sigma_{\text{Kohler}}(B), \quad (\text{A3})$$

where $\Delta\sigma_{\text{WAL}}(B)$ is the WAL contribution and $\Delta\sigma_{\text{Kohler}}(B)$ is a phenomenological classical orbital background term.

The WAL contribution is given by

$$\begin{aligned} \frac{\Delta\sigma_{\text{WAL}}(B)}{G_0} = & \Psi\left(\frac{B}{B_\phi + B_{\text{so}}}\right) \\ & + \frac{1}{2\sqrt{1-\gamma^2}} \Psi\left(\frac{B}{B_\phi + B_{\text{so}}(1 + \sqrt{1-\gamma^2})}\right) \\ & - \frac{1}{2\sqrt{1-\gamma^2}} \Psi\left(\frac{B}{B_\phi + B_{\text{so}}(1 - \sqrt{1-\gamma^2})}\right), \end{aligned} \quad (\text{A4})$$

where $\Psi(x) = \ln x + \psi(\frac{1}{2} + \frac{1}{x})$ and ψ is the digamma function.

The Kohler-type background is written as

$$\frac{\Delta\sigma_{\text{Kohler}}(B)}{G_0} = -\frac{\sigma(0)}{G_0} \frac{A_K B^2}{1 + C B^2}, \quad (\text{A5})$$

where A_K and C are fitting parameters [15].

In the present analysis, we fix $g = 2$ and $m^* = 3m_e$ [59], where g is the Landé g factor, m_e is the free-electron mass, and m^* is the effective mass. The characteristic fields associated with phase coherence and spin relaxation are defined as $B_\phi = \hbar/(4eD\tau_\phi)$ and $B_{\text{so}} = \hbar/(4eD\tau_{\text{so}})$, where τ_ϕ is the phase coherence time and τ_{so} is the spin relaxation time.

The diffusion constant is given by $D = \frac{1}{2}v_F^2\tau$, where v_F is the Fermi velocity and τ is the elastic scattering time. Using the carrier density n_s , the Fermi wave vector and Fermi velocity are written as $k_F = \sqrt{2\pi n_s}$ and $v_F = \hbar k_F/m^*$, and the elastic scattering time is obtained from the mobility μ as $\tau = \mu m^*/e$.

The Zeeman contribution is included through the dimensionless parameter $\gamma = g\mu_B B/(4eDB_{\text{so}})$, where μ_B is the Bohr magneton. Within the D'yakonov-Perel spin relaxation mechanism, the spin relaxation time τ_{so} and the Rashba coupling constant α are related by $\tau_{\text{so}} = \hbar^4/(4\alpha^2 m^{*2} D)$ [14]. The Rashba coupling constant is therefore obtained as $\alpha = \sqrt{e\hbar^3 B_{\text{so}}/m^*}$.

Using these definitions, we fit the perpendicular-field magnetoconductance and extract the phase coherence and Rashba coupling constants. The fitting results are summarized in Table II.

APPENDIX B: IN-PLANE MAGNETORESISTANCE

To analyze the in-plane magnetoresistance, we decompose the measured magnetoresistance into WAL and Kondo

TABLE II. Fitting parameters for sample A at different temperatures.

T (K)	B_{so} (T)	B_{ϕ} (T)	A_K (10^{-3} T $^{-2}$)	C (10^{-3})
2.0	0.65 ± 0.00	0.06 ± 0.00	1.57 ± 0.00	1.32 ± 0.01
2.5	0.74 ± 0.01	0.08 ± 0.00	1.55 ± 0.00	1.29 ± 0.01
3.0	0.79 ± 0.03	0.09 ± 0.00	1.54 ± 0.00	1.29 ± 0.01
3.5	0.75 ± 0.03	0.14 ± 0.01	1.56 ± 0.00	1.34 ± 0.01
4.0	0.74 ± 0.03	0.18 ± 0.01	1.58 ± 0.01	1.39 ± 0.01

contributions, which are treated here as independent field-dependent corrections. The fitting forms used for each contribution are summarized below.

1. In-plane suppression of WAL

To capture the low-temperature contribution observed in the in-plane magnetoresistance, we model the suppression of quantum interference by an in-plane magnetic field using a logarithmic field dependence. In a two-dimensional diffusive metal, an in-plane magnetic field suppresses the phase-coherent quantum interference through Zeeman-induced dephasing [69] and Aharonov-Bohm phase accumulation associated with the finite thickness of the conducting channel [70], leading to an effective dephasing rate proportional to B^2 and a logarithmic field dependence of the interference correction.

Accordingly, the resistance change can be written phenomenologically as

$$\Delta R_{\parallel}^{\text{WAL}}(B) \propto \ln(1 + aB^2), \quad (\text{B1})$$

where a sets the magnitude of the dephasing scale. This form has been used in analyses of in-plane magnetoresistance in SrTiO₃-based 2DEGs [84,85].

2. Field dependence of the Kondo contribution

To describe the field suppression of the Kondo contribution in the in-plane magnetoresistance, we use the zero-temperature Bethe ansatz result for spin-1/2 Kondo scatterers as a phenomenological description for the low-temperature regime considered here. This provides the relation between

TABLE III. Fitting parameters obtained from the in-plane magnetoresistance analysis at 2 K and 12 K.

Parameter	2 K	12 K
$R_K(0)/R(0)$	0.508	0.427
a	0.822 ± 0.039	–
B_K (T)	27.91 ± 0.14	68.35 ± 13.51

the scatterer moment and the Kondo magnetoresistance in a magnetic field. The magnetization is written in terms of $x = B/B_K$ as

$$M(x) =$$

$$\begin{cases} \frac{1}{\sqrt{2\pi}} \sum_{k=0}^{\infty} \left(-\frac{1}{2}\right)^k (k!)^{-1} \left(k + \frac{1}{2}\right)^{k-\frac{1}{2}} e^{-(k+\frac{1}{2})x} x^{2k+1}, & x \leq \sqrt{2} \\ 1 - \pi^{-3/2} \int_0^{\infty} \frac{dt}{t} \sin(\pi t) e^{-t \ln(t/2e)} x^{-2t} \Gamma\left(t + \frac{1}{2}\right), & x \geq \sqrt{2}, \end{cases} \quad (\text{B2})$$

where B_K is the characteristic field scale of the Kondo contribution, related to the Kondo temperature and the localized spin g factor. Γ denotes the gamma function [4,30,84].

The impurity magnetization then determines the field-dependent Kondo magnetoresistance through

$$\frac{R(B_{\parallel}) - R(0)}{R(0)} = -\frac{R_K(0)}{R(0)} \left[1 - \cos^2\left(\frac{\pi}{2} M\left(\frac{B_{\parallel}}{B_K}\right)\right) \right], \quad (\text{B3})$$

where $R_K(0)$ is the zero-field Kondo contribution to the resistance, and $R(0)$ is the total zero-field resistance [84]. The $R_K(0)/R(0)$ ratio is fixed to the fitting parameters obtained from the zero-field $R(T)$ analysis in Sec. III B using the Kondo + WAL model.

In fitting the in-plane magnetoresistance, this zero-temperature expression is used as a phenomenological description of the field suppression of the Kondo contribution in the low-temperature regime. The resulting fitting parameters are summarized in Table III.

Although EEI corrections are, in principle, also possible in a two-dimensional system, we do not see such an effect in the temperature range where the Kondo contribution is significant. Therefore, above 2 K, the in-plane magnetoresistance is described primarily by the magnetic field suppression of quantum interference effects and Kondo scattering.

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