Interferometry and coherent single-electron transport through hybrid superconductor-semiconductor Coulomb islands

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Majorana zero modes are leading candidates for topological quantum computation due to their non-local character and non-abelian exchange statistics 1,2. Among their attributes, spatially separated Majorana modes are expected to allow coherent single-electron transport through one-dimensional topological superconductors in the Coulomb blockade (CB) regime 3. We have investigated this feature by patterning an elongated epitaxial InAs-Al Coulomb island embedded in an Aharonov-Bohm interferometer. Using a parallel magnetic field to lower the energy of a discrete sub-gap state in the island below its charging energy, conductance oscillations in the ring were observed with a flux period of h/e (h is Planck’s constant and e is the elementary charge), indicating coherent single-electron transport through the interferometer. Oscillation amplitude was largest when CB conductance peaks in the island were 1e periodic, and suppressed when CB was 2e periodic or when superconductivity was suppressed. Oscillation phase shifts of π were observed when the charge occupancy of the island was changed by 1e, indicating that the interferometer can detect island parity. Magnetic fields oriented orthogonal to the island reduced the field at which 2e periodic peaks split and where coherent transport could also be observed, suggesting additional non-Majorana mechanisms for 1e transport through these moderately short wires.

Initial experiments reporting signatures of Majorana zero modes (MZMs) in hybrid superconductor-semiconductor nanowires focussed on zero-bias conductance peaks (ZBPs) using local tunnelling spectroscopy 4–7. Subsequently, Majorana islands in the Coulomb blockade (CB) regime provided additional evidence of MZMs based on nearly 1e-spaced CB peaks 8, and indicated a Rashba-like spin orbit coupling (SOC) with spin-orbit field lying in-plane, perpendicular to the wire axis 9. Under some circumstances, these signatures can be mimicked by trivial modes 10,11, motivating a new generation of experiments that explicitly probe non-local properties, which are more difficult to mimic. For instance, non-locality of MZMs was recently investigated by measuring the energy splitting induced by the interaction of a quantum dot and a zero-energy state in a hybrid nanowire 12.

Non-locality can also be accessed by interferometric measurements of a Majorana island, where CB couples separated MZMs and fixes fermion parity 3,13–16. Single electrons tunnelling through a Majorana island are expected to be coherently transferred between MZMs via a process referred to as electron teleportation 9. Coherent single-electron tunnelling through a CB island is a signature of non-locality, though not necessarily of MZMs, which require other signatures as well, as discussed in detail below. For example, the predicted absence of phase lapses through a CB peak in the transmission through the interferometer is a specific signature of MZMs 16. To probe the non-locality of Majorana modes, a Majorana island can be embedded in the arm of an Aharonov-Bohm (AB) interferometer. If single-electron transport in both the reference arm and the Majorana island is coherent, conductance through the interferometer is expected to show magnetic flux modulations with period h/e 17. The same approach was used to investigate coherent transport in semiconductor quantum dots 18,21.

Devices were fabricated using an InAs two-dimensional electron gas (2DEG) heterostructure covered by 8 nm of epitaxially grown Al 22. The bare 2DEG (without Al) showed a phase coherence length of lϕ ~ 4 μm (see Supplementary Fig. S.1). Figure 1 shows a micrograph of device 1 with a 1.2 μm long and 0.1 μm wide superconducting Al stripe formed by wet etching. Ti/Au top-gates were evaporated on top of a 25 nm HfO2 dielectric grown by atomic layer deposition. We studied two lithographically similar interferometers with circumferences
of 5.6 µm for device 1 and 5 µm for device 2.

Applying a negative voltage, \( V_{\text{pg}} \), to the central gate serves two purposes. It depletes the 2DEG surrounding the Al wire to form both the Coulomb island and the AB ring center and also adjusts the chemical potential and charge occupancy of the island. Energizing all exterior gates confines the 2DEG into an AB interferometer by connecting the Coulomb island to a normal conducting reference arm. The resistance of the reference arm was controlled by a negative gate voltage \( V_{\text{pg}} \) (Fig. 1a). Figure 1b shows zero-bias differential conductance \( G \) as a function of \( B_{\parallel} \) and \( V_{\text{pg}} \). Zero-bias differential conductance \( G \) appears as a function of \( V_{\text{sd}} \) and \( V_{\text{pg}} \) showing Coulomb diamonds for \( B_{\parallel} = 1 \, \text{T} \) (c), 2 T (d), 2.5 T (e), and 3.3 T (f). The measurements shown in panels b-f were taken with the reference arm closed.

At low temperatures, tunneling of single electrons onto an island with a superconducting gap \( \Delta \) is suppressed by CB, except at charge degeneracies. When the lowest sub-gap state energy, \( E_{0} \), exceeds the charging energy \( E_{c} \), ground-state degeneracies only occur between even-occupied states, resulting in \( 2e \)-periodic CB conductance peaks [23]. Odd-occupied ground states are lowered into the accessible spectrum by a Zeeman field, resulting in even-odd CB peak spacing when \( 0 < E_{0} < E_{c} \). The difference in peak spacings between even and odd states, \( S = S_{\text{e}} - S_{\text{o}} \), is proportional to \( E_{0} \) [8]. For well-separated MZMs, \( E_{0} \) tends exponentially toward zero, yielding \( 1e \) periodic Coulomb peaks [24]. When the MZMs are not widely separated, CB peak spacings oscillate with field and density [8–10].

We first investigated the Coulomb island without interferometry by depleting a segment of the reference arm with gate voltage \( V_{\text{c}} \) (see Fig. 1a). Figure 1b shows zero-bias differential conductance \( G = dI/dV \) of the island as a function of parallel magnetic field \( B_{\parallel} \) and gate voltage \( V_{\text{pg}} \), which controls the electron occupancy of the island. CB peaks are \( 2e \)-periodic at zero field and split around 2 T, becoming \( 1e \) periodic as the sub-gap state moves toward zero energy. Measuring \( G \) as a function of both source-drain bias \( V_{\text{sd}} \) and \( V_{\text{pg}} \) reveals Coulomb diamonds (Fig. 1c-f). At low \( B_{\parallel} \), diamonds are \( 2e \)-periodic with distinct negative differential conductance (Fig. 1d), which transition to an even-odd peak spacing difference at moderate fields (Fig. 1f), similar to previous work on super-
FIG. 2. Evolution of conductance oscillations in parallel magnetic field. Magnetoconductance for parallel field values $B_\parallel = 0, 2.2, 2.7$, and 3.3 T (left to right). a–d, Zero-bias differential conductance $G(B_\perp = 0)$ versus gate voltage $V_{pg}$ used to control electron occupation. e–h, Conductance $\Delta G$ as a function of $V_{pg}$ and perpendicular magnetic field $B_\perp$ controlling the flux in the interferometer with corresponding power spectrums in i–l (solid black lines indicate the frequency window bounding the Aharonov-Bohm oscillations). $\Delta G$ is the conductance with a subtracted slowly varying background. $\Delta V_{pg} = 0$ corresponds to $V_{pg} = -1.896$ V.

coupling Coulomb island [8, 9, 23, 25–28]. At high fields, the 1e periodic diamonds show a discrete zero-bias state for consecutive charge degeneracy points that is well separated from the superconducting gap (Fig. 1f). This sub-gap feature remained at zero bias until the superconducting gap closure, and persists for 3 mV in $V_{pg}$, corresponding to an energy range of 0.8 meV. Below, we show that the ZBP was sensitive to rotations of the in-plane field. Overall, these observations are consistent with a MZM picture. The magnitude of $B_\parallel$ where 1e periodicity is observed is in agreement with ZBPs measured in tunneling spectroscopy in InAs 2DEGs [6]. In contrast, as a function of $B_\parallel$ the peak spacing remained even-odd, and discrete ZBPs on consecutive peaks were not observed in spectroscopy (see Supplementary Fig. S.2). The normal state of the Coulomb island appears above $B_\parallel \sim 3$ T with $E_c = 80 \mu$eV (Fig. 1f).

The reference arm of the AB interferometer was connected by tuning $V_c$ (see in Figs. 2a–d). Figures 2b–h show the conductance $\Delta G$ through the full interferometer (with smooth background subtracted [29]) as a function of $B_\perp$ and gate voltage $V_{pg}$, which controls the occupancy of the island. The power spectrum (PS) of $\Delta G(B_\perp)$ is shown in Fig. 2i–l). Figure 2e shows small oscillations in $\Delta G(B_\perp)$ for the 2e periodic peaks at $B_\parallel = 0$ T. At $B_\parallel = 2.2$ T, where the peak spacing was even-odd (Fig. 2f), the conductance showed a moderate oscillation amplitude with a period $\Delta B_\perp = 1.5$ mT, and a PS signal peaked around $f = 0.65$ mT$^{-1}$. The period $\Delta B_\perp$ is consistent with a single flux quantum $h/e$ piercing the interferometer, indicating coherent transport of single electrons. At higher parallel field, $B_\parallel = 2.7$ T, the CB peak spacing was uniformly 1e, and oscillation amplitude was maximal (see Figs. 2g, k). The oscillation amplitude was largest on the CB peaks and smaller in the valleys. When the Coulomb island was driven normal, $B_\parallel > 3$ T, conductance oscillations were reduced, becoming comparable to oscillation amplitude seen at low parallel field (Fig. 2h, l).

It is interesting to compare the $B_\parallel$ dependence of the AB oscillation amplitude to the corresponding field dependence of the lowest sub-gap state, $E_0(B_\parallel)$, of the island. To quantify the oscillations amplitude, $\langle \tilde{A} \rangle$, we compute the integrated power spectrum over a band of frequencies marked by horizontal lines in Fig. 2 then average over Coulomb peaks and valleys. The sub-gap energy is found from the difference between even and odd CB peak spacings, averaged separately, $\langle S \rangle = \langle S_e \rangle - \langle S_o \rangle$. As seen in Fig. 3h, $\langle S \rangle$ remains constant as a function of $B_\parallel$ (indicating 2e transport) until a sub-gap state moves below $E_c$, reaching zero at 2.2 T without overshoot (as expected for MZMs in a long wire). At low fields, where
FIG. 3. Coulomb peak spacing difference and oscillation amplitude. a,c. Peak spacing difference \( \langle S \rangle \) as a function of parallel magnetic field \( B_{\parallel} \) for devices 1 and 2. b,d. Aharonov-Bohm oscillation amplitude \( \langle \tilde{A} \rangle \) as a function of \( B_{\parallel} \). The solid line is a guide to the eye. Insets show characteristic magnetoconductance \( \Delta G \) as a function of gate voltage \( V_{\text{pg}} \) controlling electron occupancy and perpendicular magnetic field \( B_{\perp} \) controlling the magnetic flux in the interferometer in the 1e regime (indicated by yellow markers in the main panel). Yellow ticks show CB peak positions. \( \Delta V_{\text{pg}} = 0 \) corresponds to \( V_{\text{pg}} = -1.896 \) V and \( -0.945 \) V for b and d, respectively.

FIG. 4. In-plane magnetic field rotations. a-c. Differential conductance \( G \) as a function of gate voltage \( V_{\text{pg}} \) controlling electron occupancy and source-drain bias voltage \( V_{\text{id}} \) showing Coulomb diamonds for in-plane rotation angles of \( \alpha = 0 \) (a), 5° (b), and 10° (c) with \( \alpha = 0 \) corresponding to \( B_{\parallel} = 2.5 \) T. d. Oscillation amplitude \( \langle \tilde{A} \rangle \) as a function of in-plane rotation angle \( \alpha \) for \( B_{\parallel} = 2.5 \) T. The solid line is a guide to the eye.

The CB periodicity was 2e, the oscillation amplitude \( \langle \tilde{A} \rangle \) is small (Fig. 3b). When \( \langle S \rangle \) approaches zero at high fields (>2 T), \( \langle \tilde{A} \rangle \) exhibits a sharp increase that coincides with the 2e to 1e transition. Above 3 T, the Coulomb island was in the normal state and \( \langle \tilde{A} \rangle \) returned to the low value found in the low-field 2e regime.

Figure 3d shows a similar study for device 2. In Fig. 3d, \( \langle S \rangle \) shows strong even-odd below 1 T, fluctuates around \( \langle S \rangle = 0 \) between 1-2 T, then settles to 1e \( \langle \langle S \rangle = 0 \rangle \) above 2 T. CB spectroscopy reveals a discrete state that oscillates around zero energy in both \( B_{\parallel} \) and \( V_{\text{pg}} \) without forming a robust zero-bias peak (see Supplementary Fig. S4). This behaviour is compatible with hybridized MZMs or extended Andreev bound states (ABS) [9]. Figure 3d shows that phase coherent transport first appears above 1 T and \( \langle \tilde{A} \rangle \) gradually increases until reaching a maximum amplitude for 1e peak spacing at 2.1 T, before diminishing in the normal state.

Upon changing the charge occupancy by one electron we found a transmission phase shift of \( \pi \) occurs in the 1e regime (see insets of Fig. 3). For device 2, we observed phase lapses in the CB valley leading to the same side of CB peaks having the same phase, while device 1 showed a more rigid phase in the valleys (also see Supplementary Fig. S5) [19–21]. The pattern of phase lapses, and the difference between devices, is currently not understood and appears to vary depending on gate voltage configurations.

We next investigate the angular dependence of the in-plane magnetic field for device 1. Figure 4a shows 1e Coulomb diamonds at \( B_{\parallel} = 2.5 \) T with a discrete state at each charge degeneracy point (see Fig. 1e). Rotating the magnetic field by an angle \( \alpha = 5^\circ \) (keeping the in-plane field magnitude \( B_{\parallel} = 2.5 \) T constant) lifted the discrete state from zero energy and transitioned the peak spacing to even-odd, whereas at \( \alpha = 10^\circ \) the periodicity was again 1e but without a discrete ZBP. Small rotations (\( |\alpha| < 7.5^\circ \)) reduced the oscillation amplitude, as expected for even-odd periodicity. However, at larger angles (\( |\alpha| > 10^\circ \)) where the discrete zero-bias peak was absent, a strong interference signal was observed (Fig. 4d). An in-plane field-angle sensitivity of the zero-energy mode is consistent with a MZM interpretation. However, the observation of coherent transport in the absence of a discrete zero-energy mode suggests trivial quasiparticles are phase coherent over the length of the island.
To further explore the dependence on in-plane field, we
study the effect of magnetic fields aligned orthogonal to
the wire axis ($B_{\perp}$ and $B_{\parallel}$) on the oscillation amplitude
in device 2. In Fig. 5d–f, the peak spacing difference ($S$)
is shown as a function of the three field directions. For
the two orthogonal fields, $\langle S \rangle$ reaches zero at $\sim 0.2$ T,
in comparison to 1 T for $B_{\parallel}$. A perpendicular field $B_{\perp}$
induces large oscillations in $\langle S \rangle$ before the normal state
transition at 0.4 T (Fig. 5c) and a transverse field $B_{t}$
causes $\langle S \rangle$ to oscillate around zero (Fig. 5d). Figures 5e–f
shows the oscillation amplitude $\langle A \rangle$ for the three field
directions. In all directions, coherent transport was ob-
erved, with oscillation amplitude first increasing as $\langle S \rangle$
approached zero. This indicates that the oscillation am-
plitude is dictated by the energy $E_0$ in all field directions
and shows that interference is not unique to a parallel
magnetic field.

Finally, we comment on the physical mechanism that
correlates the oscillation amplitude to the energy of $E_0$.
At low fields, the Coulomb island favours an even parity
where transport of electrons occurs as two sequen-
tial tunnelling events on either end of the Coulomb is-
land [23, 26]. The two electrons acquire the condens-
ate phase when forming a Cooper pair, which suppresses
single-electron coherence. At moderate fields, a discrete
sub-gap state is brought below $E_c$ and a single-electron
transport channel is opened, allowing for coherent res-
onant tunnelling through the Coulomb island. Finally,
when the Coulomb island is in the normal state, we in-
terpret the reduction in interference signal to reflect the
short coherence length in the diffusive Al wire.

In conclusion, the observation of AB interference com-
bined with a stable discrete zero-energy states is con-
sistent with current MZM predictions and allows for the
contribution of ABS localized at the ends of the wire to be
excluded [15], which have been a possible interpretation of
the previously observed ZBP's in local tunneling ex-
periments [11]. Our results also highlight the contribution
to coherent transport by electronic modes extending over
the full length of the island (at least in the length investigat-
gated here) and demonstrates how AB interference alone
is not sufficient evidence to draw conclusions about elec-
tron teleportation via MZMs. To discriminate between
trivial and topological states, stable 1e peak spacing and
discrete zero-bias conductance peaks at successive charge
dergeneracy points must accompany AB interference.
Another approach could involve the suppression of trivial
extended modes by increasing the wire length to greatly
exceed the diffusive coherence length $\xi = \sqrt{\xi_0 l_c} \sim 1$ μm
(for $\Delta = 75$ μeV at $B_{\parallel} = 2.5$ T), where $\xi_0$ is the clean
coherence length and $l_c \sim 300$ nm is the semiconducting
mean free path [30].

These results suggest that InAs-Al 2DEGs are a
promising route towards more complex experiments re-
lated to the braiding or fusion of MZMs. We have estab-
lished Coulomb islands as coherent links necessary for
topological qubit networks, as well as demonstrated that
the transmission phase can measure the parity of the is-
land. Future devices will take advantage of improved
material quality to allow for increased wire lengths to
suppress coherent trivial quasiparticle transport, allow-
ing MZM contributions to be better separated from other
potential contributions.

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**METHODS**

The devices were fabricated on wafers grown by molecular beam epitaxy on an InP substrate. The wafer stack consists of a 1 μm graded In1−xAlxAs insulating buffer, a 4 nm In0.81Ga0.19As bottom barrier, a 5 nm InAs quantum well, and a top barrier consisting of 5 nm In0.9Al0.1As for device 1 and 10 nm In0.81Ga0.19As for device 2. A 7 nm film of epitaxial Al was then grown in-situ without breaking the vacuum of the chamber. The InAs 2DEGs were characterized with a Hall bar geometry (Al removed), which showed a peak mobility of $\mu = 17,000 \, \text{cm}^2\text{V}^{-1}\text{s}^{-1}$ for an electron density of $n = 1.7 \times 10^{12} \, \text{cm}^{-2}$ and $n = 7.5 \times 10^{11} \, \text{cm}^{-2}$ for device 1 and device 2, respectively.

Devices were fabricated using standard electron beam lithography techniques. The devices were electrically isolated by etching mesa structures by first removing the top Al film with Al etchant Transene D, followed by a deep III-V chemical wet etch $\text{H}_2\text{O} : \text{C}_4\text{H}_8\text{O}_7 : \text{H}_3\text{PO}_4 : \text{H}_2\text{O}_2$ (220:55:3:3). Next, the Al film on the mesa was selectively etched with Al etchant Transene D to produce a Coulomb island. The InAs 2DEGs in a proximitized semiconductor nanowire revealed by the Coulomb blockade of Andreev reflection. *Phys. Rev. Lett.* **70**, 1862 (1993).

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Electrical measurements were performed by standard lock-in techniques at 166 Hz by applying the sum of a variable dc bias voltage $V_{sd}$ and an ac excitation voltage of 3 to 10 $\mu$V applied to one of the top ohmic contacts as shown in Fig. 1a. The resulting current across the device was recorded by grounding a bottom ohmic via a low-impedance current-to-voltage converter, and the four terminal voltage was measured by an ac voltage amplifier with an input impedance of 500 M$\Omega$. All measurements were taken in a dilution refrigerator with a base temperature of 20 mK and an electron temperature of 40 mK estimated by the temperature dependence saturation of ZBP conductance [6].
FIG. S.1. Temperature dependence of Aharonov-Bohm oscillations in a normal 2DEG. a, False-colour electron micrograph of a normal conducting AB interferometer defined in an InAs 2DEG (green) by gates (Ti/Au). The hole forming the interferometer center is created by a wet etch. b, Magnetoconductance $\Delta G$ as a function of perpendicular magnetic field $B_\perp$ controlling the flux in the interferometer for several temperatures. Periodic oscillations are observed with a frequency of $f \sim 0.26 \text{ mT}^{-1}$, which agrees with a single magnetic flux quantum $\hbar/e$ piercing the interferometer loop. c, Temperature dependence of the AB oscillations amplitude $A$ measured from the power spectrum of the curves in b. For a diffusive interferometer, the amplitude $A = A_0 \exp(-L/l_0(T))$ where $l_0(T) \propto T^{-1/2}$ is the phase coherence length and $L = 4.5 \mu m$ is the circumference of the interferometer [31]. The exponential fit $A = A_0 \exp(-a T^{1/2})$ gives a base temperature coherence length of $l_0(20 \text{ mK}) = 4 \mu m \pm 1 \mu m$. Error bars show the standard deviation between 4 data sets at each temperature.
FIG. S.2. Transverse and perpendicular fields for device 1 in the regime of Fig. 1. 

a, Zero-bias differential conductance \( G \) as a function of gate voltage \( V_{pg} \) controlling electron occupancy and perpendicular field, \( B_\perp \). 

b, Coulomb peak spacing difference \( \langle S \rangle \) as a function of \( B_\perp \). 

c, Differential conductance \( G \) as a function of source-drain bias voltage \( V_{sd} \) and \( V_{pg} \) showing Coulomb diamonds for \( B_\perp = 0.24 \) T. 

d, Zero-bias differential conductance \( G \) as a function of \( V_{pg} \) and transverse field, \( B_t \). 

e, Coulomb peak spacing difference \( \langle S \rangle \) as a function of \( B_t \). 

f, Differential conductance \( G \) as a function of \( V_{sd} \) and \( V_{pg} \) showing Coulomb diamonds for \( B_t = 1 \) T. The field directions are represented in Fig. 1.
FIG. S.3. Conductance oscillations evolutions in parallel field for Device 2. Magnetoconductance for parallel field values $B_\parallel = 0, 1.2, 2.1,$ and $3.5$ T (left to right). a-d, Zero-bias differential conductance $G(B_\perp = 0)$ versus gate voltage $V_{pg}$ controlling electron occupation. e-h, Magnetoconductance $\Delta G$ as a function of $V_{pg}$ and perpendicular field $B_\perp$ controlling the flux in the interferometer with corresponding power spectra in i-l. A single flux quantum piercing the loop area $A_{loop} \sim 1.8 \mu m^2$ corresponds to a frequency $f_{loop} = A_{loop}/(h/e) \sim 0.44$ mT$^{-1}$. e-l, a slowly varying background has been subtracted.

FIG. S.4. Coulomb blockade for device 2. a, Zero-bias differential conductance $G$ as a function of parallel magnetic field $B_\parallel$ and gate voltage $V_{pg}$ controlling the electron occupancy with the reference arm closed. b-e, Differential conductance $G$ as a function of voltage bias $V_d$ and $V_{pg}$ for $B_\parallel = 0$ T (b), 2 T (c), 2.75 T (d), and 3.5 T (e).
FIG. S.5. Second gate configuration of device 1.  
a. Peak spacing difference $\langle S \rangle$ as a function of parallel magnetic field $B_{\parallel}$. 
b. Aharonov-Bohm oscillation amplitude $\langle \tilde{A} \rangle$ as a function of $B_{\parallel}$. The solid line is a guide to the eye. 

- **c-j**: Magnetoconductance for parallel magnetic fields $B_{\parallel} = 0$ T, 1.7 T, and 2.5 T and transverse magnetic field $B_{t} = 1$ T (left to right). 

- **c-f**: Zero-bias differential conductance $G(B_{\perp} = 0)$ versus gate voltage $V_{pg}$ used to control electron occupation. 

- **g-j**: Magnetoconductance $\Delta G$ as a function of $V_{pg}$ and perpendicular magnetic field $B_{\perp}$ controlling the flux in the interferometer.