Electric field tunable superconductor-semiconductor coupling in Majorana nanowires

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Abstract

We study the effect of external electric fields on superconductor-semiconductor coupling by measuring the electron transport in InSb semiconductor nanowires coupled to an epitaxially grown Al superconductor. We find that the gate voltage induced electric fields can greatly modify the coupling strength, which has consequences for the proximity induced superconducting gap, effective g-factor, and spin-orbit coupling, which all play a key role in understanding Majorana physics. We further show that level repulsion due to spin-orbit coupling in a finite size system can lead to seemingly stable zero bias conductance peaks, which mimic the behavior of Majorana zero modes. Our results improve the understanding of realistic Majorana nanowire systems.

1. Introduction

The hybrid superconductor-semiconductor nanowire system is the prime candidate to realize, control, and manipulate Majorana zero modes (MZMs) for topological quantum information processing [1–3]. MZMs can be engineered in these hybrid nanowire systems by combining the one-dimensional nature of the nanowire, strong spin-orbit coupling, superconductivity, and appropriate external electric (to control the chemical potential) and magnetic fields (to control the Zeeman energy) to drive the system into a topologically non-trivial phase [4, 5]. To induce superconductivity in the semiconductor nanowire, it needs to be coupled to a superconductor. The electronic coupling between the two systems turns the nanowire superconducting [6], known as the proximity effect. Following this scheme, the first signatures of MZMs were observed in these hybrid systems, characterized by a zero bias peak (ZBP) in the tunneling conductance spectrum [7–10]. Since then, significant progress has been made in Majorana experiments [11–14], enabled by more uniform coupling between the superconductor and semiconductor nanowire. This has been achieved by improved interface engineering: through careful ex situ processing [15–17], by depositing the superconductor on the nanowires in situ [18, 19], and a combination of in situ and ex situ techniques [20], finally leading to the quantization of the Majorana conductance [13].

However, the treatment of the superconductor-semiconductor coupling in the interpretation of experiments is often oversimplified. This coupling has recently been predicted to depend substantially on the confinement induced by external electric fields [21]. In this work, we experimentally show that the superconductor-semiconductor coupling, as parameterized by the induced superconducting gap, is affected by...
gate induced electric fields. Due to the change in coupling, the renormalization of material parameters is altered, as evidenced by a change in the effective g-factor of the hybrid system. Furthermore, the electric field is shown to affect the spin–orbit interaction, revealed by a change in the level repulsion between Andreev states. Our experimental findings are corroborated by numerical simulations.

2. Experimental set-up

We have performed tunneling spectroscopy experiments on four InSb–Al hybrid nanowire devices, labeled A–D, all showing consistent behavior. The nanowire growth procedure is described in [20]. A scanning electron micrograph (SEM) of device A is shown in figure 1(a). Figure 1(b) shows a schematic of this device and the measurement set-up. For clarity, the wrap-around tunnel gate, tunnel gate dielectric and contacts have been removed on one side. A normal-superconductor (NS) junction is formed between the part of the nanowire covered by a thin shell of aluminum (10 nm thick, indicated in green, S), and the Cr/Au contact (yellow, N). The transmission of the junction is controlled by applying a voltage \( V_{\text{Tunnel}} \) to the tunnel gate (red), galvanimically isolated from the nanowire by 35 nm of sputtered \( \text{Si}_3\text{N}_4 \) dielectric. The electric field is induced by a global back gate voltage \( V_{BG} \), except in the case of device B, where this role is played by the side gate voltage \( V_{SG} \). Further details on device fabrication and design are included in appendices A and B. To obtain information about the density of states (DOS) in the proximitized nanowire, we measure the differential conductance \( dI/dV_{\text{Bias}} \) as a function of applied bias voltage \( V_{\text{Bias}} \). In the following, we will label this quantity as \( dI/dV \) for brevity. A magnetic field is applied along the nanowire direction (x-axis in figures 1(b), (c)). All measurements are performed in a dilution refrigerator with a base temperature of 20 mK.

3. Theoretical model

The device geometry used in the simulation is shown in figure 1(c). We consider a nanowire oriented along the x-direction, with a hexagonal cross-section in the yz-plane. The hybrid superconductor-nanowire system is described by the Bogoliubov–de Gennes (BdG) Hamiltonian.
The first term contains contributions from the kinetic energy and the chemical potential, as well as the electrostatic potential $\phi$. The second and third terms describe the Rashba spin–orbit coupling, with the coupling strength $\alpha_y (\alpha_z)$ depending on the $y$-component ($z$-component) of the electric field. The Zeeman energy contribution, proportional to $g$, the Landé $g$-factor, is given by the fourth term. Finally, the superconducting pairing $\Delta$ is included as the fifth term. All material parameters are position dependent, taking different values in the InSb nanowire and the Al superconductor. For additional details about the simulation, see appendices C and D.

If the coupling between the superconductor and semiconductor is small (compared to the bulk gap of the superconductor $\Delta$, known as weak coupling), superconductivity can be treated as a constant pairing potential term in the nanowire Hamiltonian, with the induced superconducting gap being proportional to the coupling strength [22]. However, if the coupling becomes strong, the wave functions of the two materials hybridize, and the superconductor and semiconductor have to be considered on equal footing [23]. We achieve this by solving the Schrödinger equation in both materials simultaneously. When desired, the orbital effect of the magnetic field is added via Peierls substitution [24]. The simulations are performed using the kwant package [25].

The electrostatic potential in the nanowire cross-section is calculated from the Poisson equation, assuming an infinitely long wire. We use a fixed potential $V_{\text{Gate}}$ as a boundary condition at the dielectric-substrate interface. The superconductor enters as the second boundary condition, with a fixed potential to account for the work function difference between superconductor and semiconductor [26]. We approximate the mobile charges in the nanowire by a 3D electron gas (Thomas–Fermi approximation). It has been demonstrated that the potentials calculated using this approximation give good agreement with results obtained by self-consistent Schrödinger–Poisson simulations [27]. The calculated potential for a given $V_{\text{Gate}}$ is then inserted into the Hamiltonian (1).

By solving the Schrödinger equation for a given electrostatic environment, we can see how the gate potential alters the electronic states in the nanowire, how they are coupled to the superconductor, and how this coupling affects parameters such as the induced gap, effective $g$-factor, and spin–orbit energy.

### 4. Gate voltage dependence of the induced superconducting gap

When the transmission of the NS-junction is sufficiently low (i.e., in the tunneling regime), the differential conductance $dI/dV$ is a direct measure of the DOS in the proximitized nanowire [28]. In figure 2(a), we plot $dI/dV$ measured in device A as a function of applied bias voltage $V_{\text{Bias}}$ and tunnel gate voltage $V_{\text{Tunnel}}$ for $V_{BG} = −0.6$ V. In the low transmission regime, we resolve the superconducting gap $\Delta$ around 250 $\mu$eV, indicated by the position of the coherence peaks. The ratio of sub-gap to above-gap conductance (proportional to the normal state transmission of the junction, $T$) follows the behavior expected from BTK theory [29, 30], indicating the sub-gap conductance is dominated by Andreev reflection processes (proportional to $T^2$). This is generally referred to as a hard gap. However, for more positive back gate voltages, the sub-gap conductance is larger and shows more resonances, as is illustrated in figure 2(b) for $V_{BG} = −0.3$ V. Figure 2(c) shows line traces taken at a similar transmission (above-gap conductance) for both cases. The sub-gap conductance for $V_{BG} = −0.3$ V (black line) exceeds that of the hard gap case (red line) by an order of magnitude. This is indicative of a surplus of quasi-particle states inside the gap, referred to as a soft gap.

The gate voltage induced transition from soft to hard gap is generically observed in multiple devices. To understand this phenomenology, we calculate the electron density in the nanowire cross-section for different values of $V_{\text{Gate}}$. Because the charge neutrality point in our devices is unknown, there is a difference between the gate voltages used in the experiment and the values of $V_{\text{Gate}}$ used in the simulation. By comparing the transition point between hard and soft gaps in the experiment and the simulation, we estimate that the experimental gate voltage range $−0.6$ V $< V_{BG} < −0.4$ V roughly corresponds to the simulated gate voltage range $−0.4$ V $< V_{\text{Gate}} < −0.2$ V.

For more negative $V_{\text{Gate}}$, the electric field from the gate pushes the electrons towards interface with the superconductor (inset of figure 2(a)). We solve the Schrödinger equation for the calculated electrostatic potential and find that this stronger confinement near the interface leads to a stronger coupling. This results in a hard gap, as illustrated by the calculated energy spectrum (figure 2(d), red line). However, for more positive voltages, the electrons are attracted to the back gate, creating a high density pocket far away from the superconductor (inset of figure 2(b)). These states are weakly coupled to the superconductor, as demonstrated

\[
H = \left[ \frac{\hbar^2 k^2}{2m^*} - \mu - e\phi \right] \tau_z + \alpha_y (k_x \sigma_y - k_y \sigma_y) \tau_x \\
+ \alpha_z (k_x \sigma_y - k_y \sigma_y) \tau_z + \frac{1}{2} g\mu_B \sigma_z + \Delta \tau_z.
\]
by a soft gap structure (figure 2(d), black line, see also appendix E). We can therefore conclude that the electron tunneling between the semiconductor and the superconductor is strongly affected by the gate potential.

The change in superconductor-semiconductor coupling does not just affect the hardness, but also the size of the gap. For each back gate voltage, we fit the BCS-Dynes expression [31] for the DOS in order to extract the position of the coherence peaks, giving the gap size $\Delta$. The results are shown in figure 2(e). Further details on the fitting procedure are given in appendix F. As $V_{BG}$ becomes more positive, the superconductor-semiconductor coupling becomes weaker, reducing the size of the gap. From $V_{BG} > -0.4$ V onward it becomes difficult to accurately determine the gap, as it tends to become too soft and the coherence peaks are not always clearly distinguishable. The top right inset shows the shift of the coherence peak (indicated by the arrow) for $V_{BG} = -0.6$ V (solid red line) and $V_{BG} = -0.4$ V (dashed black line). Bottom left inset: induced gap from the calculated DOS as a function of $V_{Gate}$, consistent with the experimental observation.

5. Effective $g$-factor

As the electric field induced by the back gate clearly has an important effect on the hybridization between the nanowire and the superconductor, we now look at the effect this has on the Zeeman term in the Hamiltonian. This term affects the energy dispersion of spinful states in a magnetic field. We study the dispersion of the states in the nanowire by measuring $dI/dV$ in device A as a function of applied bias voltage and magnetic field, as
shown in figures 3(a) and (b). We define the effective g-factor as $g_{\text{eff}} = \frac{2}{\rho_0} \left| \frac{\Delta E}{\Delta B} \right|$, with $\Delta E$ the absolute value of the average slope of the observed peak in the differential conductance as it disperses in magnetic field. This effective g-factor is different from the pure spin g-factor $g_{\text{spin}}$ as the dispersion used to estimate $g_{\text{eff}}$ is generally not purely linear in magnetic field, and has additional contributions from the spin–orbit coupling, magnetic field induced changes in chemical potential, and orbital effects [21, 26, 32]. The effective g-factor is the parameter which determines the critical magnetic field required to drive the system through the topological phase transition [33]. We obtain the slope $\Delta E / \Delta B$ from a linear fit (shown as black dashed lines in figures 3(a), (b), see appendix G for details) of the observed peak position. Figure 3(c) shows the extracted $g_{\text{eff}}$ for device A, with more positive back gate voltages leading to larger $g_{\text{eff}}$ (visible as a steeper slope). A similar result has recently been reported in hybrid InAs–Al nanowires [34].

We use our numerical model to calculate the DOS in the nanowire as a function of applied magnetic field, shown in figures 3(d) and (e). From the calculated spectrum, we apply the same procedure used to fit the experimental data to extract $g_{\text{eff}}$ (white dashed lines). The results for different values of $V_{\text{G}}$ are given in figure 3(f) as black circles. The applied back gate voltage changes the hybridization of the states in the InSb ($g_{\text{spin}} = 40$ [35]) and the Al ($g_{\text{spin}} = 2$). As a more positive gate voltage increases the weight of the wave function in the InSb, we expect the renormalized g-factor to increase as the gate voltage is increased, consistent with the results of figures 3(c) and (f).

To see how well $g_{\text{eff}}$ describes the Zeeman term in the Hamiltonian, we turn our attention to the energy spectrum at $k = 0$. At this point, the effect of spin–orbit coupling vanishes, if orbital effects are excluded, we can then define the absolute value of the pure spin g-factor as $g_{\text{spin}} = \frac{2}{\rho_0} \left| \frac{\Delta f(k = 0)}{\Delta B} \right|$. The resulting values for $g_{\text{spin}}$ are shown as red squares in figure 3(f). By comparing the results for $g_{\text{eff}}$ and $g_{\text{spin}}$, we can conclude that when the lowest energy state has a momentum near $k = 0$ (as is the case for $V_{\text{G}} < -0.2$ V), the effect of spin–orbit
coupling is negligible, and \( g_{eff} \) is a good proxy for the pure spin \( g \)-factor. However, when this is no longer the case, deviations can be observed, as is the case for \( V_{Gate} \geq -0.2 \) V. As we expect the experimental gate voltage range of figure 3(c) to be comparable to values of \( V_{Gate} < -0.2 \) V, we conclude that the experimentally obtained \( g_{eff} \) is a reasonable approximation of \( g_{spin} \) in this parameter regime. However, we stress once more that in general, one needs to be careful when interpreting the \( g_{eff} \) extracted from experimental data as the \( g \)-factor entering the Hamiltonian in the Zeeman term.

The increasing trend of \( g_{eff} \) does not change when the orbital effect of magnetic field is considered (see appendix G, figure G4). However, there is a significant increase in the predicted values, in agreement with previous findings for InAs nanowires [32]. The values in figure G4 are larger than the ones generally observed in our experiment (see figure 3(c)), suggesting that the orbital effect is not a dominant mechanism in determining the effective \( g \)-factor in these devices. We note that the data from device \( \Lambda \) used to make these plots was taken solely in the hard gap regime, where one expects a strong confinement near the superconductor. This suppresses the orbital contribution of the magnetic field. Another possible explanation for the discrepancy between the results of the simulation and the experimental data is an overestimation of the density in the nanowire, as higher sub-bands have a stronger contribution from the orbital effect. Minimizing the orbital effect is desirable for Majorana physics, as the orbital contributions of the magnetic field are detrimental to the topological gap [24].

6. Level repulsion due to spin–orbit coupling

The term in the Hamiltonian that remains to be explored describes the Rashba spin–orbit coupling. The strength of the spin–orbit coupling is determined by the parameter \( \alpha \), which depends on the material (and thus, on the superconductor–semiconductor coupling), and the electric field [36–38]. Therefore, we expect that this term will be affected by the gate potential as well. In finite systems, the spin–orbit interaction can couple states with different orbitals and spins [39]. These states are thus no longer orthogonal to each other, and the spin–orbit mediated overlap between them causes energy splitting, leading to level repulsion [40–42]. This level repulsion, which is generic in class D systems in the presence of superconductivity, magnetic field and spin–orbit coupling [43, 44], can be extracted from the low energy nanowire spectrum as measured by tunneling spectroscopy [45].

In figures 4(a)–(c), we show the evolution of the level repulsion between the two lowest energy sub-gap states (labeled \( L_1 \) and \( L_2 \), as indicated by the white dashed lines in panel (c)) in device B. For these measurements, the global back gate is grounded, with the electric field being induced by applying a voltage to the side gate (side gate shown in appendix B).

We parameterize the level repulsion by two quantities: the coupling strength \( \delta_{SO} \), and the splitting \( A \), defined as the maximum deviation of \( L_1 \) from zero energy after the first zero crossing. This splitting has previously been linked to the overlap between two MZM in a finite system [46]. In figure 4(e), we zoom in on the anti-crossing feature in panel figure 4(b), showing the minimum energy difference between \( L_1 \) and \( L_2 \) (given by \( 2\delta_{SO} \)) and the splitting \( A \). We extract these parameters by a fit of the anti-crossing (solid green lines, with the uncoupled states shown by the dashed black lines, details of the fitting procedure are in appendix H).

Because we expect finite size effects to be relevant, we cannot use our previous theoretical model, as it is based on an infinitely long nanowire. Therefore, we modify the model to take into account the finite size of the nanowire system, and calculate the low energy spectrum for different values of the Rashba spin–orbit strength (see appendix I). In figure 4(d), we plot the two lowest energy states in the nanowire as a function of the Zeeman energy (\( E_z = g_e \mu_B B \)), in units of the superconducting gap \( \Delta \). If \( \alpha = 0 \) (no spin–orbit coupling, dashed black lines), there is no coupling between the states, and no level repulsion occurs. However, if spin–orbit coupling is included (e.g., \( \alpha = 0.1 \) eV A, solid red lines), the levels repel each other, with the magnitude of the anti–crossing given by \( 2\delta \). The level repulsion strength scales with \( \alpha \) (inset of figure 4(d)), providing a way to estimate \( \alpha \) based on the low energy spectrum using \( 2\delta \sim \alpha \pi / l \), where \( l \) is the length of the nanowire.

In figure 4(f), we plot \( \delta_{SO} \) (black circles) and \( A \) (red squares) as a function of the applied side gate voltage. The two parameters follow opposite trends, with \( A \) being maximal when \( \delta_{SO} \) is minimal. When \( \delta_{SO} \) is larger, the levels repel each other more, leading to \( L_1 \) being pushed closer to zero energy, reducing the splitting \( A \). When \( V_{SG} < 2.0 \) V, both parameters become smaller with decreasing \( V_{SG} \). At this point, other states at higher energies become relevant for the lowest energy dispersion (a situation demonstrated in figure 4(a)), and our method to extract these parameters breaks down. We expect this method to be reliable when the energetically lowest two states can be clearly separated from the rest.

Because \( \delta_{SO} \) depends not only on \( \alpha \), but also on the details of the confinement potential, as well as the coupling to the superconductor, a precise estimate goes beyond the current approximations in our model. That being said, based on the observed magnitude of \( \delta_{SO} \) and our simulations of the finite nanowire system, we can
estimate the Rashba parameter $\alpha$ to be around 0.1 eV Å in this gate voltage range. This value is comparable to the values reported in InSb nanowire based quantum dots [47], and smaller than the values measured in weak anti-localization experiments [37]. A large value of $\alpha$ is beneficial for Majorana physics, as it determines the maximum size of the topological gap [48].

7. ZBP in extended magnetic field range

In the previous sections, we have described the effect of the gate induced electric field on the various terms in the Hamiltonian \(1\). As this Hamiltonian is known to describe Majorana physics, we now turn our attention to possible signatures of MZMs in this system. In particular, when $2\delta_{SO}$ becomes comparable to the energy of $L_2$, we find that $L_1$ can become pinned close to zero bias over an extended range in magnetic field, as demonstrated in figure 5(b) (data from device A). Figure 5(d) shows that the state stays pinned to zero energy over a range of over 0.2 T, corresponding to a Zeeman energy of over 300 $\mu$eV, which is larger than the induced gap. The stability of the ZBP in terms of the ratio of Zeeman energy to induced gap is comparable to the most stable ZBPs reported in literature [11, 12]. When we fix the magnetic field to $B = 0.26$ T and change the back gate voltage (figure 5(e)), it appears that there is a stable ZBP over a few mV as well.

We might be tempted to conclude that this stability implies this is a MZM. However, if we change either the gate voltage (figures 5(a), (c)) or the magnetic field (figure 5(f)) a little bit, we observe that this stability applies only to very particular combinations of gate voltage and magnetic field. One should keep in mind that in a finite system, MZMs are not expected to be stable with respect to local perturbations if the system size is comparable to the Majorana coherence length, which is likely the case in our devices. This further complicates the determination of the origin of the observed peaks. As we find no extended region of stability, we conclude that it is unlikely that this state pinned to zero energy is caused by a topological phase transition. Rather, this seems to be due to a fine-tuned coincidence in which the repulsion between two states combined with particle–hole

![Figure 4](image_url)
symmetry leads to one of the states being pinned to $E = 0$. We reiterate that simply having a stable zero energy state over an extended range in magnetic field is not sufficient to make claims about robust Majorana modes [49–51]. Further experimental checks, such as stability of the ZBP in an extended region of the parameter space spanned by the relevant gate voltages [11], as well as magnetic field, are required in order to assign a possible Majorana origin.

8. Conclusion and outlook

We have used InSb nanowires with epitaxial Al superconductor to investigate the effect of the gate voltage induced electric field on the superconductor–semiconductor coupling. This coupling is determined by the distribution of the wave function over the superconductor and semiconductor, and controls essential parameters of the Majorana Hamiltonian: the proximity induced superconducting gap, the effective $g$-factor, and spin–orbit coupling. Our observations show that the induced superconductivity, as parameterized by the hardness and size of the induced gap, is stronger when the electrons are confined to a region close to the superconductor. The stronger coupling leads to a lower effective $g$-factor. We also determine that the gate voltage dependence of the effective $g$-factor is dominated by the change in coupling to the superconductor, rather than by orbital effects of the magnetic field. Finally, we study the effect of level repulsion due to spin–orbit coupling. Appropriate tuning of the repulsion leads to level pinning to zero energy over extended parameter ranges, mimicking the behavior expected from MZMs. Our result deepens the understanding of a more realistic Majorana nanowire system. More importantly, it is relevant for the design and optimization of future advanced nanowire systems for topological quantum information applications.
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Author contributions

MWAdM, JDSB, DX, and HZ fabricated the devices, performed the measurements, and analyzed the data. GWW, AB, AEA, and RML performed the numerical simulations. NvL and GW contributed to the device fabrication. RLMOhV, SG, and DC grew the InSb nanowires under the supervision of EPAMB. JAL, MP, and JSL deposited the aluminum shell on the nanowires under the supervision of CJP. LPK and HZ supervised the project. MWAdM and HZ wrote the manuscript with comments from all authors. MWAdM, JDSB, and DX contributed equally to this work.

Appendix A. Fabrication procedure

1. Nanowire deposition: a SEM-based nanomanipulator is used to deterministically place the InSb–Al nanowires unto a degenerately p-doped Si substrate covered by 20 nm of LPCVD Si₃N₄ (devices A, C, and D) or 285 nm of thermal SiO₂ (device B).

2. Mask preparation and lithography: for every fabrication step, we use standard electron beam lithography techniques to create the mask. The mask consists of a layer of PMMA 950KA6 spun at 4000 rpm. After writing, the mask is developed in a solution of MIBK:IPA (1:3 ratio) for 60 s, followed by a IPA rinse for 60 s. After each deposition step, liftoff is done using acetone.

3. Contact preparation and deposition: before depositing the contact material, the Al and AlOₓ are locally removed by Ar plasma etch. The contacts are deposited by electron beam evaporation of Cr/Au (10/100–200 nm). For device B, the side gates are also evaporated in this step.

4. Dielectric deposition: as a top gate dielectric we sputter 35 nm of SiNₓ (devices A, C, and D).

5. Top gate deposition: the top gates are deposited by electron beam evaporation of Ti/Au (10/200 nm) (devices A, C, and D).
Appendix B. Device information and schematics

<table>
<thead>
<tr>
<th>Device</th>
<th>Top view SEM</th>
<th>Cross-section</th>
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<tbody>
<tr>
<td>A</td>
<td><img src="device_A.png" alt="" /></td>
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<tr>
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<tr>
<td>C</td>
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<td>D</td>
<td><img src="device_D.png" alt="" /></td>
<td><img src="cross-section_D.png" alt="" /></td>
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Figure B1. SEM images and schematic cross-sectional views of the devices used as part of this research. Data from devices A and B is presented in the main text. Data from devices C and D is presented in the supplement for completeness. Note that the data for device B is obtained by changing the voltage on the side gate ($V_{SG}$), shown in blue in the SEM image. Scale bar is 500 nm.

Appendix C. Simulation of electrostatics and nanowire spectrum

For the electrostatics simulations we use the geometry of device A (as shown in figure 1(c) of the main text). We describe the device as an infinite wire oriented along the $x$-direction, with a hexagonal cross-section in the $yz$-plane. The electrostatics are described by the Poisson equation

$$\nabla \cdot (\varepsilon_0(r) \nabla \phi(r)) = \frac{\rho_{\text{tot}}(r)}{\varepsilon_0},$$

where $\rho_{\text{tot}}(r)$ is a functional of the potential $\phi(r)$. We include four contributions to $\rho_{\text{tot}}$,

$$\rho_{\text{tot}} = \rho_e + \rho_{hh} + \rho_{lh} + \rho_{\text{fixed}},$$

where $\rho_e$, $\rho_{hh}$, and $\rho_{lh}$ are the mobile charges of the conduction band, heavy hole (hh) band and light hole (lh) band of the InSb nanowire and $\rho_{\text{fixed}}$ are the fixed charges in the system. For the mobile electron charges we assume a 3D electron gas density (Thomas–Fermi approximation).
over the whole semiconductor region.

\[ k \cdot B \times y yz \]

required for a hard induced gap in the InAs difference between the two materials. A negative band offset of the semiconductor to the superconductor is we assume the conduction band of InSb is pinned

\[ \text{where the average electric} \]

The Rashba couplings

\[ \text{Appendix D. Electric} \]

\[ \text{with} \]

\[ \text{with} \]

\[ \text{with} \]

\[ \text{with} \]

\[ \text{with} \]

\[ \text{with} \]

\[ \text{The boundary condition at the substrate-dielectric interface is set to the applied gate voltage, } V_{\text{Gate}}. \]

\[ \text{On the remaining three boundaries of the system we use Neumann} \]

\[ \text{For the dielectric constant of InSb, the sputtered SiN, and the LPCVD SiN}_x \text{ we take take 15.15, 7.5, and 8, respectively.} \]

\[ \text{After the electrostatic potential has been calculated for a given } V_{\text{Gate}}, \text{we plug it into the Schrödinger equation and solve it for the cross-section of the device. We use a Rashba Hamiltonian with a BdG superconducting term} \]

\[ \text{where } \alpha \text{ is the Rashba parameter, } \alpha_{\phi} \text{ the Rashba parameter of the orbital effect, and } \beta \text{ the magnetic flux quantum.} \]

\[ \text{The Rashba couplings } \alpha_x \text{ and } \alpha_y \text{ which are nonzero only in the semiconductor region, result from the symmetry breaking by the electrostatic potential and are obtained from} \]

\[ \text{where the average electric field in direction } i \text{ is obtained by averaging } E_i \text{ over the whole semiconductor region. Parameters for bulk InSb are used } [52]: \text{the Kane matrix element } P = 0.9641 \text{ eV nm, the bandgap } E_0 = 0.237 \text{ eV, and the spin–orbit gap } \Delta_0 = 0.810 \text{ eV. The resulting Rashba parameters } \alpha_i \text{ are plotted in figure D1(a).} \]
We define the spin–orbit energy $E_{SO} = \frac{m^* (\alpha_x^2 + \alpha_y^2)}{2}\hbar^2$. The spin–orbit energy is plotted as a function of $V_{Gate}$ in figure D1(b). The average electric field in the nanowire increases as the applied gate voltage becomes more negative, leading to an enhancement of the spin–orbit coupling. At $V_{Gate} = 0.08$ V, the average electric field in the nanowire becomes equal to 0 due to symmetry, eliminating the influence of spin–orbit coupling on the nanowire spectrum.

Appendix E. Simulated band structure

The band structure of the superconductor-semiconductor nanowire system for different values of $V_{Gate}$ is shown in figure E1. To quantify the coupling of a given state to the superconductor, we calculate the weight of the state in the semiconducting region SM as $W_{SM} = \int_{SM} |\psi(k_x)|^2 \, dy \, dz$. 

Figure D1. (a) Rashba coefficients $\alpha_y$ and $\alpha_z$ as a function of $V_{Gate}$. At $V_{Gate} = 0.08$ V, the average electric field in the wire goes to zero due to symmetry, leading to vanishing spin–orbit coupling. (b) Calculated $E_{SO}$ as a function of $V_{Gate}$. 
Appendix F. Gap fitting and additional data

To extract the gap, we measure the differential conductance $dI/dV$ as a function of $V_{\text{Bias}}$ and tunnel gate voltage $V_{\text{Tunnel}}$ for different back gate voltages $V_{\text{BG}}$. In the tunneling limit, $dI/dV$ is approximately proportional to the DOS. To ensure we are in this limit, we take only the traces where the conductance at high bias ($\sim 500 \, \mu V$) is between $0.03 \cdot e^{-0.08 \cdot \Delta^2}$. We use the BCS-Dynes expression for a dissipation broadened superconducting DOS\[31\] to arrive at the following expression for the conductance:

$$\frac{dI}{dV} = G_0 \text{Re} \left\{ \frac{V_{\text{Bias}} - i\Gamma}{\sqrt{(V_{\text{Bias}} - i\Gamma)^2 - \Delta^2}} \right\}. \quad (F.1)$$

This equation is fitted to the data (separately for positive and negative bias), as shown in figure F1 for $V_{\text{BG}} = -0.6 \, \text{V}$. We take the average of the extracted gap values for different values of $V_{\text{Tunnel}}$, with the error bar given by the standard deviation (results plotted in figure 2(e)).

Device B shows similar behavior to device A: as the side gate voltage is increased, the observed gap becomes smaller (as illustrated in figure F2).
In figure F3, we show differential conductance traces as a function of $V_{\text{Bias}}$ in device D for different values of the back gate voltage. The voltage on the tunnel gate is chosen such that the transmission through the junction (parameterized by $G_N$) is constant.

Although the sub-gap conductance is similar for all three gate voltages, there is a strong broadening of the coherence peak as the gate voltage becomes more positive. This broadening is associated with dissipation due to an increase in the number of quasiparticles, caused by pair breaking in the superconductor. We plot the extracted gap $\Delta$ and dissipation broadening $\Gamma$ in figure F4.

As in the other devices, the gap decreases for more positive gate voltages, although in this case the effect is minor. The size of the gap is quite stable over an extended range in gate voltage. We speculate that this is related to the diameter of the wire, which is smaller than in the other devices. The reduced thickness means the superconductor can screen the gate voltage more effectively throughout the wire diameter, reducing the effect of the gate on the superconductor-semiconductor coupling.
Appendix G. g-factor fitting and additional data

For each back gate voltage, we measure the $dI/dV$ as a function of $V_{\text{Bias}}$ and the magnetic field $B$. We then identify the lowest energy peak in the spectrum. The position of this peak at a given field is obtained by a peak finding algorithm, the results of which are shown as the green circles in figure G1. The slope $D_B$ is determined by a linear fit (dashed black line in figure G1). From the slope, we get $g_{\text{eff}}$ by using the relation $D_B = \frac{\Delta E}{\Delta B}$ for a spin-$\frac{1}{2}$ particle, with $\mu_B$ the Bohr magneton. This procedure is performed separately for positive and negative bias. The reported $g_{\text{eff}}$ is then calculated as a weighted average of the absolute value of the positive and negative bias results (weights determined by the variance of the fit parameters).

The effective $g$-factor for device A is reported in the main text (figures 3(c) and 4(d), respectively). In figure G2, we plot the extracted $g$-factors of both $L_1$ and $L_2$ in device B. For completeness, data from device C is shown in figure G3.

The effective $g$-factor of $L_1$ (black circles) changes appreciably when the side gate voltage is changed, with the effect comparable to the one observed in device A. In contrast, $g_{\text{eff}}$ of $L_2$ (red squares) is almost unaffected by the gate and has a lower value. This may be due to $L_2$ being closely confined near the superconductor, leading to a decreased $g$-factor due to stronger hybridization, and a weaker gate response due to enhanced screening.

To determine the importance of orbital effects, we calculate the nanowire spectrum as a function of magnetic field including this effect (figures G4(a), (b)). The orbital effect leads to an increase of the extracted values of $g_{\text{eff}}$ and $g_{\text{spin}}$ (figure G4(c)). Note that the definition of $g_{\text{spin}}$ used in the main text is no longer valid when the orbital effect is included. Nevertheless, for consistency we apply the same procedure. As we do not observe these high $g$-factors in any of our devices, we conclude that the orbital effect does not give a significant contribution to the observed changes of $g_{\text{eff}}$ with the gate voltage.

![Figure G3. Fit (red line) of equation (F.1) to conductance data from device D (black circles, squares and triangles) for different $V_{BG}$ at similar junction transparencies. The dissipation broadening sharply decreases for more negative gate voltages.](image)

![Figure G4. Extracted gap $\Delta$ as a function of $V_{BG}$ for device D. Inset: dissipation broadening $\Gamma$ as a function of $V_{BG}$. A decrease in the gap is accompanied by an increase in broadening, signaling the emergence of a soft gap.](image)
Figure G1. Differential conductance as a function of $V_{\text{bias}}$ and magnetic field. We apply a linear fit (dashed black lines) to the extracted peak positions (green circles) to obtain the average slope $\frac{\Delta g_{\text{eff}}}{\Delta V_{\text{bias}}}$. 

Figure G2. Extracted values of $g_{\text{eff}}$ as a function of $V_{SG}$ for $L_1$ (black circles) and $L_2$ (red squares) in device B. 

Figure G3. Extracted values of $g_{\text{eff}}$ for device C.
Appendix H. Anticrossing fitting

Near the anticrossing, we approximate the energy of the lowest subgap state $L_1$ as $E_i + \frac{1}{2}g_i \mu_B B + aB^2$. The linear term represents the Zeeman contribution to the energy, while the quadratic term is a correction to account for the curvature at high fields. This is possibly due to the presence of additional levels interacting with $L_1$ in this field range. As the dispersion of $L_2$ is mostly linear in the field range of interest, we approximate it as $E_2 = \frac{1}{2}g_2 \mu_B B$. Adding the coupling parameter $\delta_{SO}$, we find the energy levels of the coupled system from the eigenvalues of the matrix

$$
\begin{bmatrix}
E_1 + \frac{1}{2}g_1 \mu_B B + aB^2 & \delta_{SO} \\
\delta_{SO} & E_2 - \frac{1}{2}g_2 \mu_B B
\end{bmatrix}.
$$

By fitting the expression for the eigenvalues to the data (see figure H1), we extract the parameters $E_{1,2}, g_{1,2}, a$, and $\delta_{SO}$. To prevent overfitting, we use estimates for the uncoupled asymptotes to constrain the fit parameters. From the obtained parameters we also calculate the splitting $A$, defined as the maximum deviation from zero energy of the lowest energy state $L_1$, after the first zero energy crossing has occurred.
Appendix I. Simulation of finite size nanowire system

To simulate the finite nanowire system, we solve the Hamiltonian (1) in a simplified setup. We consider a rectangular cross-section in the yz-plane similar to the one used in reference [21], where the top facet of the rectangle is covered by the superconductor, and a uniform gate voltage $V_{\text{Gate}}$ is applied to the bottom facet, as illustrated in figure I1. First, we assume an infinitely long nanowire oriented in the x-direction, and calculate the electrostatic potential in the Thomas–Fermi approximation, similar to the procedure described in appendix C. The Fermi level in the nanowire is tuned such that it supports the same number of transverse modes at $V_{\text{Gate}} = 0$ as the hexagonal nanowire studied previously. We use the same material parameters as in the previous simulation, which can be found in table C1.

We then plug the resulting electrostatic potential into (1) and solve the Schrödinger equation to find the low energy spectrum of the finite nanowire. We take a length of 750 nm, similar to the studied devices. We calculate only the modes in the semiconductor, assuming a superconducting gap of $\Delta = 250 \mu\text{eV}$. We find that the origin of the level repulsion between states is indeed spin–orbit coupling, which couples different longitudinal (along the x-direction) states within the same transverse (y- and z-directions) subband.

The result is illustrated in figure I2, where we plot the low energy spectrum as a function of Zeeman energy $E_Z$ for a fixed value of $V_{\text{Gate}}$ and different values of $\alpha$. An increase in the spin–orbit coupling strength leads to an increase in the level repulsion.

However, even if $\alpha$ is fixed, the magnitude of the level repulsion can be changed by changing the confinement potential, as demonstrated in figure I3.

When the gate voltage is changed, it alters the confinement potential. This affects the energy of the levels coupled by the spin–orbit coupling, and as such directly influences the magnitude of $\delta$, even though the spin–orbit coupling strength itself is not changed appreciatively. In figure I4 we plot the calculated energy gap due to level repulsion, $2\delta$, and the maximum splitting from zero energy of the lowest energy state after the first zero crossing, $A$, as a function of $V_{\text{Gate}}$. The two parameters follow opposite trends, consistent with the experimental observation in figure 4(f).

However, the trend with gate voltage is opposite: $\delta$ increases with more positive gate voltage, whereas in the experiment it decreases. We note that the geometry used in this simulation is a simplified version of the one used in the experiment. The dependence of the confinement energy on gate voltage is strongly dependent on the geometry, which differs between the simulation and the experiment. It is therefore expected that the trend of $\delta$ with gate voltage is not universal, and requires the details of the systems to be very similar before comparisons can be made.
Figure 11. Schematic cross-section of the geometry used to simulate the finite nanowire system. A potential $V_{\text{Gate}}$ is applied to the bottom facet, while the potential at the top facet is fixed by the work function difference between the two materials. The magnetic field is applied in the $x$-direction, along length of the nanowire.

Figure 12. Calculated low energy spectrum of the finite size nanowire as a function of Zeeman energy for different values of $\alpha$. Values calculated for $V_{\text{Gate}} = -0.536$ V, which is also used in main text figure 4(d). All energy scales are in units of the superconducting gap $\Delta$.

Figure 13. Calculated low energy spectrum of the finite size nanowire as a function of Zeeman energy for different values of $V_{\text{Gate}}$. Values calculated $\alpha = 0.1$ eV Å, energy scales in units of $\Delta$. 
Appendix J. Additional ZBP data

Figure J1(a) shows the differential conductance measured in device A as a function of $V_{\text{Bias}}$ and $V_{\text{Tunnel}}$ for $B = 0.35 \, \text{T}$ and $V_{\text{BG}} = -0.37 \, \text{V}$. The low energy spectrum in this parameter regime does not depend on the transmission of the NS-junction. In figure J1(b), we show line traces for different values of $V_{\text{Tunnel}}$. Even though the transmission of the junction is changed by a factor of two, the peak position of the low energy states are not affected. Data from main text figure 5 was obtained for $V_{\text{Tunnel}} = -87 \, \text{mV}$.

Figure J2 shows additional data on the evolution of the level repulsion between $L_1$ and $L_2$ in device A (supplementing the data presented in main text figures 5(a)–(c)) as the back gate voltage is increased. As discussed in the main text, we do not find an extended region in parameter space with a stable zero bias conductance peak.

In figure J3 we show the low energy spectrum of device A as a function of $V_{\text{Bias}}$ and $V_{\text{BG}}$ for different magnetic fields (supplementing the data presented in main text figures 5(e), (f)). For specific combinations of magnetic field and gate voltage, we can find a zero energy state. However, as we do not find an extended region in parameter space, it is unlikely that a topological phase transition is responsible for this observation.
Figure J1. (a) $\frac{dI}{dV}$ measured in device A as a function of $V_{\text{Bias}}$ and $V_{\text{Tunnel}}$ for $B = 0.35$ T and $V_{BG} = -0.37$ V. (b) Line traces at the values of $V_{\text{Tunnel}}$ indicated by the colored lines in panel (a).
Figure J2. Differential conductance as a function of $V_{\text{bias}}$ and magnetic field. Although the lowest energy state stays near zero over an extended magnetic field range for some gate voltages, this behavior is not robust.
Figure J3. Differential conductance as a function of $V_{\text{Bias}}$ and $V_{\text{BG}}$. We find some stable ZBPs for certain ranges in back gate voltage at specific fields, but this is only true for fine tuned parameters.

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