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Published in:
Physical Review B

DOI:
10.1103/PhysRevB.103.104508

Publication date:
2021

Document version
Publisher's PDF, also known as Version of record

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Citation for published version (APA):
Topological superconductivity in semiconductor–superconductor–magnetic-insulator heterostructures

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(Received 20 November 2020; revised 20 January 2021; accepted 24 February 2021; published 12 March 2021)

Hybrid superconductor-semiconductor heterostructures are promising platforms for realizing topological superconductors and exploring Majorana bound state physics. Motivated by recent experimental progress, we theoretically study how magnetic insulators offer an alternative to the use of external magnetic fields for reaching the topological regime. We consider different setups, where (1) the magnetic insulator induces an exchange field in the superconductor, which leads to a splitting in the semiconductor by proximity effect, and (2) the magnetic insulator acts as a spin-filter tunnel barrier between the superconductor and the semiconductor. We show that the spin splitting in the superconductor alone cannot induce a topological transition in the semiconductor. To overcome this limitation, we propose to use a spin-filter barrier that enhances the magnetic exchange and provides a mechanism for a topological phase transition. Moreover, the spin-dependent tunneling introduces a strong dependence on the band alignment, which can be crucial in quantum-confined systems. This mechanism opens up a route towards networks of topological wires with fewer constraints on device geometry compared to previous devices that require external magnetic fields.

DOI: 10.1103/PhysRevB.103.104508

I. INTRODUCTION

Topological superconductivity has been predicted to appear in one-dimensional spin-orbit coupled semiconductors proximitized by an $s$-wave superconductor. In these systems, an external magnetic field causes the gap to close. In the presence of a strong spin-orbit interaction, the gap reopens, leading to a topological phase [1,2]. The quantum phase transition brings the system into the topological regime when

\[ V_Z^2 > \mu^2 + \Delta^2, \]

where $V_Z$ is the Zeeman energy, $\mu$ the semiconductor chemical potential, and $\Delta$ the superconducting pairing potential. In the topological phase, the system behaves as a spinless $p$-wave superconductor and Majorana bound states appear at the edges of the system. These states have been proposed to be used for topological quantum computation [3,4]. Following this idea, different hybrid semiconductor–superconductor (Sm–Sc) platforms have been proposed to exhibit topological superconductivity, such as proximitized nanowires [5], selective-area-grown wires [6], and two-dimensional electron gas (2DEG) systems [7–10].

One limitation of these platforms is the requirement of external magnetic fields to induce the topological phase transition. External magnetic fields have several drawbacks as they are detrimental to superconductivity and the topological phase is very sensitive to the relative orientation between the magnetic and the spin-orbit fields. This problem becomes even more evident when considering more complicated geometries featuring several nanowires that cannot be aligned in the same direction, like the ones proposed to demonstrate Majorana non-Abelian statistics in real space [11–13].

Magnetic materials are alternatives to the use of external magnetic fields. This idea was discussed by some early works in the field, where magnetic insulators (MIs) were used to induce a spin splitting by means of stray fields [14,15]. In addition, clean interfaces with a MI also lead to exchange fields in the proximitized materials, which provide a more effective way to control the local spin splitting.

Developments in the fabrication technology have enabled the integration of thin layers of the MI EuS in the hybrid Sm–Sc InAs-Al nanowires with excellent interface quality [16–18]. This material has also been tested in combination with Au showing signatures consistent with the presence of Majorana bound states [19].

In experiments with ferromagnetic hybrid nanowires, spectroscopic measurements have shown the onset of a zero-bias conductance peak which has been interpreted as a signature of localized Majorana bound states at their ends [20]. This signature has been detected only when the Al and the EuS layers overlap [Fig. 1(a)]. Samples with nonoverlapping facets, like Fig. 1(b), have shown no signatures of topological phases. This behavior is in contrast with expectations that the main effect of MIs is to induce an exchange field in the Sm [21].

In these devices, however, the induced exchange field in the Sm is weak and short range [18,20]. For this reason, the
In this paper, we provide a theory for the combined superconductivity and magnetic proximity effect in SC–MI–Sm and MI–Sc–Sm heterostructures, as illustrated in Fig. 2. We show that the combined magnetic and superconducting proximity effects of a spin-split SC cannot induce a topological transition in the Sm by themselves. To overcome this limitation, we propose a new heterostructure layout where a thin film of MI on the Sm, an alternative explanation can be that the spin-split SC induces superconductivity and exchange field at the same time in the Sm. In this case, the relevant exchange coupling would be directly between the MI and the Sc [20].

In this work we disregard the induced exchange field in the Sm as experimental evidence suggests that it is small in the tested devices [18,20]. We note that recent theoretical works suggest that it can lead to topological superconductivity by careful control of the gate voltages [22–24]. A direct MI-induced exchange field in the Sm would provide an enhancement of the Zeeman energy in the Sm, enlarging the topological region in our case. We also neglect the magnetic orbital effects induced by the stray field of the MI as they are usually weaker than the ones induced by the exchange field [16].

If the MI considered is a magnetic semiconductor like EuS, the conduction band is not accessible at low temperatures. The modest band gap can be used to fabricate spin-filter tunnel barriers, using thin films, that allows spin-dependent tunneling through it when placed between two metallic regions [33]. The proposed setup is illustrated in Fig. 2(b).

In the following, we compare the two different situations represented in Fig. 2. For the case shown in Fig. 2(a), the MI only causes a spin splitting in the SC density of states. This situation is relevant for the geometry shown in Fig. 1(a) where a thick MI layer is placed in between the SC and MI layers having a weak effect on the Sm density of states. In the experiments [20], the SC and MI layers have a width of few nanometers while the typical size of a magnetic domain in the EuS is ~10δ0 for EuS-Al heterostructures [31]. Since the dimensions of the systems under consideration are smaller than δ0, we can disregard inhomogeneities and assume that the MI induces a homogeneous exchange field in the SC, which couples to the spin of the electrons by the Zeeman term $H_Z = V^{\text{MI}} \cdot \sigma$.

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sufficiently thin and, therefore, all the interfaces contribute to the tunneling processes.

To simplify the treatment, we consider a translation-invariant system along the longitudinal direction, \( z \), consisting of a single-channel Sm coupled to a Sc. The complete Hamiltonian of the system reads

\[
H = H_{\text{Sm}} + H_{\text{Sc}} + H_t. \tag{2}
\]

The Sm is described by

\[
H_{\text{Sm}} = \sum_n \left( \frac{p_z^2}{2m_{\text{Sm}}} - \mu_{\text{Sm}} \right) c_{p_z}^{\dagger} c_{p_z} + p_z c_{p_z}^\dagger \varepsilon_{\text{Sm}} c_{p_z}, \tag{3}
\]

where we use the spinors \( c_{p_z} = (c_{p_z \uparrow}, c_{p_z \downarrow}) \) and \( c_{p_z \sigma} \) is the electron annihilation operator in the Sm, while \( \varepsilon_{\text{Sm}} \) is the spin-orbit coupling strength.

The bare Hamiltonian for the Sc is given by

\[
H_{\text{Sc}} = \sum_n \xi_{np_z} c_{np_z}^\dagger c_{np_z} + \mathcal{A}_{np_z} c_{np_z}^\dagger + \mathcal{A}_0 c_{np_z} + \mathcal{V}(\xi_{np_z}, \sigma_{np_z} \varepsilon_{\text{Sm}}), \tag{4}
\]

where \( a_{np_z} \) is the electron annihilation operator for the mode \( n \) in the Sc and \( \xi_{np_z} = \frac{p_{np_z}^2}{2m_{\text{Sc}}} + \varepsilon_{n} - \mu_{\text{Sc}} \). We neglect the possible superconductive interband coupling and we assume singlet pairing in the parent superconductor gap matrix \( \mathcal{A}_{np_z} = \Delta_{0, np_z} \sigma_0 \). As anticipated, our Sc model features a homogeneous exchange field \( \varepsilon_{\text{Sc}} = \Gamma e \), induced by the nearby MI, which we consider to be aligned to the wire. A piece of crucial information is the distribution of the transverse modes in the Sc, \( \varepsilon_n \), which strongly depends on the device geometry. The energy separation between transverse modes can range from a value larger than the superconducting gap for very thin Sc film to zero for a bulk superconductor, where the modes above the gap form a continuum.

The Sc and Sm regions are coupled through a spin-dependent, momentum conserving, tunneling Hamiltonian

\[
H_t = \sum_{np_z} (c_{np_z}^\dagger T_{np_z} c_{np_z} + \text{H.c.}), \tag{5}
\]

where the hopping matrix \( T_{np_z} \) describes the electron tunneling processes taking place at the interfaces between the two materials. We write \( T_{np_z} \), in the basis of Pauli matrices \( T_{np_z} = t_{0, np_z} \sigma_0 + \sum_i t_{i, np_z} \sigma_i \). In the following, we set \( t_{x, np_z} = t_{y, np_z} = 0 \) as they are negligible for a uniformly polarized MI layer. Moreover, we consider only the case of real and positive \( t_{0, np_z} \) and \( t_{z, np_z} \). This is justified as we are assuming that spin-orbit coupling is absent in the Sc. An extended derivation with all the terms can be found in Appendix A3.

The bare Sc retarded Green’s function in the basis of time-reversed pairs reads

\[
G^R_{\text{Sc}}(\omega, n, p_z) = \left[ (\omega + \ii \varepsilon_{np_z}) \tau_0 \sigma_0 - \xi_{np_z} \tau_z \sigma_0 - \mathcal{A}_{np_z} \tau_z \sigma_0 - \mathcal{V}(\xi_{np_z}, \tau_z \sigma_0) \right]^{-1}, \tag{6}
\]

where the \( \tau_i \) are the Pauli matrices in the particle-hole space. In this work we ignore any back action of the Sm on the Sc as the electron density in the Sm is orders of magnitude smaller than the one in the Sc.

The retarded Green’s function of the Sm reads

\[
G^R_{\text{Sm}}(\omega, n, p_z) = \left[ (\omega + \ii \varepsilon_{np_z}) \tau_0 \sigma_0 - \xi_{np_z} \tau_z \sigma_0 - \mathcal{A}_{np_z} \tau_z \sigma_0 - \mathcal{V}(\xi_{np_z}, \tau_z \sigma_0) \right]^{-1}, \tag{6}
\]

where the self-energy \( \Sigma(\omega, n, p_z) = \sum_{n'} T_{np_z} G^R_{\text{Sc}}(\omega, n, p_z) T_{np_z} \) describes the coupling to the Sc. This allows us to write a low-energy effective model for the Sm. Since we focus on the quantum phase transition characterized by the gap closing, we work with the effective Hamiltonian \( \hat{H}_0 = H_{\text{Sm}} + \hat{H}_0 \) where the induced Hamiltonian is \( \hat{H}_0 = \Sigma(\omega = 0) \). We can split the effective Hamiltonian into three different contributions:

\[
\hat{H}_0 = \sum_n \left[ c_{np_z}^\dagger \tilde{\mathcal{A}}_0(p_z) \sigma_0 c_{np_z} + c_{np_z}^\dagger \mathcal{V}_0(p_z) \sigma_0 c_{np_z} \right] + \text{H.c.}, \tag{7}
\]

with \( \tilde{\mathcal{A}}_0(p_z) = \Delta_0(p_z) \sigma_0 + \sum_i \Delta_i(p_z) \sigma_i \). These three terms describe the shift in the chemical potential, the induced exchange field and the induced superconducting gap matrix in the Sm. The explicit forms of these contributions are

\[
\mathcal{V}_0^{(1)}(p_z) = \sum_n \left( \frac{\mathcal{V}_0^2(t_{0, np_z}^2 + t_{z, np_z}^2)}{\varepsilon_{np_z}^2 + \Delta_0^2} - \mathcal{V}_0 \right), \tag{8}
\]

\[
\mathcal{V}_0^{(2)}(p_z) = \sum_n \frac{-2\mathcal{V}_0 t_{z, np_z} t_{0, np_z}}{\varepsilon_{np_z}^2 + \Delta_0^2} - \mathcal{V}_0, \tag{9}
\]

\[
\Delta_0(p_z) = \sum_n \frac{\Delta_0^2}{\varepsilon_{np_z}^2 + \Delta_0^2} - \mathcal{V}_0, \tag{10}
\]

where we have divided the induced Zeeman term into two contributions \( \mathcal{V}_0 = \mathcal{V}_0^{(1)} + \mathcal{V}_0^{(2)} \). The first one is proportional to the splitting in the parent Sc \( \mathcal{V}_0 \), while the second one is linked to the spin-asymmetric tunneling amplitude of the barrier \( t_z \).

III. RESULTS

A. Semiconductor symmetrically coupled to a spin-split superconductor

In this section, we show that the combined superconducting and exchange proximity effects induced by the coupling to a spin-split Sc cannot induce, alone, a topological transition in the Sm. We consider a system like the one sketched in Fig. 2(a), where MI induces a spin splitting in the Sc. The Sc, in turn, induces superconductivity and an exchange field in the Sm. To check the presence of topological phases, we calculate the ratio between the induced exchange field \( \mathcal{V}_0 \), and the gap, \( \Delta_0 \). The topological phases appear when the condition in Eq. (1) is met, which leads to a closing of the superconducting gap at \( p_z = 0 \). For this reason, we focus on this point in momentum space in the following. A necessary
condition to satisfy the inequality is having a gap polarization ratio $|V_z/\Delta| > 1$.

Taking the ratio of Eqs. (8) and (10) for $t_z = 0$, we see that

$$\left| \frac{\tilde{V}_z}{\Delta_0} \right| = \frac{|V_z|}{\Delta_0}, \quad (11)$$

so the induced gap polarization ratio in the Sm is the same as in the Sc.

The gap polarization ratio $|V_{Sc}/\Delta_0|$ has to be less than unity in the parent Sc, otherwise superconductivity is suppressed in the whole device. In the case of a large homogeneous Sc, this ratio is limited by the stricter Chandrasekhar-Clogston bound, which dictates that a finite superconducting gap can only be maintained if $|V_{Sc}/\Delta_0| < 1/\sqrt{2}$ [36,37]. In thin Sc films, quantum confinement causes an increase of the superconducting pairing amplitude leading to a superconductive phase that survives under stronger exchange fields. However, the Chandrasekhar-Clogston limit still holds in terms of gap polarization ratio [35]. Even below this limit, the gap parameter in a clean Sc subjected to strong exchange fields ceases to be spatially homogeneous [38,39]. Therefore, it is not possible to obtain topological phases by coupling a spin-split Sc to a Sm.

This result is independent of the mode distribution in the Sc within the constant $\Delta_0$ and $V_z$ approximation [40]. The same result is found for a continuous flat density of transverse modes in the Sc where $\Delta_0$ is taken finite and constant for a wide range of energies being zero otherwise, as shown in Appendix B. In this case, the Zeeman term can be weakly enhanced, but this effect is totally negligible in realistic systems. This result can be generalized to the case where a multimode Sm is coupled to a multimode Sc as discussed in Appendix A2.

B. Spin-dependent tunneling

We consider now the case of a spin-asymmetric tunneling between the Sc and the Sm, taken as momentum independent for simplicity and described by $T = t_0 \delta_{0} + t_\sigma$. As Eqs. (8)–(10) show, the induced terms in the effective Hamiltonian are dependent on the distribution of the transverse modes in the Sc with respect to the chemical potential. In particular, they decay with the energy difference between the bottom of the Sc subband and the chemical potential. This means that the modes closer to the Sm Fermi energy give the dominant contribution to the induced superconducting pairing and exchange fields at $p_z = 0$.

We first analyze the contribution of an isolated Sc mode to the effective Hamiltonian, as the ones from the different modes add up. The behavior of the induced term in the effective Hamiltonian is illustrated in Fig. 3. The two terms of the induced exchange field $V_z^{(1)}$ and $V_z^{(2)}$ sum constructively for $\mu_{Sc} > \varepsilon_n$. Both $V_z$ and $\Delta_0$ decay for $|\varepsilon_{n,0}| \to \infty$, leading to the existence of an optimal regime where both the induced exchange field and superconducting pairing are maximal and $|V_z/\Delta_0| > 1$. This is the ideal region for searching for topological superconductivity.

In general, a spin-dependent component $t_z$ in the tunneling matrix enhances the collinear component in the exchange field $V_z$, suppressing the noncollinear components in $\tilde{V}$ and the superconducting gap in the Sm. From Eqs. (8)–(10) we see that $\Delta$ and $V_z^{(1)}$ share the same dependence on the Sc band structure (they decay as $\sim \xi^{-2}$) while having a different prefactor which depends on the tunneling matrix. Therefore, the spin splitting in the Sc leads to an enhancement of $|V_z|$ and provides a first mechanism to induce topological superconductivity in the Sm.

In contrast to $V_z^{(1)}$, the second contribution $V_z^{(2)}$ to the induced exchange field in the semiconductor is totally independent of the spin polarization in the parent superconductor. This term depends solely on the spin-asymmetric component of the tunneling Hamiltonian. Moreover, $V_z^{(2)}$ has an energy dependence $\sim \xi^{-1}$, with a sign that depends on the relative position of the mode to the Sc Fermi energy. Since $\Delta_0$ and $V_z^{(1)}$ decay as $\sim \xi^{-2}$, $V_z^{(2)}$ dominates as the energy difference between the Sc and Sm modes increases. This contribution exhibits a sign change, which leads to a cancellation of $V_z^{(2)}$ in the limit of small energy separation between modes. However, in thin Scs, the transverse modes can exhibit a large energy separation because of quantum confinement effects.

This result suggests that there are two different mechanisms that can drive the system to the topological phase. First, the spin-dependent tunneling enhances the induced gap-polarization ratio of the spin-split superconductor. Alternatively, the spin-dependent tunneling combined with the quantum confinement in the Sc can lead to the appearance of the topological phase in regimes where, otherwise, the phase diagram would be globally trivial.

C. Combined proximity effect dominated by one transverse mode

If the separation between transverse modes in the Sc is large enough, the contribution of the mode closest to the chemical potential dominates over the other ones. We can visualize this system like two coupled one-dimensional wires, where one features both spin splitting and superconductivity.
The first conclusion is that the effect of the spin splitting of the Sc is negligible for realistic parameters, as it appears summed to the Sc Fermi energy which is many orders of magnitude bigger [Eqs. (12)]. Analyzing the behavior with respect to the tunneling parameters, it can be noticed that there exists a region corresponding to $t_0 \sim \pm \sqrt{-\xi \mu_{Sm} + \mu_{Sm\gamma} \mu_{Sc}}$, where a small polarization of the barrier can induce a topological phase with a relatively big gap. Since the chemical potential in the nanowire is tunable using electrostatic gates, it is, in principle, possible to set the operating point of the device near this optimal point.

D. Superconductor with several relevant transverse modes

In a more realistic scenario, the Sc features many transverse modes. In this case, Eqs. (8)–(10) explain how the induced terms in the Sm effective Hamiltonian are determined by the sum of the contributions of each mode. The relative energy of the transverse modes in the Sc strongly affects the gap polarization in the Sm, becoming a critical factor for the appearance of the topological phase. For the systems of interest, the dimensions of the Sc section vary from a few to hundreds of nanometers. For small length scales, the effect of quantum confinement separates the Sc subbands such that it is not possible to treat the density of state like a continuum. Both $\Delta_0$ and $V_z^{(1)}$ have a Lorentzian shape with a full width at half maximum $\Gamma = \sqrt{\Delta_0^2 - V_z^2}$. We can use $\Gamma$ as a reference to distinguish between discrete modes, $\delta \varepsilon \gg \Gamma$, and a continuum distribution of modes, $\delta \varepsilon \ll \Gamma$, where $\delta \varepsilon$ is the average separation of modes. Indeed, if the average energy separation between the transverse modes is $\delta \varepsilon \gtrsim \Gamma$, the resonance peaks do not overlap entirely and the resonance peaks of $V_z^{(2)}$ do not completely cancel out. In this case, the net exchange field experienced in the Sm will be due to the sum of the contribution of each transverse mode.

To provide a clearer picture, we consider the case where the band structure of the Sc can be described by perfectly equidistant transverse modes. In this case, $\delta \varepsilon$ is the relative separation between the Sc modes. This approximation recovers the continuum limit for a 2D system, where we expect a constant density of transverse modes $g_{\beta} = (\delta \varepsilon)^{-1}$. We also define $\delta \mu_{Sc}$ as the relative position of the Fermi energy in the Sc from the middle of the band. We proceed by summing up the modes contributions following Eqs. (8)–(10) to calculate the value of the induced terms in the effective Hamiltonian (calculation details can be found in Appendix D). The behavior of the induced terms for $p_z = 0$ is illustrated in Fig. 5.

By varying the chemical potential in the Sc, as the subbands cross the Fermi energy we see an alternation of trivial regions (gray background) and regions where topological superconductivity can appear by tuning $\mu_{Sm}$ (white background). This behavior is more clearly illustrated by the topological invariant and the effective gap, shown in Fig. 6. In this figure, the effective gap is given by the smallest eigenvalue of the effective Hamiltonian (7).

The observation of an alternation of topological and trivial phases while varying $\delta \mu_{Sc}$ is consistent with Fig. 5. The limit $g_{\beta} \ll 1/\Delta_0$ corresponds to a large Sc where the effect of quantum confinement becomes completely negligible.
This behavior can be realized if the average separation of the transverse modes is such that $\delta \varepsilon \ll \Gamma$. In this limit, the contributions from each mode exchange field, leading to a globally trivial or topological phase, depending on the polarization of the spin-filter barrier $t_z/t_0$. A quantitative criterion can be obtained by integrating Eqs. (8)–(10) over a constant density of transverse modes as shown in Appendix B. In this limit, the term in Eq. (9) vanishes and we get for the gap polarization

$$\frac{V_z}{\Delta_0} = \frac{V_z}{\Delta_0} \frac{t_z^2 + t_0^2}{t_z^2 - t_0^2}. \tag{13}$$

We note that the spin-dependent tunneling leads to an enhancement of the induced exchange field, while reducing the induced superconducting gap. This effect can be used to bring the induced gap polarization ratio above 1, closing the gap at $p_z = 0$ and inducing a quantum phase transition to the topological phase. Therefore, the spin-asymmetric tunneling provides a way to overcome the limitation of Eq. (11) and to induce a phase transition to the topological phase, Fig. 2(b). The topological phase appears for a barrier polarization

$$\frac{t_z}{t_0} > \sqrt{\frac{1 - |V_z/\Delta_0|}{1 + |V_z/\Delta_0|}}. \tag{14}$$

IV. DISCUSSION

In this section, we discuss the applicability of the spin tunneling mechanism for generating topological superconductivity to the case of EuS-Al-InAs platforms.

EuS is a magnetic semiconductor with an optical band gap of $E_g = 1.65$ eV, which can be effectively tuned by quantum confinement and, therefore, it depends on the film thickness [41]. The magnetic properties of this material are attributed to the Eu atoms which are characterized by strongly localized half-filled $4f$ shells. They behave as localized spins with a magnetic moment of around $7\mu_B$. For this reason, the EuS can be well described as a Heisenberg ferromagnet with a Curie temperature of $T_m = 16.6$ K [42,43]. The magnetization induces a splitting of the conduction band of around $\Delta E_c = 0.36$ eV [44]. These properties are very promising in the view of fabricating spin-filter tunnel barriers.

As discussed in the previous section, the behavior of the system is strongly dependent on the band structure of the EuS and the topological phase is induced by adding a spin-asymmetric tunneling to enhance $V_z$ and suppress $\Delta_0$, Eq. (13).

To check if the ferromagnetic hybrid nanowire is in this regime, we can estimate the average mode separation by a simple particle in a box model $\delta \varepsilon \simeq \frac{\hbar^2}{2m_sc^2} \frac{\pi^2}{Lc^2}$, where $Lc$ refers to the largest dimension of the cross section of the EuS shell. Assuming that $\Gamma \sim 100$ $\mu$eV, which is in line with the experimental measurements of EuS-Al heterostructures [20,31], we estimate that, in order to observe well-separated modes, the maximum dimension of the shell should be in the order of 60 nm. In experiments, the facet length is around 60 nm [20]. For this reason, we expect the contributions of the modes in the EuS to overlap significantly.

Previous measurements performed on EuS-Al heterostructures have shown a polarization of around 50% of the gap [31]. In this case, applying (14), we can estimate that a spin-dependent barrier with a 58% polarization is enough to bring the Sm to the topological regime.

![FIG. 5. Induced gap $\Delta_0$ (blue line), exchange field contributions $\tilde{V}_z^{(1)}$ and $\tilde{V}_z^{(2)}$ (dashed lines), and total spin splitting $\tilde{V}_z = \tilde{V}_z^{(1)} + \tilde{V}_z^{(2)}$ (red line) at $p_z = 0$ in the Sm in the case of a multimode Sc. Parameters: $\Delta_0 = 1$, $t_0 = 1$, $t_z = 0.4$, $V_z = 0.5$, $\delta \varepsilon = 10$. The area with gray background is topologically trivial while the one with white background satisfies the condition $|V_z| > \Delta_0$. The effect of the resonance peaks in $\tilde{V}_z^{(2)}$ cause an oscillation between topologically trivial and nontrivial phases while varying the chemical potential in the Sc.

![FIG. 6. Phase diagram for a single mode Sm coupled to a multimode Sc, showing the effective gap in the Sm $|\Delta_{eff}| = \min_{n,p} |E_{n,p}|$, where $E_{n,p}$ are the Hamiltonian eigenvalues and sign given by the topological invariant $W_{2n}$. We observe an oscillation between the topological (red) and trivial (blue) phases as we vary the chemical potential from the middle of the gap $\delta \mu$Sc. As the density of transverse modes $g_{\alpha}$ increases, the peaks get closer and overlap until they merge in the continuous limit $g_{\alpha} \gg (\Delta_0)^{-1}$. In the continuum limit, the system is in a globally trivial or topological phase, depending on the condition in Eq. (14). We assume that the chemical potential in the Sm is tuned such that $\mu_{Sm} = -\tilde{\mu}$. Parameters used: $m_{sc}^* = 1$, $m_{sm}^* = 0.2$, $\alpha_c/\Delta_0 = 20$, $V_z/\Delta_0 = 0.5$, $t_0/\Delta_0 = 1$. For the left panel we use $t_z/\Delta_0 = 0.4$, while for the right panel we use $t_z/\Delta_0 = 0.7$.}
Estimating the optimal magnetic barrier length to achieve this polarization is a complicated task as strong spin-polarized band bending is expected at the interfaces. As the barrier gets thicker, we expect a stronger polarization, but at the same time, the coupling between the two systems gets strongly suppressed. Therefore, the optimal barrier length would be determined by the trade-off between a strongly polarizing thick barrier, which, however, suffers low transparency, and a thin transparent barrier with low polarization.

The introduction of MI in 2DEG is more challenging as the geometry and fabrication constraints make the introduction of Sc–Sm, MI–Sc, and MI–Sc interfaces at the same time particularly complicated. The use of MI as a spin-filter tunnel barrier simplifies the design of topological quantum devices by eliminating the need of the Sc–Sm interface. Using this new operating principle, the geometry sketched in Fig. 2(b) can provide a viable alternative to achieve zero-field topological superconductivity in 2DEG devices.

Finally, we note that both nanowires Sc shell and Sc layers in the 2DEG system are around 5 nm thick, so we can already expect a measurable effect of quantum confinement in this direction. For this reason, we expect that the resonance effect discussed previously can be measured by scaling down the width of the devices.

V. CONCLUSIONS

In this work, we have demonstrated that a spin-split superconductor cannot induce, alone, topological superconductivity in a spin-orbit coupled semiconductor by the combined superconducting and magnetic proximity effect. This is in contradiction with the hypothesis that an exchange field in the superconductor can induce a topological transition in the semiconductor. We have shown how a spin-filter tunnel barrier, provided by a thin magnetic insulating layer between the semiconductor and the superconductor, can be used to overcome this limitation. The spin-dependent tunneling suppresses the induced superconducting pairing potential while enhancing the spin splitting, thus providing a way to bring the system to the topological phase. In the case of a distribution of discrete modes in the superconductor, the phase diagram exhibits an alternation of trivial and topological regions as a function of the chemical potential of the superconductor due to a change on the sign in the contributions to the exchange field. The total induced exchange field strongly depends on the relative energy difference between the closest levels to the Fermi energy in the superconductor. While this mechanism is unlikely to explain the results in Ref. [20] because of the relatively thick EuS layer used, the concept of spin-dependent coupling can be exploited in the next generation of topological superconducting devices without magnetic field. Spin-filter tunnel barriers can be achieved by depositing a thin film of few nanometers of a ferromagnetic insulator at the interface between the superconductor and the semiconductor. The proposed mechanism is compatible with the currently used hybrid superconductor-semiconductor platforms, including nanowires and 2DEG systems.

Note added. Recently, several independent papers on the same subject have been made publicly available as preprints. The result concerning the impossibility of topological phases by superconductor-mediated exchange field has been extended to generic systems in terms of a more general requirement about the minimal Zeeman field required for a topological phase transition [45]. The ferromagnetic hybrid nanowire physics has been explored using a self-consistent treatment of a diffusive superconductor in Ref. [46]. The results are in agreement with the ones presented in our work. Another independent work has found that perfectly clean systems with an extremely thin superconductive layer could be used to obtain topological phases in the MI–Sc–Sm stack [47].

ACKNOWLEDGMENTS

This research was supported by the Danish National Research Foundation, the Danish Council for Independent Research—Natural Sciences, and the Microsoft Corporation. This project has received funding from the European Research Council (ERC) under the European Union’s Horizon 2020 research and innovation programme under Grant Agreement No. 856526. R.S.S. and M.L. acknowledge funding from QuantERA project 2D hybrid materials as a platform for topological quantum computing and from NanoLund.

APPENDIX A: COMPLETE SELF-ENERGY CALCULATION

In this Appendix, we discuss more in detail the self-energy model employed to derive the effective Hamiltonian discussed in Eq. (7). We derive the model, discussing the frequency dependence, the presence of multiple modes in the Sm, and the effect of a general hopping matrix with arbitrary direction of the magnetization axis. Our model consists in a translation-invariant system composed by a Sm region coupled to a Sc in which an exchange field is induced by the nearby MI. This system is described by the Hamiltonian

$$ H = H_{\text{Sm}} + H_{\text{Sc}} + H_t, $$

where the three terms read

$$ H_{\text{Sm}} = \sum_{mp} \varepsilon_m(p_c) c_{m\sigma}^\dagger c_{m\sigma} + \sum_{p} c_p c_{p\sigma}^\dagger \sigma \gamma c_{p\sigma}, $$

$$ H_{\text{Sc}} = \sum_{np} \varepsilon_n(p_z) a_{n\sigma}^\dagger a_{n\sigma} + V_z a_{np}^\dagger \gamma a_{np} + [a_{np}^\dagger (i\sigma \Delta_{np}) a_{n-p\sigma}] + \text{H.c.}, $$

$$ H_t = \sum_{mnzp} \varepsilon_{mp} T_m(p_z) a_{n\sigma}^\dagger + \text{H.c.}. $$

We adopt the spinor basis $c_p = (c_p^\uparrow c_p^\downarrow)$, where $c_{p\sigma}$ is the electron annihilation operator in the nanowire for spin $\sigma$. Both the Sm and the Sc have a multimode structure. In Eq. (A2), we label the Sm modes by the index $m$ and we write the kinetic energy as $\varepsilon_m(p_c) = \varepsilon_m + \frac{p_c^2}{2m_{\text{Sm}}}$, where $m_{\text{Sm}}$ and $\sigma$ are the electron effective mass and the spin-orbit coupling parameter in the Sm. The Hamiltonian in Eq. (A3) describes the Sc, where $a_{np}$ is the electron annihilation operator for mode $n$ with kinetic energy as $\varepsilon_n(p_z) = \varepsilon_n + \frac{p_z^2}{2m_{\text{Sc}}}$, $m_{\text{Sc}}$. Here, $\Delta_{np} = \Delta_{0np}\sigma_0$ is the superconducting pairing for each band. We neglect superconducting interband coupling and the proximity-induced exchange field by MI in the Sc.
The two regions are coupled by a spin-dependent tunneling Hamiltonian with hopping matrix $T_{mn}(p_z)$, which takes into account all the different electron transmission processes taking place at the interfaces between the different materials in the device. We can write $T_{mnp}$ in the basis of Pauli matrices:

$$T_{mnp}(p_z) = t_{0,mn}(p_z)\sigma_0 + \sum_{l=xyz} t_{l,mn}(p_z)\sigma_l, \quad (A5)$$

where the $t_l$ matrix is the tunneling Hamiltonian written in the basis of Pauli matrices.

The full Green’s function for the semiconductor subspace reads

$$G^R_{Sm}(\omega, n, p_z) = \frac{1}{\omega + i0^+ - \mathcal{H}_{Sm} - \Sigma(\omega, p_z)}, \quad (A7)$$

where $\mathcal{H}_{Sm}$ is the bare Sm Hamiltonian matrix and we have defined the self-energy as

$$[\Sigma(\omega, p_z)]_{mn} = \sum_n \mathcal{H}_{l,mn}(p_z)G^R_{Sm}(\omega, n, p_z)\mathcal{H}_{l,m'n}(p_z), \quad (A8)$$

where the $\mathcal{H}_l$ matrix is the tunneling Hamiltonian written in particle-hole space

$$\mathcal{H}_{l,mn} = [t_{0,mn}(p_z)\tau_z + it_{l,mn}(p_z)\omega]\sigma_0 + \sum_l [t_{l,mn}(p_z)\tau_0 + it_{l,mn}(p_z)\tau_l]\sigma_l. \quad (A9)$$

We can see that the first-order terms are proportional to $\omega$.

$$\mathcal{H}_{l,mn} = \frac{1}{\omega + i0^+ - \mathcal{H}_{Sm} - \Sigma(\omega, p_z)}\mathcal{H}_{l,m'n}(p_z),$$

and we denote by $P(\omega, p_z)$ the remaining $\tau_\sigma\sigma_\tau$ component.

The terms in $\mathcal{H}_{l,m'n}$ parameters are, in general, complex numbers.

In the following, we will use the basis of time-reversed pairs, where $\tau_l$ are Pauli matrices for the particle-hole space, being $\mathcal{H}$ the total Hamiltonian in this basis. In this basis, the bare $\mathcal{S}$ retarded Green’s function reads

$$G^R_{Sm}(\omega, n, p_z) = \frac{1}{\omega + i0^+ - \mathcal{H}_{Sm} - \Sigma(\omega, p_z)}, \quad (A7)$$

where $\mathcal{H}_{Sm}$ is the bare Sm Hamiltonian matrix and we have defined the self-energy as

$$[\Sigma(\omega, p_z)]_{mn} = \sum_n \mathcal{H}_{l,mn}(p_z)G^R_{Sm}(\omega, n, p_z)\mathcal{H}_{l,m'n}(p_z), \quad (A8)$$

where the $\mathcal{H}_l$ matrix is the tunneling Hamiltonian written in particle-hole space

$$\mathcal{H}_{l,mn} = [t_{0,mn}(p_z)\tau_z + it_{l,mn}(p_z)\omega]\sigma_0 + \sum_l [t_{l,mn}(p_z)\tau_0 + it_{l,mn}(p_z)\tau_l]\sigma_l. \quad (A9)$$

Introducing the expansion in the self-energy term and keeping the zeroth and first order for each component, we see that the $\tau_0\sigma_0$ component reads

$$\tilde{E}(\omega, p_z) = -\omega \sum_n \left(\frac{1}{\omega} + \frac{t_z^2 + t_l^2}{\xi_{np}^2 + \Delta_{0,np}^2}\right)\sigma_0 + O(\omega^2). \quad (A11)$$

The semiconductor Green’s function can be rewritten like

$$G_{Sm}(p_z) = \frac{Z(p_z)}{\omega - Z(p_z)[\mathcal{H}_{Sm}(p_z) - \tilde{H}_0(p_z) + \tilde{H}_1(\omega, p_z)]}, \quad (A13)$$

where $Z(p_z) = [1 + C(p_z)]^{-1}$ with $C(p_z) = -\omega\tilde{C}(p_z)$.

We can see that the first-order terms are proportional to $\omega^{3/2}/\xi_{np}^2 + \Delta_{0,np}^2 - V_z^2 < \omega^2$. Therefore, provided that we are interested in the $|\omega/\Delta_0| \ll 1$ region, we can ignore these terms. This approximation breaks down in the strong-coupling

$$\mathcal{H}_{l,mn} = \frac{1}{\omega + i0^+ - \mathcal{H}_{Sm} - \Sigma(\omega, p_z)}\mathcal{H}_{l,m'n}(p_z),$$

and we denote by $P(\omega, p_z)$ the remaining $\tau_\sigma\sigma_\tau$ component.
regime, namely, when the tunneling amplitudes overcome $\Delta_0$. In this case, while the phase diagram can still be studied within the $\omega \to 0$ limit, Eq. (A12) overestimates the size of the induced gap. For a more complete reference to the strong tunneling regime, we refer to [48].

2. Multimode semiconductor

Our objective now is to generalize the previous study to the case of a multimode Sm in the case of spin-symmetric tunneling ($t_\perp = 0$). Since we focus on the phase diagram, we consider the lowest order in the effective Hamiltonian $\hat{H}_0 = \Sigma(\omega = 0, p_\perp)$. The three contributions read

\begin{equation}
\check{\mu}(p_\perp) = \sum_n \frac{\xi_{np}}{\xi_{np}^2 + \Delta_{0, np}^2} \left[ t_{0, np} f_{0, np} t_{0, mp} - V_z^2 \right].
\end{equation}

\begin{equation}
\check{V}_z^{(1)}(p_\perp) = \sum_n \frac{V_z}{\xi_{np}^2 + \Delta_{0, np}^2} \left[ t_{0, np} f_{0, np} + t_{0, n} t_{0, mp} - V_z^2 \right],
\end{equation}

\begin{equation}
\Delta_0(p_\perp) = \sum_n \left[ (t_{0, np} f_{0, np} - t_{0, n} t_{0, mp}) \tau_z - (t_{0, np} f_{0, np}^* + t_{0, n} t_{0, mp}^*) \tau_x \right] \times \frac{\Delta_{np}}{\xi_{np}^2 + \Delta_{0, np}^2 - V_z^2}.
\end{equation}

For $p_\perp = 0$, $t_0$ real and $\Delta_{0, np} = \Delta_0$, we can write the effective Hamiltonian matrix as

\begin{equation}
\hat{H}_0 = \Sigma \otimes \sigma_0 \tau_z + M \otimes (V_z \sigma_0 + \Delta_0 \sigma_0 \tau_z),
\end{equation}

where we have defined the $\Xi$ matrix as

\begin{equation}
\Xi_{mnr} = \xi_m \delta_{mn} - \sum_n \frac{t_{0, np} f_{0, np} t_{0, mp}}{\xi_{np}^2 + \Delta_{0, np}^2 - V_z^2},
\end{equation}

where $\xi_m$ are the energies of the Sm modes. The $M$ matrix is

\begin{equation}
M_{mnr} = \sum_n \frac{t_{0, np} f_{0, np} t_{0, mp}}{\xi_{np}^2 + \Delta_{0, np}^2 - V_z^2}.
\end{equation}

The $M$ matrix is Hermitian so it is diagonalizable as $M = U \Lambda U^\dagger$. We can apply the unitary transformation $U \otimes \tau_0 \sigma_0$ to the Hamiltonian matrix and get to

\begin{equation}
\hat{H}_{\text{eff}} = U^\dagger \Xi U \otimes \sigma_0 \tau_z + \Delta \otimes (V_z \sigma_0 + \Delta_0 \sigma_0 \tau_z) + \check{\mu}(p_\perp) \cdot \sigma c_p.
\end{equation}

where the last term is the diagonal element $\check{\mu}_m$. The off-diagonal hopping terms in the matrix $\check{\mu}$ come from the transformed $U^\dagger \Xi U \otimes \sigma_0 \tau_z$. This Hamiltonian is a combination of independent modes $\sum_m \xi_m \tau_z \sigma_0 + \lambda_m (V_z \sigma_0 + \Delta_0 \sigma_0 \tau_z)$ interacting through the spin-independent $\check{\mu}$ terms which acts as a tunneling Hamiltonian for these channels. The gap polarization ratio is given by $|V_z/\Delta_0|$ for these channels, which is the same as in the parent Sc. For this reason, if the parent Sc is in a trivial state, all the modes in the Sm will be in the trivial regime as well. The eigenmodes of the full effective Hamiltonian will be given by superpositions of these states. However, since we are mixing trivial states we can only obtain trivial states as a result.

3. Case of general tunnel matrix with complex element

We now move to the case of a general tunneling matrix $T_{np}$, coupling a multimode Sc to a monomodal Sm (thus we drop the index $m$). In this case we consider only the zeroth-order term in $\omega$ which is the effective Hamiltonian $\hat{H}_0$. We can split it into three different contributions:

\begin{equation}
\hat{H}_0 = \sum_{p_\perp} \check{\mu}(p_\perp) c_{p_\perp}^\dagger \sigma_0 c_{p_\perp} + \check{V}(p_\perp) \cdot \sigma c_{p_\perp} + \check{\Delta}(p_\perp) \sigma_0 + \sum_i \check{\Delta}_i(p_\perp) \sigma_i
\end{equation}

with $\check{\Delta}(p_\perp) = \check{\Delta}_0(p_\perp) \sigma_0 + \sum_i \check{\Delta}_i(p_\perp) \sigma_i$ and $\check{V} = \check{V}^{(1)} + \check{V}^{(2)}$. The three terms are given by

\begin{equation}
\check{\mu}(p_\perp) = \sum_n \frac{\xi_{np}}{\xi_{np}^2 + \Delta_{0, np}^2} \left[ |t_{0, np}|^2 + |t_{x, np}|^2 + |t_{y, np}|^2 + |t_{z, np}|^2 \right] - V_z \left[ |t_{0, np} t_{0, np}^* + \text{H.c.}| + (|t_{x, np} t_{x, np}^* + \text{H.c.}|) \right],
\end{equation}

\begin{equation}
\check{V}_z^{(1)}(p_\perp) = \sum_n \frac{V_z}{\xi_{np}^2 + \Delta_{0, np}^2} \left[ |t_{0, np}|^2 + |t_{z, np}|^2 - |t_{x, np}|^2 - |t_{y, np}|^2 \right],
\end{equation}

\begin{equation}
\check{V}_z^{(2)}(p_\perp) = \sum_n \frac{\xi_{np}}{\xi_{np}^2 + \Delta_{0, np}^2} \left[ (|t_{0, np} t_{x, np}^* + \text{H.c.}| + (|t_{x, np} t_{x, np}^* + \text{H.c.}|) \right] \epsilon_{ijk},
\end{equation}

\begin{equation}
\check{\Delta}_0(p_\perp) = \sum_n \frac{\Delta_{np}}{\xi_{np}^2 + \Delta_{0, np}^2 - V_z^2} \left[ t_{0, np}^2 - t_{x, np}^2 - t_{y, np}^2 - t_{z, np}^2 \right].
\end{equation}

where $\epsilon_{ijk}$ is the Levi-Civita symbol. We can see that the presence of transverse components in the tunneling matrices causes a reduction of both the induced exchange field and superconducting pairing potential. However, if these amplitudes are complex, a nontrivial resonance contribution can appear in the $V_z^{(1)}$ term enhancing the induced exchange field. Moreover, when complex amplitudes are present the induced gap can show interference effects due to the coupling between the various Sc bands induced by the Sm.

APPENDIX B: BCS INTEGRATION-INDUCED GAP AND EXCHANGE FIELD

In this section we integrate the induced gap and exchange field due to the coupling to the superconductor. We take the
usual BCS approximation which considers an attractive interaction between electrons with absolute energies below \( \omega_D \). Therefore, the parent gap becomes finite and equal to \( \Delta_0 \) for energies in the range \([-\omega_D, \omega_D]\), being zero otherwise. Using this approximation, the induced gap is given by
\[
\tilde{\Delta}_0 = 2 \Delta_0 g_m (t_0^2 - t_z^2) \frac{\arctan(\omega_D/\sqrt{\Delta_0^2 - V_z^2})}{\sqrt{\Delta_0^2 - V_z^2}}, \tag{B1}
\]
where \( g_m \) is the density of transverse modes, considered to be energy independent. The induced exchange field is given by
\[
\tilde{V}_z = V_z g_m (|t_0|^2 + |t_z|^2) \frac{\pi}{\sqrt{\Delta_0^2 - V_z^2}}, \tag{B2}
\]
where the superconductor bandwidth has been taken infinite for simplicity.

For \( t_z = 0 \), \( \tilde{V}_z/N_c \geq \tilde{\Delta}_0/\Delta_0 \), where the equal case corresponds to the limit \( \omega_D \to \infty \). Taking realistic numbers for bulk aluminum, where \( \omega_D/\Delta_0 \approx 200 \), we find \( \tilde{V}_z/N_c < \tilde{\Delta}_0/\Delta_0 \). Therefore, we consider that the Sc proximitized by the MI fulfills the Chandrasekhar-Clogston limit, \( V_z < \Delta_0/\sqrt{2} \) [36,37], the spin splitting cannot cause a topological transition in the Sm, unless \( \omega_D/V_z \sim 1 \). We do not expect that an energy-dependent tunneling amplitude significantly changes this conclusion.

APPENDIX C: ANALYTIC DERIVATION OF THE PFAFFIAN

In this section we analytically derive the topological invariant of the two-wires system, namely, a system composed by a single mode Sm coupled to a single mode Sc. We can write the Hamiltonian matrix as
\[
\mathcal{H} = \begin{pmatrix} \mathcal{H}_{\text{Sm}} & \mathcal{H}_t \\ \mathcal{H}_t^\dagger & \mathcal{H}_{\text{Sc}} \end{pmatrix}, \tag{C1}
\]
where the three terms read
\[
\mathcal{H}_{\text{Sm}} = \left( \frac{p_z^2}{2m^2_{\text{Sm}}} - \mu_{\text{Sm}} \right) \sigma_0 \tau_z + \alpha \sigma_p \tau_p \tau_z, \tag{C2}
\]
\[
\mathcal{H}_t = \frac{p_z}{2m_{\text{Sc}}} V_z g_m (|t_0|^2 + |t_z|^2) \cot[i \pi g_m (i \Gamma + \delta \mu - \frac{p_z^2}{2m_{\text{Sc}}})] + \frac{i \pi g_m (i \Gamma - \delta \mu + \frac{p_z^2}{2m_{\text{Sc}}})]}{i \Gamma}, \tag{D1}
\]
\[
\mathcal{H}_t^{(2)}(p_z) \simeq \frac{\pi g_m \alpha \tau_z}{2} \cot[2 \pi g_m (i \Gamma + \delta \mu - \frac{p_z^2}{2m^2_{\text{Sc}}})] - \cot[2 \pi g_m (i \Gamma - \delta \mu + \frac{p_z^2}{2m^2_{\text{Sc}}})], \tag{D2}
\]
\[
\tilde{\Delta}_0(p_z) \simeq \frac{\pi g_m \alpha \tau_z}{2} \cot[2 \pi g_m (i \Gamma + \delta \mu - \frac{p_z^2}{2m^2_{\text{Sc}}})] + \cot[2 \pi g_m (i \Gamma - \delta \mu + \frac{p_z^2}{2m^2_{\text{Sc}}})]. \tag{D3}
\]

To simplify the expression, we can define the auxiliary functions
\[
F(x, y) = \frac{\cot(ix + y) - \cot(ix - y)}{2} = \frac{\sin(2y)}{\cosh(2x) - \cos(2y)}, \tag{D4}
\]
\[
G(x, y) = i \frac{\cot(ix + y) + \cot(ix - y)}{2} = \frac{\sinh(2x)}{\cosh(2y) - \cos(2x)}. \tag{D4}
\]
We finally get to compact expressions for the effective Hamiltonian terms:

\[ \tilde{V}_z^{(1)}(p_z) \simeq \pi g_m \frac{V_z}{\Gamma} \left( t_0^2 + t_z^2 \right) G \left[ \pi g_m \Gamma, \pi g_m \left( \delta \mu - \frac{p_z^2}{2m_{Sc}} \right) \right], \]  
\[ \tilde{V}_z^{(2)}(p_z) \simeq 2\pi g_m t_0 F \left[ \pi g_m \Gamma, \pi g_m \left( \delta \mu - \frac{p_z^2}{2m_{Sc}} \right) \right], \]  
\[ \tilde{\Delta}_0(p_z) \simeq -\pi g_m \frac{\Delta_0}{\Gamma} \left( v_0^2 - t_z^2 \right) G \left[ \pi g_m \Gamma, \pi g_m \left( \delta \mu - \frac{p_z^2}{2m_{Sc}} \right) \right]. \]

This expression can be used to derive the Sm effective Hamiltonian and, with a simple linear search for the minimum eigenvalue, find the effective gap shown in Fig. 6.

[40] If we allow the pairing potential to be different in each subband, the induced gap polarization ratio could take any value in the
range of \([\min(V_i/\Delta_{0,a}), \max(V_i/\Delta_{0,a})]\) depending on the coupling terms \(T_n\). For a gapped superconductor all the subbands close to the Fermi level fulfill \(\Delta_{0,a} > V_i\). With this constraint, it is still not possible to get topological superconductivity in a semiconductor coupled to a spin-split superconductor through a spin-symmetric barrier.


