We discuss and Sec. IV V computation is braiding, which could be performed "topological" qubit operations tuning the pairwise interaction between the MFs ("non-topological" qubit operations) were also discussed in the literature, Refs. 11–16). Different structures have been proposed as possible hosts for MFs, especially in connection with the future possibility of topologically-protected quantum computations.

It is generally assumed that MFs localize at some heterogeneity separating media with different topological numbers, which may be a vortex core in a superconductor or superfluid, sample boundaries, interfaces in heterostructures, etc. (see, e.g., Refs. 11–16). The attachment of the MF to such a physical "defect" may be disadvantageous for a number of reasons: the presence of physical inhomogeneities in the system requires additional efforts at the fabrication stage, it introduces extra disorder, and creates interfaces, whose properties are difficult to control. In particular, binding a MF to some point in space limits its control and manipulation options.

However, it is not a general law that MF must be localized near some "topological defect". Below we consider two models of a two-dimensional (2D) topological superconductor hosting a vortex. Systems of this type have been studied experimentally, and hints of a Majorana state at the vortex core were reported. From the theory standpoint, these models are particularly interesting, because, as it will be demonstrated below, they may serve as a platform where a confinement mechanism unrelated to a physical heterogeneity is realized. The MFs can arise only in pairs, since only a superposition of two MFs has a physical sense. Therefore, in addition to the well-studied MF at the core, a second, exterior, Majorana state can emerge. The second MF is not necessarily pinned by some interface or sample edge. We show that it can be localized by a finite magnetic field. For a uniform field, the wave function of the exterior MF can be calculated exactly. Its weight is centered at some field-dependent mesoscopic radius away from the vortex center. In other words, the wave function is localized at a circle of radius . Varying the magnetic field, we may control . This feature could be useful for the manipulation of Majorana states. We will analyze the tunability of the pairwise splitting between MFs in the case of two vortices.

The paper is organized as follows. In Sec. II, we discuss the generation of the Majorana fermions in the topological insulator/superconductor heterostructure. In Sec. III, similar ideas are applied to p-wave superconductor. The discussion and conclusions are in Sec. IV and Sec. V respectively.

II. TI/SC HETEROSTRUCTURE

Given the recent success in the fabrication of TI/SC heterostructures, let us discuss this system first. Our model describes the 2D states at the surface of the TI. Proximity to the superconductor induces superconducting correlations in these states. An external magnetic field inserts a vortex with integer vorticity into the system, which is schematically shown in Fig. I. The cor-
The corresponding eigenenergy is zero, \( \varepsilon_{\text{MF}} = 0 \).

We introduce a spinor \( F^\mu(r) = (f_1^\mu, f_2^\mu, f_3^\mu, f_4^\mu)^T \) as
\[
\psi = \exp[-i\phi(t_\tau - \sigma_3)/2 + i\mu\phi]F^\mu(r),
\]
where the index \( \mu \) represents the total angular momentum of the state. In the case of a vortex with single vorticity \( (l = 1) \), the transformation \( \psi \) is well-defined only when \( \mu \) is an integer. We can express the equation
\[
H\psi = \varepsilon\psi \quad \text{as} \quad \varepsilon = 0 \quad \text{and} \quad \mu = 0. \]

If \( \varepsilon \) and \( \mu \) vanish, we introduce the following linear combinations
\[
X_1 = if_1^0 + f_4^0, \quad X_2 = if_1^0 - f_4^0,
Y_1 = if_2^0 + f_3^0, \quad Y_2 = if_2^0 - f_3^0,
\]
for which Eqs. \( \text{[27,28]} \) decouple into two independent systems of equations for \( (X_1, Y_2) \) and \( (X_2, Y_1) \). Each of these systems can be reduced to the following equation
\[
\chi'' + \frac{\chi'}{\bar{r}} + \chi \left( U^2 - \frac{1}{\bar{r}^2} \right) = 0,
\]
where the dimensionless variables
\[
\bar{r} = \frac{r}{l_B}, \quad \bar{U} = \frac{U}{\hbar\omega_c}
\]
are used, and the quantity
\[
\omega_c = \frac{v_F}{l_B}
\]
may be viewed as the cyclotron frequency for a massless relativistic particle. Function \( \chi \) is connected to \( Y_{1,2} \) as follows
\[
Y_1 = \chi(\bar{r}) \exp \left[ -\alpha \int_{\bar{r}_0}^{\bar{r}} f(\bar{r}')d\bar{r}' \right],
\]
\[
Y_2 = \chi(\bar{r}) \exp \left[ \alpha \int_{\bar{r}_0}^{\bar{r}} f(\bar{r}')d\bar{r}' \right],
\]
where \( \alpha = \frac{l_B}{\xi} \), and \( \xi = \frac{\hbar v_F}{|\Delta|} \).

Length scale \( \xi \) is the familiar coherence length due to the proximity effect. The substitution
\[
\chi(\bar{r}) = \frac{g(\bar{r}^2/2)}{\bar{r}}
\]
transforms Eq. (10) into a 1D Schrödinger equation describing a quantum particle in an attractive Coulomb potential. Using this analogy, it is easy to check that Eq. (10) has a normalizable solution, provided that

\[ \bar{U}^2 = 2N \geq 0, \quad N \text{ is any non-negative integer,} \quad (17) \]

When these conditions hold, we solve Eq. (10) and using Eqs. (7) (8), and (9) we obtain the solution for the MF localized near the vortex core in the form

\[ \psi_e = B_e \exp \left[ \frac{i \pi}{4} - \frac{r^2}{4l_B^2} \right] \int_0^r \frac{f(r')}{r'} \frac{dr'}{\xi} \Psi_{\text{MF}} , \quad (18) \]

\[ \Psi_{\text{MF}}(r, \phi) = \begin{cases} L_N^{(0)}(r^2/2) - i(\bar{\xi}/\bar{U}) L_N^{(1)}(r^2/2) e^{i\phi} \tau_z & \text{if } \bar{U} = N = 0, \quad L_N^{(1)}/\bar{U} = 0. \quad (19) \end{cases} \]

where \( L_m^{(n)}(y) \) are generalized Laguerre polynomials,\(^{38}\) and the normalization coefficient \( B_e \) is real. Note, if \( \bar{U} = N = 0 \), then \( L_N^{(1)}/\bar{U} \equiv 0. \) (20)

To prove that \( \psi_e \) is a MF, it is enough to check that \( \sum \psi_e = \psi_e. \)

The solution \( \psi_e \) of Eq. (18) corresponds to Eq. (13). The solution that corresponds to Eq. (14) describes yet another MF state

\[ \psi_e = B_e \exp \left[ \frac{i \pi}{4} - \frac{r^2}{4l_B^2} + \int_0^r f(r') \frac{dr'}{\xi} \tau_z \Psi_{\text{MF}} \right] . \quad (21) \]

In this expression \( \Psi_{\text{MF}} \) is defined by Eq. (19), and the coefficient \( B_e \) is real. If \( N \) is not too large, the maximum of the wave function weight \( |\psi_e|^2 \) is at

\[ r \sim r^* = \frac{2l_B^2}{\xi} = \frac{2\epsilon \Delta}{e \nu_F B}. \quad (22) \]

This means that Majorana fermion \( \psi_e \) localizes along a circle centered at the vortex core, with field-dependent radius \( r^*(B) \). We will refer to this state as the exterior MF (thus, the subscript ‘e’).

There is an important distinction between these two MFs. The state at the vortex core \( \psi_e \) is localized by ‘the vorticity’. That is, as long as the vortex is present, the wave function \( \psi_e \) remains normalizable even for zero magnetic field (\( l_B \to \infty \)), with its weight mostly confined within a circle of radius \( \xi \), the latter quantity being field-independent. Because of this, it is permissible to neglect \( A \) in the model Hamiltonian. This is a common approximation used in the literature dedicated to the core-bound Majorana fermions.\(^{20,22}\) However, including the magnetic field \( B \) into the model is of central importance to study the exterior MF \( \psi_e \), since the MF weight concentrates mostly at a field-dependent radius \( r^* \sim 1/B \). If \( B \) decreases, the radius \( r^* \) grows. Eventually, in some very weak field, the exterior state recedes to the outer boundaries of the system.

When the value of \( \bar{U} \) violates condition (17), strictly speaking, the Majorana states disappear. Let us consider a weak violation of (17): \( \bar{U} = \sqrt{2N} + \delta \bar{U} \), where \( \delta \bar{U} = \bar{U}/\nu_F \). Treating the term \( \delta \bar{U} \tau_z \) as a perturbation, one can evaluate the matrix element \( \delta \bar{U} \langle \psi_e | \tau_z | \psi_e \rangle \) and obtain the corresponding energy splitting

\[ \delta E \approx \delta \bar{U} \sqrt{\frac{l_B^2}{\xi}} \exp \left( -\frac{l_B^2}{\xi} \right) . \quad (23) \]

This value characterizes the hybridization between the core and exterior MF states for non-zero \( \delta \bar{U} \). Because of this hybridization, two MFs are replaced by a single Dirac fermion with energy \( \delta E \). The eigenenergy is an oscillating function of \( U \), vanishing each time when condition (17) is met. In the limit \( \xi \ll r^* (\xi \ll l_B) \), the overlap between the Majorana wave functions is exponentially small, and the hybridization may be neglected.

These overlap oscillations are not uncommon in the Majorana fermion physics: similar phenomena were discussed in other systems as well.\(^{39,40}\) It is interesting that single-parameters tuning is sufficient to nullify the hybridization between the MFs. This is unlike a common two-level system, which avoids level crossing unless multi-parameter fine-tuning is performed. Such a deviation from a generic behavior occurs because a single (real) parameter \( t_{12} \) is enough for the complete specification of the most general Hamiltonian

\[ H_{12} = it_{12} \gamma_1 \gamma_2 . \quad (24) \]

III. SPINLESS TWO-DIMENSIONAL \( p \)-WAVE SUPERCONDUCTOR

Two-dimensional \( p \)-wave superconductors\(^{41}\) \( \text{Sr}_2\text{RuO}_4 \) is another system where an exterior MF can emerge. The relevant description is very similar to the calculations presented above. We write down the model’s Hamiltonian in the form:\(^{1,39}\)

\[ H = \left( \frac{i \hbar \nabla - \epsilon A/\sqrt{2m}}{2m} - U - \frac{\hbar}{2m} \left\{ \Delta, \partial_z \right\} - \frac{\hbar}{2m} \left\{ \Delta^*, \partial_z \right\} \right) \left( \frac{i \hbar \nabla + \epsilon A/\sqrt{2m}}{2m} + U \right) . \quad (25) \]

Here \( m \) is the electron mass, \( p_F = mv_F \) is the Fermi momentum, the anticommutator is defined in the usual manner: \( \left\{ \Delta, \partial_z \right\} = \Delta \partial_z + \partial_z \Delta \), where

\[ \partial_z = e^{i\phi} (\partial_r + ir^{-1} \partial_\phi) . \quad (26) \]
The order parameter, as before, is given by Eq. (30). For the case of a single vortex \((l = 1)\) we seek zero-energy eigenfunctions in the form

\[
\psi_{\text{MF}} = \begin{pmatrix} e^{i(\phi - \pi/4)} u_{\text{MF}}, e^{-i(\phi - \pi/4)} v_{\text{MF}} \end{pmatrix}^T.
\] (27)

Similar to previous consideration, we derive a system of differential equations for the radial part of \(\psi_{\text{MF}}\).

Introducing the particle-hole conjugation operator for the \(p\)-wave superconductor as \(\Xi = \tau_y K\), one can prove that Hamiltonian in Eq. (28) possesses particle-hole symmetry: \(\Xi H \Xi = -H\). This property, together with the fact that the differential equations for the radial part of the wave function are real, implies that the MF radial wave function satisfies the conditions \(u_{\text{MF}} = \lambda v_{\text{MF}}\), where \(\lambda = \pm 1\). Using these relations we derive the differential equation for function \(\chi\)

\[
\chi'' + \frac{\chi'}{r} + \chi \left( \bar{U} - \alpha^2 f^2 - \frac{1}{r^2} \frac{r^2}{4} - \frac{1}{2} \right) = 0,
\] (28)

which is connected to the MF wave function as follows

\[
u_{\text{MF}} = \chi(\bar{r}) \exp \left[ \lambda \alpha \int_0^{\bar{r}} f(\bar{r}') d\bar{r}' \right],
\] (29)

where \(\bar{r} = \frac{r}{R_B}\), \(\bar{U} = \frac{2m \hbar^2 U}{\hbar^2} = \frac{2U}{\hbar \omega_c}\),

and the cyclotron frequency for a massive particle equals to \(\omega_c = eB/mc\).

Approximating \(f(r) \approx 1\) we notice that Eqs. (28) and (10) have the same structure. Exploiting this, one can prove that the Hamiltonian (28) admits two MF solutions

\[
\bar{U} - \alpha^2 - \frac{1}{2} = \frac{2U}{\hbar \omega_c} - \frac{r_B^2}{\xi^2} - \frac{1}{2} = 2N,
\] (31)

where \(N\) is non-negative integer. From Eq. (29) it is easy to see that the core-localized MF corresponds \(\lambda = -1\), while the exterior MF corresponds \(\lambda = 1\). Up to a normalization coefficient, the MF wave functions are

\[
\psi_{\nu, c} = \frac{r}{2r_B} \exp \left( \lambda \Delta \int_0^r f d\bar{r}' - \frac{r^2}{4R_B^2} \right) \left( \frac{e^{i(\phi - \pi/4)}}{\lambda \alpha e^{-i(\phi - \pi/4)}} \right).
\] (32)

One can verify that \(\Xi \psi_{\nu, c} = \lambda \psi_{\nu, c}\), and both MF wave functions are orthogonal to each other. Hence, \(\psi_{\nu, c}\) are two independent MFs. When the condition (31) is violated, the MFs are hybridized, forming a single Dirac electron as in the previously discussed case of the TI/SC heterostructure.

We assumed above that \(f(r) \approx 1\), while the function \(f\) deviates from unity near the core. This deviation may be accounted using perturbation theory. As a result, the condition (31) and the wave functions (32) will be slightly corrected.

IV. DISCUSSION

It is typically assumed that, in order to localize a MF, one needs a boundary separating two parts of the system with different topological numbers or some “topological” defect. Two models considered above serve as counterexamples to this statement: we have shown that the exterior Majorana wave function is not “latched” to any inhomogeneity. Instead, the localization radius \(r^*\) is a magnetic field-dependent quantity, and may be manipulated in real time.

The latter feature allows us to control the coupling between the MFs. Let two vortices are pinned at distance \(R\) from each other. Each vortex hosts a core-localized state and an exterior Majorana states. Straightforward calculations show that the splitting between the exterior MF of the first vortex and the core MF of the second vortex is zero. However, the splitting between the exterior MFs of different vortices is non-zero and depends on \(B\). When the magnetic field is high, such that \(R \gg r^* = 2R_B/\xi \sim 1/B\), the coupling between two exterior MFs is exponentially weak, and can be neglected. With the decrease of \(B\) the localization radius \(r^*\) grows, and so does the coupling between the MFs.

Changing the hybridization between the MFs allows one to perform braiding without moving the vortices. While the topological computations are not the main topic of this letter, we briefly outline a possible braiding protocol similar to one proposed in Ref. 3. Consider first the motion of a single MF. Initially \((t = 0)\) we have one separate MF \(\gamma_1\) and one Dirac fermion, the latter consisting of two coupled MFs \(\gamma_2\) and \(\gamma_3\). Transportation of the MF \(\gamma_1\) from its initial position to the position of \(\gamma_3\) can be described by the Hamiltonian

\[
H(t) = \zeta_{12} \alpha(t) \gamma_1 \gamma_2 + \zeta_{23} [1 - \alpha(t)] \gamma_2 \gamma_3,
\] (33)

where \(\zeta_{ij}\) are the tunneling amplitudes between the \(i\)-th and \(j\)-th MFs. The coefficient \(\alpha(t)\) changes adiabatically from \(\alpha(0) = 0\) to \(\alpha(t_1) = 1\). The equation of motion

\[
\gamma_i = i[H(t), \gamma_i(t)]
\] (34)

can be written as

\[
\gamma_a = 2 \epsilon_{abc} B^b \gamma_c,
\] (35)

where the vector \(B\) is equal to

\[
B(t) = [1 - \alpha(t)] \zeta_{23} (1, 0, 0) + \alpha(t) \zeta_{12} (0, 0, 1),
\] (36)

and \(\epsilon_{abc}\) is the antisymmetric Levi-Civita tensor. Straightforward calculations give the solution

\[
\gamma_3(t_1) = \text{sgn} (\zeta_{12} \zeta_{23}) \gamma_1 (0).
\] (37)

Thus, the MF can be transported by tuning the interaction between MFs. For exterior MFs, this interaction may be varied by adjusting local magnetic field.

The procedure described above can be used to implement braiding of two exterior MFs. As a result of
FIG. 2: (Color online) Braiding of two MFs. The consecutive steps are shown in panels (1-4). Each step or shift is a transportation of a single MF between two positions [see discussion for Eqs. 88 and 87]. This move (or shift) is represented by a dashed arrow. Orange circles (points) correspond to an exterior (core) MF. When two MFs couple and form a single Dirac fermion, the color is changed to blue. Initially [panel (1)] we have two exterior MFs (1 and 2) on vortices A and B. The MFs of vortex C are coupled into a single Dirac state. In panel (4) everything is the same except that MFs 1 and 2 switched their positions.

We analyzed the generation of MFs in topological superconductors in the presence of vortices. Discussing both (1) the topological insulator - superconductor heterostructure and (2) a spinless p-wave superconductor, we established that the combined effect of the vortex and the transverse magnetic field $B$ gives rise to the generation of two MFs. The first MF is well-known: it localizes near the vortex core. The second, exterior, MF is localized by the magnetic field. Its wave function weight is centered along a circle of radius $r^* \propto 1/B$. Varying the magnetic field we can change the positions of the MFs and tune the splitting between the core and exterior MFs of the same vortex or between exterior MFs of different vortices in real time. The tunability of the pairwise couplings between MFs by a magnetic field may open a new route for future topological quantum operations.

**Acknowledgements**

We acknowledge partial support from the Dynasty Foundation and ICFPM (MMK), the Ministry of Education and Science of the Russian Federation Grant No. 14Y26.31.0007, RFBR Grant No. 15-02-02128. F.N. was partially supported by: the RIKEN iTHESS Project, the MURI Center for Dynamic Magneto-Optics via the AFOSR Award No. FA9550-14-1-0040, the Japan Society for the Promotion of Science (KAKENHI), and a grant from the John Templeton Foundation.

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1. J. Alicea, “New directions in the pursuit of Majorana fermions in solid state systems”, Rep. Prog. Phys. 75, 076501 (2012).
2. M. Leijnse and K. Flensberg, “Introduction to topological superconductivity and Majorana fermions”, Semicond. Sci. Tech. 27, 124003 (2012).
3. C. Beenakker, “Search for Majorana Fermions in Superconductors”, Annu. Rev. Condens. Matter Phys. 4, 113-136 (2013).
4. S.R. Elliot, M. Franz, “Colloquium: Majorana fermions in nuclear, particle, and solid-state physics”, Rev. Mod. Phys. 87, 137 (2015).
5. M. Sato, S. Fujimoto, “Majorana Fermions and Topology in Superconductors”, preprint [arXiv:1601.02726] (unpublished).
6. A. Yu. Kitaev, “Fault-tolerant quantum computation by anyons”, Ann. Phys. (NY) 303, 2 (2003).
7. C. Nayak, S.H. Simon, A. Stern, M. Freedman, S. Das
M. Burrello, B. van Heck, A.R. Akhmerov, “Braiding of non-Abelian anyons using pairwise interactions”, Phys. Rev. A 87, 022343 (2013).

J.D. Sau, D.J. Clarke, and S. Tewari, “Controlling non-Abelian statistics of Majorana fermions in semiconductor nanowires”, Phys. Rev. B 84, 094505 (2011).

B. van Heck, A.R. Akhmerov, F. Hassler, M. Burrello, C.W.J. Beenakker, “Coulomb-assisted braiding of Majorana fermions in a Josephson junction array”, New J. Phys. 14, 035019 (2012).

S. Bravyi, “Universal quantum computation with the $\nu = 5/2$ fractional quantum Hall state”, Phys. Rev. A 73, 042313 (2006).

P. Bonderson, D.J. Clarke, C. Nayak, and K. Shtengel, “Implementing Arbitrary Phase Gates with Ising Anyons”, Phys. Rev. Lett. 104, 180505 (2010).

K. Flensberg, “Non-Abelian Operations on Majorana Fermions Via Single-Charge Control”, Phys. Rev. Lett. 106, 090503 (2011).

T.L. Schmidt, A. Nunnenkamp, and C. Bruder, “Majorana Qubit Rotations in Microwave Cavities”, Phys. Rev. Lett. 110, 107006 (2013).

A.A. Kovalev, A. De, and K. Shtengel, “Spin Transfer of Quantum Information Between Majorana Modes and a Resonator”, Phys. Rev. Lett. 112, 106402 (2014).

P. Zhang, F. Nori, “Coherent manipulation of a Majorana qubit by a mechanical resonator”, Phys. Rev. B 92, 115303 (2015).

K. Sato, Y. Takahashi, S. Fujimoto, “Non-Abelian topological order in s-Wave Superfluids of Ultracold Fermionic Atoms”, Phys. Rev. Lett. 103, 020401 (2009).

J.D. Sau, R.M. Lutchyn, S. Tewari, S. Das Sarma, “Generic New Platform for Topological Quantum Computation Using Semiconductor Heterostructures”, Phys. Rev. Lett. 104, 040502 (2010).

J.D. Sau, R.M. Lutchyn, S. Tewari, and S. Das Sarma, “Robustness of Majorana fermions in proximity-induced superconductors”, Phys. Rev. B 82, 094522 (2010).

M. Sato, Y. Takahashi, S. Fujimoto, “Non-Abelian topological orders and Majorana fermions in spin-singlet superconductors”, Phys. Rev. B 82, 134521 (2010).

A.L. Rakhamanov, A.V. Rozhkov, F. Nori, “Majorana fermions in pinned vortices”, Phys. Rev. B 84, 075141 (2011).

P.A. Ioselevich, M.V. Feigelman, “Anomalous Josephson Current via Majorana Bound States in Topological Insulators”, Phys. Rev. Lett. 106, 077003 (2011).

T.D. Stanescu, R.M. Lutchyn, and S. Das Sarma, “Majorana fermions in semiconductor nanowires”, Phys. Rev. B 84, 144522 (2011).

R.P. Tiwari, U. Zülicke, and C. Bruder, “Majorana Fermions from Landau Quantization in a Superconductor and Topological-Insulator Hybrid Structure”, Phys. Rev. Lett. 110, 186805 (2013).

J.Q. You, Z.D. Wang, W. Zhang, F. Nori, “Encoding a qubit with Majorana modes in superconducting circuits”, Sci. Rep. 4, 5535 (2014).

R.S. Akzyanov, A.V. Rozhkov, A.L. Rakhamanov, F. Nori, “Tunneling spectrum of a pinned vortex with a robust Majorana state”, Phys. Rev. B 89, 085409 (2014).

R.S. Akzyanov, A.L. Rakhamanov, A.V. Rozhkov, F. Nori, “Majorana fermions at the edge of superconducting islands”, Phys. Rev. B 92, 075432 (2015).

V. Mourik, K. Zuo, S.M. Frolov, S.R. Plissard, E.P.A.M. Bakkers, and L.P. Kouwenhoven, “Spectra of Majorana fermions in hybrid superconductor-semiconductor nanowire devices”, Science 336, 1003 (2012); M.T. Deng, C.L. Yu, G.Y. Huang, M. Larsson, P. Caroff, and H.Q. Xu, “Anomalous Zero-Bias Conductance Peak in a Nanoribbon-InSb Nanowire—Nanoribbon Hybrid Device”, Nano Lett. 12, 6414 (2012); L.P. Rokhinson, X. Liu, and J.K. Furdyna, “The fractional a.c. Josephson effect in a semiconductor-superconductor nanowire as a signature of Majorana particles”, Nat. Phys. 8, 795 (2012); A. Das, Y. Ronen, Y. Most, Y. Oreg, M. Heiblum, and H. Shtrikman, “Zero-bias peaks and splitting in an AlInAs nanowire topological superconductor as a signature of Majorana fermions”, Nat. Phys. 8, 887 (2012).

P. Zhang, F. Nori, “Coherent manipulation of a Majorana qubit by a mechanical resonator”, Phys. Rev. B 92, 115303 (2015).

L. Fu and C.L. Kane, “Superconducting Proximity Effect and Majorana Fermions at the Surface of a Topological Insulator”, Phys. Rev. Lett. 100, 096407 (2008).

M. Sato, Y. Takahashi, S. Fujimoto, “Non-Abelian topological order in s-Wave Superfluids of Ultracold Fermionic Atoms”, Phys. Rev. Lett. 103, 020401 (2009).

J.D. Sau, R.M. Lutchyn, S. Tewari, S. Das Sarma, “Generic New Platform for Topological Quantum Computation Using Semiconductor Heterostructures”, Phys. Rev. Lett. 104, 040502 (2010).

J.D. Sau, R.M. Lutchyn, S. Tewari, and S. Das Sarma, “Robustness of Majorana fermions in proximity-induced superconductors”, Phys. Rev. B 82, 094522 (2010).
Confined Superconducting Condensates", Phys. Rev. Lett. 107, 097202 (2011).