Topological solid states of matter [1, 2] have the property that their bulk states are gapped, while states at the boundaries are gapless and protected by some discrete quantum symmetry. The topological aspect emerges when considering the transport properties of the boundary states, where the transport current happens to also be a topological current. The most well established topological solid states of matter are topological insulators (TIs), which are gapped in the bulk and have helical (or chiral) gapless states at the boundaries which are protected by time-reversal symmetry. Helical here means that the electronic spin is locked to momentum due to strong spin-orbit coupling. Thus, the boundary states have an helicity determined by the eigenvalues of $\sigma \cdot k / k$ at each boundary. Topological insulators have been predicted to exist [3] and confirmed experimentally in subsequent papers [4]. Although many of the materials investigated experimentally are not perfect insulators in the bulk, the observed boundary helical states are robust features of these materials.

The experimental situation is less clear in another predicted topological solid state of matter, namely topological superconductors (TSCs) [1, 2]. TSCs follow a symmetry classification scheme closely related to TIs, as far as Hamiltonians of the Bogoliubov-de Gennes type are concerned [5]. Just like TIs, TSCs have gapped states in the bulk and symmetry-protected gapless states at the boundaries. Unlike TIs, in TSCs the $U(1)$ symmetry is broken, either spontaneously or by proximity effect. The gapless boundary states are Majorana fermions, which are fermionic particles that are their own antiparticles. In order to support such states at the boundaries, the topological superconductivity must feature a $p$-wave type of gap. Parity is the underlying symmetry protecting the boundary Majorana states. In one dimension a paradigmatic simple model for topological superconductivity has been proposed by Kitaev [7] where the Majorana zero-energy states live at the ends of a quantum wire. An experimental way of realizing a superconducting state in a quantum wire is by proximity effect. In this case a semiconducting wire with strong spin-orbit coupling is deposited on the surface of an $s$-wave superconductor in the presence of an external perpendicular magnetic field. Then by proximity effect $p$-wave like superconductivity is induced on the wire for a certain range of parameters. [8, 9]. There are some experimental signatures of Majorana modes in Indium antimonide nanowires in contact with normal and superconducting electrodes [10]. More recently, strong support for Majorana boundary zero modes have been reported in an experiment with a ferromagnetic chain of Iron fabricated on the surface of Lead [11].

Three-dimensional topological superconductors have also been discussed theoretically, in particular focusing on vortex physics [12] and possible topological phase transitions [12]. A distinctive feature of both three-dimensional TIs and TSCs with respect to their non-topological counterparts is the topological magnetoelectric response induced by a so-called chiral anomaly [13]. When fermions in topological materials interact with the electromagnetic field, a Berry phase mixing electric and magnetic fields is induced [14]. In TIs this occurs due to strong spin-orbit coupling that locks spin to momentum. The resulting Berry phases combine in the form of a so-called axion term, which is a magnetoelectric term $\sim E \cdot B$ with a periodic field, $\theta$, appearing as a coefficient [14]. This coefficient corresponds to a topological invariant implied by the band structure. In the case of TSCs a topological magnetoelectric contribution also arises, but now $\theta$ corresponds to the phase difference between order parameters of opposite chirality [15]. Recently, axionic superconductivity has been also discussed in the context of doped narrow-gap semiconductors [16].

In this paper we investigate the Higgs mechanism and anomalous Hall effect of three-dimensional TSCs following the model introduced recently in Ref. [15]. In the simplest case, the model features two superconducting order parameters associated with left and right fermion chiralities interacting with the electromagnetic field, and a topological magnetoelectric term in the form we described in the previous paragraph. We
will show that the Higgs mechanism in such a TSC, quantum fluctuations imply a topologically non-trivial phase that cannot be continuously deformed into the topologically trivial one. This result is not directly obvious from the classical Lagrangian of the system derived earlier in Ref. [15]. Another consequence of the topological magnetoelectric term is the occurrence of an anomalous Hall effect when an electric field is applied parallel to the surface of a TSC. Due to the magnetoelectric effect, a transverse current is generated due to a Lorentz-like force involving the relative superfluid velocity \( \sim \nabla \theta \) and the applied electric field. In this Hall effect the generated transverse current is negative, a situation reminiscent from the anomalous Hall effect in superconductors predicted long time ago by Josephson [17], and observed later in a high-\( T_c \) cuprate superconductor [18]. However, in the latter case the anomalous Hall effect occurs due to vortex motion induced by the Faraday law, and is typically a very small effect. Furthermore, the Lorentz force acts in this case directly on the vortex core, and therefore on the normal components of the superconductor. For this reason, it automatically leads to dissipation. In the case of three-dimensional TSCs, on the other hand, the anomalous Hall effect occurs even in absence of vortices and is induced solely by an external electric field via the magnetoelectric effect, generating in this way a dissipationless anomalous Hall current on the surface.

The effective Lagrangian for a three-dimensional TSC featuring two Fermi surfaces is given by [15].

\[
\mathcal{L}_{\text{eff}} = \frac{\theta e^2}{4\pi^2} e^\mu\nu\sigma\tau \mathcal{F}_{\mu\sigma} \mathcal{F}_{\tau\nu} - \frac{1}{4} F_{\mu\nu} F^{\mu\nu} + \sum_{\nu=L,R} \left[ \frac{i}{2} (\partial_{\mu} - q A_{\mu}) \phi_{\nu} \mathcal{F}_{\mu\nu} \right] - \frac{u}{2} (|\phi_L|^2 + |\phi_R|^2)^2 + 2 J (\phi^*_L \phi_R + \phi^*_R \phi_L),
\]

where \( q = 2e \). The Greek indices run from 0 to 3 and \( F_{\mu\nu} = \partial_{\mu} A_{\nu} - \partial_{\nu} A_{\mu} \), with \( (A_\mu) = (A_\theta, \mathbf{A}) \). The Lagrangian [1] corresponds to an abelian Higgs model with a two-component scalar field. In contrast to the standard Higgs model, Eq. (1) features a so-called axion term [19], which is the first term appearing in the Lagrangian above. The term is topological in nature and contains a scalar field (the axion) \( \theta = \theta_L - \theta_R \), where \( \theta_L \) and \( \theta_R \) are the phases of \( \phi_L \) and \( \phi_R \), respectively. In terms of electric and magnetic field components, it becomes \( e^2 \theta (\mathbf{E} \cdot \mathbf{B})/(8\pi^2) \), having precisely the magnetoelectric form mentioned in the introductory paragraphs. A Josephson coupling term \( \propto \phi^*_L \phi_R + \phi^*_R \phi_L \) accounts for the interference between the two superconducting order parameter fields. This is a characteristic feature in superconductors with two or more components of the order parameters, and is absent only if prohibited by symmetries of the problem [20]. For \( J < 0 \) the Josephson coupling implies \( \theta = 0 \) in the mean-field ground state, yielding a topologically trivial superconductor. For \( J > 0 \), on the other hand, \( \theta \) is locked to \( \pi \), thus leading to a topologically non-trivial superconducting ground state. Furthermore, since \( \theta \) is periodic, \( \theta = \pi \) corresponds to a situation where the time-reversal symmetry is preserved [13]. Thus, at the mean-field level, \( J = 0 \) separates a topologically trivial ground state from a non-trivial one. Thus, varying \( J \) from positive to negative values induces a topological quantum phase transition.

In the \( U(1) \) Higgs mechanism the phases disappear from the spectrum due to spontaneous breaking of the local \( U(1) \) symmetry, being transmuted into the longitudinal mode for the photon, which becomes gapped. Thus, only amplitude modes remain in the spectrum of scalar particles. The Higgs mechanism is equivalent to integrating out the phase degrees of freedom, which in the case of the Higgs model automatically leads to a massive gauge particle. This point of view of integrating out the phases is particularly appealing in the case where a Josephson coupling is present. However, additional non-linearities arise in the presence of the axion term. To see this, let us first consider the Higgs mechanism in Eq. (1) for the case where the axion term is absent. In this case we can simply write \( A_\mu \rightarrow A_\mu + 1/q \left( \rho^2 \partial_\mu \partial_\nu + \rho^2 \partial_\nu \partial_\mu \right) \),

\[
A_\mu \rightarrow A_\mu + \frac{1}{q} \left( \frac{\rho^2 \partial_\mu \partial_\nu + \rho^2 \partial_\nu \partial_\mu}{\rho^2 + \rho^2} \right),
\]

which immediately decouples the phases from the gauge potential. Thus, the phases formally can be integrated out such that the resulting effective action contains only a massive photon coupled to amplitude degrees of freedom. However, in the presence of the axion term, the shift (2) will not immediately decouple the phases from the gauge field. Thus, integrating out the phases generate direct interactions between photons, even in a mean-field regime where the amplitudes are assumed to be uniform. The integration of the phases can be performed exactly at the mean-field critical point \( J = 0 \) for the topological phase transition, yielding,

\[
\mathcal{L}_{\text{Higgs}} |_{J=0} = -\frac{1}{4} F^2 + \frac{q^2 (\rho^2_L + \rho^2_R)}{2} A^2 + \left( \frac{e^2}{16\pi^2} \right) \left( \rho^2_L + \rho^2_R \right) \left( A^2 F^2 + 2A_\mu A_\nu F_{\mu\nu} F^{\mu\nu} \right) + \frac{1}{2} \left( |\partial_\mu \rho L|^2 + |\partial_\mu \rho R|^2 \right) - \frac{m^2}{2} (\rho^2_L + \rho^2_R) - \frac{u}{8} \left( \rho^2_L + \rho^2_R \right)^2.
\]

The Higgs mechanism will occur for \( m^2 < 0 \), leading to the condensation of the amplitudes and generating as usual a mass for the photon. In addition, Eq. (3) has a direct interaction between the photons. Therefore, the Higgs mechanism yields in this case an interacting photon theory at the mean-field level (or tree level, in quantum field theory language), in contrast to the usual Abelian Higgs mechanism, where the photon action is Gaussian.

For the topological phase \( J > 0 \) we have to integrate at lowest order the Gaussian phase fluctuations around \( \theta = \pi \). This renders the induced photon-photon interaction non-local, with the topological term appearing as a total derivative. Thus, for
uniform amplitudes we now obtain the effective Lagrangian,

\[ \mathcal{L}_{\text{Higgs}}|_{J>0} = -\frac{1}{4} F^2 + \frac{q^2 (\rho_\perp^2 + \rho_\parallel^2)}{2} A^2 + \frac{e^2}{64\pi^2} e^{i\theta(p\cdot F)} F_{\mu\nu} F_{\sigma\tau} + \frac{1}{2} \left( \frac{e^2}{16\pi^2} \right)^2 \left( \frac{\rho_\perp^2 + \rho_\parallel^2}{\rho_\perp^2 + \rho_\parallel^2} \right) \times \int d^4x V_{\mu\nu}(x - x') e^{i\epsilon_{\mu\nu\rho}\epsilon^{\pm\rho\lambda\eta} A^\lambda(x) A_\lambda(x')} (x) A_\delta(x') F^{\lambda\delta}(x) F_{\omega\eta}(x'), \]  

where

\[ V_{\mu\nu}(x) = \int \frac{d^3p}{(2\pi)^3} \frac{p_\mu p_\nu}{p^2 + m_p^2} e^{ip\cdot x}, \]  

with \( m_p^2 = 2J(\rho_\perp/\rho_\parallel + \rho_\parallel/\rho_\perp) \). Appearances notwithstanding, there exists a regime where the Lagrangian (4) can be viewed as simpler than (3). In fact, since the phase fluctuations in the regime \( J > 0 \) are gapped, the non-local photon-photon interaction can be neglected in the long wavelength limit. Thus, neglecting the non-local interaction and minimizing the resulting effective Lagrangian with respect to both \( \rho_\perp \) and \( \rho_\parallel \) yields the result \( \rho_\perp^2 = \rho_\parallel^2 = -(J + m^2) \), implying the condition \( m^2 < -J < 0 \).

Now we turn to the discussion of the following crucial aspect of the topological phase with respect to the surface states. It turns out that due to the axion term the topological surface states cannot be continuously deformed into topologically trivial ones in the long wavelength limit when crossing the critical point. To see this we note that since \( e^{i\theta(p\cdot F)} F_{\mu\nu} F_{\sigma\tau} = 2\partial^\lambda(e_{\mu\nu\rho\lambda} A^\rho F^{\lambda\sigma}) \), each surface contains a Chern-Simons term \[ \int \frac{d^3q}{(2\pi)^3} \delta_{\mu\nu} (p + q) \Delta_{\nu\rho}^\pm(q) = 0, \]  

while the result is divergent for \( \theta = 0 \). The above result would also be valid for any other fixed \( \theta \neq 0, \pi \). Thus, after other one-loop scattering amplitudes are taken into account to obtain the full four-Higgs vertex, we see that at the critical point the topological field theory cannot be continuously deformed into a topologically trivial one. Such a continuous deformation can be done in the Higgs phase, where there are no gapless modes. However, at the critical point such a continuous deformation is not possible. Thus, quantum critical fluctuations in this system will govern topologically stable universal behavior in physical quantities, such as for instance critical exponents and amplitude ratios.

Now we turn to the Meissner effect aspects of a TSC, which as we will now show, implies an anomalous Hall effect even in the absence of vortices. This is more conveniently done by rewriting the Lagrangian in a London limit exhibiting explicitly electric and magnetic fields, i.e.,

\[ \mathcal{L}_{\text{eff}} = \frac{1}{8\pi} \left( \epsilon e^2 - \frac{1}{\mu} B^2 \right) \epsilon \cdot B + \frac{e^2}{8\pi^2} E \cdot B \]

\[ + \frac{1}{2} \sum_{l = L, R} \rho_\perp^2 (\partial_\mu \theta_l - q A_\mu)^2 + J \rho_\perp \rho_\parallel \cos \theta \]

\[ - \frac{m^2}{2} (\rho_\perp^2 + \rho_\parallel^2) - \frac{u}{8} (\rho_\perp^2 + \rho_\parallel^2)^2. \]

From the effective Lagrangian we obtain that the electric displacement and magnetic fields are given respectively by

\[ \mathbf{D} = \epsilon \mathbf{E} + e^2 \mathbf{B}/(2\pi) \]  

and \[ \mathbf{H} = \mu^{-1} \mathbf{B} - e^2 \mathbf{E}/(2\pi) \], while the superconducting current is given by

\[ j = q(\rho_\perp^2 (\nabla \theta_L - q A) + \rho_\parallel^2 (\nabla \theta_R - q A)). \]

From Eq. (9) we obtain the usual London equation in absence of vortices, \( \nabla \times j = -(1/\lambda^2) \mathbf{B} \), where \( \lambda^2 = 1/M^2 \) is the square of the penetration depth. Thus, the Maxwell equation in the presence of the axion field,

\[ \frac{1}{\mu} (\nabla \times \mathbf{B}) = 4\pi j + e\partial_\mu \mathbf{E} + e^2/\pi (\nabla \theta \times \mathbf{E} + \partial_\mu \theta \mathbf{B}), \]  

FIG. 1. Difference in behavior for surface photon-mediated Higgs scattering at the critical point in the long wavelength limit. For the topologically trivial case (\( \theta = 0 \)) the corresponding Feynman diagram diverges. On the other hand, for the topologically non-trivial case (\( \theta = \pi \)), the same diagram vanishes. The same result would hold for any other non-vanishing value of \( \theta \). This shows that a topologically non-trivial one in the long wavelength limit.

\[ \phi_\mu = \cdots \]

\[ \theta = 0 : \quad \lim_{|p| \to 0} \theta = \infty \]

\[ \theta = \pi : \quad \lim_{|p| \to 0} \theta = 0 \]
yields the equation determining the London electrodynamics of the TSC in the form,
\[ \mu \varepsilon \mu \partial_t^2 \mathbf{B} - \nabla^2 \mathbf{B} + m_A^2 \mathbf{B} = \]
\[ + \frac{e^2 \mu}{\pi} \left[ \mathbf{E} \times (\nabla \times \mathbf{E}) + \nabla \times (\partial_t \theta \mathbf{B}) \right], \]
where we have defined \( m_A^2 \equiv 4\pi \mu / \lambda^2 = 4\pi \mu M^2 \). For the axion term, retaining its traditional form,
\[ - \nabla^2 \theta + m_\theta^2 \sin \theta = \frac{e^2}{8\pi^2} \left( \frac{1}{\rho_L} + \frac{1}{\rho_R} \right) \mathbf{E} \cdot \mathbf{B}. \]
In the low frequency regime and in absence of vortices, the London equation (11) simplifies to,
\[ - \nabla^2 \mathbf{B} + m_A^2 \mathbf{B} = \frac{e^2 \mu}{\pi} \left[ \mathbf{E} \times (\nabla \times \mathbf{E}) \right], \]
while the current satisfies,
\[ - \nabla^2 \mathbf{j} + m_A^2 \mathbf{j} = - \frac{e^2 m_A^2}{4\pi^2} (\nabla \theta \times \mathbf{E}). \]
The London equation for the electric field is unaffected by the axion term, retaining its traditional form, \( - \nabla^2 \mathbf{E} + m_E^2 \mathbf{E} = 0 \). This result is closely related to the fact that the electromagnetic energy density does not contain a magnetoelectric term.

We now consider a solution with a simple geometry, namely, a semi-infinite TSC \((z \geq 0)\) with a surface at \( z = 0 \) at an external electric field \( \mathbf{E}_0 = E_0 \hat{x} \) parallel to the surface. We obtain,
\[ - \frac{d^2 \theta}{dz^2} + m_\theta^2 \sin \theta = \frac{e^2}{8\pi^2} \left( \frac{1}{\rho_L} + \frac{1}{\rho_R} \right) E_A(z) B_A(z), \]
where \( E_A(z) = E_0 e^{-m_A z} \), and
\[ - \frac{d^2 B_A}{dz^2} + m_A^2 B_A = - \frac{e^2}{\pi} \frac{d}{dz} \left[ E_A(z) \frac{d \theta}{dz} \right], \]
\[ - \frac{d^2 j_y}{dz^2} + m_A^2 j_y = - \frac{e^2 m_A^2}{4\pi^2} E_A(z) \frac{d \theta}{dz}. \]
The solution for Eq. (17) in terms of the axion is,
\[ j_y(z) = \frac{e^2 m_A^2 E_0}{4\pi^2} \left[ \frac{\theta(0)}{2m_A} e^{-m_A z} - e^{-m_A z} \int_{z}^{\infty} e^{-2m_A z'} \theta(z') \right], \]
where \( \theta(0) = \pi \). Since \( \mathbf{E} \cdot (\nabla \times \mathbf{j}) = - \lambda^2 \mathbf{E} \cdot \mathbf{B} \), we obtain the following relation,
\[ \frac{1}{\lambda^2} B_A(z) = m_A j_y(z) + \frac{e^2 m_A^2 E_0}{4\pi^2} e^{-m_A z} [\theta(z) - \theta(0)], \]
which implies \( m_A B_A(0) = 4\pi \mu j_y(0) \). Thus, we obtain that the usual boundary condition of the London theory, \( d j_y/dz_{z=0} = m_A \dot{j}_y(0) \), is obviously fulfilled by the solution (18) in the presence of the axion field. However, the Maxwell equation (10) in the static regime implies a boundary condition that deviates from the usual one in the London theory of non-topological superconductors,
\[ \frac{dB_A}{dz} \bigg|_{z=0} = \frac{e^2 \mu E_0}{\pi} \frac{d \theta}{dz} \bigg|_{z=0}. \]
From Eqs. (18) and (19) we see that an approximate solution can be obtained by considering terms proportional to \( e^4 \) as being of higher order, which amounts to approximating Eq. (15) as being homogeneous. In this case we can use the domain wall solution \( \theta(z) = \pi + 2 \arcsin \left[ \tanh(m_\theta z) \right] \) in Eq. (18), which yields \( j_y(z) \) explicitly. The explicit solution for \( j_y(z) \) with \( m_\theta \neq m_A \) in terms of hypergeometric and Lerch transcendent, is not very illuminating. Instead, we plot it in Fig. 2 for four different values of the ratio \( m_\theta/m_A \). It has a negative sign, just like in the case of the anomalous Hall effect in high-\( T_c \) superconductors [18]. As emphasized in the introductory paragraphs, the anomalous Hall effect in non-topological superconductors has a quite different origin from the one discussed here. In three-dimensional TSCs the anomalous Hall current arises independently of vortex motion and is associated with a dissipationless current.

For \( m_\theta = m_A \) (blue curve in Fig. 2), the expression for \( j_y(z) \) does not involve special functions, reading,
\[ j_y(z) = \frac{e^2 m_A E_0}{4\pi^2} \left[ (\pi/2) e^{-m_A z} - 2 \right. \]
\[ - 2 \left. [e^{-m_A z} \arctan(e^{m_A z}) - e^{m_A z} \arctan(e^{-m_A z})] \right]. \]
In conclusion, we have shown that due to quantum electromagnetic fluctuations, the Higgs mechanism in three-dimensional TSCs implies a robust topological state of matter, since it cannot be continuously deformed into a topologically non-trivial one. This is an example of a topological state that is protected due to the coupling of phase and electromagnetic fluctuations via the axion term. In the low frequency limit this implies a London regime that leads to the generation of an
anomalous Hall current having a negative sign. This anomalous Hall current is dissipationless and is the consequence of a Lorentz-like force involving the relative superfluid velocity which is simply given by the gradient of the phase difference between the chiral superconducting components.

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[1] M. Z. Hasan and C. L. Kane, Rev. Mod. Phys. 82, 3045 (2010).
[2] X.-L. Qi and S.-C. Zhang, Rev. Mod. Phys. 83, 1057 (2011).
[3] C. L. Kane and E. J. Mele, Phys. Rev. Lett. 95, 226801 (2005); B. A. Bernevig, T. L. Hughes, S.-C. Zhang, Science 314, 1757 (2006); L. Fu, C. L. Kane, Phys. Rev. B 76, 045302 (2007); H. Zhang, C.-X. Liu, X.-L. Qi, X. Dai, Z. Fang, S.-C. Zhang, Nat. Phys. 5, 438 (2009).
[4] M. König, S. Wiedmann, C. Brüne, A. Roth, H. Buhmann, L. W. Molenkamp, X.-L. Qi, S.-C. Zhang, Science 318, 766 (2007); D. Hsieh, D. Qian, L. Wray, Y. Xia, Y. S. Hor, R. J. Cava, and M. Z. Hasan, Nature 452, 970 (2008); Y. Xia, D. Qian, D. Hsieh, L. Wray, A. Pal, H. Lin, A. Bansil, D. Grauer, Y. S. Hor, R. J. Cava, M. Z. Hasan, Nat. Phys. 5, 398 (2009); Y. L. Chen, J. G. Analytis, J.-H. Chu, Z. K. Liu, S.-K. Mo, X. L. Qi, H. J. Zhang, D. H. Lu, X. Dai, Z. Fang, S. C. Zhang, I. R. Fisher, Z. Hussain, Z.-X. Shen, Science 10, 178 (2009).
[5] Y. Xu, I. Miotkowski, C. Liu, J. Tian, H. Nam, N. Alidoust, J. Hu, C.-K. Shih, M. Z. Hasan, and Y. P. Chen, Nat. Phys. 10, 956 (2014).
[6] A. P. Schnyder, S. Ryu, A. Furusaki, and A. W. W. Ludwig Phys. Rev. B 78, 195125 (2008).
[7] A. Yu. Kitaev, Phys. Usp. 44, 131 (2001).
[8] J. Alicea, Y. Oreg, G. Refael, F. von Oppen, and M. A. Fisher, Nat. Phys. 7, 412 (2011).
[9] For a review on the theoretical and experimental situation, see T. D. Stanescu and S. Tewari, J. Phys.: Condens. Matter 25, 233201 (2013).
[10] V. Mourik, K. Zuo, S. M. Frolov, S. R. Plissard, E. P. A. M. Bakkers, and L. P. Kouwenhoven, Science 336, 1003 (2012).
[11] S. Nadj-Perge, I. K. Drozdov, J. Li, H. Chen, S. Jeon, J. Seo, A. H. MacDonald, B. A. Bernevig, and A. Yazdani, Science 346, 602 (2014).
[12] P. Hosur, P. Ghaemi, R. S. K. Mong, and A. Vishwanath, Phys. Rev. Lett. 107, 097001 (2011); H.-H. Hung, P. Ghaemi, T. L. Hughes, and M. J. Gilbert, Phys. Rev. B 87, 035401 (2013).
[13] An excellent textbook discussion of the chiral anomaly can be found in M. Srednicki, Quantum Field Theory (Cambridge University Press, Cambridge, 2007).
[14] X.-L. Qi, T. L. Hughes, and S.-C. Zhang, Phys. Rev. B 78, 195424 (2008).
[15] X.-L. Qi, E. Witten, and S.-C. Zhang, Phys. Rev. B 87, 134519 (2013).
[16] P. Goswami and B. Roy, Phys. Rev. B 90, 041301(R) (2014).
[17] B. D. Josephson, Phys. Lett. 16, 242 (1965).
[18] S. J. Hagen, C. J. Lobb, and R. L. Greene, M. G. Forrester, and J. H. Kang, Phys. Rev. B 41, 11630 (1990).
[19] F. Wilczek, Phys. Rev. Lett. 58, 1799 (1987).
[20] E. Babaev and M. Speight, Phys. Rev. B 72, 180502(R) (2005); J. Smiseth, E. Smørgrav, E. Babaev, and A. Sudbø, Phys. Rev. B 71, 214509 (2005).
[21] S. Deser, R. Jackiw, and S. Templeton, Phys. Rev. Lett. 48, 975 (1982); A. J. Niemi and G. W. Semenoff, Phys. Rev. Lett. 51, 2077 (1983); A. N. Redlich, Phys. Rev. D 29, 2366 (1984).