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Wave functions of bosonic symmetry protected topological phases

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We study the structure of the ground-state wave functions of bosonic symmetry protected topological (SPT) insulators in three space dimensions. We demonstrate that the differences with conventional insulators are captured simply in a dual vortex description. As an example, we show that a previously studied bosonic topological insulator with both global U(1) and time-reversal symmetry can be described by a rather simple wave function written in terms of dual “vortex ribbons.” The wave function is a superposition of all the vortex-ribbon configurations of the boson, and a factor \((-1)\) is associated with each self-linking of the vortex ribbons. This wave function can be conveniently derived using an effective field theory of the SPT phase in the strong-coupling limit, and it naturally explains all the phenomena of this SPT phase discussed previously. The ground-state structure for other three-dimensional (3D) bosonic SPT phases are also discussed similarly in terms of vortex loop gas wave functions. We show that our methods reproduce known results on the ground-state structure of some 2D SPT phases.

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I. INTRODUCTION

The disordered ground states of strongly interacting quantum many-body systems can have much richer structures compared with classical disordered states. The quantum richness of a system is encoded in the entanglement of its ground-state wave function, and without assuming any symmetry of the Hamiltonian, the ground-state wave function of a quantum many-body state can have long-range entanglement, which implies that the system has a “topological order.” In the last few years, motivated by the discovery of free fermion topological insulators protected by time-reversal symmetry,\(^1\)\(\text{--}^6\) it was realized that a short-range entangled state can still be fundamentally distinct from trivial product states, as long as the system preserves certain global symmetry \(G\). These nontrivial quantum disordered phases with short-range entanglement are called “symmetry protected topological” (SPT) phases. They are separated from the trivial product state through sharp quantum phase transitions in the bulk, which are either continuous or first order. In space dimension \(d = 1\), the Haldane spin chain provides an old and nice example of an SPT phase.\(^7\)\(\text{--}^8\) It has a bulk gap and no fractional excitations, but nevertheless has dangling symmetry protected spin-1/2 moments at the edge.\(^9\)\(\text{--}^{12}\) The Haldane chain thus provides an early example of an interacting topological insulator.

In this paper, we are mainly concerned with three-dimensional symmetry protected topological insulators of bosons and spin systems. A formal mathematical classification\(^13\) of SPT phases based on group cohomology allows a number of such phases to exist (depending on the global symmetry), but sheds little light on the physical properties. The latter have been discussed recently in Ref.\(^14\). A characteristic feature of all SPT phases is the presence of nontrivial surface states at the interface with a trivial insulator. Indeed, though the bulk is gapped and has no fractional excitations or topological order, such an interface cannot be in a trivial insulating state. Reference\(^14\) described the effective surface theory for a number of three-dimensional bosonic topological insulators and determined the structure of the allowed nontrivial phases. These either spontaneously break the defining global symmetry or, if gapped, have surface topological order. Exotic symmetry preserving gapless states were also shown to be possible. A key feature is that the surface effective field theory realizes symmetry in a manner not possible in strictly two-dimensional systems. Bulk topological field theories and effective field theory descriptions have also been provided.

In this work, we will study the structure of the ground-state wave function of various such three-dimensional (3D) bosonic SPT insulating phases with global U(1) and time-reversal \((\mathbb{Z}_T^2)\) symmetries. The differences with conventional Mott insulators are conveniently captured in a dual description in terms of closed vortex loops. In Mott insulating phases (conventional or topological), the vortex loops have proliferated and the ground-state wave function can be described as a vortex loop gas (see Sec. II). We show that when compared with the conventional insulator, this vortex loop gas has extra phase factors depending on the topology of the vortex loop configuration. We demonstrate that these wave functions simply capture all of the major phenomena of the SPT phases, both in the bulk and at the boundary, that were discussed in Ref.\(^14\). As a key example, in Sec. III we discuss a nontrivial SPT phase with symmetry either a direct \([U(1) \times \mathbb{Z}_T^2]\) (as is appropriate for a spin model realization of an interacting boson system) or semidirect \([U(1) \rtimes \mathbb{Z}_T^2]\) product. Here the vortex lines should be viewed as ribbons with a nonzero thickness and there is a phase \(-1\) associated with each self-linking of a vortex ribbon. We briefly also discuss a different SPT phase that occurs for \(U(1) \rtimes \mathbb{Z}_T^2\) where each vortex loop can be viewed as a 1D Haldane spin chain. These results are obtained by analyzing both the \(\sigma\)-model effective field theory and the topological “BF” effective field theories proposed in Ref.\(^14\) for these phases.

In 2D, a result with a similar flavor has been derived by Levin and Gu\(^15\) for an SPT phase with Ising, i.e., \(\mathbb{Z}_2\), which is symmetry in terms of a domain-wall loop gas with phase
factors. In Sec. IV, we reproduce this result using our methods. We also discuss the ground-state wave-function structure of the 2D boson topological insulator with $U(1) \times Z_T^2$ symmetry. We use this to obtain a dual vortex description of this state and show that the physics is correctly captured.

II. WAVE FUNCTION OF TRIVIAL 3D BOSE-MOTT INSULATOR

Let us start with briefly reviewing the trivial Mott insulating phase of bosons. This is conveniently modeled by a quantum disordered phase of interacting $U(1)$ rotors on a 3D lattice, which is described by the Hamiltonian $H = \sum_{(i,j)} -t \cos(\theta_i - \theta_j) + U(\bar{\theta}_i)^2$. The boson creation operator $b_i = e^{i\theta_i}$ and $n_j$ is the corresponding $U(1)$ charge at site $i$. $\theta_i$ and $n_i$ are canonically conjugate. The quantum disordered phase of the rotors is equivalent to the familiar Mott insulator phase and occurs when $t/U \ll 1$. In the strong-coupling limit $t \rightarrow 0$, the ground-state wave function is a trivial direct product state:

$$|\Psi\rangle = \prod_i |\theta_i = 0\rangle \sim \prod_i \int_0^{2\pi} d\theta_i |\theta_i\rangle. \quad (1)$$

The wave function of the quantum disordered phase with finite but small $t/U$ can be derived through perturbation on the wave function given by Eq. (1). For our purposes, it is useful to consider a simple approximate form of the wave function that captures the physics of the Mott phase,

$$|\Psi\rangle \sim \int_0^{2\pi} \prod_i d\theta \exp \left[ \sum_{(i,j)} -K \cos(\theta_i - \theta_j) \right] \prod_i |\theta_i\rangle, \quad (2)$$

where $K \sim t/U \ll 1$. This wave function is a superposition of configurations of $\theta_i$ with a weight that is the same as the Boltzmann weight of the 3D classical rotor model. The standard duality formalism of the 3D classical rotor model leads to the dual representation of this wave function,

$$|\Psi\rangle \sim \int D\vec{A} \sum_f \exp \left[ -\int d^3x \frac{1}{2K} (\vec{\nabla} \times \vec{A})^2 + i2\pi \vec{A} \cdot \vec{J} \right] \times |\vec{A}(x), \vec{J}(x)\rangle. \quad (3)$$

Vector field $\vec{J}$ takes only integer values on the dual lattice and it represents the vortex loop in the phase $\theta$. In order to guarantee the gauge invariance of $A$, $J$ must have no source in the bulk: $\vec{\nabla} \cdot \vec{J} = 0$. The vortex loop $\vec{J}$ can only end at the boundary, which corresponds to a 2D vortex. The $U(1)$ gauge field $\vec{A}$ induces long-range interactions between vortex loops with coupling strength $K$. In the limit $K \rightarrow 0$, i.e., the strong-coupling limit of the original rotor, the wave function given by Eq. (3) for quantum disordered lattice bosons becomes an equal-weight superposition of all vortex loop configurations, with a weak long-range interaction.

Quite generally, the Mott insulating phase is obtained when the vortex loops have proliferated. Consequently, the ground-state wave function can be described as a loop gas of oriented interacting vortex loops. The discussion above provides a derivation of this loop gas wave function starting from a simple but approximate microscopic boson wave function. A crucial point about the structure of the loop gas wave function for the trivial Mott insulator is that it has positive weight for all loop configurations.

III. WAVE FUNCTION OF 3D BOSONIC SPT PHASES

A 3D SPT phase with $U(1)$ symmetry is also a quantum disordered phase of rotor $\theta$, thus it is expected that its wave function is still a superposition of vortex loop configurations. However, more physics needs to be added to the vortex loops in order to capture the novel physics of the SPT phase. One of the central results of this paper is to determine the structure of this vortex loop gas wave function for the 3D SPT phases with $U(1)$, and time-reversal symmetry, as discussed in Ref. 14. We first focus on one example which occurs for both $U(1) \times Z_T^2$ and $U(1) \times Z_T^2$. We show that the ground state is described by a superposition of vortex loop configurations $|C_v\rangle$, but each vortex loop should be viewed as a “ribbon” rather than a line, and a self-linking of this ribbon contributes exactly factor $(-1)$ (Fig. 1):

$$|\Psi\rangle \sim \sum_{C_v} (-1)^{N_t} \psi_0[C_v] |C_v\rangle, \quad (4)$$

where $N_t$ is the number of self-linkings. Here, $\psi_0[C_v]$ is the weight of that vortex configuration in a trivial Mott insulator. The self-linking of a vortex ribbon is the linking number between the loops defined by the two ends of the ribbon.

This wave function given by Eq. (4) already implies that the vortex at the boundary must be a fermion. In Ref. 14, using a 2+1 boundary field theory, it was proved that the vortex of the $U(1)$ rotor at the boundary of this SPT is a fermion. (This means that the dual effective field theory of the surface is in terms of a fermionic vortex field rather than the usual dual vortex theory which is in terms of a bosonic vortex.) A vortex at the boundary is the end (source) of the vortex ribbon in the bulk. As was discussed in Ref. 16, and as shown below, exchanging the ends of ribbons is equivalent to twisting one of the ribbons by $2\pi$, which according to Eq. (4) should contribute factor $(-1)$. Thus, the bulk wave function given by Eq. (4) already implies that the vortex at the boundary must be a fermion.

\[ \Psi = \begin{array}{c} \ast \ast \ast \end{array} \]
The ground-state wave function of Eq. (5) can be derived either using a bulk nonlinear $\sigma$-model effective field theory for the SPT phase or using a bulk topological “BF” field theory. We will present both these derivations below. In Appendix B, we present a lattice regularized space-time path integral for these SPT phases, which is equivalent to these effective field theories and which may be preferred by some readers.

In order to describe the 3D SPT with either $\text{U}(1) \times Z_2^T$ or $\text{U}(1) \times Z_2^T$ symmetry, Ref. 14 proposed the following nonlinear $\sigma$ model that involves a five-component unit vector $\vec{n} = (n^1, \ldots, n^5)$, with a topological term $\Theta = 2\pi$:

$$ S = \int d^3x d\tau \frac{1}{g} \left(\partial_\mu \vec{n}\right)^2 + \frac{i\Theta}{\Omega_4} \epsilon_{abcdn^a \partial_\mu n^b \partial_\nu n^c \partial_\rho n^d \partial_\tau n^e}, $$

(5)

where $\Omega_4$ is the volume of a four-dimensional sphere with unit radius. Equation (5) has an enlarged SO(5) symmetry, but later we will reduce this symmetry down to physical $\text{U}(1) \times Z_2^U$ or $\text{U}(1) \times Z_2^T$.

In $3+1D$, an order-disorder phase transition occurs while tuning $g$. We will focus on the quantum disordered phase with strong coupling $g$. Since $\Theta = 2\pi$ in Eq. (5), its quantum disordered phase has the same bulk spectrum as the case with $\Theta = 0$. Thus, coupling constant $g$ flows to infinity in the quantum disordered phase, and this is the limit we will focus on in this paper.

The physical meaning of the $\Theta$ term in a nonlinear $\sigma$ model (NLSM) is usually interpreted as a factor $\exp(\Theta)$ attached to every instanton event in the space-time. Then this interpretation would lead to the conclusion that $\Theta = 2\pi$ is equivalent to $\Theta = 0$. However, this interpretation is very much incomplete because it only tells us that theories with $\Theta = 2\pi$ and 0 have the same partition function when the system is defined on a compact manifold. These two theories actually have very different ground-state wave functions. In order to expose the wave function, we need to keep an open boundary condition of time. In this case, the wave function can be derived using the following path integral:

$$ \langle \vec{n}(0) | \Psi \rangle \langle \Psi | \vec{n}'(x) \rangle \sim \int D\vec{n}(x,\tau) \exp(-S) \delta_{\vec{n}(\infty, 0) = \vec{n}, \vec{n}(0, \infty) = \vec{n}}. $$

(6)

The ground-state wave function $|\Psi\rangle$ can then be obtained straightforwardly in the strong-coupling limit $g \to +\infty$:

$$ |\Psi\rangle \sim \int D\vec{n}(x) W[\vec{n}] |\vec{n}(x)\rangle, $$

$$ W[\vec{n}] = e^{\frac{i\Theta}{\Omega_4} \int d^3x \int_0^\infty du \epsilon_{abcdn^a \partial_\mu n^b \partial_\nu n^c \partial_\rho n^d \partial_\tau n^e}}. $$

(7)

Here, $\vec{n}(x,u)$ is an extension of the real-space configuration $\vec{n}(x)$ that satisfies $\vec{n}(x,0) = (0,0,0,0,1)$ and $\vec{n}(x,1) = \vec{n}(x)$. Equation (7) is a superposition of all the configurations of the O(5) vector field $\vec{n}(x)$, with a weight that is proportional to the real-space Wess-Zumino-Witten (WZW) term $W[\vec{n}]$ at level 1. Thus the ground-state wave function of Eq. (5) with $\Theta = 2\pi$ is fundamentally different from the case with $\Theta = 0$. A similar relation between the bulk $\Theta$ term of $1+1D$ O(3) NLSM and its ground-state wave function was discussed previously in Ref. 17, in the context of a 1D spin chain.

Now let us reduce the artificial SO(5) symmetry of Eq. (5) to $\text{U}(1)$ and $Z_2^T$. We decompose the five-component vector $\vec{n}$ as $\vec{n} = [\sin(\alpha)\vec{\phi}, \cos(\alpha)\vec{\phi}]$, where $\vec{\phi}$ is a unit four-component vector and $\phi_0 = \pm 1$ is an Ising order parameter. We further define two bosonic rotor operators $b_1 \sim \phi_1 + i\phi_2$, $b_2 \sim \phi_3 + i\phi_4$. Under the $\text{U}(1)$ and $Z_2^T$, we take these variables to transform as

$$ Z_2^T : b_1, b_2 \to b_1, b_2 \left[\text{U}(1) \times Z_2^T\right], $$

$$ b_1, b_2 \to -b_1^*, -b_2^* \left[\text{U}(1) \times Z_2^T\right], $$

$$ \phi_0 \to -\phi_0, $$

$$ \text{U}(1) : b_1 \to e^{i\theta} b_1, \ b_2 \to e^{i\theta} b_2. $$

(8)

We assume the system favors $\vec{\phi}$ over $\phi_0$. If the time-reversal symmetry is preserved, namely, $\langle \phi_0 \rangle = 0$, then the WZW term in the wave function given by Eq. (7) reduces to a $\theta$ term for the four-component unit vector $\vec{\phi}$ in 3 + 0 dimensions. Thus, we get the following wave function:

$$ |\Psi\rangle \sim \int \mathcal{D}\vec{\phi}(x) $$

$$ \times \exp \left( \int d^3x \frac{i\pi}{12\pi^2} \epsilon_{abcd} \epsilon_{\mu\nu\rho} \phi^a \partial_\mu \phi^b \partial_\nu \phi^c \partial_\rho \phi^d \right) |\vec{\phi}(x)\rangle $$

$$ = \int \mathcal{D}\vec{\phi}(x) (-1)^N |\vec{\phi}(x)\rangle. $$

(9)

This wave function is a superposition of all configurations of $\vec{\phi}(x)$ in real space, with a $\theta$ term defined in 3D real space, at precisely $\theta = \pi / N_c$ is the Skyrmion number of the four-component vector $\vec{\phi}$, since $\pi / N_c \in \mathbb{Z}$. The value $\theta = \pi$ in the wave function is protected by time-reversal symmetry $Z_2^T$. If this $Z_2^T$ symmetry is broken, then $\theta$ in this wave function will be tuned away from $\pi$.

For our purposes, we need to introduce anisotropies that reduce the symmetry from O(4) to $\text{U}(1) \times \text{U}(1)$. Then the $\theta$ term (at $\theta = \pi$) implies that there is a phase factor of $-1$ each time the vortex loops of the two boson species link. Thus, this two-species boson Mott insulator has a wave function which is a superposition of all vortex loops of the two species with a crucial factor of $(-1)^k$, where $L$ is the total number of linked opposite species vortex loops. In contrast, for the trivial Mott insulator of the two boson species system, the weight for all vortex loop configurations can be taken to be positive.

It is implicit in the discussion in terms of a four-component unit vector $\vec{\phi}$ that classical configurations of the $b_{1,2}$ fields are always such that $b_{1,2}$ cannot simultaneously vanish. As the amplitude of either of these fields vanishes in their vortex core, this implies that the vortex loops of the two species cannot intersect. Thus, a configuration with a linking of the two vortex loops cannot be deformed to one without a linking.

This bulk wave function given by Eq. (9) also implies that at the 2D boundary, the vortex of $b_1$ and vortex of $b_2$ (sources of vortex loops) have a mutual semion statistics because when one flavor of vortex encircles another flavor through a full circle, the bulk vortex loops effectively acquire one extra linking (Fig. 2), which according to the bulk wave function would contribute the factor $(-1)$.

Let us now provide an alternate derivation of this result using the bulk topological BF theory for the SPT phase, also
two flavors of vortices at the boundary have mutual semion statistics.
−
the bulk wave function is a superposition of two flavors of vortex
wave function would contributes factor (9) reduces to wave function (4). As we mentioned
proposed in Ref. 14. This theory takes the form

\[ 2\pi \mathcal{L}_{\text{3D}} = \sum_{I} \epsilon^{\mu\nu\lambda\sigma} B_{\mu\sigma}^{I} \partial_{\nu} a_{\lambda}^{I} + \Theta \sum_{I,J} K_{IJ} \frac{1}{4\pi} \epsilon^{\mu\nu\lambda\sigma} \partial_{\mu} a_{\nu}^{I} a_{\lambda}^{I} a_{\sigma}^{J}, \]

(10)

Here, \( B_{\mu\sigma}^{I} \) is a rank-2 antisymmetric tensor that is related to the current of the bosons of species \( I = 1,2 \) through \( j_{\mu}^{I} = \frac{1}{2} \epsilon_{\mu\nu\lambda} \partial_{\nu} B_{\lambda\sigma}^{I} \). \( a_{\lambda}^{I} \) is a 1-form gauge field which describes the vortices of the bosons. Specifically, the magnetic field lines of \( a_{\lambda}^{I} \) are identified with the vortex lines of the boson of species \( I \). For the SPT state of interest, the \( K \) matrix is simply \( \sigma_{I} \). The parameter \( \Theta = \pi \) (not to be confused with the \( \theta \) parameter in the \( \sigma \)-model description). The crucial difference with the trivial Mott insulator is the second \( \Theta \) term. To get the ground-state wave function, we again evaluate the Euclidean path integral with open temporal boundary conditions. Using the well-known fact that the \( \Theta \) term is the derivative of a Chern-Simons term, we end up with the following ground-state wave functional:

\[ \psi[a_{\lambda}^{I}, B_{\mu\sigma}^{IJK}] \sim e^{i \int d^{3}\mathbf{x} K^{IJK} \epsilon^{\mu\nu\lambda} \partial_{\mu} a_{\nu}^{I} a_{\lambda}^{J} / a_{\nu}^{K}} \psi_{0}[a_{\lambda}^{I}, B_{\mu\sigma}^{IJK}] . \]

(11)

Here, \( \psi_{0} \) is the wave functional for the trivial Mott insulator. The wave functional for the SPT insulator is thus modified by a phase factor given by a \( 3 + 0 \)-dimensional Chern-Simons term. As is well known, the Chern-Simons term is related to a counting of the total linking number of the configuration of the magnetic flux lines of the gauge fields. Specializing to the case at hand, we see that in the presence of a \( 2\pi \) flux line of \( a_{1} \), there is a phase factor \( \Theta = \pi \) whenever a \( 2\pi \) flux line of \( a_{2} \) links with it. Thus, we reproduce the result that there is a phase of \( \pi \) associated with each linking of opposite species vortex lines.

Finally, if the \( U(1) \times U(1) \) symmetry is broken down to diagonal \( U(1) \), then the vortex loops of the two species will be confined to each other. The resulting common vortex loop of the rotor \( b \sim b_{1} \sim b_{2} \) becomes a ribbon, whose two edges are the vortex loops of \( b_{1} \) and \( b_{2} \). Further, for simplicity, we assume that there is an energetic constraint at short distances that prevents two vortex lines of the same species from approaching each other. In particular, we assume that the binding length scale of the opposite-species vortex loops is smaller than the allowed separation between same-species vortex loops.

Then the vortex ribbons cannot intersect each other. Such a “hard-core” constraint on the short-distance physics should not affect the universal long-distance behavior of the wave function. [Indeed, if we do allow intersection of the ribbons, then the phase factors are obtained by going back to the original \( U(1) \times U(1) \) theory with the two species of vortex loops. The embedding of the \( U(1) \) in the higher \( U(1) \times U(1) \) symmetry gives a short-distance “regularization” of ribbon intersections. Banning intersections enables us to discuss the essential physics without complications.] Note that the binding of the two species of vortex loops gives a physical implementation of the mathematical concept of “framing” used to describe the topology of knots. The linking between the two flavors of vortex loops becomes a self-linking of the ribbon. Thus, wave function (9) reduces to wave function (4). As we mentioned before, this bulk wave function given by Eq. (4) implies that the end point of a vortex ribbon at a 2D boundary is a fermion.

Similarly, bulk external sources for vortex ribbons will also be fermions. Such sources are points in three-dimensional space where we force the vortex lines to end. In the original boson Hilbert space, vortex lines cannot, of course, end. So these bulk vortex sources are to be regarded as probes of the system where we locally modify the Hilbert space. The statistics of these bulk external vortex sources is readily understood from the bulk wave function and is discussed in Appendix C. A quick hint of the fermionic statistics comes from asking about the behavior under \( 2\pi \) spatial rotations. The vortex ribbon emanating from a vortex source is twisted by \( 2\pi \), and this has an extra phase \( -1 \) compared to the untwisted ribbon. Thus, the vortex source has “topological” spin-1/2, as would be expected if it were a fermion.

All of these results concur with the boundary theory discussed in Ref. 14. There, a boundary field theory for the SPT is derived, which is a \( 2 + 1 \) NLSM with four-component vector \( \phi \), and there is a \( 2 + 1 \) space-time \( \Theta \) term at precisely \( \Theta = \pi \). This space-time \( \Theta \) term implies that the vortex of \( b_{2} \) carries 1/2 charge of \( b_{1} \), and vice versa. Thus, vortices of \( b_{1} \) and \( b_{2} \) have a mutual semion statistics. When the symmetry is broken down to one single \( U(1) \), the 2D vortex at the boundary becomes a bound state of the two flavors of vortices: thus, eventually this bound vortex becomes a fermion.

Reference 14 also described a different interesting SPT phase with \( U(1) \times Z_{2}^{T} \) symmetry. There the surface theory is such that the surface vortex carries a Kramer’s doublet in its core. A bulk effective field theory of this phase is also obtained \( 14 \) by starting with the \( O(5) \) nonlinear \( \sigma \) model [Eq. (5)] with anisotropies, but with a different realization of symmetry from the one described above. For instance, we can decompose five-component vector \( \vec{n} \) in a different way: \( \vec{n} = (\text{Re}[b], \text{Im}[b], N^{x}, N^{y}, N^{z}) \), where \( b \) is a rotor field that transforms under \( Z_{2}^{T} : b \rightarrow -b^{*} \). \( \vec{N} \) is a three-component vector that changes sign under \( Z_{2}^{T} \), but is uncharged under the.
IV. GROUND-STATE WAVE FUNCTION OF 2D SPT PHASES

A. 2D SPT phase with $Z_2$ symmetry

Let us now switch gears to SPT phases in 2D. We begin by making contact with the work of Levin and Gu\textsuperscript{15} on the Ising SPT phase. The simplest SPT phase in 2D has a $Z_2$ global symmetry. In Ref. 15, a lattice model for this phase has been discussed. The ground-state wave function was argued to be a superposition of all configurations of $\vec{n}$. The crucial ingredient is the $\Theta$ term for the four-component unit vector $\vec{\phi}$. The field theory action reads

$$S = \int d^2x d\tau \left( \frac{1}{g} (\partial_\mu \vec{\phi})^2 + \frac{i 2\pi}{12\pi^2} \epsilon_{abcd} \epsilon_{\mu
u\rho\sigma} \phi^a \partial_\mu \phi^b \partial_\nu \phi^c \partial_\rho \phi^d \right).$$

The crucial ingredient is the $\Theta$ term for the four-component unit vector at $\Theta = 2\pi$. This action given by Eq. (12) has an $SO(4)$ symmetry, and this $SO(4)$ symmetry contains a subgroup $Z_2$ symmetry $\vec{\phi} \rightarrow -\vec{\phi}$. Eventually, we will break the artificial $SO(4)$ symmetry of Eq. (12) down to this $Z_2$ subgroup.

In the paramagnetic phase, i.e., in the limit $g \rightarrow +\infty$, the bulk ground-state wave function is

$$|\Psi\rangle \sim \int D\vec{\phi}(x) \exp \left\{ \frac{i 2\pi}{12\pi^2} \int d^2x \int d^2x \epsilon_{abcd} \epsilon_{\mu
u\rho\sigma} \phi^a \partial_\mu \phi^b \partial_\nu \phi^c \partial_\rho \phi^d \right\} |\vec{\phi}(x)\rangle,$$  

which is a superposition of all configurations of $\vec{\phi}(x)$ with a real-space WZW weight.

Now let us decompose the four-component vector into $\vec{\phi} = [\cos(\alpha) \phi_0, \sin(\alpha) \vec{n}]$, where $\phi_0 = \pm 1$ is an Ising order parameter, and $\vec{n}$ is a unit three-component vector: $(\vec{n})^2 = 1$. We also break the $SO(4)$ symmetry down to $Z_2 \times SO(3)$ symmetry:

$$Z_2 : \phi_0 \rightarrow -\phi_0, \quad \vec{n} \rightarrow -\vec{n}, \quad SO(3) : \text{rotation of } \vec{n}.$$  

Under this symmetry reduction, if the system energetically favors vector $\vec{n}$ over $\phi_0$ (favors $\alpha = \pi/2$), then the wave function given by Eq. (13) reduces to

$$|\Psi\rangle \sim \sum_{n_\phi} (-1)^{N_\phi} |\vec{n}(x)\rangle,$$

where $N_\phi$ is the number of Skyrmions of O(3) vector $\vec{n}$ in the 2D space. As long as we keep the $Z_2$ symmetry $\vec{\phi} \rightarrow -\vec{\phi}$, and the expectation value of $\phi_0$ is zero, then each Skyrmion will contribute a phase factor of exactly $(-1)$. Now let us further break $Z_2 \times SO(3)$ symmetry down to $Z_2 \times SO(2)$:

$$Z_2 : \phi_0 \rightarrow -\phi_0, \quad \vec{n} \rightarrow -\vec{n}, \quad SO(2) : \text{rotation of } n^x, n^y.$$  

Also, we assume that the system favors $n^z$ over $n^x$ and $n^y$; then each Skyrmion becomes a domain wall of $Z_2$ order parameter $n^z(x)$ (Fig. 3). Now the wave function (15) reduces to a superposition of configurations of $Z_2$ order parameter $n^z(x)$:

$$|\Psi\rangle \sim \sum_{n^z} (-1)^{N_{n^z}} |n^z(x)\rangle,$$
where $N_d$ is the number of closed domain-wall loops of $n^i(x)$. Eventually, we can also break the residual SO(2) symmetry, and the wave function given by Eq. (17) is unchanged.

The wave function (17) is exactly the one derived from the lattice model of 2D SPT phase with $Z_2$ symmetry.\(^\footnote{In the appendix, we will also demonstrate that the effective field theory (12) implies that after coupling this SPT phase to a dynamical $Z_2$ gauge field, the $\pi$ flux of this $Z_2$ gauge field has a semion statistics, which is consistent with the result in Ref. 15.}

With a full SO(4) symmetry, the edge states of Eq. (12) with precisely $\Theta = 2\pi$ is the nonchiral SU(2)$_1$ conformal field theory [or, equivalently, as an SO(4) nonlinear $\sigma$ model with a level-1 WZW term]. Since the original SO(4) symmetry is reduced to its $Z_2$ subgroup $\phi \rightarrow -\phi$, we have to argue that the edge state of Eq. (12) survives under this symmetry reduction. Because the $Z_2$ symmetry acts on all four components of $\phi$, in the boundary WZW model, terms allowed by the $Z_2$ symmetry are $\sum_{i,j} g_{ij} \phi_i \phi_j$ ($i,j = 0,1,2,3$). If these terms are relevant, it leads to spontaneous $Z_2$ symmetry breaking and twofold degeneracy at the boundary. Thus, the edge state cannot be completely trivial (gapped and nondegenerate) as long as the $Z_2$ symmetry is preserved.

\textbf{B. 2D SPT phase with $U(1) \times Z_2^2$ symmetry}

Finally, we consider the 2D bosonic topological insulator which occurs when the global symmetry is $U(1) \times Z_2^2$. This may be described by starting again with the same four-component nonlinear $\sigma$ model, but with the following implementation of the physical symmetry. We write $\phi_2 = i\phi_3 = b$ and let $b$ have charge $1$ under the global $U(1)$. Under $Z_2^1$, we demand $b \rightarrow b, \phi_0 \rightarrow -\phi_0, \phi_1 \rightarrow -\phi_1$. As before, we again assume first an anisotropy that prefers $i\bar{n}$ over $n$ so that the ground-state wave function is given by Eq. (15). Now we introduce further anisotropy to reduce to the desired $U(1) \times Z_2^2$. The defects of the charged field $b$ are, of course, point vortices. In the core of these vortices, the amplitude of $b$ is suppressed and the $\phi$ points entirely in the $\phi_1$ direction. There are two different vortices—known as merons—depending on whether in the core $\phi_1 = \pm 1$. Each meron may be viewed as half a Skyrmion and has $N_r = \text{sgn}(\phi_1)/1/2$. Thus, the ground-state wave function is then a sum over all possible configurations of the two kinds of meron vortices with phase factors $e^{\pm i\pi/2}$ for the two kinds of vortex. Let $n_{\pm}$ be the vortex number of either species at site $i$ in a lattice description. Then, we require that the total vorticity $N_c = \sum_i n_{+} + n_{-} = 0$. Then, the phase factor in the wave function is $e^{\pm i\pi\sum(n_{+} - n_{-})} = (-1)^{\sum n_{+}}$. Thus, there is a relative phase of $-1$ associated with $-1$ vortices compared to $+1$ vortices.

We now argue that this structure of the wave function matches what is known about the 2D bosonic topological insulator. Consider a dual description of such an insulator. From our arguments above, there are two kinds of vortex fields $\Phi_{+}$ corresponding to the two meron vortices. The dual vortex theory will have a Lagrangian

$$L_d = \sum_{\mu=\pm} |(\partial_{\mu} - ia_{\mu})\Phi_{\pm}\phi^\ast_{\pm}|^2 + \cdots + \frac{k}{2}(\epsilon_{\mu\nu\lambda}\partial_{\mu}a_{\nu})a_{\lambda}^2.$$

Here, $a_{\mu}$ is the usual fluctuating noncompact $U(1)$ gauge field whose flux density is the original boson number. There must, in addition, be terms where the meron cores tunnel into each other:

$$\lambda(\Phi_{+}^\ast, \Phi_{-}\phi^\ast_{+} + H.c.).$$

We begin by ignoring these and we will reinstate them later. Under time reversal, a vortex must go to an antivortex (as the boson phase is odd) and the meron cores flip into each other. Thus, under $Z_2^1$,

$$\Phi_{+} \rightarrow -\Phi_{-}^\ast, \Phi_{-} \rightarrow -\Phi_{+}^\ast.$$  \footnote{In the trivial insulator, all vortex configurations contribute with the same sign and we must choose $\Phi_{+} \rightarrow +\Phi_{-}^\ast$. But for the topological insulator, there is a relative $-\text{sign}$ between the two vortex species. Thus, we must choose $\Phi_{+} \rightarrow -\Phi_{-}^\ast$. Condensing vortices that transform in this manner will give us the boson topological insulator. Now let us include the meron core tunneling term. Then the two vortex species mix with each other so that we identify a single vortex, $\Phi_v = \Phi_{+} \sim \Phi_{-}^\ast$. Its transformation under time reversal is $\Phi_v \rightarrow \pm \Phi_v^\ast$, where the $+\text{sign}$ describes the trivial insulator and the $-\text{sign}$ describes the topological insulator. This is exactly the same transformation law for the vortices that is dictated by the edge theory analysis of the 2D boson topological insulator.\footnote{Thus, the wave-function description we developed captures the physics of this state, and further gives a bulk dual vortex description.}  \footnote{In summary, we have demonstrated in this work that although most of the novel phenomena of a SPT phase occur at its boundary, its bulk ground-state wave function is indeed drastically different from a trivial direct product disordered phase. This bulk wave function can be conveniently derived from the effective field theory of the SPT phase. The structure of the ground-state wave functions in terms of dual vortex configurations derived in this work provide a simple physical picture of the phenomena associated with these SPT phases. The dual ground-state bulk wave function provides a nice intuitive understanding of the differences between ordinary and topological boson insulators. However, in this paper, we have not attempted to explicitly construct microscopic models for these phases. Progress in this direction is reported very recently in Refs. 23 and 24, which appeared after this paper was submitted.}

Finally, we note that for the 3D boson topological insulators, there is some superficial similarity with the wave functions of Walker and Wang,\footnote{though a detailed understanding of the relationship is presently not clear to us. The Walker-Wang models also have ground-state wave functions as string net configurations with amplitude determined by a $2+1$-dimensional topological quantum field theory. In some cases, these can correspond to SPT phases (see Ref. 24). However, it is not clear how the strings in the Walker-Wang models are related to the physical bosons; in particular, they are not to be identified with physical vortex loops. The exploration of the connections between the wave functions in these Walker-Wang constructions of SPT phases and our dual wave functions is an interesting avenue for future research.} though a detailed understanding of the relationship is presently not clear to us. The Walker-Wang models also have ground-state wave functions as string net configurations with amplitude determined by a $2+1$-dimensional topological quantum field theory. In some cases, these can correspond to SPT phases (see Ref. 24). However, it is not clear how the strings in the Walker-Wang models are related to the physical bosons; in particular, they are not to be identified with physical vortex loops. The exploration of the connections between the wave functions in these Walker-Wang constructions of SPT phases and our dual wave functions is an interesting avenue for future research.
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APPENDIX A: DYNAMICAL $Z_2$ GAUGE FIELDS IN THE 2D ISING SPT

In this Appendix, we demonstrate that the effective field theory given by Eq. (12) not only gives us the correct ground-state wave function [Eq. (17)] of the 2D Ising SPT phase, but, after coupling the SPT phase to a $Z_2$ state wave function [Eq. (17)] of the 2D Ising SPT phase, gives the effective theory of Ref. 14. This enables a further elaboration of the considerations leading to the bulk effective “BF phases discussed in this paper. We will also briefly review the considerations leading to the bulk effective “BF phases discussed in this paper. We will also briefly review the considerations leading to the bulk effective “BF phases discussed in this paper. We will also briefly review the considerations leading to the bulk effective “BF phases discussed in this paper. We will also briefly review the considerations leading to the bulk effective “BF phases discussed in this paper.

First of all, Eq. (12) can be rewritten as a SU(2) principle chiral model, by introducing SU(2) matrix field $G = \phi^0\sigma^0 + i\phi\cdot\sigma$. $G$ has SU(2)-left and SU(2)-right transformations: $G \to V^L G V_R$. Let us “gauge” SU(2)-left and SU(2)-right transformations with dynamical U(1) gauge fields $a_{\mu}\sigma^z$ and $b_{\mu}\sigma^\pm$, i.e., replace $\partial_{\mu}G$ with $\partial_{\mu}G + i a_{\mu}\sigma^z G + ib_{\mu}G \sigma^\pm$. According to Refs. 26 and 27, after integrating out matrix field $G$, gauge fields $a_{\mu}$ and $b_{\mu}$ both acquire a Chern-Simons term:

$$S_{CS} = \frac{\kappa}{4\pi} \int d^2 x d\tau \left[ \frac{i}{2} \epsilon_{\mu\nu\rho} \partial_\mu a_{\nu} \partial_\rho a_{\mu} - \frac{i}{4\pi} \epsilon_{\mu\nu\rho} b_{\mu} \partial_\nu b_{\rho} \right].$$

This is because Eq. (12) also describes a U(1) bosonic SPT with Hall conductivity 2.

A dynamical U(1) gauge field with a Chern-Simons action at level-k has the following properties: its charged quasiparticle carries gauge flux $2\pi/k$, and this quasiparticle has a statistics angle $\pi/k$. Thus, the Chern-Simons action given by Eq. (A1) gives the $\pi$ flux of U(1) gauge field $a_\mu$ and $b_\mu$ a semion statistics, with statistics angle $+\pi/2$ and $-\pi/2$, respectively. Notice that the two U(1) gauge groups share the same statistics, with statistics angle $\pi/k$, and this quasiparticle has a statistics angle $\pi/k$. Thus, the Chern-Simons action given by Eq. (A1) gives the $\pi$ flux of U(1) gauge field $a_\mu$ and $b_\mu$ a semion statistics, with statistics angle $+\pi/2$ and $-\pi/2$, respectively. Notice that the two U(1) gauge groups share the same statistics, with statistics angle $+\pi/2$ and $-\pi/2$, respectively. Notice that the two U(1) gauge groups share the same statistics, with statistics angle $+\pi/2$ and $-\pi/2$, respectively. Notice that the two U(1) gauge groups share the same statistics, with statistics angle $+\pi/2$ and $-\pi/2$, respectively. Notice that the two U(1) gauge groups share the same statistics, with statistics angle $+\pi/2$ and $-\pi/2$, respectively. Notice that the two U(1) gauge groups share the same statistics, with statistics angle $+\pi/2$ and $-\pi/2$, respectively.

In Ref. 15, using their lattice model, the authors concluded that the dynamical $\pi$ flux of this $Z_2$ gauge field has a semion statistics. Here we have derived the same result using our field theory (12).

APPENDIX B: LATTICE VERSION OF EFFECTIVE FIELD THEORY

In this Appendix, we briefly discuss a lattice regularized version of the bulk effective field theory for the 3D SPT phases discussed in this paper. We will also briefly review the considerations leading to the bulk effective “BF + FF” effective theory of Ref. 14. This enables a further elaboration of the discussion in the main text on the properties of external sources of bulk vortex lines.

To set the stage, first consider the Euclidean lattice action for a 3 + 1-dimensional XY model in Villain form,

$$S_0 = \sum_{j_\mu} \left[ \frac{g}{2} (j_\mu)^2 + ij_\mu (\nabla_\mu \theta - 2\pi m_\mu) \right].$$

Here, $m_\mu$ is an integer defined on the links. Physically, it defines the integer vortex current

$$J_{\mu\nu} = \epsilon_{\mu\nu\lambda\kappa} \nabla_\lambda m_\kappa.$$

It is slightly more convenient to go to a “gauche” where we explicitly sum over $m_0$, which has the effect of forcing $j_0$ to be an integer. We can then drop $m_0$ from the action. Let us also define $a_\mu = 2\pi m_\mu$. Now consider two species of bosons, i.e., two XY models. We put one XY model on a 4D cubic lattice and the other on a different 4D cubic lattice such that the spatial links of one lattice penetrate the spatial plaquettes of the other. (This different treatment of time and space is not necessary, but it helps one to visualize.) Formally, the lattice sites of one lattice are $(l_0, l_1, l_2, l_3)$ with $l_0 = \text{integer}$, and for the other lattice we have $(l_0, l_1 + \frac{1}{2} l_2, l_3 + \frac{1}{2})$. Then, let us write the action for the coupled XY models:

$$S = \sum_{l=1,2} S_{0l} + S_{top},$$

$$S_{top} = i \frac{\Theta}{8\pi} K_{1,2} \epsilon_{\mu\nu\lambda\kappa} \nabla_\mu a_\nu \nabla_\lambda a_\kappa,$$

with the matrix $K = \sigma_0$. The first term $S_{0l}$ is just the sum of the above $XY$ actions for each species. The second “topological” term enforces the phases associated with the vortex world sheet configurations. In the trivial boson insulator, $\Theta = 0$, while in the boson SPT phase, $\Theta = \pi$. Thus, the Boltzmann weight for the SPT phase differs from that of the trivial insulator only through phase factors that depend on the vortex world sheet configurations.

The action in Eq. (B3) can be regarded as a lattice version of the continuum nonlinear $\sigma$ model used in the main paper. Indeed, the term $S_{top}$ correctly captures the physics of the $\theta$ term of the $\sigma$ model. We now show the relation to the BF + FF effective theory. As usual, a dual description of the boson system is obtained by writing the conserved boson 4-currents in terms of dual 2-form fields $B^{I\mu\nu}$:

$$j_{I\mu} = \frac{1}{2\pi} \epsilon_{\mu\nu\lambda\kappa} \nabla_\nu B^{I\lambda}_\kappa.$$

The term $j_{I\mu}$ in the lattice action above then leads to the familiar BF term. The integer constraint on $m_\mu = \frac{a_\mu}{2\pi}$ can be implemented softly by including a term $-\lambda \cos(a_\mu - \nabla_\mu \theta^I)$. The $\theta^I$ is just the original boson phase. As explained in Ref. 14, in the Mott insulator at energies below the boson gap, the bosons may be integrated out to leave behind just a Maxwell term for the $a_\mu$ fields. The term $S_{top}$ in the lattice action given by Eq. (B3) simply becomes the FF term of the BF + FF action.

The properties of external sources of vortex lines can be readily discussed in terms of these bulk effective field theories. Since the difference with the trivial insulator comes entirely from the $S_{top}$ term, it is appropriate to focus first on the vortex world sheet configurations, take care of the consequences of the
which according to the bulk wave function should acquire factor to creating one extra vortex ribbon with 2π. Interchanging two vortex sources is also homotopically equivalent to self-twisting one of the two ribbons by 2π in (e), which according to the bulk wave function should acquire factor −1. (f) Interchanging two vortex sources is also homotopically equivalent to creating one extra vortex ribbon with 2π self-twist in the bulk.

\[ S_{\text{top}} \]

and then put back the interaction with the smooth part of the boson phase fields represented by the coupling to \( B_{\mu \nu} \). External bulk sources of vortex lines are simply monopole sources of 2\( \pi \) magnetic flux of the internal gauge fields \( a_\mu \). We may now specialize to the SPT phase for \( U(1) \times Z_2 \) discussed in the main text where the boundary vortex is a fermion. Now consider letting the two boson species tunnel into each other so that the corresponding vortex lines are bound together to form the vortex ribbon. Formally, in a coarse-grained description, this is obtained by setting \( a_{1\mu} \sim a_{2\mu} \). Bulk external sources of the vortex ribbon are simply monopoles of this common internal gauge field. The wave-function-based discussion of this paper shows that these vortex sources are fermions. Alternately, this follows from the observation of Ref. 14 that the vortex field of the boundary dual vortex theory is a fermion.

**APPENDIX C: STATISTICS OF VORTEX SOURCES**

According to the main text, a vortex at the 2D boundary of the \( U(1) \times Z_2 \) SPT phase is a fermion. Here we consider bulk vortex sources. Consider the vortex loop left behind by creating a vortex source and antivortex source together, moving the vortex source around and then annihilating with the antivortex source. Now consider creating another such loop. We can envisage two situations depending on whether or not the two vortex sources were exchanged with each other during the process of forming the two closed loops. Pictures corresponding to these two situations may be found in Fig. 2 of the well-known paper by Wilczek and Zee. As argued in that paper, the difference between the two pictures corresponds to a self-linking of one of the ribbons by 2\( \pi \). This means that the process of exchange of the vortex sources has introduced a phase −1, and thus the vortex sources are fermions.

Here we provide further pictures to illustrate this in Fig. 4. Let us consider two vortex sources in the bulk [Fig. 4(a)]. After interchanging two vortex sources [Fig. 4(b)], the vortex-ribbon configurations can be continuously deformed into Fig. 4(e) \([b \rightarrow c \rightarrow d \rightarrow e]\), which is simply self-twisting one of the ribbons by 2\( \pi \). In Fig. 4(b), the ribbon (red and green vortex lines) connecting to the right vortex source is on the top. Step 1: (b) \(\rightarrow (c) \), connect the two red (green) lines in (b) and reopen in the horizontal direction. Step 2: (c) \(\rightarrow (d) \), deform the red line on the right into a circle and a straight line; (d) \(\rightarrow (e) \), reconnect the red circle to the red line on the left, and now the ribbon on the left has a 2\( \pi \) self-twist.

Figure 4(b) can also be continuously deformed into Fig. 4(f), which compared with Fig. 4(a) has created another ribbon with a 2\( \pi \) self-twist, or, equivalently, two different vortex loops with linking number 1. Both configurations 4(e) and 4(f) introduce factor \((-1)\) compared with 4(a).

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