Topological photonic bands in two-dimensional networks of metamaterial elements

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Abstract. We show that topological frequency band structures emerge in two-dimensional electromagnetic lattices of metamaterial components without the application of an external magnetic field. The topological nature of the band structure manifests itself by the occurrence of exceptional points in the band structure or by the emergence of one-way guided modes. Based on an electromagnetic network with nearly flat frequency bands of nontrivial topology, we propose a coupled-cavity lattice made of superconducting transmission lines and cavity QED components, which is described by the Jaynes–Cummings–Hubbard model and can serve as a simulator of the fractional quantum Hall effect.

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1. Introduction

The topological description of the quantum states of matter sets a new paradigm in the description and classification of atomic solids. Namely, atomic solids whose energy band structure possesses nontrivial topological properties constitute a new class of materials whose salient properties are robust to phase transitions which modify the symmetry order of the atomic solid. Prominent examples of such topological atomic solids are integer/fractional quantum Hall effect (I/FQHE) systems and topological insulators (TIs). Well-known examples of topological properties are the existence of chiral edge states in QHE systems and the presence of gapless surface states in TIs, both of which are immune to order–disorder phase transitions.

The advent of artificial electromagnetic (EM) structures such as photonic crystals and metamaterials has established over the years a continuous conveyance of ideas and methods from atomic solids to their EM counterparts. Quite naturally, the concept of topological order has been adapted to photonic crystals starting with the QHE: a two-dimensional (2D) lattice of gyromagnetic/gyroelectric cylinders is a system with broken time-reversal symmetry \cite{1} with frequency bands characterized by a nonzero Chern number, allowing for the emergence of unidirectional (one-way) edge states \cite{2} in analogy with the chiral edge states in QHE systems such as a 2D electron gas or graphene nanoribbons under a magnetic field. Anomalous QHEs can also be simulated with artificial chiral metamaterials of gyromagnetic components \cite{3}.

In TIs \cite{4} and quantum spin Hall systems \cite{5} in 2D, the presence of magnetic field is not a prerequisite for the appearance of topological electron states. In analogy with atomic TIs, in certain 3D photonic crystals and metamaterials with proper design, topological frequency bands appear without including gyromagnetic/gyroelectric materials, which require the application of an external magnetic field in order to break time-reversal symmetry \cite{6–8}.

In this work, we propose a class of 2D EM networks possessing topological frequency bands without the application of an external magnetic field. Namely, we show that topological bands emerge in 2D lattices of EM resonators connected with left- and right-handed metamaterial elements such as transmission lines (TLs) or waveguides loaded with a negative refractive-index medium. The topological nature of the corresponding frequency bands is manifested by the emergence of an exceptional point for transverse electric (TE) and by the generation of one-way modes for transverse magnetic (TM) waves. In the latter case, the system can be viewed as a simulator of the FQHE for polaritons.

2. Lattice of coupled dipoles

The EM crystals under study here are amenable to a photonic tight-binding (TB) description within the framework of the coupled-dipole method \cite{9}. The latter is an exact means of solving Maxwell’s equations in the presence of nonmagnetic scatterers. We consider a lattice of cavities within a lossless metallic host. The $i$th cavity is represented by a dipole of moment $\mathbf{P}_i = (P_{ix}, P_{iy}, P_{iz})$, which stems from an incident electric field $\mathbf{E}^\text{inc}$ and the field that is scattered by all the other cavities of the lattice. This way the dipole moments of all the cavities are coupled to each other and to the external field, leading to the coupled-dipole equation

$$\mathbf{P}_i = \alpha_i(\omega) \left[ \mathbf{E}^\text{inc} + \sum_{i' \neq i} G_{ii'}(\omega) \mathbf{P}_{i'} \right]. \quad (1)$$
\( G_{i,i'}(\omega) \) is the electric part of the free-space Green’s tensor and \( \alpha_i(\omega) \) is the polarizability of the \( i \)th cavity. Equation (1) is a \( 3N \times 3N \) linear system of equations, where \( N \) is the number of cavities of the system.

For a particle/cavity of electric permittivity \( \epsilon \) embedded within a material host of permittivity \( \epsilon_h \), the polarizability \( \alpha \) is provided by the Clausius–Mossotti formula \( \alpha = (3V/4\pi)(\epsilon - \epsilon_h)/(\epsilon + 2\epsilon_h) \), where \( V \) is the volume of the particle/cavity. For a lossless Drude-type (metallic) host, i.e. \( \epsilon_h(\omega) = 1 - \omega_p^2/\omega^2 \) (where \( \omega_p \) is the bulk plasma frequency), the polarizability \( \alpha \) exhibits a pole at \( \omega_0 = \omega_p\sqrt{2/(\epsilon + 2)} \) (surface plasmon resonance). By making a Laurent expansion of \( \alpha \) around \( \omega_0 \) and keeping the leading term [8], we may write \( \alpha = F/(\omega - \omega_0) \equiv F/\Omega \), where \( F = (27V/8\pi)\omega_0\epsilon/(2\epsilon + 4) \). For a sufficiently high value of the permittivity of the dielectric cavity, the electric field of the surface plasmon is much localized at the surface of the cavity. As a result, in a periodic lattice of cavities, the interaction of neighboring surface plasmons is very weak, leading to much narrow frequency bands. By treating such a lattice in a TB manner [8], we may assume that the Green’s tensor \( G_{i,i'}(\omega) \) does not vary much with frequency and therefore \( G_{i,i'}(\omega) \approx G_{i,i'}(\omega_0) \).

In this case, equation (1) becomes an eigenvalue problem

\[
\sum_{i \neq i'} G_{i,i'}(\omega_0) P_i = \Omega P_i, \tag{2}
\]

where \( F \) has been absorbed within the definition of \( G_{i,i'}(\omega_0) \) and we have set \( E^{inc} = 0 \) in equation (1) as we are seeking the eigenmodes of the system of cavities. In the following, we will be dealing with 2D lattices of cavities. We can, therefore, treat the case where the electric field lies within the plane of cavities (TE modes) and that where the electric field is perpendicular to the plane (TM modes) separately.

### 2.1. Transverse electric modes

In this case, \( P_i = (P_{i,x}, P_{i,y}) \) and the Green’s tensor \( G_{i,i'}(\omega_0) \) is given by

\[
G_{i,i'}(\omega_0) = F q_0^3 \begin{bmatrix} C(q_0|r_{i,i'|}) I_2 + J(q_0|r_{i,i'|}) \end{bmatrix} \begin{pmatrix} x_{i,i'}^2 & x_{i,i'} y_{i,i'} \ x_{i,i'} y_{i,i'} & y_{i,i'}^2 \ x_{i,i'} & y_{i,i'} \ x_{i,i'} & y_{i,i'} \ x_{i,i'} & y_{i,i'} \ \end{pmatrix}, \tag{3}
\]

with \( r_{i,i'} = r_i - r_i' \), \( q_0 = \omega_0\sqrt{\epsilon_h(\omega_0)}/c \) and \( I_2 \) is the \( 2 \times 2 \) unit matrix. Since we focus our attention on the surface plasmon frequency \( \omega_0 \), we operate in the subwavelength regime where \( q_0|r_{i,i'|}| \ll 1 \). In this regime, the functions \( C(q_0|r_{i,i'|}) \), \( J(q_0|r_{i,i'|}) \) are written as

\[
q_0^3 FC(q_0|r_{i,i'|}) \simeq -q_0^3 F J(q_0|r_{i,i'|}) \simeq q_0^3 F \exp(iq_0|r_{i,i'|}|/(q_0|r_{i,i'|}|)) = t_{i,i'} \exp(i\phi_{i,i'}), \tag{4}
\]

where \( t_{i,i'} \) and \( \phi_{i,i'} \) are real numbers. In what follows, the cavities are connected via coupling elements, i.e. waveguides or TLs, in which case the phase factors \( \phi_{i,i'} \) are not necessarily related with the wavevector of the host medium \( \epsilon_h \) and can therefore be considered as independent parameters.

For a 2D lattice of cavities, we assume the Bloch ansatz for the polarization field, i.e.

\[
P_i = P_{n\beta} = \exp(ik \cdot R_n) P_{0\beta}. \tag{5}
\]
Figure 1. (a) Square lattice of EM resonators connected with metamaterial-based coupling elements. The arrows denote NN hoppings and the solid lines NNN hoppings. The direction of the arrow shows whether wave propagation in the coupling element is left or right handed. (b) Frequency band structure corresponding to the lattice of the left panel for \( t = t' = 1, \phi = \pi/3 \).

The cavity index \( i \) becomes composite, \( i = n\beta \), where \( n \) enumerates the unit cell and \( \beta \) the positions of inequivalent cavities in the unit cell. Also, \( R_n \) denotes the lattice vectors and \( k = (k_x, k_y) \) is the Bloch wavevector. By substituting equation (5) into equation (2), we finally obtain

\[
\sum_{\beta'} \tilde{G}_{\beta\beta'}(\omega_0, k) P_{0\beta'} = \Omega P_{0\beta},
\]

where

\[
\tilde{G}_{\beta\beta'}(\omega_0, k) = \sum_{n'} \exp[ik \cdot (R_n - R_{n'})] G_{n\beta; n'\beta'}(\omega_0).
\]

Solution of equation (6) provides the TE frequency band structure of a periodic system of cavities.

In order to seek topological Bloch modes in a 2D lattice, we need at least two distinct frequency bands. Since TE modes correspond to two degrees of freedom for the polarization field, i.e. \( (P_x, P_y) \), we may consider a 2D lattice with one cavity per unit cell. Namely, we consider the square lattice of figure 1(a) where we consider nearest-neighbor (NN) and next-nearest-neighbor (NNN) hoppings of the EM field among the cavities. The NN hopping carries a nonzero phase \( t \exp(\pm i\phi) \) whose signs are denoted by arrows in figure 1(a). The NNN hopping is denoted by \( t' \). For the lattice of figure 1(a), the Green’s tensor \( \tilde{G} \) of equation (7) becomes

\[
\tilde{G} = t[\cos \phi \cos(k_x a/2) \cos(k_y a/2) - i \sin \phi \sin(k_x a/2) \sin(k_y a/2)] \begin{pmatrix} 3 & -1 \\ -1 & 3 \end{pmatrix} + 2t' \begin{pmatrix} \cos(k_x a) & 0 \\ 0 & \cos(k_y a) \end{pmatrix}.
\]

Figure 1(b) shows the frequency band structure derived from equation (8) for \( t = t' = 1, \phi = \pi/3 \). We observe that at some point along the \( XM \) symmetry line the frequency bands coalesce into a single band. The point beyond which the bands coalesce is an exceptional point.
and has been observed in $\mathcal{PT}$-symmetric lattices [10]. In general, exceptional points emerging in parameter space are associated with topological charge and geometric (Berry) phase [11]. The topological properties of an exceptional point have been revealed by encircling it in parameter space [12], as was demonstrated in a microwave cavity experiment [13]. Although a proper theory of the topological properties of the exceptional points in lattices is still lacking, based on previous work [11] we can indirectly assign topological charge and geometric phase to the exceptional point appearing in the frequency band structure of figure 1(b).

2.2. Transverse magnetic modes

Next, we assume that the polarization at each dipole is oriented in the $z$-axis. In this case, the eigenvalue problem of equation (2) becomes scalar and

$$G_{ii'} = F q_0^3 C(q_0 |r_{ii'}|) \simeq q_0^3 F \exp(i q_0 |r_{ii'}|) / (q_0 |r_{ii'}|) = t_{ii'} \exp(i \phi_{ii'}).$$  \hspace{1cm} (9)

The same applies to equations (6) and (7) and the TM problem becomes equivalent to the electronic case. Since the minimal model to have topological frequency bands is a two-band model, we adopt the checkerboard lattice of [14] (see figure 2(a)). Namely, apart from considering NN and NNN hoppings as in figure 1(a), we also consider next-next-nearest-neighbor (NNNN) hoppings (denoted by arcs in figure 2(a)) with strength $t''$. The NN hoppings are, again, complex, $t \exp(\pm i \phi)$, where the sign is denoted by arrows in figure 2(a). The NNN hopping strength is $t'_1$ ($t'_2$) if two sites are connected by a solid (dashed) line. We note that in EM lattices such as those considered here, a negative phase $-\phi$ can be easily achieved when the cavities are connected, e.g., by 1D left-handed transmission lines (LHTLs), i.e. transmission lines supporting backward-propagating waves where the phase velocity is opposite to the group velocity [15, 16]. Alternatively, the cavities may be connected by waveguides loaded with a left-handed (LH) metamaterial. Obviously, a positive phase $+\phi$ can be achieved by a similar means (right-handed transmission lines (RHTLs)).

For the checkerboard lattice of figure 2(a), the Green’s tensor $\tilde{G}_{\beta\beta'}$ becomes

$$\tilde{G} = \begin{pmatrix} G_{11} & G_{12} \\ G_{21} & G_{22} \end{pmatrix},$$  \hspace{1cm} (10)

where

$$G_{11} = 2t'_1 \cos(k_\alpha) + 2t'_2 \cos(k_\alpha) + 4t'' \cos(k_\alpha) \cos(k_\alpha),$$

$$G_{12} = G_{21} = 4t \cos \phi \cos(k_\alpha/2) \cos(k_\alpha/2) - 4i t \sin \phi \sin(k_\alpha/2),$$

$$G_{22} = 2t'_1 \cos(k_\alpha) + 2t'_2 \cos(k_\alpha) + 4t'' \cos(k_\alpha) \cos(k_\alpha).$$  \hspace{1cm} (11)

Atomic lattices with exotic hoppings such as those considered here have been used for simulating the FQHE states at zero magnetic field, as nearly flat topological bands can emerge which simulate the Landau levels associated with a uniform magnetic field [17–19]. By taking as TB parameters [14] $t = 1$, $\phi = \pi/4$, $t'_1 = -t'_2 = 1/(2 + 2 \sqrt{2})$, $t'' = 1/(2 + 2 \sqrt{2})$, a nearly flat band emerges as is evident from the frequency band structure of figure 3(a). Based on the equivalence of the Green’s tensor of equations (10) and (11) to the Hamiltonian of the electronic problem [14], each of the two bands of figure 3(a) carries a Chern number $\pm 1$. The topological nature of the frequency bands of the lattice of figure 2(a) is also manifested by the emergence of one-way bands (the photonic counterpart of the electron chiral edge states [1, 2]) in the frequency band structure for a slab geometry of figure 3(b).
Figure 2. (a) Checkerboard lattice of EM resonators connected with metamaterial-based coupling elements. The arrows denote NN hoppings and the solid and broken lines NNN hoppings. The direction of the arrow shows whether wave propagation in the coupling element is left or right handed. Two of the NNNN hoppings are shown as dotted arcs. (b) The EM resonator of the lattice is a superconducting circuit QED system consisting of an LC resonator coupled to a Cooper-pair box (CPB). Typical TL for (c) right- ((d) left-) handed coupling elements along with the corresponding dispersion relation.

The occurrence of topological properties such as the exceptional point in figure 1(b) and the one-way modes in figure 3 are a result of the synthetic gauge field, which is generated by the geometry of the metamaterial-based coupling elements (the formation of closed flux loops of the phase of the EM field).

3. Simulation of the fractional quantum Hall effect

Having established a nearly flat topological frequency band of the EM field for the lattice of figure 2(a), we are able to design a system for creating a (semi-classical) simulator of the FQHE. The most natural choice would be to consider a coupled cavity array (CCA) wherein polaritons propagate through a hopping mechanism (as in our case) and interact strongly with the reservoir
Figure 3. (a) Frequency band structure for the infinite checkerboard lattice of figure 2 and (b) the frequency band structure for a finite slab consisting of 20 unit planes (parameters: $t = 1$, $\phi = \pi/4$, $t'_1 = -t'_2 = 1/(2 + \sqrt{2})$, $t'' = 1/(2 + 2\sqrt{2})$).

of modes when they reside within the cavity [20, 21]. The FQHE with a magnetic field can also be simulated by atoms confined in a 2D CCA [22]. A way to realize a CCA (in the microwave regime) would be to follow the remedy of [8], wherein a lattice of dielectric cavities is immersed within an artificial, lossless, plasmonic host medium [23] such as a network of thin metallic wires of a few tens of $\mu$m in diameter and spaced a few mm apart. The dielectric cavities are to be connected with each other with waveguides (see figure 4 of [8]), which are the coupling elements in our TB formalism. This way, the interaction between cavities is achieved only via the waveguiding elements since the plasmonic host (a network of metallic wires) does not allow wave propagation within its volume. LH propagation inside the waveguides can be achieved by loading the (hollow) waveguide with a proper microwave LH metamaterial, e.g. a lattice of split-ring resonators and rods [24] or a lattice of metal-coated ferroelectric spheres [24, 25]. Of course, ordinary, right-handed (RH) propagation within the waveguide elements can be achieved by an empty waveguide or by one loaded with a plain dielectric material.

In this paper, we present a different design which is much easier to realize in the laboratory. Namely, we consider a square lattice wherein the sites are connected with TLs. This design has
two main advantages over the lattice of cavities within the artificial plasma described above and in [8]: (a) there is no need for an artificial plasmonic host, as in TL networks the wave propagation takes place only along the TLs, and (b) it is much more feasible to fabricate microwave LHTLs and RHTLs than to load a waveguide with a (miniaturized) LH metamaterial. We therefore pursue the design of a topological photonic lattice with a 2D network of TLs. The positive (negative) phases of the TB formalism of section 2 can be realized with RHTLs (LHTLs) as shown in figure 2(c) (figure 2(d)), in which case the topological bands lie in the GHz regime.

In order to simulate the FQHE for microwave photons, we need to implement a cavity QED scheme in this regime. This can be achieved by considering a superconducting circuit cavity QED system [26, 27] consisting of a Cooper-pair box (CPB) coupled to a TL resonator (see the equivalent circuit of figure 2(b)). The CPB operates as an artificial atom (a two-level system) [28, 29] and couples to the microwave photons of a superconducting TL resonator, which plays the role of an on-chip cavity reservoir. The microwave response of a superconducting circuit cavity QED system is described by the Jaynes–Cummings Hamiltonian [26, 27]

\[ H^{\text{JC}} = \hbar \omega_t (a_i^+ a_i + 1/2) - \frac{1}{2} (E_{\text{el}} \sigma_+^i + E_j \sigma_-^j) + \hbar g (a_i^+ \sigma_-^i + a_i \sigma_+^i), \]

where \( \omega_t = 1/\sqrt{LC} \) is the frequency of the superconducting resonator, \( a_i^+ (a_i) \) creates (annihilates) a microwave photon in the TL resonator (cavity), \( \sigma_+^i (\sigma_-^i) \) creates (annihilates) an excitation in the CPB. \( g \) is the coupling parameter between the CPB and the TL resonator, \( E_{\text{el}} \) is the electrostatic energy and \( E_j = E_{1,\text{max}} \cos(\pi \Phi_b) \) is the Josephson energy of the CPB. \( \Phi_b = \Phi / \Phi_0 \) is a flux bias applied by a coil to the CPB and controls the Josephson energy \( E_j \).

A superconducting circuit cavity QED system where microwave photons propagate in the lattice of figure 2(a) is described by a Jaynes–Cummings–Hubbard Hamiltonian of the form [30]

\[ H^{\text{ICH}} = \sum_i H_i^{\text{JC}} + H^{\text{TB}}, \]

where \( H^{\text{TB}} \) is the tight-binding form of the Hamiltonian of the microwave photons propagating within the checkerboard lattice of figure 2(a), i.e.

\[ H^{\text{TB}} = -t \sum_{(i,j)} \exp(i\phi_{ij}) (a_i^+ a_j + \text{H.c.}) - \sum_{(i,j)} t'_{ij} (a_i^+ a_j + \text{H.c.}) - t'' \sum_{\langle\langle i,j\rangle\rangle} (a_i^+ a_j + \text{H.c.}), \]

which is the direct-space representation of the Green’s tensor of equation (10). An important ingredient that gives rise to the FQHE is the presence of repulsive interactions among the microwave photons and this is inherently present in equation (13) as photon blockade [30]. The latter phenomenon has recently been observed experimentally in the GHz regime for microwave photons and this is inherently present in equation (13) as photon blockade [30]. The different FQHE phases can be calculated by direct diagonalization of the Hamiltonian of equation (13). We note that the proposed quantum simulator for the FQHE differs fundamentally from previous proposals [33, 34] since it essentially constitutes a passive design requiring no externally applied electric or magnetic fields.

Some typical values of the Hamiltonian of equation (13) are [26]: \( \omega_t = 38 \text{ GHz}, E_{1,\text{max}} = 8 \text{ GHz}, E_C = 5.2 \text{ GHz}, g \approx 0.314 \text{ GHz}. \) The latter parameter, \( g, \) is much larger than the loss rate of the TL resonator \( (\sim 0.005 \text{ GHz}) \) and the decoherence rate of the CPB \( (\sim 0.004 \text{ GHz}) \). The frequency \( \omega_t \) of the TL resonator should fall within the operating bandwidth of the LH- and RHTLs. In figures 2(c) and (d) we have considered the ideal TL, which has infinite bandwidth.

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However, actual LH- and RHTLs have very large bandwidth, which is a distinctive feature of nonresonant metamaterials compared to resonant \[16\]. Lastly, if the superconducting QED chip has a thickness of 1 mm, the coupling TLs (RH or LH) have 10 mm length and the TL resonator covers an area of 30 mm\(^2\), for the given resonator frequency (38 GHz), the hopping strength \(t\) is about 0.5 GHz, which is also significantly larger than both the TL loss and CPB decoherence rates. For the above TL length (10 mm) and the given operating frequency (\(\omega_r = 38\) GHz), the subwavelength condition of equation (4) is sufficiently met.

In order to probe experimentally the FQHE with the proposed structure, one needs to create the phase diagram of the spectrum gap between the FQHE ground-state manifold and the lowest excited states, as a function of the coupling parameters \(g\) for NN and NNN hopping when the latter lie in the photon blockade regime. Generally speaking, in the FQHE state the spectral gap assumes much larger values than in superfluid and solid phases \[18\]. The frequencies of the ground and excited states (and thus their corresponding gaps) can be measured by microwave transmission experiments.

4. Conclusion

In conclusion, we have shown that topological frequency bands emerge in 2D EM lattices of metamaterial components in the absence of an applied magnetic field. The topological nature of the corresponding band structures gives rise to important phenomena such as one-way waveguiding and coalescence of EM modes. The above lattices can be the basis for realizing a simulator for the FQHE based on superconducting TLs and circuit cavity QED systems.

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