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Published in:
Physical Review B

DOI:
10.1103/PhysRevB.99.115114

Citation for published version (APA):
Visualizing the connection between edge states and the mobility edge in adiabatic and nonadiabatic topological charge transport

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(Received 15 November 2018; revised manuscript received 26 February 2019; published 11 March 2019)

The ability to pump quantized amounts of charge is one of the hallmarks of topological materials. An archetypical example is Laughlin’s gauge argument for transporting an integer number of electrons between the edges of a quantum Hall cylinder upon insertion of a magnetic flux quantum. This is mathematically equivalent to the equally famous suggestion of Thouless that an integer number of electrons is pumped between two ends of a one-dimensional quantum wire upon sliding a charge-density wave over a single wavelength. We use the correspondence between these descriptions to visualize the detailed dynamics of the electron flow during a single pumping cycle, which is difficult to do directly in the quantum Hall setup because of the gauge freedom inherent in its description. We find a close correspondence between topological edge states and the mobility edges in charge-density wave, quantum Hall, and other topological systems. We illustrate this connection by describing an alternative, nonadiabatic mode of topological transport that displaces precisely the opposite amount of charge compared to the adiabatic pump. We discuss possible experimental realizations in the context of ultracold atoms and photonic waveguide experiments.

DOI: 10.1103/PhysRevB.99.115114

I. INTRODUCTION

The experimental discovery of the integer quantum Hall effect (IQHE) in GaAs [1] revealed the first example of a two-dimensional (2D) material that has an insulating bulk but metallic edge states. The IQHE state is now understood to be part of a much larger class of topological insulators [2–9], which can be labeled by topological invariants such as the Chern number or spin-Chern number [10]. Although these indices determine the number and types of states localized on the edges of a finite sample, they may be computed entirely from bulk quantities and depend on the symmetries of the bulk Hamiltonians. Different topological classes emerge depending on whether time-reversal, particle-hole, and chiral symmetries are present or not [11,12] and may be further refined using lattice symmetries [13–17].

A correspondence between the bulk topological invariants and the dynamics of edge states in the IQHE was put forward by Laughlin [10,18–20]. It shows that adiabatically threading a single Aharonov-Bohm flux quantum $\phi_0 = h/e$ through the interior of an IQHE cylinder [18,19] results in a quantized number of electrons moving from one edge of the cylinder to the opposite. The amount of charge transported is given precisely by the bulk Chern number, defined as an integral over the 2D Brillouin zone [21]. The resulting cyclic charge transfer may be regarded as a dynamical manifestation of the IQHE [22] and has been observed in a Corbino geometry [23,24]. More generally, it is an example of the type of quantized adiabatic particle transport or topological charge pumping first proposed by Thouless [10,21,25,26]. Experimentally, topological pumps have been realized in cold-atom [27,28] and single-spin [29] systems and have attracted attention in both the adiabatic [30,31] and nonadiabatic [32,33] regimes.

Despite its status as an archetype of topological transport, the detailed dynamics of exactly how electrons are transferred between edges as a quantum of flux is threaded through Laughlin’s IQHE cylinder remains difficult to visualize. The straightforward comparison of electronic wave functions for values of the flux that differ by fractions of the flux quantum is hampered by the flux dependence of the canonical momentum. In this paper, we circumvent this problem by using a well-known family of one-dimensional (1D) charge-ordered systems that can be mathematically mapped onto the 2D IQHE setup [34]. This allows us to directly visualize the topological transport in real space and time, as electrons flow from one edge of the system to the other, and thus to clarify the nature of the transport in both the charge-ordered and IQHE systems.

The rendering of edge state dynamics additionally reveals the in-gap states are connected to the so-called mobility edge, which is the critical state separating extended from localized states in the bulk of a disordered system [35,36]. In the case of quantum Hall systems, the bulk extended states are known to be constricted to just a single energy in the center of each (impurity-broadened) Landau level, as long as inter-Landau-level scattering is negligible [35,37]. Both edge states and the mobility edge are topological in nature [38], which we confirm by considering their common insensitivity to the presence of weak impurities. Together, the edge states and the mobility edges make up a single set of connected states winding around the electronic spectrum. We illustrate...
this connection of all the topological states by suggesting an alternative form of topological charge transport based on a nonadiabatic variation of externally applied fields. The nonadiabatic charge pump transfers precisely the opposite amount of charge from its traditional adiabatic counterpart and thus acts as an “anti-Thouless” pump. We suggest a possible experimental realization in cold atomic gases in an optical lattice.

II. MAPPING BETWEEN ONE AND TWO DIMENSIONS

Harper’s equation, and, more generally, a Mathieu equation [39–41], can be interpreted both as a tight-binding model for the IQHE in two dimensions and as a family of mean-field descriptions for monatomic commensurate charge-density waves (CDWs) in one dimension [34,42–45]. The filling of the charge-ordered system, and hence its CDW periodicity, for example, can be taken to correspond to the magnetic field strength perpendicular to the surface of a quantum Hall cylinder, which determines the filling of its Landau levels [34,41,43]. Under the same mapping, the phase \( \phi \) of the CDW order parameter then corresponds to a flux threading the quantum Hall cylinder [18], while the spatial coordinate of the CDW chain is directly related to the spatial coordinate parallel to the axis of the 2D cylinder. The mapping is indicated schematically in Fig. 1.

Upon varying the parameter \( \phi \) adiabatically from zero to \( 2\pi \), the CDW slides along the 1D chain over precisely one wavelength. The flux threading the IQHE system is increased by one quantum under the same variation. When introducing edges, we thus expect to find adiabatic transport of a quantized number of electrons from one side of the system to the other in both cases. In this way, the mapping relates the quantized adiabatic particle transport (Thouless pumping [21]) in charge-ordered systems to the topological transport between edge states upon insertion of a flux quantum in a quantum Hall cylinder (Laughlin’s gauge argument [18]).

The bulk of both the CDW and IQHE systems is insulating, and only the edge states cross the Fermi level. By solving Harper’s equation on a cylinder, the spectrum can be plotted as a function of momentum in the periodic direction. It consists of bulk bands and isolated topological edge states crossing the gaps between them. The edge states are protected in the sense that the number of edge channels cannot be modified as long as the bulk of the system remains gapped [18]. In a CDW system the periodic direction is given by the mean-field value of the phase variable \( \phi \), as shown in Fig. 1. Each value of the phase then corresponds to a single realization of a CDW on the chain, which may or may not host edge states. The combined spectrum of the family of CDWs containing all values of \( \phi \) coincides with that of the IQHE cylinder.

To be concrete, consider a CDW on a finite chain of \( N \) sites described by the Hamiltonian:

\[
H = -t \sum_{j=1}^{N-1} (c_j^\dagger c_{j+1} + c_{j+1}^\dagger c_j) + V \sum_{j=1}^{N-1} c_j^\dagger c_j^\dagger c_{j+1}^\dagger c_{j+1} + \sum_{j=1}^{N} \zeta_j c_j^\dagger c_j. \tag{1}
\]

Here, \( t > 0 \) is the tunneling amplitude, \( V \) is the strength of the nearest-neighbor Coulomb interaction between electrons, and \( \zeta_j \) is a random on-site potential that describes the effect of impurities. The operator \( c_j^\dagger (c_j) \) creates (annihilates) a spinless electron at position \( x = ja \), where \( a \) is the lattice constant. The filling factor \( n \) is a common fraction, so that \( Nn \) is the total number of electrons in the system. As usual, we assume the system is charge neutral in total and ignore any ionic charges.

Using the mean-field ansatz \( \langle c_{j+1}^\dagger c_j + c_{j-1}^\dagger c_{j+1} \rangle \approx 2(c_j^\dagger c_j) = A \cos(Qja + \phi) \) for the particle density, the interaction term becomes \( \sum_j VA \cos(Qja + \phi)c_j^\dagger c_j \) [34]. This ansatz defines the propagation vector \( Q = 2\pi n/a \) of the charge-density modulations, as well as the phase \( \phi \), which
determines the position of the CDW with respect to the lattice. Substituting the interaction back into Eq. (1) yields the full mean-field Hamiltonian for the 1D CDW chain. For periodic boundary conditions and $\zeta_j = 0$, this form of the Hamiltonian will coincide with Hofstadter’s tight-binding description of the IQHE [41] if we make the identifications $k \rightarrow k_x$, $Q \rightarrow eB/(\hbar c)$ and $\phi \rightarrow k_x a$. Here, $k$ is the momentum in the CDW, and $k_{x, y}$ are the $x$, $y$ components of momentum in the IQHE.

The radial magnetic field $B$ in the IQHE cylinder determines the filling of its Landau levels. Under this mapping, adding flux along the central axis of the IQHE cylinder corresponds to changing the value of the phase $\phi$ in the CDW [18] and thus to sliding the charge-density wave along the chain.

The Hamiltonian itself may be realized as a mean-field description of charge order in 1D chains of aligned orbitals within a three-dimensional material [46] or as an effective description of cold atoms in an optical lattice [34,47]. Here, we focus first on the topological properties of the theoretical mean-field model, before discussing possible experimental probes of the various emerging modes of topological transport in more realistic settings.

III. VISUALIZING CHARGE TRANSFER BETWEEN EDGE STATES

In Laughlin’s argument for quantized transport across the IQHE cylinder [18], the electromagnetic gauge structure plays a central role. Although helpful in establishing why the charge transferred between the ends of the cylinder must be an integer, the presence of gauge freedom makes it hard to directly visualize the precise dynamics. We know and understand what happens upon insertion of a single flux quantum, but questions like at which value of the flux one edge state becomes unoccupied and the opposite one becomes occupied and what the wave function looks like after insertion of only half a flux quantum are difficult to answer in a gauge-independent fashion. For the CDW system, this problem does not exist. The process corresponding to a flux insertion is the sliding of the CDW by precisely one wavelength, and this results in a precisely quantized amount of charge being transferred from one end of the chain to the other (Thouless pumping [21]). The quantized conductivity is determined by the sum of Chern numbers for all occupied bands (in $k$, $\phi$ space), as in the IQHE. However, in this case we can plot the charge distribution for both bulk and edge states for any value of $\phi$ and hence visualize the topological transport as a continuous process.

As shown in Fig. 1, the spectrum and eigenstates can be computed for a specific choice of parameter values, including any value for the phase $\phi$. As an example, we show the results for a period-3 CDW with $n = 1/3$, $N = 63$, and $AV/\tau = 5$. We use this specific example throughout the paper but note that taking other parameter choices does not qualitatively affect any of our results. Among the numerically obtained wave functions, bulk and edge states can easily be distinguished. The bulk wave functions for a CDW with open boundary conditions can be understood as products of plane waves (solutions of a CDW system with periodic boundary conditions) and the eigenstates of a particle in a box. The edge states, on the other hand, are exponentially localized on one side of the chain.

Upon varying $\phi$, the highest occupied state changes from an edge state localized on one side to that on the opposite side. In the example of 1/3 filling shown in Fig. 2, the highest state with $E < 0$ at $\phi = 0$ is an edge state localized on the right side of the chain ($j = 63$). In the region $\pi/3 < \phi < \pi$, the edge state is adiabatically transformed into a bulk state, until it emerges again as an edge state on the opposite side of the system. At $\phi \approx 5\pi/3$, the highest occupied state discontinuously changes from being localized on the left side of the chain ($j = 1$) to being localized on the right. From the spectrum, it is clear that this behavior stems from the two edge states crossing in energy at this point.

Notice that the states displayed in Fig. 2 are the highest occupied states at zero temperature for a given value of the phase variable $\phi$. Starting from a phase value $\phi < 5\pi/3$ and adiabatically sliding the CDW forward, the system will, in fact, not stay in the instantaneous ground state. The two edge states crossing within the bulk gap are located at opposite edges of the chain in real space, and any matrix element of local operators that could assist in tunneling across is exponentially small in the chain length. A state on one end of the system can therefore not simply jump to the other end in the way suggested by Fig. 2. For a sufficiently long CDW, adiabatic variation of $\phi$ causes the system to end up in an excited state at $\phi = 2\pi$, with a high-energy edge state occupied and a lower-energy edge state empty. The system can only return to its instantaneous ground state, and the topological material can only function as an adiabatic charge pump if the two edges of the CDW chain are connected to one another through some external coupling. In fact, in any experimental implementation of a topological quantum pump, one would indeed include a wire connecting the two sides of
the Thouless pump or, in the IQHE interpretation of the same
to the other, through the connecting wire, and thus realizes
φ
respectively. The leftmost plot shows the state which in (b) starts out as a right edge state (blue line) for low
j
bulk gaps are edge states localized at the left (j = 1) end of the chain, blue lines are right (j = 63) edge states, and purple lines in the middle
of bulk bands represent mobility edges. (b) Schematic close-up view of the crossing of two edge states and of an edge state entering the
band. If the two edges are connected, they undergo an avoided crossing, which can be traversed either adiabatically by tuning φ very
slowly or nonadiabatically by increasing φ more rapidly in the region of the avoided crossing. As the edge state enters the bulk band, it
necessarily undergoes a series of avoided crossings, which ultimately connect the edge state to the mobility edge. (c) Intensity plot of selected
wave function amplitudes |ψ(φ)|\(^2\) for different values of the phase φ. In all plots, dark blue and white indicate low and high amplitudes,
respectively. The leftmost plot shows the state which in (b) starts out as a right edge state (blue line) for low φ, then becomes a left edge state
(red line), and after the avoided crossing at φ \simeq \pi becomes a bulk state (gray line). The orange line in (c) indicates that separate color scales
were used for the edge and bulk states. The other two panels show the states that start out at low φ as the topmost state in the lower bulk band
and the one below it. The avoided crossings at φ \simeq \pi can be identified by the changing number of nodes in the wave function.

the system and would typically measure the current through
the wire as the phase is being varied. Such a connection
can easily be included in the simulation as a weak hopping
element between the opposing sides of the chain. It allows
the two edge states to interact and turns their intersection
in the spectrum into an avoided crossing by opening up
a small energy gap. Adiabatic evolution, which retains the
instantaneous ground state throughout a pumping cycle, is
once again possible. Varying φ slowly or nonadiabatically by increasing φ more rapidly in the region of the avoided crossing. As the edge state enters the bulk band, it
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were used for the edge and bulk states. The other two panels show the states that start out at low φ as the topmost state in the lower bulk band
and the one below it. The avoided crossings at φ \simeq \pi can be identified by the changing number of nodes in the wave function.

IV. THE CONNECTION OF EDGE STATES TO THE
MOBILITY EDGE

Inspecting the bulk wave functions in the spectrum of
Fig. 1, two features stand out. First of all, as the edge state
takes over its character. The existing bulk state takes over
as a right edge state at φ \simeq \pi/3, it goes through a second series of avoided crossing and reemerges
as a right edge state at φ \simeq \pi. The edge states are of the
form ψ(x) \propto e^{iλx}, with λ being negative and positive for edge
states localized at small and large x, respectively. The mobility
edge naturally connects these two states and has the same
shape of the wave function, with λ = 0. This special, totally
delocalized wave function is thus a plane wave connecting
both sides of the sample. Within the CDW, the plane wave is
again dressed by charge modulations, as shown in Fig. 4.
Note that for even Nn, it is possible to have two orthogonal,
delocalized wave functions. We consider chains with odd Nn
here, so that a single mobility edge connects to edge states in
both bulk gaps.

The topological nature of the mobility edge becomes ap-
parent upon adding a random impurity potential ζ_j. Notice
that in contrast to the IQHE, where the presence of weak
impurities is necessary to observe the quantization of the Hall
conductivity, in the CDW system topological transport may be
observed even with only the Coulomb interaction localizing
charges within the unit cell. Nevertheless, the presence of weak impurities (modeled here with |ζ_j|/t < 0.2) may be
considered, and their effect is much the same as in the IQHE,
localizing electronic states at specific locations in the chain.
The difference between the mobility edge and other bulk
states is now immediately obvious. As shown in Fig. 4, typical
bulk states are severely affected by the impurity potential, being amplified and suppressed at random locations. The mobility edge, on the other hand, retains a more or less constant amplitude along the entire length of the chain, as long as the impurity strength is weak. This clearly indicates the topological nature of the mobility edge and hence its connection with the topological edge states.

V. NONADIABATIC QUANTIZED PARTICLE TRANSPORT

The connection between edge states and the mobility edge may be illustrated by considering a nonadiabatic mode of quantized transport, which exists in addition to the well-known adiabatic Thouless pump. Starting, for example, with the (red) left edge state just below the (red) left edge state just below known adiabatic Thouless pump. Starting, for example, with quantized transport, which exists in addition to the well-may be illustrated by considering a nonadiabatic mode of transport will equal

\[ C \sin \phi = |\zeta_j|/t < 0.2 \]

\[ \approx \zeta_j/|t| \]

At this wavelength, one could use, for example, Sr atoms to load the lattice. Gaussian beams with a spacing of 0.7 μm and a short wavelength of 532 nm can be observed by injecting photons into the waveguide array. Waveguide arrays may directly simulate the CDW Hamiltonian [55]. Variations of the mean-field solutions even under nonadiabatic driving.

In such protocols, however, the highest occupied state is in nonadiabatic, localized bulk states for part of the pumping cycle. The purely nonadiabatic process suggested here, in which the highest occupied state jumps all avoided crossings that it encounters, is special because it depends entirely on the special nature and connectivity of the topological edge states and mobility edge states.

The nonadiabatic topological charge pump presented here is intended purely as an illustration of the connection between different types of topological states. Nevertheless, nonadiabatic or hybrid protocols may in some aspects be preferable to purely adiabatic ones in practical implementations of topological pumping. As was recently pointed out, the presence of a spectrum of unoccupied states imposes much more stringent conditions for achieving adiabaticity than just the timescale of the driving being longer than the inverse gap size [48–50]. Nonadiabatic driving across an avoided crossing, on the other hand, may be accomplished by sufficiently fast changes in the driving parameters, as long as the overlap between initial and target states is large. We confirmed that this is the case in the model CDW system by achieving an almost complete transfer of occupation across avoided crossings upon ramping up the driving speed. For any experimental implementation, a practical restriction on parameter values may arise from a preference for maintaining the validity of the mean-field solution even under nonadiabatic driving.

VI. A POSSIBLE IMPLEMENTATION

To realize the proposed illustration of a nonadiabatic pumping cycle in experiments, several techniques for imaging the real-space structure of CDW systems can be used. First of all, the static real-space structure of a CDW can be imaged by scanning tunneling microscopy (STM), which directly observes the amplitude and phase of charge-density modulations induced by a CDW [51–54], as well as excess charge at the end of a chain [55]. Straightforward experimental realizations of the dynamic quantized transport protocols, on the other hand, may use ultracold atoms in an optical lattice or photons in a waveguide array. Waveguide arrays may directly simulate the CDW Hamiltonian [55]. Variations of the mean-field CDW phase are then implemented by appropriate variations of the index of refraction. The quantized transport of charge can be observed by injecting photons into the waveguide at one edge and observing the intensity distribution in the waveguides after various propagation distances. The required values for experimental parameters are similar to those used in the literature [55] and can be realistically obtained with existing technology.

Using ultracold atoms, the mean-field Hamiltonian can be constructed by projecting a circular optical dipole lattice of 63 sites through a microscope objective [47,56]. This can be done, for example, using an objective with a numerical aperture of 0.8 [57] and a short wavelength of 532 nm [58], obtaining a waist of the projected Gaussian beams of the order of 0.4 μm. To ensure that the projected potential is attractive at this wavelength, one could use, for example, Sr atoms to load the lattice. Gaussian beams with a spacing of 0.7 μm then result in a potential that closely resembles a sine wave. The lattice depth will be approximately 0.3 times the depth
of one Gaussian beam. The lattice sites can be created by imaging a mask or a pattern from a digital mirror device (DMD) through the objective [47]. In order to confine the atoms in the direction orthogonal to the plane of the ring, an optical lattice can be applied in that direction.

Appropriate values of the model parameters, establishing a Mott-insulator regime with a weak residual tunneling, may be achieved, for example, by choosing a lattice depth of about eight recoil energies, or approximately 10 kHz. The mean-field energies are then $V = 270$ Hz and $t = 45$ Hz. A weak link may be introduced between two particular sites in the lattice to allow them to act as edges. The strength of the tunneling across the weak link can be adjusted arbitrarily far down from 45 Hz by increasing the spacing between the two selected sites.

To create and move the CDW, an additional lattice with larger spacing may be added [28]. The required specifics for this additional lattice are much more relaxed than the ones for the primary lattice and can be freely adjusted within a sizable range. The rotation of the secondary lattice could be achieved by imaging a DMD or a rotating mask onto the atoms. The visualization of the topological transport brings to the fore an explicit connection between edge states localized at the ends of the chain or cylinder and the mobility edges residing in the very center of their bulk states. To illustrate the connection between these two types of topological states, localized and extended, we formulate a purely nonadiabatic protocol that nonetheless results in quantized, topological transport of charge. The number of electrons pumped in a single cycle is precisely equal to that of the well-known adiabatic pump, but they flow in the opposite direction, creating a sort of anti-Thouless pump. We expect our conclusions, in particular the connection between edge states and the bulk mobility edge, to apply more generally to disordered topological systems. Any system with a nontrivial Chern or $Z_2$ invariant is guaranteed to have edge states, and similarly, mobility edges arise generically in models with disorder-induced localization. A simplified intuitive picture of the connection between the two types of states, suggested by the CDW system studied here, can then be drawn starting from an edge state localized on one side of the system. By evolving this state as a function of some system parameter, such as the CDW phase, it may be adiabatically connected to an edge state on the opposite side of the system. If the evolution is adiabatic throughout, however, a fully delocalized state must generically exist between the two edge-localized extremes.

Finally, we suggest experiments on ultracold atoms in an optical lattice, as well as in photonic waveguides, which can test the proposed connection between edge states and mobility edges, as well as the nonadiabatic pumping cycle, using realistic parameter values.

**ACKNOWLEDGMENTS**

This work is part of the Delta Institute for Theoretical Physics (DITP) consortium, a program of the Netherlands Organization for Scientific Research (NWO) that is funded by the Dutch Ministry of Education, Culture and Science (OCW). J.v.W. acknowledges support from a VIDI grant financed by NWO. This project has received funding from the European Research Council (ERC) under the European Union’s Seventh Framework Programme (FP7/2007-2013; Grant Agreement No. 615117 QuantStro).