Interplay of Chiral and Helical States in a Quantum Spin Hall Insulator Lateral Junction


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We study the electronic transport across an electrostatically-gated lateral junction in a HgTe quantum well, a canonical 2D topological insulator, with and without applied magnetic field. We control carrier density inside and outside a junction region independently and hence tune the number and nature of 1D edge modes propagating in each of those regions. Outside the bulk gap, magnetic field drives the system to the quantum Hall regime, and chiral states propagate at the edge. In this regime, we observe fractional plateaus which reflect the equilibration between 1D chiral modes across the junction. As carrier density approaches zero in the central region and at moderate fields, we observe oscillations in resistance that we attribute to Fabry-Perot interference in the helical states, enabled by the broken time reversal symmetry. At higher fields, those oscillations disappear, in agreement with the expected absence of helical states when band inversion is lifted.

Above a certain critical thickness, the 2-dimensional electron gas (2DEG) of a HgTe quantum well presents an inverted band structure characteristic of a 2D topological insulator (2D-TI) [1, 2]. At the edge of the topological insulator, quantum spin Hall (QSH) helical states propagate [3, 4]. When the Fermi level lies in the bulk gap of a 2D-TI, conduction is dominated by those edge states [5, 6] and is in principle protected by time-reversal symmetry (TRS) against single-electron backscattering processes. The application of magnetic field is expected to lift such protection. Nonetheless, band inversion and counterpropagating QSH-like edge states are predicted to persist up to a critical magnetic field $B_c$. Above this field, band inversion should disappear, leaving a 2D band structure identical to that of a topologically trivial [7] semiconductor in the Quantum Hall (QH) regime [5, 6, 8].

Previous experiments on HgTe quantum wells in this thickness range show that the resistance in the bulk gap increases in the presence of moderate magnetic fields [6, 9, 10] as predicted [8, 11], but our understanding of the evolution of edge conduction with magnetic field is incomplete. For the related problem of assigning quantum numbers to different chiral quantum Hall (QH) modes, a fruitful approach has been to study scattering between those modes by measuring transport through junctions between regions of different carrier density. This approach has been widely applied in GaAs quantum wells (for a review see Ref. [12]) and more recently in the Dirac 2DEG of graphene [13–16]. Thus, to characterize helical modes under broken TRS, studying their interplay with quantum Hall chiral states could be a promising strategy.

In this work, we explore electronic transmission across a lateral heterojunction fabricated on a HgTe quantum well with inverted band structure. Above a critical field, our results are consistent with expectations for equilibration of QH edge modes. Results are similar below the critical field for high carrier densities, but clearly differ when the junction is tuned through zero density. There, we first observe how the maximum of resistance associated with the bulk gap narrows and shifts towards lower values of carrier density. We find this to be a consequence of the remaining band inversion and the existence of helical edge states. Over the density regime corresponding to the peak shift, the resistance of our device presents oscillations which we attribute to Fabry-Perot interference of helical states enabled by the lifting of TRS protection.

Fig. 1(a) presents the geometry of our device. A Hall bar mesa is defined following the method described in Refs. [17, 18] on a HgTe quantum well with inverted band structure. The 2DEG is formed in a quantum well epitaxially grown over a conductive substrate [18, 19], allowing for the application of an overall back-gate voltage. A narrow top gate electrode is placed at the center of the device, defining two regions with separately-tunable density (inset of Fig. 1(c)). Central region refers to the area covered by the top-gate electrode and outer region to its surroundings. The carrier density in the outer region $n$ can be tuned by the applied back-gate voltage $V_{bg}$, while the density $n'$ in the central region depends on both back- and top-gate ($V_{tg}$) voltages. Both $n$ and $n'$ can be estimated using a simple capacitor model [18].
FIG. 1. (a) Electron micrograph of Hall bar device with narrow top gate. The geometry of the device is sketched in the inset of panel (c). The width ($W$) of the Hall bar is 10 $\mu$m and the physical length ($L$) of the top gate is 2 $\mu$m. Carrier density in the outer and central regions of the junction are denoted by $n$ and $n'$ respectively. (b) 4-terminal resistance ($R$) across top-gated region measured as a function of top gate and back gate voltages at zero applied magnetic field. The diagonal black line marks the estimated position of $V_{bg} = 0$. (c) Selected linecuts extracted from panel (b): Linecut 1 (blue) shows the evolution of $R$ as a function only of $V_{tg}$, that is, as a function of $n'$ while the outer region is kept at constant $n = 10^{13}$ cm$^{-2}$. Linecut 2 shows $R$ as function solely of $V_{bg}$, which changes the density of the whole device, for a value of $V_{bg} = V_{bg0}$ such that the device density is homogeneous, that is, $n' \simeq n$.

The evolution of the four-terminal resistance $R$ measured across the junction at 2.1 K (inset of Fig. 1(c)) as a function of both $V_{tg}$ and $V_{bg}$ at zero applied magnetic field is shown in Fig. 1(b). As previously reported [5], the resistance presents a finite maximum when the chemical potential lies in the bulk gap and conduction is dominated by the QSH edge states. As a function of spatially-uniform carrier density, four-terminal resistance (Curve 2 in Fig. 1(c)) shows a maximum value higher than $h/e^2$, the value associated with the ballistic Quantum Spin Hall regime. This is expected: the edge mean free path in similar heterostructures has been reported to be a few microns [5], substantially less than the edge length between contacts in the present geometry, so backscattering should result in increased resistance. In contrast, as a function only of density in the central region $n'$ (Curve 1 in Fig. 1(c)) the maximum resistance is lower than $h/e^2$, suggesting the presence of bulk conduction in parallel to the QSH edge states. A detailed look at the data (Fig. S2 of [18]) reveals oscillations in the resistance which likely arise from the Fabry-Perot like interference of bulk conduction paths (for a detailed analysis see [18]).

The locations of resistance maxima in the $(V_{tg}, V_{bg})$ parameter space (Fig. 1(a)) fall along two lines: a horizontal line around $V_{bg} = V_{bg0} = 0$ V representing zero density in the outer region, and a diagonal line representing zero density in the central region of the junction ($n' = 0$) (see [18]). The two lines define four quadrants of electron and hole densities in the central and outer regions of the junction, labeled in Fig. 1(b).

At finite fields, four-terminal resistance $R$ in the $n$-$n'$-$n$ quadrant shows a sharp tiled pattern of fractional resistance values ranging from 0 to $h/e^2$ (Fig. 2(a,b), $B = 3$ and 5 T respectively.) The Landauer-Büttiker formalism describes those fractional values as the result of full equilibration between co-propagating edge states in the junction (cf. [12, 14, 20].) In the unipolar regime and in a 4-terminal configuration, the predicted resistance across the junction is given by

$$ R = \frac{h}{e^2} \frac{N - N'}{NN'} \quad (1) $$

where $N$ and $N'$ are the number of quantum Hall states propagating in the outer and central regions of the junction respectively. The bottom row of tiles in Fig. 2(a,b) corresponds to $N = 1$, with $N'$ increasing from left to right. Our data show a good agreement with the fractional plateaus at $R = 0, 1/2, 2/3, 3/4,...$ expected for $N' = 1, 2, 3, 4,...$ ($N = 1$ linecut in Fig. 2(c)). Similarly, a second row of tiles appears for $N = 2$, and in the corresponding linecut of Fig. 2(a), plateaus of resistance can be observed near the expected values for $N = 2$ paired...
with a range of $N'$, although deviations from the ideal behavior are larger here than for $N = 1$. Similarly, plateaus are observed near but not exactly at the expected values for the $p-p'-p$ and the $n-p-n$ quadrants (see [18] for details.)

These results highlight the role of the strong spin-orbit interaction and inversion symmetry breaking in HgTe. If the $s_z$ component of spin were conserved, transmission would then be spin-selective and only those states with the same spin polarization would equilibrate with each other, leading to fewer plateaus in resistance $R$. In contrast, our data suggest that full equilibration occurs for all possible values of $N$ and $N'$. Inversion symmetry breaking provides a mechanism spin mixing that allows this to happen [21]. Adapting the theory in Ref. 22, we estimate that the equilibration length for our system is around 2 μm at 2 T, which is indeed smaller than the junction width.

In the $n-n'-n$ quartet, the tiled structure of fractional resistance values associated to a given pair of values $(N,N')$ is similar at both 3 and 5 T (Fig. 2(a) and (b) respectively). A contrasting behavior emerges around zero density. First, the zero density resistance values associated to a given pair of values $(n,n')$ and $(N,N')$ are very different, with a range of $N'$, although deviations from the ideal behavior are larger here than for $N = 1$. Similarly, plateaus are observed near but not exactly at the expected values for the $p-p'-p$ and the $n-p-n$ quadrants (see [18] for details.)

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protected in the presence of extra symmetries such as mirror symmetry. In the experiment, such symmetries are absent, so there should always be a small minigap. (Figs. 4a-c). The location of this edge minigap within the bulk gap depends on details such as the potential at the edge. Therefore the edge and bulk charge neutrality points do not necessarily occur for the same value of gate voltage, and accordingly the center of the resistance maximum originating from the minigap at finite field does not necessarily align with the center of the bulk gap. The observed gate voltage position shift and narrowing of the resistance maximum at 3 T compared to zero or 5 T (see Fig. 2(d)) is consistent with an origin in the edge state minigap.

In the 1-0-1 configuration, when the incoming chiral edge mode from the outer region reaches the junction it can scatter into two possible outgoing modes: the copropagating helical edge mode or the chiral mode parallel to the junction. When the chemical potential in the central region is very close to the bottom of the lowest Landau Level, the incoming edge mode is almost perfectly matched to the copropagating helical one, while the counterpropagating edge mode forms a loop spanning the whole junction (Fig. 4(a)). This must be so because the counterpropagating mode has smaller momentum and therefore is located farther from the edge. Transport in this scenario is almost equivalent to the 1-1-1 situation, seen in the experimental data as an extension of the whole junction (Fig. 4(a)). This must be so because the counterpropagating edge mode forms a loop spanning the whole junction (Fig. 4(a)). This must be so because the counterpropagating edge mode forms a loop spanning the whole junction (Fig. 4(a)).

As the chemical potential approaches the crossing of the helical edge modes, the chiral mode connects to the one parallel to the junction, while the helical modes form a loop at either edge (Fig. 4(b)). These loops should disappear at the minigap, and reappear with opposite orientation below it (Fig. 4(c)). The existence of these loops, allowed because the protection from backscattering is lifted by \( B \), implies that coherent transport should be affected by multiple reflections at the interfaces. This effect should manifest in Fabry-Perot type oscillations as a function of chemical potential, because the accumulated phase \( \delta = kL \) depends smoothly on chemical potential. This explains the oscillations observed at 3 T in Fig. 2(d), in the density range assigned to the bulk gap and adjacent to the edge minigap, and their disappearance beyond \( B \), (i.e. at 5 T) where no helical modes exist.

Resistance oscillations are also present at zero field in a similar density regime – these we associate with bulk states. However, at 3T, oscillations present an amplitude about an order of magnitude higher than their zero field counterparts. Moreover, our data (Fig. 2(c) and Fig. S4(d)) indicate that the bulk states causing the oscillations at zero field must be fully localized at 3T (see [18]), further suggesting that oscillations at 3T have their origin in edge rather than bulk states.

Furthermore, oscillations are periodic in \( n' \), with no substantial dependence on \( n \) (Fig. 4(d)). This is consistent with our interpretation: changing \( n' \) will change the edge momentum of the loop modes and hence the phase, while \( n \) only determines the momentum of the incoming modes, which should have no effect on the phase of the oscillations. Fig. 4(d) also shows that oscillations are present for \( n \) values corresponding to \( N=1 \), but they disappear when approaching \( N=2 \). This is consistent with the \( N=2 \) Quantum Hall chiral edge modes not being fully established at \( B = 3 \) T, as already seen in the imperfectly-quantized equilibration plateaus in Fig. 2(c).

Taken together, the amplitude, position and field- and carrier density-dependence of resistance oscillations support our interpretation of their origin in the constructive interference of helical edge states.

Finally, assuming a Fabry-Perot scenario yields some quantitative estimates for system parameters. Based on bulk 2D Fabry-Perot oscillations at \( B = 0 \) we estimate the effective length \( L^* \) of the central region to be 0.6 \( \mu \)m (see section 4 of [18] for details). This length need not be the same as the physical top-gate length, due to the smooth shape of the gate-induced potential. Given that we observe coherent FP oscillations from the QSH edge states as well, a lower bound on the edge localization length can be set at \( L^* \). We expect that at larger channel lengths or with higher disorder, the QSH loop responsible for interference will break up into more loops and coherence will gradually be lost, in a way similar to Ref. 11. We may also estimate the 1D edge mode carrier density: given the above value of \( L^* \) we infer \( C_{1D} = \delta(n_{1D})/\delta(V_{bg}) \approx 1.4 \times 10^6 \text{ cm}^{-1} \text{V}^{-1} \) (see section 6 of [18] for details).

While our results below critical field are compatible with those in Ref. [23], we present here evidences for physical scenarios that were not accessible in that work. More specifically, the dual-gate configuration of our device allows us to perform a detailed study of the equilibration of QH states in HgTe QWs and to infer the role.
played here by spin-orbit interaction. More importantly, we present one of very few evidences for a transition between QSH and QH regimes in a 2D-TI. Finally, we observe signatures of coherent interference on helical states, likely due to the geometry of our junction. Our results suggest that valuable information about the QSH state under broken TRS can be inferred from the electronic transmission across a QH-QSH-QH heterojunction.

While the present work was under review, a related theoretical work by Nanclares et al. [24] has been published.

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[7] The electronic structure of a 2DEG in the Quantum Hall regime is also topologically nontrivial. Here, however, we use trivial to describe only the lowest energy bands characterized by the Z2 topological invariant.
[18] See Supplemental Material at [URL will be inserted by publisher], which contains Refs. 25–39, for details on device fabrication, theoretical calculations, and further data at zero and finite fields.

