Topological protection brought to light by the time-reversal symmetry breaking

S.U. Piatrusha,1, 2 E.S. Tikhonov,1, 2 Z.D. Kvon,3, 4 N.N. Mikhailov,3, 4 S.A. Dvoretsky,3 and V.S. Khrapai1, 2

1Institute of Solid State Physics, Russian Academy of Sciences, 142432 Chernogolovka, Russian Federation
2Moscow Institute of Physics and Technology, Dolgoprudny, 141700 Russian Federation
3Institute of Semiconductor Physics, Novosibirsk 630090, Russian Federation
4Novosibirsk State University, Novosibirsk 630090, Russian Federation

Recent topological band theory distinguishes electronic band insulators with respect to various symmetries and topological invariants, most commonly, the time reversal symmetry and the $Z_2$ invariant. The interface of two topologically distinct insulators hosts a unique class of electronic states – the helical states, which shortcut the gapped bulk and exhibit spin-momentum locking. The magic number of defects, the free propagation of a wave turns into a diffusion. Given the phase coherence is preserved, the constructive interference between the time-reversed random diffusion paths gives rise to a coherent backscattering of a wave. This genuine quantum effect is observed as a narrow resonant in the intensity of light backscattered off a milky solution [21] and as a weak localization correction to the conductance of a diffusive metal [3]. Eventually, in a sufficiently disordered system, the coherent propagation gets suppressed, be it light [6, 7], electron [8], or even sound [9] or matter waves [10, 11] – the phenomenon known as Anderson localization [12, 13]. For the helical states, however, the situation inverts thanks to a destructive interference of the time-reversed paths. Thereby the free propagation is maintained in the presence of a symmetry-conserving disorder, which is known as topological protection [14].

Helical edge states represent a unique example of a 1D electronic system, that can only be realized at the interface of a two-dimensional (2D) $Z_2$ topological and trivial insulators [15]. In this work – the quantum spin-Hall (QSHI) insulator in HgTe/CdHgTe quantum wells (QWs) with the inverted band structure and vacuum [16]. Spin-momentum locking is manifested in two counter-propagating opposite-spin species at each edge of a planar device, which provide the only transport channel when the Fermi level is tuned within the 2D bulk energy gap [17]. Numerous direct consequences of this physical picture are corroborated experimentally, including the observation of quantized conductance $G \approx G_0 = e^2/h$ of the shortest edge channels [17, 18], non-local transport in zero magnetic field [18, 21], positive magnetoresistance [17, 21], the spin-charge sensitivity [25] and the unconventional behavior in lateral p-n junctions [25, 26].

In spite of the impressive progress, the mean-free path of the helical electrons is typically disappointingly small [17, 22, 27, 30], even compared to the conventional high-purity 1D conductors [31]. The puzzles of the trivial ohmic behavior and weak or even absent temperature dependence [17, 18, 27, 30, 32, 33] along with the nearly universal partition noise [33, 34] further indicate that the edge transport beyond the mean-free path is classical, rather than quantum coherent, by nature. As a matter of fact, the advantages offered by the concept of topological protection, which are of paramount importance for numerous applications [38, 37], were so far hidden by an extremely efficient phase breaking mechanism of a debated origin [38, 44]. Here, we approach this problem from a different perspective, using coherent backscattering as a marker for a breakdown of the topological protection. We observe that the time-reversal symmetry breaking by magnetic field restores the coherent backscattering and drives the Anderson localization of the helical edge channels. This behavior is in stark contrast with the zero magnetic field scenario, unveiling the actual strength of the topological protection in the time-reversal symmetric case.

We investigate two QSHI devices of different crystallographic orientation and QW thickness, $d$. In device D1, the QW with $d = 8.3\text{ nm}$ resides in a (013) plane, while the device D2 is based on the (112) QW with $d = 14\text{ nm}$. We note that the QWs of both devices are similar in design to the QWs studied in [18, 24].

* e-mail: tikhonov@issp.ac.ru
Both devices are shaped as multi-terminal Hall bars with Ti/Au metallic top gates, see insets of Fig. 1a and Fig. 1b for the schematic representation of the device and measurement configuration. The microscope image of one of the devices may be found in Supplemental Material Fig. 1. For further device fabrication details see methods section. Using the gate voltage $V_g$, the Fermi level can be tuned within the bulk energy gap, as large as 30 meV in D1 and 3 meV in D2 [16] [18]. In this way the QSHI regime is realized, with the predominant edge conduction confirmed by transport measurements in similar devices [18] [23] [27] [33] and, independently, here via nonlocal resistance measurements (see Supplemental Material Fig. 2). Various distances between the neighboring ohmic contacts allow us to choose the different lengths of the edge channels, spanning the range between 2 $\mu$m and 38 $\mu$m in each device. The device D1 demonstrates resistance $R \approx R_0 \equiv \hbar/e^2$ for the shortest edges, while in D2 the resistance is about twice as large for the same edge length (see Supplemental Material Figs. 5, 7).

Fig. 1a shows the two-terminal resistance of the 38 $\mu$m-long edge in D1 as a function of the gate voltage, measured at $T = 800$ mK and $T = 50$ mK, with and without magnetic field $B_\perp = 50$ mT. By contrast, in a small magnetic field of $B_\perp = 50$ mT the resistance increases dramatically up to $R \sim 1$ G$\Omega$ at $T = 50$ mK and drops down again by more than a factor of 10$^3$ as the temperature is raised to $T = 800$ mK. Similar yet less pronounced effect of the magnetic field is observed in device D2, as shown in Fig. 1b for the 12 $\mu$m-long edge. Here the edge conduction dominates in the range $-2.5 \, V < V_g < -3.5 \, V$ and the resistance increases by almost two orders of magnitude for $B_\perp = 200$ mT. The straightforward crosscheck demonstrates that in all our measurements the current flows along the edges of the device, while the bulk conduction contribution remains negligible even for $R \sim 1$ G$\Omega$ (see Supplemental Material Fig. 3). All the edges of both our devices D1 and D2 exhibit the reported resistance increase in a small magnetic field, including those with $R \approx R_0$ (see Supplemental Material Figs. 5–7 for additional data).

Figs. 1a and 1b highlight our main result that a tiny magnetic field gives rise to the dramatic increase of the resistance of the helical edge states accompanied by strong $T$-dependence and giant mesoscopic fluctuations. Altogether, this behavior is a hallmark of the Anderson localization of the electronic states and manifests a transition from the topologically protected phase to the trivial insulator in a magnetic field as anticipated in various scenarios [45–47]. The underlying microscopic explanation is depicted in Fig. 1c. In $B = 0$, the dispersion relation of the helical electrons consists of two opposite spin counter-propagating branches, the coherent backscatter-
FIG. 2. Temperature and magnetic field dependence of the edge conductance in the localized regime. Log-averaged conductance $G_{\text{typ}}$ (see text) of the 38 $\mu$m-long edge in D1, as a function of inverse temperature in (a) $B_{\perp}$ and (b) $B_{\parallel}$, oriented at 45° to the edge. Application of magnetic field leads to the activation-like $G_{\text{typ}}(T)$, in contrast to the weak metallic $G_{\text{typ}}(T)$ in $B = 0$. The strong anisotropy with respect to the $B$-orientation is evident. (c) Activation energy $\Delta$ (see text) as a function of $B_{\perp}$, extracted from measurements in (a) and the similar ones for the 6 $\mu$m-long edge. The dashed lines are $\Delta = g \mu_B B_{\perp}$ with the specified values of $g$. (d) The same as (c), but for the 38 $\mu$m-long edge in $B_{\parallel}$, extracted from measurements in (b).

We now discuss in detail the $T$-dependence of the edge conductance $G = 1/R$ for the 38 $\mu$m-long edge in D1. As seen from the red and green lines in Fig. 1a, in the absence of magnetic field the $T$-dependence within the charge neutrality point (CNP) region may be of both metallic and insulating type. Within a range of gate voltages $-7.4 < V_g < -8.3$ we observe a weak metallic $T$-dependence, with $G$ increasing by as much as a factor of 2 as $T$ is reduced from 800 mK to 50 mK. We are not aware of similar observations in HgTe QWs, where the reported $T$-dependencies are usually either completely absent or weakly insulating. We note, however, that most of the studies discuss $R(V_g)$-dependencies at temperatures above 1 K. Additionally, we note that in D2 the $R(T)$-dependence in $B = 0$ is of weakly insulating type at any $V_g$ (see Supplemental Material Fig. 4).

The $G(T)$-dependence in a magnetic field is much more impressive. In the presence of strong fluctuations, we analyze the log-averaged (typical) conductance $G_{\text{typ}} \equiv G_0 \exp(\ln G/G_0)$, with the averaging performed over the small gate-voltage region $-7.65 \text{ V} < V_g < -7.55 \text{ V}$ within the resistance maximum in Fig. 2a. The resulting data is shown in Fig. 2a. Strikingly, already for $B_{\perp} \geq 1 \text{ mT}$ the trend of $G_{\text{typ}}(T)$ changes from metallic to activated insulating dependence $G_{\text{typ}} \propto \exp(-\Delta/k_B T)$ (the remnant $B$-field did not exceed 2 mT and was compensated in the experiment with 0.1 mT precision). The activation energy reaches about $\Delta \approx 25 \mu$eV in 10 mT, shows a sub-linear increase with $B_{\perp}$ and increases with the length of the edge, Fig. 2b). Thus, it is difficult to make an obvious relation of the activated behavior with the single-particle spectrum of the helical edge states, e.g. with a Zeeman gap opening at the Dirac point. Similar observations hold for the in-plane orientation of the magnetic field, see Fig. 2b for the case of $B_{\parallel}$ directed at about 45° with respect to the edge under study. Here, the magnetic fields $B_{\parallel}$ roughly an order of magnitude stronger are required to observe the activated behavior comparable to the $B_{\perp}$ case. This might be a consequence of the Lande $g$-factor anisotropy predicted for HgTe QWs in some works.

We now quantitatively analyze the observed giant conductance fluctuations, which are the distinctive feature of the Anderson localized phase. Here, in contrast to the metallic phase, the conductance is exponentially sensitive to the minor variations of disorder, or equivalently, to the Fermi energy. As a result, the fluctuations of the conductance are as large as the average value and obey the log-normal distribution, i.e. it is the quantity $\ln G/G_0$ which is Gaussian-distributed in $B_{\perp}$ case. In Fig. 2a we study the conductance fluctuations of a 38 $\mu$m-long edge in D1 as a function of magnetic field at various $T$. The normalized variance of the logarithm of the conductance, given by $-\text{Var}(\ln G/G_0)\ln(G/G_0)^{-1}$, is plotted in panels (a) and (b) for $B_{\perp}$ and $B_{\parallel}$, respectively. At the lowest available $T = 50 \text{ mK}$, the normalized variance of the logarithm of the conductance increases by more than one order of magnitude, compared to the $B = 0$ case, for $B_{\perp} \approx 20 \text{ mT}$ and $B_{\parallel} \approx 500 \text{ mT}$. Still, the observed values are substantially below the theoretical value of 2, which would correspond to a quantum-coherent Anderson localized phase. We attribute the difference to the impact of the averaging in the presence of dephasing, which...
also qualitatively explains the strong $T$-dependence of the fluctuations in Fig. 3.

To further characterize the transport properties of the localized states, we study the non-linear transport regime originating from the delocalization of the edge states by the electric field $\mathbf{E}$ for the 38$\mu$m-long edge in D1 at $T = 50$ mK. The dependence of $I$ on $V_{sd}$ and $V_g$ is plotted in panels (a) and (b) of Fig. 4 for $B_\perp = 100$ mT and $B_{||} = 100$ mT, respectively. Within the CNP region one can see the dramatic changes in $I$ at increasing $|V_{sd}|$ (note the log-scale). Below the certain threshold value of $V_{sd}$, which depends on $V_g$ with pronounced reproducible fluctuations above that, we observe only the negligible current $|I| < 1$ pA through the edge. For the magnetic field of 100 mT, the typical bias range of suppressed conduction changes from about 0.5 mV in $B_\perp$ to about 50 $\mu$V in $B_{||}$. For this threshold, the corresponding energy scale is considerably higher than the activation energy extracted from the $T$-dependencies similar to that of Fig. 2, indicating that the applied bias is shared among a few strongly localized electronic states along the edge. Above the threshold, the conduction reasonably comparable to the $B = 0$ case is restored. Two representative $I-V_{sd}$ cuts of Figs. 4a and 4b at $V_g = -7.3$ V are detailed in the corresponding insets, along with the $I-V_{sd}$ curves at $V_g = 0$. The observed highly non-linear transport behavior is yet another evidence of the Anderson localization of the helical states driven by the $B$-field and contrasts with the almost linear current-voltage response in the time-reversal symmetric $B = 0$ case (see Supplemental Material Figs. 3, 5–7).

In conclusion, through the low-temperature magneto-resistance measurements we were able, for the first time, to directly demonstrate the actual strength of topological protection in one-dimensional helical edge states of HgTe-based topological insulators. Breaking the time-reversal symmetry with an external magnetic field allowed us to expose the hallmark Anderson localization features for the edge states: the exponential $T$-dependence, giant reproducible mesoscopic fluctuations and the gap-opening-like features in the $I-V_{sd}$ characteristics. Concerning the admittedly hard task of achieving the conductance quantum for the edges longer than several micrometers, our observations reveal how topological protection in the time-reversal symmetric case still almost perfectly sustains the edge transport from localization.

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FIG. 4. Gap opening in the edge by external magnetic field. (a) Current, flowing through the 38 µm-long edge of D1 as a function of gate and bias voltages in $B_{\perp} = 100$ mT. Measured current is plotted as a logarithm of its absolute value (see text). The inset demonstrates $I - V_{sd}$ curves, measured at $V_{g} = -7.3$ V in $B_{\perp} = 0$ (red line) and 100 mT (blue line). (b) The same as (a) for $B_{\parallel} = 100$ mT, oriented at 45° to the edge.

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Supplemental Material

Device fabrication and measurement techniques

The single HgTe QWs as heterojunction Cd$_{0.7}$Hg$_{0.3}$Te/HgTe/Cd$_{0.7}$Hg$_{0.3}$Te were grown by molecular beam epitaxy on SI GaAs substrates with buffer ZnTe and CdTe layers CdTe with the thickness of 0.1 $\mu$m and 5-7 $\mu$m, respectively. The mesa was fabricated via plasma chemical etching followed by covering with SiO$_2$/Si$_3$N$_4$ insulating layers, 200 nm thick in total. All devices are equipped with evaporated Au/Ti metallic gates.

All devices were measured in a BlueFors-LD250 dilution refrigerator with a base temperature of 17 mK equipped with a 9 T superconducting solenoid. The lowest obtainable electronic temperature was $T \approx 50$ mK, verified by the Johnson-Nyquist thermometry. The two-terminal linear response resistances were obtained by differentiating of the $I$-$V$ curves, measured via the transimpedance amplifier with 1 G$\Omega$ coefficient.

Supplemental Material Fig. 1. Optical microphotograph of the 14 nm HgTe/CdHgTe device D2. The regions marked with white dots correspond to the contact leads with the hall bar in between, while in the other regions the QW is completely etched. The central hall-bar region is covered with the Ti/Au top gate (yellow).

Supplemental Material Fig. 2. Non-local resistance measurements in the devices D1 and D2. Non-local resistance as a function of gate voltage, measured via conventional lock-in technique in the configurations indicated in the insets for devices (a) D1 and (b) D2. The data was taken at $T = 50$ mK.
Supplemental Material Fig. 3. **Edge transport in localized mode in D1.** (a) Two-terminal conductance measurements in $B_\perp = 50 \text{ mT}$ at $T = 50 \text{ mK}$ with two different contact configurations (see inset schematic). Bias voltage $V_{sd}$ is applied to the contact C1, where the total flowing current $I$ is measured. The solid blue line corresponds to the configuration, when only C2, positioned at 38 $\mu$m from C1, is grounded. The solid orange line corresponds to the situation, when C3 is additionally grounded. The dashed green curve is the repeated measurement of the blue curve directly after the orange one. The insignificant difference between all three curves verifies the edge transport domination. (b) Two-terminal conductance measurements in different magnetic fields for the configuration indicated in the inset. (c) Non-local voltage $V_{nl}$, measured in different magnetic fields for the configuration indicated in the inset. The presence of finite slope of $V_{nl}$ inside the transport gap indicates the edge transport presence.

Supplemental Material Fig. 4. **Temperature dependence of the 6 $\mu$m-long edge resistance in device D2 in zero magnetic field.** (a) Edge resistance as a function of the gate voltage $V_g$ for four $T$ values. The measurements were performed in the configuration of the inset. (b) Conductance of the same edge, as in (a) versus $1/T$ for fixed gate voltages (see legend).
Supplemental Material Fig. 5. **The effect of perpendicular magnetic field in the 8.3 nm QW.** (a) The two-terminal $R(V_g)$-dependence of a 6 $\mu$m-long edge in D1 at various $B_\perp$, and (b,c) the gap-opening in the same edge at $B_\perp = 10$ mT, all measured at $T = 50$ mK. The insets demonstrate the corresponding $I$-$V_{sd}$ curves at the specified $V_g$ values. (d-f) The same measurements for the 38 $\mu$m-long edge. Note that the data was taken in another cooling compared to Fig. 1a hence a horizontal shift of $R(V_g)$-dependence.
Supplemental Material Fig. 6. **The effect of in-plane magnetic field in the 8.3 nm QW.** (a) The two-terminal $R(V_g)$-dependence of a 2 μm-long edge in D1 at various $B_∥$ and (b,c) the gap-opening in the same edge at $B_∥ = 500$ mT, all measured at $T = 50$ mK. The insets demonstrate the corresponding $I$-$V_{sd}$ curves at the specified $V_g$ values. (d-f) The same measurements for the 38 μm-long edge.
Supplemental Material Fig. 7. The effect of perpendicular magnetic field in the 14 nm QW. (a) The two-terminal $R(V_g)$-dependence of a 6 \( \mu \)m-long edge in D2 at various $B_z$, and (b,c) the gap-opening in the same edge at $B = 200 \text{ mT}$, all measured at $T = 50 \text{ mK}$. The insets demonstrate the corresponding $I$-$V_{sd}$ values. (d-f) The same measurements for the 12 \( \mu \)m-long edge. (g-i) Similar measurements for the 38 \( \mu \)m-long edge with panels (h,i) demonstrating the change in the gap with increasing magnetic field.