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Electrostatic Coupling between Two Surfaces of a Topological Insulator Nanodevice

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We report on electronic transport measurements of dual-gated nanodevices of the low-carrier density topological insulator (TI) Bi1.5Sb0.3Te1.3Se1.3. In all devices, the upper and lower surface states are independently tunable to the Dirac point by the top and bottom gate electrodes. In thin devices, electric fields are found to penetrate through the bulk, indicating finite capacitative coupling between the surface states. A charging model allows us to use the penetrating electric field as a measurement of the intersurface capacitance $C_{TI}$ and the surface state energy-density relationship $\mu(n)$, which is found to be consistent with independent angle-resolved photoemission spectroscopy measurements. At high magnetic fields, increased field penetration through the surface states is observed, strongly suggestive of the opening of a surface state band gap due to broken time-reversal symmetry.

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Three dimensional topological insulators (3D TIs) have been undergoing intense theoretical and experimental research on the properties of their unique surface states [1,2]. The presence of bulk carriers has hampered experimental progress, so a variety of crystal growth [3–8] and in situ charge displacement techniques [9–13] have been applied to suppress bulk conductivity. For example, quaternary TI materials of the form Bi$_{2-x}$Sb$_{x}$Te$_{y}$Se$_{1-y}$ have a significantly suppressed bulk contribution to transport, reaching large bulk resistivities and insulating-like temperature dependence [6,14,15]. Furthermore, exfoliation or growth of thin crystals has been used to achieve surface-dominated transport [7,8,12,16,17]. However, amid the extensive experimental effort on TI device transport, there is no study reporting independent control over the densities of both the upper and lower surface states in a single TI device. A full understanding of transport phenomena in TIs, such as the quantum Hall [18,19] and Josephson effects [20–22], will require independent tuning of the density of each surface state. Additionally, proposals for topological exciton condensates explicitly require fine tuning the density of both surfaces [23], and finite displacement fields from two gates can affect the quantum anomalous Hall effect in TI-based systems [24,25].

In this Letter, we report electronic transport measurements of exfoliated Bi$_{1.5}$Sb$_{0.5}$Te$_{1.7}$Se$_{1.3}$ (BSTS) nanodevices with top and bottom gate electrodes. We show, for the first time, that the chemical potential of the upper and lower surface states can be controlled independently, resulting in different resistance peaks when either surface chemical potential crosses the Dirac point. For thin devices, we find signatures of finite capacitative coupling between the surface states, consistent with fully depleted bulk states. We explain the data through a charging model which incorporates the finite density of states of the surface bands. Using angle-resolved photoemission spectroscopy (ARPES) as a control measurement of the surface state, this model allows us to measure the chemical potential $\mu$ and charge density $n$ of a topological surface state as well as the intersurface capacitance $C_{TI}$. At high magnetic fields, increased field penetration through the surface states is observed, strongly suggestive of the opening of a surface state band gap.

BSTS was prepared by melting high purity samples of the constituent elements in a sealed quartz ampoule under inert atmosphere. Sample structure was confirmed by x-ray powder diffraction, and large single crystals showed similar bulk transport behavior to previous reports [6]. Static ARPES shows that the chemical potential is inside the bulk band gap and that the Dirac point energy is above the bulk valence band edge (see Supplemental Material [26]). Pump-probe time-resolved ARPES (TRARPES) allows access to unoccupied states as shown in Fig. 1(b) [27,28]. The Fermi velocity near the Dirac point is $v_F \approx 3.2 \times 10^5$ m/s, and the band gap at room temperature is $E_g \approx 240$ meV. Note that the surface state dispersion is strongly electron-hole asymmetric. These data are consistent with previous experiments [14,29].

Thin flakes for transport studies were obtained by mechanical exfoliation onto a doped silicon wafer with a
easily tune the device through a resistance peak (the 2D resistivity of the BSTS and h-BN layers. For all data presented here, microscopy (AFM) was used to determine the thickness of Ohmic contacts and top gate electrodes. Atomic force microscopy (AFM) was used to determine the thickness of the BSTS and h-BN layers. For all data presented here, a four-probe voltage measurement was used to determine the 2D resistivity $R_{SQ}$. Here, we report results measured on BSTS devices of different thicknesses: device $A$ is 42 nm, and device $B$ is 82 nm. The behavior of device $A$ was reproduced in a third device [26]. All three devices were fabricated from flakes from the same exfoliation and, therefore, from the same region of the bulk crystal. Figure 1(a), shows an AFM image of device $A$.

On devices $A$ and $B$, both the top and bottom gates easily tune the device through a resistance peak ($R_{peak}$) by adjusting the applied voltages $V_T$ and $V_B$, respectively, as shown in Figs. 2(a) and 2(c). $R_{peak}$ is associated with a minimum in carrier density (i.e., the surface Dirac point), as confirmed via the Hall effect [26]. Interestingly, the top-gate $R_{peak}$ is observable up to room temperature; in contrast, for the bottom gate, $R(V_B)$ changes into a broad $S$ shape, consistent with gating studies of other TIs using SiO$_2$ gate dielectrics [7,10,17]. The disappearance of a distinct resistance peak in the limit of strong disorder was predicted by recent theories for TI surface states with electron-hole asymmetry [31], suggesting that the difference in the field-effect behavior may be related to the disorder profile at the interface. Strong differences in the disorder profile at SiO$_2$ and h-BN interfaces have been observed in graphene [32].

Two-dimensional maps of the resistivity with respect to both top and bottom gate voltage reveal a distinct difference in the behavior of devices $A$ and $B$, shown in Figs. 2(d) and 2(c), respectively. The black dots identify $V_{peak}(V_B)$, the top gate voltage at which $R_{peak}$ is found, at each $V_B$. We associate $V_{peak}$ with charge neutrality of the upper surface state: $n_U = 0$. For device $B$, $V_{peak}$ is independent of $V_B$, demonstrating no capacitive coupling between the upper surface and the bottom gate electrode. The fact that thicker devices do not have this capacitive coupling suggests that mobile bulk electronic states exist in the interior. By contrast, $V_{peak}$ in device $A$ is dependent on $V_B$. The observed relationship $V_{peak}(V_B)$ means that there exists a finite and non-constant capacitive coupling between the upper surface and the bottom gate. This capacitive coupling requires field penetration through the lower surface state and the interior of the thinner crystal, which fail to completely screen electric fields. The contrasting gating behavior of the devices is corroborated by the temperature dependence of their resistivities (see Supplemental Material [26]). We also note that while dual-gated TI devices have been previously reported [12,33], the devices reported here are unique in that the two surface states are tuned independently and separately observed.

Here, we focus on the capacitive coupling between the bottom gate and the upper surface in the thin crystal, and data regarding coupling of the top gate and lower surface are presented in the Supplemental Material [26]. The slope of $V_{peak}(V_B)$ is a measure of the ratio of the capacitive coupling of the bottom and top gates to the upper surface, which includes partial screening of electric fields by the lower surface state. At $V_B \sim -20$ V the slope of $V_{peak}(V_B)$

FIG. 1 (color online). (a) Colorized AFM image of device $A$, including schematic circuit elements describing the transport measurement, where $V_{XX}$ is the longitudinal voltage drop and $I_{SD}$ is the source-drain voltage. Red (dark horizontal bar) is BSTS, blue (wide vertical bar, variable brightness) is h-BN, and gold (bright) is Ti/Au (contacts and gate electrode). The scale bar is 2 microns. (b) TRARPES measurement of a BSTS crystal. The white line indicates the chemical potential.

FIG. 2 (color online). Gate dependence of the resistivity of devices $A$ and $B$. (a) Bottom gate dependence of resistivity at $V_T = 0$ at low temperature (blue, green) and 270 K (dashed line) from cooldown 2. (c) Top gate dependence of resistivity at $V_B = 0$ at low temperature (blue, green) and 270 K (dashed line) from cooldown 2. (b), (d) 2D map of resistivity while modulating both gate electrodes for devices $B$ and $A$, respectively, from cooldown 1. The black dots identify $V_{peak}(V_B)$, the top gate voltage at which $R_{peak}$ is found as a function of $V_B$.\n
and the resistance of the lower surface are simultaneously determined at a maximum, i.e., near the Dirac point [see Fig. 3(b)]. This is consistent with a minimum in the screening effectiveness of the lower surface state at the Dirac point. Understanding this behavior quantitatively requires a detailed charging model, which we discuss below.

By considering the BSTS surface states as a grounded pair of 2D electronic states, the general gating behavior can be understood via a charging model construction originally developed for parallel graphene layers [34]. This model is schematically represented in Fig. 3(a), where the important quantities are the applied gate voltages \( V_T \), \( V_B \), the geometric capacitances per unit area of the gates \( C_B \), \( C_T \), the intersurface capacitance per unit area \( C_{TI} \), the charge densities of the gate electrodes \( n_T \), \( n_B \), and the charge density and chemical potentials of the lower \( n_L, \mu_L \) and upper \( n_U, \mu_U \) surface states. Four coupled equations completely describe the charging of the system: one from charge neutrality and three from Faraday’s law, which restricts the sum of voltage drops around a loop to equal zero, including the change in chemical potential of the surface states \( \Delta \mu_j = \mu_j - \mu_0^j \), where \( \mu_0^j \) is the initial Fermi energy relative to the Dirac point for surface state \( j = U, L \). A detailed derivation is provided in the Supplemental Material [26]. For this Letter, we are interested in the condition that the chemical potential at the upper surface is at the Dirac point. By setting \( n_U = 0 \) and \( \mu_U = 0 \), a useful pair of equations can be derived

\[
\mu_L = -\frac{C_T}{C_{TI}} eV'_T, \tag{1}
\]

\[
\frac{1}{C_B} e n_L = V'_B + \left( \frac{1}{C_B} + \frac{1}{C_{TI}} \right) C_T V'_T, \tag{2}
\]

where \( V'_{T,B} = V_{T,B} - V_{0,T,B} \), and \( V_{0,T,B} \) are constants that depend on the initial densities and chemical potentials of the two surfaces (see Supplemental Material [26]). Equations (1) and (2) serve as a linear transformation from a trajectory in gate voltage space [Fig. 3(b)] to a relationship between chemical potential and density for the lower surface state [Fig. 3(c)].

Experimentally, three unknowns remain: the intersurface capacitance \( C_{TI} \) and the initial offset carrier densities of the upper and lower surfaces \( n^0_{L,U} \). To constrain these parameters, an independent measurement of \( \mu(n) \) is required. ARPES measurements of the surface state band structure can be easily converted to a model for \( E(n) \), including an explicit treatment of the bulk states [26]. A three-parameter least-squares fit between the transformation of the transport data and the ARPES model is performed and shown in Fig. 3(c) [26]. The interlayer capacitance from this fit is \( C_{TI} = 740 \pm 20 \text{nF/cm}^2 \), corresponding to an effective bulk permittivity of \( \kappa_{TI} \approx 32 \), comparable to values for similar compounds [35–37]. The initial electron densities of the upper and lower surface states are found to be \( n^0_L \approx 0.1 \times 10^{12} \text{ cm}^{-2} \) and \( n^0_U \approx 1.2 \times 10^{12} \text{ cm}^{-2} \), which agrees well with values simply calculated from the magnitude of \( V_T \) and \( V_B \) necessary to reach the resistance peaks.

It is important to note that \( C_{TI} \) can be affected in a few ways. For example, localized electronic states could polarize, increasing \( C_{TI} \). As another possibility, low-density, poorly conducting bulk states could weakly screen electric fields, reducing \( C_{TI} \). However, in the thin limit, the surface states should efficiently screen charged bulk impurities, resulting in an absence of charged puddles of bulk states at charge neutrality for crystals of thickness \(< 70 \text{ nm} \) [38]. This length scale is consistent with the observation that device \( B \) (82 nm thick) appears to have conducting states screening the two surfaces from each other.

We now turn to the behavior of the thin device in high magnetic fields. The Hall mobility of this sample is low, of order 200 cm$^2$/Vs; as a result, no evidence of Landau levels is found, and a clear \( R_{\text{peak}} \) remains. Nevertheless,
Assuming $C_{\text{ARPES model}}$ (red curve). Arrows indicate the increase in the total density of states. Detailed Shubnikov–de Haas analysis of band gap opening in the lower surface state. This work was partly supported by the DOE, Basic Energy Sciences Office, Division of Materials Sciences and Engineering, under Award No. DE-SC0006418 (V. F., S. E., H. S., and P. J. H.), by the Gordon and Betty Moore Foundation Grant No. GBMF2931 and the STC Center for Integrated Quantum Materials, NSF Grant No. DMR-1231319 (B. H. and R. C. A.), and by an MIT MRSEC Initiative under NSF Grant No. DMR-0819762 (F. M. and N. G.). This work made use of the Materials Research Science and Engineering Center Shared Experimental Facility Grant No. GBMF2931 and the STC Center for Integrated Quantum Materials, NSF Grant No. DMR-1231319 (B. H. and R. C. A.), and by an MIT MRSEC Initiative under NSF Grant No. DMR-0819762 (F. M. and N. G.). This work made use of the Materials Research Science and Engineering Center Shared Experimental Facility Grant No. GBMF2931 and the STC Center for Integrated Quantum Materials, NSF Grant No. DMR-1231319 (B. H. and R. C. A.), and by an MIT MRSEC Initiative under NSF Grant No. DMR-0819762 (F. M. and N. G.). This work made use of the Materials Research Science and Engineering Center Shared Experimental Facility.
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