Quantum tunneling through planar p–n junctions in HgTe quantum wells

L B Zhang¹, Kai Chang¹,⁴, X C Xie², H Buhmann³ and L W Molenkamp³

¹ SKLSM, Institute of Semiconductors, Chinese Academy of Sciences, PO Box 912, Beijing 100083, People’s Republic of China
² Department of Physics, Oklahoma State University, Stillwater, OK 74078, USA
³ Physikalisches Institut (EP3), Universitat Wuerzburg, Am Hubland, D-97074 Wuerzburg, Germany
E-mail: kchang@red.semi.ac.cn

New Journal of Physics 12 (2010) 083058 (10pp)
Received 3 May 2010
Published 27 August 2010
Online at http://www.njp.org/
doi:10.1088/1367-2630/12/8/083058

Abstract. We demonstrate that a p–n junction created electrically in HgTe quantum wells with inverted band structure exhibits interesting intraband and interband tunneling processes. We find a perfect intraband transmission for electrons injected perpendicularly to the interface of the p–n junction. The opacity and transparency of electrons through the p–n junction can be tuned by changing the incidence angle, the Fermi energy and the strength of the Rashba spin–orbit interaction (RSOI). The occurrence of a conductance plateau due to the formation of topological edge states in a quasi-one-dimensional (Q1D) p–n junction can be switched on and off by tuning the gate voltage. The spin orientation can be substantially rotated when the samples exhibit a moderately strong RSOI.

⁴ Author to whom any correspondence should be addressed.
1. Introduction

Electrical manipulation of the transport property of topological insulators (TIs) is becoming one of the central issues in condensed matter physics and has attracted rapidly growing interest. Instead of symmetry breaking and a local order parameter, TIs are characterized by a topological invariant \([1]–[3]\). Originally, graphene was proposed to be a quantum spin Hall (QSH) insulator \([4]\), but this turned out to be difficult to realize because the spin–orbit interaction (SOI) in this material is too small to create a gap observable in transport experiments. The QSH insulator state was proposed independently to occur in HgTe quantum wells (QWs) with inverted band structure \([5, 6]\) and was demonstrated in a recent experiment \([7]\). This nontrivial TI \([8]–[14]\) can be distinguished from an ordinary band insulator (BI) by the existence of a \(Z_2\) topological invariant \([8, 9]\). It possesses an insulating bulk and metallic edges, where electrons with up and down spins counter-propagate along opposite edges, and consequently it shows a quantized conductance when the Fermi energy is swept across the bulk gap \([7]\). These topological edge states are robust against local perturbations, e.g. impurity scattering \([8, 15, 16]\) processes and the Coulomb interaction \([17, 18]\), and they lead to novel transport properties. The low-energy spectrum of carriers in QWs with inverted band structure, e.g. TIs, can be well described by a four-band Hamiltonian \([5]\). It is highly desirable to study how to manipulate the transport properties electrically, from both a basic physics and a device application perspective.

In this work, we investigate quantum tunneling through planar p–n junctions in HgTe QWs with inverted band structure. Perfect tunneling transmission is found for electrons incident at normal angles. The opacity and transparency of the tunnel barrier can be controlled by tuning the angle of incidence of the charge carriers, the gate voltage (Fermi energy) and the SOI. An interesting spin refraction effect is found utilizing the strong Rashba spin–orbit interaction (RSOI) in HgTe QWs, i.e. the tunnel barrier is transparent for one spin orientation and opaque for the other. It is difficult to observe this phenomenon in a conventional semiconductor two-dimensional electron gas (2DEG) due to the weak SOI \([19]\). The quantum states in HgTe QWs consist of electron and heavy hole states and therefore show a very strong SOI that is one and even two orders of magnitude larger than that in the conventional 2DEG of a semiconductor. This therefore makes it possible to observe spin refraction. The topological edge channels lead to
Figure 1. (a) Schematic structure of a planar p–n junction in a HgTe QW with inverted band structure. The electric gate is positioned above the shaded region. (b) Schematic representation of the potential profile of the n–p junction and the band structure of HgTe QWs without RSOI (red solid lines) and with RSOI (black dashed lines). The green dashed line indicates the Fermi energy. $E_g^T$ indicates the gap in the transmission. (c) The transmission as a function of incident angle $\theta$ with RSOI, $\alpha = 50$ meV nm (the dashed curves), and without RSOI, $\alpha = 0$ (the solid curves) for a given Fermi energy $E_F = 20$ meV. $V_g = 5$, 0 and $-5$ mV correspond to the blue, red and black lines, respectively. The inset shows the spin refraction schematically. (d) Transmission as a function of Fermi energy $E_F$ for a fixed incident angle $\theta = 0$, $V_g = 5$ mV (black line) and $V_g = -5$ mV (blue line). $E_g$ denotes the bulk gap of the HgTe QW. The other parameters used in the calculation are $A = 364.5$ meV nm, $B = -686$ meV nm$^2$, $C = 0$, $D = -512$ meV nm$^2$ and $M = -10$ meV.

a plateau of the quantized conductance in the p–n junction sample with narrow transverse width, but the spin orientation can be rotated significantly by the RSOI. The conductance plateau can be destroyed when the QSH topological insulator is driven into the normal insulator by tuning the external electric field.

2. Theory

2.1. Two-dimensional (2D) system

A planar p–n junction, as schematically shown in figure 1, can be fabricated using top gates inducing an electrostatic potential in the electron gas underneath. Electrons are injected from a nearby quantum point contact (QPC) and transmitted and/or reflected at the interface of the p–n junction. The transmitted and/or reflected electrons can be collected by other QPCs to provide further angular information. The angle of incidence angle of the electrons can be tuned by applying a perpendicular magnetic field.
The low-energy spectrum of carriers in an HgTe QW with inverted band structure, including the RSOI, is

\[
H = \begin{pmatrix}
H(k) & H_{\text{RSOI}}(k) \\
H_{\text{RSOI}}(-k) & H^*(-k)
\end{pmatrix}
\]

\[
= \begin{pmatrix}
\epsilon_+ + M(k) & Ak_- & i\alpha k_- & 0 \\
Ak_+ & \epsilon_- - M(k) & 0 & 0 \\
-i\alpha k_+ & 0 & \epsilon_k + M(k) & -Ak_+ \\
0 & 0 & -Ak_- & \epsilon_k - M(k)
\end{pmatrix},
\]

(1)

where \( k = (k_x, k_y) \) is the carriers’ in-plane momentum, \( \epsilon_\pm = C - D(k_x^2 + k_y^2), \) \( M(k) = M - B(k_x^2 + k_y^2), \) \( k_\pm = k_\pm \pm i\alpha \) and \( A, B, C, D \) and \( M \) are the parameters describing the band structure of the HgTe QWs. Note that the insulator state is characterized by the parameter \( M \), which is determined by the thickness of the HgTe QW \( \frac{M}{M} \) (for negative (positive) \( M \), the QW is a TI (BI), respectively). Furthermore, \( H(k) = \epsilon_k I_2 + d_1(k)\sigma_x \), with \( d_1 = Ak_x, d_2 = Ak_y, d_3 = M(k) \). Without RSOI, i.e. \( \alpha = 0 \), the eigenvalues and eigenvectors are

\[
\epsilon_\pm = \epsilon_k \pm \sqrt{M(k)^2 + A^2k^2},
\]

\[
\chi_\pm = N_\pm (Ak e^{-i\theta_\pm}, \pm d(k) - M(k))^T,
\]

respectively, where \( N_\pm = (A^2k^2 + (d(k) - M(k))^2)^{-1/2} \) are the normalization constants, \( d(k) = \sqrt{M^2(k) + A^2k^2} \) and \( \theta \) is the incident angle. When the RSOI is included, i.e. \( \alpha \neq 0 \), the eigenvalues become

\[
\epsilon_+^\alpha = \epsilon_k \pm \frac{k\alpha}{2} + \sqrt{\left(M(k) \pm \frac{k\alpha}{2}\right)^2 + A^2k^2},
\]

\[
\epsilon_-^\alpha = \epsilon_k \pm \frac{k\alpha}{2} - \sqrt{\left(M(k) \pm \frac{k\alpha}{2}\right)^2 + A^2k^2},
\]

(3)

respectively. The corresponding analytical expression for the eigenvectors with RSOI is too lengthy to be presented here for simplicity.

We first model the electrostatic potential of a p–n junction by a step-like potential, i.e. \( V(x) = 0 \), for \( x < 0 \); \( V \) for \( x \geq 0 \). This idealized model can give the essential features of the quantum tunneling. A more realistic smooth potential will be used in calculations aiming at comparison with experiments (see figure 2). We assume that electrons are injected from the QPC on the left side of the junction with wave vector \( k_0^L = k_F \). The wave functions on the left/right side are then

\[
\psi_L(x < 0) = \left[ \chi_n^L e^{ik_0^L x \cos(\theta_n^L)} + \sum_{m=1}^{4} r_m \chi_m^L e^{-ik_0^L x \cos(\theta_m^L)} \right] e^{ik_0^L y \sin(\theta_n^L)},
\]

\[
\psi_R(x > 0) = \sum_{m=1}^{4} r_m e^{ik_0^R x \cos(\theta_m^R)} e^{ik_0^R y \sin(\theta_m^R)},
\]

(4)

where \( \chi_{\lambda}^L (\lambda = L, R) \) is a four-component vector. The transmission \( T(\theta) \) is obtained using scattering matrix theory by matching the wave functions and the currents at the interface of the
Figure 2. (a, b) Contour plot for the electron transmission as a function of incidence angle $\theta$ and gate voltage $V_g$ for Fermi energies $E_F = \pm 20$ meV in the ungated region, respectively. (c, d) Gate voltage $V_g$ dependence of conductance $G$ for the Fermi energy $E_F = \pm 20$ meV, respectively. Here the potential profile of the p–n junction is modeled by a smooth potential.

p–n junction with conserved $p_y = k_F \sin \theta_n^L$, i.e. $\psi_L(x < 0) = \psi_R(x > 0)$ and $j_L(x)\psi_L(x < 0) = j_R(x)\psi_R(x > 0)$, where $j_{L,R}(x)$ are the current operators along the propagation direction, i.e. the $x$-axis,

$$j_{L,R}(x) = \begin{pmatrix} -2D_xk_x & A & i\alpha & 0 \\ A & -2D_xk_x & 0 & 0 \\ -i\alpha & 0 & -2D_xk_x & -A \\ 0 & 0 & -A & -2D_xk_x \end{pmatrix},$$

where $D_+ = D + B$ and $D_- = D - B$. Lengthy analytic expressions for the transmission $T(\theta)$ and reflection coefficients can be obtained but are omitted here for brevity. The conductance can be calculated by $G = \int_{-\pi/2}^{\pi/2} T(\theta) f(E_F) \, d\theta$.

2.2. Quasi-one-dimensional (Q1D) system

For the Q1D system shown in figure 4(a), the traveling-wave-like or evanescent-wave-like eigenstates of the Schrödinger equation $H\psi = E\psi$ in a given region $\lambda$ can be written as

$$\psi_\lambda(x, y) = \exp(ik^L_x x) \sum_n \chi^L_n \phi_n(y),$$

where $\phi_n(y) = \sqrt{\frac{2}{W}} \sin \frac{n\pi y}{W}$ with $n = 1, 2, \ldots, N$ and $\{\chi^L_{m,n}\} (\lambda = L, R)$ represents the expanded coefficients. The longitudinal wave vector $k_x = k_0^L$ and the eigenvector $\chi^L_n$ (with $4 \times 1$ blocks $\chi^L_n; n = 1, 2, 3, \ldots$) are determined from the generalized eigenvalue problem,

$$\begin{pmatrix} 0 & 1 \\ S & T \end{pmatrix} \begin{pmatrix} \chi \\ F \end{pmatrix} = k_x \begin{pmatrix} 1 & 0 \\ 0 & X \end{pmatrix} \begin{pmatrix} \chi \\ F \end{pmatrix},$$

New Journal of Physics 12 (2010) 083058 (http://www.njp.org/)
where $F = k \chi$, $S$, $T$ and $X$ consist of $4 \times 4$ submatrices with the form

$$
S_{mn} = \begin{pmatrix}
\Delta^+_{mn} & -iA\eta_{mn} & \alpha\eta_{mn} & 0 \\
-iA\eta_{mn} & \Delta^-_{mn} & 0 & 0 \\
\alpha\eta_{mn} & 0 & \Delta^+_{mn} & -iA\eta_{mn} \\
0 & 0 & iA\eta_{mn} & \Delta^-_{mn}
\end{pmatrix},
$$

$$
T_{mn} = \begin{pmatrix}
0 & A & i\alpha & 0 \\
A & 0 & 0 & 0 \\
-i\alpha & 0 & 0 & -A \\
0 & 0 & -A & 0
\end{pmatrix},
$$

$$
X_{mn} = \begin{pmatrix}
D_+ & 0 & 0 & 0 \\
0 & D_- & 0 & 0 \\
0 & 0 & D_+ & 0 \\
0 & 0 & 0 & D_-
\end{pmatrix},
$$

and $\Delta^+_{mn}$, $\Delta^-_{mn}$ and $\eta_{mn}$ are given by

$$
\Delta^+_{mn} = C + V - E \pm M - D \pm \langle \varphi_m(y) | k^2 | \varphi_n(y) \rangle,
$$

$$
\eta_{mn} = \langle \varphi_m(y) | k | \varphi_n(y) \rangle.
$$

Assuming hard-wall confinement and an electron incident from the eigenstates in the left lead with wave vector $k^2_L$, and emerging on the right side, the wave functions in the left/right leads of the structure can be written as

$$
\psi_L(x < 0) = e^{ik^2_L x} \sum_n \chi^L_{ln} \varphi_n(y) + \sum_{mn} r_{mn} \chi^L_{m,n} e^{-ik^2_L x} \varphi_n(y),
$$

$$
\psi_R(x > 0) = \sum_{mn} t_{mn} \chi^R_{m,n} e^{-ik^2_R x} \varphi_n(y).
$$

For the Fermi energy lying in the bulk gap, electrons can only be injected from the edge states in the left lead, while for the Fermi energy lying above the bulk gap, electrons can be injected from both the bulk and the edge states. The number of the basis $N$ is chosen as large as it is necessary to obtain a desired convergent numerical result. Using scattering matrix theory and the Landauer–Büttiker formula, we can obtain the total transmission $T$,

$$
T = \sum_{m,n} \frac{|v^R_{m,n}|^2}{|v^L_{m,n}|^2} |t_{m,n}|^2,
$$

where RM denotes the summation over all right-moving modes in the left and right leads.

3. Numerical results and discussions

3.1. 2D system

We first investigate the tunneling behavior of the electron through a planar p–n junction in the 2D system. It is interesting to find that electrons incident perpendicularly on the interface can be perfectly transmitted by the p–n junction, without any backscattering, due to the unique band dispersion of HgTe QWs. This feature is caused by the helicity of the band structure and is very

New Journal of Physics 12 (2010) 083058 (http://www.njp.org/)
similar to Klein tunneling in graphene. However, there are a couple of important differences between a TI and graphene [20]: (i) the spin $\sigma$ in the Dirac Hamiltonian of graphene denotes a pseudospin referring to the sublattices, while in a TI, $\sigma$ is a genuine electron spin; (ii) in the TI, interband tunneling is suppressed because of the different energy dispersions in conduction and valence bands, i.e. different group velocities. In fact, the carriers will be completely reflected when the incident angle is larger than the critical angle $\theta_c$. The critical angle $\theta_c$ can be determined by Snell’s law, $\sin \theta_c / \sin \theta_R = k_R^2 / k_c^2 = n_R / n_L$, and thus $\theta_c = \arcsin(k_R^2 / k_c^2)$. An electron in the TI thus behaves as a photon that is injected from a material with a larger refractive index $n_2$ into a medium with a smaller refractive index $n_1$. This critical angle for electrons can be tuned significantly by changing the gate voltage $V_g$ (see the solid curves in figure 1(c)), while the critical angle is difficult to change for photons. As a next step, we now include the RSOI generated by a perpendicular electric field under the electric gate (in the gray region of figure 1(a)). It is in principle possible to tune the potential height and the strength of the RSOI independently utilizing top- and back-gates [21]. Interestingly, the RSOI leads to a spin-dependent change in the critical angle $\theta_\sigma$ (see figure 1(c)). This is because the RSOI induces a spin splitting in the band structure, resulting in different Fermi wave vectors $k_\sigma^F (\sigma = \uparrow, \downarrow)$ (see the solid curves in figure 1(b)) for the two subbands. This implies that it is possible to induce spin-dependent reflection at the p–n junction, which allows for a fully spin-polarized tunneling current when the incident angle $\theta$ of the electrons is tuned properly (see the inset of figure 1(c)), i.e. $\theta_\sigma > \theta > \theta_1$, where $\theta_1 (\theta_\uparrow)$ is the critical angle for the spin-up (down) electron. The difference between the critical angles $\theta_c$ with and without the RSOI can be tuned significantly by adjusting the gate voltage. This allows the realization of a spin refractive device, as the building block for spin optics. When we tune the Fermi energy for a fixed RSOI strength and incidence angle $\theta$, the transmission exhibits a gap in which electron tunneling is forbidden. The width of the gap $E_g^\uparrow$ corresponds to the total bulk gap of the inverted HgTe QW (see figures 1(b) and (d)). Figure 1(d) shows that an initially opaque medium can suddenly become transparent by adjusting the Fermi level slightly. This rapid switching between opacity and transparency is a unique feature of the HgTe planar p–n junction.

In order to see how the tunneling behavior can be controlled by adjusting the incidence angle and the Fermi energy on both sides of the junction, we show the $\theta$- and $V_g$-dependent transmission in figure 2. One observes clearly that perfect transmission occurs at small incidence angles $\theta_c \rightarrow 0$ for the intraband tunneling process, e.g. the n–n and p–p processes (see figures 2(a) and (b)). In the vicinity of the gap where the tunneling is forbidden, the critical angle $\theta_c$ depends strongly on the gate voltage $V_g$ that tunes the Fermi level on the right-hand side of the junction. The perfect intraband transmission could thus provide us with a new experimental method to map the band structure in the vicinity of the bulk gap. As is evident from figures 2(c) and (d), the conductance shows a strongly asymmetric behavior when comparing intraband and interband tunneling. For the n–n or p–p processes, the perfect transmission leads to a large conductance, while the small interband transmission, i.e. the n–p process, gives rise to a small conductance. This interesting feature is in good agreement with a very recent experiment (figure 2.4 of [22]). The quantitative difference comes from the simplified model used in the calculation in which the contribution from the light hole is neglected when the Fermi energy is located at the bulk states.

We further investigate the spin transport properties of the p–n junction by calculating the spin projection $\langle \Sigma^\sigma \rangle = \langle S^\sigma_{c,e} \rangle + \langle S^\sigma_{c,h} \rangle$, where $\langle S^\sigma_{c,e} \rangle (\langle S^\sigma_{c,h} \rangle)$ denotes the z-component of the electron (heavy hole) spin. When the RSOI is present in the system, which, like the Fermi
Figure 3. The spin projection $\langle \Sigma^z \rangle$ as a function of gate voltage $V_g$ (in panels a and b) for a fixed Fermi energy $E_F = 20$ meV and spin orientation on the left side of the p–n junction, which is indicated by the dashed blue lines. The red and black lines correspond to the spin projections $\langle \Sigma_1^z \rangle$ and $\langle \Sigma_1^+ \rangle$ of the transmitted electron, respectively. The insets schematically show the changes in the spin orientations for the cases without (panel a) and with (panel b) RSOI. The other parameters are (a) $\theta = 0$, $\alpha = 0$ and (b) $\theta = 0$, $\alpha = 5$ meV nm.

energy, can be tuned by the voltage imposed on the structure (in the right gray region of figure 1(a)), the spin orientation of transmitted carriers changes greatly even for very weak RSOI ($\alpha = 5$ meV nm). For example, consider an electron injected from the left side carrying up-spin through the p–n junction. When the RSOI is absent (see figure 3(a)), the gate voltage $V_g$ only slightly changes the spin projection $\langle \Sigma^z \rangle$. In the presence of a moderate RSOI ($\alpha = 5$ meV nm), the spin orientation of transmitted electrons becomes in-plane, since the out-of-plane component $\langle \Sigma^z \rangle = \langle \Sigma_1^+ \rangle + \langle \Sigma_1^- \rangle$ vanishes due to the effective magnetic field resulting from the RSOI (see figure 3(b)). This demonstrates that the RSOI can substantially rotate the spin orientation.

3.2. QID system

Up to now, we have considered only the 2D case and ignored the contribution from the helical edge states. Next, we turn to discuss the boundary effect on the quantum tunneling through a p–n junction in a Q1D QSH bar. Such a Q1D QSH bar structure, which could be fabricated using standard lithographic techniques, is shown schematically in the inset of figure 4(a).

The helical edge states can lead to a conductance plateau at $G = 2e^2/h$ when the Fermi energy is located within the bulk gap. When electrons tunnel through a p–n junction including edge states, perfect transmission can be observed for the intra- and inter-band tunneling processes due to the conservation of the helicity of the edge states. This perfect transmission in a Q1D system differs from that in two dimensions discussed above (see figure 2). The minigaps in figure 4(a) indicated by the vertical arrows in the plateau are caused by the coupling between the counter-propagating edge states in the n- and p-regions with opposite spins stemming from the opposite sides of the QSH bar. The most interesting aspect revealed by our calculation is that the conductance plateau due to the edge states is independent of gate voltage $V_g$ due to the
Figure 4. (a, b) Fermi energy dependence of the transmission \( T \) (red line) in a QSH bar with gate voltage \( V_g = -10 \) mV for TI–TI and TI–BI tunneling. \( E_g(\text{eV}_g) \) denotes the gap in the bulk (the height of the potential profile of the p–n junction), ranging from −11 to 13 meV. The insets show schematically the edge channels in the TI/TI and TI/BI p–n junctions in the QSH bars. (c) The spin projections \( \langle \Sigma^z \rangle \) as a function of gate voltage \( V_g \) for a fixed Fermi energy \( E_F = 10 \) meV. The dashed blue lines denote the spin orientation of the incident electrons. The solid and dashed lines correspond to the cases with and without RSOI, respectively, where the red and black curves represent the spins \( \langle \Sigma^z_1 \rangle \) and \( \langle \Sigma^z_2 \rangle \). (d) Contour plot of the transmission as a function of \( M \) and \( V_g \) for a fixed Fermi energy \( E_F = 10 \) meV. The width of the QSH bar is set at \( W = 200 \) nm.

conservation of the helicity of the edge states. A pronounced perfect transmission feature can also be observed in the other tunneling processes, i.e. p–p, n–n, p–n and n–p junctions. In a Q1D system, this feature is quite distinct from that in a 2D system (see figure 2). Surprisingly, the tunneling process between BI and TI, which is shown in figure 4(b), is very different. By tuning the band parameter \( M \) electrically [6], the transport through the topological edge channels can be blocked for the TI/BI (\( M > 0 \)) hybridized structure and becomes vanishingly small for the topological insulator (\( M < 0 \)) (see the insets of figures 4(a) and (b)). When the electrons are transported in a TI system (\( M < 0 \)), the conductance plateau does not change with gate voltage \( V_g \) except for the region of the minigaps, whereas for \( M > 0 \), the conductance plateau disappears totally because the BI blocks the edge channels. This means that one can switch on/off the transport property electrically in a hybrid BI/TI system. Interestingly, the RSOI would not affect the transmission plateau since the RSOI preserves time reversal symmetry and would not destroy the edge channels. The spin orientation of the transmitted electrons can, however, be changed significantly by the RSOI, just as in the 2D case. In figure 4(c), one observes that the spin projection \( \langle \Sigma^z \rangle \) of the transmitted electrons vanishes for spin-up injection because the RSOI behaves like an in-plane magnetic field leading to a giant spin rotation.
4. Conclusion

In summary, we have studied quantum tunneling through p–n junctions in HgTe QWs with inverted band structures in 2D and Q1D systems. An interesting perfect transmission of the quantum tunneling process is found for electrons injected normal to the interface of the p–n junction in the 2D system. The opacity and transparency of the p–n junction can be tuned by changing the incidence angle of the incoming carriers, the gate voltage (which determines the Fermi level on the right-hand side of the junction) and the strength of the RSOI. Spin-up and spin-down electrons can be separated spatially utilizing the RSOI, which could result in a building block for spin optics. In the Q1D system, tunneling through topological edge states can be switched on and off by tuning the band parameters electrically. This provides an efficient means of controlling the transport properties of topological edge channels electrically.

Acknowledgments

This work was supported by NSFC grant numbers 60525405 and 10874175. XCX is supported by the US DOE and the US NSF. KC thanks Professor R B Tao for a helpful discussion.

References

[1] Thouless D J, Kohmoto M, Nightingale M P and den Nijs M 1982 Phys. Rev. Lett. 49 405
[2] Haldane F D M 1988 Phys. Rev. Lett. 61 2015
[3] Wen X G and Niu Q 1990 Phys. Rev. B 41 9377
[4] Kane C L and Mele E J 2005 Phys. Rev. Lett. 95 226801
[5] Bernevig B A, Hughes T L and Zhang S C 2006 Science 314 1757
[6] Yang W, Chang K and Zhang S C 2008 Phys. Rev. Lett. 100 056602
[7] Li J and Chang K 2009 Appl. Phys. Lett. 95 222110
[8] König M, Wiedmann S, Brune C, Roth A, Buhmann H, Molenkamp L W, Qi X L and Zhang S C 2007 Science 318 766
[9] Kane C L and Mele E J 2005 Phys. Rev. Lett. 95 146802
[10] Fu L and Kane C L 2007 Phys. Rev. B 76 045302
[11] Hsieh D, Qian D, Wray L, Xia Y, Hor Y S, Cava R J and Hasan M Z 2008 Nature 452 970
[12] Zhang H J, Liu C X, Qi X L, Dai X, Fang Z and Zhang S C 2009 Nat. Phys. 5 438
[13] Xia Y et al 2009 Nat. Phys. 5 398
[14] Ran Y, Zhang Y and Vishwanath A 2009 Nat. Phys. 5 298
[15] Bernevig B A and Zhang S C 2006 Phys. Rev. Lett. 96 106802
[16] Li J, Chu R L, Jain J K and Shen S Q 2009 Phys. Rev. Lett. 102 136806
[17] Xu C and Moore J E 2006 Phys. Rev. B 73 045322
[18] Wu C J, Bernevig B A and Zhang S C 2006 Phys. Rev. Lett. 96 106401
[19] Khodas M, Shekhter A and Finkel’stein A M 2004 Phys. Rev. Lett. 92 86602
[20] Beenakker C W J 2008 Rev. Mod. Phys. 80 1337
[21] Nitta J, Akazaki T, Takayanagi H and Enoki T 1997 Phys. Rev. Lett. 78 1335
[22] König M 2007 Spin-related transport phenomena in HgTe-based quantum well structures PhD Thesis Universität Würzburg

New Journal of Physics 12 (2010) 083058 (http://www.njp.org/)