Topological edge states and transport properties in zigzag stanene nanoribbons with magnetism

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Abstract
In this work, we investigate the topological phase transitions and corresponding transport properties in zigzag stanene nanoribbon with different magnetism. The results show that the off-resonant circularly polarized (ORCP) light may induce anisotropic chiral edge state with a magnetic phase transition from antiferromagnetic state to nonmagnetic state. In combination with the ORCP light and electric field, the 100% spin-polarized edge state can be induced with some magnetic orders. The finite-size effect is also an important factor for the magnetic phase transitions, which in turn induces topological phase transitions from the band insulator to topological phases. By constructing the topological-insulator junctions with some topological edge states, we further study the Fabry–Perot resonant, where multiple reflection edge states cause strong current loops. By modulating the ORCP and electric field, the system can also be regarded as a switcher, to control the charge current or spin polarized current. These findings pave a way for designing topological device with magnetic edges in the future nano spintronics.

1. Introduction
Stanene, synthesized recently [1, 2], is a two-dimensional topological insulator with robust edge states on the surface [3–6]. Compared with graphene’s weak–orbit coupling and plane geometry, the stronger spin–orbit coupling and buckled geometry in silicene and germanene as well as stanene can lead to abundant topological phases [7–13]. Also because of the buckled geometry, topological edge states in these other versions of graphene can be easily modulated by external fields. Without considering spontaneous edge magnetism, there are some recent typical applications in topological transistor by using robust edge states. By modulating the vertical field, Zheng et al [14] theoretically designed multichannel depletion-type field-effect transistor in the zigzag germanene nanoribbon, which is important for topological device with low dissipation. They also used the off-resonant circularly polarized (ORCP) light to obtain the switching between a 100% spin-polarized weak current and non-polarized strong current [15].

For the realistic applications of the silicene-like zigzag nanoribbons, the spontaneous edge magnetism should be considered, where the antiferromagnetic state is the ground state due to its stability [10, 16–18]. It should be noted that these prior works [19–21] has shown the effect of magnetic order on topological and transport properties, which gives insight into the related research in our system. In particular, this work [22] displayed that the magnetization has a remarkable influence on the transport properties of the edge channel. Without the spontaneous edge magnetism, some deep properties cannot be theoretically predicted [19, 20, 22–27]. Especially in narrow zigzag nanoribbons, the Coulomb interaction from the Hubbard term is relatively strong [18, 23, 27, 28]. In the mean-field Kane–Mele–Hubbard approximation, some studies have investigated the topological and magnetic phase transitions in silicene-like nanoribbons [27, 29–31]. But all these studies do not consider the ORCP light. Actually, the off-resonant light is not only important for the topological phase transition, but also for the magnetic phase transition. According to bulk-edge
correspondence [10, 11, 31, 32], some distinct topological phases are reported in the zigzag nanoribbons with the edge magnetism. And for different topological edge states, some works [4, 15, 33, 34] found interesting transport mechanisms. But the transport properties of the topological edge states with the edge magnetism have not been investigated. It should be noted that the electronic structure and low-energy effective model of stanene are first discussed in this work [35]. Here, we focus on the transport properties of zigzag stanene nanoribbon (zStNR) with the spontaneous edge magnetism, and the influence of the ORCP light.

In this work, we give detailed investigations on the topological and magnetic phase transitions and the transport properties of some edge states by using the mean-field Kane–Mele–Hubbard model and the nonequilibrium Green’s function (NEGF) theory. Compared with the few current researches of the off-resonant light on the topological materials with magnetism, we find that the off-resonant light and electric field may induce many phase transitions. In addition, the finite-size effect is also a major factor. With these effects, the initial band-insulator (BI) is transformed into the quantum anomalous Hall (QAH) and spin-polarized quantum anomalous Hall (SQAH) states, while the initial antiferromagnetic edges are turned into the reverse antiferromagnetic edge (with the magnetic direction opposite to the antiferromagnetic state) or nonmagnetic edge. In order to further investigate the transport properties of these magnetic topological edge states, we also study the junction systems with the spontaneous edge magnetism.

2. Model and methods

We used the Kane–Mele–Hubbard model with NEGF theory to investigate the electronic and magnetic properties in zStNR with edge magnetism. With the ORCP light and electric field, the Hamiltonian in the mean-field Hubbard approximation can be expressed as

\[
H = -t \sum_{<ij,\alpha} \sigma^z_{\alpha,\beta} c^\dagger_{i\alpha} c_{j\beta} + i \frac{\lambda}{3\sqrt{3}} \sum_{\langle\langle ij\rangle\rangle,\alpha,\beta} \nu_{ij\alpha} \sigma^x_{\alpha,\beta} c^\dagger_{i\alpha} c_{j\beta} + i \frac{\lambda}{3\sqrt{3}} \sum_{\langle\langle ij\rangle\rangle,\alpha,\beta} \nu_{ij\alpha} c^\dagger_{i\alpha} c_{j\beta} + U \sum_{i\alpha} \left( \langle n_{i\alpha\uparrow} \rangle - 1/2 \right) n_{i\alpha} + \Delta \sum_{i\alpha} \xi_i c^\dagger_{i\alpha} c_{i\alpha},
\]

where \(c^\dagger_{i\alpha}(c_{i\alpha})\) denotes the creation (annihilation) operator for an electron at the lattice site \(i\) with spin \(\alpha\), \(\sigma^z_{\alpha,\beta}\) is the spin Pauli matrix element with respect to the \(z\) component. The first term presents the hopping between nearest-neighbor sites \(<ij\rangle\) with the strength \(t = 1.3\) eV. The second term presents the intrinsic spin–orbit coupling between next-nearest-neighbor sites \(<\langle ij\rangle\rangle\) with the strength \(\lambda = 0.1\) eV [24, 36], where \(\nu_{ij} = +1(-1)\) if the next-nearest-neighbor hopping is anticlockwise (clockwise) with respect to opposite \(z\) axis. The third term stands for the ORCP light with the strength \(\lambda_{\Omega}\), where \(\lambda_{\Omega}\) is set as negative (positive) value for left (right) circulation [37, 38]. The fourth term shows the \(e-e\) interaction in mean-field Hubbard approximation, where \(n_{i\alpha}\) denotes the particle number operators at site \(i\) with spin \(\alpha\), \(\langle n_{i\alpha\uparrow} \rangle\) is the corresponding expectation value called as the averaged electron density. \(U\) is the on-site Coulomb repulsion constant, based on the first-principles calculation and experimental fittings (here we set \(U = t\) [1, 17, 36, 39]). The last term is for the staggered electric field, with \(\xi_i = +1(-1)\) for the A (B) sublattice.

For the periodic system, the averaged electron density in the first Brillouin zone can be expressed as

\[
\langle n_{i\alpha\uparrow} \rangle = \frac{a}{2\pi} \int_0^{2\pi/a} n_{i\alpha}(k) dk,
\]

where \(n_{i\alpha}(k) = \phi^\dagger_{i\alpha}(k)\phi_{i\alpha}(k)\) is the wavevector-dependent electron density at site \(i\), \(a\) is the distance between two nearest unit cells. In addition, the Fermi–Dirac distribution is written as \(f(E, \mu) = 1 / \{ \exp[(E - \mu)/k_B T] + 1 \}\), where the Fermi energy is set as zero for simplicity. Following the details of our previous work [27], the averaged electron density can be used to calculate the band structure with the self-consistent calculations. For the open system, such as a nanojunction system, we employ the NEGF method to evaluate the averaged electron density. For the two-terminal system, the regarded spin-dependent Green’s function of the device is expressed as [40]

\[
G^{\alpha\bar{\alpha}}(E) = \left[ (E + i\eta)I - H^\alpha_0(E) - \sum_L \sum_R \left( E - \sum_L \langle E \rangle \right) \right]^{-1},
\]

where \(\eta\) is the infinitesimal real number and \(I\) is the unit matrix sharing the same dimension of the device. \(H^\alpha_0\) is the spin-dependent Hamiltonian of the device with the spin index \(\alpha\), \(\sum_{L/R}^{\alpha\bar{\alpha}}\) is the self-energy matrix,
describing the coupling between the left/right lead and device, which can be obtained from the surface Green’s function \([41, 42]\). Moreover, the spin-dependent Hamiltonian \(H^\alpha_i\) contains the averaged density \(\langle n_{\alpha i} \rangle\) with opposite spin, that is related by the local density of state

\[
\rho_{\alpha i}(E) = -\frac{1}{\pi} \text{Im}[G^\alpha_i(E)].
\]  

(4)

Then, the averaged electron density \(\langle n_{\alpha i} \rangle\) at each site is calculated by

\[
\langle n_{\alpha i} \rangle = \int_{-\infty}^{+\infty} \rho_{\alpha i}(E)f(E, \mu)dE.
\]  

(5)

By combining this equation with equations (2)–(4), the self-consistent solution \(\langle n_{\alpha i} \rangle\) can be obtained. Thereby, one can calculate the magnetic moment at each site, given by this formula \(M_2 = \frac{1}{2}(\langle n_{\uparrow i} \rangle - \langle n_{\downarrow i} \rangle)\). In addition, the residue theorem is the efficient and quick method for calculating equation (5) rather than the direct integration \([43–45]\). With the self-consistent calculation, one can evaluate the transmission coefficient from the left lead to the right lead by the following formula

\[
T_{L-R}^\alpha(E) = \text{Trace}[\Gamma^\alpha_R(E)G^\alpha(E)\Gamma^\alpha_L(E)G^\alpha(E)],
\]  

(6)

where \(\Gamma^\alpha_R(E) = -2 \text{Im}[\sum_{\alpha R} \Gamma^\alpha(E)]\) are the linewidth functions, and the advanced Green’s function reads \(G^\alpha(E) = [G^\alpha(a)]^*\). At zero temperature, by assuming the linear response limit, the spin-dependent conductance is given by \(G^\alpha(\mu) = (e^2/h)T^\alpha(\mu)\). In our calculation, we set a sufficiently low temperature \(k_B T = 0.001\) eV throughout our paper, in which the conductance formula is approximately valid. In order to study the detailed transport mechanism of the zStNR, we would present the local bond currents. With the help of the Heisenberg equation of motion and NEGF formulism, the formula can be expressed as \([46, 47]\)

\[
f_{ij}^\alpha(E) = -\frac{1}{\hbar} H_{ij} \text{Im}[G^\alpha(E)\Gamma^\alpha_j(E)G^\alpha_i(E)]_{ji},
\]  

(7)

where \(f^\alpha_{ij}(E)\) is the spin-resolved local bond current from site \(i\) to \(j\). Here, we assume that the local bond current propagates from the left lead.

### 3. Results and discussions

In order to simulate the transport properties in the topological-insulator junctions made of zStNR with the edge magnetism, we choose a fixed low temperature \(k_B T = 0.001\) eV. Actually, we have checked that this low temperature has no effect on the edge magnetism (not shown here), so we can regard this system as zero temperature case. The SOC constant in this paper is chosen as \(\lambda_{\Omega} = 0.1\) eV for StNR system. The unit cell of zStNR consists of \(N_0\) sub-unit and each contains four atoms shown in figure 7(f). In addition, except for the color labels in a few figures, the red (blue) line denotes the spin-up (spin-down) mode throughout this paper.

#### 3.1. Topological and magnetic edges and phase transitions in StNR

Here, we employ the mean-field Kane-Mele-Hubbard model to calculate the band structures and magnetic moment with an initial antiferromagnetic order. In our previous work \([27]\), we found that the large width of the system with the edge magnetism and strong spin–orbit coupling lead to the quantum spin Hall state. It is well known that the distribution of the spin-up and spin-down densities in zStNR is different from those in 2D infinite stanene due to the modified boundaries, this quick and simple method \([31]\) is used to evaluate topological properties based on the number of the edge states. From this method, we further study the QAH and SQAH states with external fields. In present of the ORCP light, the band structures are shown in figures 1(a) and (b), corresponding to the right and left circulation, respectively. And the band structures in figures 1(a) and (b) correspond to the anisotropic chiral edge states in figures 1(c) and (d), respectively. It is obvious that these band structures are very like the ones without the edge magnetism \([4, 7]\), but the origin is different. From figures 1(a) and (b), we see the band slopes, or the group velocities differ a lot between the two spin components, which means these two chiral edge states are anisotropic, not as those in common QAH state.

In order to investigate the topological and magnetic phase transitions with the off-resonant light \(\lambda_{\Omega}\), we present the energy gap and magnetic moment in figures 2(a) and (b), respectively. It is shown that the spin-up (spin-down) gap closes in the region of \(\lambda_{\Omega} \gtrsim 0.27\) eV, corresponding to the anisotropic chiral edge state in figure 1(c) (figure 1(d)). And the maximum energy gaps of the spin-up and spin-down lie at the values of \(\lambda_{\Omega} = 0.1\) eV and \(\lambda_{\Omega} = 0.1\) eV, respectively. In figures 2(a) and (b), it is shown that with
Figure 1. The band structures (a) and (b) and corresponding edge states (c) and (d) of zStNR under light. (a) $\lambda_{\Omega} = 0.3$ eV and (b) $\lambda_{\Omega} = -0.3$ eV correspond to (c) and (d), respectively. The fixed parameters are set as $N_y = 3$, $\lambda_{so} = 0.1$ eV, $t = 1.3$ eV, and the thick lines denote the edge states.

Figure 2. The energy gap and magnetic moment of zStNR as a function of the off-resonant circularly polarized light intensity. The parameters of the system are set as $N_y = 3$, $\lambda_{so} = 0.1$ eV.

Increasing the strength of $\lambda_{\Omega}$ in the region $[0, 0.35 \text{ eV}]$, the spin-up gap gradually decreases to zero, the spin-down gap gradually increases to the peak at $\lambda_{\Omega} = 0.1$ eV and then to zero, and $M_Z$ decreases slowly, then sharply to zero. For the region $-0.35 \text{ eV} \leq \lambda_{\Omega} \leq 0$, the tendency of the band gaps with opposite spin is symmetric with the case in the region $0 \leq \lambda_{\Omega} \leq 0.35 \text{ eV}$. In addition, we see that in the region of $|\lambda_{\Omega}| < 0.27 \text{ eV}$, the system keeps steady in the antiferromagnetic state. Beyond this region, the antiferromagnetic state is broken into non-magnetic state, which indicates the strong off-resonant light can break the edge magnetism. Moreover, we find that the regions of non-magnetic state exactly correspond to the anisotropic chiral edge states.

Based on the energy gap shown in figure 2(a), the staggered electric field is used to find out 100% spin-polarized edge state. For this purpose, we choose the maximum gap of the spin-up or spin-down mode in figure 2(a) as our initial condition. Figures 3(a) and (b) show that the QAH state can be translated into the SQAH state in the appropriate electric field, corresponding to the 100% spin-polarized edge states.
Figure 3. The band structures and corresponding edge states with different parameters of (a) and (c) $\lambda = 0.1 \text{eV}$, $\Delta = 0.14 \text{eV}$; (b) and (d) $\lambda = -0.1 \text{eV}$, $\Delta = -0.14 \text{eV}$. The fixed parameters are set as $N_y = 3$, $\lambda_{so} = 0.1 \text{eV}$, and the thick lines present the edge states.

Figure 4. The energy gaps and magnetic moments as a function of different electric fields. The fixed parameters are set as $N_y = 3$, $\lambda_{so} = 0.1 \text{eV}$, other parameters are set as: (a) and (c) $\lambda = 0.1 \text{eV}$; (b) and (d) $\lambda = -0.1 \text{eV}$. 
Figure 5. The band structures, energy gap and magnetic moment as a function of the ribbon width $N_y$. The fixed parameters of zSInR are set as $\lambda_{so} = 0.1\, eV$, $\lambda_{13} = 0.15\, eV$.

Figure 6. The band structures, energy gap and magnetic moment as a function of the ribbon width $N_y$. The fixed parameters of zSInR are set as $\lambda_{so} = \lambda_{13} = 0.1\, eV$, $\Delta = -0.14\, eV$. 
These edge states with opposite spin modes and transport directions in figures 3(c) and (d) correspond to the band structures in figures 3(a) and (b), respectively. It is well known that as the gaps of the valleys \( K \) and \( K' \) with spin modes are symmetrical with respect to the energy axis at \( k = \pi/a \), the crossing point of the edge states would be at \( k = \pi/a \). Compared with the case in figure 1, the crossing points of the edge states in figures 1(a) and (b) are not at \( k = \pi/a \), as the symmetry of the valleys are broken down by the electric field and magnetism.

In order to investigate the trends of energy gaps and magnetic moment with the electric field, we present the corresponding pictures in figures 4(a)–(d). It is noted that the system experiences the topological and magnetic phase transitions under the electric field. In figure 4(a), we note that the gaps of the spin-up mode disappear in the regions of \(-0.25 \, \text{eV} \leq \Delta \leq -0.17 \, \text{eV} \) and \( 0.11 \, \text{eV} \leq \Delta \leq 0.25 \, \text{eV} \), corresponding to the 100% spin-polarized edge state in figure 3(c). Based on this work [22], the AF (FM) state mostly acts as antiferromagnetic (ferromagnetic) exchange field \( M \). We approximatively regard the effect of the AF state as an average antiferromagnetic exchange field \( \bar{M} = \frac{1}{N} \sum_{i=1}^{N} M_i \), where \( N \) is the total atom number of a unit cell. In addition, this effective \( M_i \) are mainly distributed in the boundaries of the system. It should be noted that the band gaps in figures 4(a) and (b) both have an abrupt change due to the corresponding magnetic abrupt change (or the change of the average antiferromagnetic exchange field \( \bar{M} \)). Compared to the two topological phase transitions (at \( \Delta = -0.17 \, \text{eV} \) and \( \Delta = 0.11 \, \text{eV} \)) in figure 4(a), the corresponding edge magnetic moment indicates the magnetic phase transition in figure 4(c). We see although there only exist single-spin bands near the Fermi level, there still exist antiferromagnetic edges with two different magnetisms. These two regions, crossing the magnetic phase transition point \( \Delta = -0.17 \, \text{eV} \), correspond to the antiferromagnetic state and reverse antiferromagnetic state (figure 4(c)), respectively. The tendency of the band gap and magnetic moment become antisymmetric as the right circulation light is changed into the left circulation light, as shown in figures 4(b) and (c), respectively. Compared to the edge state in figure 4(c), the one in figure 4(d) possesses the opposite spin mode and direction.

In the following, we will take into account the finite-size effect on the topological and magnetic phase transitions. Actually, the magnetic change that is affected by the finite-size effect is the direct factor for the topological phase transition. For the purpose of finding out general results, typical cases in figures 2(a) and
Figure 8. Different types of the topological-insulator junction systems and corresponding transmission spectra. (a)–(c) correspond to (d)–(f), respectively, all the topological edge states and corresponding parameters are from figures 1(c) and (d) and 3(c), and the lengths of the device in different system are set as $N_x = 40$.

4(a) are selected as our examples to investigate the finite-size effect. In figures 5(a)–(c), it is shown that the spin-up and spin-down gaps sequentially become closed with increasing the ribbon width. The band structure with the small spin-down gap in figure 5(b) corresponds to the 100% spin-polarized edge state (not shown here); and the one in figure 5(c) without gaps corresponds to the anisotropic chiral edge state. These results display that with increasing the width the system experiences the topological phase transitions from the band insulator (BI) to the SQAH state and then the QAH state, for which the finite-size effect is not direct factor. In figure 5(e), although the $M_Z$ almost keeps steady, the $M$ decreases with increasing the width, which is direct factor for topological phase transitions. We also present the picture to illustrate the details of the tendency of the energy gap with the width in figure 5(d). It is noted that the spin-down gap gradually decreases to zero when increasing the width, and then the zero gaps keep steady for the region of $N_y > 19$. We also note that the spin-up gaps maintain steady for the region of $N_y > 5$. In figure 5(d), the finite-size effect almost has no influence on the magnetic moment, which means the process of the topological phase transition is still under the antiferromagnetic state. This is not like the cases in figures 2 and 4, where the topological and magnetic phase transitions may simultaneously appear.

Here, we choose the case in figure 4(a) as our example to investigate the finite-size effect. From the band structures in figures 6(a)–(c), we can see that the spin-up gap gradually decreases to zero as the width reaches the value $N_y = 5$, while the spin-down gap closes and then opens. Actually, the band structure in figure 6(c) corresponds to the 100% spin-polarized edge state. In addition, there is no hidden topological phase in the band structure in figure 6(b). In order to investigate the detail of the finite-size effect on the topological and magnetic phase transitions, we present the tendency of the energy gap and magnetic moment in figures 6(d) and (e), respectively. It is noted that the spin-down gap suddenly opens at the value $N_y = 7$, while the spin-up gap still closes in the region of $N_y > 7$ in figure 6(d). Moreover, the value $N_y = 6$ of the rapid change of the spin-down gap is also the critical value of the magnetic phase transition in figure 6(e). Actually, the magnetic transition $M_Z$ means an abrupt change $M$, causing that the band gap experiences an abrupt change.
3.2. Transport properties of the topological edges with magnetism in StNR

In our previous work [4], we used the topological edge states to design different types of the filters and switches in the zigzag silicene nanoribbon without the edge magnetism. In the following, the focus of our discussions is to investigate the transport properties of these edge states (shown in figures 2 and 4) in the zigzag StNR with the edge magnetism.

In figure 7(a), the system is composed of different topological edge states, where the leads (from figures 1(c) and (d)) have no magnetic state with strong off-resonant light and the device (from figure 3(d)) is in the antiferromagnetic state. And this magnetic moment distribution is shown in figure 7(b). In this system, there is no spin-up edge state in the device, we can expect the spin-up transmission is zero. For the spin-down edge states, as there exist two paths for transmitting and reflecting the device, we expect the transmission is between 0 and 1. In figures 7(c)–(e), there are a lot of spin-down transmission peaks due to some resonance. It is shown that the number of the transmission peaks increase as the length of the device increases. For the larger length of the device, the tunneling effect almost disappears, leading to zero transmissions of the spin-down mode in figures 7(d) and (e). Actually, these phenomena, where the number of the spin-down transmission peaks increase with increasing the length of the device, satisfy this empirical formula of the Fabry–Perot (FP) interference \[ T = g/L \] [48]. \( g \) is a constant denoting the state degeneracy, and the separation period \( T \) becomes smaller for the longer device length \( L \).

In order to further study the FP interference, we present the local current distributions at the first peak near zero energy in figure 7(c). In figure 7(f), the spin-down local current from the lead L reflects, due to the opposite transport direction between the lead L and device, and then propagates downward along the
interface between the lead and device. Afterwards, this local current propagates along the lower edge of the device and forward along the interface between the device and lead R. At last, the local current propagates along the upper edge of the device, leading to the current loop. After multiple times of this reflected process of local current, the current loop will become so strong. At this time, as much the local edge current is incident from the left, as much the local current is transmitted along the upper edge of the lead R, this leads to the 100% spin-down transmission. In figure 7(g), due to the gap effect of the device, the spin-up current from the lead L is reflected and propagates downward along the interface between the lead L and device, and then propagate along the lower edge of the lead L, this process leads to zero transmission.

Based on these edge states presented above, we construct other topological-insulator junction systems for different types of the FP resonance in figure 7. According to the same analysis in figure 7(a), the spin-down transmission is zero due to the gap effect, while the spin-up transmission corresponds to the FP resonant due to the mismatch of the edge states between the leads and device. In addition, the spin-up current loop is clockwise not shown here, which is different from the counterclockwise current loop in figure 7(f). In figures 8(b) and (c), it is shown that both spin-up and spin-down edge states between the leads and device are mismatched. Following the same analysis, we expect that both spin-up and spin-down transmissions simultaneously have so many peaks shown in figures 8(e) and (f). Beyond our expectation, the separation period $T$ of these transmissions peaks is different between the spin-up and spin-down currents in figures 8(d) and (e). This may result from the anisotropic property of these two chiral edge states: they have different group velocities. So the group velocity mismatch of these edge states (between the device and leads) also has an effect on the separation period.

Here, we give a detailed discussion about the influence of the mismatched group velocity on the separation period. We only compare the spin-down transmission curves from the system in figure 8(c) as our example to illustrate this mechanism. In figure 9(c), it can be easily seen that the separation period $T$ gets smaller with increasing the strength of the off-resonant light in the device, which can be roughly explained by the view of the mismatch of the group velocity. In figures 9(a)–(d), it is shown that the slope values of the band structures in the device get bigger as the strength of the off-resonant light gets stronger. Compared to the band structures of the leads shown in figure 1(c) in this system, the mismatch degree of the group velocity between the device and leads becomes larger with increasing the light intensity, as a result of increasing the separation period $T$. This is a rough interpretation. The detailed mechanism is needed for our further study with the help of the effective Hamiltonian of the zigzag edge states.

In the following, we demonstrate some applications for a switcher by modulating the external fields in zStNR. In figure 10(a), it can be seen that the leads and device share common edge states with the same propagation direction, causing the spin-up and spin-down currents freely pass through the device. This system apparently can transport charge current. Beyond this case, the system can be further transformed

**Figure 10.** Simple applications for the charge and spin current switchers in the zStNR junction system. (a) The edge states are from figures 1(c); (b) the edge states of the leads (device) are also from figure 1(c) (figure 3(c)); (c) the insulator system without the off-resonant light is from figure 2(a).
into a spin filter (figure 10(b)) by applying the staggered electric/ORCP light field on the device, where the leads are nonmagnetic states and the device is antiferromagnetic state. And then the charge current is switched into the 100% spin-polarized current. Finally, by turning off all the external fields in the leads and device, there are no topological edge states in the system, and become the antiamagnetic state with a gap (shown in figure 10(c)). There is no current in the system which leads to an off state.

4. Conclusions

We used the Kane-Mele-Hubbard model to investigate the topological and magnetic phase transitions and transport properties in $z$StNR. Under the ORCP light field, the BI with the edge magnetism is transformed into the QAH state, corresponding to the anisotropic chiral edge state, while the antiferromagnetic state is broken into the nonmagnetic state. Besides the off-resonant light, the electric field induces the topological phase transition from the BI state to SQAH state, corresponding to the 100% spin-polarized edge state, while the antiferromagnetic state is transformed into the reverse antiferromagnetic state. We also investigated the finite-size effect induces magnetic phase transitions, which in turn induces plentiful transitions from the BI state to SQAH state and then QAH state. With the suitable off-resonant light and electric field, the BI state is transformed into the SQAH state, while the antiferromagnetic state is transformed into the reverse case. In addition, we also constructed the topological-insulator junctions with different edge states to investigate the topological transport properties. With the help of the local current distributions, we found that the length of the FP device and mismatch of the group velocities both have an influence on the separation period $T$. The system has also a function of a switcher under the off-resonant light and electric field.

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References

[1] Xu Y, Yan B, Zhang H-J, Wang J, Xu G, Tang P, Duan W and Zhang S-C 2013 Large-gap quantum spin Hall insulators in tin films Phys. Rev. Lett. 111 136804
[2] Zhu F-F, Chen W-J, Xu Y, Gao C-L, Guan D-D, Liu C-H, Qian D, Zhang S-C and Jia J-F 2015 Epitaxial growth of two-dimensional stanene Nat. Mater. 14 1020–5
[3] Qiao Z H, Yang S A, Wang B, Yao Y G and Niu Q 2011 Spin-polarized and valley helical edge modes in graphene nanoribbons Phys. Rev. B 84 035431
[4] Lü X L and Xie H 2019 Spin filters and switchers in topological-insulator junctions Phys. Rev. Appl. 12 064040
[5] Sternativo P and Dolcini F 2014 Effects of disorder on electron tunneling through helical edge states Phys. Rev. B 90 125135
[6] Colomes E and Franz M 2018 Antichiral edge states in a modified Haldane nanoribbon Phys. Rev. Lett. 120 086603
[7] Ezawa M 2013 Photoinduced topological phase transition and a single Dirac-cone state in silicene Phys. Rev. Lett. 110 026603
[8] Ezawa M 2015 Monolayer topological insulators: silicene, germanene, and stanene J. Phys. Soc. Japan 84 1211003
[9] Ezawa M 2012 Topological phase transition and electrically tunable diamagnetism in silicene Eur. Phys. J. B 85 363
[10] Matthes L and Bechtstedt F 2014 Influence of edge and field effects on topological states of germanene nanoribbons from self-consistent calculations Phys. Rev. B 90 165431
[11] Kane C L and Mele E J 2005 Quantum spin Hall effect in graphene Phys. Rev. Lett. 95 226801
[12] Pan H, Li Z, Liu C-C, Zhu G, Qiao Z and Yao Y 2014 Valley-polarized quantum anomalous Hall effect in silicene Phys. Rev. Lett. 112 106802
[13] Liu C C, Feng W X and Yao Y G 2011 Quantum spin Hall effect in silicene and two-dimensional germanium Phys. Rev. Lett. 107 076802
[14] Zheng J, Xiang Y, Li C L, Yuan R Y, Chi F and Guo Y 2021 Multichannel depletion-type field-effect transistor based on ferromagnetic germanene Phys. Rev. Appl. 16 024046
[15] Zheng J, Xiang Y, Li C L, Yuan R Y, Chi F and Guo Y 2020 All-optically controlled topological transistor based on xenes Phys. Rev. Appl. 14 034027
[16] Xu C, Luo G, Liu Q, Zheng J, Zhang Z, Nagase S, Gao Z and Lu J 2012 Giant magnetoresistance in silicene nanoribbons Nanoscale 4 3311–7
[17] Lado J L and Fernandez-Rossier J 2014 Magnetic edge anisotropy in graphenelike honeycomb crystals Phys. Rev. Lett. 113 027203
[18] Lü X-L, Zhang C-X, Wang W-J, Cheng X and Xie H 2019 Excitation and phase transitions of spin density waves in graphene nanoribbons J. Phys.: Condens. Matter 31 455501
[19] Modarressi M, Kuang W B, Kaloni T P, Roknabadi M R and Schreckenbach G 2016 Topological phase in oxidized zigzag stanene nanoribbons AIP Adv. 6 095919
[20] Wong B M, Ye S H and O’Bryan G 2012 Reversible, opto-mechanically induced spin-switching in a nanoribbon-spiropyran hybrid material Nanoscale 4 1321–7
[21] Son Y-W, Cohen M L and Louie S G 2006 Energy gaps in graphene nanoribbons Phys. Rev. Lett. 97 216803
[22] Fu B, Abid M and Liu C-C 2017 Systematic study on stanene bulk states and the edge states of its zigzag stanene New J. Phys. 19 103040
[23] Ildarabadi F and Farghadan R 2021 Fully spin-valley-polarized current induced by electric field in zigzag graphene and germanene nanoribbons Phys. Chem. Chem. Phys. 23 6084–90
[24] Ildarabadi F and Farghadan R 2021 Edge magnetization and spin-valley-caloritronics in germanene and stanene nanoribbons J. Magn. Magn. Mater. 529 167870
[25] Wierzbicki M, Barna´s J and Swirkowicz R 2015 Thermoelectric properties of silicene in the topological- and band-insulator states Phys. B 91 165417
[26] Qi I S, Hu K G and Li X 2018 Electric control of the edge magnetization in zigzag stanene nanoribbons from first principles Phys. Rev. Appl. 10 034048
[27] Lü X-L, Xie Y and Xie H 2018 Topological and magnetic phase transition in silicene-like zigzag stanene nanoribbons New J. Phys. 20 043054
[28] Xie H, Gao J-H and Han D 2018 Excited spin density waves in zigzag graphene nanoribbons New J. Phys. 20 013035
[29] Luo M 2020 Topological edge states of a graphene zigzag nanoribbon with spontaneous edge magnetism Phys. Rev. B 102 075421
[30] Lee K W and Lee C E 2018 Transverse field-induced quantum valley Hall effects in zigzag-edge graphene nanoribbons Phys. Lett. A 382 2137–43
[31] Cao J and Xiong S J 2013 Topological phase transition in a graphene system with a coexistence of Coulomb interaction, staggered potential, and intrinsic spin–orbit coupling Phys. Rev. B 88 085409
[32] Knez I, Du R-R and Sullivan G 2011 Evidence for helical edge modes in inverted InAs/GaSb quantum wells Phys. Rev. Lett. 107 136603
[33] Shevtsov O, Carmier P, Petitjean C, Groth C, Carpentier D and Waintal X 2012 Graphene-based heterojunction between two topological insulators Phys. Rev. X 2 031004
[34] Tian H, Wang S, Hu J and Wang J 2015 The chirality dependent spin filter design in the graphene-like junction J. Phys.: Condens. Matter 27 125005
[35] Liu C-C, Jiang H and Yao Y 2011 Low-energy effective Hamiltonian involving spin-orbit coupling in silicene and two-dimensional germanium and tin Phys. Rev. B 84 195430
[36] Rachel S and Ezawa M 2014 Giant magnetoresistance and perfect spin filter in silicene, germanene, and stanene Phys. Rev. B 89 195303
[37] Kitagawa T, Oka T, Brataas A, Fu L and Demler E 2011 Transport properties of nonequilibrium systems under the application of light: photoinduced quantum Hall insulators without Landau levels Phys. Rev. B 84 235108
[38] Lü X-L and Xie H 2020 Bipolar and unipolar valley filter effects in graphene-based P/N junction New J. Phys. 22 073003
[39] Wierzbicki M, Barna´s J and Swirkowicz R 2015 Zigzag nanoribbons of two-dimensional silicene-like crystals: magnetic, topological and thermoelectric properties J. Phys.: Condens. Matter 27 485301
[40] Farghadan R and Yoosefi M 2016 Magnetism and spin transport of carbon chain between armchair graphene nanoribbon electrodes Physica E 83 414–9
[41] Lee D H and Joannopoulos J D 1981 Simple scheme for surface-band calculations. I Phys. Rev. B 23 4988–96
[42] Sancho M P L, Sancho J M L and Rubio J 1984 Quick iterative scheme for the calculation of transfer-matrices—application to Mot(100) J. Phys. F: Met. Phys. 14 1205–15
[43] Zhang Y, Chen S G and Chen G H 2013 First-principles time-dependent quantum transport theory Phys. Rev. B 87 085110
[44] Wang Y, Yan C-Y, Frauenheim T, Chen G H and Niehaus T A 2011 An efficient method for quantum transport simulations in the time domain Chem. Phys. 391 69–77
[45] Xie H, Cheng X and Lü X-L 2020 Steady and dynamic magnetic phase transitions in the open interacting quantum dots arrays J. Magn. Magn. Mater. 497 165867
[46] Power S R, Thomsen M R, Jauho A P and Pedersen T G 2017 Electron trajectories and magnetotransport in nanopatterned graphene under commensurability conditions Phys. Rev. B 96 075425
[47] Stegmann T and Szpak N 2019 Current splitting and valley polarization in elastically deformed graphene 2D Mater. 6 015024
[48] Calvo M R et al 2017 Interplay of chiral and helical states in a quantum spin Hall insulator lateral junction Phys. Rev. Lett. 119 226401