Electrostatically Controlled Magnetization Rotation in Ferromagnet-Topological Insulator Planar Structures

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Abstract

An approach to the electrostatic control of $90^\circ$ magnetization rotation in the hybrid structures composed of topological insulators (TIs) and adjacent ferromagnetic insulators (FMI) is proposed and studied. The concept is based on TI electron energy variation with in-plane to put-of-plane FMI magnetization turn. The calculations explicitly expose the effect of free energy variability in the form of the electrically controlled uniaxial magnetic anisotropy, which depends on proximate exchange interaction and TI surface electron density. Combining with inherent anisotropy, the magnetization rotation from in-plane to out-of-plane direction is shown to be realizable for 1.7 $\sim$ 2.7 ns under the electrical variation of TI chemical potential in the range $\pm$ 100 meV around Dirac point. When bias is withdrawn a small signal current can target the out-of-plane magnetization instable state to the desirable direction of in-plane easy axis, thus the structure can lay the foundation for low energy nonvolatile memory prototype.

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Because of the absence of the natural magnetic charges and thus the leakage conditions, magneto-electronic automata presents a great advantage of the magnetic devices compared with their electrical counterparts in terms of the non-volatility. Up to date the magnetization driving force in such devices is realized via a magnetic field of the channel current or current-induced angular momentum transfer in magnetic tunnel junctions. The current-assisted data processing in both techniques consumes much energy, the reduction of which is one of the major concerns for present industrial developments and academic researches.

Electric field control in addition to the current induced effects would has been pursued with multiferroic materials. The tension variation in some FM materials can rotate their easy axis 90° jointly with magnetization vector $\mathbf{M}$. In FM/piezoelectric hybrid structures such effect evokes the new paradigm of magnetization rotation in the electrical field. However it should be noted that the desirable piezo-effect in such structures requires large gate electric field, which indicates the need of searching for much stronger magnetoelectric effects in modern magneto-electronics.

Another approach relies on the effective magnetic fields $\mathbf{H}_{\text{ex}}$ mediated by the proximate exchange interactions between surface magnetic ions and itinerant electrons (or holes) in semiconductors, graphene, or topological insulators. A distinctive feature of the $\mathbf{H}_{\text{ex}}$ consists in its dependence on electron density at the interface of the electric channel and magnetic layers, rather than on an electric current. This property of the $\mathbf{H}_{\text{ex}}$ combines the magnetic and electric responses inside the new "meta-materials" in a strongly correlated manner that stems from the electrostatic and quantum mechanical nature of the exchange interaction.

Recent studies have demonstrated the electrical control efficiency of the proximate exchange interaction (PEI) between a magnetic layer and graphene for different spintronic applications. As an example, the conductivity of the graphene sandwiched between ferromagnetic insulators (FMIs) reveals the giant magnetic resistance. On the other hand, the magnetization of FMI can be turned and even switched by electrical variation of the carriers density in adjacent graphene.

In present letter we explore the new prospects granted by topological insulators (TIs) and composite structures (Fig. 1) based on TIs. Specific spin-momentum coupling in 3D TI’s surface Dirac fermions can be naturally incorporated into magneto-electric effects of TI/FMI composite structures. As a result, a lot of unusual physical phenomena have been
predicted in TIs affected by adjacent FMIs. Here we focus on converse TI effect on adjacent FMI. Namely, the PEI of TI surface electrons is shown to appear in the form of FMI effective magnetic anisotropy electrically driven without any structure tension. Combining with inherent anisotropy and anisotropy of the sample shape, the in-plane to out-of-plane orientation transition can be electrically stimulated by variance of TI-mediated anisotropy. This effect is quantitatively demonstrated for uniaxial nanomagnet deposited on the surface of TI (e.g., Bi$_2$Se$_3$). It is remarkable that the system with perpendicular $M$ appears in neutral position with respect to $M$ returning to either of the in-plane states along easy axis. This instability can be used for directional targeting of $M$ by a current induced torque. Small signal current is sufficient to tilt the magnetization to the desired direction and the following switch will be driven by the shape and intrinsic anisotropy terms. The possible design of nonvolatile memory and logic devices can be realized on basis of these magnetic units.

Qualitatively the easy axis rotation can be thought as an effect of the electron energy reducing under exchange interaction with proximate FMI. When the magnetization is in the in-plane direction, the PEI results in a shift of the Dirac cones in the momentum space whereas an energy gap can be generated and the energy dispersion becomes non-linear for the out-of-plane orientation with $M_z \neq 0$. From the thermodynamic point of view, the case with a band gap ($M_z \neq 0$) is expected to be energetically more favorable with lower electronic free energy than that with the shifted bands ($M_z = 0$) following a similar consideration demonstrated in bilayer graphene. A departure of chemical potential $\mu$ from the energy level of the Dirac point $\varepsilon_D$ suppresses this tendency because of energy increase of the conduction electrons partially compensates valence band energy decrease while gap opening for $\mu > \varepsilon_D$ and fewer active electrons for $\mu < \varepsilon_D$.

To proceed the quantitative analysis let us consider a thermodynamical potential of TI electrons interacting with FMI,

$$\Phi_e(\mu) = -T \sum_{b,k} \ln \left( 1 + \exp \frac{\mu - \varepsilon_D - \varepsilon_{b,k}}{T} \right)$$

where the sum involves electronic bands $b$ and wave vector $k$, $T$ is the temperature in energy units. We are interested in alteration of $\Phi_e(\mu)$ with the Fermi level varying near the Dirac point of the surface states. Thus the rest of electronic states in Eq. (1) are irrelevant and summation in $\Phi_e(\mu)$ can be safely restricted by the spectrum of surface Dirac fermions with
2D effective Hamiltonian

\[ H = \hbar v_F [\sigma \times \mathbf{k}] \hat{z} + D k^2 + \mathbf{G} \sigma \]  

(2)

where \( v_F \) is the electron Fermi velocity near Dirac point, \( \sigma \) the Pauli matrix of electron spin, \( D \) is the material parameter for quadratic term. The remaining term in Eq. (2) describes the energy of an electron spin in the exchange effective fields (in units of energy) \( \mathbf{G} = G \mathbf{m} \) of the proximate FMI, \( \mathbf{m} = \mathbf{M} / |\mathbf{M}| \) ensures the collinearity of the effective field and FMI magnetization. Introducing the angles \( \theta \) between \( \hat{z} \) and \( \mathbf{m} \) and \( \varphi \) between \( \mathbf{k} \) and planes \( (\mathbf{m}, \hat{z}) \) (Fig. 1) the energy spectrum of Hamiltonian (2) can be expressed as

\[ \varepsilon_{b,k} = D_0 p^2 + b \sqrt{p^2 + G^2} + 2Gp \sin \theta \sin \varphi, \]  

(3)

where \( p = \hbar v_F k \) is the electron surface momentum in energy unit, \( D_0 = D / \hbar^2 v_F^2 \) and \( b = \pm 1 \) marks the conduction and valence bands. Equation (3) clearly illustrates the band structure reconstruction lowering the energies of electrons near the top of valence band when \( \mathbf{M} \) turns from in-plane direction (\( \theta = \pi/2 \)) to vertical position (\( \theta = 0 \)). The valence band electrons \( (b = -1) \) contribute to this effect while conduction band electrons \( (b = 1) \) may reduce it. Thus the proximity with TI plays the role of a mediator enhancing the uniaxial FMI anisotropy along the normal \( \hat{z} \).

Following Eqs. (3) and (1), \( \Phi_e(\mu) \) is an even function of \( G_x = G \sin \theta \) that determines the dependence \( \Phi_e(\mu) = -K_{\text{eff}}(\mu) \sin^2 \theta \) and proportionality \( K_{\text{eff}}(\mu) \sim G^2 \) in low-order expansion on \( G \). Similar relation remains for \( \Delta \Phi_e(\mu) = \Phi_e(\varepsilon_D) - \Phi_e(\mu) \), which reflects the variable part of the TI mediated anisotropy energy

\[ \Delta \Phi(\mu) = A \Delta K_{\text{eff}}(\mu) \sin^2 \theta. \]  

(4)

Here \( A \) is the area of FMI/TI interface and

\[ \Delta K_{\text{eff}}(\mu) = f(\mu, T) D_0 G^2, \]  

(5)

where \( D_0 \) is introduced for convenience sake. Computation of Eq. (1) with Eq. (3) confirms the validity of the definition in Eqs. (4) and (5). Function \( f(\mu, T) \) can be numerically found in terms of Eqs. (1) and (3); Fig. 2 depicts the results of calculations. The results indicates the amplitude of anisotropy changes with surprisingly minor temperature effect but significantly asymmetry in chemical potential variation over conduction band \( (\mu - \varepsilon_D > 0) \) and valence
band \((\mu - \varepsilon_D < 0)\). Consequently, Eq. (5) predicts the variation range of anisotropy as large as 0.5 meV/nm² provided the strength of exchange interaction is \(G = 40\) meV.\(^{10}\) In such a case a contact area section of tens nanometers may supply the anisotropy energy change of eV-scale that guarantees the nonvolatile magnetic bi-stability at room temperature.

Whether the variation of \(\Delta K_{\text{eff}}(\mu)\) is capable to cant over \(M\) depends on symmetry and magnitude of total magnetic energy \(F\) of the nanomagnet. It is convenient to approximate the actual shape of the FMI (e.g. rectangular parallelepiped) by triaxial ellipsoid with homogeneous magnetization over the whole volume \(V\). At such approach the free energy can be simplified to normal form

\[
F = 2\pi V \sum_{i,j} N_{ij} M_i M_j + U_{an}(\mathbf{m}),
\]

where first term describes the energy of demagnetizing field, and last one represents the magnetic anisotropy, which includes independent on \(\mu\) part \(\Phi_e(\varepsilon_D)\) of PEI and tension-mediated term with uniaxial anisotropy, \(N_{ij}\) the components of demagnetizing tensor. In the principal coordinates \(i, j = x, y, z\), this tensor becomes diagonal, \(N_{ij} = N_i \delta_{i,j}\), with principal values

\[
N_i = \frac{a_x a_y a_z}{2} \int_0^\infty dq \frac{dq}{Q_i \sqrt{Q_x Q_y Q_z}}
\]

where \(\sum_i N_i = 1\) and \(Q_i = q + a_i^2\) with lengths \(a_x, a_y, a_z\) of the ellipsoid principal axes. If they coincide with crystalline axes, the free energy per unit volume can be brought into canonical form

\[
\frac{F}{V} = 2\pi M^2 \sum_{i=x,y,z} N_i m_i^2 + u_{an}(\mathbf{m})
\]

with \(u_{an}(\mathbf{m}) = U_{an}(\mathbf{m})/V\). In the case of planar FMI/TI structure (i.e. \(N_z >> N_x, N_y\)) the demagnetizing field tends to establish the hard axis along normal \(\hat{z}\) to interface. So, in soft magnetic materials like Yttrium Iron Garnet, the proper choice of FMI sizes can ensure the electrical switch between in-plane to out of plane \(M\) direction. This condition, however, is not enough to guarantee the nonvolatile bistability because there is a path that connects the states \(M_x\) and \(-M_x\) through the saddle points \(\pm M_y\). Apparently demagnetizing field can also establish the in-plane anisotropy in the structures with shorter FMI layer length in \(y\)-direction, but desirable retention time with saddle point height \(\sim 2\pi M^2(N_y - N_x)V = 1\) eV can be reached at least in sub-micron device sizes (200 -300 nm) as \(N_y\) is still relatively small compared with \(N_z\).
An example of the alternative approach assumes the FMI with intrinsic uniaxial crystalline anisotropy in form of hard $y$-axis. In such a case the anisotropy term is

$$u_{an}(m) = K_1 m_y^2 + [K_u + \frac{1}{\hbar^2} \Delta K_{\text{eff}}(\mu)](1 - m_z^2), \quad (9)$$

where $K_1 > 0$, $K_u$ accumulates the effects of crystal lattice distortion, which, in turn, vary with crystal doping, the strain at interfaces with TI and spacer layer attached on opposite sides, and PEI accounted for $\Phi_e(\varepsilon_D)$. Thus, this parameter depends on particular device implementation.

The total magnetic energy expressed in spherical coordinate of $m$ (i.e., $m_x = \sin \theta \cos \varphi; m_y = \sin \theta \sin \varphi; m_z = \cos \theta$, see Fig. 1) is shown in Fig. 3(a) and 3(b) for two chemical potentials. The graph clearly illustrates the alternation of energy valleys and energy hills located along in-plane directions, $\theta = \pi/2$ and $\varphi = 0$ or $\pi$, [Fig. 3(a)] and in the normal directions, $\theta = 0$ or $\pi$, [Fig. 3(b)], as chemical potential changes. A small perturbation of unstable states at tops of energy hills may guide the following relaxation into desirable valley. For such perturbation we suggest to evoke a proximity induced effective magnetic field $B_{\text{eff}} = [\mathbf{J} \times \hat{z}] G/ev_F L_z M$ ($e$ is the electron charge, $L_z$ is the FMI thickness) which accompany the spin polarized TI surface current $\mathbf{J}$. This field can be used to choose the path to the desired valley by the current direction. Figure 3(c) and (d) illustrates the distortion of energy relief under small electric current. All together Fig. 3(c)-(d) exhibit the possibility of the electric field control of magnetization turn and even switch provided appropriate handling by bias voltage and signal current.

The dynamics of magnetization turn determines the switching speed. A common approach to this problem is based on the solution $\mathbf{m} = \mathbf{m}(t)$ of the Landau-Lifshitz-Gilbert equation, which, in the case of mono-domain approximation with homogeneous magnetization, takes the form:

$$\frac{\partial \mathbf{m}}{\partial t} = -\gamma \mathbf{m} \times \mathbf{H}_{\text{eff}} + \alpha \mathbf{m} \times \frac{\partial \mathbf{m}}{\partial t}. \quad \text{(10)}$$

where $\gamma$ is the gyromagnetic ratio, $\alpha$ the Gilbert damping factor and the effective field $\mathbf{H}_{\text{eff}} = (\mu_0 M_0 V)^{-1} \partial F/\partial \mathbf{m}$ stems from the total free energies $F$ [Eq. (6)]. Note that $\mathbf{H}_{\text{eff}}$ crucially depends on the magnetization. For example, if the initial magnetization $\mathbf{m}(0)$ perfectly aligned with $x$ axis and ignoring fluctuation, the $\mathbf{H}_{\text{eff}}$ and $\mathbf{m}$ are collinear and the system comes into frustrating (metastable) state with respect to subsequent relaxation to the $z$ or $-z$ direction. However any thermal dispersion of $\mathbf{m}(0)$ may work as a start point.
for magnetization switch. To show this effect we consider evolution $\mathbf{m}(t)$ starting from $\theta$ and $\varphi$ that deviate from idealistic case $\theta = 90^\circ$ and $\phi = 0^\circ$ on only $0.1^\circ$. The resulting switch is shown in Fig. 4. The switch is gradually accelerated till $t = 1$ $ns$, after which the $\mathbf{H}_{\text{eff}}$ gathers strength and steep switch occurs with oscillation. The actual switch time is about $1.7$ $ns$ marked by $m_z$ changing from 0.1 (point B) to 0.9 (point A), while it may take around $1$ $ns$ for the switch to be initiated.

To see the dispersion of switch time, we initializes the random array of magnetization states according to Boltzmann distribution at room temperature. The dispersion of the shift is marked by the square point A (when $m_z \geq 0.9$) in Fig.4 and the range is indicated by the corresponding arrow. This means the total $90^\circ$ switch time varies from $1.7$ $ns$ to $2.7$ $ns$ because of thermal fluctuation. One possible solution to the undetermined switch time is to use a TI surface current to initiate the switch by current induced torque. Thus $180^\circ$ switches can be achieved.

In conclusion, a magneto-electric effect in form of electrically controlled magnetic anisotropy has been presented for the structure taking advantage of the proximity exchange interaction between a TI and a FMI layer. The effect is based on exchange energy correlation with weak temperature dependence and it is more respondent for the change in valence band than that in the conduction band. This interaction effectively induces a change of the perpendicular anisotropy energy of the FMI layer which results in a thermally initiated magnetization rotation from in-plane to out-of-plane. The initiating time can be as long as $1.4$ $ns$ if well aligned and the switch time is about $1.7$ $ns$. A small signal current can be applied to both initiating the switch and to determine the magnetization after withdrawing the bias by tilting the magnetization to the desired direction at the beginning of the relaxation.

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FIG. 1. (Color online) Schematic illustration of the structure composed of ferromagnetic insulator layer (FMI) deposited on topological insulator (TI). The controlling gate is not shown. The $x, y, z$ are the reference frame used in the paper. The easy axis (two-sided arrow) is along $x$-direction without gate voltage. The $\theta$ and $\varphi$ are the spherical coordinates of the off-balance magnetization vector $\mathbf{M}$. A signal current can be applied to the TI surface to tilt the magnetization.
FIG. 2. (Color online) Parameter $f(\mu, T)$ reflecting the TI mediated anisotropy as a function of chemical potential $\mu$ at three temperatures 300 K (curve 1), 200 K (curve 2) and 77 K (curve 3). Calculations use the parameters of Bi$_2$Se$_3$ for Hamiltonian (2): $v_F = 6 \cdot 10^7$ cm/s; $D = 13.0$ eV·Å$^2$. 
FIG. 3. (Color online). Energy landscape in polar coordinates $\theta$ (radial variable) and $\phi$ (azimuthal variable, Fig. 1) of magnetization direction. The equal-energy curves are calculated for sizes of FMI planar layer $50 \times 50 \times 2$ nm$^3$, $M = 4\pi \times 160$ G, $K_1 = K_u = 0.05$ meV/nm$^3$ (a) Chemical potential and signal current are $\mu - \varepsilon_D = -90$ meV and $J = 0$. Two symmetrical valleys correspond to $\mathbf{M}$ orientations along easy axis $x$. (b) Easy axis is turned to $z$-direction with $\mu - \varepsilon_D = 0$ and $J = 0$. (c) Signal current shifts the valleys and tilts magnetization to $-x$ direction ($\mu - \varepsilon_D = 0$ meV and $J = 1 \mu$A/nm) that drives the $\mathbf{M}$ relaxation to the reversal direction pointed at figure (d) with ($\mu - \varepsilon_D = -90$ meV and $J = 1 \mu$A/nm).
FIG. 4. (Color online) Anisotropy driven switching trajectory. The initial magnetization is set to $\theta = 89.9^\circ$, $\phi = 0.1^\circ$ to enable the initiation of the switch. Square B ($m_z = 0.1$) and square A ($m_z = 0.9$) mark the starting and ending of the switch respectively. The switching time is around $1.7 \text{ ns}$ while the initiating time can be as long as $1.4 \text{ ns}$ due to thermal fluctuation at $300 \text{ K}$. The finishing point A thus varies in the range indicated by the double arrow.