Dissipationless Multiferroic Magnonics

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We propose that the magnetoelectric effect in multiferroic insulators with coplanar antiferromagnetic spiral order, such as BiFeO₃, enables electrically controlled dissipationless magnonics. Applying an oscillating electric field in these materials with frequency as low as household frequency can activate Goldstone modes that manifests fast planar rotations of spins, whose motion is not obstructed by crystalline anisotropy. Combining with spin ejection mechanisms, such a fast planar rotation can deliver electricity at room temperature over a distance of the magnetic domain, which is free from the energy loss due to Gilbert damping.

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Introduction.- A primary goal of spintronic research is to seek for mechanisms that enable electric (E) field controlled spin dynamics, since, in practice, E fields are much easier to manipulate than magnetic (B) fields. As spins do not directly couple to E field, incorporating spin-orbit coupling seems unavoidable for this purpose. Along this line came the landmark proposals such as spin field effect transistor [1] and spin-orbit torque [2–5], the realizations of which suggest the possibility of spin dynamics with low power consumption. On the other hand, in another major category of spintronics, namely magnonics, which can deliver electricity with frequency as low as household frequency up to the range of magnetic resonance, which aims at the generation, propagation, and detection of magnons, a mechanism that enables electrically controlled magnonics without the aid of a magnetic field has yet been proposed.

Raman scattering experiments on BiFeO₃ (BFO) shed light on this issue. The magnetic order of BFO is a canted antiferromagnetic (AF) spiral on the plane spanned by the electric polarization P along [111] and one of the three symmetry-equivalent wave vectors on a rhombohedral lattice [6, 7]. The spins have only a very small out-of-plane component [8, 10]. Applying a static E field ~ 100kV/cm significantly changes the cyclon (in-plane) and extra-cyclon (out-of-plane) magnons because of the magnetoelectric effect [7]. Indeed, spin-orbit coupling induced magneto-electric effects are a natural way to connect E field to the spin dynamics of insulators [12, 13]. Motivated by the Raman scattering experiments on BFO, in this Letter we propose that applying an oscillating E field to a coplanar multiferroic insulator (CMI) that has AF spiral order can achieve electrically controlled dissipationless magnonics, which can deliver electricity with frequency as low as household frequency up to the range of magnetic domains. Compared to the magnonics that uses B field, microwave, or spin torques to generate spin dynamics in prototype Y₃Fe₅O₁₂ (YIG) [14–16], the advantage of using CMI is that a single domain sample up to mm size is available [17], and Raman scattering data indicate well-defined magnons in the absence of B field [7]. The dipole-dipole interaction is negligible because of the AF order, and we show that crystalline anisotropy is generally reduced.

Spin dynamics in CMI.- We start from the AF spiral on a square lattice shown in Fig. I(a), described by

\[ H = \sum_{i,\alpha} J S_i \cdot S_{i+\alpha} - D_\alpha \cdot (S_i \times S_{i+\alpha}) \]  (1)

where \( \alpha = \{a,c\} \) are the unit vectors defined on the \( xz \)-plane, \( J > 0 \), and \( D_\alpha = D_\alpha \hat{y} > 0 \) is the Dzyaloshinskii-Moriya (DM) interaction. The staggered moment \( (-1)^{i} S_i \) in the ground state shown in Fig. I (a) is characterized by the angle \( \theta_\alpha = \hat{Q} \cdot \hat{\alpha} = -\sin^{-1} (D_\alpha / \hat{J}_\alpha) \) between neighboring spins, where \( \hat{J}_\alpha = \sqrt{J^2 + D_\alpha^2} \). The DM interaction

\[ D_\alpha = D_\alpha^0 + \pi \alpha E \times \alpha \]  (2)

can be controlled by an E field [18], where \( D_\alpha^0 \) represents the intrinsic value due to the lack of inversion symmetry of the \( \alpha \)-bond. In the rotated reference frame \( S' \) defined by

\[ S'^{\beta}_{ix} = S_i \cos Q \cdot r_i + S_i^z \sin Q \cdot r_i, \]

\[ S'^{\beta}_{iy} = -S_i^z \sin Q \cdot r_i + S_i^x \cos Q \cdot r_i, \]

and \( S'^{\beta}_{iz} = S_i^y \), the Hamiltonian is

\[ H = \sum_{i,\alpha} \hat{J}_\alpha (S'^{\beta}_{ix} S'^{\beta}_{ix} + S'^{\beta}_{iy} S'^{\beta}_{iy}) + J S'^{\alpha}_{iz} S'^{\alpha}_{iz} . \]  (4)

Since \( \hat{J}_\alpha = J \), the spins have collinear AF order and all \( S'^{\beta}_{ix} = (-1)^i S \) lie in \( xz \)-plane.

The spin dynamics in the absence of B field is governed by the Landau-Lifshitz-Gilbert (LLG) equation

\[ \frac{dS'_i}{dt} = \frac{\partial H}{\partial S'_i} \times S'_i + \alpha G S'_i \times \left( \frac{\partial H}{\partial S'_i} \times S'_i \right) \]  (5)

expressed in the \( S' \) frame, where \( \alpha G \) is the phenomenological damping parameter. Eq. (5) can be solved by the
Ignoring the damping term in Eq. (5) yields eigenenergies
\[ \frac{\omega_k \pm 1}{2S} = \left[ \left( \sum_\alpha \tilde{J}_\alpha \pm \gamma_\alpha (k) \right)^2 - \left( \sum_\alpha \gamma_\alpha (k) \right)^2 \right]^{1/2}, \tag{7} \]
where \( \gamma_\alpha (k) = \left( \tilde{J}_\alpha / 2 \pm J / 2 \right) \cos k \cdot \alpha \). Their eigenvalues and eigenvectors near \( k = (0,0) \) and \( k = (\pi, \pi) \) are summarized below
\[ \left\{ \omega_k^{\pm (0,0)}, \omega_k^{\pm (\pi, \pi)} \right\} = 2S \sqrt{2(D_a^2 + D_c^2)}, \tag{8} \]

The in-plane magnon \( dS_t^i / dt = (dS_t^x / dt, 0, 0) = 0 \) is gapless, while the out-of-plane magnon \( dS_t^y / dt = (0, dS_t^y / dt, 0) \) develops a gap, as displayed in Fig. 1 (c). Even including the damping term in Eq. (4), the in-plane magnons very near the Goldstone modes \( \omega_{-k}^{(0,0)} \) and \( \omega_{k}^{(\pi, \pi)} \) remain unchanged and damping-free. Away from the Goldstone limit, the eigenenergies become complex, hence the magnons are subject to the damping and decay within a time scale set by \( \gamma_\alpha G^{-1} \). Note that Eq. (4) is identical to the easy-plane XXZ model, so the magnons therein should have the same properties.

**Spin dynamics induced by oscillating E field.** We analyze now the spin dynamics in the damping-free in-plane magnon channel induced by magnetoelctric effects (Eq. 1). Unlike the spin injection by using the spin Hall effect (SHE) to overcome the damping torque \( \Omega \), our design does not require an external B field, and is feasible over a broad range of frequencies. Consider the device shown in Fig. 2 where an oscillating electric field \( E = E^0 \cos \omega t \) is applied parallel to the ferroelectric moment over a region of length \( L = Na \), such that the DM interaction in Eq. (2) oscillates in this region. Thus, the wave length of the spiral changes with time yielding an oscillation of the number of spirals inside this region,
\[ n_Q = \frac{L}{2\pi / |Q|} = \frac{N}{2\pi} \left[ D_0^0 + \mp E^0 a \cos \omega t \right], \tag{9} \]
assuming \( D_0 = D_a^0 + \mp E^0 a \ll J, D_c = 0 \), and \( E \perp a \). Suppose the spin \( S_0 \) at one boundary is fixed by, for instance, surface anisotropy because of specific coating. Then \( S_N \) at the other boundary rotates by
\[ \frac{\partial \theta_N}{\partial t} = -\frac{N}{J} E^0 a \omega \sin \omega t, \tag{10} \]
because whenever the number of waves \( n_Q \) changes by 1, \( S_N \) rotates \( 2\pi \) in order to to wind or unwind the spin texture in the E field region. The significance of this mechanism is that although the E field is driven by a very small frequency \( \omega \), the spin dynamics \( \partial \theta_N \) at the boundary is many orders of magnitude enhanced because of the winding process. The rotation of \( S_N \) serves as a driving force for the spin dynamics in the field-free region from \( S_N \) to \( S_{N+M} \). As long as the spin dynamics is slower than the energy scale of the DM interaction \( \partial \theta_N \sim |D_0| / h \sim THz \), one can safely consider the E field region as adiabatically changing its wave length but remaining in the ground state. The spins in the field-free region rotate coherently \( \partial \theta_N = \partial \theta_{N+1} = ... = \partial \theta_{N+M} \), synonymous to exciting the \( \omega_{-k}^{(0,0)} \) mode in Eq. (4), hence the spin dynamics in the field-free region remains damping-free.

We now examine the effect of crystalline anisotropy on the spin dynamics. The square lattice has 4-fold degenerate anisotropy energy \( H_{an} = \sum_i K \cos 4\theta_i \). When the \( k = (0,0) \) Goldstone mode is activated with an angular velocity \( \Omega \) (time-dependent if one uses Eq. (11)), the planar angle of the spin is \( \theta_i = Q \cdot r_i + \Omega t + (-1)^i \pi \). The total anisotropy energy per cross section channel at
where $\lambda = 8\pi a$, which is a periodic function of the number of longitudinal sites $N_L$, and is bounded by $-K\lambda/4\pi a < \langle H_{an}(t) \rangle < K\lambda/4\pi a$ at any time $t$. This result suggests that the spiral order acts to reduce the effect of the crystalline anisotropy such that $\langle H_{an} \rangle$ is not extensive in $N_L$. Including a random fluctuation in the spin texture $\delta\theta \rightarrow \theta + \delta\theta$, where $\delta\theta$ is a random small angle, results in fluctuations in Eq. (11), but the conclusion remains valid. Since typically $K\lambda/4\pi a \ll \delta\lambda$, the crystalline anisotropy will not obstruct the coherent rotation of the spin texture driven by the boundary spin dynamics in Eq. (11).

\[
\langle H_{an}(t) \rangle = K \sum_{i=1}^{N_L} \cos(4Q \cdot r_i + 4\Omega t)
\]

\[
\approx \frac{K\lambda}{8\pi a}[\sin(4QN_L a + 4\Omega t) - \sin(4Q a + 4\Omega t)]
\]

where $J_{sf}$ is the spin relaxation time in the NM. In equilibrium, we assume $\mathbf{m}$ hybridizes with each $\mathbf{S}_i$ on the spiral texture locally. If $\mathbf{S}_i$ varies slowly, $\mathbf{m}$ follows $-\mathbf{S}_i$ with a very small deviation $\mathbf{m} = \mathbf{m}_0 + \delta\mathbf{m} = -n_0 \hat{S}_i + \delta\mathbf{m}$, where $n_0$ is the local equilibrium spin density. The spin current tensor is $\mathbf{J}_s = -D_0 \nabla \times \delta\mathbf{m}$, where $D_0$ is the spin diffusion constant. One obtains

\[
\delta\mathbf{m} = -\frac{\tau_{ex}}{1 + \xi^2} \left\{ -\xi n_0 \frac{\partial \hat{S}_i}{\partial t} - n_0 \hat{S}_i \times \frac{\partial \hat{S}_i}{\partial t} \right\}
\]

FIG. 2: (color online) Experimental proposal of using oscillating $\mathbf{E}$ field to induce spin dynamics in CMI. The AF spiral order is shown in the $S'$ frame. The $\mathbf{E}$ field is applied between $S'_0$ and $S'_N$, causing dynamics in the whole spin texture. Two ways for spin ejection out of $S'_{N+M}$ are proposed: (a) Using SHE to convert it into a charge current. (b) Using time-varying spin accumulation and inductance.

Spin ejection and delivery of electricity.- We now address the spin ejection from the CMI to an attached normal metal (NM). A spin current is induced in the NM when a localized spin $\mathbf{S}_i$ at the NM/CMI interface rotates [16, 26]. Defining the conduction electron spin $\mathbf{m}(r, t) = -\langle \sigma \rangle / 2$, the s-d coupling at the interface $H_{sd} = \Gamma \sigma \cdot \mathbf{S}_i$ defines a time scale $\tau_{ex} = \hbar/2S|\Gamma|$, with $\Gamma < 0$ [26]. The Bloch equation in the NM reads

\[
\frac{\partial \mathbf{m}}{\partial t} + \nabla \cdot \mathbf{J}_s = \frac{1}{\tau_{ex}} \mathbf{m} \times \mathbf{S}_i - \frac{\delta \mathbf{m}}{\tau_{sf}}
\]

where $\mathbf{J}_s = J_{s}^{NM} \otimes \sigma \hbar/2$ is the spin current tensor, and $\tau_{sf}$ is the spin relaxation time in the NM. In equilibrium,

\[
\mathbf{m} = \mathbf{m}_0 + \delta\mathbf{m} = -n_0 \hat{S}_i + \delta\mathbf{m}
\]

where $\xi = \tau_{ex}/\tau_{sf} < 1$ so one can drop the first term on the right hand side, and replace $\hat{S}_i \times \partial_t \hat{S}_i \rightarrow \delta(r) \hat{S}_i \times \partial_t \hat{S}_i$ since $\hat{S}_i$ is located at the NM/CMI interface $r = 0$ ($r$ as coordinate perpendicular to the interface). The resulting equation solves the time dependence of $\delta\mathbf{m}$. Away from $r = 0$, Eq. (12) yields $D_0 \nabla^2 \delta\mathbf{m} = \delta\mathbf{m}/\tau_{sf}$, which solves the spatial dependence of $\delta\mathbf{m}$. The spin current caused by a particular $\mathbf{S}_i$ then follows

\[
J_{s}^{NM} \cdot \delta\mathbf{m} = \delta\mathbf{m} \cdot \mathbf{D}_0 - \frac{\tau_{ex} n_0 D_0}{(1 + \xi^2) \lambda N} \hat{S}_i \times \frac{\partial \hat{S}_i}{\partial t} e^{-r/\lambda N}
\]

and is located at the NM/CMI interface

\[
\lambda_N = \sqrt{D_0 \tau_{sf}}, \text{ similar to results obtained previously [16]}.
\]

If only the in-plane Goldstone mode is excited, as shown in Fig. 2, it is equivalent to a global rotation of spins $\mathbf{S}_i = (-1)^i(\sin(\theta(t) + Q \cdot r_i), 0, \cos(\theta(t) + Q \cdot r_i))$ in the field-free region. Thus the time dependence in Eq. (13), $\mathbf{S}_i \times \partial_t \mathbf{S}_i = \dot{\mathbf{y}} \partial \theta/\partial t$, is described by Eq. (10), and is the same for every $\mathbf{S}_i$ at the NM/CMI interface, even though each $\mathbf{S}_i$ point at a different polar angle. In other words, the spin current ejected from each $\mathbf{S}_i$ of the AF spiral, described by Eq. (12), is the same, so a uniform spin current flows into the NM.

We propose two setups to convert the ejected spin current into an electric signal. The first device uses inverse SHE in a NM deposited at the side of the spiral plane, yielding $\delta\mathbf{m}$ perpendicular to $\mathbf{J}_s^{NM}$ and consequently a voltage in the transverse direction, as shown in Fig. 2 (a). The second design ejects spin into a NM film deposited on top of the spiral plane, as shown in Fig. 2 (b), causing $\delta\mathbf{m}$ parallel to $\mathbf{J}_s^{NM}$.
in the NM develops and oscillates with time, producing an oscillating magnetic flux $\Phi_d$ through a coil that wraps around the NM, hence a voltage $E = -\partial \Phi_d / \partial t$.

Experimental realizations. - The Raman scattering data on BFO [2] show that applying $|E| \sim 100$ kV/cm can change the spin wave velocity by $\delta v_0 / v_0 \sim 1\%$. We can make use of this information to estimate the field-dependence $\pi$ in Eq. (2). The $\omega_k$ mode in Eq. (17) near $k = (0, 0)$ is

$$\omega_{k=0}^{-} = 2\sqrt{2} SJk a \left[ 1 + \frac{5 \left( D_0^2 k^2 + D_1^2 k^2 \right)}{16 J^2 k^2} \right]$$

$$= (v_0 + \delta v_0) k,$$

where $v_0 = 2\sqrt{2}SJa$ is the spin wave velocity in the absence of DM interaction. Assuming $D_0 \neq 0$, $D_1 = 0$, and $E \perp a$, the Raman scattering data gives $\pi \sim 10^{-18}$ C/m. We remark that a coplanar magnetic order can be mapped into a spin superfluid [21, 22] by

$$(S_i) = S (\sin \theta_0, 0, \cos \theta_0) = \sqrt{E} (\text{Im} \psi_i, 0, \text{Re} \psi_i),$$

where $v$ is the volume of the 3D unit cell. Within this formalism, the E field can induce quantum interference of the spin superfluid via magnetoelectric effect, in which the electric field vector $\Phi_E = \oint E \times dl$ is quantized [23, 24]. The flux quantum is $\Phi_E^0 = 2\pi J / \pi$, which is $\Phi_E^0 \sim 1$ V for BFO, close to that ($\sim 10$ V) obtained from current-voltage characteristics of a spin field-effect transistor [22], indicating that strong spin-orbit interaction reduces the flux quantum to an experimentally accessible regime. For instance, BFO has a spin wave length $2\pi / \lambda \sim 100$ nm, so in a BFO ring of $\mu$m size, the number of spirals at zero field is $n_Q \sim 10$, and applying $|E| \sim 1$ kV/cm can change $n_Q$ by 1. Beside changing the winding number, we remark that the magnetoelectric effect can also be used to affect the topological properties of a magnet in a different respect [25]. Table I lists the parameters and their order of magnitude values by assuming CMI has similar material properties as other magnetic oxide insulators such as YIG, and we adopt lattice const $a \sim 1$ nm for both CMI and the NM for simplicity.

For the device in Fig. 2, consider the field $|E|^0 \sim 100$ kV/cm oscillating with a household frequency $\omega \sim 100$ Hz is applied to a range $L \sim 1$ mm. This region covers $N = L/a \sim 10^6$ sites with a number of spirals $n_Q \sim 10^4$ at zero field. The $E$ field changes the number of spirals to $n_Q \sim 10^5$ within time period $1/\omega \sim 0.01$s, so the spins at the boundary $S_N$ wind with angular speed $\partial_\theta \theta_N \sim 10^5 \sin \omega t$ which is enhanced by 5 orders of magnitude from the driving frequency $\omega$. To estimate the ejected spin current in Eq. (14), we use the typical spin relaxation time $\tau_{sl} \sim 10^{-12}$ s and length $\lambda_N \sim 10$ nm for heavy metals [16]. The $s$-$d$ coupling can range between $10^{-2}$ 0.01 eV to 1 eV. We choose $\Gamma \sim 0.1$ eV, which gives $\tau_{ex} \sim 10^{-12}$ s. The spin Hall angle $\theta_H \sim 0.1$ has been achieved [27, 28]. To estimate $n_0$, we use the fact that the $s$-$d$ hybridization $\Gamma r \cdot \Sigma$, is equivalent to applying a magnetic field $H = 2\Gamma S_i / \hbar \mu_B$ locally at the interface atomic layer of the NM. Given the typical magnetic susceptibility $\chi_m \sim 10^{-4}$ cm$^3$/mol and molar volume $V_m \sim 10$ cm$^3$/mol, the interface magnetization of the NM is $n_0 \mu_B = \chi_m H / V_m \sim 10^4$ C/m, thus $n_0 \sim 10^{27}$/m$^3$. The oscillating E field gives $\hat{S}_i \times \partial_\theta \hat{S}_i = \partial_\theta_N \hat{y} \sim \hbar \times 10^7$ Hz, so the ejected spin current is $J_{sNM} \sim 10^{24} A$/m$^2$. Using the design in Fig. 2(a) to convert $J_{sNM}$ into a charge current via inverse SHE yields $J_{sNM} \sim 10^4 A$/m$^2$, hence a voltage $\sim \mu$V oscillating with $\omega$ in a mm-wide sample. To use the setup in Fig. 2(b), a NM film of area $\sim 1$ mm$^2$ and thickness $\sim 10$ nm yields $E \sim \mu$V oscillating with $\omega$. 

In summary, we propose that an oscillating E field can generate significant spin dynamics in multiferroic insulators. For multiferroics that have coplanar AF spiral order, such as BFO, applying an oscillating E field with frequency as low as household frequency generates a coherent planar rotation of the spin texture whose frequency is many orders of magnitude enhanced. This coherent rotation activates the Goldstone mode of multiferroic insulators, which is not obstructed by crystalline anisotropy. The Goldstone mode can be used to deliver electricity at room temperature up to the extensions of magnetic domains, in a way that would be essentially free from the energy loss due to Gilbert damping.

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