Spin-Orbit Torque Magnetization Switching in MoTe₂/Permalloy Heterostructures

Shiheng Liang, Shuyuan Shi, Chuang-Han Hsu, Kaiming Cai, Yi Wang, Pan He, Yang Wu, Vitor M. Pereira, and Hyunsoo Yang*

The ability to switch magnetic elements by spin-orbit-induced torques has recently attracted much attention for a path toward high-performance, nonvolatile memories with low power consumption. Realizing efficient spin-orbit-based switching requires the harnessing of both new materials and novel physics to obtain high charge-to-spin conversion efficiencies, thus making the choice of spin source crucial. Here, the observation of spin-orbit torque switching in bilayers consisting of a semimetallic film of 1T'-MoTe₂ adjacent to permalloy is reported. Deterministic switching is achieved without external magnetic fields at room temperature, and the switching occurs with currents one order of magnitude smaller than those typical in devices using the best-performing heavy metals. The thickness-dependence can be understood if the interfacial spin-orbit contribution is considered in addition to the bulk spin Hall effect. Further threefold reduction in the switching current is demonstrated with resort to dumbbell-shaped magnetic elements. These findings foretell exciting prospects of using MoTe₂ for low-power semimetal-material-based spin devices.

Spintronics relies on both the electron’s charge and spin degrees of freedom to develop functional electronic devices for the future needs in information and communication technologies. A crucial milestone in this direction is the manipulation of magnetic moments by planar electrical currents.\(^1\)\(^-\)\(^3\) The simplest heterostructures interface a ferromagnetic strip with a conducting, nonmagnetic film such that the spin Hall effect (SHE)\(^4\)\(^-\)\(^6\) and/or Rashba–Edelstein effects\(^6\)\(^-\)\(^8\) induce a torque in the adjacent magnet. This has spurred magnetic memory concepts based on planar heterostructures.\(^9\)\(^-\)\(^11\) Although different microscopic mechanisms contribute to the magnetic torque,\(^10\) the central ingredient is the development of a robust non-equilibrium spin texture and accumulation at the interface between the conducting and recording layers in response to a charge current. This relies on the spin-orbit coupling (SOC) of conduction electrons in the nonmagnetic layer, and the nonmagnet is thus the critical component: not only should it feature a high charge-to-spin conversion efficiency, it must desirably have a low resistivity as well to enable low-power operation.\(^13\)

Spin-orbit torques (SOT) primarily due to SOC have recently become a topic of intense focus.\(^10\)\(^,\)\(^11\) On the one hand, torques arising from the SHE in adjacent heavy metals such as Ta, Pt, and W\(^1\)\(^2\),\(^12\)\(^-\)\(^14\) appear to require large switching currents; this drives an impetus to discover new materials capable of overcoming this limitation. On the other hand, a number of emerging 2D crystals and topological materials are promising platforms for spintronics\(^15\)\(^,\)\(^16\) due to their large intrinsic SOC and nontrivial spin textures.\(^17\)\(^-\)\(^19\) In addition, crystal symmetries can be explored to engineer charge-to-spin conversion and SOT in device configurations that have been so far inaccessible.

Beyond large intrinsic SOC, some 2D materials display unique electronic structures, nontrivial spin textures, and spin-polarized surface states.\(^20\)\(^-\)\(^22\) Most importantly, in the specific case of 2D semimetals, the combination of large conductivity and large SOC suggests that high charge-to-spin conversion efficiencies might be possible. Compared with topological insulators based on Bi\(_2\)Se\(_3\),\(^18\) for example, semimetals display ten times smaller electrical resistivities, suggesting that SOT-based magnetization switching could operate with smaller power requirements. Previous studies of SOT based on 2D materials have been reported in semiconducting MoS\(_2\) and WSe\(_2\),\(^23\)\(^,\)\(^24\) metallic NbSe\(_2\),\(^25\) semimetallic WTe\(_2\),\(^26\)\(^-\)\(^28\) and MoTe\(_2\).\(^29\)\(^,\)\(^30\)

Especially, the novel observations of out-of-plane spin induced anti-damping torque in WTe\(_2\) and MoTe\(_2\) hold a great promise for field-free switching of the perpendicularly magnetized system. However, a direct demonstration of SOT-induced...
magnetization switching utilizing a 2D material is still at its infancy and the thickness-dependence of SOTs in 2D materials is not well understood.

Here, we employ the semimetal 1T'-MoTe2 as a spin source and observe a large charge-to-spin conversion efficiency attributed to an underlying electronic structure that generates both a robust spin Hall conductivity (SHC) and interfacial spin accumulation. We report the direct observation of current-driven magnetization switching via SOT in an MoTe2/Py heterostructure without magnetic fields at room temperature. The threshold current density required for complete magnetization reversal is found to be \(J_{\text{th}} \approx \) the ST-FMR signal. An in-plane external magnetic field is swept over an angle \(\theta\) with respect to the current direction \(x\) to match the ferromagnetic resonance condition. A typical \(V_{\text{max}}\) signal shown in Figure 1b is analyzed by \(V_{\text{mix}}(H) = V_S F_S + V_A F_A\), where \(F_S\) and \(F_A\) are symmetric and antisymmetric Lorentzian functions, respectively. The value of the effective spin Hall angle (\(\xi_{\text{SH}}\)) from our ST-FMR is 0.27 for the MoTe2 (83.1 nm)/Py sample (Section S2, Supporting Information), which is similar to the spin-to-charge conversion efficiency of 0.21 reported in the MoTe2 stacked with graphene lateral spin valves.[20]

The fitting of the ST-FMR signal is also used to estimate the amplitudes of the in-plane and out-of-plane torques (cf. Figure 1a), whose dependence on the field orientation should be \(\tau_{ij} = \tau_{ij}\cos(\theta)\) and \(\tau_{ij} = \tau_{ij}\cos(\theta) + \tau_{ij}\) respectively.[26] We obtain a ratio \(\tau_{ij}/\tau_{ij} = 0.55\). The contribution \(\tau_{ij}\) is an out-of-plane anti-damping-like torque which has been previously detected in WTe2/Pt with currents applied along the low-symmetry axis[26] (Section S3, Supporting Information). Since our current channels are defined along the low-symmetry axis of MoTe2, we investigate the presence of such torque contribution in our MoTe2/Py structures by fitting the angular-dependence of the antisymmetric component to \(V_{S}(\theta) = A \times \cos(\theta) \times \sin(2\theta) + B \times \sin(2\theta)\), as shown by the red curve in Figure 1c. For thin films of \(\beta\)-MoTe2 as reported by Stiehl et al.[29] there is an out-of-plane antidamping torque present in MoTe2/permalloy heterostructures, while in our case, we obtain \(B = 0\), which indicates that \(\tau_{ij}\) could be negligible for relatively thicker films of MoTe2.

To characterize the magnetic anisotropy of the Py layer, the ST-FMR data have been analyzed for the dependence of the

Figure 1. Spin torque ferromagnetic resonance measurements. a) Schematic of the ST-FMR measurement setup and illustration of the device with the SOT-induced magnetization dynamics. The directions \(a\) and \(b\) refer to the basal-plane crystal axes of monoclinic 1T'-MoTe2 and are, respectively, oriented parallel to the \(x\) (longitudinal) and \(y\) (transverse) directions defined with respect to the current flow. b) A representative ST-FMR signal (open symbols) from a MoTe2 (83.1 nm)/Py (6 nm) device with fits of the symmetric Lorentzian \((V_S F_S)\) component (red lines), and the anti-symmetric Lorentzian \((V_A F_A)\) component (green lines). The microwave frequency is 5 GHz and the applied microwave power is 14 dBm. An in-plane external magnetic field \(B_{\text{ext}}\) is applied at an angle \(\theta = 36^\circ\) with respect to \(I_{RF}\). c) Amplitude of the anti-symmetric component of the resonance as a function of the angle of the in-plane field with the current applied parallel to the \(a\) axis. d) Ferromagnetic resonance field as a function of the in-plane magnetization angle, with the current applied parallel to the \(a\) axis. All measurements are done at room temperature.
magnetic resonance frequency on the direction of in-plane magnetization, as shown in Figure 1d, with current along the a axis. Even though the shape anisotropy alone would determine the easy axis of Py to lie along the x or a axis direction, the magnetic anisotropy induced by the proximal MoTe₂ results in the easy axis along y, parallel to the b crystallographic direction of MoTe₂. Direct measurements of the magneto-optical Kerr effect (MOKE) on our MoTe₂/Py devices further corroborate that the b axis of MoTe₂ is the easy axis in the Py layer (Section S4, Supporting Information).

We have studied the magnetization switching response to pulsed dc currents (I) at room temperature. The magnetization of Py is collinear with the spin polarization in the semimetal, which implies that an anti-damping torque can switch the Py magnetization without external magnetic fields. The switching events have been captured by MOKE imaging in Figure 2a–f. The top row (Figure 2a–c) demonstrates that the magnetic moment of Py can be switched by pulsed current (30μs pulse width) applied between the two electrodes outlined in green (the MoTe₂/Py bilayer is outlined in red). Prior to each set of width) applied between the two electrodes outlined in green moment of Py can be switched by pulsed current (30 μs). The top row (Figure 2a–c) demonstrates that the magnetic moment of Py can be switched by pulsed current (30 μs). The top row (Figure 2a–c) demonstrates that the magnetic moment of Py can be switched by pulsed current (30 μs). The top row (Figure 2a–c) demonstrates that the magnetic moment of Py can be switched by pulsed current (30 μs).

Figure 2. SOT-driven magnetization switching in MoTe₂ (108.2 nm)/Py heterostructures. a–c) MOKE images of magnetization switching (at zero magnetic field) under a pulsed dc current I along the −x direction (leftward in these figures; the magnetic easy axis of Py is along the b axis). Current densities increase from (a) to (c) and, when not zero, their values in the bilayer (J_{total}) and those estimated for the MoTe₂ layer alone (J_{MoTe₂}) are stated underneath each panel. The red/dashed rectangles outline the 12 μm wide channel and the green/dashed ones mark the electrodes. The dark (light) contrast signals the local magnetization along +y (+y). White and black block-arrows are drawn to indicate the magnetization direction and its reversal in the panels with fully switched states. d–f) MOKE images of the reverse switching process, where the initially magnetization state along +y is reversed by a current along −x. g) Thickness-dependence of the switching current densities J_{total} and J_{MoTe₂}. h) Thickness-dependence of the experimental SOT efficiency and the intrinsic bulk spin Hall angle (SHA) of MoTe₂ slabs calculated at the Fermi level. All measurements are done at room temperature.

The current-driven switching has been reproduced on a number of different devices with different MoTe₂ thickness (Figure S6, Supporting Information). Figure 2g summarizes the critical switching current extracted for the bilayer (J_{total}) as a function of MoTe₂ thickness. Both critical current densities increase with thickness. We note that the current density required for complete magnetization reversal is low, at J_{MoTe₂} = 6.71 × 10^7 A cm⁻² for a 66.1-nm-thick MoTe₂ film. This value is one order of magnitude below that needed if the best heavy metals are used as spin sources. For a comparison, we also perform current-driven switching measurements on control samples of Pt/Py and Cu/Py (Section S7, Supporting Information). The magnetization switching takes place at J_{total} = 2.8 × 10^7 A cm⁻² in the Pt (6 nm)/Py (6 nm) sample, which is about one order larger than the J_c in MoTe2/Py samples. We do not observe any switching
behavior in Cu (12 nm)/Py (6 nm) with a current density up to $J_{\text{total}} = 2.97 \times 10^7$ A cm$^{-2}$, which is one order of magnitude larger than the $J_C$ for our MoTe$_2$/Py samples.

Figure 2h displays the estimated SOT efficiency, which we quantify as an effective $\xi_{\text{SH}}$ on the basis of the conventional anti-damping spin-torque-driven model of magnetization switching\cite{1,31} (Section S5, Supporting Information). Consistently with the evolution of the switching currents, the SOT efficiency is smaller for thicker MoTe$_2$ films. We determine the SOT efficiency of our thinnest device (66.1 nm) to be $\xi_{\text{SH}} = 0.35$, which is larger than those of Ta ($\approx 0.15$)\cite{31} or Pt ($\approx 0.06$)\cite{12}.

The extracted $\xi_{\text{SH}}$ from the switching measurements shows a good agreement with that from the ST-FMR data of $\xi_{\text{SH}} = 0.27$ for 83.1 nm MoTe$_2$. The decreasing trend in Figure 2h as the thickness increases is opposite to that seen in magnetization switching driven by the SHE in the heavy metal layer and WTe$_2$\cite{28}, where the efficiency increases with thickness and saturates when $t$ exceeds the spin diffusion length, $\lambda_{sf}$\cite{5}. These observations point to a possible influence of the interface between MoTe$_2$ and Py. This is supported by the fact that, if one considers only the contribution of the intrinsic SHE in MoTe$_2$, the theoretical value of the spin Hall (SH) angle computed for bulk MoTe$_2$ ($\theta_{\text{SH}} \approx 0.13$) is smaller than the experimentally estimated values, as indicated by the dashed line in Figure 2h. This theoretical value was obtained as the ratio $\theta_{\text{SH}} = \left(\frac{2e}{\hbar}\right)\frac{J_{\mu}}{J_{\text{iso}}} = \left(\frac{2e}{\hbar}\right)\frac{\sigma_{\text{SH}}^{\mu}}{\sigma_{\mu}}$, between the calculated SHC ($\sigma_{\mu}^{\mu}$) and the experimental charge conductivity ($\sigma_{\text{exp}}^{\mu}$). From experiments, we extract the SHC of $\approx 250 (\hbar/\Omega \cdot \text{cm}^{-1})$, which is similar to the value of $\approx 280 (\hbar/e \cdot \Omega^{-1} \cdot \text{cm}^{-1})$ reported by Safeer et al.\cite{30} but larger than the previously reported value of $\approx 29 (\hbar/e \cdot \Omega^{-1} \cdot \text{cm}^{-1})\cite{29}$.

To understand the possible microscopic origin of the large experimental SOT efficiency, we calculate both the intrinsic SHC and non-equilibrium spin accumulation of MoTe$_2$ from first-principles (Sections S9,S10, Supporting Information). The crystal structure of monoclinic 1T-MoTe$_2$ (space group P2$_1$/m, #11) is shown in Figure 3a, where the box highlights the 3D unit cell that encompasses two inequivalent MoTe$_2$ layers (the commonly referred zigzag chains of the transition metal ion run along the $a$ direction). Figure 3b shows a representative band structure computed for a 20-monolayer slab. The energy dispersion displays an underlying band inversion centered at the $\Gamma$ point in the Brillouin zone. Both few-layer and bulk crystals are semimetals with a low carrier density that arises from small Fermi surface pockets centered along the $\Gamma X$ line in reciprocal space, with holes (electrons) nearer (farther) to the $\Gamma$ point.

Figure 3c summarizes our results of the SHC, where the SH tensor component $\sigma_{\alpha\beta}^{\mu}$ quantifies the generation of a spin current $J_{\mu}^{\alpha}$ that propagates along the Cartesian direction $u_{\alpha}$ carrying spins polarized along $u_{\alpha}$, as a result of an electric field pointing along $u_{\beta}$: $J_{\mu}^{\alpha} = \sigma_{\alpha\beta}^{\mu} E_{\beta}$. The experimental measurement layout with charge flowing along the $x$ direction restricts the analysis to the SH component $\sigma_{x\beta}^{\mu}$\cite{31}. The component $\sigma_{x\beta}^{\mu}$ is not relevant, since our measurements reveal $\tau_{\beta} = 0$. Figure 3c

Figure 3. Electronic structure and spin response of MoTe$_2$. a) Crystal structure of the monoclinic 1T' phase of MoTe$_2$. b) Band structure of a MoTe$_2$ slab with 20 monolayers. c) The component $\sigma_{x\beta}^{\mu} (= \sigma_{x\beta}^{\mu})$ of the SHC calculated for different thicknesses, at three different chemical potentials $\mu$ ($\mu = 0$ in undoped MoTe$_2$). d) The non-equilibrium spin density $\delta S_{\alpha}^{\mu}$, which results from the accumulation of a finite spin density pointing along the $b$ crystalline direction under a charge current parallel to $a$ using $\tau_{\beta} = \tau$ and $\lambda_{sf} = 10$ nm.
shows the intrinsic SHC calculated in Kubo linear response (Section S9, Supporting Information) for charge neutrality (i.e., for chemical potential \( \mu = 0 \)) and varying thickness up to our numerical limit of 40 monolayers (277 nm). While the variation is somewhat pronounced between 1–3 layers, it settles to an asymptotic value \( \sigma_{s+}^{a} = 60 \,(h/e)^{-1} \, \text{cm} \, \text{m}^{-1} \) within a few nm. By repeating the calculation at different chemical potentials, we have confirmed that the sign of the SH response is robust both as a function of thickness and doping, and also that its magnitude can be increased (decreased) under electron (hole) doping (Figure S9, Supporting Information). Even though the range of thicknesses covered in the calculations is thinner than that of our experimental devices, these theoretical results unveil that the SHC of MoTe\(_2\), \( \sigma_{s+}^{a} \), has a considerable magnitude.

The inversion symmetry of the \( \text{IT}^{\prime} \) crystalline structure is broken in our bilayer devices. Combined with the relatively large SOC that one can infer from recent experimental values of \( \xi_{\text{SH}} \) in MoTe\(_{2}\)[29,30] in addition to the SHE, we inquire a non-equilibrium spin accumulation arising from the Rashba–Edelstein effect.[7] Figure 3d shows the thickness-dependence of spin accumulation collinear with the magnetization in Py, \( \delta_{s}^{\alpha} \), computed at charge neutrality as well as \( \mu = \pm 60 \, \text{meV} \).

For a direct comparison with SHC, we introduce a spin diffusion current as \( J_{\text{sd}}^{\alpha} = D \delta_{s}^{\alpha} J_{\text{tot}} = \delta_{s}^{\alpha}E_{z} \), where \( D = \lambda_{sd}^{2}/\tau_{s} \) is a spin diffusion constant, \( \lambda_{sd} \) and \( \tau_{s} \) are the spin diffusion length and relaxation time, respectively. This allows us to present the results in terms of the quantity \( \delta_{s}^{\alpha} \), which has dimensions of a spin conductivity. We used \( \tau_{s} = \tau \) and \( \lambda_{sd} = 10 \, \text{nm} \), but different values can be considered simply rescaling the curves shown by the desired prefactor \( \lambda_{sd}/\tau_{s} \), where \( \tau \) is the Drude relaxation time; the value of \( \lambda_{sd} \) was estimated combining a recent experiment[30] with our calculated value of \( \theta_{\text{SH}} \) in the bulk; the assumption \( \tau_{s} = \tau \) is expected for Elliot–Yaffet relaxation with strong spin-orbit coupling.[34] As anticipated, the spin accumulation decays quickly with increasing the thickness (note that \( \delta_{s}^{\alpha} \) is normalized by the thickness), reflecting the surface nature of this contribution. More importantly, we can see that the magnitudes of the accumulation-derived spin conductivity, \( \delta_{s}^{\alpha} \), shown in Figure 3d for the thinnest slabs (<4 nm) match those of the SHC plotted in Figure 3c. It indicates that interfacial spin accumulation contributes, at least, on equal footing with the SHE to the torque, and might even be the dominant effect if disorder is found to significantly reduce the intrinsic SHC from the values shown in Figure 3c.

We now demonstrate that the threshold switching current can be further decreased by engineering the shape of the heterostructure. Figure 4a shows dumbbell-patterned heterostructures of MoTe\(_2\) (109.9 nm)/Py (6 nm), where the diameters of the left (right) disks are 16.60 µm (15.60 µm) and the channel length (width) is 11.50 µm (9.11 µm). This geometry reduces the domain nucleation energy that must be overcome to initiate the magnetization switching process by localizing the domain wall within the central channel. If one begins from the fully saturated state, a current density of \( J_{\text{MoTe2}} = 1.32 \times 10^{6} \, \text{A} \, \text{cm}^{-2} \) is needed to fully reverse the initial Py magnetization (Section S6, Supporting Information). Figure 4b–h shows MOKE images of the SOT-driven domain wall motion, at zero magnetic field, driven by a pulsed dc current along the a axis of MoTe\(_2\). Beginning with the domain wall at the edge of the rightmost disk (Figure 4b), a leftward (–x) current displaces the domain wall to the left. Complete displacement of the wall to the leftmost edge occurs for \( J_{\text{MoTe2}} \) between 4.78 \( \times \) 10\(^{5} \) and 5.45 \( \times \) 10\(^{5} \) A cm\(^{-2} \). This is a near threefold reduction in comparison with the value needed for complete switching from full saturation (no domain wall, 1.32 \( \times \) 10\(^{6} \) A cm\(^{-2} \)).

As we see in Figure 4e, the domain wall stops at the edge of the leftmost disk, which happens because of

---

Figure 4. SOT-driven domain wall motion in dumbbell-patterned heterostructures of MoTe\(_2\) (109.9 nm)/Py. a) Optical image of the dumbbell-shaped device. b–h) MOKE images of domain wall motion (at zero magnetic field) under a pulsed dc current \( I \) along the \(-x\) (b) to (e), red arrows) or \(+x\) directions (f) to (h), blue arrows) of MoTe\(_2\). The current densities in the bilayer \( J_{\text{tot}} \) and in the MoTe\(_2\) layer \( J_{\text{MoTe2}} \) are stated underneath each panel. The arrows indicate the current direction. i–l) Micromagnetic simulation results of the SOT-driven domain wall motion. (i) and (j) show the domain wall being blocked at, respectively, the left and right edges of the disks under an applied dc current. (k) and (l) show the relaxed stable states after removal of the current.
the domain wall surface energy and the smaller current density in the disks compared with the channel. Micromagnetic simulations shown in Figure 4i–l confirm this behavior (details in Section S8, Supporting Information) and show that the disk regions effectively oppose the penetration of the domain wall. Finally, reversing the current direction from the state in Figure 4e, one can displace the domain wall in the opposite direction until it stops at the edge of the rightmost disk (Figure 4f–h). Consequently, in this shape-engineered geometry, the magnetization switching, semimetal MoTe$_2$, spin-orbit torque, spintronics.

Keywords
magnetization switching, semimetal MoTe$_2$, spin-orbit torque, spintronics.

Supporting Information
Supporting Information is available from the Wiley Online Library or from the author.

Acknowledgements
S.L. and S.S. contributed equally to this work. The authors thank Prof. G. Liang for discussion, G. Yang for helping experiments, and J. Wang for Raman measurements. This research was supported by SpOT-LITE programme (A$^*$STAR Grant no. 18A6600057) through RIE2020 funds from Singapore, Singapore Ministry of Education AcRF Tier 1 (R-263-000-D61-114), and Samsung Electronics’ University R&D program (Exotic SOT materials/SOT characterization). S.L. also thanks the support of NNSFC (No. 11904088).

Conflict of Interest
The authors declare no conflict of interest.

Keywords
magnetization switching, semimetal MoTe$_2$, spin-orbit torque, spintronics

Received: April 25, 2020
Revised: June 29, 2020
Published online: August 2, 2020

[1] L. Liu, C.-F. Pai, Y. Li, H. W. Tseng, D. C. Ralph, R. A. Buhrman, Science 2012, 336, 555.
[2] I. M. Miron, K. Garello, G. Gaudin, P.-J. Zermatten, M. V. Costache, S. Aufrert, S. Bandiera, B. Rodmacq, A. Schuhl, P. Gambardella, Nature 2011, 476, 189.
[3] A. Brataas, A. D. Kent, H. Ohno, Nat. Mater. 2012, 11, 372.