Deterministic switching of a perpendicularly polarized magnet using unconventional spin–orbit torques in WTe$_2$

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Spin–orbit torque (SOT)-driven deterministic control of the magnetic state of a ferromagnet with perpendicular magnetic anisotropy is key to next-generation spintronic applications including non-volatile, ultrafast and energy-efficient data-storage devices. However, field-free deterministic switching of perpendicular magnetization remains a challenge because it requires an out-of-plane antidamping torque, which is not allowed in conventional spin-source materials such as heavy metals and topological insulators due to the system’s symmetry. The exploitation of low-crystal symmetries in emergent quantum materials offers a unique approach to achieve SOTs with unconventional forms. Here we report an experimental realization of field-free deterministic magnetic switching of a perpendicularly polarized van der Waals magnet employing an out-of-plane antidamping SOT generated in layered WTe$_2$, a quantum material with a low-symmetry crystal structure. Our numerical simulations suggest that the out-of-plane antidamping torque in WTe$_2$ is essential to explain the observed magnetization switching.

Spin–orbit torque (SOT) has emerged as an efficient means of manipulating the magnetic state of ferromagnetic (FM) materials$^{1-4}$. The potential technological applications of SOT-driven magnetization manipulation include energy-efficient non-volatile magnetic memories$^{5}$ and spin–torque oscillators$^{6}$. In SOT-induced magnetic switching, a charge-current density $j_c$ (x direction) flowing in the plane of a bilayer structure consisting of spin-source material and a ferromagnet results in a spin current flowing in the out-of-plane direction (z direction) via spin galvanic effects$^{7-9}$. This spin current in turn exerts a torque on the magnetization of a nearby magnetic layer. This torque has an antidamping component $\tau_{z0} \propto m \times p \times m$ and a field-like component $\tau_{z1} \propto \hat{p} \times \hat{m}$ where $\hat{m}$ is the magnetization direction and $\hat{p}$ is the spin polarization direction. Due to the symmetry of heavy-metal (HM)/FM bilayer heterostructures, the spin is polarized in the $y$ direction and the in-plane antidamping torque ($\tau_{z1}^{IP}$) has the form of $\hat{m} \times \hat{z} \times \hat{m}$ and can only deterministically switch the magnetization of a magnet that has an in-plane magnetic anisotropy$^{10}$. However, for memory applications, magnets with perpendicular magnetic anisotropy (PMA) are highly desired because this allows for ultra-compact packing and thermally stable, nanometre-sized magnetic bits$^{1}$. In conventional HM/FM systems, an external magnetic field is applied along the direction of the charge current to break the in-plane symmetry of the system, allowing for deterministic switching of the magnetization state of a PMA magnet$^{11}$. Previously, a structural asymmetry$^{9,10}$, tilted anisotropy of the nanomagnet$^{11}$, an in-plane magnetized layer$^{11-14}$, an in-plane effective magnetic field in FM/ferroelectric structures$^{15}$, anomalous Hall effect (AHE) and planar Hall effect in FM/HM/FM trilayers$^{16}$, and interplay of SOT and spin-transfer torque$^{17}$ have been explored to achieve a field-free deterministic switching of PMA magnets using SOTs. Another approach for this is to explore the utility of emergent quantum systems as a spin-source material wherein SOTs can be controlled by crystal symmetries. Recently, transition-metal dichalcogenides with low-symmetry crystal structures, such as WTe$_2$, have been shown to exhibit an out-of-plane antidamping torque ($\tau_{z1}^{OP}$) of the form $\hat{m} \times \hat{z} \times \hat{m}$ when a charge current is applied along the low-symmetry axis of the WTe$_2$/FM bilayer system$^{18,19}$. In addition, an efficient field-free SOT switching, due to a strong charge-to-spin conversion efficiency in WTe$_2$, of an in-plane ferromagnet has been demonstrated in WTe$_2$/permalloy bilayer systems with a very small charge density (~2.96 × 10$^9$ A m$^{-2}$)$^{20}$. However, it has not yet been established experimentally whether the unconventional out-of-plane antidamping SOT in layered materials with low-crystal symmetry is strong enough to enable the deterministic switching of a PMA magnet.

Here we experimentally realize field-free deterministic switching of a perpendicularly polarized magnet using SOTs in a quantum material with a low-symmetry crystal structure. For this, we fabricate SOT devices using a van der Waals-based layered quantum material platform. Thin films of WTe$_2$ are used as a spin-source material for generating SOTs$^{11,20-22}$. In addition to the generation of

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an out-of-plane antidamping torque, WTe₂ also exhibits properties that are highly relevant to achieving a large charge-to-spin conversion efficacy: strong spin–orbit coupling, non-trivial band dispersion, topologically protected spin-polarized bulk and surface states, a pronounced Edelstein effect and an intrinsic spin Hall effect. WTe₂ is a low-symmetry system, the plane of which is schematically depicted in Fig. 1a (with the axis normal to the plane). The surface of WTe₂ only has mirror symmetry with respect to the plane but not about the plane. Thus, the system is asymmetric relative to a 180° rotation about the axis. For the PMA magnet, we utilize Fe₂.₇₈GeTe₂ (FGT), which is a layered van der Waals FM material. Mechanical dry-transfer techniques to assemble WTe₂/FGT bilayers and standard device fabrication techniques are used to prepare the SOT devices (Methods and Supplementary Note 1).

In total, we have measured six WTe₂/FGT devices and here we present data from devices A and F. The heterostructures of devices A and F consist of WTe₂ (25.8)/FGT (4.1)/h-BN (39.6) and WTe₂ (10.2)/FGT (9.0)/h-BN (31.3) respectively, where the numbers in parentheses indicate the layer thicknesses in nanometres and h-BN is hexagonal boron nitride. An optical micrograph of device A is shown in Fig. 1b along with a side-view schematic of the device (top panel). The crystallographic and axes are labelled and are confirmed by polarized Raman spectroscopy. Raman spectra are collected by rotating the polarization of the incident laser for different angles relative to the axis of WTe₂ and the integrated intensities under each peak are calculated for the contour plot shown in Fig. 1c. The polarization angle dependence of the device A Raman peak at 212 cm⁻¹ (corresponding to the black dashed line in Fig. 1c) exhibits minimum intensity when the excitation laser polarization is along the straight edge of the WTe₂ flake (Fig. 1b). This is consistent with previous reports and clearly distinguishes the axis of WTe₂. We have also carried out atomic structure characterization of device A using scanning transmission electron microscopy (STEM). Figure 1d shows a STEM cross-sectional image, confirming that the lateral direction is the axis of WTe₂. Inset: the atomic arrangement in WTe₂ in the plane.

**Fig. 1 | Crystal structure, polarized Raman and electron microscopy of the SOT devices.** a, A model showing the crystal structure of WTe₂ with and axes labelled. The crystal is invariant (non-invariant) upon a mirror operation. b, Top: a schematic showing the side view of the device. Bottom: an optical image of device A. The electrodes are set up to allow the application of current pulses along the and axes of WTe₂. The FGT, WTe₂ and h-BN flakes are clearly labelled. c, The angle-dependent polarized Raman spectral intensity of a WTe₂ flake used in device A. The measurements were obtained by rotating the laser polarization with respect to the long edge (axis) of the WTe₂ flake. The dashed line indicates the Raman peak at 212 cm⁻¹. d, Cross-sectional STEM image of the device viewed along the -axis orientation of WTe₂, confirming that the lateral direction is the axis of WTe₂. Inset: the atomic arrangement in WTe₂ in the plane.
The presence of a strong $\tau_{op}^{AD}$ in a WTe$_2$/permalloy heterostructure was previously probed by spin–torque ferromagnetic resonance. The $\tau_{op}^{AD}$ is independent of the reversal of magnetization, reverses with current direction and can efficiently switch a perpendicular magnetization. The $\tau_{op}^{AD}$ is not allowed in conventional spin-source systems, such as HMs and topological insulators, due to two-fold rotational symmetry. However, this symmetry is broken in WTe$_2$. Specifically, WTe$_2$ has no mirror symmetry in the ac plane (Fig. 1a), allowing for a non-zero $\tau_{op}^{AD}$ when a current charge is applied along its a axis (Fig. 2a). On the other hand, when a current charge is applied along the b axis of WTe$_2$, the preserved mirror symmetry in the bc plane requires that $\tau_{op}^{AD} = 0$ as depicted in Fig. 2d. Essentially, when a current is applied along the b axis and a bc plane mirror operation is performed, the current does not switch direction (that is, $I \rightarrow (-I)$, but the out-of-plane antidamping torque (a pseudovector) will change sign (that is, $\tau_{op}^{AD} \rightarrow -\tau_{op}^{AD}$). But a current-induced torque must change sign with current reversal, requiring that $\tau_{op}^{AD} = 0$ when the current is along the b axis.

To examine the presence of $\tau_{op}^{AD}$ we perform AHE loop shift measurements (Methods). An $\tau_{op}^{AD}$ can abruptly shift the AHE hysteresis loop once the current passes a threshold value such that the intrinsic damping is compensated. When a current charge pulse ($I_p$) is applied along the a axis (that is, $I \parallel \hat{a}$), AHE hysteresis loops measured at low pulse currents ($I_p = \pm 2$ mA) look identical for different current polarity as shown in the top panel of Fig. 2b. However, when $I_p$ increases to $\pm (+)7$ mA (Fig. 2b, bottom), the AHE loop shifts to the negative (positive) field values. On the other hand, when $I_p$ is applied along the b axis, no notable loop shift is observed for both low and high current pulse magnitudes (Fig. 2e). Note that there is also a reduction in the coercive field at high current densities that is mainly caused by Joule heating (Supplementary Note 5). The measurements are repeated at different charge current magnitudes for $I \parallel \hat{a}$ and $I \parallel \hat{b}$ and are shown in Fig. 2c,f, respectively. When $I \parallel \hat{a}$, we observe a threshold current effect in $H_{sh}(|I_p|)$ and a clear splitting between $H_{sh}^+ (H_{sh}^-)$ at higher currents, where $H_{sh}^+ (H_{sh}^-)$ is the loop shift field measured at positive (negative) current. We speculate that the observed asymmetry in $H_{sh}(|I_p|)$ for positive and negative currents may be due to the fact that the device is not patterned into an ideal Hall cross geometry. In contrast, for $I \parallel \hat{b}$, $H_{sh}(|I_p|)$ shows very weak $I_p$ dependence and remains close to zero for low and high current values, which is similar to what has been reported previously for conventional HM/FM systems. These results clearly show the presence of a non-zero $\tau_{op}^{AD}$ when current is applied along the a axis. The abrupt shift, instead of a linear shift, in $H_{sh}$ for $I \parallel \hat{a}$ also indicates that no significant out-of-plane field-like torque is present in the system.

Next, we demonstrate field-free switching of perpendicular magnetization of FGT employing charge-current-induced SOTs in WTe$_2$ (see Methods for measurement details). When $I \parallel \hat{a}$, we observe a clear deterministic switching (Fig. 3a). Figure 3a shows that the terminal state is determined by the polarity of $I_p$, that is, a positive current favours magnetization pointing up and a negative current favours magnetization pointing down. We attribute the observed deterministic switching of FGT to $\tau_{op}^{AD}$ generated by the out-of-plane spin polarization, $p_{s\perp}$, injected from WTe$_2$. On the other hand, the behaviour is completely different when the current is applied along the b axis. There is no deterministic switching observed in the absence of an external in-plane magnetic field (Fig. 3b). This behaviour is similar to SOT switching in conventional HM/FM bilayer systems, that is, the terminal state is close to the demagnetized state $m_r = 0$, indicating that $\tau_{IP}^{AD}$, generated by spin polarization pointing in the y direction, drives the magnetic moment of FGT to the device plane. Without the presence of $p_{s\perp}$ (or a symmetry-breaking in-plane magnetic field), multidomains
and the temperature of the device remains below the Curie temperature (Tc). The black arrow shows the difference between m_z = +1 and m_z = −1.

The thermal effect due to Joule heating in our devices is estimated by comparing the temperature and charge-current-dependent longitudinal resistance of the WTe2/FGT bilayer (Supplementary Note 5). For this, we look at a device F in which we have thinner WTe2 (higher resistance), such that a larger current flows through the FGT layer giving us a more accurate estimation of the temperature profile of magnetic layer due to Joule heating (Supplementary Note 3). The threshold current (Ith), defined as the current that switches the magnetic state across m_z = 0, is obtained at different temperatures and translated to the threshold current density (Jth), considering only current flowing in WTe2, for both the a axis and the b axis (Fig. 4a). The inset of Fig. 4a shows the field-free SOT switching observed in device F when a charge current is applied along the a axis. We also estimate ΔT, which is the change in temperature of the device when the Ith current pulse is applied. We observe that ΔT is higher (lower) when we perform SOT-switching experiments at lower (higher) temperatures. At lower temperatures, FGT coercive fields are higher (Supplementary Note 4) and will require larger Joule heating (a larger ΔT value) and hence large current densities to switch the magnetization as observed in Fig. 4b. These results suggest that Joule heating assists the SOT-induced magnetization reversal by lowering the energy barrier for magnetization switching in our low-temperature SOT-switching experiments. In addition, it is clear from Fig. 4b that the temperature of the device remains below the Curie temperature (Tc). The red dashed line in Fig. 4b indicates the ΔT that is required to surpass the Tc when switching experiments are performed at various temperatures (140–200 K).

To better distinguish the roles of τop and τab, we perform numerical simulations of magnetization switching for a monodomain FGT. While the actual switching processes are likely to involve domain-wall motions, it is customary to use a monodomain simulation that captures the essential physics behind the switching process (Supplementary Note 6), which is simple enough for an intuitive understanding. To this end, we use the coercive field as the effective anisotropy (Supplementary Note 7). We also take thermal effects into account by considering the temperature dependence of the coercive field in the presence of current-induced Joule heating (Supplementary Notes 4 and 5), which yields the effective anisotropy in our simulations as a function of the applied current and the setting temperature. We first investigate the case of I||a at 160 K. As shown in Fig. 4c, by varying the current density in WTe2 (j) and the ratio of τop/j, we find that the initial state (m_z = −1) can be switched to an in-plane direction (m_z = 0) that is energetically metastable. The threshold decreases with an increasing τop/j, as reflected by the phase boundary. If the driving current is turned off, the metastable state (m_z = 0) will equally probably relax towards m_z = −1 and m_z = 1. In reality, the metastable state will


Fig. 4 | Simulation of SOT-induced switching, Joule heating and temperature dependences. a. The threshold current density $J_{th}$ as a function of temperature obtained for a charge current applied along the $a$ and $b$ axes for device F. For both cases, a larger $J_{th}$ is required to switch the magnetization of FGT at lower temperatures. Inset: SOT-induced-field-free deterministic switching for $I \parallel \hat{a}$ at 160 K. b. The temperature rise, $\Delta T$, from the setting temperature (that is, the temperature at which switching experiments is performed) ascribed to Joule heating (when $J_{th}$ is reached) for a charge current applied along the $a$ and $b$ axes. At lower temperatures, $\Delta T$ is higher, suggesting a larger Joule heating (and hence a higher current density) is associated with the switching process. The red dashed line depicts the $\Delta T$ that is required to exceed the $J_{th}$ at different temperatures where switching experiments are performed. c,d. Simulations for monodomain magnetization switching when a charge current is applied along the $b$ axis (c) or $a$ axis (d) at 160 K. The initial magnetic state is set to be $m_z = -1$, and the blue, red and white regions correspond to unswitched, switched ($m_z = 1$) and metastable ($m_z = 0$) states, respectively. When a current is applied along the $b$ axis (c), $\tau_{OP}^{AD}$ switches the magnetization to a metastable state with $m_z = 0$ and $m_y > 0$ ($m_x < 0$) for positive (negative) current. When a charge current is applied along the $a$ axis (d), an out-of-plane antidamping torque $\tau_{IP}^{AD}$ is generated, and the magnetization can only be switched for a positive current direction. If $\tau_{IP}^{AD}$ is not strong enough, however, $\tau_{OP}^{AD}$ will dominate and switch the magnetization only to the metastable magnetic states, that is, no deterministic switching occurs. e,f. The phase boundaries at various temperatures for FGT monodomain switching for a charge current applied along the $b$ axis (e) or $a$ axis (f). The $\tau_{IP}^{AD}$ used at 140, 160 and 180 K is 0.86, 0.59 and 0.41 mT ($\times 10^{10}$ A m$^{-2}$)$^{-1}$, respectively, in f. The threshold current density becomes larger at lower temperatures, consistent with the experimental results in a.

This is consistent with Fig. 3b, justifying that an in-plane torque $\tau_{IP}^{AD}$ alone is unable to switch the magnetization from $m_z = -1$ to $m_z = 1$. Using the threshold $J_{th}$ for the $b$ axis at $T = 160$ K read off from Fig. 4a, we can evaluate in Fig. 4c the value of $\tau_{IP}^{AD}/j$. We emphasize that in the effective monodomain picture, this value does not have quantitative correspondence but only serves as an intermediate quantity in the simulation. With this $\tau_{IP}^{AD}/j$, we can then simulate the switching process for $I \parallel \hat{a}$ by varying $j$, and $\tau_{OP}^{AD}/j$, at 160 K. As shown in Fig. 4d, we find three different phases characterized by the final state of $m_z$: the unswitched phase ($m_z = -1$), the switched phase ($m_z = 1$) and the in-plane metastable phase ($m_z = 0$). These phases intersect at a tricritical point (Fig. 4d, green circle), which does not exist when $j$ flips sign. It is clear from Fig. 4d that the switched phase only exists when the ratio of $\tau_{OP}^{AD}/j$ is above the tricritical point, and the phase boundary of magnetization switching continuously shifts towards smaller $j$. This striking feature strongly indicates that the out-of-plane antidamping torque $\tau_{OP}^{AD}$ is essential for the magnetization switching in the perpendicular direction, consistent with the earlier micromagnetic simulations that consider both $\tau_{IP}^{AD}$ and $\tau_{OP}^{AD}$ (ref. 29). As a comparison, we also simulate the switching processes at 140 and 180 K and plot the phase boundaries along with that for 160 K in Fig. 4c,f, where we see that the threshold current for magnetization switching becomes larger (smaller) at lower (higher) temperatures, agreeing qualitatively with the observation in Fig. 4a. This behaviour is physically reasonable because at lower temperatures, the coercive field (which we adopt as the effective anisotropy) becomes larger so that the energy barrier for switching becomes higher. These results qualitatively explain the observed SOT switching in WTe$_2$/FGT bilayers by showing that $\tau_{IP}^{AD}$ and $\tau_{OP}^{AD}$ coexist, and in particular, that $\tau_{OP}^{AD}$ needs to be large enough to achieve field-free deterministic switching as the terminal state will otherwise be demagnetized from the metastable phases.

We have experimentally demonstrated that an out-of-plane antidamping SOT in WTe$_2$ is strong enough to enable thermally assisted field-free magnetization switching of a magnet with PMA. The result from monodomain simulation excludes the possibility of deterministic switching without a strong enough out-of-plane antidamping SOT. There are still open questions that need to be addressed in the future, including the exact mechanism of magnetization


switching observed in our devices, more comprehensive understanding of thermal effects on SOT switching, and the potential of realizing WTe₂-based scalable and room temperature functional SOT switching devices. However, the presence of a strong out-of-plane antidamping torque enabled by the special crystal symmetry in WTe₂ and its unique physical consequences make transition metal dichalcogenides with lower-symmetry crystal structure an appealing spin-source material for future SOT-related magnetic memory technologies.

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References
**Methods**

**Device fabrication.** WTe2, FGT and h-BN crystal flakes were prepared by previously published procedures. WTe2, h-BN and FGT were mechanically exfoliated on separate Si/SiO2(300 nm) substrates inside an argon-filled glovebox. The flakes were then optically inspected and selected using a fully automated microscope inside the glovebox. On a separate substrate, electrodes were defined using standard electron beam lithography with methyl methacrylate/poly(methyl methacrylate) bilayer resist; electron beam deposition was used for the platinum electrodes. The heterostructure was made with a custom-built transfer tool inside the glovebox using a polydimethylsiloxane stamp and thin film of polycarbonate. The final device consists, from top to bottom, of h-BN/FGT/WTe2/Pt with Cr(5 nm)/Au(110 nm) bond pads.

**SOT measurements.** The electrical measurements were performed at variable temperatures in high vacuum (<10⁻⁸ mtorr) conditions. An electromagnet was rotated such that the magnetic field can be applied in both in- and out-of-plane directions of the device. A Keithley 6221 current source and a Keithley 2182A nanovoltmeter are used for AHE hysteresis loop and current pulse-induced SOT-switching experiments. The current pulse used is a square current pulse with varying magnitude and a width of 100 μs (unless stated otherwise). The transverse resistance, Rxy, is measured with a smaller-magnitude current (50 μA) in delta mode (unless stated otherwise) to determine the magnetization state of the FGT. We define the normalized perpendicular magnetization m = M = M0, where M0 is the magnitude of the FGT, M is the magnetization magnitude and M0 is the saturation magnetization. In the SOT measurements, the initial magnetic state is prepared in m = ±1 (using an external magnetic field) and Ic is swept from zero to positive (negative) maximum values in steps of 250 μA. We define the threshold current Ith = (Ic+ + Ic−)/2, where Ic+ (Ic−) is the pulse current magnitude that drives the magnetization state across m = 0 at positive (negative) pulse sides. For SOT switching with pulse trains, a series of read current pulses of 500 μA were applied to read the magnetization state before and after the write current pulse was applied. AHE loop shift measurements are performed by measuring m(Hc), where at each perpendicular field Hc, a pulse current Ic(500 μs long) is applied and m is measured simultaneously in pulse delta mode. We define the loop shift field of the loop by Hc = (Hc+ + Hc−)/2, where Hc+ (Hc−) is the magnetic field for which the magnetization switches from down to up (up to down).

**Numerical simulations.** The effective single-domain dynamics is simulated by solving the Landau–Lifshitz–Gilbert equation in the presence of in-plane and out-of-plane antidamping torques as

\[ \dot{m} = -\gamma H_0 (J_c, t) \times m \times m + \alpha m \times (\hat{z} \times m + c_{ad} (\hat{y} \times \hat{z} \times m) + c_{ad} (\hat{x} \times \hat{z} \times m) + c_{ad} (\hat{y} \times m)) \]

where \( m \) is the unit vector of magnetization of FGT, \( t \) is time, \( H_0 \) is the effective easy-axis anisotropy (taken as the coercive field, \( H_c = H_0 \)), \( J_c \) is the charge current density, \( c_{ad} (\hat{y} \times m) \) and \( c_{ad} (\hat{x} \times m) \) are the amplitudes of the Oersted field, the in-plane antidamping torque, the out-of-plane antidamping torque and the out-of-plane field-like torque with respect to \( J_c \), respectively, and \( \alpha \) is the Gilbert damping parameter. To incorporate the heating effects, we determine the dependences of the coercive field on the current and the setting temperature from the hysteresis loop and the resistance measurements (Supplementary Notes 4 and 5). Based on the materials parameters and the device geometry, we take and set \( c_{ad} = 6.4 \times 10^{-2} \text{ mT} \times 10^6 \text{ A m}^{-2} \). In the simulations, we use \( c_{ad} = 0.5 \text{ mT} \times 10^6 \text{ A m}^{-2} \), but we find that \( c_{ad} = 0.5 \text{ mT} \times 10^6 \text{ A m}^{-2} \) has negligible impact on the switching process; the result does not show visible differences if \( c_{ad} = 0 \) or \( c_{ad} = 4 \text{ mT} \times 10^6 \text{ A m}^{-2} \). The current density \( j \), is applied along the \( \hat{x} \) direction, so the in-plane spin accumulation and the Oersted field are along the \( \hat{y} \) direction. The initial state \( \hat{m} \) is known to be along the \( \hat{z} \) direction with including a negligibly small deviation to reduce the incubation time, and the final state is obtained when the system reaches equilibrium.

**Data availability**

All the data supporting the findings of this study are available in the article and its Supplementary Information. Further information is available from the corresponding author on reasonable request.

**References**


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**Author contributions**

S.S. and J.K. designed the experiments and supervised the research. J.H.K., R.M., S.Y. and J.G. provided the bulk crystals and performed the experiments. H.Z. and R.C. prepared the bulk FGT crystals. J.L. and J.E.G. provided the h-BN/FGT/WTe2/Pt heterostructures. All authors contributed to writing the manuscript.

**Competing interests**

The authors declare no competing interests.

**Additional information**

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