Ferroelectric control of magnetic anisotropy of 5d transition metal monolayers

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Electric field control of magnetism using ferroelectric materials offers promising applications in low-power spintronics. We have conducted systematic density functional theory calculations to investigate the magnetic properties of epitaxial 5*d* transition metal monolayers on a ferroelectric substrate: PbTiO₃ (PTO). Our study reveals that the magnetocrystalline anisotropy energy of the osmium monolayer is significantly enhanced to 18.1 meV/Os in the upward polarization state and, moreover, the anomalous Hall coefficient of osmium/PTO film is tunable by reversing the electric polarization of the PTO substrate. Additionally, the electric polarization reversal in PTO rotates the easy magnetization axis of the iridium and platinum monolayers by 90°, which is attributed to the rearrangement of interfacial charges. Our findings suggest an efficient approach to control the magnetization direction of monoatomic layers and provide valuable insights for the development of low-energy spintronics devices.

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I. INTRODUCTION

The control of magnetic properties of ultrathin films with electric field has attracted intensive research interest, as it involves new fundamental physics about electromagnetic coupling in materials and is promising for technological development of nonvolatility and energy-efficient spintronics devices [1-6]. To realize robust responses, numerous prototypical material systems have been proposed, primarily based on either single-phase multiferroic materials [7] with several coexisting ferroic orders, or composite multiferroics [8] that combine ferroelectric (FE) and ferromagnetic (FM) materials. Single-phase multiferroic materials usually have weak magnetoelectric coupling due to their low magnetization and antiferromagnetic order, or the low temperature required for the multiferroicity. These issues can be somewhat overcome by using composite multiferroic systems such as FE-FM heterostructures [9], in which the magnetization, magnetic order, magnetic anisotropy, and even the domain structure can be manipulated by electric field. The application of electric biases across ultrathin FE-FM films may not only change the strain or electric polarization in the FE layers [10], but also may induce charge transfer and orbital rehybridization [11] between the FE and FM components, and hence directly change the magnetic properties of the FM layers.

In particular, magnetic anisotropy energy (MAE) determines the direction of easy magnetization, which may further alter other magnetic properties such as the tunnel magnetoresistance [12] in magnetoresistive random access memory devices [13]. Therefore, the control of MAE by ferroelectrics is one of the frontiers for the design of low-dimensional multifunctional materials and devices. Many composite multiferroic heterostructures have been attempted either theoretically or experimentally, such as Ni/Pb($Zn_{1/3}Nb_{2/3}$)O₃ -PbTiO₃(PZN – PT), CoFe₂O₄/PZN-PT [14], Ni/BaTiO₃ [15], La_{0.88}Sr_{0.1}MnO₃/BaTiO₃ [16], MnPt/BaTiO₃ [17], 5d capping metal/CoFe/BaTiO₃ [18], FePt/PbTiO₃ [19], and IrCo/PbTiO₃ [20]. Although these studies cover rather broad FM materials, from ferromagnetic oxides, heavy-metal FM metal bilayers, to single magnetic molecules [21] or atoms [22,23], the effect of electric field or electric polarization on the magnetic properties of 5d transition metal (TM) monolayers has rarely been investigated. It is interesting to examine the interplay between strong, intrinsic spin-orbital coupling (SOC), low-dimensional magnetization, and local electric polarization at interfaces of 5d TM layers and FE substrates for the design of new materials with strong magnetoelectric coupling.

In this work, we systematically investigate the effect of the electric polarization of the PbTiO₃ (PTO) substrate on magnetic properties of different 5d TM monolayers. Through first-principles calculations for TM/PTO/palladium (TM = Hf, Ta, W, Re, osmium, Ir, Pt) films, we find that a large magnetoelectric coupling effect may present in at last three such systems, for which both the sign and magnitude of MAE can be tuned by reversing the electric polarization in PTO. For example, MAE of the osmium monolayer is significantly enhanced to 18.1 meV/osmium, as the PTO is in the upward FE polarization state. Furthermore, the FE reversal changes the sign of MAE for the iridium or platinum monolayer. The large magnetoelectric coupling arises from the field-induced charge rearrangement at the interface. Our work shows the controllability of MAE and other magnetic properties of 5dthin layers with reversible FE polarization, and paves the way for the development of energy-efficient spintronic devices.

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FIG. 1. (a) Side view and (b) top view of the TM/PbTiO₃/ palladium system with the electric polarization of PbTiO₃ pointing upward (P_{\uparrow}) and downward (P_{\downarrow}), respectively. The orange spheres represent 5*d* TM atoms.

II. MODELS ANDS METHODS

To simulate the FE control of the magnetic properties of 5d TM monolayers, we constructed a slab model for the TM/PTO/palladium heterostructures as shown in Fig. 1. The PTO substrate was modeled by three cubic cells with TiO₂ as the outmost layers, because the TiO₂ terminated surface is the most stable one for PTO [11,20]. A palladium layer was added on the bottom surface to mimic the electrode and to eliminate unwanted magnetization of the bottom TiO₂ surface. One TM layer was stacked on the PTO substrate, with the TM atoms sitting above the topmost oxygen atoms. A vacuum space of 16 Å was inserted between adjacent slabs to avoid artificial interactions between them.

All calculations were performed based on density functional theory, using the projector augmented wave method [24,25] as implemented in the Vienna Ab initio Simulation Package [26,27]. The exchange-correlation potential was described within the framework of the generalized gradient approximation with the Perdew-Burke-Ernzerhof functional [28]. The cutoff energy for the plane-wave basis expansion was set to 550 eV. A convergence criterion of 0.01 eV/Å for the forces was used during the structural relaxation procedure. The two-dimensional Brillouin zone of the TM/PTO/palladium slab was sampled with a $13 \times 13 \times 1$ and a denser $23 \times 23 \times 1$ k-point mesh for the colinear and MAE calculations, respectively. The strong correlation effect for the 5d electrons was treated by an effective onsite Hubbard term [29] of $U_{\text{eff}} = 3 \text{ eV}$. Different effective Hubbard U values (0 to 5 eV) were tested and no qualitative changes were found (see Supplemental Material Fig. S3 [30]). The dipole correction [31] was considered to correct the errors introduced by the periodic boundary conditions along the surface normal for the nonsymmetric slabs.

Our optimized in-plane lattice constant of PTO (a = b = 3.84 Å) is slightly smaller than the experimental lattice size of 3.90 Å [32,33]. We used the same lattice constant for all different 5*d* TM monolayers, because the substrate in experiments is typically more than a few

nanometers thick and the metal monolayers are epitaxially grown on it, such as lead on Si(111) [34], silver on Ge(111) [35] and indium on Si(111) [36]. The optimized lattice constants of isolated 5d TM monolayers are smaller than that of PTO, except hafnium (see Supplemental Material Table SI [30]). The binding energies of these 5*d* monolayers on PTO are also large (see Supplemental Material Table SI [30]), so it is reasonable to assume that PTO effectively grasps adatoms, allowing for the formation of stable epitaxial structures even without perfect lattice matching. Taking osmium/PTO/palladium as an example, we performed ab initio molecular dynamics simulations with AN Andersen thermostat [37] at 300 K using a $3 \times 3 \times 1$ supercell, and demonstrated the dynamical stability of these 5d monolayers on the PTO substrate (see Supplemental Material Fig. S1 [30]). We also considered the other stacking model and calculated the MAEs for the osmium, iridium and platinum monolayers. The detailed comparisons are shown in Supplemental Material Fig. S2, and Supplemental Material Tables SII and SIII [30]). In the FE phase, the FE polarization of PTO points along the [001] direction. To preserve the FE phase, atomic positions of the two lower PTO cubic cells were fixed during the structural optimization. For the convenience of following discussions, we denote the upward electric polarization as P_{\uparrow} and downward as P_{\downarrow} .

In general, MAE results from the magnetocrystalline anisotropy (MCA) given by SOC and the shape anisotropy energy (SAE) given by the magnetostatic dipole-dipole interactions. The MCA term can be calculated using the total-energy method or the torque method proposed by Wang *et al.* [38,39]:

$$\mathbf{MCA} = E_{\mathbf{SOC}}^{x} - E_{\mathbf{SOC}}^{z} = \sum_{\mathbf{occ}} \langle \psi_{i,k} | \frac{\partial H_{\mathbf{SOC}}}{\partial \theta} | \psi_{i,k} \rangle |_{\theta = \mathbf{45}^{\circ}} \quad (1)$$

where $\psi_{i,k}$ is the wave function of the *i*th state at the *k*-point, H_{SOC} is the SOC Hamiltonian, and θ is the polar angle of the spin. The SAE term (SAE = $E_{\text{dipole}}^x - E_{\text{dipole}}^z$) typically exists in ultrathin films, and the energy of magnetic dipole interactions can be expressed as

$$E_{\text{dipole}} = \frac{1}{2} \frac{\mu_0}{4\pi} \sum_{i,j} \frac{m^2}{r_{ij}^3} (1 - 3\cos^2\theta_{ij}), \qquad (2)$$

where θ_{ij} is the angle between magnetic moments with the separation \vec{r}_{ij} . The SAE is typically a few tenths of millielectron volts per atom or smaller [40,41], so we mainly focus on MCA in the following discussions.

III. RESULTS AND DISCUSSION

The magnetic properties of TM/PTO/palladium (TM = Hf, Ta, W, Re, Os, Ir, Pt) are shown in Fig. 2. From Fig. 2(a) and (b), we can see that all 5*d* elements except tungsten have nonzero spin moments. In particular, osmium and iridium have spin moments larger than 1.0 μ_B , and their MAEs are also large: up to 18 meV per atom for Os. As shown in Fig. 2(a), their orbital magnetic moments are considerably larger than those of 3*d* TMs, because of the strong SOC. As magnetic moments and MAEs of hafnium, tantalum, and rhenium are noticeably small, we believe that they are not



FIG. 2. The magnetic properties per TM atom in the TM/PTO/palladium (TM = Hf, Ta, W, Re, Os, Ir, Pt) system. (a) Spin moments and orbital moments. The insets show the spin density isosurfaces of osmium/PTO/palladium in the P_{\uparrow} and P_{\downarrow} states. (b) MAEs were calculated by the torque method.

suitable for the development of magnetic materials. The spin and orbital magnetic moments of platinum are also small, so platinum might not be a good candidate element either, but we will include platinum in following discussions due to its interesting MAE results.

Interestingly, osmium has the largest spin moments and MAEs in the two polarization states. The spin density of osmium/PTO/palladium is shown as the insets in Fig. 2(a). It is obvious that the isosurfaces show contributions from different *d*-states to the magnetization when the FE state of PTO is switched from P_{\downarrow} to P_{\uparrow} . Accordingly, the spin moment of osmium increases from 1.56 $\mu_{\rm B}$ to 2.40 $\mu_{\rm B}$ with the reversal of polarization. The easy magnetization axis of osmium remains out of plane, but its MAE increases from 4.89 meV to 18.09 meV. Moreover, the reversal of electric polarization changes the sign of MAEs of the iridium and

platinum monolayers from -4.40 meV to 2.74 meV for iridium and from -1.32 meV to 0.91 meV for platinum, respectively. This means that the direction of their easy magnetization axis can be switched between in-plane and out-of-plane when the external electric field is reversed. The results of fourfold enhancement of MAE and the magnetic reorientation reveal the large magnetoelectric coupling effect in these heterostructures. For comparison, the easy axis of the isolated osmium and platinum monolayers is out of plane, and the easy axis of the isolated iridium monolayer is in plane (see Supplemental Material Table SIV [30]).

The remarkable magnetoelectric coupling effects depend on the charge transfer at the interface, which can be significantly affected by FE reversals. As an example, we use the iridium/PTO/palladium heterostructure as an example to demonstrate this point. The charge density difference and



FIG. 3. Charge density difference and layer-resolved PDOS of iridium/PTO/palladium for the (a) P_{\uparrow} and (b) P_{\downarrow} states. The charge density difference is shown in the iridium-titanium plane for the P_{\uparrow} state and the iridium-oxygen plane for the P_{\downarrow} state. The red and blue regions represent charge accumulation and depletion, respectively.

TABLE I. Bader charges of atoms in TM/PTO/palladium ($TM = Os$, Ir, Pt) at the upper interface. The values in parentheses are the charge
difference between the TM/PTO/palladium heterostructure, the bare PTO/palladium surface, and the freestanding TM monolayer. Positive and
negative signs indicate gain and loss of electrons.

TM/PTO/palladium		0	Ti	TM
Osmium/PTO/palladium	P_{\uparrow}	7.164 (-0.057)	2.101 (-0.142)	8.124 (+0.124)
	$P_{\downarrow} \ P_{\uparrow} - P_{\downarrow}$	+0.135	+0.209 (+0.055)	+0.338 (-0.214)
Iridium/PTO/palladium	P_{\uparrow}	7.145 (-0.070)	2.000 (-0.287)	9.212 (+0.212)
	P_{\downarrow}	6.995 (+0.115)	1.878 (+0.039)	8.835 (-0.165)
Platinum/PTO/palladium	$P_{\uparrow} - P_{\downarrow}$ P_{\uparrow} P_{\downarrow}	+0.130 7.137 (-0.080) 6 994 (+0 109)	+0.122 1.968 (-0.322) 1.908 (+0.032)	+0.377 10.240 (+0.240) 9 850 (-0.150)
	$P_{\uparrow} \stackrel{r_{\downarrow}}{-} P_{\downarrow}$	+0.143	+0.060	+0.390

layer-resolved projected density of states (PDOS) for the P_{\uparrow} and P_{\perp} states are shown in Fig. 3. It is obvious that the charge transfer is centered around iridium and titanium atoms for the P_{\uparrow} state, but it is around the topmost oxygen atoms for the P_{\downarrow} state. The Bader charges in Table I show that each iridium atom gains 0.213 e and each titanium atom loses 0.287 e in the P_{\uparrow} state compared to the reference systems, i.e., the bare PTO/palladium surface and freestanding iridium monolayer. On the contrary, the iridium atom loses 0.165 e and the oxygen atom gains 0.115 e in the P_{\downarrow} state. A similar charge transfer also exists for osmium and platinum systems, as shown in Table I. Such charge redistributions partially compensate the polarization effect from PTO as they create opposite electric dipoles at the interface against the original polarization, as marked by the arrows in Fig. 3. The induced dipoles also somewhat reduce the internal electric polarization in PTO, as demonstrated by the layer-resolved PDOS curves that show very mild misalignment in the PTO1 and PTO2 layers. The macroscopic average of electrostatic potential [42] also confirms the fading of the local electric field in PTO after the iridium monolayer is stacked on the PTO/palladium surface (see Supplemental Material Fig. S4 [30]). This phenomenon is consistent with previous theoretical results for the platinum/PTO/platinum capacitor [43]. At the interface, the opposite charge transfers for the P_{\uparrow} and P_{\downarrow} states affect the electronic structure very differently, as seen from the PDOS of iridium in Fig. 3(a) and 3(b). Therefore, PTO alters the magnetic properties of 5d TM monolayers via field-driven charge transfer at the interface.

We further plot the PDOS curves of *d*-orbitals of osmium, iridium, and platinum on PTO/palladium and their dependences of MAE on the Fermi level from the rigid band model analysis in Figs. 4(a)–4(c). The splitting of states in two spin channels tapers off from osmium to platinum for both polarization states. As seen in Figs. 4(d)–4(f), MAEs are sizeable and oscillate in the energy range. For each system, the polarization reversal significantly modifies the electronic structure and, consequently, the MAE, as seen from the noticeable separation between the red and blue lines. This implies the large magnetoelectric coupling. With a small shift of E_F , the MAE of osmium may remain roughly unchanged. In contrast, the MAE of iridium may double in magnitude in both polarization states if we shift the Fermi level down by as little as 0.2 eV. The MAE of platinum is also sensitive to the downward shift of E_F in the P_{\uparrow} case. It is worth noting that these systems have large SOC, and the perturbative rigid band model may only provide qualitative guidance for the manipulation of their MAEs. Nevertheless, these trends provide useful guidance for experiments, as the variation of the Fermi level can be achieved by doping or applying a bias field.

Among these three systems, iridium/PTO/palladium appears to be the most promising one for the implementation of a controllable spin device. Thus, we will delve on the mechanism of the sign change of the MAE in the iridium/PTO/palladium heterostructure. Within the second-order perturbation approach [44], the MAE can be split into three parts:

$$E_{\rm MCA} = E_{\rm MCA}^{\uparrow\uparrow} + E_{\rm MCA}^{\downarrow\downarrow} + E_{\rm MCA}^{\uparrow\downarrow+\downarrow\uparrow}, \qquad (3)$$



FIG. 4. PDOS of *d*-orbitals of the TM atom (a)–(c) and dependences of MAE on the Fermi level from the rigid band model (d)–(f) for three TM/PTO/palladium systems (TM = Os, Ir, Pt from top to bottom).



FIG. 5. The MAE differences of iridium/PTO/palladium in the first Brillouin zone between the P_{\uparrow} state and the P_{\downarrow} state. (a) The total MAE difference. The contribution from (b) the occupied spin-up state and unoccupied spin-up state, (c) the spin-down state and spin-down state, and (d) the spin-up state and spin-down state. The red and blue solid circles represent the positive and negative values. The area of circles denotes the magnitude of the difference. (e) The MAE difference is denoted by the black line along high symmetric lines. The colored line and regions represent contributions to MAE from different spin states.

with

$$E_{\text{MCA}}^{\uparrow\uparrow} = \xi^{2} \sum_{o^{\uparrow}, u^{\uparrow}} \frac{|\langle o^{\uparrow} | L_{z} | u^{\uparrow} \rangle|^{2} - |\langle o^{\uparrow} | L_{x} | u^{\uparrow} \rangle|^{2}}{\varepsilon_{u}^{\uparrow} - \varepsilon_{o}^{\uparrow}},$$

$$E_{\text{MCA}}^{\downarrow\downarrow\downarrow} = \xi^{2} \sum_{o^{\uparrow}, u^{\uparrow}} \frac{|\langle o^{\downarrow} | L_{z} | u^{\downarrow} \rangle|^{2} - |\langle o^{\downarrow} | L_{x} | u^{\downarrow} \rangle|^{2}}{\varepsilon_{u}^{\downarrow} - \varepsilon_{o}^{\downarrow}},$$

$$E_{\text{MCA}}^{\uparrow\downarrow\downarrow} = \xi^{2} \sum_{o^{\uparrow}, u^{\downarrow}} \frac{|\langle o^{\uparrow} | L_{x} | u^{\downarrow} \rangle|^{2} - |\langle o^{\uparrow} | L_{z} | u^{\downarrow} \rangle|^{2}}{\varepsilon_{u}^{\downarrow} - \varepsilon_{o}^{\uparrow}},$$

$$E_{\text{MCA}}^{\downarrow\uparrow\uparrow} = \xi^{2} \sum_{o^{\uparrow}, u^{\downarrow}} \frac{|\langle o^{\downarrow} | L_{x} | u^{\downarrow} \rangle|^{2} - |\langle o^{\downarrow} | L_{z} | u^{\downarrow} \rangle|^{2}}{\varepsilon_{u}^{\downarrow} - \varepsilon_{o}^{\uparrow}}.$$
(4)

Here, ξ is the strength of SOC, arrows represent spins, and ε_o^{\uparrow} and $\varepsilon_u^{\downarrow}$ are the energy levels of the occupied spin-up state and unoccupied spin-down state, respectively. In order to confirm which part dominates the change of MAE, we plot the distribution of MAE(P_{\uparrow})-MAE(P_{\downarrow}) in the first Brillouin zone, as seen in Fig. 5(a). It shows that the strong electromagnetic coupling of iridium/PTO/palladium stems from contributions in a pocket between the Γ and X' points. From the spin channel splittings in Figs 5(b) and 5(c), we may see that contributions from the $\Delta E_{MCA}^{\uparrow\uparrow}$ and $\Delta E_{MCA}^{\downarrow\downarrow}$ terms are approximately canceled in the entire Brillouin zone. Thus, the $\Delta E_{MCA}^{\uparrow\downarrow+\downarrow\uparrow\uparrow}$ term plays the leading role in the sign change of MAE induced by FE reversal, as shown in Fig. 5(d). Furthermore, the MAE differences along high symmetric lines are plotted in Fig. 5(e). We may reach the same conclusion, as the dominant peaks occur on the $\Gamma - X'$ path.

the dominant peaks occur on the $\Gamma - X'$ path. As a further step, the $E_{\text{MCA}}^{\uparrow\downarrow+\downarrow\uparrow}$ term and corresponding spin-resolved band structures of iridium/PTO/palladium along $\Gamma - X'$ are shown in Fig. 6 for two FE states. The *d*-orbitals that contribute to the $E_{MCA}^{\uparrow\downarrow\downarrow\downarrow\downarrow\uparrow}$ term are marked on band structures (nonzero matrix elements only include $\langle z^2 | L_x | yz \rangle =$ $\sqrt{3}$, $\langle xy|L_x|xz\rangle = 1$, $\langle x^2 - y^2|L_x|yz\rangle = 1$, $\langle xz|L_z|yz\rangle = 1$, and $\langle x^2 - y^2 | L_z | xy \rangle = 2$). As shown in Fig. 6(a), the large negative MAE in the P_{\uparrow} state results from the matrix element $\langle x^2 - y^2, \uparrow | L_z | xy, \downarrow \rangle$ at the 2/3($\Gamma - X'$) path (region II), where $|xy, \downarrow\rangle$ becomes occupied and the two states have a small energy difference. In the P_{\downarrow} state, the negative contributions in region II significantly decrease because the unoccupied spin-up $d_{x^2-y^2}$ shifts away from Fermi level. On the contrary, the MAE is greatly enhanced in region I because the matrix element $\langle xz, \downarrow | L_x | xy, \uparrow \rangle$ provides large positive contributions, with the occupied spin-down d_{xz} approaching the Fermi level. Therefore, the total MAE changes from negative to positive as the P_{\uparrow} state switches to the P_{\perp} state.

Finally, we noted that osmium has a large magnetic moment and positive MAE in both FE states. This inspires us to investigate further the existence of the anomalous Hall effect and its controllability with the reversal of the FE polarization. Here, we focus on the intrinsic anomalous Hall conductivity (AHC) and calculate the Hall conductivity σ_{xy} from the Berry



FIG. 6. The MAE from the contribution between different spins and the spin-resolved band structure along $\Gamma - X'$ in (a) the P_{\uparrow} state and (b) the P_{\downarrow} state of iridium/PTO/palladium. The red and blue points indicate spin up and spin down, respectively. Point size represents the weight of the *d*-orbitals, but for clarity, only orbitals dominating the change of MAE are marked.

curvature over the Brillouin zone [45] as

$$\sigma_{xy} = -\frac{e^2}{h} \int_{\text{BZ}} \frac{d^3k}{(2\pi)^3} \Omega(k), \qquad (5)$$

where $\Omega(k)$ is determined as

-1.0

-1.5

-2.0

Μ

Here, f_n is the Fermi occupation factor, v_x and v_y are velocity operators, and ψ_{nk} and $\hbar\omega_n$ are the spinor Bloch wavefunction and eigenvalue of the *n*th band, respectively.

The calculated AHC of osmium/PTO/palladium is 1995 $(\Omega \text{ cm})^{-1}$ and 664 $(\Omega \text{ cm})^{-1}$ for the P_{\uparrow} state and the P_{\downarrow} state, respectively. The large variation also gives an efficient means to control the transport properties of osmium/PTO/palladium with electric polarization. For the P_{\uparrow}



FIG. 7. Berry curvature and spin-resolved band structure (with SOC involved) along high symmetric lines. (a) and (b) are for the P_{\uparrow} state, and (c) and (d) are for the P_{\downarrow} state of the osmium/PTO/palladium heterostructures.

state, the AHC is mainly attributed to the large Berry curvature near 1/2(M-X), as seen in Fig. 7(a). The spin-resolved band structure with SOC is shown in Fig. 7(b). The spin-up d_{xy} and spin-down d_{xz} states interact across the Fermi level, as marked by red circles, resulting in the large Berry curvature because the small energy splitting induced by SOC causes small denominators in Eq. (6). While for the P_{\downarrow} state, the negative Berry curvature appears near $1/3(X-\Gamma)$ and makes the AHC decrease, as seen in Fig. 7(c). We found that the negative Berry curvature results from the intermixing of two spin components in the d_{xy} -orbital, as marked by the blue circle in Fig. 7(d). This reveals the reason why the AHC of the osmium monolayer in the P_{\downarrow} state is smaller than that in the P_{\uparrow} state.

IV. CONCLUSION

In summary, we have conducted systematic density functional theory calculations to investigate the influence of electric polarization switching in perovskite titanate on the magnetic properties of 5d transition metal monolayers. We

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found that the magnetocrystalline anisotropy energy of the osmium monolayer is significantly enhanced from 4.9 meV per osmium in the P_{\downarrow} state to 18.1 meV per osmium in the P_{\uparrow} state. Additionally, the osmium monolayer exhibits a large anomalous Hall conductivity, the magnitude of which can be tuned by ferroelectric reversals of the PTO substrate. The FE reversal may also change the sign of the MAEs for the iridium and platinum monolayers, primarily resulting from the polarization-induced charge rearrangements at the interface. Our findings present an efficient approach to modulating the magnetic properties of two-dimensional systems and provide valuable insights for further development of TM/FE heterostructures in spintronic applications.

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