Half-Dirac semimetals and the quantum anomalous Hall effect in Kagome Cd$_2$N$_3$ lattices

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Half-Dirac semimetals (HDSs), which possess 100% spin-polarizations for Dirac materials, are highly desirable for exploring various topological phases of matter as low-dimensionality opens unprecedented opportunities for manipulating the quantum state of low-cost electronic nanodevices. The search for high-temperature HDSs is still a current hotspot and yet challenging experimentally. Herein based on first-principles calculations, we propose the realization of Half Dirac semimetals (HDS) in two-dimensional (2D) Kagome transition-metal nitride Cd$_2$N$_3$, which is robust against strain engineering. Monte Carlo simulations reveal that Cd$_2$N$_3$ possesses a Curie temperature reaching up to $T_C = 225$ K, which is much higher than that of the reported monolayers CrI$_3$ ($T_C = 45$ K) and Cr$_2$Ge$_2$Te$_6$ ($T_C = 20$ K). The band crossings in Cd$_2$N$_3$ are gapped out by the spin–orbit coupling, which brings about the quantum anomalous Hall (QAH) effect with a sizeable band gap of $E_g = 4.9$ meV, characterized by the nonzero Chern number ($C = 1$) and chiral edge states. A tight-binding model is further used to clarify the origin of HDSs and nontrivial electronic properties. The results suggest monolayer transition-metal nitrides as a promising platform to explore fascinating physical phenomena associated with novel 2D emergent HDSs and QAH insulators toward realistic spintronics devices, thus stimulating experimental interest.

1. Introduction

Spintronics, using an electron’s spins instead of its charge to carry information, has attracted extensive attention because of faster transport and dissipation-less properties in magnetic materials.\textsuperscript{1–2} To build spintronic devices, the selection of ferromagnetic (FM) materials and control of ferromagnetism, such as the generation and injection of spin-polarized carriers, the manipulation and detection of the spin direction, and the long-distance spin-polarized transport, are crucial and yet challenging. Many concepts of spintronic materials have been put forward to solve these problems, including the use of half-metals (HMs),\textsuperscript{3–5} topological insulators (TIs),\textsuperscript{6–8} magnetic semiconductors,\textsuperscript{9–12} topological nodal line semimetals (TNLSs),\textsuperscript{13–15} and so on. Interestingly, two-dimensional (2D) FM ordering has been experimentally discovered in van der Waals crystals CrI$_3$ ($T_C = 45$ K)\textsuperscript{16} and Cr$_2$Ge$_2$Te$_6$ ($T_C = 20$ K),\textsuperscript{17} though their $T_C$ is lower than room temperature, opening the door for the application of 2D spintronics at the nanoscale. Spin-gapless semiconductors (SGSs), as shown in Fig. 1(a), possessing a band gap in one of the spin channels while the others keeping a sizeable zero gap in the Fermi level, are more interesting and have been proposed by Wang in 2008.\textsuperscript{18} Owing to the special zero band-gap character of SGSs, these types of materials exhibit many novel properties, including (i) only a tiny amount of energy is required to excite electrons from the valence band (VB) to the conduction band (CB); (ii) the excited charge carriers, both electrons and holes, can be 100% spin-polarized simultaneously; (iii) by using the Hall effect, 100% spin-polarized electrons and holes can be easily separated. Many 2D and 3D SGSs,\textsuperscript{19–26} including TMF$_3$ (ref. 15) and MX (M = Li, Na, K, Rb, Cs; X = S, Se, Te)\textsuperscript{28} exhibit various topological phases. Also, some examples have been synthesized such as Mn$_2$CoAl,\textsuperscript{19} and CoFeMnSi.\textsuperscript{20} All these provide a new playground and opportunities for spintronic, electronic and optic devices.

Dirac semimetal is a new type of quantum state with nontrivial topological properties. Its body energy band generally has...
a three-dimensional Dirac point near the Fermi surface. The representative example is monolayer graphene composed of a purely 2D honeycomb network of carbon atoms. Low-energy excitation in graphene is governed by π electrons in 2p orbitals, whose energy dispersion has linear band crossings with the Dirac-like node at the Fermi level, called the Dirac cones. It exhibits unique electronic properties such as ballistic charge transport, high carrier mobility, and the quantum Hall effect. Subsequently, many 2D lattices such as germanene, carbon allotropes, boron allotropes, SiC$_3$, and C$_3$N$_5$ have been identified as Dirac materials. However, the Dirac cone in these materials is intrinsically spin-degenerated, which actually limits its utilization in massless spintronics. Thus, the investigation of the controllability of the band crossings and topological phases by using the p-orbital degrees of freedom is desirable, but such an interesting possibility has rarely been investigated.

The ultimate goals for future spintronic devices are to require ultra-fast transport and ultra-lower energy consumption, that is, eliminating the effective mass of electrons or holes and enabling massless charges to be completely spin-polarized. In comparison with parabolic-like SGSs (Fig. 1(a)) and four-degenerated band Dirac cones (Fig. 1(b)), Dirac-like SGSs, termed HDSs, as illustrated in Fig. 1(c), are the better class of magnetic materials for use in spintronic devices as HDSs possess more exotic spin and charge states. Thus, it is highly desirable to find stable HDSs with an intrinsic 100% spin-polarized Dirac state. Normally, HDSs can be divided into two types: type-I, the d-state HDSs, in which the Dirac state is contributed by the d-orbital of transition-metal atoms, and type-II, the p-state HDSs, whose states are from the p-orbital of transition-metal atoms. To date, the majority of reported HDS members belong to type-I HDSs, while there has been barely any reports on 2D type-II HDS materials. It is highly desirable to design type-II HDSs with robust spin ordering and a high Curie temperature simultaneously, which may make experimental synthesis largely accessible.

Meanwhile, among the 2D materials family, 2D Kagome lattices have been attracting great interest due to their rich properties. Most members of the family are carbides and oxides (e.g., Be$_2$C$_2$, V$_2$O$_3$, Nb$_2$O$_3$, and LiFeSe), consisting of one flat layer with $D_{6h}$ symmetry. Considering that carbon has one electron more than the oxygen atom, it is natural to ask whether transition-metal nitrides can also be stabilized in a 2D crystal Kagome lattice. If so, this may open the door to a new family of 2D materials, i.e., 2D Kagome transition-metal nitrides, which will also grow into a big family with fascinating quantum properties.

In this work, we try to address the above challenge and question by proposing a new family of 2D transition-metal nitrides with the Kagome structure. Based on first-principles calculations, we show that Cd$_2$N$_3$ can be stabilized in this 2D lattice as shown in Fig. 2(a), which is dynamically and thermally stable. Monte Carlo simulations on the basis of the 2D Heisenberg Hamiltonian model show that the Curie temperature $T_C = 225$ K, having potential for high temperature applications. The coexistence of fully spin-polarization and linear dispersion Dirac states can be realized in Cd$_2$N$_3$, which features ultra-clean energy dispersion and ultra-high Fermi velocity. Distinguished from most proposed nonmagnetic Dirac materials, the four-fold degeneracy neck crossing-point traces out two-fold degeneracy lines emerged in the single-spin channel and the profile of HDSs is revealed by a tight-binding (TB) model. To the best of our knowledge, the Cd$_2$N$_3$ monolayer represents the first example that hosts the 2D HDS in the equilibrium state. When turning on the spin–orbit coupling (SOC), Cd$_2$N$_3$ becomes a quantum anomalous Hall (QAH) insulator with its energy gap of $E_g = 5.0$ meV, larger than recently predicted $E_g = 2.3$ meV in Mn–DCA, which is characterized by the nonzero Chern number ($C = 1$) and chiral edge states. In comparison to Cr- or V- doped (Bi, Sb)$_2$Te$_3$ films such a lattice without any magnetic doping is easier to synthesize and has much higher homogeneity. Thus, our discovery not only predicts a family of new 2D HDS materials, which vastly enriches the SGS family, but also provides an experimentally feasible platform to explore the new emergent QAH effect and their fascinating fundamental physical effects.

2. Computational details and method

All the spin-polarized calculations are performed by using density-functional theory (DFT), as implemented in the Vienna $ab$ initio simulation package (VASP).

The projector-augmented-wave (PAW) potential, Perdew–Burke–Ernzerhof (PBE) exchange-correlation functional plane-wave basis with a kinetic energy cutoff of 500 eV are employed. An all-electron description, the projector augmented wave method, is used to describe the electron-ion interaction. The Brillouin zone is sampled by using a $9 \times 9 \times 1$ Gamma-centered Monkhorst–Pack grid, and the SOC is included by a second
variational procedure on a fully self-consistent basis. During structural optimization, all atomic positions and lattice parameters are fully relaxed, and the maximum force allowed on each atom is less than 0.02 eV Å\(^{-1}\). Spin polarization is included through all the calculations. Since standard DFT may fail to describe the magnetism, the Heyd–Scuseria–Ernzerhof (HSE06) functional\(^*\) has been carried out to examine the magnetism and electronic structures. Phonon dispersion is calculated using the finite displacement method, as implemented in the PHONOPY code.\(^{25}\)

3. Results and discussion

Because most elements and compounds are not favorable to directly form a 2D Kagome lattice, we propose a new transition-metal nitride, Cd\(_2\)N\(_3\), to substitute the Kagome lattice, as illustrated in the top and side views of Fig. 2(a). The Cd\(_2\)N\(_3\) monolayer forms a honeycomb lattice with a point symmetry of \(D\text{}_{\text{6h}}\), by sharing one nitrogen (N) bridge with each neighbor, reminiscent of the crystal structure in the 2D organic Mn-DCA lattice.\(^{39-43,47,48}\) The optimized lattice parameters for Cd\(_2\)N\(_3\) are \(a_1 = a_2 = 7.29\) Å, and three N atoms locate around each Cd atom and the bond angle N–Cd–N is 120°, with a Cd–Cd distance of \(d = 4.21\) Å. To characterize the thermodynamic stability of Cd\(_2\)N\(_3\), we calculate the cohesive energy \(\Delta E_{\text{c}}\) defined as

\[
\Delta E_{\text{c}} = E(\text{Cd}_2\text{N}_3) - 2\mu(\text{Cd}) - 3/2\mu(\text{N}_2)
\]

where \(E(\text{Cd}_2\text{N}_3)\) is the total energy, and \(\mu(\text{Cd})\) and \(\mu(\text{N}_2)\) are the chemical potentials taken from bulk Cd and gas phase N\(_2\). The calculated result for Cd\(_2\)N\(_3\) is \(\Delta E_{\text{c}} = -0.67\) eV per atom, indicating an exothermic chemical reaction. These values are comparable to those of monolayers Cr\(_3\) (−0.903 eV per atom)\(^{39}\) and Cr\(_2\)-Ge\(_2\)Te\(_6\) (−0.552 eV per atom).\(^{40}\) The moderate formation energy also indicates that Cd\(_2\)N\(_3\) can be realized with a strong bonding network.

It is important to examine whether the experimentally fabricated Cd\(_2\)N\(_3\) is stable and whether the magnetic state survives at room temperature. To test the kinetic stability of Cd\(_2\)N\(_3\), we perform phonon spectrum calculations for Cd\(_2\)N\(_3\), as shown in Fig. 1(b). Apparently, there is no imaginary mode in the whole Brillouin zone, indicating that the Cd\(_2\)N\(_3\) lattice is dynamically stable. Additionally, the thermal stability of the Cd\(_2\)N\(_3\) lattice is assessed by performing \textit{ab initio} molecular dynamics (MD) simulations. We use a \(4 \times 4\) supercell to carry out the individual simulations at a temperature of 300 K. As seen from Fig. 2(c), the crystal structure of Cd\(_2\)N\(_3\) does not collapse throughout 8 ps MD simulation. In fact, Wang \textit{et al.}\(^{46}\) reported that the Y\(_2\)O\(_3\) lattice forms a complete monolayer on platinum-supported graphene. The Y\(_2\)O\(_3\) monolayer interacts weakly with graphene, but is stable at high temperature. Scanning tunnelling microscopy reveals that Y\(_2\)O\(_3\) exhibits a 2D hexagonal lattice rotated by 30° relative to the hexagonal graphene lattice. The above results reveal that the Cd\(_2\)N\(_3\) monolayer has very good thermal stability and could possibly be prepared in the experiment and survive at room temperature.

We now focus on the magnetic structures of Cd\(_2\)N\(_3\). Spin-polarized results reveal a spin moment of 5.0 \(\mu_B\) per cell as a result of large Hund’s rule, which can be further illustrated by spin-polarized electron density, \(\Delta \rho = \rho_1 - \rho_2\), the difference between the electron densities of two spin channels, shown in the inset in Fig. 3(c). A large amount of spins is localized at the N site (−1.21 \(\mu_B\)) with a slight contribution from the Cd ion (−0.08 \(\mu_B\)), indicating that the N ions mainly contribute to magnetism, also verified by Bader charge analysis.\(^{29}\) As is known, the Mermin–Wagner theorem\(^{61}\) reveals that the long-range magnetic ordering is absent by thermal fluctuations in 2D lattices. However, it doesn’t hold true for all 2D structures, especially for strong anisotropy cases.\(^{52,57}\) Thus, to determine the magnetic ground state of Cd\(_2\)N\(_3\), we consider several possible configurations, including, FM, Neel antiferromagnetic (NAFM), stripe AFM (SAFM), and zigzag AFM (ZAFM), as shown in Fig. 3(a). The calculated results indicate that the FM state (\(\Delta E_{\text{FM}} = E_{\text{FM}} - E_{\text{AFM}}\)) is energetically 14.43, 12.41, 9.75, and 9.32 meV N\(-1\) lower than the NAFM, SAFM, and ZAFM states, respectively. Furthermore, we also check the spin configuration of the coplanar 120° AFM state, and find that the Cd\(_2\)N\(_3\) monolayer prefers energetically FM (\(\Delta E_{\text{FM}} = 16.58\) meV N\(^{-1}\)) states.

Magnetocrystalline anisotropy energy (MAE), defined as the required energy to rotate the magnetization from an easy axis to a hard axis of a ferromagnet and determines the difficulty of spin flipping, plays a key role in establishing 2D FM ordering. For a Kagome lattice, MAE (\(\theta, \varphi\)) at an arbitrary direction (\(\theta, \varphi\)) follows

\[
\text{MAE}(\theta, \varphi) = K_1 \cos^2 \theta + K_2 \cos^4 \theta + K_3 \cos^4 3\varphi
\]

where \(K_1\) and \(K_2\) are system-dependent anisotropy constants and \(\varphi\) is the azimuthal angle of rotation. Fig. 3(b) shows MAE (\(\theta, \varphi\)) as a function of the azimuth angle \(\theta\) measured from the out-
of-plane axis and the horizontal azimuth angle \( \varphi \) on the in-plane axis. We find that the energy difference MAE \((\theta, \varphi)\) is almost independent of in-plane azimuthal angle \( \varphi \), i.e., \( K_3 = 0 \), and MAE \((\theta, \varphi)\) reaches its minimum in the out-of-plane axis, indicating an out-of-plane magnetization. According to the fits of angular dependence of MAE \((\theta, \varphi)\), both \( K_1 \) and \( K_2 \) are positive with a negligible dependence on \( \varphi \), and MAE \((\theta, \varphi)\) reaches a maximum value of 15 \( \mu \text{eV} \, \text{N}^{-1} \) at \( \theta_{xz} = \theta_{yz} = \pi/2 \), belonging to the family of 2D Ising magnets. Consequently, \( \text{Cd}_2\text{N}_3 \) is a 2D Ising magnet, clearly different from the cases of 1T- and 2H-VSe\(_2\) with the in-plane magnetization.\(^{33}\) The out-of-plane anisotropy of \( \text{Cd}_2\text{N}_3 \) could make the FM phase susceptible to thermal fluctuation and correspondingly feasible to realize a 2D magnet at a finite temperature.

The Curie temperature \((T_C)\) is a key parameter for the practical application of spintronic devices. Therefore, it is necessary to understand the behavior of magnetism with temperature before implementing \( \text{Cd}_2\text{N}_3 \) into practical spintronic devices. On the basis of the 2D Heisenberg model, the spin Hamiltonian can be expressed as

\[
H = -J_1 \sum_i S_i \cdot S_j - J_2 \sum_{i<k} S_i \cdot S_k - A S_z^2 \tag{3}
\]

where \( J_1 \) and \( J_2 \) are the nearest and next-nearest magnetic exchange interaction parameters, \( S_i \) is the spin vector of each atom, \( A \) is the anisotropy parameter, and \( S_z \) is the \( z \) component of the spin vector. A supercell of 100 \( \times \) 100 with periodic boundary conditions is used here. Fig. 3(c) shows the temperature-dependent magnetic moment per unit cell. The magnetic moment begins to drop dramatically at \( T_C = 225 \text{ K} \), implying the formation of the FM state. Additionally, heat capacity \((C_v)\) has been calculated, as shown in Fig. 3(c), further confirming the FM-PM phase transition temperature at critical temperature \( T_C \). The high \( T_C = 225 \text{ K} \) indicates that the \( \text{Cd}_2\text{N}_3 \) monolayer may be a promising high-temperature spintronic material.

Having determined the magnetic ground state of the \( \text{Cd}_2\text{N}_3 \) monolayer, we present the representative band structures without SOC, as shown in Fig. 4(a), where the blue and red colors display the spin-up and spin-down channels, respectively. One can see that both the spin-up and spin-down channels are completely split away from each other. Remarkably, the spin-up channel possesses a very large band gap 2.88 \( \text{eV} \), whereas the spin-down one did not show a gap. Thus the charge transport is dominated by the spin-down electron, and the current flow in such a system should be fully spin-polarized, which is consistent with the case in Fig. 1(a). The band character can be further confirmed by the HSE06 method, which has a larger band gap of 3.56 \( \text{eV} \) in the spin-up channel. Though the bulk gap is enhanced for HSE06, the HDS bands are not altered near the Fermi level, indicating the robustness of the band character against the computational method. As \( E_F \) lies almost in the middle of the spin gap, fully spin-polarized spin-filter efficiency can be maintained in a wide positive or negative bias range, rendering \( \text{Cd}_2\text{N}_3 \) an attractive candidate for spin-injection.

The most prominent findings on band dispersion are that in the spin-up channel, there are two several energy band crossings near \( E_F \), as shown Fig. 5(a), which mainly originates from the representation of the \( \text{D}_{3h} \) point group. From the 3D band profile near the \( K \) point, accompanied by symmetry analysis, we can obtain the six Dirac-like nodes located exactly at the Fermi level. Notably, this is a typical Kagome band around the Fermi level, consisting of a completely flat band above two Dirac bands remaining at the \( K \) point. From the two linear bands in Fig. 5(a), the Fermi velocity \( v_F \) of the carriers can be evaluated using linear fitting, \( v_F = \frac{\text{d}E}{\text{d}k} \). The obtained \( v_F = 4.3 \times 10^5 \text{ m s}^{-1} \) for \( \text{Cd}_2\text{N}_3 \) is comparable to that of graphene (8.2 \( \times \) 10\(^5\) \( \text{m s}^{-1} \)) and silicene (5.3 \( \times \) 10\(^5\) \( \text{m s}^{-1} \)).\(^{65}\) To understand the origin of Dirac states in \( \text{Cd}_2\text{N}_3 \), we calculate the orbital-resolved band structures, as shown in Fig. 6. One can see that the wave functions at the Fermi level are mainly from the N-2pz orbital with a slight contribution from Cd atoms, so there is no sp\(^3\) hybridization occurring on \( \text{Cd}_2\text{N}_3 \). The linear crossing Dirac states by the N-2pz orbital is consistent with the p\(_z\) Dirac state due to the nature of weak \( \pi \) bonds.\(^{66}\) Namely, \( \text{Cd}_2\text{N}_3 \) is a superior candidate of single-spin p-band Dirac materials for high-speed spintronic devices.
The peculiar band crossings that we found can host topologically nontrivial states in the presence of effective SOC. Fig. 3(b) shows the band structure with SOC, with a gap of \( E_g = 5.0 \text{ meV} \) at the Dirac-like node. The band gap is similar to the case of 2D Mn-DCA \((E_g = 2.3 \text{ meV})\) but smaller than recently predicted \( E_g = 20 \text{ meV} \) in Cs\(_2\)Mn\(_3\)F\(_{12}\). It is noteworthy that, due to the small atomic numbers, SOC has negligible effects on the band structure of Cd\(_2\)N\(_3\). Such band gap opening at the K point suggests a topologically nontrivial feature, which can be confirmed by the Berry curvature \( \Omega(k) \) from the Kubo formula\(^{47-49} \) expressed as

\[
\Omega(k) = \sum_n f_n \Omega_n(k)
\]

where the summation is over all of the occupied states, \( E_n \) is the eigenvalue of the Bloch function \( |\Psi_{nk}\rangle \), \( f_n \) is the Fermi-Dirac distribution function, and \( v_x \) and \( v_y \) are the velocity operators. Fig. 5(b) and (c) show the Berry curvature in the BZ around the K and \( K' \) points, which are nonzero with the same sign. By integrating the Berry curvatures over the first BZ, we calculated the Chern number \( C \), expressed as

\[
C = \frac{1}{2\pi} \sum_n \int_{BZ} d^2 k \Omega_n.
\]

One can see that \( C = 1 \) with each Dirac cone contributing 0.5, characterizing the topological nature of the gaped state. So the anomalous Hall conductivity \( \sigma_{xy} = \sigma_y^z \) shows a quantized charge Hall plateau at a value of \( e^2/h \) located in the insulating gap of the spin-up Dirac-like node. Such a nonvanishing Chern number \( C = 1 \) and quantized Hall conductivity in Cd\(_2\)N\(_3\) characterize the QAH effect.

The two band degeneracy point and the associated band topology can be well modeled by the TB model, taking the \( D_{2h} \) lattice symmetry constraint into account. We introduce a Slater-Koster type TB model\(^{51} \) with the p orbital on the Kagome lattice formed by N atoms. The energy bands near the Dirac points can be expressed by the effective Hamiltonian

\[
H = \hbar v_F (k_x \sigma_x + k_y \sigma_y) + m_1 \sigma_z \tau_z,
\]

where \( v_F \) is the Fermi velocity, \( \sigma \) are the Pauli matrices, \( \tau = \pm 1 \) refers to inequivalent N sublattices, and \( \tau_z = \pm 1 \) refers to the valley pseudo-spin of the Cd\(_2\)N\(_3\) monolayer. Wave vector \( k_{xy} \) is referenced from \( k \) or \( k' \).

SOC connects the electron spin and orbital degrees of freedom. So the crystal potential in 2D materials can be approximated by the spherical atomic potential, which gives on-site contribution to the TB Hamiltonian. By averaging the radial degree of freedom, the SOC reads

\[
H_{eff}^{SO} = \xi \mathcal{L} \cdot \mathcal{L}',
\]

where \( \mathcal{L} \) is the angular momentum operator. The matrix element \( \xi \mathcal{L} \cdot \mathcal{L}' \) is given on the basis of directed atomic orbitals \( \mu, \nu \) and \( \ell \) is the angular momentum resolved atomic SOC strength with \( \ell = \{s, p, d, \ldots \} \). Obviously, the energy bands are gapped by the effective SOC, which requires nonzero \( s_z \) due to its out-of-plane magnetic moment in the Cd\(_2\)N\(_3\) monolayer. Thus, our TB analysis provides a systematic understanding of the existing first-principles results, and furthermore, a useful guide for further material exploration.

The derived Haldane effective model describes the QAH effect,\(^{70} \) and the existence of topologically protected chiral edge states is one of the most important signatures of the QAH effect. To reveal the nontrivial topological nature of the Cd\(_2\)N\(_3\) lattice, we construct Green’s functions\(^{71} \) for the semi-infinite boundary based on the maximally localized Wannier function method.\(^{72,73} \) Fig. 4(d) shows the local density of states (LDOSs) of the edge states. Obviously, the nontrivial edge states connecting the CB and VB cross the insulating gap of the spin-up Dirac cone, consistent with the Chern number \( C = 1 \). The spin-polarized Dirac-fermion mediated topological character suggests that the Cd\(_2\)N\(_3\) monolayer is intrinsic, holding the potential QAH phase.

As is well known, the physical properties of 2D materials can be effectively tuned by strain, like the band gap engineering of 2H\(^2\)TMD semiconductors.\(^{74} \) Strain has also been proposed to effectively modulate dielectric properties,\(^{75} \) SOC,\(^{76} \) thermal conductivity\(^{77} \) and interlayer coupling in vdW heterostructures (HTs)\(^{78} \) in 2D lattices. In the following, we will show that strain can induce SOC induced bulk gaps and spin-exchange interactions. Fig. 7 presents the evolution of direct gap \( E_g \) and FM exchange energy as a function of strain \( \varepsilon \), defined as \( \varepsilon = (a - a_0)/a_0 \), where \( a_0 \) is the strained (equilibrium) lattice constants. Note that the nontrivial topological states remain within the strain range of 10%. This suggests that Cd\(_2\)N\(_3\) maintains a topologically nontrivial state, which is stable against external strains. The gap \( E_g \) decreases slightly with tensile strains, becoming \( E_g = 4.95 \text{ meV} \) under strain 10%. Additionally, the variation of total energies of the FM and AFM states (\( \Delta E_{FM} = \))
Spin relaxation times due to small SOC, to less changes of strains, for example, when ε = 10%, ΔE_{FM} = 14.43 meV, leading to less changes of T_C. The dependency of T_C on strain is clearly different from the cases of 2D FeCl_2 (ref. 30) and CdSe_2,31 where T_C changes sensitively with respect to external strains. So we can infer that Cd_2N_3 may be a promising intrinsic high-temperature spintronic material. Here we must note that previously proposed strategies to realize p-band FM ordering may be hard to be experimentally controllable, because a strong external electric field or carefully selective doping is required. Our work thus reports the first intrinsic p-band HDS material, even though the experimental synthesis of Cd_2N_3 may still remain a challenge. As these magnetic materials potentially offer large spin relaxation times due to small SOC,79 our results highlight a new promising material for experimental validation studies toward the realistic p-band spintronics applications.

Finally, one critical point is whether the QAH effect of 2D Cd_2N_3 can remain on a substrate, since the substrates are inevitable in device applications.80,81 As is known, the BN monolayer is chemically inert and does not easily bond strongly with other atoms, and thus may be adopted as a protective film to grow Cd_2N_3. So, the Cd_2N_3/BN HTS has been constructed, as shown in Fig. 8(a). Structural optimization indicates that the distance of Cd_2N_3 and BN layers is 3.10 Å with a binding energy of −34 meV per unit cell, a typical van der Waals structure. In this case, the main features of the QAH effect in the Cd_2N_3 lattice remain intact. Fig. 8(b) presents the calculated band structure with SOC. As expected, here there is still a SOC-induced gap at the Dirac point around the Fermi level, and the states around the Fermi level are dominantly contributed by the Kagome Cd_2N_3 band. Given that the 2D BN substrate electrically insulates the adjacent QSH layer of Cd_2N_3, protecting parallel helical edge channels from being gapped by interlayer hybridization, Cd_2N_3/BN is an ideal HTS to support the dissipationless charge/spin transport in the quantum device. This demonstrates the feasibility of constructing the quantum device with the Cd_2N_3/BN heterostructure, as illustrated in Fig. 8(c).

### 4. Conclusion

In summary, we employ first-principles calculations to demonstrate the possibility of realizing HDSs in Kagome lattices with T_C = 225 K, which is much higher than that of the reported monolayers CrI_3 (T_C = 45 K) and Cr_2Ge_2Te_6 (T_C = 20 K). The inclusion of SOC can turn Cd_2N_3 into a QAH insulator, with a sizable band gap of E_g = 5.0 meV. Notably, the nontrivial topology is robust against biaxial strain with its band gap reaching up to E_g = 4.95 meV under strain 10%. Also, the TB model is constructed to clarify the origin of the HDS and nontrivial electronic properties. Considering that Cd_2N_3 enjoys good stability and excellent flexibility, our findings would stimulate further material exploration toward the HDS matter in Kagome materials. These outstanding properties of Cd_2N_3 indicate that it is a promising 2D material for potential application in designing energy efficient spintronic devices.

Additionally, the proposed Cd_2N_3 is likely to be synthesized by molecular beam epitaxy (MBE) or by the chemical vapor deposition method. We can synthesize Cd_2N_3 with a synthesis method similar to that of Nb_2O_3.82 The QAH effect can be probed by scanning probe spectroscopy, ARPES, and ultimately by device fabrications and charge transport measurements. Thus, it offers a new platform for the study of the QAH effect in Kagome lattices for realizing low-dissipation spintronic devices, which is of significant fundamental importance.

### Conflicts of interest

There are no conflicts to declare.

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References

1 S. A. Wolf, D. D. Awschalom, R. A. Buhrman, J. M. Daughton, S. von Coudnár, M. L. Roukes, A. Y. Chtchelkanova and D. M. Treger, Science, 2001, 294, 1488.
2 C. W. Zhang and S. S. Yan, Appl. Phys. Lett., 2009, 95, 232108.
3 R. A. de Groot, F. M. Mueller, P. G. van Engen and K. H. J. Buschow, Phys. Rev. Lett., 1983, 50, 2024.
4 H. Van Leuken and R. A. De Groot, Phys. Rev. Lett., 1995, 74, 1171.
5 K. W. Lee and C. E. Lee, Adv. Mater., 2012, 24, 2019–2023.
6 D. Hsieh, D. Qian, L. Wray, et al., Nature, 2008, 452, 970.
7 Y. Xia, D. Qian, D. Hsieh, et al., Nat. Phys., 2009, 5, 398.
8 Y. S. Hor, A. Richardson, P. Roushan, et al., Phys. Rev. B, 2009, 79, 195208.
9 D. Hsieh, Y. Xia, D. Qian, et al., Phys. Rev. Lett., 2009, 103, 146401.
10 H. Ohno, A. Shen, F. Matsukura, et al., Appl. Phys. Lett., 1996, 69, 363–365.
11 H. Kimura, T. Fukumura, M. Kawasaki, et al., Appl. Phys. Lett., 2002, 80, 94–96.
12 A. Ramasubramaniam and D. Naveh, Phys. Rev. B, 2013, 87, 195201.
13 K. Mullen, B. Uchoa and D. T. Glatzhofer, Phys. Rev. Lett., 2015, 115, 026403.
14 C. Fang, Y. Chen, H.-Y. Kee and L. Fu, Phys. Rev. B, 2015, 92, 081201(R).
15 R. W. Zhang, Z. Zhang, C. C. Liu, et al., Phys. Rev. Lett., 2020, 124, 016402.
16 B. Huang, G. Cheng, E. N. Moratalla, D. R. Klein, R. Cheng, K. L. Seyler, D. Zhong, E. Schmidgall, M. A. McGuire and D. H. Cobden, Nature, 2017, 546, 270.
17 C. Gong, L. Li, Z. Li, H. Ji, A. Stern, Y. Xia, T. Cao, W. Bao, C. Wang and Y. Wang, Nature, 2017, 546, 265.
18 X. L. Wang, Phys. Rev. Lett., 2008, 100, 156404.
19 G. D. Liu, X. F. Dai, H. Y. Liu, et al., Phys. Rev. B, 2007, 77, 014424.
20 S. Skafetosouros, K. Ozdogan, et al., Appl. Phys. Lett., 2013, 102, 022402.
21 G. Y. Gao and K. L. Yao, Appl. Phys. Lett., 2013, 103, 232409.
22 Z. F. Wang, S. Jin and F. Liu, Phys. Rev. Lett., 2013, 111, 096803.
23 Y. Pan and Z. Yang, Phys. Rev. B, 2010, 82, 195308.
24 L. Bainsla, A. I. Cdallick, M. M. Raja, et al., Phys. Rev. B, 2015, 92, 045201.
25 F. Zheng, C. Zhang, P. Wang, et al., J. Appl. Phys., 2013, 113, 154302.
26 G. Gao, G. Ding, J. Li, et al., Nanoscale, 2016, 8, 8986–8994.
27 C. L. Kane and E. J. Mele, Phys. Rev. Lett., 2005, 95, 226801.
28 X. Zhou, R. W. Zhang, Z. Zhang, et al., J. Phys. Chem. Lett., 2019, 10(11), 3101.
29 A. H. Castro Neto, F. Guinea, N. M. R. Peres, K. S. Novoselov and A. K. Geim, Rev. Mod. Phys., 2009, 81, 109–162.
30 Y. P. Bliokh, V. Freilikhver and F. Nori, Phys. Rev. B, 2013, 87, 2451.
31 K. S. Novoselov, A. K. Geim, S. V. Croll, D. Jiang, M. L. Katsnelson, I. V. Grigorieva, S. V. Dubonos and A. A. Firsov, Nature, 2005, 438, 197–200.
32 Y. Zhang, Y. W. Tan, H. L. Stormer and P. Kim, Nature, 2005, 438, 197–204.
33 S. Cahangirov, M. Topsakal, E. Aktürk, H. Şahin and S. Ciraci, Phys. Rev. Lett., 2009, 102, 236804.
34 J. Chen, J. Xi, D. Wang and Z. Shuai, J. Phys. Chem. Lett., 2013, 4, 1443.
35 C. L. Xu, R. Z. Wang, M. S. Miao, X. L. Wei, Y. P. Chen, H. Yan, W. M. Lau, L. M. Liu and Y. M. Ma, Nanoscale, 2014, 6, 1113–1118.
36 Z. Wang, X. F. Zhou, X. Zhang, Q. Zhu, H. Dong, M. Zhao and A. R. Oganov, Nano Lett., 2015, 15, 6182–6186.
37 H. Zhang, Y. Xie, Z. Zhang, C. Zhong, Y. Li, Z. Chen and Y. Chen, J. Phys. Chem. Lett., 2017, 8, 1707.
38 X. Qin, Y. Wu, Y. Liu, B. Chi, X. Li, Y. Wang and X. Zhao, Sci. Rep., 2017, 7, 10546.
39 C. Pu, D. Zhou, Y. Li, H. Liu, Z. Chen, Y. Wang and Y. Ma, J. Phys. Chem. C, 2017, 121, 2669.
40 S. J. Zhang, C. W. Zhang, S. F. Zhang, W. X. Ji, P. Li, P. J. Wang, S. S. Li and S. S. Yan, Phys. Rev. B, 2017, 96, 205433.
41 P. Zhou, Z. S. Mab and L. Z. Sun, J. Mater. Chem. C, 2018, 6, 1206–1214.
42 Y. P. Wang, W. X. Ji, C. W. Zhang, P. Li, P. J. Wang, B. Kong, S. S. Li, Z. S. Yan and K. Liang, Appl. Phys. Lett., 2017, 110, 233107.
43 M. H. Zhang, X. L. Chen, W. X. Ji, P. J. Wang and C. W. Wang, Appl. Phys. Lett., 2020, 116, 172105.
44 Y. Y. Deng, Y. Song, Y. Zhang, J. Wang, N. Z. Sun, Z. Yi, Y. Wu, Y. Z. Wu, S. Zhu, J. Wang, J. Chen and X. H. Zhang, Nature, 2018, 563, 94–99.
45 C. Gong, L. Li, Z. Li, H. Ji, A. Stern, Y. Xia, T. Cao, W. Bao, C. Wang, Y. Wang, Z. Qiu, R. J. Cava, S. G. Louie, J. Xia and X. Zhang, Nature, 2017, 546, 265.
46 B. Huang, G. Clark, E. N. Moratalla, D. R. Klein, R. Cheng, K. L. Seyler, D. Zhong, E. Schmidgall, M. A. McGuire, D. H. Cobden, W. Yao, D. Xiao, P. J. Herrero and X. Xu, Nature, 2017, 546, 270.
47 M. H. Zhang, C. W. Zhang, P. J. Wang and S. S. Li, Nanoscale, 2018, 10, 20226.
48 X. J. Ni, W. Jiang, H. Q. Huang, K. H. Jin and F. Liu, Nanoscale, 2018, 10, 11901.
49 L. Zhang, S. F. Zhang, W. X. Ji, C. W. Zhang, P. Li, P. J. Wang, S. S. Li and S. S. Yan, Nanoscale, 2018, 10, 20748.
50 G. Kresse and J. Furthmüller, Phys. Rev. B, 1996, 54, 134179.
51 P. E. Blöchl, Projector augmented-wave method, Phys. Rev. B, 1994, 50, 17953.
52 G. Kresse and D. Joubert, Phys. Rev. B, 1999, 59, 13169.
53 J. P. Perdew, K. Burke and M. Ernzerhof, Phys. Rev. Lett., 1996, 77, 3865.
54 J. P. Perdew, K. Burke and M. Ernzerhof, Phys. Rev. Lett., 1996, 77, 3865.
55 A. Togo and I. Tanaka, Scr. Mater., 2015, 108, 1.
56 Z. F. Wang, Z. Liu and F. Liu, Nat. Commun., 2013, 4, 1471.
57 X. Zhang, Z. Wang, M. Zhao and F. Liu, Phys. Rev. B, 2016, 93, 165401.
58 R. Addou, A. Dahal and M. Batzill, Nat. Nanotechnol., 2013, 8, 41.
59 B. Huang, G. Clark, E. Navarro-Moratalla, D. R. Klein, R. Cheng, K. L. Seyler, D. Zhong, E. Schmidgall, M. A. Mcguire and D. H. Cobden, Nature, 2017, 546, 270.
60 C. Gong, L. Li, Z. Li, H. Ji, A. Stern, Y. Xia, T. Cao, W. Bao, C. Wang and Y. Wang, Nature, 2017, 546, 265.
61 G. Henkelman, A. Arnaldsson and H. Jonsson, Comput. Mater. Sci., 2016, 36, 354.
62 N. D. Mermin and H. Wangner, Phys. Rev. Lett., 1966, 17, 1133.
63 M. Esters, R. G. Hennig and D. C. johnson, Phys. Rev. B, 2017, 96, 235147.
64 Y. Jiao, F. Cda, J. Bell, A. Bilic and A. Du, Angew. Chem., 2016, 128, 10448–10451.
65 Y. L. Wang and Y. Ding, Solid State Commun., 2013, 155, 6–11.
66 Z. X. Guo, S. Furuya, J. Iwata and A. Oshiyama, Phys. Rev. B, 2013, 87, 235435.
67 Y. G. Yao, L. Kleinman, A. H. MacDonald, J. Sinova, T. Jungwirth, D. S. Wang, E. Wang and Q. Niu, Phys. Rev. Lett., 2004, 92, 037204.
68 Y. G. Yao and Z. Fang, Phys. Rev. Lett., 2005, 95, 156601.
69 J. C. Slater and G. F. Koster, Phys. Rev., 1954, 94, 1498.
70 S. Konschuh, M. Gmitra and J. Fabian, Phys. Rev. B, 2010, 82, 245412.
71 J. L. Lado and J. Fernández-Rossier, 2D Mater., 2017, 4, 035002.
72 C. L. Kane and E. J. Mele, Phys. Rev. Lett., 2005, 95, 226801.
73 C. C. Liu, W. X. Feng and Y. G. Yao, Phys. Rev. Lett., 2011, 107, 076802.
74 P. Johari and V. B. Shenoy, ACS Nano, 2012, 6, 5449.
75 A. Kumar and P. K. Ahluwalia, Phys. B, 2013, 419, 66.
76 P. Koskinen, I. Fampiou and A. Ramasubramaniam, Phys. Rev. Lett., 2014, 112, 186802.
77 Z. Ding, Q.-X. Pei, J.-W. Jiang and Y.-W. Zhang, J. Phys. Chem. C, 2015, 119, 16358.
78 M. Sharma, A. Kumar, P. K. Ahluwalia and R. Pandey, J. Appl. Phys., 2014, 116, 063711.
79 S. Sanvito, Molecular spintronics, Chem. Soc. Rev., 2010, 40, 3336.
80 S. S. Li, W. X. Ji, S. J. Hu, C. W. Zhang and S. S. Yan, ACS Appl. Mater. Interfaces, 2017, 9, 41443.
81 M. Zhou, Z. Liu, W. M. Ming, Z. F. Wang and F. Liu, Phys. Rev. Lett., 2014, 113, 236802.
82 S. Q. Wang, C. Noguera and M. R. Castell, Phys. Rev. B, 2019, 100, 125408.