Non-coplanar spin textures with finite scalar spin chirality can be artificially induced at surfaces and interfaces through the interfacial Dzyaloshinskii-Moriya interaction. However, stabilizing a proper magnetic skyrmion crystal via this route remains elusive. Here, using an epitaxial bilayer of platinum and geometrically frustrated kagome-lattice ferromagnet Fe$_3$Sn, we show the possible formation of a two-dimensional skyrmion crystal under well-regulated Fe$_3$Sn thickness conditions. Magnetization measurements reveal that the magnetic anisotropy is systematically varied from an inherent in-plane type to a perpendicular type with the thickness reduction. Below approximately 0.5 nm, we clearly detect a topological Hall effect that provides evidence for finite scalar spin chirality. Our topological Hall effect analysis, combined with theoretical simulations, not only establishes its interfacial Dzyaloshinskii-Moriya interaction origin, but also indicates the emergence of a stable skyrmion crystal phase, demonstrating the potential of kagome-lattice ferromagnets in spin chirality engineering using thin-film nanostructures.
Scalar spin chirality, $\mathbf{S}_1 \cdot (\mathbf{S}_2 \times \mathbf{S}_3)$, is a quantity that corresponds to a solid angle subtended by three spins $\mathbf{S}_1$, $\mathbf{S}_2$, and $\mathbf{S}_3$ on neighboring triangular sites\textsuperscript{1-3}. Electrons hopping through non-coplanar spin textures with finite scalar spin chirality acquire Berry phase equivalent to a fictitious magnetic field, leading to topological Hall effect (THE)\textsuperscript{4,5} that is distinct from other Hall effects. Among various non-coplanar spin textures, a swirling spin texture called the magnetic skyrmion\textsuperscript{4-6} is particularly interesting because it can not only behave like an individual particle-like object but also can form an ordered skyrmion crystal (SKX) under the delicate competition of magnetic interactions and external perturbations. Understanding the stability and controllability of SKX phase is currently at the frontier of condensed matter research.

The established approach for the formation of SKX phase is to use noncentrosymmetric bulk crystals\textsuperscript{4,5} in which finite contributions of spin-orbit coupling (SOC) induce the Dzyaloshinskii-Moriya interaction (DMI)\textsuperscript{7,8}, expressed as $H_{\text{DMI}} = - D \cdot (\mathbf{S}_i \times \mathbf{S}_j)$ with $D$ being the DMI vector. This interaction favors a relative twist between otherwise parallel or anti-parallel spins, giving rise to complex magnetic states such as canted ferromagnetism\textsuperscript{7,8}, non-coplanar and helical/spiral spin states\textsuperscript{4,5}. In thin-film heterostructures, the role of SOC in bulk crystals can be activated by fabricating an asymmetric interface with individual magnetic and SOC layers, which is known as the interfacial DMI\textsuperscript{9,10}. In various interface systems including metals\textsuperscript{11-16}, perovskite-type oxides\textsuperscript{17,18}, and topological insulators\textsuperscript{19,20}, the formation of isolated skyrmions and the detection of THE have been reported. However, the two-dimensional SKX has rarely been achieved by the interfacial DMI mechanism\textsuperscript{20}. Given the crucial role of crystal symmetry in the bulk approaches, a choice of specific lattice systems that enable the tuning of complex magnetic interactions is essential.

With this perspective, we focused on kagome-lattice magnets with a triangular-based lattice, which is suitable for inducing non-coplanar magnetic states required for the SKX phase. We selected a ferromagnetic intermetallic compound Fe$_3$Sn with the $D_{3h}$ structure. Bulk Fe$_3$Sn is ferromagnetic below Curie temperature $T_C$ of 743 K with in-plane magnetic anisotropy in the Fe$_3$Sn kagome plane ($ab$ plane)\textsuperscript{21}. Very recently, it has been reported that Fe$_3$Sn can be epitaxially grown as a thin film on Pt(111) (ref.\textsuperscript{22}). The combination of Fe$_3$Sn and Pt offers the following advantages for spin chirality insulators\textsuperscript{18,19}, the formation of isolated skyrmions and the detection of THE revealed by the systematic measurements are consistent with the formation of SkX in a theoretical model based on a ferromagnetic kagome lattice. Thickness-dependent magnetic anisotropy in Fe$_3$Sn/Pt bilayers. By comparing $M$ versus $\mu_0H$ curves in out-of-plane and in-plane configurations, we detected the expected variation of magnetic anisotropy, as displayed in Fig. 2a-c. For $t = 0.64$ nm (Fig. 2a), the $M$ saturates easily under the application of in-plane $H$ (blue curve), showing the in-plane magnetic anisotropy consistent with the bulk behavior\textsuperscript{21}. In the sample with $t = 0.48$ nm (Fig. 2b), however, the $M$ saturates at almost comparable $\mu_0H$ values in the two measurement configurations. The even thinner condition of $t = 0.24$ nm appears to favor perpendicular magnetic anisotropy (Fig. 2c). From the $M/M_{\text{sat}}$ versus $\mu_0H$ curves, we calculated the magnetic anisotropy energy as the effective anisotropy field $\mu_0H_{\text{eff}}$ in the unit of $T$ ($\mu_0H_{\text{eff}} > 0$ for perpendicular magnetic anisotropy and $< 0$ for in-plane magnetic anisotropy). As plotted in Fig. 2d, $\mu_0H_{\text{eff}}$ gradually varies from negative to positive with decreasing $t$, changing its sign at approximately 0.5 nm (~1 unit cell). This result captures that the inherent in-plane magnetic anisotropy turns to perpendicular magnetic anisotropy at a crossover thickness of $t_{\text{ana}} \sim 0.5$ nm. To gain insight into spin textures in these ultrathin Fe$_3$Sn bilayer samples, we performed a numerical simulation for a single kagome plane that took into account the ferromagnetic Heisenberg interaction $J$, magnetic anisotropy $A$ (in-plane and perpendicular magnetic anisotropy for $A < 0$ and $> 0$, respectively), and DMI (see Methods). As schematically shown in Fig. 2e, this model considers spin interactions between the neighboring spins on the kagome plane in an out-of-plane magnetic field. Bulk and interfacial DMI contribute to the out-of-plane and in-plane components of $\mathbf{D}$ ($D_x$ and $D_y$), respectively, which modifies the local spin interactions to produce a finite spin solid angle. Figure 2f displays the total spin solid angle per the unit cell of Fe$_3$Sn, $|\Omega_{\text{SKX}}|$, as a function of $A$ and $D_y$ ($J = D_x = 1$ and out-of-plane magnetic field $g\mu_0H = 0.5$, where $g$ is the $g$-factor and $\mu_0$ is the Bohr magnetron). In the upper-left region of large negative $A$ and small $D_y$, the spins basically lie in the kagome plane
Fig. 1 Kagome-lattice ferromagnet Fe₃Sn/Pt bilayer. a Schematic structure of Fe₃Sn(t nm)/Pt(10 nm) bilayer on Al₂O₃(0001). An SiOₓ cap layer is not drawn for simplicity. b, c Concept of the scalar spin chirality driven by the interfacial DMI. Above t_Ma, bulk-like in-plane magnetic anisotropy stabilizes a coplanar spin state with no scalar spin chirality. Below t_Ma, the interfacial DMI modifies the magnetic anisotropy, inducing a non-coplanar spin state with finite scalar spin chirality. d Cross-sectional TEM image of a bilayer sample with t = 8 nm viewed along Al₂O₃[1120]. Scale bar, 5 nm. e Selected-area electron diffraction pattern including the film and substrate regions. Scale bar, 5 nm⁻¹. f M versus μH curves for t = 0.8 nm measured at T = 400, 350, 300, and 250 K in an out-of-plane H configuration. The solid and dashed curves correspond to the decreasing-field and increasing-field scans, as shown by the black solid and dashed arrows, respectively. g t dependence of M_sat at T = 400 K in an out-of-plane H configuration. The M_sat is averaged over μH = 4–7 T. The bulk M_sat is calculated from the mass magnetization data in ref. 23, using the reported density of 8.45 g cm⁻³ (JCPDS PDF 01-074-5857).

Verifying scalar spin chirality via the THE analysis. The arrow indicates the direction of negative Ω_tot, we found that a moderate contribution of interfacial DMI would be canceled by opposite contributions from the top Pt/Fe₃Sn and bottom Fe₃Sn/Pt interfaces. Although there is no direct correspondence between the actual samples and simulation conditions, the thick-bilayer, thin-bilayer, and thin-trilayer structures can be compared with the simulation conditions of negative A and small D⊥, small positive A and large D⊥, and small positive A and small D⊥, respectively. Note here that a large part of conduction in these heterostructures is governed by the highly conducting Pt layer (see Supplementary Figs. 9 and 10 for the t-dependence of sheet resistance and the magnetoresistance data, respectively). Figure 3d shows Hall resistance Rₓᵧ versus μH curves of the thick-bilayer structure at T = 400 K. In addition to the linear ordinary Hall effect of Pt (Supplementary Fig. 11k), a nonlinear response in the Rₓᵧ is clearly discernable. A fit using an empirical relation, Rₓᵧ = Rₓᵧ H + Rₓᵧ M + Rₓᵧ xᵧ, where Rₓᵧ and Rₓᵧ are the ordinary and anomalous Hall coefficients in the unit of Ω T⁻¹, and Rₓᵧ M (= Rₓᵧ xᵧ) and Rₓᵧ are the anomalous and topological Hall resistances, reveals the dominant contribution of Rₓᵧ xᵧ to the Rₓᵧ (Fig. 3e inset), which is attributed to the anomalous Hall effect (AHE) of Fe₃Sn. The negligibly small residual Rₓᵧ xᵧ (Fig. 3e) indicates the absence of THE in the thick-bilayer structure. This is consistent with the bulk-like in-plane magnetic anisotropy favors the coplanar spin state without scalar spin chirality (Fig. 1b). Contrastingly, in the thin-bilayer structure, the extracted Rₓᵧ (Fig. 3f) at T = 400 K ~T_c (Supplementary Fig. 5) overwhelms both ordinary and anomalous Hall resistances (Fig. 3f and the inset of Fig. 3g), indicating the presence of finite scalar spin chirality that contributes to THE. The addition of top Pt layer in the thin-trilayer structure, as intended, completely diminishes the THE (Fig. 3h). On the other hand, the comparable AHE for the thin-trilayer and thin-bilayer structures at T = 300 K (the insets of Fig. 3f, h; also see Supplementary Fig. 11 for the data at various T) indicates their macroscopically similar ferromagnetic states. In recent studies on SrRuO₃ ultrathin films...
and SrRuO₃-based perovskite multilayers and superlattices, the impact of inhomogeneity on the occurrence of THE-like $R_{yx}$ anomalies has been argued. When AHE changes its sign depending on thickness as in the present system (Supplementary Fig. 11), local thickness fluctuation could give rise to hump-like $R_{yx}$ behavior via the superposition of AHE components with different signs. To understand the origin of the observed $R_{yx}^T$, we evaluated asymmetric Ta/Fe₃Sn/Pt and W/Fe₃Sn/Pt trilayer structures (Supplementary Fig. 12 and Supplementary Note 1). The results were consistently explained by considering different magnitudes of the interfacial DMI contributions from the top and bottom interfaces. In conjunction with the insignificant thickness fluctuation suggested by the TEM and magnetization measurements, this strongly supports that the interfacial DMI plays a more decisive role in the occurrence of $R_{yx}^T$ than inhomogeneity. Furthermore, the superposition of $R_{yx}$ versus $\mu H$ curves for different $t$ values at $T = 400$ K (Supplementary Fig. 11) cannot reproduce the sharp $R_{yx}$ peaks detected in the thin-bilayer structure (Fig. 4f). These observations are fully consistent with the generation of finite scalar spin chirality owing to the local modification of spin interactions by the interfacial DMI.

**Analysis of THE and possible formation of SkX phase.** Having verified the interfacial DMI origin, we demonstrate the $t$-controlled variation of scalar spin chirality via the detection of THE. As shown in Fig. 4a–c, a slight reduction of $t$ effectively lowers the $T$ range where THE appears (see Supplementary Fig. 13 for the corresponding $R_{yx}^T$ data). Upon comparing the $R_{yx}$ data including those of other $t$ values (Supplementary Fig. 11), we noticed the sign reversal of $R_{yx}$ at approximately $T = 370$ K for $t = 0.40$ nm (Fig. 4b) and at $T = 340$ K for $t = 0.32$ nm (Fig. 4c). Figure 4d shows a contour plot of $\Delta R_{yx}(\pm 3 T) = R_{yx}(+3 T) - R_{yx}(-3 T)$ on the $t$-$T$ plane; its sign reversal occurs in close proximity to the THE region (surrounded by dashed lines). According to recent band structure calculation, Fe₃Sn is classified as a magnetic Weyl semimetal with Weyl nodes near the Fermi level. The $T$-induced shift of the Fermi level and the resulting change in the intrinsic AHE contribution may play a role in the $T$-induced sign reversal of AHE.

More importantly, as shown in Fig. 4e, the magnitude of $R_{yx}^T$ increases with increasing $T$ (up to our measurement limit of 400 K), concomitantly with the increase of the peak magnetic field at which $R_{yx}$ shows local maxima/minima due to THE, $\mu H_{\text{peak}}$ (black circles).
Using the standard linear-response theory for a Kondo-lattice model, we performed the analysis of finite-temperature Hall conductivity \(\sigma_{xy}\) of the intrinsic Berry phase mechanism (Fig. 4f, also see the \(\Omega_{\text{ext}}\) result for Supplementary Fig. 14). Overall, the simulation results reproduce the observed experimental trend, indicating the contribution of thermal fluctuation\(^{18,31}\) to the intrinsic (topological) Hall effect. These excellent agreements between the experiments and simulation suggest the formation of SkX phase in the THF region. To support this, we attempted to extract real-space features from the THF data using the relation\(^{5,16,17}\): 
\[
\rho_{ij}^T = \frac{P \rho_{ij} \phi_i \phi_j}{\Omega_{\text{ext}}},
\]
where \(\rho_{ij}^T\) is the topological Hall resistivity, \(P\) is the spin polarization, \(\rho_{ij}\) is the ordinary Hall coefficient, \(\phi_i\) is the skyrmion density, and \(\phi_j\) is one magnetic flux quantum \(= \hbar/e\) with \(e\) being the Planck constant and \(\hbar\) being the elementary charge. The \(n_{sk}^{0.5}\) corresponds to the average separation of skyrmions. Assuming a parallel circuit consisting of two conducting layers, we calculated the resistivity and Hall resistivity (Supplementary Fig. 15) of the Fe\(_3\)Sn\(_2\) layer from the 0.40-nm-thick bilayer data and reference Pt monolayer data at \(T = 350 K\), yielding \(\rho_{sk}^{0.5} = 37.4 \ \Omega \ \text{cm}\) and \(R_0 = 8.48 \times 10^{-5} \ \text{cm}^3 \ \text{C}^{-1}\). These values give \(n_{sk}^{0.5} = 9.68 \ \text{nm}\) and \(30.6 \ \text{nm}\) for \(P = 0.1\) and 1, respectively, which are reasonable as compared with the size of skyrmions reported for other bilayer systems (Supplementary Table 1). We therefore think that densely arranged skyrmions like SkX exist in the Fe\(_3\)Sn\(_2\)/Pt bilayer samples. Direct observation of the spin textures using microscopy could be an interesting future study.

**Conclusion**

The epitaxial interface of kagome-lattice ferromagnet Fe\(_3\)Sn\(_2\) and Pt enables the rational control of magnetic and electrical properties based on the interfacial DMI. Considering the rich variety of kagome-lattice magnets such as Fe\(_3\)Sn\(_2\), FeSn, and Co\(_3\)Sn\(_2\), that have been discovered from the aspect of topological physics, the development of heterointerfaces and superlattices is worthy of investigation. These fascinating features of kagome-lattice magnets will offer tremendous opportunities for exploring new functionalities of SkX-based phenomena.

**Methods**

**Thin-film growth.** The films were fabricated on Al\(_2\)O\(_3\) substrates by radio-frequency magnetron sputtering at an Ar gas pressure of 0.5 Pa. The Pt, Fe\(_3\)Sn, and SiO\(_2\) layers were in situ deposited at 600, 400, and 100°C with Pt, Fe\(_3\)Sn [ref. 12], and SiO\(_2\) targets, respectively. The 2-nm-thick Pt top layer of the thin-bilayer structure (Fig. 3c) was deposited at 100°C before the SiO\(_2\) capping. The crystal structure of the films was characterized by TEM and X-ray diffraction using Cu K\(_\alpha\) radiation. The \(t\) values were calculated based on the sputtering rate that was calibrated with the cross-sectional TEM image shown in Fig. 1d.

**Measurements.** The magnetization was measured with a superconducting quantum interference device magnetometer (MPMS3, Quantum Design) upon decreasing \(\mu_0H\) from 7 T to −7 T and increasing \(\mu_0H\) from −7 T to 7 T. By subtracting a diamagnetic contribution from Al\(_2\)O\(_3\) substrate, which was estimated from the high-field data at \(\mu_0H = 4\) T, the \(M\) was calculated. By antisymmetrizing the decreasing-field and increasing-field \(M\) data, the two antisymmetric \(M\) curves shown in Figs. 1f and 2a and Supplementary Figs. 4 and 8 were obtained. The electrical properties were measured with a physical property measurement system (PPMS, Quantum Design). The films were patterned into a Hall-bar shape by mechanical scratch, and electrical contacts were made with indium solder. For the analysis of the Hall effect, the decreasing-field and increasing-field \(R_{xy}\) data were anti-symmetrized against \(\mu_0H\) to eliminate spurious contributions arising from thermoelastic effect and misalignment of the Hall voltage probes.

**Simulation.** We considered a classical Heisenberg model on the kagome lattice to understand the magnetism of Fe\(_3\)Sn\(_2\) thin films. The Hamiltonian reads

\[
H = \sum_{\langle i,j \rangle} [-\mathbf{S}_i \cdot \mathbf{S}_j - D_{xy} - S_z \cdot S_z] + \sum_{i} -\mathbf{A}(\mathbf{S}_i)^2 - \mu_0 \mathbf{H} \mathbf{S}_i^2.
\]

where \(\mathbf{S}_i\) represents a classical spin with fixed length \(|\mathbf{S}_i| = 1\) on \(i\)th site. The first sum of \(\langle i,j \rangle\) runs over all the nearest neighbor sites; \(D_{xy}\) and \(D_z\) represent the ferromagnetic...
Heisenberg coupling and the DMI, respectively. The second sum represents the single ion anisotropy $A$ and the Zeeman coupling to the external magnetic field. $H$ perpendicular to the kagome plane. We set the direction of $D_3$, as shown in Fig. 2e, from the symmetry point of view. The out-of-plane component $D_3$ is inherent to the inversion symmetry breaking on the bond centers on the kagome lattice (bulk DMI), while in-plane component $D_1$ arises from the breaking of mirror symmetry due to the attached Pt layer (interfacial DMI). Note that the latter satisfies $C_c$ and $C_s$ rotational symmetries around the center of the triangular and hexagonal plaquettes of the kagome lattice, respectively. We adopted $J = D_1$; similar conditions have been used in previous studies. The magnetic field $H$ was set to be weak so that the magnetization is not forced to be parallel to the applied $H$ direction. In the finite-temperature analysis shown in Fig. 4f, we set $D_1 = 0.25 < J = D_2$. Because $D_1$ is a contribution that is effective only at the interface, this assumption is qualitatively valid. Positive and negative $D_1$ gave the identical simulation results. As for the single ion anisotropy $A$, we set $A = 0.1$ by considering the saturation field $0.5 T$ and $T_c > 400 K$ (i.e., $A << J$) in the experimental magnetization data.

To obtain the ground-state spin configuration of (1), we used a combined method of simulated annealing and local optimization. The annealing was performed from the temperature 1 to 0.001 with 100 steps in the logarithmic scale. In each step, we spent 1000 Monte Carlo (MC) sweeps. After the annealing, we optimized the spin directions one by one so as to minimize the local energy with fixed surrounding spins. We repeated 20000 sweeps of this optimization process. For analyzing the finite-temperature properties, we performed MC simulations with 100000 MC sweeps after 100000 thermalization at each temperature. In all the calculations, we considered $N = 32^2$ spins with $L = 48$ under the periodic boundary condition.

To detect the non-coplanar spin structure in the SkX, we computed the solid angle defined as

$$
\Omega_{\text{sol}} = \frac{1}{N} \sum_{i,j,k} \left( \Omega_{\text{sk}} + \frac{1}{2} \left( \Omega_{\text{sk}} + \Omega_{\text{sk}} + \Omega_{\text{sk}} + \Omega_{\text{sk}} + \Omega_{\text{sk}} + \Omega_{\text{sk}} \right) \right)
$$

with the solid angle

$$
\Omega_{\text{sk}} = 2 \arctan \frac{\left( S_i \cdot S_j \right)}{1 + S_i \cdot S_j + S_j \cdot S_k + S_k \cdot S_i}
$$

which is defined in the range of $(-2\pi, 2\pi)$. The sums of $(i, j, k)$ and $(i, j, k, l, m, n)$ run over all the triangular and hexagonal plaquettes of the kagome lattice, respectively, in which $i, j, \ldots$ were assigned in the counterclockwise order when viewed from the $z$ direction. We also computed the magnetizations (Supplementary Fig. 6):

$$
\rho_{\text{KLM}} = -\frac{1}{N} \sum_{i,j} \left( \frac{c_i c_j + h.c.}{2} \right) - J_{\text{KLM}} \left( \sum_{i,j,k} \sigma_i \sigma_j \sigma_k \right) \cdot S_i
$$

with the given spin configurations $S_i$, and computed $\alpha_{\text{H}}$ by using the standard Kubo formula. The first term represents the kinetic energy of itinerant electrons with the nearest neighbor hopping $t_k$, and the second term represents the Hund’s coupling $J_{\text{KLM}}$ between itinerant electron spins defined by the Pauli matrices $\sigma$ and localized spins $S_i$. In the calculations, we took $J_{\text{H}} = 4t_k$ and fixed the electron density at $n_0 = \sum_{i,j} c_i c_j / N = 0.1$. This low $n_0$ is regarded as a dilute limit, which is appropriate to capture the generic feature of the single kagome plane. For simplicity, we set the temperature of the itinerant electrons sufficiently low as $T_{\text{KLM}} = t_k/40$ to focus on the fluctuations of spins.

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Fig. 4 Analysis of $T$- and $t$-dependent THE. a–c $R_{\text{xy}}$ versus $\mu_B H$ curves measured at $T = 400$–250 K for $t = 0.48, 0.40$ and 0.32 nm, respectively. The two curves corresponding to the decreasing-field and increasing-field scans are shown, which are represented by black solid and red dotted curves, respectively. The data are shifted vertically for clarity. The gray-shaded regions represent the $T$ range where no clear THE signals are discerned. d Contour plot of $\Delta R_{\text{xy}}$ (±3 $T$) as a function of $t$ and $T$, which is produced from all data for $t = 0.80$–0 nm. The THE region indicated by the dashed lines corresponds to the unshaded data in Fig. 4a–c. e Contour plot of $R_{\text{xy}}$ as a function of $\mu_B H$ and $T$ for $t = 0.40$ nm. The $\mu_B H_{\text{peak}}$ values at which $R_{\text{xy}}$ becomes maximal are shown by black circles. f Finite-temperature Hall conductivity $\alpha_{\text{H}}$ calculated using the Monte Carlo simulation and the linear-response theory with a Kondo lattice model. The used parameters are $J = D_3 = 1.0$, $A = 0.1$, and $D_1 = 0.25$. Here, $T_{\text{eff}}$ is the effective temperature against $J$. The electron temperature and the electron filling number (up to 2 per site) are set to 0.025 and 0.1, respectively. Black spheres represent the temperature at which $|\alpha_{\text{H}}|$ becomes maximal, and error bars correspond to the full width at half maximum.
Data availability

The data that support the findings of this study are available from the corresponding author upon reasonable request.

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Author contributions

K.F. grew the films and characterized their structural, magnetic, and electrical properties. T.S. and K.T. contributed to the magnetization measurements and analysis. Y.K. and Y.M. performed the Monte Carlo simulation and the calculation using the linear-response theory. K.N. contributed to theoretical interpretations of the experimental results. K.F., Y.K., Y.M. and A.T. wrote the manuscript with input from other authors. A.T. supervised the project.

Competing interests

The authors declare no competing interests.

Additional information

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Correspondence and requests for materials should be addressed to Kohei Fujiwara.

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