

Unraveling proton spin polarization through hypertriton production in heavy-ion collisions

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Ultrarelativistic nuclear collisions create the quark-gluon plasma known as the hottest, least viscous, and most vortical fluid ever produced in terrestrial laboratories. Its vortical structure has been uncovered through the spin polarization of Lambda (Λ) hyperons, attributed to the spin-orbit coupling that transfers the system's orbital angular momentum to the quark spin, which is then inherited by hadrons via quark recombination or coalescence. However, Λ polarization reflects primarily the strange-quark component, leaving the spin dynamics of the up and down quarks largely unexplored. Although the proton is an ideal probe, its stability makes direct measurements experimentally challenging. Here, we propose to unravel proton spin polarization via hypertriton ($^3_\Lambda\text{H}$) measurements, exploiting the fact that spin information is preserved when polarized nucleons and Λ coalesce to form hypertriton. We show that, over a broad range of collision energies, the polarizations of proton, Λ , and hypertriton are related by a simple linear scaling law. Since both Λ and hypertriton polarizations can be measured via their self-analyzing weak decays, this linear relation provides a practical experimental avenue for accessing spin polarizations of protons and neutrons—the dominant baryonic degrees of freedom in nuclear collisions.

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I. INTRODUCTION

Quantum chromodynamics (QCD), the fundamental theory of the strong interaction, predicts that high-energy nuclear collisions can create an exotic state of matter composed of deconfined quarks and gluons, known as the quark-gluon plasma (QGP) [1]. This exotic phase evolves as a nearly perfect fluid and is believed to have existed in the first few microseconds after the Big Bang.

In semicentral collisions, the QGP carries substantial orbital angular momentum, on the order of $10^3\hbar$, oriented perpendicular to the reaction plane [2]. Through the spin-orbit or spin-vorticity coupling, part of this angular momentum can be transferred to the spin degrees of freedom of quarks, and subsequently inherited by hadrons via quark coalescence or recombination at hadronization [3]. This mechanism gives

rise to the phenomenon of global spin polarization [3–6], analogous to the Barnett and Einstein-de Haas effect in condensed matter [7,8]. Experiments at the Relativistic Heavy Ion Collider (RHIC) have observed global spin polarization of Λ and $\bar{\Lambda}$ hyperons, confirming that the QGP is the most vortical fluid ever observed [2]. Since then, significant progress has been made to understand the role of spin degrees of freedom in the dynamics of QGP [9–21].

Despite these advances, a crucial piece in our understanding of spin dynamics in hot and dense QCD matter remains missing: The spin polarization of protons and neutrons—the fundamental constituents of finite nuclei and the dominant baryonic degrees of freedom in nuclear collisions—has not been measured. In the SU(6) quark model, the proton polarization ($\mathcal{P}_p$) is directly related to the u and d quark polarizations and can be approximated as $\mathcal{P}_p \approx (4\mathcal{P}_u - \mathcal{P}_d)/3$ [3]. Measuring $\mathcal{P}_p$ is therefore essential for constraining the polarization of light quarks in the QGP, complementing Λ measurements that primarily probe strange quarks since $\mathcal{P}_\Lambda \approx \mathcal{P}_s$ in the quark model considerations [3]. Moreover, proton polarization measurements are indispensable for understanding spin-dependent transport phenomena [22,23], including the hadronic spin Hall effect [24]. Unlike hyperons, however, protons and neutrons do not decay within the detector, rendering conventional decay-based polarization measurements infeasible. A direct measurement would require external polarimeters based on secondary interactions, which are not yet implemented in current heavy-ion detectors [25].

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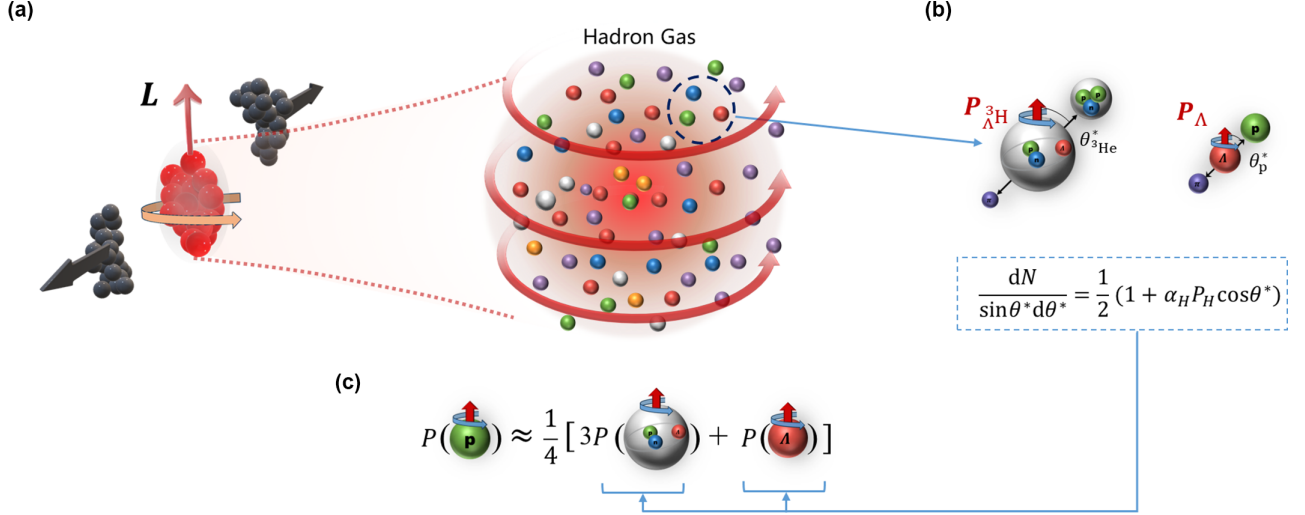


FIG. 1. Schematic illustration of decoding proton spin polarization in heavy-ion collisions. (a) In noncentral heavy-ion collisions, the overlap region of colliding nuclei carries a large orbital angular momentum L , leading to spin polarization of quarks through the spin-orbit interaction. Hadrons emerging from the QGP at hadronization acquire spin polarization via coalescence of polarized quarks, and subsequently evolve through the hadronic phase until the kinetic freeze-out. (b) Hypertritons (${}^3_\Lambda\text{H}$) are polarized when formed via the coalescence of polarized nucleons and Λ hyperons at the kinetic freeze-out. Both ${}^3_\Lambda\text{H}$ and Λ undergo self-analyzing weak decays (e.g. $\Lambda \rightarrow \pi^- + p$ and ${}^3_\Lambda\text{H} \rightarrow \pi^- + {}^3\text{He}$), allowing their spin polarizations to be determined from the angular distributions ($dN/\sin\theta^* d\theta^*$) of their decay products. (c) The proton spin polarization (P_p) can then be reconstructed from the polarizations of ${}^3_\Lambda\text{H}$ and Λ using a model-validated linear relation.

Here, we propose an approach to access proton spin polarization by combining independent measurements of Λ and ${}^3_\Lambda\text{H}$ polarizations, both of which can be determined from their parity-violating weak decays, e.g., $\Lambda \rightarrow \pi^- + p$ and ${}^3_\Lambda\text{H} \rightarrow \pi^- + {}^3\text{He}$. This approach is based on the idea that when polarized protons, neutrons, and Λ hyperons coalesce into ${}^3_\Lambda\text{H}$ (a loosely bound state of a spin-1 deuteron and a spin-1/2 Λ hyperon [26]) at the kinetic freeze-out in nuclear collisions, their spin information is preserved [22,27]. Based on comprehensive model simulations for noncentral Au + Au collisions over a broad range of collision energies ($\sqrt{s_{\text{NN}}} = 7.7\text{--}200$ GeV), including spin-dependent hadronic interactions and resonance decays, we show that proton spin polarization (P_p) can be reliably reconstructed via a simple linear relation:

$$P_p \approx \frac{1}{4} (3P_{{}^3_\Lambda\text{H}} + P_\Lambda). \quad (1)$$

This relation is a direct reflection of the fact that in the spin wave function of the ${}^3_\Lambda\text{H}$ ($J^\pi = \frac{1}{2}^+$ [28]), the Λ spin aligns with the total spin with 1/3 probability and antialigns with 2/3 probability. A schematic illustration of the proposed method for decoding proton spin polarization via Λ and ${}^3_\Lambda\text{H}$ production is shown in Fig. 1.

II. RESULTS

Global spin polarizations of Λ hyperon and proton in heavy-ion collisions. To test the relation in Eq. (1), we investigate the global spin polarizations of the Λ hyperon, proton, and hypertriton in noncentral Au + Au collisions. The analysis is performed within the multiphase transport (AMPT) model [29], which simulates the collision dynamics through four stages: initialization, partonic evolution, hadronization via quark coalescence, and subsequent hadronic rescatterings

and decays. We extend the model to incorporate spin evolution from hadronization to kinetic freeze-out. At hadronization, particles with nonzero spin, such as nucleons and Λ hyperons, acquire polarization through spin-vorticity coupling, computed via the Cooper-Frye formula [30,31]. Their polarizations are then modified during the hadronic phase by spin-dependent processes, including resonance decays and spin-flip scatterings. After kinetic freeze-out, when interactions cease, spin states remain fixed, allowing the final polarization to be determined. Using this extended framework (see Appendix), we first examine the global polarizations of nucleons and Λ hyperons in noncentral Au + Au collisions at $\sqrt{s_{\text{NN}}} = 7.7\text{--}200$ GeV.

Figure 2(a) shows the energy dependence of the global polarization of midrapidity ($|y| < 1$) Λ hyperons along the direction of total angular momentum at hadronization (blue band) and kinetic freeze-out stages (orange band) in collisions at impact parameter $b = 8$ fm, corresponding to 20%–50% centrality. For comparison, experimental data from the STAR Collaboration are represented by filled symbols for Λ hyperons [2,32] with combined statistical and systematic uncertainties. The global spin polarization of Λ hyperons decreases with increasing beam energies, in good quantitative agreement with the STAR data as well as other theoretical calculations summarized in Refs. [6,33,34]. Note that the experimental data points are rescaled using the updated decay parameter of Λ hyperon $\alpha_\Lambda = 0.732 \pm 0.014$ [35].

Our results show that the Λ polarization at kinetic freeze-out is reduced by approximately 20% relative to its value at hadronization. This suppression arises from combined effects of resonance feed-down and spin-flip scatterings during the hadronic phase. Specifically, the feed-down of primary Σ^* resonances tends to enhance the Λ polarization [36], while processes such as $\pi\Lambda \rightarrow \Sigma^*(1385) \rightarrow \pi\Lambda$ scatterings and

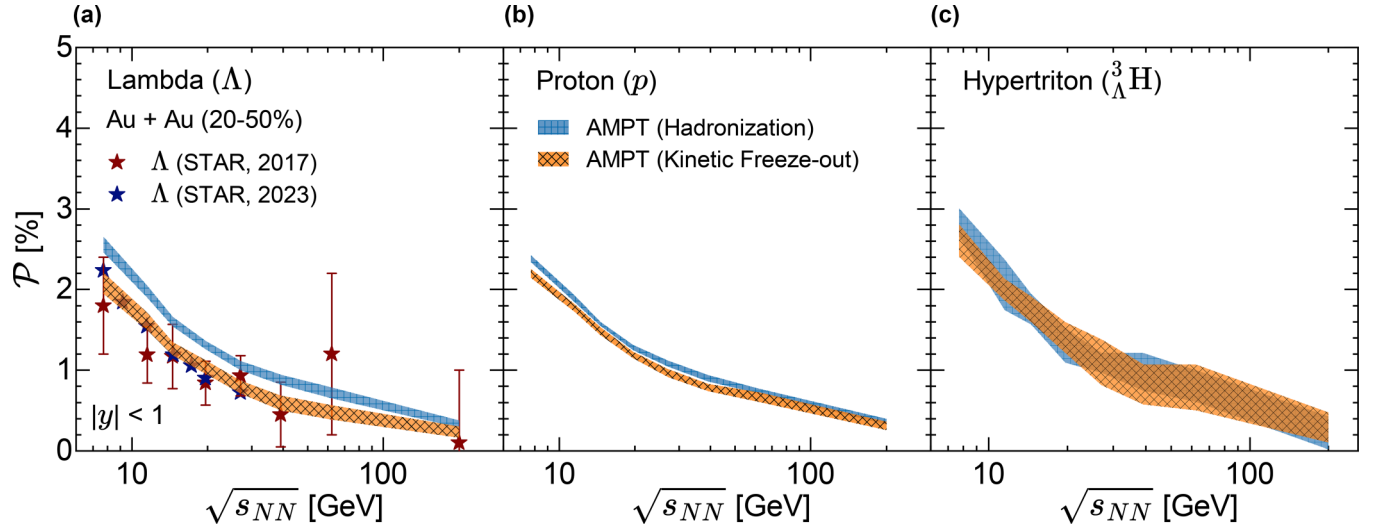


FIG. 2. Collision energy dependence of global spin polarizations of Λ hyperons, protons, and hypertritons. (a) Global polarization of midrapidity ($|y| < 1$) Λ hyperons in Au + Au collisions at $\sqrt{s_{NN}} = 7.7\text{--}200$ GeV and 20%–50% centrality. The blue and orange bands represent model predictions from an extended AMPT model at hadronization and kinetic freeze-out stages, respectively, using an impact parameter $b = 8$ fm. Experimental data for Λ hyperons from the STAR Collaboration are shown with error bars, where the uncertainties combine statistical and systematic contributions [2,32]. (b) Corresponding global polarization of proton at hadronization (blue band) and kinetic freeze-out (orange band). (c) Same as in panel (b), but for hypertritons. Uncertainties in theoretical results are from statistics.

$\Sigma^0 \rightarrow \gamma \Lambda$ decays lead to a reduction. The obtained magnitude of suppression from hadronic rescatterings and decays is consistent with recent studies using rate equations [37,38].

Figure 2(b) depicts the beam energy dependence of proton spin polarization at the hadronization stage (blue band) and at kinetic freeze-out (orange band). At hadronization, the spin polarizations of nucleons and Λ hyperons are found to be close to each other. This similarity arises because, in the Cooper-Frye formalism for polarization (see Appendix), particles with the same spin quantum number and similar masses exhibit similar polarization values, largely independent of their quark content. Unlike the case for Λ hyperons, whose polarization is slightly reduced during the hadronic phase, the nucleon polarization remains almost unchanged between hadronization and kinetic freeze-out stages. This difference results from a near cancellation between two competing effects on proton spin polarization: an enhancement from feed-down decays of primary $\Delta(1232)$ resonances (with spin 3/2), and a suppression from spin depolarization induced by quasielastic scatterings with pions, particularly through the process $\pi N \rightarrow \Delta \rightarrow \pi N$.

Global spin polarization of hypertriton. We next investigate the spin polarization of the ${}^3_{\Lambda}\text{H}$, assuming that it is formed by coalescence of freeze-out nucleons and Λ hyperons. In this case, the nucleon polarization is carried over and encoded in the polarization of the resulting ${}^3_{\Lambda}\text{H}$. The ${}^3_{\Lambda}\text{H}$ can be viewed as a loosely bound state of a spin-1 deuteron and a spin-1/2 Λ with total spin-parity $J^\pi = \frac{1}{2}^+$. This is supported by several experimental observations including its comparable lifetime with free Λ [26,39], extremely small Λ separation energy of $B_\Lambda \approx 0.164$ MeV [40–42], the decay branching ratio [28], and recent calculations based on chiral hyperon-nucleon and hyperon-nucleon-nucleon interactions [43]. The possibility of other spin configurations is less likely but has recently been explored in Ref. [27].

The coalescence model based on the density matrix formalism [44] has been shown to successfully describe ${}^3_{\Lambda}\text{H}$ production in high-energy nuclear collisions [45–47]. The number of ${}^3_{\Lambda}\text{H}$ produced in the coalescence model is given by [27] (see Appendix for more details)

$$N_{{}^3_{\Lambda}\text{H}, \pm \frac{1}{2}} = \sum_{np\Lambda} g_{{}^3_{\Lambda}\text{H}, \pm \frac{1}{2}} W_{{}^3_{\Lambda}\text{H}}(\mathbf{x}_n, \mathbf{x}_p, \mathbf{x}_\Lambda; \mathbf{p}_n, \mathbf{p}_p, \mathbf{p}_\Lambda), \quad (2)$$

where the summation is over all possible combinations of neutron, proton, and Λ hyperon, and $W_{{}^3_{\Lambda}\text{H}}$ is the Wigner function representing the internal structure of the ${}^3_{\Lambda}\text{H}$ in phase space. The $g_{{}^3_{\Lambda}\text{H}, \pm \frac{1}{2}}$ in the above equation denotes the spin statistical factor,

$$\begin{aligned} g_{{}^3_{\Lambda}\text{H}, \pm \frac{1}{2}} &= \frac{2}{3} w_{n, \pm \frac{1}{2}} w_{p, \pm \frac{1}{2}} w_{\Lambda \mp \frac{1}{2}} \\ &\quad + \frac{1}{6} w_{n, \mp \frac{1}{2}} w_{p, \pm \frac{1}{2}} w_{\Lambda \pm \frac{1}{2}} \\ &\quad + \frac{1}{6} w_{n, \pm \frac{1}{2}} w_{p, \mp \frac{1}{2}} w_{\Lambda \pm \frac{1}{2}}, \end{aligned} \quad (3)$$

where $w_{i, \pm \frac{1}{2}} = \frac{1}{2}(1 \pm \mathcal{P}_i)$ are the diagonal elements of the spin density matrix, with $\mathcal{P}_i$ denoting the spin polarization of particle i (nucleon or Λ hyperon). For the case of unpolarized nucleons and Λ hyperons (i.e., $\mathcal{P}_i = 0$), the statistical factor simplifies to $g_{{}^3_{\Lambda}\text{H}, \pm \frac{1}{2}} = g_{{}^3_{\Lambda}\text{H}, -\frac{1}{2}} = 1/8$. In terms of the numbers of ${}^3_{\Lambda}\text{H}$ in spin-up and spin-down states, the ${}^3_{\Lambda}\text{H}$ spin polarization can be determined by $\mathcal{P}_{{}^3_{\Lambda}\text{H}} = (N_{{}^3_{\Lambda}\text{H}, \frac{1}{2}} - N_{{}^3_{\Lambda}\text{H}, -\frac{1}{2}}) / (N_{{}^3_{\Lambda}\text{H}, \frac{1}{2}} + N_{{}^3_{\Lambda}\text{H}, -\frac{1}{2}})$.

Figure 2(c) shows the collision energy dependence of ${}^3_{\Lambda}\text{H}$ polarization, which also decreases with increasing beam energies, resembling the trends for the polarizations of protons and Λ hyperons. The polarization of ${}^3_{\Lambda}\text{H}$ at the kinetic freeze-out stage is almost identical to that at the hadronization stage. Besides, the polarization of ${}^3_{\Lambda}\text{H}$ is slightly larger than

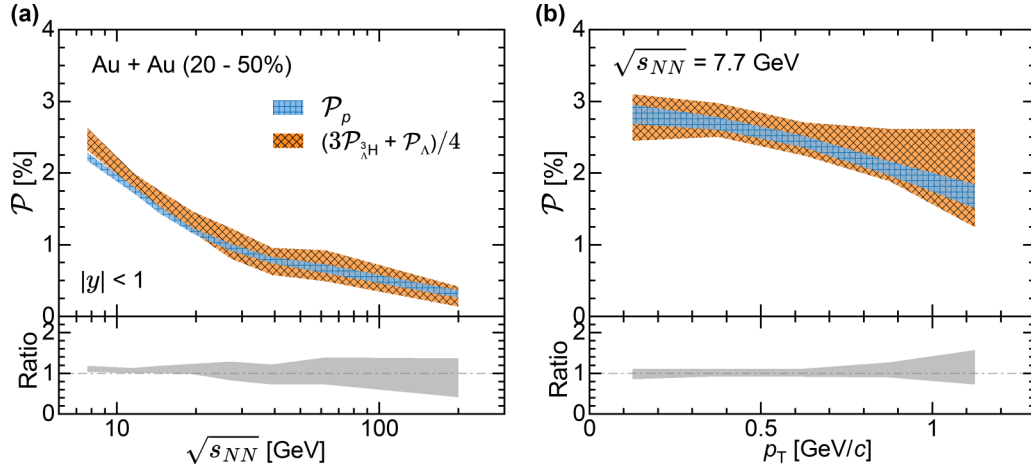


FIG. 3. Decoding proton spin polarization in heavy-ion collisions. (a) Collision energy dependence of global proton spin polarization in noncentral Au + Au collisions at $\sqrt{s_{NN}} = 7.7\text{--}200$ GeV, evaluated at midrapidity ($|y| < 1$). (b) Transverse momentum dependence of proton spin polarization at the energy of 7.7 GeV. The lower panels show the ratio of the extracted proton polarization (orange bands) to its true value (blue bands), defined as the direct AMPT kinetic freeze-out result in the model calculations. Colored bands represent statistical uncertainties in theoretical results.

that of Λ hyperon, which is advantageous for experimental measurements.

Decoding proton spin polarization from Λ and ${}^3_\Lambda\text{H}$ polarizations. We now demonstrate how to unravel the proton spin polarization from polarizations of hypertriton and Λ hyperon. The spin polarization of ${}^3_\Lambda\text{H}$ inherits those of nucleons and Λ hyperons at their kinetic freeze-out. This becomes evident if one assumes that the polarizations of Λ hyperons and nucleons are independent of their momenta, in which case the polarization of ${}^3_\Lambda\text{H}$ from Eq. (2) can be simplified as [27]

$$\mathcal{P}_{{}^3_\Lambda\text{H}} \approx \frac{2}{3}\mathcal{P}_n + \frac{2}{3}\mathcal{P}_p - \frac{1}{3}\mathcal{P}_\Lambda \approx \frac{1}{3}(4\mathcal{P}_p - \mathcal{P}_\Lambda), \quad (4)$$

where we have neglected higher-order terms (since polarizations are typically only a few percent) and assumed isospin symmetry, i.e., $\mathcal{P}_p \approx \mathcal{P}_n$. This relation arises from spin algebra within the coalescence framework, obtained by projecting the spin density matrix of the n - p - Λ system onto the hypertriton wave function. A similar relation can also be derived in a thermal framework under local equilibrium, where $\mathcal{P}_p = \mathcal{P}_\Lambda = \mathcal{P}_{{}^3_\Lambda\text{H}} = \frac{\omega}{2T}$, with ω and T denoting the vorticity and temperature, respectively.

Possible isospin-breaking effects, such as electromagnetic fields or baryon stopping at lower collision energies, may lead to a splitting between proton and neutron polarizations. In such cases, a simultaneous determination of both would require direct measurements of their polarizations using a nucleon polarimeter in future experiments [25].

The dominance of proton polarization in Eq. (4) explains why the ${}^3_\Lambda\text{H}$ polarization changes only slightly from hadronization to kinetic freeze-out in Fig. 2(c), in contrast to the larger change for Λ polarization in Fig. 2(a). Inverting Eq. (4) yields the expression for proton spin polarization given in Eq. (1).

Figure 3(a) compares the model-calculated proton spin polarization (blue band) with the value of $(3\mathcal{P}_{{}^3_\Lambda\text{H}} + \mathcal{P}_\Lambda)/4$ (orange band), obtained from the calculated polarizations of ${}^3_\Lambda\text{H}$ and Λ hyperons in Au + Au collisions. The two

results are in excellent agreement, and their ratio $((3\mathcal{P}_{{}^3_\Lambda\text{H}} + \mathcal{P}_\Lambda)/4\mathcal{P}_p)$, shown in the lower panel, remains consistent with unity within uncertainties. This confirms the validity of Eq. (1) as a reliable method to infer nucleon polarization from ${}^3_\Lambda\text{H}$ and Λ measurements.

Considering that the mass difference between nucleons and the Λ hyperon is small, and that their momenta must be nearly identical when forming a loosely bound hypertriton via coalescence, Eq. (1) can be generalized to include momentum dependence as

$$\mathcal{P}_p(p_T) \approx \frac{1}{4}(3\mathcal{P}_{{}^3_\Lambda\text{H}}(p_T) + \mathcal{P}_\Lambda(p_T)), \quad (5)$$

where p_T denotes the transverse momentum.

Figure 3(b) shows the p_T dependence of $\mathcal{P}_p(p_T)$ and $(3\mathcal{P}_{{}^3_\Lambda\text{H}}(p_T) + \mathcal{P}_\Lambda(p_T))/4$ in Au + Au collisions at $\sqrt{s_{NN}} = 7.7$ GeV. Both quantities exhibit a decreasing trend with p_T and remain in close agreement throughout. Their ratio, shown in the lower panel, stays near unity within uncertainties, further validating the effectiveness of Eq. (1) in capturing both the magnitude and momentum dependence of proton spin polarization. The same agreements hold for other collision energies.

While we employ the AMPT model to provide realistic phase-space distributions and freeze-out conditions, the resulting relation between proton, Λ , and hypertriton polarizations is expected to be largely model independent. Similar conclusions should hold in other frameworks, such as hydrodynamic models with a partonic equation of state coupled to the same freeze-out prescription, provided that small spin correlations and smooth phase-space distributions are maintained. A systematic comparison with such approaches would help quantify possible model dependence, which we leave for future work.

As mentioned earlier, both $\mathcal{P}_\Lambda$ and $\mathcal{P}_{{}^3_\Lambda\text{H}}$ can be determined from their parity-violating weak decays, e.g., $\Lambda \rightarrow \pi^- + p$ and ${}^3_\Lambda\text{H} \rightarrow \pi^- + {}^3\text{He}$. In these decays, the angular distribution of the decay products (p and ${}^3\text{He}$) in the mother's rest

frame follows $dN/\sin\theta^*d\theta^* = (1 + \alpha_H \mathcal{P}_H \cos\theta^*)/2$ [2,27], where α_H (subscript H stands for Λ and ${}^3_\Lambda\text{H}$) is the decay parameter, the polarization $\mathcal{P}_H = (N_{H,\frac{1}{2}} - N_{H,-\frac{1}{2}})/(N_{H,\frac{1}{2}} + N_{H,-\frac{1}{2}})$ characterizes the imbalance between the number of mother particle ($N_{H,\pm\frac{1}{2}}$) in different spin states, and θ^* denotes the angle between the daughter's momentum and the spin direction of the mother particle. The decay parameter of ${}^3_\Lambda\text{H}$ is linearly proportional to that of Λ , i.e., $\alpha_{{}^3_\Lambda\text{H}} \approx -\alpha_\Lambda/2.58$ [27]. These decays are self-analyzing, as the momenta of the daughters carry direct information about the mother's polarization. With the measurements of $\mathcal{P}_\Lambda$ and $\mathcal{P}_{{}^3_\Lambda\text{H}}$, Eqs. (1) and (5) provide a practical method to unravel the proton spin polarization.

We note that hypertriton measurements are statistically challenging because their yields are significantly lower than that of Λ hyperons. However, hypertriton production is expected to reach maximum at $\sqrt{s_{\text{NN}}} = 3\text{--}7$ GeV [48], where future heavy-ion experiments are anticipated to achieve high statistics.

III. CONCLUSION AND OUTLOOK

The spin information of protons and neutrons—the dominant baryonic constituents in nuclear collisions—is largely inaccessible due to their stabilities. To enable such measurements, we have proposed a promising route to unravel the proton spin polarization through independent measurements of the polarizations of Λ hyperons and ${}^3_\Lambda\text{H}$. To validate this method, we perform transport model calculations incorporating spin-dependent scatterings and resonance feed-down effects during the hadronic phase of Au + Au collisions at $\sqrt{s_{\text{NN}}} = 7.7\text{--}200$ GeV. We find that the proton polarization is related to the measurable Λ and ${}^3_\Lambda\text{H}$ polarizations through a simple and robust linear relation given by $\mathcal{P}_p \approx (3\mathcal{P}_{{}^3_\Lambda\text{H}} + \mathcal{P}_\Lambda)/4$.

Accessing the proton spin polarization with our proposed method would open a window on the spin dynamics of up and down quarks, complementing existing insights from Λ polarization that primarily probes strange quarks [3]. It further enables studies of spin-dependent transport phenomena in hadronic matter [22,23] and provides constraints on the equation of state of spin-polarized nuclear matter [49]. Importantly, this method can be readily applied to ongoing and upcoming experiments at RHIC, the High Intensity Heavy-Ion Accelerator Facility [50], the Nuclotron-Based Ion Collider Facility [51], and the Facility for Antiproton and Ion Research [52].

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D.-N.L. performed the numerical simulations and prepared the figures. D.-N.L. and Y.-P.Z. conducted the analytical analysis. K.-J.S. and Y.-G.M. supervised the project and contributed to the interpretation of the results. All authors participated in discussions and contributed to the preparation of the manuscript.

The authors declare no competing interests.

DATA AVAILABILITY

All the data supporting the findings in this work are available within the manuscript and any additional data and code are available from the corresponding authors upon reasonable request.

APPENDIX

Spin polarization of hadrons at hadronization. We calculate the spin polarization of hadrons in Au + Au collisions at $\sqrt{s_{\text{NN}}} = 7.7\text{--}200$ GeV using the AMPT model [29], which has been extensively employed to describe a broad range of observables in relativistic heavy-ion collisions. To describe the spin dynamics of hadrons during the hadronic matter expansion, we extend the AMPT model to include the spin degree of freedom of hadrons and further implement the spin-dependent scatterings and decays.

At hadronization, the spin polarization of hadrons is commonly determined by the thermal vorticity of the QGP under the assumption of local thermal equilibrium [5,31]. The leading-order spin polarization vector of a particle with nonzero spin s , mass m , four-momentum p^μ , and produced at space-time point x^μ is given by [31,53–56]

$$\mathcal{P}^\mu(x, p) = -\frac{s+1}{6m}(1 - n_F)\epsilon^{\mu\nu\rho\sigma}p_\nu\varpi_{\rho\sigma}(x), \quad (\text{A1})$$

where $n_F = 1/[1 + \exp(\beta \cdot p \mp \mu/T)]$ is the Fermi-Dirac distribution function for particles (−) and antiparticles (+). In heavy-ion collisions, $n_F \approx 1$ due to the high temperature and low baryon and strangeness densities. In Eq. (A1), the thermal vorticity tensor $\varpi_{\mu\nu}$ is defined as $\varpi_{\mu\nu} = -\frac{1}{2}(\partial_\mu\beta_\nu - \partial_\nu\beta_\mu)$, with $\beta^\mu = u^\mu/T$ being the four-velocity (satisfying the normalization condition $u^\mu u_\mu = 1$) of the fluid divided by temperature. The temperature field and velocity field are obtained by solving the eigenequation $T^{\mu\nu}u_\mu = \epsilon u^\mu$, where $T^{\mu\nu}$ is the energy-momentum tensor and ϵ denotes the energy density of fluid cells in its local rest frame [57]. The temperature field is estimated as $T = 0.199 \text{ GeV}(\frac{\epsilon/\gamma_q}{1+3N_f/4})^{1/4}$ [58]

with a fugacity factor $\gamma_q = \frac{1}{2}((\frac{N_+}{N_-})^{3/2} + (\frac{N_-}{N_+})^{3/2})$, where N_+ (N_-) represents the number of (anti)quarks in a cell summed over all events. Note that these fields are determined by averaging over a sufficiently large number (typically 10^5) of events

from AMPT and the spin polarization vector is calculated in the particle's rest frame with a Lorentz boost [57].

Effects of hadronic rescatterings. During the subsequent expansion of hadronic matter, the spin polarizations of nucleons and Λ hyperons are modified by resonance feed-down and spin-dependent scatterings, with dominant contributions from $\Delta(1232)$ and $\Sigma^*(1385)$ resonances [37,38]. For a two-body decay of the form $M \rightarrow D + X$, the spin polarization of the daughter particle ($\mathcal{P}_D$) is linearly proportional to that of the mother particle ($\mathcal{P}_M$), i.e.,

$$\mathcal{P}_D = C_{M \rightarrow D} \mathcal{P}_M, \quad (\text{A2})$$

where $C_{M \rightarrow D}$ is the spin transfer coefficient for the decay [59–61]. For strong decays such as $\Delta \rightarrow N\pi$, $\Sigma^* \rightarrow \Lambda\pi$, and $N^* \rightarrow N\pi$, the spin transfer coefficient is $C_{M \rightarrow D} = 1$. In contrast, for the electromagnetic decay $\Sigma^0 \rightarrow \Lambda\gamma$, the spin transfer coefficient is negative, $C_{\Sigma^0 \rightarrow \Lambda} = -\frac{1}{3}$ [59]. As a result, Λ polarization is enhanced by feed-down from Σ^* decays but suppressed by contributions from Σ^0 . Likewise, the feed-down from primary Δ resonances leads to an enhancement of nucleon polarization. For weak decays $\Xi \rightarrow \Lambda\pi$ and $\Omega \rightarrow \Lambda K$, which are not included in the present study, their effects on Λ polarization are negligible because of the relative small number of Ξ and Ω [38,62].

For the inverse process of resonance formation, such as $\pi + D \rightarrow R$, a similar linear relation holds, i.e.,

$$\mathcal{P}_R = C_{D \rightarrow R} \mathcal{P}_D. \quad (\text{A3})$$

For example, in the formation of spin-3/2 resonance Σ^* via $\pi\Lambda \rightarrow \Sigma^*$, the spin transfer coefficient is $C_{D \rightarrow R} = 5/9$ [59,63], corresponding to a Λ spin flip to nonflip ratio of 2/7 in the s channel [37]. The same value, $C_{D \rightarrow R} = 5/9$, applies to Δ formation via $\pi N \rightarrow \Delta$. In contrast, for the formation of a spin-1/2 resonance such as N^* via $\pi N \rightarrow N^*$, the spin transfer coefficient is $C_{D \rightarrow R} = 1$. Consequently, quasielastic scatterings such as $\pi\Lambda \rightarrow \Sigma^* \rightarrow \pi\Lambda$ and $\pi N \rightarrow \Delta \rightarrow \pi N$ reduce the spin polarization of primary Λ hyperons and nucleons to approximately 5/9 of their initial values at hadronization. As shown in Fig. 2(a), the combined effects of resonance

formation and decay suppress the Λ polarization by about 20%, consistent with results obtained from solving rate equations with a parametrized fireball expansion [37,38].

Hypertriton production within a spin-dependent coalescence model. In the covariant coalescence model, the invariant momentum distribution of hypertritons formed from the coalescence of polarized protons, neutrons, and Λ hyperons is given by [27,63]

$$\begin{aligned} E \frac{d^3 N_{\Lambda^3\text{H}, \pm\frac{1}{2}}}{d\mathbf{P}^3} &= E \int \prod_{i=n,p,\Lambda} p_i^\mu d^3\sigma_\mu \frac{d^3 p_i}{E_i} \tilde{f}_i(\mathbf{x}_i, \mathbf{p}_i) \\ &\times \left(\frac{2}{3} w_{n, \pm\frac{1}{2}} w_{p, \pm\frac{1}{2}} w_{\Lambda, \mp\frac{1}{2}} \right. \\ &+ \frac{1}{6} w_{n, \pm\frac{1}{2}} w_{p, \mp\frac{1}{2}} w_{\Lambda, \pm\frac{1}{2}} \\ &+ \left. \frac{1}{6} w_{n, \mp\frac{1}{2}} w_{p, \pm\frac{1}{2}} w_{\Lambda, \pm\frac{1}{2}} \right) \\ &\times W_{\Lambda^3\text{H}} \delta^{(3)}\left(\mathbf{P} - \sum_i \mathbf{p}_i\right), \end{aligned} \quad (\text{A4})$$

where p^μ is the four-momentum, $d\sigma_\mu$ represents the normal vector of the hypersurface of freeze-out particles, $\tilde{f}_i(\mathbf{x}_i, \mathbf{p}_i)$ denotes their spin-averaged phase-space distributions with the coordinate $\mathbf{x}_i$ and momentum $\mathbf{p}_i$ in the frame of the emission source, and $W_{\Lambda^3\text{H}}$ is the Wigner function of ${}^3\Lambda\text{H}$, which we take to have a Gaussian form, $W_{\Lambda^3\text{H}} = 8^2 \exp[-\frac{\rho^2}{\sigma_\rho^2} - \frac{\lambda^2}{\sigma_\lambda^2} - p_\rho^2 \sigma_\rho^2 - p_\lambda^2 \sigma_\lambda^2]$, with ρ , λ , p_ρ , and p_λ denoting the relative distance in coordinate and momentum spaces [64]. The parameters used in the Wigner function are $\sigma_\rho = 2.26$ fm to reproduce the empirical deuteron root-mean-squared radius of $\sqrt{\langle r_d^2 \rangle} = 1.96$ fm [44] and $\sigma_\lambda = 6.52$ fm, which corresponds to the world average Λ separation energy of $B_\Lambda = 0.164$ MeV in ${}^3\Lambda\text{H}$. Using the phase-space information of nucleons and Λ hyperons at kinetic freeze-out (i.e., their positions and momenta at the last interactions) from the extended AMPT model, the integrated yield of hypertriton from Eq. (A4) is simplified to Eq. (2).

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