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Investigating different Λ and $\bar{\Lambda}$ polarizations in relativistic heavy-ion collisions

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A B S T R A C T

Based on the chiral kinetic equations of motion, spin polarizations of various quarks, due to the magnetic field induced by spectator protons as well as the quark-antiquark vector interaction, are studied within a partonic transport approach. Although the magnetic field in QGP enhances the splitting of the spin polarizations of partons compared to the results under the magnetic field in vacuum, the spin polarizations of s and $\bar{s}$ quarks are also sensitive to the quark-antiquark vector interaction, challenging that the different Λ and $\bar{\Lambda}$ spin polarization is a good measure of the magnetic field in relativistic heavy-ion collisions. It is also found that there is no way to obtain the large splitting of the spin polarization between Λ and $\bar{\Lambda}$ at $\sqrt{s_{NN}} = 7.7$ GeV with partonic dynamics.

Understanding the properties of the quark-gluon plasma (QGP) is one of the main purposes of relativistic heavy-ion collision experiments. In noncentral heavy-ion collisions, QGP is expected to be polarized perpendicular to the reaction plane [1–3] due to the large angular momentum as well as the strong magnetic field. Theoretical studies predict that the strong vorticity and magnetic field lead to a series of chiral effects (see, e.g., Ref. [4] for a review) as well as the spin polarizations of hyperons and vector mesons [5–8], which are experimentally measurable through their decays. On the experimental side, continuous efforts have been made on measuring the spin polarization of these particles [9–12]. In the collision systems at higher energies with nearly zero baryon chemical potential, shorter duration of the magnetic field, and smaller angular velocity, the spin polarizations of Λ and $\bar{\Lambda}$ are found to be very small [9,11]. Recently, the finite spin polarizations of Λ and $\bar{\Lambda}$ at lower collision energies have been observed experimentally [12], with the $\bar{\Lambda}$ spin polarization slightly larger than that of Λ. Considerable efforts have been devoted to understanding the polarization of Λ [13–17] but few of them try to address the different spin polarizations of Λ and $\bar{\Lambda}$.

The studies in Refs. [13–17] attribute the hyperon polarization to the coupling to the vorticity field of the QGP, and the spin polarizations of quarks and antiquarks are affected in a similar way. On the other hand, the vector potentials, including those from the quark-antiquark vector interaction and the electromagnetic field, are expected to be responsible for the different polarizations for Λ and $\bar{\Lambda}$ at lower collision energies. Due to the finite baryon chemical potential, quarks and antiquarks are affected by different spin-dependent interactions in the baryon-rich matter. It was also proposed that the difference of the spin polarization between Λ and $\bar{\Lambda}$ can be used as a measure of the magnetic field in relativistic heavy-ion collisions (see, e.g., Ref. [18]), with the strength of the later suffering from the uncertainty of the electrical conductivity of the QGP. The strength of the vector potentials, especially the magnetic field, is responsible for the occurrence of the chiral magnetic effect and the chiral magnetic wave.

In the present study, we investigate the different spin polarizations of Λ and $\bar{\Lambda}$ in Au + Au collisions at $\sqrt{s_{NN}} = 39$ and 7.7 GeV as a result of the vector potentials with partonic transport simulations based on the chiral kinetic equations of motion. The vector potentials include the dominating magnetic field contribution from the spectator protons in the QGP with a temperature-dependent electrical conductivity, and the space component of the quark-antiquark vector potential related to the net quark flux. We found that the s and $\bar{s}$ quark spin polarizations, which are responsible for the Λ and $\bar{\Lambda}$ spin polarizations via the coalescence model, are sensitive to the strength of both the magnetic field and the quark-antiquark vector potential. In addition, there is no way to

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generate a large splitting of the spin polarization between Λ and $\bar{\Lambda}$ at $\sqrt{s_{NN}} = 7.7$ GeV with partonic dynamics.

We generate the initial phase-space information of partons from a multiphase transport (AMPT) model [19], with the momenta of initial partons from melting hadrons produced by the Heavy-Ion Jet INteraction Generator (HIJING) model [20], and their coordinates in the transverse plane (x, y) set to be the same as those of the colliding nucleons that produce their parent hadrons. In order to take into account the finite thickness of the QGP medium at $\sqrt{s_{NN}} = 39$ and 7.7 GeV, the longitudinal coordinates (z) of initial partons are sampled uniformly within $(-lm_N/\sqrt{s_{NN}}, lm_N/\sqrt{s_{NN}})$, where $l = 14$ fm is approximately the diameter of the Au nucleus, and $m_N = 0.938$ GeV is the nucleon mass. Each parton is given a formation time related to the energy and the transverse mass of its parent hadron [19]. Afterwards, the evolution of these partons is described by transport simulations with elastic scatterings between all partons, with the isotropic cross section of 3 mb at 7.7 GeV and 10 mb at 39 GeV, as well as the partonic mean-field potentials. In our previous studies [21,22], these mean-field potentials are taken from a 3-flavor Nambu–Jona–Lasinio (NJL) model, leading to almost zero dynamical mass for partons at high energy densities due to the scalar potential. Since we are only interested in the different spin polarizations of quarks and antiquarks due to vector potentials in the present study, we employ the Lagrangian with only the quark-antiquark vector interaction as well as the external magnetic field for massless partons as follows:

$$\mathcal{L} = \bar{\psi}\gamma_\mu(i\partial^\mu - Q A^\mu_{ext} - \frac{2}{3}G_V\langle\bar{\psi}\gamma^\mu\psi\rangle)\psi. \tag{1}$$

In the above, $\psi = (\psi_u, \psi_d, \psi_s)^T$ is respectively the quark field for u, d, and s quark, $Q = \mathrm{diag}(q_u e, q_d e, q_s e)$ represents respectively their electric charges, and $A^\mu_{ext} = (\varphi, \vec{A}_m)$ is the external electromagnetic potential. The $-\frac{2}{3}G_V\langle\bar{\psi}\gamma^\mu\psi\rangle$ term represents the flavor-singlet quark-antiquark vector interaction after the mean-field approximation [21,22]. The vector coupling constant G_V, whose value affects the critical point of the chiral phase transition in the phase diagram [23–26], is chosen to be 0 or 1.1 times the scalar coupling constant in the original NJL model. The vector density can be expressed as

$$\langle\bar{\psi}\gamma^\mu\psi\rangle = 2N_c \sum_{i=u,d,s} \int \frac{d^3k}{(2\pi)^3 E_i} k^\mu (f_i - \bar{f}_i), \tag{2}$$

where $N_c = 3$ is the color degeneracy, $E_i = k$ is the energy for massless quarks (antiquarks), and f_i and $\bar{f}_i$ are respectively the phase-space distribution functions of quarks and antiquarks of flavor i, which are calculated from the test-particle method [27] by averaging parallel events in the dynamical simulation. As in the original NJL model, the above momentum integration is cut off at 750 MeV [26,40].

The Euler–Lagrange equation for each quark flavor i can be obtained from the Lagrangian [Eq. (1)] as

$$[\gamma^\mu(i\partial_\mu - A_\mu)]\psi_i = 0. \tag{3}$$

In the above, $A_\mu = (A_0, -\vec{A})$ contains the time and space components of the vector potential expressed respectively as

$$A_0 = b_i g_V \rho_0 + q_i e\varphi, \tag{4}$$

$$\vec{A} = b_i g_V \vec{\rho} + q_i e\vec{A}_m, \tag{5}$$

with $g_V = \frac{2}{3}G_V$, $\rho_0 = \langle\bar{\psi}\gamma^0\psi\rangle$ and $\vec{\rho} \equiv \langle\bar{\psi}\vec{\gamma}\psi\rangle$ being respectively the time and space components of the vector density, the baryon charge number $b_i = 1$ for quarks and -1 for antiquarks, and q_i being the electric charge number of the quark species i. φ and $\vec{A}_m$ are the scalar and vector potential of the real external electromagnetic field, and their expressions in vacuum are respectively

$$\varphi(t, \vec{r}) = \frac{e}{4\pi}\sum_n Z_n \frac{1}{R_n - \vec{v}_n \cdot \vec{R}_n}, \tag{6}$$

$$\vec{A}_m(t, \vec{r}) = \frac{e}{4\pi}\sum_n Z_n \frac{\vec{v}_n}{R_n - \vec{v}_n \cdot \vec{R}_n}, \tag{7}$$

where Z_n is the charge number of the nth spectator nucleon, $\vec{v}_n$ is its velocity at the retarded time $t'_n = t - |\vec{r} - \vec{r}_n|$ when the radiation is emitted, and $\vec{R}_n = \vec{r} - \vec{r}_n$ is the relative position of the field point $\vec{r}$ with respect to the nucleon position $\vec{r}_n$. Considering the finite electrical conductivity of the QGP, the vector potential of the electromagnetic field induced by a point particle with charge e moving in the $+z$ direction at the velocity v along the trajectory $z = vt + z_0$ is expressed as [28]

$$\begin{aligned}\vec{A}^e_m = {} & \frac{\hat{z}e}{4\sigma_{con}[(z-z_0)/v]} \times \frac{\exp\left\{\frac{-b^2}{4\{\lambda(t)-\lambda[-(z-z_0)/v]\}}\right\}}{4\{\lambda(t)-\lambda[-(z-z_0)/v]\}} \\ & \times \theta[vt-(z-z_0)]\theta[(z-z_0)-vt_0] \\ & + \frac{\hat{z}ev\gamma}{4\pi}\int_0^{+\infty} dk_\perp J_0(k_\perp b) \\ & \times \exp[-k_\perp^2\lambda(t) - k_\perp\gamma|(z-z_0)-vt_0|].\end{aligned} \tag{8}$$

In the above, t_0 is the time when the QGP emerges, $\sigma_{con}(t)$ is the electrical conductivity of the QGP and $\lambda(t) = \int_{t_0}^t dt'/[\sigma_{con}(t')]$ is related to its time evolution, $\gamma = 1/\sqrt{1-v^2}$ is the Lorentz factor, b is the distance between the field point and the point particle with charge e perpendicular to the z direction, J_0 is the zeroth-order Bessel function of the first kind, and θ is Heaviside step function. Equation (7) is used to calculate $\vec{A}_m$ in vacuum before t_0, and Eq. (8) is used to calculate $\vec{A}_m$ from the summation of $\vec{A}^e_m$ after t_0 when the QGP is produced. Since partons are continuously produced and σ_{con} increases gradually from 0 to finite, t_0 should in principle to be set as early as possible. In the present study we choose $t_0 \sim 0.09$ fm/c, before which there are too few partons leading to large fluctuations.

After decoupling the 4×4 Eq. (3) into the 2×2 Schrödinger equation, the single-particle Hamiltonian can be obtained as

$$H = c\vec{\sigma} \cdot \vec{k} + A_0, \tag{9}$$

where $\vec{k} = \vec{p} - \vec{A}$ is the real momentum of the particle with $\vec{p}$ being its canonical momentum, c is the helicity of the particle, and $\vec{\sigma}$ are the Pauli matrics. In the semiclassical limit by considering $\vec{\sigma}$ as the expectation value of the particle spin, the canonical equations of motion from the above single-particle Hamiltonian are

$$\dot{\vec{r}} = c\vec{\sigma}, \tag{10}$$

$$\dot{\vec{k}} = c\vec{\sigma} \times \vec{B} + \vec{E}, \tag{11}$$

$$\dot{\vec{\sigma}} = 2c\vec{k} \times \vec{\sigma}, \tag{12}$$

where $\vec{B} = \nabla \times \vec{A}$ and $\vec{E} = -\nabla A_0 - \frac{\partial\vec{A}}{\partial t}$ are the total space and time components of the vector potential, including the contributions from the real electromagnetic field originated from the spectator protons and the effective electromagnetic field originated from the quark-antiquark vector interaction. Using the adiabatic approximation $\vec{\sigma} \approx c\hat{k} - \frac{\hbar}{2k}\hat{k} \times \dot{\hat{k}}$ that satisfies $\hat{k} \cdot \dot{\vec{r}} \approx 1 + O(\hbar^2)$, the chiral kinetic equations of motion can be obtained as [29–31]