Hypertriton production in relativistic heavy ion collisions

Zhen Zhang *, Che Ming Ko

Cyclotron Institute and Department of Physics and Astronomy, Texas A&M University, College Station, TX 77843, USA

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Based on the phase-space distributions of freeze-out nucleons and Λ hyperons from a blast-wave model, we study hypertriton production in the coalescence model. Including both the coalescence of Λ with proton and neutron as well as with deuteron, which is itself formed from the coalescence of proton and neutron, we study how the production of hypertriton is affected if nucleons and deuterons are allowed to stream freely after freeze-out. Using central Pb+Pb collisions at √sNN = 2.76 as an example, we find that this only reduces slightly the hypertriton yield, which has a value consistent with the experimental data, even if the volume of the system has expanded to a size similar to the freeze-out volume for a hypertriton if its dissociation cross section by pions in the system is given by its geometric size. Our results thus suggest that the hypertriton yield in relativistic heavy ion collisions is essentially determined at the time when nucleons and deuterons freeze out, although it still undergoes reactions with pions.

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1. Introduction

Besides allowing the opportunity to study the properties of strongly interacting matter at extreme temperature and density, high energy heavy-ion collisions also provide the possibility to produce hypernuclei that contain strange baryons [1–3]. Because of the abundant anti-strange quarks produced in ultrarelativistic heavy ion collisions, anti-hypernuclei can also be produced in these collisions. Indeed, both hypertriton, which is a bound state of proton, neutron and Λ hyperon, and its anti-nucleus, i.e., the antihypertriton, have been detected at the Relativistic Heavy Ion Collider (RHIC) by the STAR Collaboration [4] and at the Large Hadron Collider (LHC) by the ALICE Collaboration [5] through their weak decays $ \frac{3}{1} H \rightarrow \frac{3}{1} He + \pi^- \text{ and } \frac{3}{1} \Lambda \rightarrow \frac{3}{1} He + \pi^+$, respectively, with the same branch ratio of about 25% [6]. The study of hypernuclei production in high-energy heavy-ion collisions is also of interest because it can provide information on local baryon and strangeness correlations in the collisions [7], if their production is through the coalescence of protons, neutrons and Λ hyperons at the final stage of the collisions [8]. For example, the ratio $S_3 = \frac{3}{1} H / (\frac{3}{1} He \times \frac{1}{1} p)$ [9] has been suggested as a possible probe of the onset of deconfinement in high-energy heavy-ion collisions [10].

In addition to the coalescence model mentioned in the above, hypernuclei production in relativistic heavy ion collisions has also been studied using the statistical model [3], in which their abundances are determined by assuming that they are in chemical equilibrium with other hadrons and nuclei at a chemical freeze-out temperature that is close to that for the quark-gluon plasma to hadronic matter phase transition. This is in stark contrast to the coalescence model, which assumes that hypernuclei are formed from the coalescence of protons, neutrons and Λ hyperons at the kinetic freeze-out of heavy ion collisions [8,11]. Both models are, however, quite successful in describing the experimental data, although slightly earlier freeze out of Λ hyperons than nucleons is introduced in the coalescence model study of Ref. [11]. The reason for this may be due to the fact that their numbers remain unchanged during the hadronic evolution from the chemical to the kinetic freeze out, like the deuteron [12] and other particles, which has recently been shown to be associated with the conservation of entropy per particle [13].

Given its very small binding energy of about 130 keV [14] and large root-mean-square radius of about 4.9 fm [15], the hypertriton is, however, expected to be formed later than the kinetic freeze out time for nucleons and Λ hyperons due to its larger dissociation cross section by pions and thus shorter mean-free path than these hadrons. To illustrate the effect of the large hypertriton size on its production in relativistic heavy ion collisions, we use the coalescence model based on the phase-space distributions of freeze-out nucleons and Λ hyperons from a blast-wave model. In particularly, we use the blast-wave models FOA-N and FOA-Λ of Refs. [11,16] for central (0%-10% centrality) Pb+Pb collisions at √sNN = 2.76 TeV [5]. Besides the coalescence of proton, neutron and Λ hyperon...
we also include the coalescence process \( d + \Lambda \rightarrow ^3\text{H} \) between
the deuteron \( d \) and \( \Lambda \), which is found to enhance the hypertriton yield
by about a factor of two, to take account of the fact that the hypertriton wave function is
dominated by a \( \Lambda \) hyperon that is loosely bonded to a deuteron [5,15,17]. We find that the yield of hypertriton
obtained from the coalescence model does not change much
if it is calculated with nucleons and \( \Lambda \) that are allowed to stream freely after freeze out to a volume that is nineteen times larger,
similar to the hypertriton freeze-out volume that is obtained with a
dissociation cross section given by its geometrical size.

This paper is organized as follows. In Secs. 2 and 3, we briefly
introduce the blast-wave fireball model and the coalescence model,
respectively. Our results and related discussions on the hypertriton
yield in central Pb–Pb collisions at \( \sqrt{s_{NN}} = 2.76 \text{ TeV} \) are presented
in Sec. 4. Finally, we give the conclusion in Sec. 5.

2. The blast wave model

Following Refs. [11,16], we use a blast-wave model to describe the phase-space distributions of nucleons and \( \Lambda \) hyperons at the kinetic freeze out of a heavy ion collision. With the assumption
that the longitudinal proper time \( \tau = \sqrt{t^2 - z^2} \) for the freeze-out
hypersurface \( \Sigma^\mu \) has a Gaussian distribution,
\[
J(\tau) = \frac{1}{\sqrt{2\pi} \Delta \tau} \exp \left[ -\frac{(\tau - \tau_0)^2}{2(\Delta \tau)^2} \right],
\]
with a mean value \( \tau_0 \) and a dispersion \( \Delta \tau \), the single-particle invariant
momentum spectrum of nucleons or \( \Lambda \) is then given by
\[
\frac{d^3N}{d^3p} = \int d\tau J(\tau) \int \frac{d^3\sigma_\mu}{\Sigma^\mu} f(x,p),
\]
where \( \sigma_\mu \) is the covariant normal vector to \( \Sigma^\mu \) and \( p^\mu \) is the four-momentum
of the emitted particle. The invariant distribution of these emitted particles from the hyper-surface is
\[
f(x,p) = \frac{g}{(2\pi)^3} \left[ \exp(-p^\mu u_\mu /T) / \xi \pm 1 \right]^{-1},
\]
where \( g \) is the spin degeneracy factor of the particle, \( u_\mu \) is the
flow four-velocity, \( \xi \) is the fugacity parameter determined by the
number of emitted particles, and \( T \) is the temperature of the fireball.
Taking the longitudinal flow velocity to be \( v_L = z/t \), the longitudinal
flow rapidity is then \( \eta_{\text{flow}} = \frac{1}{2} \ln(1 + v_L) / (1 - v_L) \) and is identical
to the space–time rapidity \( \eta = \frac{1}{2} \ln((t + z) / (t - z)) \). In terms of \( \eta \), the transverse flow rapidity \( \rho \equiv \frac{1}{2} \ln(1 + \beta) / (1 - \beta) \) with \( \beta \) being the magnitude of the transverse flow velocity, one then has
\[
p^\mu u_\mu = m_T \cosh \eta \cosh(\eta - \phi) - p_T \sinh \eta \cos(\phi - \phi_0),
\]
\[
p^\mu d\sigma_\mu = m_T \cosh(\eta - \phi) d\eta d\phi.
\]
In the above, \( p_T \) and \( m_T = \sqrt{m^2 + p_T^2} \) are the transverse
momentum and mass with \( m \) being the particle mass, \( \phi_0 \) and \( \phi_0 \) are the azimuthal angles of the particle transverse momentum and the transverse flow velocity with respect to the reaction plane, and \( r \) and \( \phi \) are the radial and angular coordinates of the particle in the transverse plane. For central heavy-ion collisions considered in
the present study, the azimuthal angle of the transverse flow velocity \( \phi_0 \) is equal to \( \phi \) and the transverse flow rapidity of the fluid element in the fireball can be parametrized \( \rho = \rho_0 \rho / R_0 \), with \( \rho_0 \) being the maximum transverse flow rapidity and \( R_0 \) the transverse radius of the fireball. The phase-space distributions of freeze-out
particles are thus determined by the parameters \( T, \rho_0, R_0, \rho, \Delta \tau, \xi_\rho, \xi_\phi, \xi_\Lambda \) in the blast-wave model.

We note that by assuming freeze-out at constant longitudinal proper,
we have neglected the effect that the edge of the fireball freezes out much earlier than the center. Although this effect can be included by free-streaming the frozen-out particles from the
realistic decoupling surface to one of constant longitudinal proper
time, which is expected to render these particles no longer perfectly thermal, it is not expected to significantly change our results
as the parameters used in the present study are tuned to reproduce
the experimental data on the transverse momentum spectra
of nucleons and \( \Lambda \) hyperons.

3. The coalescence model

In the coalescence model, the production probability of a cluster is determined by the overlap of the Wigner function of its internal wave function with the phase-space distributions of the constituent particles at the kinetic freeze-out of heavy ion collisions. The multiplicity of a cluster containing \( M \) particles can then
be written as [16,18–23]
\[
M = g_M \int \prod_{i=1}^{M} \frac{d\tau_i}{\Gamma_1} f(\tau_i, p_i)
\]
\[
\times f_M(x_1, \ldots, x_M; p_1, \ldots, p_M)
\]
where \( f_M(x_1, \ldots, x_M; p_1, \ldots, p_M) \) is the Wigner function of the cluster and \( g_M = (2S + 1) \left( \prod_{i=1}^{M} (2I_i + 1) \right) \) is the statistical factor
with \( S_i \) and \( S \) being the spins of the \( i \)th constituent particle and the cluster, respectively. The coordinate \( x_i \) and momentum \( p_i \) are those of the \( i \)th particle in the fireball frame, while \( x'_i \) and \( p'_i \) are the corresponding ones after Lorentz transforming to the rest
frame of the produced cluster and then propagating earlier freeze-out
particles freely to the time when the last particle in the cluster
freezes out [20,22].

The Wigner function of a cluster is obtained from the Wigner transform of its internal wave function, which is usually taken to
be the product of harmonic oscillator wave functions [16,20–23]. For example, using the ground-state wave function of a harmonic oscillator for the deuteron, its Wigner function is given by
\[
f_d(\rho, \rho_p) = 8 \exp \left[ -\frac{\rho^2}{\sigma^2} - \frac{\rho_p^2}{\sigma^2} \right].
\]
where the relative coordinate \( \rho \) and momentum \( \rho_p \) between the
proton and neutron are defined, respectively, by
\[
\rho = \frac{1}{\sqrt{2}}(x'_p - x'_n),
\]
\[
\rho_p = \sqrt{2 \frac{m_n p'_n - m_p p'_p}{m_p + m_n}}.
\]
with \( m_p \) and \( m_n \) being the masses of proton and neutron,
respectively. The width parameter \( \sigma_\rho \) in Eq. (7) is related to the
root-mean-squared charge radius of deuteron by
\[
\langle r^2_\rho \rangle = \frac{3 m_p^2}{(m_p + m_n)^2} \sigma_\rho^2.
\]
Using the deuteron charge radius \( \sqrt{\langle r^2_\rho \rangle} = 2.142 \text{ fm} \) [24] and \( m_p = m_n = 0.939 \text{ GeV} \), we find the width parameter to have the value
\( \sigma_\rho = 2.473 \text{ fm} \). In the present study, we will use the Wigner
function given by Eq. (7) in the coalescence model to study deuteron
production.
Because of the long tail of the $\Lambda$ wave function in the hypertriton [15], we use instead the following product of sums of Gaussian functions for the hypertriton wave function:

$$\Psi(\rho, \lambda) = \left( \sum_i a_i e^{-u_i r_i^2} \right) \left( \sum_j b_j e^{-v_j \lambda^2} \right).$$

In the above, $\rho$ is similarly defined as in Eq. (8) and $\lambda$ is defined by

$$\lambda = \sqrt{\frac{7}{3}} \left( \frac{m_p x_p + m_n x_n}{m_p + m_n} - x'_\Lambda \right).$$

The coefficients $a_i$ and $b_j$ as well as the width parameters $u_i$ and $v_j$ can be determined by fitting the hypertriton rms radius and the density distributions of proton or neutron and $\Lambda$ in hypertriton obtained from the stochastic variational calculations [15]. We list their values in Table 1 and show the density distributions of nucleons and $\Lambda$ in $^3\Lambda$H as functions of the distance from the center-of-mass of proton and neutron in Fig. 1.

The Wigner function or phase-space density of $^3\Lambda$H can be straightforwardly evaluated and is given by

$$f_{^3\Lambda} = 8 \left[ \sum_{i,j} \frac{\pi}{u_i + u_j} \right]^{3/2} a_i a_j \exp \left( - \frac{p_{\rho}^2}{u_i + u_j} \right) \times \exp \left( - \frac{4u_i u_j p_{\rho}^2}{u_i + u_j} \cos \left( \frac{2u_i - u_j}{u_i + u_j} \mathbf{p}_{\rho} \cdot \mathbf{\rho} \right) \right) \times 8 \left[ \sum_{m,n} \frac{\pi}{v_m + v_n} \right]^{3/2} b_m b_n \exp \left( - \frac{p_{\lambda}^2}{v_m + v_n} \right),$$

with $p_{\rho}$ similarly defined as in Eq. (9) and $p_{\lambda}$ defined by

$$p_{\lambda} = \sqrt{\frac{3}{2}} \frac{m_\Lambda (p_{\rho} + p_{\rho}'')}{m_p + m_n + m_\Lambda}.$$

4. Results

In the present study, we consider, as an example, the production of hypertriton in central Pb+Pb collision at $\sqrt{s_{NN}} = 2.76$ TeV. For the parameters in the blast-wave model, we use the FOPI-N and FOPI-Λ parameter sets reported in Refs. [11,16], i.e. $T = 121.1$ MeV, $\rho_0 = 1.215$, $\rho_9 = 19.7$ fm, $\rho_9 = 15.5$ fm/c, $\Delta \tau = 1.0$ fm/c, $\xi_p = \xi_n = 3.72$ and $\xi_\Lambda = 9.54$, which are obtained by fitting the measured spectra of $p$, $d$, $^3$He and $\Lambda$ from the ALICE Collaboration.

Using the hypertriton Wigner function in Eq. (13), we find that the coalescence of $n + p + \Lambda$ gives the $^3\Lambda$H yield $\frac{dN}{dy} |_{y=0} = 6.7 \times 10^{-5}$, which is almost a factor of two smaller than the measured yield of $1.54 \pm 0.41 \times 10^{-4}$ [5]. Compared to the hypertriton yield of $5.72 \times 10^{-5}$ in Ref. [11], which is calculated from the hypertriton Wigner function by assuming that the proton, neutron and $\Lambda$ are all in the ground state of an isotropic harmonic oscillator, our $^3\Lambda$H number is only 17% larger, indicating that details of the hypertriton wave function do not significantly affect the hypertriton yield in the coalescence model.

Including also the $d + \Lambda \rightarrow ^3\Lambda$H coalescence by treating deuteron as a point particle as introduced in Sec. 3 and with the deuteron phase-space distribution obtained via the coalescence of proton and neutron, we obtain a $^3\Lambda$H yield of $7.9 \times 10^{-5}$ from the $d + \Lambda$ coalescence, which is close to that from the coalescence of neutron, proton and $\Lambda$. The total $^3\Lambda$H yield from the two contributions is $1.46 \times 10^{-4}$ and is consistent with the experimental value within its uncertainty.

By producing hypertritons from the coalescence of freeze-out nucleons, $\Lambda$ hyperons, and deuterons as in the above calculations, it is assumed that hypertritons also freeze out at similar times. Since the hypertriton is weakly bounded ($\sim 130$ keV) and has a size comparable to the volume of the kinetically freeze-out fireball, particularly for the distance between the $\Lambda$ and the two nucleons, it can be easily dissociated by collisions with the abundant pions in heavy ion collisions at the LHC energies [25] and is thus unlikely to freeze out at the same time as nucleons and $\Lambda$ hyperons. To determine the freeze-out time for the hypertriton relative to that of nucleons, one needs its scattering cross section with pion compared to that of nucleon. Although the isospin averaged pion–nucleon scatter cross section is known to have a value of about 100 mb [26], there is no empirical information on the hypertriton-pion scattering cross section. Since the typical momentum of a pion at the kinetic freeze-out temperature of about 120 MeV is about 300 MeV, corresponding to a wavelength of less than 1 fm, which is much smaller than the size of hypertriton,

\[ \frac{dN}{dy} |_{y=0} = 6.7 \times 10^{-5} \]
one can approximate the hypertriton dissociation cross section by its quasi-elastic scattering cross section of about three times the pion–nucleon scattering cross section, if one takes the scattering cross section of pion and $\Lambda$ to have the same value as that for pion and nucleon. According to Ref. [27] based on the comparison of the particle scattering rate with the expansion rate of a fireball that undergoes the boost invariant longitudinal expansion with a constant transverse expansion velocity, the volume of the fireball at which a particle freezes out is related to its scattering cross section $\sigma$ by

$$V_f = \frac{1}{\sqrt{2\pi}} \left( \frac{N_\pi \sigma}{3} \right)^{3/2},$$

with $N_\pi$ being the number of pions in the fireball. The freeze-out volume for a hypertriton in a fireball with a fixed number of pions is thus about five times larger than that for a nucleon.

If the hypertriton-pion scattering is coherent, which would be the case for very low momentum pions, its dissociation cross section would be nine times larger than that for the pion–nucleon scattering, resulting in a hypertriton freeze-out volume about 27 times larger than that for nucleons. On the other hand, using the geometric cross section based on the hypertriton rms radius of $R = 4.9$ fm, the hypertriton dissociation cross section would be about 750 mb, and the resulting hypertriton freeze-out volume would be about eighteen times larger than that of protons.

To study the dependence of the hypertriton yield on its freeze-out time, we consider the free expansion of the fireball by letting nucleons, $\Lambda$ and deuterons to stream freely with constant velocities from their initial positions given by the blast-wave model and the coalescence model. The size of the fireball can then be determined from the rms transverse radius $r_{\text{rms}}(\Delta \tau) = \sqrt{(x^2 + y^2)}$ and longitudinal length $z_{\text{rms}}(\Delta \tau) = \sqrt{z^2}$ after the proper free-streaming time $\Delta \tau$. In Fig. 2, we show the ratios $r_{\text{rms}}(\Delta \tau)/r_{\text{rms}}(0)$ and $z_{\text{rms}}(\Delta \tau)/z_{\text{rms}}(0)$ as functions of $\Delta \tau$. It is seen that after letting all particles stream freely by $\Delta \tau = 35$ fm/c, the $r_{\text{rms}}$ increases by a factor of 2.77 while the $z_{\text{rms}}$ increases by a factor 2.47, corresponding to an effective volume ($\propto r_{\text{rms}}^2z_{\text{rms}}$) increase of about nineteen times larger than its initial value.

Fig. 3 shows the dependence of the hypertriton yield on the free streaming time $\Delta \tau$. It is seen that the number of hypertritons from the $d + \Lambda$ coalescence almost remains unchanged during the free expansion of the fireball, and that from the $n + p + \Lambda$ coalescence also decreases slowly. As a result, the total hypertriton number only decreases by about 20% at $\Delta \tau = 35$ fm/c, corresponding to an increase of the fireball volume by a factor of 19. This is due to the fact that the constituent nucleons and $\Lambda$ hyperons that are likely to form hypertritons have small relative velocities, and they thus are most likely to move together, which is especially the case for the process $d + \Lambda \rightarrow ^3\text{H}$ in which the deuteron is treated as a point particle and has a long tail in its relative wave function with respect to the $\Lambda$. The weak dependence of the hypertriton yield on its freeze-out time suggests that although the hypertriton should be produced at a much later time than the freeze-out times of nucleon, $\Lambda$ hyperons, and deuterons, the hypertriton abundance is essentially determined when nucleons and $\Lambda$ hyperons freeze out from the fireball, and the coalescence calculations based on particles at initial freeze-out hypersurface can still reasonably describe the production of hypertriton in relativistic heavy ion collisions.

5. Conclusions

Using the coalescence model based on the phase-space distributions of kinetically freeze-out nucleons and $\Lambda$ hyperons from a blast-wave model, we have studied hypertriton production in a coalescence model, which includes not only the coalescence process $p + n + \Lambda \rightarrow ^3\text{H}$ but also the coalescence process $d + \Lambda \rightarrow ^3\text{H}$ as the hypertriton can be considered as a loosely bound state of deuteron and $\Lambda$, to study the dependence of the $^3\text{H}$ yield on its freeze-out time by letting nucleons, $\Lambda$ hyperons, and deuterons to stream freely after they have frozen out from the initial fireball, and then carrying out the coalescence calculations for different free streaming times. We have found that the $^3\text{H}$ yield, which reproduces the experimental data from central Pb+Pb collisions at $\sqrt{s_{\text{NN}}} = 2.76$ TeV at the LHC with the two coalescence processes giving similar contributions, decreases slowly with the free expansion of the fireball, especially for those produced from the $d + \Lambda$ coalescence. Our result thus indicates that the hypertriton yield in relativistic heavy ion collisions is essentially determined when nucleons and $\Lambda$ hyperons freeze out kinetically, although it still undergoes scattering with the freeze-out pions. The present conclusion is based on freeze-out nucleons and $\Lambda$ hyperons from a simple blast-wave model. It will be of great interest to carry out similar studies based on more realistic models, such as the hybrid hydrodynamical+Boltzmann approach [28] and transport models [29,30], which treat properly the freeze-out of nucleons and $\Lambda$ hyperons, to obtain a more accurate description of hypertriton production in relativistic heavy ion collisions.

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References

[1] C.M. Ko, Phys. Rev. C 32 (1985) 326.
[2] P. Koch, B. Müller, J. Kafelski, Phys. Rep. 142 (1986) 167.
[3] A. Andronic, P. Braun-Munzinger, J. Stachel, H. Stöcker, Phys. Lett. B 697 (2011) 203.
[4] B.I. Abelev et al., STAR, Science 328 (2010) 58.
[5] J. Adam, et al., ALICE, Phys. Lett. B 754 (2016) 360.
[6] H. Kamada, J. Golak, K. Miyagawa, H. Witala, W. Gloeckle, Phys. Rev. C 57 (1998) 1595.
[7] J. Steinheimer, M. Mitrovski, T. Schuster, H. Petersen, M. Bleicher, H. Stöcker, Phys. Lett. B 676 (2009) 126, arXiv:0811.4077.
[8] H.H. Gutbrod, A. Sandoval, P.J. Johansen, A.M. Poskanzer, J. Gosset, W.G. Meyer, G.D. Westfall, R. Stock, Phys. Rev. Lett. 37 (1976) 667.
[9] T.A. Armstrong, K.N. Barish, S. Batsouli, S.J. Bennett, M. Bertaina, A. Chikanian, S.D. Coe, T.M. Cormier, R. Davies, C.B. Dover, et al., E864 Collaboration, Phys. Rev. C 70 (2004) 024902.
[10] S. Zhang, J. Chen, H. Crawford, D. Keane, Y. Ma, Z. Xu, Phys. Lett. B 684 (2010) 224.
[11] K.-J. Sun, L.-W. Chen, Phys. Rev. C 93 (2016) 064909.
[12] Y. Oh, C.M. Ko, Phys. Rev. C 76 (2007) 054910.
[13] J. Xu, C.M. Ko, Phys. Lett. B 772 (2017) 290.
[14] M. Jurić, G. Bohm, J. Klubuk, U. Koecker, F. Wysotski, G. Coremans-Bertrand, J. Sacton, G. Wilquet, T. Cantwell, F. Esmail, et al., Nucl. Phys. B 52 (1973) 1.
[15] H. Nemura, Y. Suzuki, Y. Fujiwara, C. Nakamoto, Prog. Theor. Phys. 103 (2000) 929.
[16] K.-J. Sun, L.-W. Chen, Phys. Lett. B 751 (2015) 272.
[17] J.G. Gorgletoń, J. Phys. G 18 (1992) 339.
[18] C.B. Dover, U. Heinz, E. Schniedermann, J. Zimányi, Phys. Rev. C 44 (1991) 1636.
[19] R. Mattiello, H. Sorge, H. Stöcker, W. Greiner, Phys. Rev. C 55 (1997) 1443.
[20] Y. Oh, Z.-W. Lin, C.M. Ko, Phys. Rev. C 80 (2009) 064902.
[21] L. Zhu, C.M. Ko, X. Yin, Phys. Rev. C 92 (2015) 064911.
[22] X. Yin, C.M. Ko, Y. Sun, L. Zhu, Phys. Rev. C 95 (2017) 054913.
[23] L. Zhu, H. Zheng, C.M. Ko, Y. Sun, arXiv:1710.05139 [nucl-th], 2017.
[24] L. Angeli, K.P. Marinova, At. Data Nucl. Data Tables 99 (2013) 69.
[25] B. Abelev, J. Adam, D. Adamová, A.M. Adare, M.M. Aggarwal, G. Aglieri Rinella, M. Agnello, A.G. Agocs, A. Agostinelli, Z. Ahmed, et al., ALICE Collaboration, Phys. Rev. C 88 (2013) 044910.
[26] K.A. Olive, et al., Particle Data Group, Chin. Phys. C 38 (2014) 090001.
[27] S. Cho, T. Hyodo, D. Jido, C.M. Ko, S.H. Lee, S. Maeda, K. Miyahara, K. Morita, M. Nielsen, A. Ohnishi, et al., Prog. Part. Nucl. Phys. 95 (2017) 279.
[28] H. Song, S.A. Bass, U. Heinz, Phys. Rev. C 83 (2011) 024912.
[29] Z.-W. Lin, C.M. Ko, B.-A. Li, B. Zhang, S. Pal, Phys. Rev. C 72 (2005) 064901.
[30] W. Cassing, E.L. Bratkovskaya, Nucl. Phys. A 831 (2009) 215.