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Relativistic transport theory for N, Δ and $N^*(1440)$ system

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Abstract

A self-consistent relativistic Boltzmann-Uehling-Uhlenbeck equation for the $N^*(1440)$ resonance is developed based on an effective Lagrangian of baryons interacting through mesons. The equation is consistent with that of nucleon's and delta's which we derived before. Thus, we obtain a set of coupled equations for the N, Δ and $N^*(1440)$ distribution functions. All the $N^*(1440)$-relevant in-medium two-body scattering cross sections within the N, Δ and $N^*(1440)$ system are derived from the same effective Lagrangian in addition to the mean field and presented analytically. Medium effects on the cross sections are discussed.

I. Introduction

It was recognized twenty years ago that particles emitted in the collisions contain important information about the equation of state of hot and dense nuclear matter. Since most of the particles such as pion, kaon, dilepton, anti-proton, anti-kaon are mainly produced through resonances, the inclusion of resonance degrees of freedom in transport theories is essential for any realistic description of relativistic heavy-ion collisions. Recent theoretical calculations [1, 2] and experimental data [3] indicated that a gradual transition to *resonance matter* would occur in the collision zone at kinetic energy ranging from SIS ($\sim$ 1AGeV) up to AGS ($\sim$ 15AGeV). At an incident energy of 2 GeV/nucleon more than 30% of the nucleons are excited to resonance states [4]. At intermediate- and high-energy range the most important baryonic resonances are $\Delta(1232)$, $N^*(1440)$ and $N^*(1535)$. Theoretical models extended to describe relativistic heavy-ion collisions at this energy range should include these resonance degrees of freedom explicitly and treat them self-consistently. The heart of the problem is to determine quantitatively all possible binary collisions relating to resonances, such as $N\Delta$, $\Delta\Delta$, $NN^*(1440)$, $NN^*(1535)$ ... collisions. Unfortunately, very little is known about resonance-relevant in-medium cross sections in high-density nuclear matter since the experimental determination of them is inaccessible yet.

A	m_A	g^A_{NN}	$g^A_{N^*N^*}$	$g^A_{\Delta\Delta}$	$g^A_{NN^*}$	γ_A	τ_A	T_A	$\Phi_A(x)$	D^μ_A	D^i_A
σ	m_σ	g^σ_{NN}	$g^\sigma_{N^*N^*}$	$g^\sigma_{\Delta\Delta}$	$g^\sigma_{NN^*}$	1	1	1	$\sigma(x)$	1	1
ω	m_ω	$-\ g^\omega_{NN}$	$-\ g^\omega_{N^*N^*}$	$-\ g^\omega_{\Delta\Delta}$	$-\ g^\omega_{NN^*}$	γ_μ	1	1	$\omega^\mu(x)$	$-\ g^{\mu\nu}$	1
π	m_π	g^π_{NN}	$g^\pi_{N^*N^*}$	$g^\pi_{\Delta\Delta}$	$g^\pi_{NN^*}$	$\not{k}\gamma_5$	$\boldsymbol{\tau}$	$\mathbf{T}$	$\boldsymbol{\pi}$(x)	1	δ_{ij}

Table 1: Some symbols and notation used in this paper, k_μ is the transformed four-momentum.

In this contribution, we will derive the self-consistent RBUU equation for the $N^*(1440)$ distribution function within the framework we have done for the nucleon's and Δ's. Special attentions will be paid to the $N^*(1440)$-relevant in-medium cross sections. Through construction the collision term of $N^*(1440)$'s RBUU equation we will give the analytically expressions for calculating all the $N^*(1440)$-relevant in-medium two-body scattering cross sections within the N, Δ and $N^*(1440)$ system. The presented cross sections are consistent with the other ingredients of the transport model and can be used directly in the study of relativistic heavy-ion collisions.

II. RBUU-type transport equation for the N^*(1440) distribution function

The effective Lagrangian which considers the N, Δ and $N^*(1440)$ system interacting through σ, ω and π mesons can be written as

$$\begin{aligned}\mathcal{L}_I &= g^A_{NN}\bar{\psi}(x)\Gamma^N_A\psi(x)\Phi_A(x) + g^A_{N^*N^*}\bar{\psi}^*(x)\Gamma^{N^*}_A\psi^*(x)\Phi_A(x) + g^A_{\Delta\Delta}\bar{\psi}_{\Delta\nu}(x)\Gamma^\Delta_A\psi^\nu_\Delta(x)\Phi_A(x) \\ &+[g^A_{NN^*}\bar{\psi}^*(x)\Gamma^{N^*}_A\psi(x)\Phi_A(x) - g^\pi_{\Delta N}\bar{\psi}_{\Delta\mu}(x)\partial^\mu\boldsymbol{\pi}(x)\cdot\mathbf{S}^+\psi(x) \\ &-g^\pi_{\Delta N^*}\bar{\psi}_{\Delta\mu}(x)\partial^\mu\boldsymbol{\pi}(x)\cdot\mathbf{S}^+\psi^*(x) + h.c.] \qquad (1)\end{aligned}$$

where $\psi_{\Delta\mu}$ is the Rarita-Schwinger spinor of the Δ-baryon. $\Gamma^N_A = \Gamma^{N^*}_A = \gamma_A\tau_A$, $\Gamma^\Delta_A = \gamma_A T_A$, A=$\sigma$, ω, π; $g^\pi_{NN} = f_\pi/m_\pi$, $g^\pi_{\Delta N} = f^*/m_\pi$, the symbols and notation are defined in Table I.

In the language of the closed time-path Green's function technique the $N^*(1440)$ Green's function in the interaction picture can be defined in the same way as for nucleon's by

$$iG_{N^*}(1,2) =< T[exp(-i\oint dx H_I(x))\psi^*(1)\bar{\psi}^*(2)] > . \qquad (2)$$

Here T is the time ordering operator defined on a time contour. The corresponding Dyson equation for $iG_{N^*}(1,2)$ can be written as

$$iG_{N^*}(1,2) = iG^0_{N^*}(1,2) + \oint dx_3\oint dx_4 G^0_{N^*}(1,4)\Sigma_{N^*}(4,3)iG_{N^*}(3,2). \qquad (3)$$

Here $G^0_{N^*}(1,2)$ is the zeroth-order Green's function of $N^*(1440)$, which is similar to that of the nucleon's zeroth-order Green's function [5, 6]. The only difference between the nucleon and the $N^*(1440)$ is the mass and the coupling strengths! As in most/all presently used RBUU-type transport models, also here we do not take into account the temperature degree of freedom. Furthermore, in our theoretical framework the negative-energy states are neglected. $\Sigma_{N^*}(4,3)$ is the N^* self-energy. The lowest-order self-energies contributing to the collision term come from the Born diagrams. Through considering the N^* self-energy up to the Born approximation and adopting the semi-classical approximation and quasi-particle approximation the self-consistent RBUU equation for the $N^*(1440)$ can be derived in the same way as that of the nucleons. The RBUU equation for the $N^*(1440)$ distribution function reads

$$\{p_\mu[\partial_x^\mu - \partial_x^\mu \Sigma_{N^*}^\nu(x)\partial_\nu^p + \partial_x^\nu \Sigma_{N^*}^\mu(x)\partial_\nu^p] + m_{N^*}^* \partial_x^\nu \Sigma_{N^*}^S(x)\partial_\nu^p\}\frac{f_{N^*}(\mathbf{x},\mathbf{p},\tau)}{E_{N^*}^*(p)} = C^{N^*}(x,p), \tag{4}$$

where $f_{N^*}(\mathbf{x},\mathbf{p},\tau)$ is the single-particle distribution function of the $N^*(1440)$. The left-hand side of Eq. (4) is the transport part and the right-hand side is the collision term. Here we have dropped the contribution of the Fock term, since it usually has only small effects on the mean field. The above equation is derived within the framework as we have done for the nucleon's [6, 7]

$$\{p_\mu[\partial_x^\mu - \partial_x^\mu \Sigma^\nu(x)\partial_\nu^p + \partial_x^\nu \Sigma^\mu(x)\partial_\nu^p] + m^* \partial_x^\nu \Sigma^S(x)\partial_\nu^p\}\frac{f(\mathbf{x},\mathbf{p},\tau)}{E^*(p)} = C(x,p). \tag{5}$$

and delta's [5]

$$\{p_\mu[\partial_x^\mu - \partial_x^\mu \Sigma_{\Delta}^\nu(x)\partial_\nu^p + \partial_x^\nu \Sigma_{\Delta}^\mu(x)\partial_\nu^p] + m_{\Delta}^* \partial_x^\nu \Sigma_{\Delta}^S(x)\partial_\nu^p\}\frac{f_{\Delta}(\mathbf{x},\mathbf{p},\tau)}{E_{\Delta}^*(p)} = C^{\Delta}(x,p). \tag{6}$$

RBUU equations. Therefore, Eqs. (4), (5) and (6) stand in a consistent form and they are coupled together through the mean field and collision term(i.e., in-medium scattering cross sections of different channels). The $\Sigma_{N^*}^S(x)$ and $\Sigma_{N^*}^\mu(x)$ are the Hartree terms of the scalar and vector $N^*(1440)$ self-energies. After taking into account the self-interaction of the σ, ω fields the field equations of the σ and ω mesons can be written as

$$m_\sigma^2\sigma(x) + b(\mathrm{g}_{NN}^\sigma)^3\sigma^2(x) + c(\mathrm{g}_{NN}^\sigma)^4\sigma^3(x) = \mathrm{g}_{NN}^\sigma\rho_S(N) + \mathrm{g}_{N^*N^*}^\sigma\rho_S(N^*) + \mathrm{g}_{\Delta\Delta}^\sigma\rho_S(\Delta), \tag{7}$$

$$m_\omega^2\omega^\mu(x) + \frac{(\mathrm{g}_{NN}^\omega)^2 m_\omega^2}{Z^2}(\omega^\mu(x))^3 = \mathrm{g}_{NN}^\omega\rho_V^\mu(N) + \mathrm{g}_{N^*N^*}^\omega\rho_V^\mu(N^*) + \mathrm{g}_{\Delta\Delta}^\omega\rho_V^\mu(\Delta). \tag{8}$$

The effective four momentum and effective mass of the $N^*(1440)$ are defined as

$$m_{N^*}^*(x) = M_{N^*} - \mathrm{g}_{N^*N^*}^\sigma\sigma(x) \tag{9}$$

$$p^\mu(x) = P^\mu - \mathrm{g}_{N^*N^*}^\omega\omega^\mu(x). \tag{10}$$

Here ρ_S(i) and ρ_V^μ(i) are the scalar and vector densities of the nucleon, $N^*(1440)$ and delta:

$$\rho_S(i) = \frac{\gamma(i)}{(2\pi)^3}\int d\mathbf{q}\frac{m_i^*}{\sqrt{\mathbf{q}^2 + m_i^{*2}}} f_i(\mathbf{x},\mathbf{q},\tau), \tag{11}$$

$$\rho_V^\mu(i) = \frac{\gamma(i)}{(2\pi)^3}\int d\mathbf{q}\frac{q^\mu}{\sqrt{\mathbf{q}^2 + m_i^{*2}}} f_i(\mathbf{x},\mathbf{q},\tau). \tag{12}$$

The abbreviations i=N, N^*, Δ, and γ(i)= 4, 4, 16, correspond to nucleon, $N^*(1440)$ and delta, respectively. The effective four-momenta and effective masses of nucleon and delta can be defined through substituting the appropriate nucleon and delta labels in Eqs. (9) and (10), respectively.

The collision term can be expressed according to the transition probability, which reads as

$$\begin{aligned} C^{N^*}(x,p) &= \frac{1}{2}\int \frac{d^3p_2}{(2\pi)^3}\int \frac{d^3p_3}{(2\pi)^3}\int \frac{d^3p_4}{(2\pi)^3}(2\pi)^4\delta^{(4)}(p+p_2-p_3-p_4) \\ &\times W^{N^*}(p,\, p_2,\, p_3,\, p_4)(F_2-F_1), \end{aligned} \tag{13}$$

where F_2, F_1 are the Uehling-Uhlenbeck factors of the gain (F_2) and loss (F_1) terms, respectively:

$$F_2 = [1-f_{N^*}(\mathbf{x},\mathbf{p},\tau)][1-f_{B_2}(\mathbf{x},\mathbf{p}_2,\tau)]f_{B_3}(\mathbf{x},\mathbf{p}_3,\tau)f_{B_4}(\mathbf{x},\mathbf{p}_4,\tau), \tag{14}$$

$$F_1 = f_{N^*}(\mathbf{x},\mathbf{p},\tau)f_{B_2}(\mathbf{x},\mathbf{p}_2,\tau)[1-f_{B_3}(\mathbf{x},\mathbf{p}_3,\tau)][1-f_{B_4}(\mathbf{x},\mathbf{p}_4,\tau)], \tag{15}$$

B_2, B_3, B_4 can be N, Δ and $N^*(1440)$. $W^{N^*}(p,\, p_2,\, p_3,\, p_4)$ is the transition probability of different channels, which has the form

$$W^{N^*}(p,\, p_2,\, p_3\, ,p_4) = \frac{1}{16E^*_{N^*}(p)E^*_{B_2}(p_2)E^*_{B_3}(p_3)E^*_{B_4}(p_4)}\sum_{AB}(T_D\Phi_D - T_E\Phi_E) + p_3 \longleftrightarrow p_4. \tag{16}$$

Here T_D, T_E are the isospin matrices and Φ_D, Φ_E are the spin matrices, respectively. D denotes the contribution of the direct diagrams and E is that of the exchange diagrams. A, $B = \sigma$, ω, π represent the contributions of different mesons. The exchange of p_3 and p_4 is only for the case of identical particles in the final state. The two-body scattering reactions relevant to the $N^*(1440)$ in the N, Δ and $N^*(1440)$ system are follows:

(1) Elastic reactions:

$NN^* \longrightarrow NN^*$, $\Delta N^* \longrightarrow \Delta N^*$, $N^*N^* \longrightarrow N^*N^*$.

(2) Inelastic reactions:

$NN \longleftrightarrow NN^*$, $N\Delta \longleftrightarrow NN^*$, $\Delta\Delta \longleftrightarrow NN^*$,
$NN^* \longleftrightarrow \Delta N^*$, $NN^* \longleftrightarrow N^*N^*$, $NN \longleftrightarrow \Delta N^*$,
$N\Delta \longleftrightarrow \Delta N^*$, $\Delta\Delta \longleftrightarrow \Delta N^*$, $N^*N^* \longleftrightarrow \Delta N^*$,
$NN \longleftrightarrow N^*N^*$, $N\Delta \longleftrightarrow N^*N^*$, $\Delta\Delta \longleftrightarrow N^*N^*$.

For the inelastic case we only calculate the $N^*(1440)$-incident cross sections, its vice versa cross sections can be obtained by means of the detailed balance [8]. We have derived the analytical expressions for the above differential cross sections through calculating the Born term of the $N^*(1440)$ self-energies. By means of the relation between the transition probability $W^{N^*}(p, p_2, p_3, p_4)$ and the differential cross section [9], Eq. (13) can be rewritten as

$$C^{N^*}(x,p) = \frac{1}{2}\int \frac{d^3p_2}{(2\pi)^3}\upsilon\sigma_{N^*}(s,t)(F_2-F_1)d\,\Omega. \tag{17}$$

Here υ is the Møller velocity, $\sigma_{N^*}(s,t)$ is the differential cross section of different N^*-incident channels. The concrete expression of $\sigma_{N^*}(s,t)$ is given in Ref. [10].

III. The centroid $N^*(1440)$ mass, coupling strengths and form factors

In order to take into account the broad decay width of the $N^*(1440)$ resonance, we introduce a centroid $N^*(1440)$ mass $\langle M_{N^*}\rangle$ in the same way as we did for the delta [6, 7]. $\langle M_{N^*}\rangle$ is defined as

$$\langle M_{N^*}\rangle = \frac{\int_{M_N+m_\pi}^{\sqrt{S}-M_N} f(M_{N^*})M_{N^*} d\, M_{N^*}}{\int_{M_N+m_\pi}^{\sqrt{S}-M_N} f(M_{N^*}) d\, M_{N^*}}, \tag{18}$$

$f(M_{N^*})$ is the Breit-Wigner function

$$f(M_{N^*}) = \frac{1}{2\pi}\frac{\Gamma(q)}{(M_{N^*}-M_0)^2+\frac{1}{4}\Gamma^2(q)}, \tag{19}$$

here $M_0 = 1440$ MeV and $\Gamma(q)$ is the momentum-dependent decay width [12]

$$\Gamma(q) = \Gamma_0 \frac{M_0}{M_{N^*}}(q/q_0)^3 \frac{1.2}{1+0.2(\frac{q}{q_0})^2}, \tag{20}$$

q_0 is related to the case of $M_{N^*} = M_0$ and $\Gamma_0 = 200$ MeV. The effect of the decay width of $N^*(1440)$ is taken into account through replacing M_{N^*} in Eq. (9) with $\langle M_{N^*}\rangle$. The in-medium $N^*+N \to N+N$ and $N^*+N \to N^*+N$ cross sections can then be calculated by means of the equations

$$\sigma^*_{N^*N\to NN} = \frac{1}{16N}\int \sigma_{N^*N\to NN}(s,t) d\,\Omega, \tag{21}$$

$$\sigma^*_{N^*N\to N^*N} = \frac{1}{8}\int \sigma_{N^*N\to N^*N}(s,t) d\,\Omega, \tag{22}$$

here N is the normalization factor stemming from the decay width of the $N^*(1440)$ [8, 11, 7]. The in-medium $N^*(1440)$ production cross section can be obtained from Eq. (21) through detailed balance [8]

$$\sigma^*_{NN\to NN^*} = \frac{1}{8}\int \frac{p^2_{NN^*}}{p^2_{NN}} \sigma_{N^*N\to NN}(s,t) d\,\Omega, \tag{23}$$

where p_{NN}, p_{NN^*} denote the c. m. three momentum of the NN and NN^* states, respectively.

For the coupling strength of g^π_{NN}, we take the most commonly used value $f^2_\pi/4\pi = 0.08$. The coupling strengths of g^σ_{NN} and g^ω_{NN} are determined by fitting the known ground-state properties for infinite nuclear matter. For the coupling strengths of nucleon-$N^*(1440)$ coupling we follow the arguments of Ref. [13]. The following relation is expected to be valid

$$\frac{g^\pi_{NN^*}}{g^\pi_{NN}} = \frac{g^\sigma_{NN^*}}{g^\sigma_{NN}} = \frac{g^\omega_{NN^*}}{g^\omega_{NN}}. \tag{24}$$

$g^\pi_{NN^*}$ is determined from the width of pion decay of the $N^*(1440)$-resonance $g^\pi_{NN^*}/g^\pi_{NN} = 0.351$. For the N^*N^* coupling strengths, unfortunately, there is no information directly available from experiments. A similar situation exists for the $\Delta\Delta$ coupling strengths. Based on the quark model and mass splitting arguments several different choices for the delta coupling strengths have been discussed in Refs. [14, 15], which will cause strong influence on the nuclear equation of state in relativistic mean field calculations [15]. We assume that the arguments of Refs. [14, 15] apply to the N^*N^* coupling strengths and mainly consider the following three cases:

$$\alpha(N^*) = \frac{g^{\omega}_{N^*N^*}}{g^{\omega}_{NN}} = 1, \qquad \beta(N^*) = \frac{g^{\sigma}_{N^*N^*}}{g^{\sigma}_{NN}} = 1, \tag{25}$$

$$\alpha(N^*) = \beta(N^*) = \frac{M_{N^*}}{M_N} \approx 1.5. \tag{26}$$

and

$$\alpha(N^*) = 1, \qquad \beta(N^*) \approx 1.5 \tag{27}$$

The influence of different choices of $\alpha(N^*)$ and $\beta(N^*)$ on the predicted in-medium cross sections will be checked. For simplicity, an universal coupling strength of $g^{\pi}_{\Delta\Delta} = g^{\pi}_{N^*N^*} = g^{\pi}_{NN}$ is always assumed.

To take account of the effects stemming from the finite size of hadrons and a part of the short range correlation, a phenomenological form factor is introduced at each vertex. Here we distinguish the form factor Λ^*_A for the nucleon-N^*(1440)-meson vertex to the Λ_A for the nucleon-nucleon-meson vertex. S. Huber and J. Aichelin claimed that Λ^*_A is about 40% of Λ_A [13]. We adopt this argument in the numerical calculations. The form factor of the N^*(1440)-N^*(1440)-meson vertex is taken the same as that of corresponding nucleon-nucleon-meson vertex. The cut-off masses Λ_σ=1200 MeV, Λ_ω=808 MeV and Λ_π=500 MeV fixed in Refs. [6, 7] will be used, which are obtained by fitting the experimental data of nucleon mean free path and the free NN scattering cross section. According to the above argument, Λ^*_σ=480 Mev, Λ^*_ω=323 MeV. But we still take $\Lambda^*_\pi = \Lambda_\pi$=500 MeV, since this value is already comparable to the Λ^*_π=400 MeV used in Ref. [13].

IV. Numerical results and discussions

In Fig. 1 we compare our theoretical predications of free $pp \to pp^*(1440)$ cross section to the available experimental data [16]. The results of the one-boson-exchange model computed by Huber and Aichelin [13] are also presented in this figure as dashed line. Our results are consistent with that of Ref. [13]. Both of them are in good agreement with the experimental data. Here we should point out that our calculations are almost parameter free. We do not fit any parameters to the predicted cross section. Only the argument of $\Lambda^*_\sigma/\Lambda_\sigma = \Lambda^*_\omega/\Lambda_\omega$=40% is taken from Ref. [13]. If $\Lambda^*_\sigma = \Lambda_\sigma$ and $\Lambda^*_\omega = \Lambda_\omega$ are adopted, the cross section will be three times larger than the empirical value at higher energy as indicated by the dotted line in the figure.

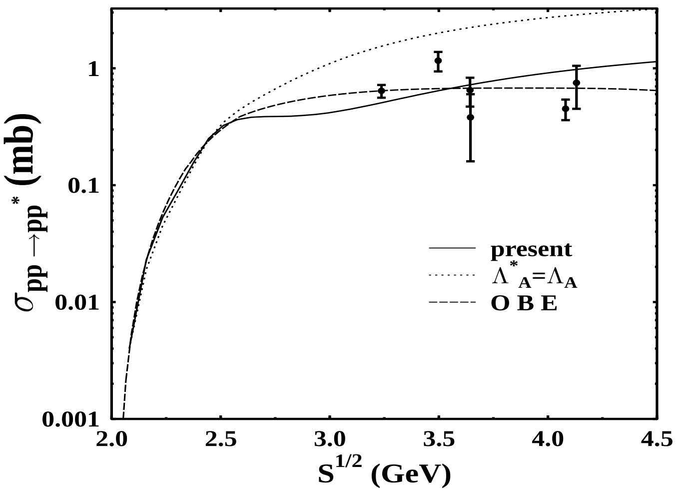

Figure 1: Free scattering cross section for reaction $pp \to pp^*(1440)$. Solid line represents the results of this work, and dashed line denotes the results of Ref. [13]. The experimental data are taken from Ref. [16]. The unitary form factor for NN and NN^* vertex, i.e., $\Lambda^*_A = \Lambda_A$ is also tested in the calculations, which is depicted by the dotted line.

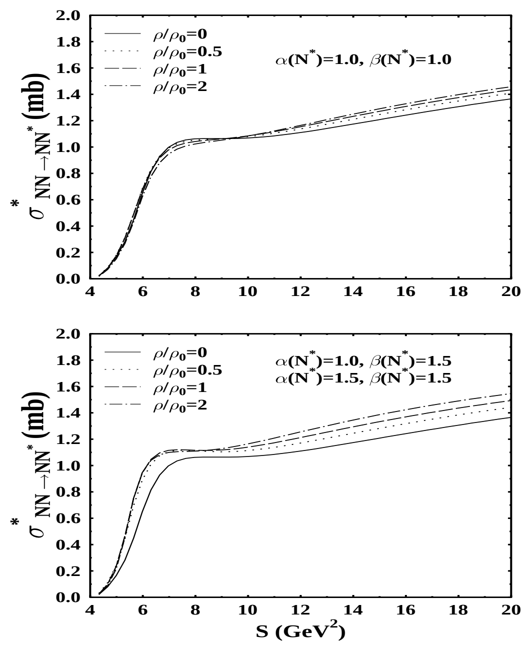

Fig. 2: The in-medium $NN \to NN^*$ cross section at different densities and energies. The calculations are performed with different sets of $\alpha(N^*)$ and $\beta(N^*)$.

Fig. 2 displays the in-medium $N^*(1440)$ production cross sections at different densities and energies. It is shown from the figure that the $\sigma^*_{NN\to NN^*}$ increases with the increase of density. When the universal coupling strengths are used, only a mild dependence on the density is exhibited. The density dependence, however, will become evident if one uses a larger scalar-$N^*(1440)$ coupling strength. The choice of the $\alpha(N^*)$ has no influence on the predicted cross sections. The reason is as follows: firstly, $g^{\omega}_{N^*N^*}$ does not enter the expressions of the $\sigma^*_{NN\to NN^*}$ explicitly; secondly, we always calculate the in-medium total energy of two particle system (small s) from the incident two particles, i.e., two nucleons in the case of the $\sigma^*_{NN\to NN^*}$. The situation will change if one considers the $\sigma^*_{N^*N\to NN}$, where the influence of $\alpha(N^*)$ will enter in the calculations of in-medium total energy (small s) from free total energy (capital S), and then affects the in-medium cross section.

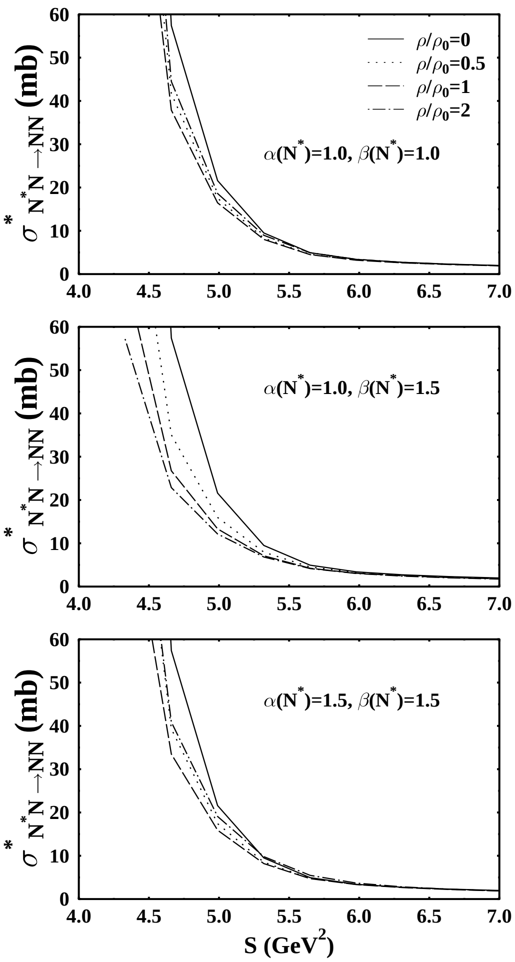

Fig. 3: The same as Fig. 2, but for an in-medium $N^*N \to NN$ cross section.

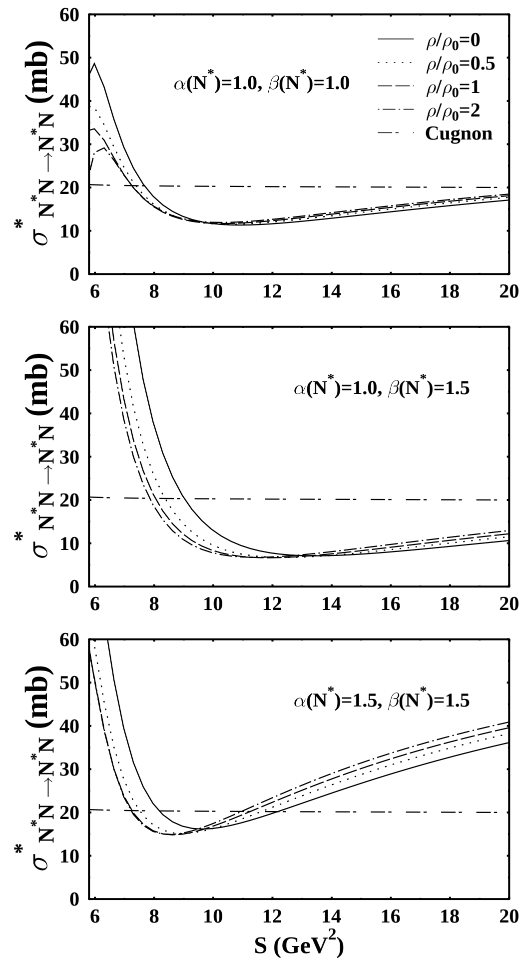

Fig. 4: The same as Fig. 2, but for an in-medium $N^*N \to N^*N$ cross section.

Fig. 3 depicts the in-medium $N^*(1440)$ absorption cross section. The cross sections drop very rapidly when the energy exceeds the threshold. That means that the absorption process are most important at energy close to the threshold as in the case of Δ absorption. When

$\alpha(N^*) = 1$, $\beta(N^*) = 1.5$ is used as the $N^*(1440)$ coupling strengths, the $\sigma^*_{N^*N\to NN}$ exhibits a clear density dependence. It decreases with the increase of density. In other two cases, i.e., $\alpha(N^*) = \beta(N^*) = 1$ and $\alpha(N^*) = \beta(N^*) = 1.5$, the dependence of the $\sigma^*_{N^*N\to NN}$ on the density becomes weaker and less explicit.

In Fig. 4 we show the in-medium $N^*N \to N^*N$ cross section at different densities and energies. As can be found from the figure, the cross sections now become very sensitive to the $\alpha(N^*)$ and $\beta(N^*)$ used because $g^\sigma_{N^*N^*}$ and $g^\omega_{N^*N^*}$ enter the expressions of the $\sigma^*_{N^*N\to N^*N}$ explicitly. Generally speaking, the density dependence of the cross section is not very evident when $\alpha(N^*) = \beta(N^*) = 1$ and $\alpha(N^*) = \beta(N^*) = 1.5$ are used, mainly due to the strong cancelation effects from the $\sigma + \omega$ mixed term. A strong density dependence appears when the set of $\alpha(N^*) = 1$, $\beta(N^*) = 1.5$ is used as the $N^*(1440)$ coupling strengths. The in-medium cross section decreases with the increase of density at lower energy and increases at higher energy. As the energy changes, the cross section firstly decreases and then increases, especially in the case of $\alpha(N^*)$=1.5. It is mainly caused by the contribution of the ω term. The ω term approaches a saturation with the increase of energy while all other terms (especially, the important cancelation term of the $\sigma + \omega$ mixed term) decrease. The Cugnon's parameterization [17] for free NN elastic cross section, which is commonly used in the transport models for the N^*N elastic cross section, is also plotted in Fig. 4. One can find an evident difference between the in-medium $N^* + N \to N^* + N$ cross section and the Cugnon's parameterization. It is therefore important to take the in-medium cross sections into account in the study of heavy-ion collisions.

V. Summary

Starting from the effective Lagrangian describing baryons interacting through mesons, using the closed time-path Green's function technique and adopting the semi-classical, quasi-particle and Born approximations we have developed a RBUU-type transport equation for the $N^*(1440)$ distribution function. The equation is derived within the same framework which was successfully applied to the nucleon [6] and delta [5] and thus we obtained a set of self-consistent equations for the N, Δ and $N^*(1440)$ system. Three equations are coupled through the self-energy terms and collision terms and should be solved simultaneously in a numerical simulation of heavy-ion collisions. Both the mean field and collision term of the $N^*(1440)$'s RBUU equation are derived from the same effective Lagrangian and given explicitly, so the medium effects on the two-body scattering cross sections are addressed automatically and can be studied self-consistently. Therefore, this approach provides a promising way to reach a covariant description of the $N^*(1440)$ in relativistic heavy-ion collisions.

Based on this approach, we have studied the in-medium two-body scattering cross sections. Since there is no information about the N^*N^* coupling strengths available, several different choices for $\alpha(N^*) = g^\omega_{N^*N^*}/g^\omega_{NN}$ and $\beta(N^*) = g^\sigma_{N^*N^*}/g^\sigma_{NN}$ are investigated. The results turn out to be sensitive to the $\alpha(N^*)$ and $\beta(N^*)$ used. Generally speaking, only a mild density-dependence of in-medium cross sections is found in the cases of $\alpha(N^*) = \beta(N^*) = 1$ and $\alpha(N^*) = \beta(N^*) = 1.5$. The situation, however, is changed when the set of $\alpha(N^*) = 1$, $\beta(N^*) = 1.5$ is adopted. An evident density-dependence appears. Qualitatively, the $\sigma^*_{NN\to NN^*}$ are found to increase with the increase of density while the $\sigma^*_{N^*N\to NN}$ near the threshold energy decreases. For the $\sigma^*_{N^*N\to N^*N}$, the situation is a little complicated. It decreases with the increase of density at lower energy and increases at higher energy. Because we have not included the screening and anti-screening effects of the medium on the interaction in the present calculations, the above arguments should be viewed with caution. Further investigations are needed.

References

[1] W. Ehehalt, et al., Phys. Rev. **C47**, R2467 (1993).

[2] M. Hofmann, et al., Phys. Rev. **C51**, 2095 (1995).

[3] V. Metag, Nucl. Phys. **A553**, 283c (1993).

[4] R. Averbeck, R. Holzmann, A. Schubert et al., GSI Sci. Rep., 80 (1994).

[5] G. Mao, Z. Li, Y. Zhuo, Phys. Rev. **C53**, 2933 (1996).

[6] G. Mao, Z. Li, Y. Zhuo, Y. Han, and Z. Yu, Phys. Rev. **C49**, 3137 (1994); G. Mao, Z. Li, Y. Zhuo, and Z. Yu, Phys. Lett. **B327**, 183 (1994).

[7] G. Mao, Z. Li, Y. Zhuo, and E. Zhao, Phys. Rev. **C55**, 792 (1997).

[8] P. Danielewicz and G. F. Bertsch, Nucl. Phys. **A533**, 712 (1991).

[9] S. R. de Groot, W. A. van Leeuwen, and Ch. G. van Weert, *Relativistic Kinetic Theory* (North-Holland, Amsterdam, 1980).

[10] G. Mao, L. Neise, H. Stöcker, W. Greiner, Z. Li, to appear in Phys. Rev. C.

[11] Gy. Wolf, W. Cassing, and U. Mosel, Nucl. Phys. **A545**, 139c (1992); Nucl. Phys. **A552**, 549 (1993).

[12] S. A. Bass, PhD thesis, Frankfurt university, unpublished.

[13] S. Huber and J. Aichelin, Nucl. Phys. **A573**, 587 (1994)

[14] S.A. Moszkowski, Phys. Rev. **D9**, 1613 (1974); S.I.A. Garpman, N.K. Glendenning and Y.J. Karant, Nucl. Phys. **A322**, 382 (1979).

[15] B.M. Waldhauser, J. Theis, J.A. Maruhn, H. Stöcker and W. Greiner, Phys. Rev. **C36**, 1019 (1987); P. Lévai, B. Lukács, B. Waldhauser and J. Zimányi, Phys. Lett. **B177**, 5 (1986).

[16] CERN-HERA Report 83-01.

[17] J. Cugnon, T. Mizutani and J. Vandermeulen, Nucl. Phys. **A352**, 505 (1981).