Dynamically generated resonances from the vector meson-octet baryon interaction in the strangeness zero sector

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Abstract

The interaction potentials between vector mesons and octet baryons are calculated explicitly with a summation of $t-$, $s-$, $u-$ channel diagrams and a contact term originating from the tensor interaction. Many resonances are generated dynamically in different channels of strangeness zero by solving the coupled-channel Lippman-Schwinger equations with the method of partial wave analysis, and their total angular momenta are determined. The spin partners $N(1650)1/2^-$ and $N(1700)3/2^-$, $N(1895)1/2^-$ and $N(1875)3/2^-$, and the state $N(2120)3/2^-$ are all produced respectively in the isospin $I = 1/2$ sector. In the isospin $I = 3/2$ sector, the spin partners $\Delta(1620)1/2^-$ and $\Delta(1700)3/2^-$ are also associated with the pole in the complex energy plane. According to the calculation results, a $J^P = 1/2^-$ state around 2000 MeV is predicted as the spin partner of $N(2120)3/2^-$. Some resonances are well fitted with their counterparts listed in the newest review of Particle Data Group(PDG)\cite{1}, while others might stimulate the experimental observation in these energy regions in the future.

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I. INTRODUCTION

The combination of chiral Lagrangian with nonperturbative unitary techniques in coupled channels of mesons and baryons has become a powerful method to study the meson-meson and meson-baryon interactions and new states in the resonance region, which are not easily explained using the conventional constituent quark model. The analysis of the meson baryon scattering matrices shows poles in the second Riemann sheet that can be associated to known resonances or new ones. In this way, $J^P = 1/2^-$ resonances have been generated dynamically in the interaction of pseudoscalar mesons with octet baryons, which fit quite well the spectrum of the known low-lying resonances with these quantum numbers $^2[10]$. Similarly, the $J^P = 3/2^-$ resonances are obtained in the interaction of pseudoscalar mesons with decuplet baryons $^{11, 12}$.

The theory on the hidden gauge symmetry supplies a mechanism to include vector mesons in the chiral Lagrangian $^{13-16}$. Therefore, the study on the interaction between vector mesons and other hadrons becomes possible. Along this clue, the interactions between vector mesons and pseudoscalar mesons are studied and the radiative decays of some axial-vector mesons generated dynamically are discussed $^{17}$. Similarly, the scattering amplitudes of two $\rho$ mesons are calculated within the framework of coupled-channel Lippman-Schwinger equations, and two resonances $f_0(1370)$ and $f_2(1270)$ are generated dynamically $^{18}$. This work has been extended to the $SU(3)$-space of vector mesons in Ref. $^{19}$, where several known resonances are also dynamically generated. In the baryon sector, the interaction of vector mesons and decuplet baryons is addressed in Refs. $^{20, 21}$, where only $t-$ channel amplitudes are analyzed in the S-wave approximation. This method is also extended to study the interaction of vector mesons and octet baryons, and several baryon resonances have been found as a result of solving the coupled-channel Lippman-Schwinger equations $^{22-24}$. However, the resonances generated dynamically are spin degenerate since the amplitude obtained from $t-$ channel interaction is independent of spin. Because the masses of vector mesons are comparable to those of baryons, only $t-$ channel diagrams might be incomplete to obtain a reliable interaction of vector mesons and baryons. Thus in Ref. $^{25}$, the $t-, s-, u-$ channel diagrams and a contact diagram originating from the tensor term of the vector meson-octet baryon interaction are all taken into account, and four spin-determined resonances are found in a non-relativistic approximation of the coupled-channel Lippman-Schwinger equations. In the
present work, we deduce the interaction kernel of vector mesons and octet baryons including a vector meson exchange in $t$-channel, octet baryon exchange in $s-$, $u-$ channels, and a contact diagram related only to the tensor interaction term in a fully relativistic framework, and then calculate the scattering amplitudes of vector mesons and octet baryons by solving the coupled-channel Lippman-Schwinger equations. The amplitude of the vector mesons and octet baryons will be expanded in terms of partial waves, and then the poles of the amplitudes in different partial waves are detected in the complex energy plane in center of mass system, which can be associated to some well-known resonances.

In Sect. II, we will show the basic Lagrangian obtained with the hidden gauge symmetry of $SU(3)$ group, where a tensor interaction term is included. Then the framework on the coupled-channel Lippman-Schwinger equations will be summarized briefly. In Sect. III the parameters in the Lagrangian which are fitted with the experimental data on coupling constants of octet baryons to vector mesons are determined. In Sect. IV the amplitudes are expanded in terms of partial waves, and the formula on the partial wave analysis is displayed. In Sect. V we will analyze the resonances found in the complex energy plane for the vector meson-octet baryon system with total strangeness zero, and the properties of these resonances and their possible PDG counterparts are discussed. Finally, we will present a summary on this article.

II. FORMALISM

We follow the formalism of the hidden gauge interaction for vector mesons of [13–16] in this manuscript. The Lagrangian involving the interaction of vector mesons among themselves is given by

$$\mathcal{L}_V = -\frac{1}{2}\langle V^{\mu\nu}V_{\mu\nu}\rangle, \quad (1)$$

where the symbol $\langle \rangle$ stands for the trace in the $SU(3)$ space and the tensor field of vector mesons is given by

$$V^{\mu\nu} = \partial^\mu V^\nu - \partial^\nu V^\mu - ig[V^\mu, V^\nu], \quad (2)$$

where $g$ is

$$g = \frac{M_V}{\sqrt{2}f_\pi}, \quad (3)$$
with $f_\pi = 93\text{MeV}$ the pion decay constant and $M_V$ the mass of the $\rho$ meson. The magnitude $V_\mu$ is defined by the matrix

$$V_\mu = \frac{1}{2} \begin{pmatrix} \rho^0 + \omega & \sqrt{2} \rho^+ & \sqrt{2} K^*+ \\ \sqrt{2} \rho^- & -\rho^0 + \omega & \sqrt{2} K^*0 \\ \sqrt{2} K^*0 & \sqrt{2} K^*+ & \sqrt{2} \phi \end{pmatrix}_\mu.$$  

The interaction of $\mathcal{L}_V$ gives rise to a three-vector vertex form

$$\mathcal{L}_{(3V)} = i2g\langle(\partial_\mu V_\nu - \partial_\nu V_\mu)V^\mu V^\nu\rangle,$$

which will contribute to the $t-$ channel of the vector meson-octet baryon interaction.

The tensor interaction term can be included in the vector meson-octet baryon Lagrangian as is done in Ref. [25]. It is no doubt that the tensor term also satisfies the $SU(3)$ hidden gauge symmetry and might make an amendment to the vector meson-octet baryon interaction. Therefore, the lagrangian for the vector meson-octet baryon interaction due to the $SU(3)$ hidden gauge symmetry can be written as

$$\mathcal{L}_{VB} = -g \left\{ F_V\langle \bar{B}\gamma_\mu [V^\mu, B]\rangle + D_V\langle \bar{B}\gamma_\mu \{V^\mu, B\}\rangle + \langle \bar{B}\gamma_\mu B\rangle\langle V^\mu\rangle \\
+ \frac{1}{4M} \left( F_T\langle \bar{B}\sigma_{\mu\nu} [V^{\mu\nu}, B]\rangle + D_T\langle \bar{B}\sigma_{\mu\nu} \{V^{\mu\nu}, B\}\rangle \right) \right\},$$

where $B$ is the $SU(3)$ matrix of octet baryons

$$B = \begin{pmatrix} \frac{1}{\sqrt{6}}\Lambda + \frac{1}{\sqrt{2}}\Sigma^0 & \Sigma^+ & p \\
\Sigma^- & \frac{1}{\sqrt{6}}\Lambda - \frac{1}{\sqrt{2}}\Sigma^0 & n \\
\Xi^- & \Xi^0 & -\sqrt{\frac{2}{3}}\Lambda \end{pmatrix},$$

and $M$ is the mass of the nucleon.

In order to obtain the correct couplings to the physical $\omega$ and $\phi$ meson, the mixing of their octet and singlet components must be considered. Under the ideal mixing assumption,
we have
\[ \omega = \sqrt{1/3}\omega_8 + \sqrt{2/3}\omega_0, \]
\[ \phi = -\sqrt{2/3}\omega_8 + \sqrt{1/3}\omega_0, \]
and only the octet parts of these wave functions are included in Eq. (6). For the singlet states we have
\[ \mathcal{L}_{V_0BB} = -g \left\{ \frac{C_V}{3} \langle \bar{B}\gamma_\mu B \rangle \langle V_0^\mu \rangle + \frac{C_T}{4M} \langle \bar{B}\sigma_{\mu\nu}V_0^{\mu\nu}B \rangle \right\}. \]
Thus, the Lagrangian for the linearly coupling to vector mesons can be explicitly written as
\[ \mathcal{L}_{VBB} = -g \left\{ F_V \langle \bar{B}\gamma_\mu [V_8^\mu, B] \rangle + D_V \langle \bar{B}\gamma_\mu \{ V_8^\mu, B \} \rangle \right\} + \frac{1}{4M} \left( F_T \langle \bar{B}\sigma_{\mu\nu} [\partial^\mu V_8^\nu - \partial^\nu V_8^\mu, B] \rangle + D_T \langle \bar{B}\sigma_{\mu\nu} \{ \partial^\mu V_8^\nu - \partial^\nu V_8^\mu, B \} \rangle \right) + \frac{C_V}{3} \langle \bar{B}\gamma_\mu B \rangle \langle V_0^\mu \rangle + \frac{C_T}{4M} \langle \bar{B}\sigma_{\mu\nu}V_0^{\mu\nu}B \rangle \right\}, \]
which will contribute to the \( t-, s-, u- \) channel interactions between vector mesons and octet baryons. Moreover, the self-coupling terms of vector mesons in Eq. (2) lead to a contact interaction between vector mesons and octet baryons, which is trivially null for the singlet vector meson, thus
\[ \mathcal{L}_{V_{VBB}} = \frac{g}{4M} \left\{ F_T \langle \bar{B}\sigma_{\mu\nu} [ig [V_8^\mu, V_8^\nu], B] \rangle + D_T \langle \bar{B}\sigma_{\mu\nu} \{ ig [V_8^\mu, V_8^\nu], B \} \rangle \right\}. \]
From Eqs. (5), (10) and (11), we can obtain the potentials for the \( t-, s-, u- \) channel and contact interactions between vector mesons and octet baryons. If the momentum of the initial vector meson is similar to that of the final vector meson, i.e., \( q_2 \approx q_1 \), the momentum transfer \( k = q_2 - q_1 \) is trivial null approximately, and then the \( t- \) channel interaction can be written as
\[ V_{ij}^t = -\frac{g}{\mu^2} \left( \bar{U}(p_2, \lambda_2)\Gamma_\mu(p_2, p_1)U(p_1, \lambda_1)(q_1^\mu + q_2^\mu) \right) \varepsilon(q_1, \delta_1) \cdot \varepsilon^*(q_2, \delta_2), \]
where
\[ \Gamma^\mu(p_2, p_1) = g_1 \gamma^\mu + g_2(p_2^\mu + p_1^\mu) \]
is the vertex of two baryons and a vector meson when the tensor term is taken into account in the Lagrangian of Eq. (6), and the coupling constants \( g_1 \) and \( g_2 \) for different octet baryons
and vector mesons are attached in Appendix I. In the above equation, \( U(p_1, \lambda_1) \) and \( U(p_2, \lambda_2) \) are the wave functions of the incoming and outgoing baryons, and \( \varepsilon(q_1, \delta_1) \) and \( \varepsilon^*(q_2, \delta_2) \) are polarization vectors of the initial and final mesons, respectively [20]. The formulas of them can be found in Appendix II. However, if the difference between \( q_2 \) and \( q_1 \) is taken into account, an additional part in Eq. (14) must be supplemented in the \( t \)-channel interaction of vector mesons and octet baryons.

\[
V_{supp,ij}^t = \frac{2g}{\mu^2} \left( \{ \hat{U}(p_2, \lambda_2) \Gamma^\mu(p_2, p_1) \varepsilon^\ast\mu(q_2, \delta_2) U(p_1, \lambda_1) \} q_2 \cdot \varepsilon(q_1, \delta_1) \right) + \left\{ \hat{U}(p_2, \lambda_2) \Gamma^\mu(p_2, p_1) \varepsilon^\ast\mu(q_1, \delta_1) U(p_1, \lambda_1) \right\} q_1 \cdot \varepsilon^*(q_2, \delta_2) \right) .
\]

In addition to the \( t \)-channel mechanism, the \( u \)-channel and \( s \)-channel mechanisms depicted in Fig. 1 are also considered in this work. The \( s \)-channel interaction of vector mesons and octet baryons can be written as

\[
V_{ij}^s = \hat{U}(p_2, \lambda_2) \Gamma^\mu(p_2, p_1 + q_1) \varepsilon^\ast\mu(q_2, \delta_2) \left( \frac{\hat{p}_1 + \hat{p}_2 + M}{s - M^2} \right) \Gamma^\nu(p_1 + q_1, p_1) \varepsilon_\nu(q_1, \delta_1) U(p_1, \lambda_1) \quad (15)
\]

with \( s = (p_1 + q_1)^2 \), while the \( u \)-channel interaction is

\[
V_{ij}^u = \hat{U}(p_2, \lambda_2) \Gamma^\mu(p_2, p_1 - q_2) \varepsilon_\mu(q_1, \delta_1) \left( \frac{\hat{p}_1 - \hat{p}_2 + M}{(p_1 - q_2)^2 - M^2} \right) \Gamma^\nu(p_1, p_1 - q_2) \varepsilon^\ast\nu(q_2, \delta_2) U(p_1, \lambda_1) .
\]

(16)

Since the three-momenta of vector mesons and octet baryons are far smaller than their masses in the concerned energy region, we make an approximation of \( (p_1 - q_2)^2 \approx M_1^2 + m_2^2 - 2M_1 m_2 \) in the propagator in Eq. (16), where \( M_1 \) and \( m_2 \) are the masses of initial octet baryons and final vector mesons, respectively.

From Eq. (11), the contact interaction of vector mesons and octet baryons is obtained

\[
V_{ij}^{CT} = -iC_{IS}^{CT} 2\hat{U}(p_2, \lambda_2) \varepsilon^\ast\mu(q_2, \delta_2) \sigma^{\mu\nu} \varepsilon_\nu(q_1, \delta_1) U(p_1, \lambda_1) .
\]

(17)

Altogether, the total kernel of the vector meson-octet baryon interaction can be written as

\[
V_{ij}(s,t) = V_{ij}^t + V_{supp,ij}^t + V_{ij}^s + V_{ij}^u + V_{ij}^{CT} ,
\]

(18)

which is a summation of the \( t \)-, \( s \)-, \( u \)- channels and contact interaction. Now the kernel in Eq. (18) is a function of the total energy in the center of mass system \( \sqrt{s} \) and the scattering angle \( \theta \).

The coefficients in Eqs. (12), (14), (15), (16) and (17) can be obtained both in the physical basis of states or in the isospin basis, and this is discussed in Appendix I. In this article,
we will directly study the interaction in the isospin basis. Especially, we will concentrate on the states of isospin $I = \frac{1}{2}$ and $\frac{3}{2}$ with strangeness zero.

The scattering matrix implies solving the coupled-channel Lippman-Schwinger equations in the on-shell factorization approach of [5, 6]

$$T(\sqrt{s}, \cos \theta) = [1 - V(\sqrt{s}, \cos \theta) G(s)]^{-1} V(\sqrt{s}, \cos \theta)$$  \hspace{1cm} (19)

with $G(s)$ being the loop function of a vector meson and a baryon which we calculate in dimensional regularization using the formula of [6]

$$G_l(s) = \frac{i2M_l}{16\pi^2} \int \frac{d^4q}{(2\pi)^4} \frac{1}{(P - q)^2 - M_l^2 + i\epsilon} q^2 - m_l^2 + i\epsilon$$

$$= \frac{2M_l}{16\pi^2} \left\{ a_l(\mu) + \ln \frac{M_l^2}{\mu^2} + \frac{m_l^2 - M_l^2 + s}{2s} \ln \frac{m_l^2}{M_l^2} + \frac{\bar{q}_l}{\sqrt{s}} \left[ \ln(s - (M_l^2 - m_l^2) + 2\bar{q}_l\sqrt{s}) + \ln(s + (M_l^2 - m_l^2) + 2\bar{q}_l\sqrt{s}) - \ln(-s + (M_l^2 - m_l^2) + 2\bar{q}_l\sqrt{s}) - \ln(-s - (M_l^2 - m_l^2) + 2\bar{q}_l\sqrt{s}) \right] \right\}$$  \hspace{1cm} (20)

with $\mu$ a regularization scale, which we take to be 630 MeV, and with a natural value of the subtraction constant $a_l(\mu)$ of $-2$, as determined in [6, 22].
In eq. (20), $\bar{q}_l$ denotes the three-momentum of the vector meson or the octet baryon in the center of mass frame and is given by

$$\bar{q}_l = \frac{\lambda^{1/2}(s, m_l^2, M_l^2)}{2\sqrt{s}} = \sqrt{s - (M_l + m_l)^2} \sqrt{s - (M_l - m_l)^2},$$

(21)

where $\lambda$ is the triangular function and $M_l$ and $m_l$ are the masses of octet baryons and vector mesons, respectively.

In order to find the pole corresponding to the resonance, we must extend our calculation of the $T$ matrix to the complex plane of $\sqrt{s}$. Because the physical G-propagators in Eq. (20) have cut-lines in the real axis, they are valid only in the first Riemann sheet. However, the poles corresponding resonances can only be found in the second Riemann sheet, and the G-propagators in the second Riemann sheet must be studied.

The G-propagator near the real axis in the second Riemann sheet is defined as

$$G^{II}_l(s) = \begin{cases} G_l(s) & \text{at } \text{Re}[\sqrt{s}] \leq \sqrt{s^0_l} \\ G_l(s) - 2i\text{Im}[G_l(s)] & \text{at } \text{Re}[\sqrt{s}] > \sqrt{s^0_l} \end{cases},$$

(22)

with $\sqrt{s^0_l} = M_l + m_l$ the $l$-th channel threshold energy.

According to Eq. (20), the imaginary part of the G-propagator in the dimensional regularization scheme can be written as

$$\text{Im}[G^I_l(s)] = \frac{M_l q_l}{4\pi \sqrt{s}}.$$

(23)

Since the widths of the $K^*$ and $\rho$ mesons are considerably large, their mass distributions must be taken into account. We follow the traditional method of convoluting the G-function with the mass distributions of the $K^*$ and $\rho$ mesons, as is used in Ref. [22].

$$\tilde{G}(s) = \frac{1}{N} \int_{(m-2\Gamma_i)^2}^{(m+2\Gamma_i)^2} d\tilde{m}^2 \left( -\frac{1}{\pi} \right) \text{Im} \frac{1}{\tilde{m}^2 - m^2 + i\tilde{m}\Gamma(\tilde{m})} G(s, \tilde{m}^2, \tilde{M}_B^2),$$

(24)

where the normalization factor is written as

$$N = \int_{(m_\rho-2\Gamma_i)^2}^{(m_\rho+2\Gamma_i)^2} d\tilde{m}^2 \left( -\frac{1}{\pi} \right) \text{Im} \frac{1}{\tilde{m}^2 - m_\rho^2 + i\tilde{m}\Gamma(\tilde{m})}$$

(25)

with

$$\Gamma(\tilde{m}) = \Gamma_i \left( \frac{m^2}{\tilde{m}^2} \right) \left( \frac{\lambda^{1/2}(\tilde{m}^2, m_d^2, m_d'^2)/2\tilde{m}}{\lambda^{1/2}(m^2, m_d^2, m_d'^2)/2m} \right)^3 \theta(\tilde{m} - m_d - m_d'),$$

(26)

and $\Gamma_i$ the decay width of the meson ($i = \rho, K^*$), which we take to be 149.1 MeV and 49.75 MeV for the $\rho$ and $K^*$ meson, respectively. In Eq. (26), $m_d, m_d'$ denote the masses of the
decay products of the vector mesons, i.e., pion masses in case of $\rho$ for the decay mode of $\rho \to \pi\pi$, and kaon and pion masses in case of $K^*$ for $K^* \to K\pi$.

As far as the mass distributions of the vector mesons are concerned, the $G$-function in the coupled-channel Lippman-Schwinger equations (19) should be replaced by $\tilde{G}(s)$ in Eq. (24).

III. PARAMETERS

| Exp. | $\omega$ | $\rho$ | $K^*$ | $\phi$ |
|------|---------|--------|-------|--------|
| $g^2/4\pi$ | 2.4     | 2.4    | 1.39  | 12.0   |

TABLE I: The experimental values on the coupling constants of vector mesons to octet baryons [27].

| $F_V$ | $D_V$ | $F_T$ | $D_T$ | $C_V$ | $C_T$ |
|-------|-------|-------|-------|-------|-------|
| 1.6405 | 0.2225 | 1.6405 | 0.2225 | -5.144 | -5.144 |

TABLE II: The parameters used in the calculation.

The experimental values on coupling constants of vector mesons to octet baryons are listed in Table I, which are taken from Ref. [27]. Thus we can fit the parameters $F_V$, $D_V$ and $C_V$ in the Lagrangian with these data. There is no doubt that the tensor terms are related to the form factors of baryons, which can be deduced from the linear couplings of vector mesons to octet baryons. However, the contribution from tensor terms to the kernels at tree diagram levels might increase the precision of the calculation. Therefore, we set the parameters $F_T$, $D_T$ and $C_T$ related to the tensor terms to be equal to $F_V$, $D_V$ and $C_V$, respectively. The values of these parameters are listed in Table II.

IV. PARTIAL WAVE ANALYSIS

If we set $l$, $j$ and $\mu$ the orbital angular momentum, the total angular momentum and the orientation of total angular momentum of the initial vector meson, $l'$, $j'$ and $\mu'$ those of the final vector meson, and $J$ the total angular momentum of the system, the scattering
amplitudes of vector mesons and octet baryons can be expanded in terms of partial waves. We have

\[< \rho', \nu'|T|\rho, \nu > = \sum_{j', j, \mu, \mu'} < \hat{q}_2; 1, 1/2; \rho', \nu'|\hat{q}_1; 1, 1/2; \rho, \nu > \]

\[< j', 1/2; J, M|T|j, 1/2; J, M > < j, 1/2; J, M|j, 1/2; \mu, \nu > < j', 1/2; \rho, \nu |\hat{q}_1; 1, 1/2; \rho, \nu >, \]

with

\[< \hat{q}_2; 1, 1/2; \rho', \nu'|j', 1/2; \mu', \nu' > = \sum_{\nu'} Y_{l', \mu' - \rho'}(\hat{q}_2) < l', 1; \mu' - \rho', \rho'|j', \mu' > \]

and

\[< j, 1/2; \mu, \nu|\hat{q}_1; 1, 1/2; \rho, \nu > = \sum_{l} < l, 1; \mu - \rho, \rho|j, \mu > Y_{l, \mu - \rho}^*(\hat{q}_1). \]

In Eq. (27), \(\rho\) and \(\rho'\) denote the orientations of spins for the initial and final vector mesons, and \(\nu\) and \(\nu'\) the orientations of spins for the initial and final baryons, respectively.

The spherical harmonics \(Y_{l,\mu - \rho}^*(\hat{q}_1)\) may be simplified by choosing the axis of quantization along the momentum of the initial vector meson \(\hat{q}_1\), and then \(Y_{l,\mu - \rho}^*(\hat{q}_1) = \delta_{\mu, \rho}[2l + 1]/4\pi^{1/2}\). In the S-wave approximation of partial wave analysis, only the final state with the orbital angular momentum \(l' = 0\) is studied. Since the parity is conserved, the contribution from the initial state with the orbital angular momentum \(l = 0, 2\) must be taken into account. Thus the scattering amplitudes can be expanded in terms of the total angular momentum \(J\) of the system, the orbital angular momentum \(l\) and the total angular
momentum $j$ of the initial vector meson.

\[< \rho', \nu'|T|\rho, \nu>\]

\[= \sum_{l,j,J,M} < 1, 1/2; \rho', \nu'|J, M > T^J_{l,j,0,1} < J, M | j, 1/2, \rho, \nu > < l, 1; 0, \rho | j, \rho > \left( \frac{2l + 1}{4\pi} \right)^{1/2}\]

\[= < 1, 1/2; \rho', \nu'|1/2, \rho' + \nu' > T^{J=1/2}_{1,1,0,0} < 1, 1/2; \rho, \nu|1/2, \rho + \nu > \left( \frac{1}{4\pi} \right)^{1/2}\]

\[+ < 1, 1/2; \rho', \nu'|3/2, \rho' + \nu' > T^{J=3/2}_{1,1,0,0} < 1, 1/2; \rho, \nu|3/2, \rho + \nu > \left( \frac{1}{4\pi} \right)^{1/2}\]

\[+ < 1, 1/2; \rho', \nu'|2, 1/2; \rho' + \nu' > T^{J=1/2}_{1,1,0,2} < 1, 2; \rho, \nu|1/2, \rho + \nu > < 2, 1; 0, \rho | 1, \rho > \left( \frac{5}{4\pi} \right)^{1/2}\]

\[+ < 1, 1/2; \rho', \nu'|3/2, \rho' + \nu' > T^{J=3/2}_{1,1,0,2} < 1, 2; \rho, \nu|3/2, \rho + \nu > < 2, 1; 0, \rho | 2, \rho > \left( \frac{5}{4\pi} \right)^{1/2}\]

\[+ < 1, 1/2; \rho', \nu'|2, 1/2; \rho' + \nu' > T^{J=3/2}_{1,2,0,2} < 2, 1; \rho, \nu|3/2, \rho + \nu > < 2, 1; 0, \rho | 2, \rho > \left( \frac{5}{4\pi} \right)^{1/2}\]

with $< ...|... >$ Clebsch-Gordan coefficients.

In Eq. (30), the first and second terms are related to the S-wave part of the initial vector meson, i.e., the case of $l = 0$, while the other terms correspond to the contributions from the case of $l = 2$, the D-wave part of the initial vector meson.

It is apparent that five amplitudes with different spin states of the initial and final vector mesons and octet baryons need to be calculated in order to obtain the values of amplitudes $T^{J=1/2}_{1,1,0,0}$, $T^{J=3/2}_{1,1,0,0}$, $T^{J=1/2}_{1,1,0,2}$, $T^{J=3/2}_{1,1,0,2}$, and $T^{J=3/2}_{1,2,0,2}$. We choose the amplitudes $< 1, 1/2|T|1, 1/2 >$, $< 0, 1/2|T|0, 1/2 >$, $<- 1, 1/2|T|1, 1/2 >$, $< -1, -1/2|T|0, 1/2 >$, and $< 0, -1/2|T|1, 1/2 >$ to obtain these values when the spin symmetry is taken into account.

We search for poles of these amplitudes at the complex plane of $\sqrt{s}$. If the $\text{Re}(\sqrt{s})$ is above the threshold of the channel, the pole might correspond to a resonance state of vector mesons and octet baryons. Otherwise, if $\text{Re}(\sqrt{s})$ of the pole is less than the threshold, it is more possible to be a bound state.

When a pole of the amplitude $T^{J}_{j,j',l',l}$ is produced in the complex plane of $\sqrt{s}$, not only the mass, the decay width, the parity and the total angular momentum $J$ of the corresponding resonance are determined, but the detailed information on the orbital and total angular momenta $l, j$ and $l', j'$ of the initial and final vector mesons to generate this resonance can also be obtained.
V. RESULTS

In this section we show our results obtained in the channels of strangeness zero and different isospins, respectively.

A. Isospin I=1/2

\begin{figure}
\centering
\includegraphics[width=0.5\textwidth]{fig2}
\caption{$|T|^2$ for different channels with $I = 1/2$, $J = 1/2$, $j' = 1$, $j = 1$, $l' = 0$ and $l = 0$.}
\end{figure}

\begin{figure}
\centering
\includegraphics[width=0.5\textwidth]{fig3}
\caption{$|T|^2$ for different channels with $I = 1/2$, $J = 1/2$, $j' = 1$, $j = 1$, $l' = 0$ and $l = 2$.}
\end{figure}

In the channel of Isospin $I = 1/2$, the isospin states for $\rho N$, $\omega N$, $\phi N$, $K^*\Lambda$, $K^*\Sigma$ can be written as

\[ |\rho N; \frac{1}{2}, \frac{1}{2} \rangle = -\sqrt{\frac{1}{3}}|\rho_0 p\rangle - \sqrt{\frac{2}{3}}|\rho_+ n\rangle, \]  

(31)
FIG. 4: $|T|^2$ for different channels with $I = 1/2$, $J = 3/2$, $j' = 1$, $j = 1$, $l' = 0$ and $l = 0$.

FIG. 5: $|T|^2$ for different channels with $I = 1/2$, $J = 3/2$, $j' = 1$, $j = 1$, $l' = 0$ and $l = 2$.

$$|\omega N; \frac{1}{2}, \frac{1}{2}\rangle = |\omega p\rangle, \quad (32)$$

$$|\phi N; \frac{1}{2}, \frac{1}{2}\rangle = |\phi p\rangle, \quad (33)$$

$$|K^* \Lambda; \frac{1}{2}, \frac{1}{2}\rangle = |K^{*+} \Lambda\rangle, \quad (34)$$

and

$$|K^* \Sigma; \frac{1}{2}, \frac{1}{2}\rangle = \sqrt{\frac{1}{3}}|K^{*+} \Sigma^0\rangle + \sqrt{\frac{2}{3}}|K^{*0} \Sigma^+\rangle, \quad (35)$$
respectively. We have used the phase convention $\rho^+ = -|1,1>$ and $\Sigma^+ = -|1,1>$ for the isospin states in Eqs. ($31$) and ($35$), which is consistent with the structure of the $V^\mu$ and $B$ matrices. It is apparent that the interaction between isospin states with isospin orientation $I_z = -1/2$ would generate the same resonances, so only the isospin states with $I_z = 1/2$ are discussed in this section.

In Fig. 2 we show the results of $|T_{ii}|^2$ as a function of $\sqrt{s}$ for the different channels in the $(I, J, j', j, l', l) = (1/2, 1/2, 1, 1, 0, 0)$ sector. We can see two peaks for $\phi N \rightarrow \phi N$ around 1700 MeV, a few MeV above the $\rho N$ threshold. These two peaks around 1700 MeV are also seen in the $K^*\Lambda$ channel but are absent or barely visible in the $\rho N$, $\omega N$ and $K^*\Sigma$ channels. On the other hand, with the same quantum numbers $(I, J, j', j, l', l)$, one finds other two peaks around 1860 MeV, which are clearly visible in the $K^*\Lambda$, $K^*\Sigma$ and $\phi N$ channels but not visible in the $\rho N$ and $\omega N$ ones. Moreover, there are also two peaks appeared around 2000 MeV in the $K^*\Sigma$ channel, which are not visible in the other channels. $|T_{ii}|^2$ as a function of $\sqrt{s}$ for the different channels with $(I, J, j', j, l', l) = (1/2, 1/2, 1, 1, 0, 2)$ is depicted in Fig. 3. The peaks appear almost at the same positions in the complex plane of $\sqrt{s}$ as those in Fig. 2.

The couplings of these resonances to vector mesons and octet baryons are different when the quantum numbers $(j', j, l', l)$ take different values. In follows, we only calculate their coupling constants to different vector mesons and octet baryons in the channel of $l = l' = 0$, i.e., $j = j' = 1$. The couplings of the resonances to different channels for the $(I, J, j', j, l', l) = (1/2, 1/2, 1, 1, 0, 0)$ sector, obtained from the residues at the poles are shown in Table III.
It is apparent that the poles at $1715 + i4$ MeV and $1728 + i0$ MeV couple strongly to the $\phi N$ channel, while the poles at $1855 + i1$ MeV and $1868 + i6$ MeV couple strongly to the $K^*\Lambda$ channels. For the two poles around 2000 MeV, at the positions of $1982 + i4$ MeV and $1999 + i5$ MeV, mainly interact with the $K^*\Sigma$ channel.

| Pole positions   | $\rho N$      | $\omega N$     | $\phi N$      | $K^*\Lambda$ | $K^*\Sigma$ |
|------------------|---------------|----------------|---------------|--------------|-------------|
| $1715 + i4$ MeV  | $-0.60 - i0.28$ | $0.43 - i0.03$ | $4.29 - i0.05$ | $-1.98 - i0.08$ | $0.18 + i0.03$ |
| $1728 + i0$ MeV  | $0.0 - i0.63$  | $0.0 + i0.40$  | $0.0 + i7.42$  | $0.0 - i3.03$  | $0.0 + i0.37$  |
| $1855 + i1$ MeV  | $-0.14 - i0.14$ | $-0.31 - i0.32$ | $1.52 - i0.01$ | $4.50 + i0.00$ | $1.85 + i0.03$ |
| $1868 + i6$ MeV  | $-0.13 - i0.14$ | $-0.31 - i0.37$ | $1.26 - i0.07$ | $2.99 - i0.06$ | $1.53 + i0.02$ |
| $1982 + i4$ MeV  | $0.10 + i0.27$ | $-0.03 + i0.02$ | $-0.28 - i0.19$ | $-1.01 - i0.17$ | $5.62 - i0.18$ |
| $1999 + i5$ MeV  | $0.07 + i0.21$ | $-0.02 + i0.02$ | $-0.20 - i0.16$ | $-0.59 - i0.12$ | $2.75 - i0.12$ |

Table III: Pole positions and coupling constants to various channels of resonances found in the isospin $I = 1/2$ and spin $J = 1/2$ sector.

The case of $J = 3/2$ also shows clear double peaks around 1700 MeV, 1860 MeV and 2000 MeV in the $(j', j', l', l) = (1, 1, 0, 0)$ and $(j', j', l', l) = (1, 1, 0, 2)$ sectors, which are depicted in Figs. 4 and 5 respectively. Similarly to the case of $J = 1/2$, the double peaks around 1700 MeV are visible in the $\rho N$ and $K^*\Lambda$ channels, and the double peaks around 1860 MeV appear mainly in the $K^*\Lambda$, $\phi N$ and $K^*\Sigma$ channels. Moreover, in Fig. 4 a lower peak can be seen around 2050 MeV near the double peaks around 2000 MeV in the $K^*\Sigma$ channel. In Fig. 6 the sector $(j', j', l', l) = (1, 2, 0, 2)$ is displayed, and we can see only one peak is left around 1700 MeV, 1860 MeV and 2000 MeV, respectively.

The couplings of the resonances to different channels for the $(I, J, j', j, l', l) = (1/2, 3/2, 1, 1, 0, 0)$ sector are listed in Table IV. Comparing with the $(I, J, j', j, l', l) = (1/2, 1/2, 1, 1, 0, 0)$ sector, we can see both $J = 3/2$ and $J = 1/2$ resonances are generated dynamically at the same positions of the complex plane of $\sqrt{s}$, i.e., $1715 + i4$ MeV, $1868 + i6$ MeV and $1982 + i4$ MeV. Furthermore, the couplings of some resonances to different channels take the same values as those of the $J = 1/2$ case, especially at the positions of $1715 + i4$ MeV and $1868 + i6$ MeV.

The poles at $1999 + i5$ MeV and $2045 + i25$ MeV all couple strongly to the $K^*\Sigma$ channel, and their couplings to different channels are similar to each other. It implies
TABLE IV: Pole positions and coupling constants to various channels of resonances found in the isospin $I = 1/2$ and spin $J = 3/2$ sector.

| Pole positions | $\rho N$ | $\omega N$ | $\phi N$ | $K^*\Lambda$ | $K^*\Sigma$ |
|----------------|----------|------------|----------|--------------|--------------|
| 1715 + $i4$ MeV | $-0.60 - i0.28$ | $0.43 - i0.02$ | $4.29 - i0.05$ | $-1.98 - i0.08$ | $0.18 + i0.03$ |
| 1868 + $i6$ MeV | $-0.13 - i0.14$ | $-0.31 - i0.37$ | $1.26 - i0.07$ | $2.99 - i0.06$ | $1.53 + 0.02$ |
| 1982 + $i4$ MeV | $0.05 + i0.13$ | $-0.01 + i0.01$ | $-0.14 - i0.10$ | $-0.51 - i0.08$ | $2.81 - i0.09$ |
| 1999 + $i5$ MeV | $0.07 + i0.21$ | $-0.02 + i0.02$ | $-0.19 - i0.16$ | $-0.59 - i0.12$ | $2.75 - i0.12$ |
| 2045 + $i25$ MeV | $0.02 + i0.48$ | $-0.25 + i0.10$ | $-0.21 - i0.48$ | $-0.66 - i0.59$ | $3.03 - i0.46$ |

TABLE V: The properties of the dynamically generated resonances with isospin $I = 1/2$ and their possible PDG counterparts.

| $J$ | Theory | PDG data |
|-----|--------|----------|
| Pole positions | Name and $J^P$ | Status | Mass | Width |
| 1/2 | 1715 + $i4$ MeV | $N(1650)1/2^-$ | **** | 1645 − 1670 MeV | 120 − 180 MeV |
| 1/2 | 1728 + $i0$ MeV | | | | |
| 1/2 | 1855 + $i1$ MeV | | | | |
| 1/2 | 1868 + $i6$ MeV | $N(1895)1/2^-$ | ** | $\approx$ 2090 MeV | 100 − 400 MeV |
| 3/2 | 1868+ $i6$ MeV | $N(1875)3/2^-$ | *** | 1820 − 1920 MeV | 160 − 320 MeV |
| 1/2 | 1982 + $i4$ MeV | | | | |
| 3/2 | 1982 + $i4$ MeV | $N(2120)3/2^-$ | ** | $\approx$ 2120 MeV | $\approx$ 300 MeV |
| 1/2 | 1999 + $i5$ MeV | | | | |
| 3/2 | 1999 + $i5$ MeV | | | | |
| 3/2 | 2045 + $i25$ MeV | | | | |

In Table V we show a summary of the results obtained and the tentative association to known states.

For the $(I, J) = (1/2, 3/2)$ $N^*$ states there is the $N(1700)$ with $J^P = 3/2^-$, which could be associated with the state we find with the same quantum number at 1715 + $i4$ MeV.
There are two resonances $N(1700)_{3/2}^-$ and $N(1685)^?_{1/2}^-$ in the same energy region of PDG data, while the total angular momentum $J$ and parity of the latter are not determined. Since the state $N(1685)^?_{1/2}^-$ does not gain status by being a sought-after member of a baryon anti-decuplet, we tend to assume the pole appeared at $1715+i4$ MeV might be the resonance $N(1700)_{3/2}^-$. We also find the resonance at $1715+i4$ MeV in the $J^P=1/2^-$ sector, and this $J^P=1/2^-$ states could correspond to the $N(1650)_{1/2}^-$, which could be the spin partner of the $N(1700)_{3/2}^-$. Our calculation shows the $N(1650)_{1/2}^-$ couples strongly to the $\phi N$ and $K^*\Lambda$ channels. At this point, it is different from the results obtained in Ref. [22], where the coupling constant to the channel $\rho N$ is the largest.

In the $J=3/2$ sector the poles at $1868+i6$ MeV and $1982+i4$ MeV could correspond to the resonances $N(1875)_{3/2}^-$ and $N(2120)_{3/2}^-$, respectively. These two resonances had been labeled as one resonance $N(2080)$ before the 2012 PDG Review [1].

In the region above 1800 MeV, only one resonance $N(1895)$ is listed with $J^P=1/2^-$ in the PDG data, which appears in the PDG review as $N^*(2090)_{S_{11}}$ before 2012 [1]. Although an estimated mass value of the state $N(1895)_{1/2}^-$ about 2090 MeV is listed in the PDG review [1], the newest multichannel analysis manifests the mass of this state is 1895 MeV in Ref. [28], which is close to the mass of the state $N(1875)_{3/2}^-$. Thus we treat the $N(1895)_{1/2}^-$ as a spin partner of the $N(1875)_{3/2}^-$, and assume the pole at $1868+i6$ MeV in the $J=1/2$ sector could correspond to the $N(1895)_{1/2}^-$. In addition, a high peak is also found at $1982+i4$ MeV in the $J=1/2$ sector, and no counterpart is listed in the PDG data. It might stimulate the experimental research to look for the resonance around 2000 MeV, which might be a spin partner of the $N(2120)_{3/2}^-$ and couple strongly to the $K^*\Sigma$ channel as listed in Table III.

B. Isospin $I=3/2$

The isospin wave functions for the $I=3/2$ case are written as

$$|\rho N; \frac{3}{2}; \frac{1}{2}\rangle = \sqrt{\frac{2}{3}} |\rho_0 p\rangle - \sqrt{\frac{1}{3}} |\rho_+ n\rangle,$$  \hspace{1cm} (36)

and

$$|K^*\Sigma; \frac{3}{2}; \frac{1}{2}\rangle = \sqrt{\frac{2}{3}} |K^*\Sigma^0\rangle - \sqrt{\frac{1}{3}} |K^*\Sigma^+\rangle,$$ \hspace{1cm} (37)
respectively. The coupling constants of octet baryons and vector mesons take the same values when the isospin orientation $I_z$ is different from one another. Especially, the $s$–channel coupling constants tend to zero, and it is apparent that the $s$–channel interaction is forbidden in the case of isospin $I = 3/2$.

![Diagram](image)

**FIG. 7:** $|T|^2$ for different channels with $I = 3/2$, $J = 1/2$, $j' = 1$, $j = 1$, $l' = 0$ and $l = 0$.

![Diagram](image)

**FIG. 8:** $|T|^2$ for different channels with $I = 3/2$, $J = 1/2$, $j' = 1$, $j = 1$, $l' = 0$ and $l = 2$.

We show the $|T_{ii}|^2$ as a function of $\sqrt{s}$ for different partial waves in Figs. 7, 8, 9, 10 and 11 respectively, and we find four peaks in these figures can be seen, which correspond to four poles near the real axis in the complex plane $\sqrt{s}$, i.e., $1654 + i11$ MeV, $1726 + i9$ MeV, $1758 + i0$ MeV and $1780 + i14$ MeV. Moreover, there are two other poles at $1619 + i160$ MeV.
FIG. 9: $|T|^2$ for different channels with $I = 3/2$, $J = 3/2$, $j' = 1$, $j = 1$, $l' = 0$ and $l = 0$.

FIG. 10: $|T|^2$ for different channels with $I = 3/2$, $J = 3/2$, $j' = 1$, $j = 1$, $l' = 0$ and $l = 2$.

and $1680 + i163$ MeV, which are far from the real axis and do not show clearly in the figures. All these states couple strongly to the $K^*\Sigma$ channel. The couplings of these resonances to the $\rho N$ and $K^*\Sigma$ channels are listed in Table VI for the $J = 1/2$ case and Table VII for the $J = 3/2$ case.

For the case of $(I, J) = (3/2, 1/2)$ there is one state in the PDG, the $\Delta(1620)1/2^-$, which has the spin partner $\Delta(1700)3/2^-$ with $J^P = 3/2^-$. These two states could be associated with the pole at $1654 + i11$ MeV, which appeared both in the $J = 1/2$ and $J = 3/2$ sectors by the method of partial wave analysis in Eq. (30).

We find a narrow peak at $1758 + i0$ MeV in the $J = 1/2$ and $J = 3/2$ sectors, and its
FIG. 11: $|T|^2$ for different channels with $I = 3/2$, $J = 3/2$, $j' = 1$, $j = 2$, $l' = 0$ and $l = 2$.

| Pole positions        | $\rho N$     | $K^*\Sigma$ |
|------------------------|--------------|--------------|
| 1619 + i160 MeV        | 0.87 − i2.64 | 4.93 − i4.95 |
| 1654 + i11 MeV         | 3.04 + i0.40 | 12.35 + i0.04 |
| 1680 + i163 MeV        | −1.14 − i0.87| 4.09 − i3.72 |
| 1726 + i9 MeV          | 1.34 + i0.41 | 8.89 + i0.01 |
| 1758 + i0 MeV          | 0.00 + i0.57 | 0.00 + i2.08 |
| 1780 + i14 MeV         | 0.68 + i0.27 | 6.45 + i0.05 |

TABLE VI: Pole positions and coupling constants to various channels of resonances found in the isospin $I = 3/2$ and spin $J = 1/2$ sector.

couplings to $\rho N$ and $K^*\Sigma$ are smaller by far than those of other states. We suspect it could be a cusp.

In Ref. [22], no resonance is found in the channel of isospin $I = 3/2$ and strangeness zero since only $t$− channel is taken into account, which supply a repulsive interaction between vector mesons and octet baryons. However, our calculation results manifest some resonances can be produced dynamically in the $I = 3/2$ sector when we take into account the other interaction modes besides $t$− channel, especially the contact term between vector mesons and octet baryons. At this point, our results are consistent to those in Ref. [25].

The states found in the $I = 3/2$ sector are summarized in Table VIII where the properties of the possible counterparts are also listed. Except the states at 1654+i11 MeV with $J = 1/2$
TABLE VII: Pole positions and coupling constants to various channels of resonances found in the isospin $I = 3/2$ and spin $J = 3/2$ sector.

| $J$ | Theory Pole positions | PDG data Name and $J^P$ Status | Mass | Width |
|-----|----------------------|---------------------------------|------|-------|
| 1/2 | $1619 + i160$ MeV     |                                 |      |       |
| 3/2 | $1619 + i160$ MeV     |                                 |      |       |
| 1/2 | $1654 + i11$ MeV      | $\Delta(1620)1/2^-$ ****       | 1600 – 1660 MeV | 130 – 150 MeV |
| 3/2 | $1654 + i11$ MeV      | $\Delta(1700)3/2^-$ ****       | 1670 – 1750 MeV | 200 – 400 MeV |
| 1/2 | $1680 + i163$ MeV     |                                 |      |       |
| 1/2 | $1726 + i9$ MeV       |                                 |      |       |
| 3/2 | $1726 + i9$ MeV       |                                 |      |       |
| 1/2 | $1758 + i0$ MeV       |                                 |      |       |
| 3/2 | $1758 + i0$ MeV       |                                 |      |       |
| 1/2 | $1780 + i14$ MeV      |                                 |      |       |

TABLE VIII: The properties of dynamically generated resonances with isospin $I = 3/2$ and its possible PDG counterparts.

VI. SUMMARY

We have studied the interaction between vector mesons and octet baryons using a unitary framework in coupled channels. In addition to the vector interaction term, a tensor term is included in the Lagrangian of vector mesons and octet baryons. With this interaction Lagrangian, we obtain the kernels between vector mesons and octet baryons from a
summation of diagrams corresponding to a vector meson exchange in the $t-$ channel, octet baryon exchange in $s-, u-$ channels, and a contact interaction related only to the tensor term. The scattering amplitudes are calculated by solving the coupled-channel Lippman-Schwinger equations, and then the results are studied with the method of partial wave analysis.

In the $I = 1/2$ sector, the double-peak structure of $|T_{ii}|^2$ is found around 1700 MeV, 1860 MeV and 2000 MeV respectively in the different partial waves, and these double-peaks could correspond to the states in the PDG data. The pole found at $1715 + i4$ MeV with different $J$ in the complex energy plane might correspond to the states $N(1650)1/2^-$ in the $J = 1/2$ case, and its spin partner $N(1700)3/2^-$ in the $J = 3/2$ case. Similarly, the pole at $1868 + i6$ MeV could correspond to another pair of spin partners, $N(1895)1/2^-$ and $N(1875)3/2^-$ of the PDG data. The $N(2120)3/2^-$ could be associated with the pole at $1982 + i4$ MeV in the $J = 3/2$ case. However, the spin partner of $N(2120)3/2^-$ is absent in the PDG data, and we predict that there should be a state around 2000 MeV in the $J = 1/2$ case, which should be associated with the pole at $1982 + i4$ MeV with $J = 1/2$ as a spin partner of the state $N(2120)3/2^-$. 

In the $I = 3/2$ sector, we also find some poles in the complex energy plane, and we assume the spin partners $\Delta(1620)1/2^-$ and $\Delta(1700)3/2^-$ could be associated with the pole at $1654 + i11$ MeV in the $J = 1/2$ and $J = 3/2$ cases.

Furthermore, it manifests that the interaction between vector mesons and octet baryons in the $I = 3/2$ case can be attractive when the $u-, s-$ channels and the contact term are taken into account besides $t-$ channel interaction. In addition, there are not PDG counterparts for some resonances in our prediction, and it might give a stimulus to search experimentally for these resonance states.

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Appendix I: Vertices used in the calculation

(a) Three-vector-meson vertex, (b) Baryon-baryon-meson vertex, where the momentum of the vector meson $q = p_2 - p_1$. (c) Contact vertex of octet baryons and vector mesons.

A. Vertices for three vector mesons

The vertex for three vector mesons can be deduced from the Lagrangian

$$L_{3V} = \frac{ig}{2} \sum_{l,m,n=1}^{8} C_{3V}(l,m,n) M_l^\mu M_m^\nu \left( \partial_\mu M_{n\nu} - \partial_\nu M_{n\mu} \right)$$

with

$$C_{3V}(l,m,n) = \sum_{i,j,k=1}^{8} X_{il} X_{jm} X_{kn} (d_{ijk} + if_{ijk}).$$

In Eq. (39), the matrix $X$ is indicated as

$$X = \frac{1}{\sqrt{2}} \begin{pmatrix}
1 & 1 & 0 & 0 & 0 & 0 & 0 & 0 \\
0 & 0 & i & -i & 0 & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & \sqrt{2} & 0 & 0 \\
0 & 0 & 0 & i & -i & 0 & 0 & 0 \\
0 & 0 & 0 & 0 & 0 & 1 & 1 & 0 \\
0 & 0 & 0 & 0 & 0 & 0 & i & -i \\
0 & 0 & 0 & 0 & 0 & 0 & 0 & \sqrt{2}
\end{pmatrix},$$

and the values of $d_{ijk}$ and $f_{ijk}$ can be found in Ref. [1]. Moreover, the coefficient in Eq. (39) must be multiplied by a factor of $\frac{1}{\sqrt{3}}$ for each $\omega$ meson line, and a factor of $-\sqrt{\frac{3}{5}}$ for each $\phi$ meson line. The vertices for three vector mesons are depicted in Fig. 12(a).
B. Baryon-baryon-meson Vertices

The interaction vertices for two baryons and one vector meson depicted in Fig. 12(\(b\)) can be obtained according to \(SU(3)\) symmetry

\[
\Gamma^\mu(p_2, p_1) = g_1 \gamma^\mu + g_2 (p_2^\mu + p_1^\mu),
\]

where

\[
g_1(l, m, n) = g ((F_V + D_V)C_{BBV}(l, n, m) + (-F_V + D_V)C_{BBV}(l, m, n))
+ \frac{g}{2M_N} ((F_T + D_T)C_{BBV}(l, n, m) + (-F_T + D_T)C_{BBV}(l, m, n)) (M_l + M_m),
\]

and

\[
g_2(l, m, n) = -\frac{g}{2M_N} ((F_T + D_T)C_{BBV}(l, n, m) + (-F_T + D_T)C_{BBV}(l, m, n))
\]

with

\[
C_{BBV}(l, n, m) = \frac{1}{2} \sum_{i,j,k=1}^8 X_i X_j X_k (d_{ijk} + i f_{ijk}).
\]

C. Contact term

From the lagrangian in Eq. (11), the coupling constant of the contact term is written as

\[
C(i', j', k', l') = \frac{ig^2}{32M_N} \sum_{i,j,k,l=1}^8 \{ F_T[\langle \lambda_i \lambda_j \lambda_k \lambda_l \rangle - \langle \lambda_i \lambda_k \lambda_j \lambda_l \rangle - \langle \lambda_i \lambda_l \lambda_k \lambda_j \rangle + \langle \lambda_i \lambda_l \lambda_j \lambda_k \rangle] \\
+ D_T[\langle \lambda_i \lambda_j \lambda_k \lambda_l \rangle - \langle \lambda_i \lambda_k \lambda_l \lambda_j \rangle + \langle \lambda_i \lambda_l \lambda_k \lambda_j \rangle - \langle \lambda_i \lambda_l \lambda_j \lambda_k \rangle] \}
X_{i',i}^\dagger X_{j,j'} X_{k,k'} X_{l,l'},
\]

where

\[
\langle \lambda_i \lambda_j \lambda_k \lambda_l \rangle = \frac{4}{3} \delta_{ij} \delta_{kl} + 2 (d_{klm} + i f_{klm}) (d_{ijm} + i f_{ijm})
\]

with \(\lambda_i\) the hermitian generator of group \(SU(3)\). The vertex of contact terms is depicted in Fig. 12(c). We can obtain the coupling constants \(C_{CT}^{IS}\) of the contact term for different isospins with Clebsch-Gordan coefficients.
Appendix II: Polarization vectors and Dirac spinors

D. Polarization vectors of vector mesons

The polarization vector of the massive vector field satisfies the constraint

\[ k \cdot \varepsilon_{k\lambda} = 0, \tag{47} \]

with \( k \) the momentum of the vector meson.

For each \( k \), a set of three linearly independent polarization vectors satisfying Eq. (47) can be constructed as in Ref. [26]. If \( \varepsilon_{k\lambda} (\lambda = 1, 2, 3) \) is any triad of three-vectors satisfying the orthonormality relations

\[ \varepsilon_{k\lambda}^* \cdot \varepsilon_{k\lambda'}^* = \delta_{\lambda\lambda'}, \tag{48} \]

then the three four-vectors can be written as

\[ \varepsilon_{k\lambda}^a = \begin{cases} \frac{(k \cdot \varepsilon_{k\lambda})}{\mu}, & (\alpha = 0), \\ \varepsilon_{k\lambda}^* + \frac{k}{\mu(\omega_k + \mu)} (k \cdot \varepsilon_{k\lambda}), & (\alpha = 1, 2, 3), \end{cases} \tag{49} \]

with \( \mu \) the mass and \( \omega_k = \sqrt{k^2 + \mu^2} \) the energy of the vector meson. The four-vectors in Eq. (49) satisfy both Eq. (47) and the orthonormality relations

\[ \varepsilon_{k\lambda}^a \cdot \varepsilon_{k\lambda'}^* = -g_{\lambda\lambda'}, \tag{50} \]

with the metric tensor

\[ g_{\lambda\lambda'} = \begin{pmatrix} 1 & 0 & 0 & 0 \\ 0 & -1 & 0 & 0 \\ 0 & 0 & -1 & 0 \\ 0 & 0 & 0 & -1 \end{pmatrix}. \tag{51} \]

Therefore, the notations used in this article are different from those in Ref. [26].

A convenient choice for the triad of orthogonal unit vectors \( \varepsilon_{k\lambda}^a \) is to take \( \varepsilon_{k3}^a \) pointing along the three-momentum \( \vec{k} \) with \( \varepsilon_{k1}^a \) and \( \varepsilon_{k2}^a \) orthogonal both to \( \varepsilon_{k3}^a \) and to each other, i.e.

\[ \varepsilon_{k3} = \frac{\vec{k}}{|\vec{k}|}, \tag{52} \]
and

\[ \vec{\varepsilon}_{k1} \cdot \vec{k} = \vec{\varepsilon}_{k2} \cdot \vec{k} = \vec{\varepsilon}_{k1} \cdot \vec{\varepsilon}_{k2} = 0. \]

The polarization vectors \( \vec{\varepsilon}_{k1} \) and \( \vec{\varepsilon}_{k2} \) represent states of \textit{transverse} polarization, while \( \vec{\varepsilon}_{k3} \) represents \textit{longitudinal} polarization.

In the scattering process of the vector meson and octet baryon, we can choose the \textit{transverse} polarization vectors of the initial vector meson to be

\[ \vec{\varepsilon}_{k1} = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 \\ i \\ 0 \end{pmatrix}, \quad \vec{\varepsilon}_{k2} = \frac{1}{\sqrt{2}} \begin{pmatrix} 1 \\ -i \\ 0 \end{pmatrix}, \quad (53) \]

and the \textit{longitudinal} polarization vector

\[ \vec{\varepsilon}_{k3} = \begin{pmatrix} 0 \\ 0 \\ 1 \end{pmatrix}, \quad (54) \]

with \( \vec{k} \) the three-momentum of the initial vector meson. For the final vector meson, the \textit{transverse} and \textit{longitudinal} polarization vectors take the same formula as those of the initial vector meson in Eqs. (53) and (54).

The polarization vectors \( \varepsilon_{k1}, \varepsilon_{k2}, \varepsilon_{k3} \) and \( k/m \) form a quartet of orthonormal four-vectors. From the completeness relation for this quartet we deduce that

\[ \sum_{\lambda=1,2,3} \varepsilon_{k\lambda}^{\mu} \varepsilon_{k\lambda}^{*\nu} = -g^{\mu\nu} + \frac{k^{\mu}k^{*\nu}}{\mu^2}. \quad (55) \]

E. Dirac spin wave function of octet baryons

The Dirac equation implies

\[ (\not{p} - M) U(p, \lambda) = 0, \quad (56) \]

where the spinor of octet baryons can be written as

\[ U(p, \lambda) = \frac{\not{p} + M}{\sqrt{2M(M + \not{E})}} U(M, \vec{0}, \lambda), \quad (57) \]
with
\[
U(M, \vec{0}, 1) = \begin{pmatrix}
1 \\
0 \\
0 \\
0
\end{pmatrix}
\] (58)

and
\[
U(M, \vec{0}, 2) = \begin{pmatrix}
0 \\
1 \\
0 \\
0
\end{pmatrix}.
\] (59)

The conjugate spinor of octet baryons is obtained as
\[
\bar{U}(p, \lambda) = \bar{U}(M, \vec{0}, \lambda) \frac{\not{p} + M}{\sqrt{2M(M + E)}},
\] (60)

and the normalization factors have been chosen in order that
\[
\bar{U}(p, \lambda)U(p, \lambda') = \delta_{\lambda, \lambda'}.
\] (61)

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