Interactions of some decuplet baryons in the quark cluster model

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In the framework of the resonating-group method, the interactions of decuplet baryon-baryon systems with strangeness numbers $s=-1$ and $s=-5$ are investigated by employing the chiral SU(3) quark model. The effective baryon-baryon interactions deduced from quark-quark interactions and scattering cross sections of the $\Sigma^*\Delta$ and $\Xi^*\Omega$ systems are calculated.

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I. INTRODUCTION

With the success of quantum chromodynamics (QCD) in describing hadronic structure in past few years, it has been widely believed that QCD is the fundamental theory of the strong interaction and that quarks and gluons are elementary building blocks of hadronic systems. However, most of the studies of hadronic structure and strong interactions rely on the theory of QCD in a relatively-low-energy regime where QCD has to be treated nonperturbatively. Because of the complexity of nonperturbative QCD (NPQCD), a variety of QCD-inspired models have been developed for this purpose. In some models, besides the quark-gluon degrees of freedom, other effective degrees of freedom, such as constituent quarks [1], chiral fields [2] etc., are included to describe perturbative QCD (PQCD) and NPQCD effects in an efficient way. The extent to which these models can be used should be tested by comparing their results to a variety of available experimental data. It is well known that low-energy scattering data mainly reflect the long-distance behavior of the strong interaction and are not sensitive to the detailed form of the short-range interaction. In order to examine the interaction at medium and short distances, the structures of hadrons, especially the possible existence of dibaryons [3–19], should intensively be studied. The reason is clear. As a result of the relatively smaller size of dibaryons, the quark-gluon structure of the dibaryon becomes apparent, so that PQCD and NPQCD effects should be carefully studied. No doubt, the existence of such a system in the real world would definitely support the NPQCD approach used to calculate these structures. As a direct benefit, it would also help us to understand the baryon-baryon (*B*-*B*) interaction, in particular its short- and medium-range parts.

Although abundant data of *NN* scattering can be used to refine the strong interactions used in the models either in the hadronic degrees of freedom or in the quark-gluon degrees of freedom, the scarce data of strange hadron scattering are not adequate to fix the free parameters in the strange sector of *B*-*B* interaction uniquely [18]. Along with the operations of AGS, SPS, TJNAF, RHIC, etc. [20–23], strange baryons, such as Λ, Σ, Ξ, and Ω, and strange mesons, for instance K and others, can be produced. It enables us to investigate strangeness-rich systems, which are of fundamental importance in establishing a reliable *B*-*B* interaction, in understanding the existence of hypernuclei [20], and in studying relativistic heavy-ion collisions [21] and some astrophysical problems [22]. It is especially noticed that almost all projects at RHIC have, among their own major aims, the aim of searching for strange pieces of nuclear matter: either dibaryons or larger objects with strangeness. Such studies provide an exciting perspective for a multistrange system investigation which often cannot be done by conventional methods. The data would give certain guidance to phenomenological models.

Various theoretical studies of such systems have been done in the past few years (see [18,23,24,33] and references therein). Stoks and Rijken extended the Nijmegen soft-core potential to the $\Sigma\Sigma$, the $\Sigma\Xi$, and the $\Xi\Xi$ systems [18]. Fujita and co-workers studied baryon-baryon interactions containing strangeness [33]. In terms of the chiral SU(3) quark model, Yu *et al.* predicted $\Omega\Omega_{(0,0)}$ and $\Xi\Omega_{(1,1/2)}$ dibaryons [25], Li *et al.* predicted dibaryons with various strangeness [26], and Zhang *et al.* studied light Λ hypernuclei [27] (also references therein).

It should be emphasized that if one wishes to predict new physics in terms of a theoretical model, the model must satisfy two requirements: (1) most of the available data must be explained by the model, and (2) when the model is extrapolated to predict new observables, no additional model parameters are required. The chiral SU(3) quark model is one of the models that meet these conditions. By employing such a model, most of the relevant data, such as the ground-state masses and electromagnetic properties of single baryons [28], *N*-*N* scattering phase shifts, and *N*-Y scattering cross sections [29], and the deuteron (*d*) and the H particle binding energies [13], can successfully be reproduced. Moreover, light hypernuclei ${}^5_\Lambda$He, ${}^4_\Lambda$He, and ${}^4_\Lambda$H can simultaneously be described [27]. When this model is applied to the study of multistrange systems, no additional parameters are needed. It should especially be noted that although relativistic effects have not been fully included in models of this kind, these models are still widely employed in preliminary investigations of hadronic physics, because they are relatively simple, the picture of underlying physics is clear, and many data can be explained. To compensate inadequate relativistic effects effectively, a few model parameters are commonly introduced by fitting empirical data. In this paper, we will employ the chiral SU(3) quark model to study further scattering and binding behaviors and effective interactions between baryons in the $\Sigma^*\Delta$ system ($s=-1$) and in the $\Xi^*\Omega$ system ($s=-5$) where, up to now, only limited data are available.

This paper is arranged as follows. In Sec. II, we briefly describe the chiral SU(3) quark model. In Sec. III, we present results and discussions. The conclusions are drawn in Sec. IV.

II. MODEL

Following Georgi's idea [30], the coupling of chiral fields to quark fields can be written as

$$\mathcal{L}_I=-g_{ch}(\overline{\psi_L}\Sigma\psi_R-\overline{\psi_R}\Sigma^\dagger\psi_L), \tag{1}$$

where g_{ch} denotes the quark-chiral field coupling constant, ψ_L and ψ_R represent the left- and right-quark spinors, respectively, and

$$\Sigma=\exp[i\pi_a\lambda_a/f], \quad a=1,2,\ldots,8, \tag{2}$$

with π_a being the Goldstone boson field and λ_a the Gell-Mann matrix of the flavor SU(3) group. As an extrapolation of the linear realization of Σ in the SU(2) case, one can approximately write Σ in a linearized form

$$\Sigma=\sum_{a=0}^{8}\sigma_a\lambda_a-i\sum_{a=0}^{8}\pi_a\lambda_a \tag{3}$$

and the interactive Lagrangian as

$$\mathcal{L}_I=-g_{ch}\overline{\psi}\left(\sum_{a=0}^{8}\sigma_a\lambda_a+i\sum_{a=0}^{8}\pi_a\lambda_a\gamma_5\right)\psi, \tag{4}$$

where λ_0 is a unitary matrix. In Eq. (4), σ_a are scalar nonet fields and correspond to scalar mesons $\sigma(a=0)$, $\sigma'(a=1,\ldots,3)$, $\kappa(a=4,\ldots,7)$, and $\epsilon(a=8)$, respectively. π_a are pseudoscalar nonet fields and describe pseudoscalar mesons $\pi(a=1,\ldots,3)$, K $(a=4,\ldots,7)$, $\eta_0(a=0)$, and $\eta_8(a=8)$, respectively. Linear combination of η_0 and η_8 forms η and η'. Clearly, $\mathcal{L}_I$ is invariant under the infinitesimal chiral $\mathrm{SU}(3)_L\times\mathrm{SU}(3)_R$ transformation. Consequently, one can write the interactive Hamiltonian as

$$H_{ch}=g_{ch}F(q^2)\overline{\psi}\left(\sum_{a=0}^{8}\sigma_a\lambda_a+i\sum_{a=0}^{8}\pi_a\lambda_a\gamma_5\right)\psi, \tag{5}$$

with ψ being the quark field. In this expression, a form factor $F(q^2)$ is phenomenologically introduced to describe the internal structure of the chiral field [31]. As usual,

$$F(q^2)=\left(\frac{\Lambda_{CSB}^2}{\Lambda_{CSB}^2+q^2}\right)^{1/2}, \tag{6}$$

and the cutoff mass Λ_{CSB} represents the breaking scale of chiral symmetry [29]. By using standard field theory method, one obtains V_{ij}^{ch}, a set of potentials between the ith and jth quarks due to exchanges of pseudoscalar and scalar mesons, as the following:

$$V_{ij}^{ch}=\sum_{a=0}^{8}V_{\sigma_a}(\vec{r}_{ij})+V_{\pi_a}(\vec{r}_{ij}), \tag{7}$$

$$V_{\sigma_a}(\vec{r}_{ij})=V_{\sigma_a}^{Cen}(\vec{r}_{ij})+V_{\sigma_a}^{LS}(\vec{r}_{ij}), \tag{8}$$

$$V_{\pi_a}(\vec{r}_{ij})=V_{\pi_a}^{SS}(\vec{r}_{ij})+V_{\pi_a}^{Ten}(\vec{r}_{ij}), \tag{9}$$

with central force

$$V_{\sigma_a}^{Cen}(\vec{r}_{ij})=-C(g_{ch},m_{\sigma_a},\Lambda_{CSB})X_1(m_{\sigma_a},\Lambda_{CSB},r_{ij})\times[\lambda_a(i)\lambda_a(j)], \tag{10}$$

spin-spin force

$$V_{\pi_a}^{SS}(\vec{r}_{ij})=C(g_{ch},m_{\pi_a},\Lambda_{CSB})\times\left[\frac{m_{\pi_a}^2}{12m_{qi}m_{qj}}X_2(m_{\pi_a},\Lambda_{CSB},r_{ij})(\vec{\sigma}_i\cdot\vec{\sigma}_j)\right]\times[\lambda_a(i)\lambda_a(j)], \tag{11}$$

spin-orbit force

$$V_{\sigma_a}^{LS}(\vec{r}_{ij})=-C(g_{ch},m_{\sigma_a},\Lambda_{CSB})\frac{m_{\sigma_a}^2}{4m_{qi}m_{qj}}\times\left[G(m_{\sigma_a}r_{ij})-\left(\frac{\Lambda_{CSB}}{m_{\sigma_a}}\right)^3G(\Lambda_{CSB}r_{ij})\right]\times[\vec{L}\cdot(\vec{\sigma}_i+\vec{\sigma}_j)][\lambda_a(i)\lambda_a(j)], \tag{12}$$

tensor force

$$V_{\pi_a}^{Ten}(\vec{r}_{ij})=C(g_{ch},m_{\pi_a},\Lambda_{CSB})\frac{m_{\pi_a}^2}{12m_{qi}m_{qj}}\times\left\{\left[H(m_{\pi_a}r_{ij})-\left(\frac{\Lambda_{CSB}}{m_{\pi_a}}\right)^3H(\Lambda_{CSB}r_{ij})\right]S_{ij}\right\}\times[\lambda_a(i)\lambda_a(j)], \tag{13}$$

and

$$S_{ij}=3(\vec{\sigma}_i\cdot\hat{r})(\vec{\sigma}_j\cdot\hat{r})-(\vec{\sigma}_i\cdot\vec{\sigma}_j), \tag{14}$$

$$C(g_{ch},m,\Lambda)=\frac{g_{ch}^2}{4\pi}\frac{\Lambda^2}{\Lambda^2-m^2}m, \tag{15}$$

$$X_1(m,\Lambda,r)=Y(mr)-\frac{\Lambda}{m}Y(\Lambda r), \tag{16}$$

$$X_2(m,\Lambda,r)=Y(mr)-\left(\frac{\Lambda}{m}\right)^3Y(\Lambda r), \tag{17}$$

$$H(x)=\left(1+\frac{3}{x}+\frac{3}{x^2}\right)Y(x), \tag{18}$$

$$G(x)=\frac{1}{x}\left(1+\frac{1}{x}\right)Y(x), \tag{19}$$

$$Y(x)=\frac{1}{x}e^{-x}, \tag{20}$$

and m_{σ_a} being the mass of the scalar meson and m_{π_a} the mass of the pseudoscalar meson.

The above-derived interactions describe the NPQCD effect in the low-momentum transfer region, which is very important in explaining the short- and medium-range forces between two baryons. However, to study the hadronic structure and the baryon-baryon dynamics, one still needs an effective one-gluon-exchange interaction V_{ij}^{OGE} which dominates the short-range PQCD behavior and a phenomenological confinement potential V_{ij}^{conf} which provides the NPQCD effect at large distances and confines three quarks in a baryon.

V^{OGE} can be derived from the Lagrangian

$$\mathcal{L}_I^{OGE}=-\frac{g_s}{2}\bar{\psi}\lambda_a^c\gamma_\mu A_a^\mu\psi, \tag{21}$$

where A_a^μ is the gluon field, λ_a^c denotes the generator of the SU(3) group in color space, and g_s represents the quark-gluon coupling constant. The explicit form is as follows:

$$V_{ij}^{OGE}=\frac{1}{4}g_ig_j(\lambda_i^c\cdot\lambda_j^c)\left\{\frac{1}{r_{ij}}-\frac{\pi}{2}\delta(\vec{r}_{ij})\left(\frac{1}{m_{qi}^2}+\frac{1}{m_{qj}^2}+\frac{4}{3}\frac{1}{m_{qi}m_{qj}}(\vec{\sigma}_i\cdot\vec{\sigma}_j)\right)\right\}+V_{OGE}^{Ten}+V_{OGE}^{LS}, \tag{22}$$

with

$$V_{OGE}^{Ten}=-\frac{1}{4}g_ig_j(\lambda_i^c\cdot\lambda_j^c)\frac{1}{4m_{qi}m_{qj}r_{ij}^3}S_{ij}, \tag{23}$$

$$V_{OGE}^{LS}=-\frac{1}{16}g_ig_j(\lambda_i^c\cdot\lambda_j^c)\frac{3}{m_{q_i}m_{q_j}r_{ij}^3}\vec{L}\cdot(\vec{\sigma}_i+\vec{\sigma}_j). \tag{24}$$

The confinement can be chosen in a quadratic form:

$$V_{ij}^{conf}=-a_{ij}^c(\lambda_i^c\cdot\lambda_j^c)r_{ij}^2-a_{ij}^{c0}(\lambda_i^c\cdot\lambda_j^c), \tag{25}$$

where a_{ij}^c is the confinement strength between quarks i and j. The introduction of a_{ij}^{c0}, which is related to the "zero-point energies" of various single baryons, is to fit the masses of octet and decuplet baryons to the empirical data. It should be noted that the form of the confinement potential—for instance, the quadratic form or the linear form—would not seriously affect lower-lying spectra of single baryons and would give negligible contribution to the interbaryon interactions.

Then, the total Hamiltonian of the six-quark system reads

$$H=\sum_i T_i-T_G+\sum_{i<j}\{V_{ij}^{OGE}+V_{ij}^{conf}+V_{ij}^{ch}\}. \tag{26}$$

TABLE I. Model parameters.

	Set I	Set II
m_u (MeV)	313	313
m_s (MeV)	470	470
b_u (fm)	0.505	0.505
g_u	0.936	0.936
g_s	0.924	0.781
a_{uu}^c (MeV/fm^2)/a_{uu}^{c0} (MeV)	54.3/−47.7	55.0/−48.9
a_{us}^c (MeV/fm^2)/a_{us}^{c0} (MeV)	65.8/−41.7	66.5/−50.6
a_{ss}^c (MeV/fm^2)/a_{ss}^{c0} (MeV)	103.0/−50.6	115.4/−73.7
m_π (fm^{-1})/Λ_π (fm^{-1})	0.7/4.2	0.7/4.2
m_K (fm^{-1})/Λ_K (fm^{-1})	2.51/4.2	2.51/4.2
m_η (fm^{-1})/Λ_η (fm^{-1})	2.78/5.0	2.78/5.0
$m_{\eta'}$ (fm^{-1})/$\Lambda_{\eta'}$ (fm^{-1})	4.85/5.0	4.85/5.0
m_{σ_0} (fm^{-1})/Λ_{σ_0} (fm^{-1})	3.17/4.2	3.17/4.2
$m_{\sigma'}$ (fm^{-1})/$\Lambda_{\sigma'}$ (fm^{-1})	4.85/5.0	4.85/5.0
m_κ (fm^{-1})/Λ_κ (fm^{-1})	4.85/5.0	7.09/7.6
m_ϵ (fm^{-1})/Λ_ϵ (fm^{-1})	4.85/5.0	7.09/7.6

Model parameters should be fixed at the beginning of the calculation. The coupling constant g_{ch} is fixed by

$$\frac{g_{ch}^2}{4\pi}=\left(\frac{3}{5}\right)^2\frac{g_{NN\pi}^2}{4\pi}\frac{m_q^2}{M_N^2}, \tag{27}$$

with $g_{NN\pi}^2/4\pi$ taken as the empirical value of 13.67. The remaining model parameters are determined in the following way: $m_u=313$ MeV and $m_s=470$ MeV are taken as usual, and the value of b_u is chosen in order to be consistent with the empirical value of the nucleon radius. m_{π_a} and m_{σ_a} are taken to be the masses of corresponding mesons, and Λ_{CSB} is taken to be the breaking scale of chiral symmetry. Then, we use eight constrains to fix the remained eight model parameters: strong-coupling constants g_u and g_s can be determined by the mass splittings between N, Δ and Σ, Λ, respectively; confinement strengths a_{uu}^c, a_{us}^c, and a_{ss}^c are fixed by the stability conditions of N, Λ, and Ξ, respectively; and a_{uu}^{c0}, a_{us}^{c0}, and a_{ss}^{c0} are fixed by the masses of N, Λ, and $\overline{\Sigma+\Omega}$, respectively. The resultant model parameters (set I and set II) are tabulated in Table I.

In Table I, the parameters in set I were frequently used in our previous works [13,26,29]. However, in multistrange systems, mesons with strangeness may play important roles. Sensibility to the masses of strange mesons can be tested by increasing the masses of strange mesons κ and ϵ to 1400 MeV and the cutoff masses to 1500 MeV, respectively. Corresponding model parameters are those as tabulated in set II. It has been shown that with either set, most of the existent data, including the masses of ground-state baryons (see Table II), the N-N and N-Y scattering data [29], the binding energy of deuteron, etc., can be well reproduced, the resultant mass of the H particle agrees with the experimental finding, and the mass spectra of light Λ hypernuclei can generally be reproduced [27]. It is natural to apply this model with these

TABLE II. The masses of octet and decuplet baryons.

	Theor.	Expt.		Theor.	Expt.
N	939	939	Δ	1237	1232
Λ	1116	1116	Σ^*	1375	1385
Σ	1194	1194	Ξ^*	1515	1530
Ξ	1334	1319	Ω	1657	1672

sets of parameters to the systems with higher strangeness, in particular, to the baryon-baryon interactions in the $s=-1$ to -5 sectors.[1]

The effective baryon-baryon interaction, the scattering cross section, and the binding energy can dynamically be calculated by using the resonating group method (RGM). In this method, the orbital wave function of the ith quark (i can be either an up, down, or strange quark) can be written as

$$\Phi_i(\vec{r}_i)=(1/\pi b_i^2)^{3/4}\exp[-(\vec{r}_i-\vec{R})^2/2b_i^2], \tag{28}$$

where $\vec{r}_i$ and $\vec{R}$ are the coordinate vectors of the ith quark and the center-of-mass motion of the baryons, respectively. The width parameter b_i is associated with the oscillator frequency ω by the constituent quark mass m_i:

$$\frac{1}{b_i^2}=m_i\omega. \tag{29}$$

The orbital wave function of a baryon can be expressed as

$$\hat{\phi}=\mathcal{A}\Pi_{i=1}^{3}\Phi_i(\vec{r}_i). \tag{30}$$

Then, the trial wave function of the six-quark system can be written as

$$\Psi=\mathcal{A}[\hat{\phi}_A(\vec{\xi}_1,\vec{\xi}_2)\hat{\phi}_B(\vec{\xi}_4,\vec{\xi}_5)\chi_{rel}(\vec{R})\chi_{\text{c.m.}}(\vec{R}_{\text{c.m.}})]_{ST}, \tag{31}$$

where $\hat{\phi}_{A(B)}$ denotes the antisymmetrized wave function of the baryon cluster $A(B)$ with $A(B)$ further specifying all the quantum numbers of the relevant cluster, $\chi_{rel}(\vec{R})$ the trial wave function of the relative motion between interacting clusters A and B, $\chi_{\text{c.m.}}(\vec{R}_{\text{c.m.}})$ the wave function of the motion of the total center of mass (c.m.), and $\vec{\xi}_i$ the Jacobi coordinate with $i=1$ and 2 for cluster A and $i=4$ and 5 for cluster B, respectively. The symbol $\mathcal{A}$ describes the operation of antisymmetrization between quarks in two interacting clusters. S and T denote the total spin and isospin of the two-baryon system, respectively. Substituting Ψ into the projection equation

$$\langle\delta\Psi|(H-E)|\Psi\rangle=0, \tag{32}$$

where

$$E=E_A+E_B+E_{rel}, \tag{33}$$

with E, E_A, E_B, and E_{rel} being the total energy, the inner energies of clusters A and B, and relative energy between clusters A and B, respectively, we obtain RGM equation [32]

$$\int[\mathcal{H}(\vec{R}',\vec{R}'')-E\mathcal{N}(\vec{R}',\vec{R}'')]\chi_{rel}(\vec{R}'')d\vec{R}''=0, \tag{34}$$

where the Hamiltonian kernel $\mathcal{H}$ and normalization kernel $\mathcal{N}$ can, respectively, be calculated by

$$\left\{\begin{matrix}\mathcal{H}(\vec{R}',\vec{R}'')\\ \mathcal{N}(\vec{R}',\vec{R}'')\end{matrix}\right\}=\left\langle[\hat{\phi}_A(\vec{\xi}_1,\vec{\xi}_2)\hat{\phi}_B(\vec{\xi}_4,\vec{\xi}_5)\delta(\vec{R}-\vec{R}')\chi_{\text{c.m.}}(\vec{R}_{\text{c.m.}})]_{ST}\left|\left\{\begin{matrix}H\\ 1\end{matrix}\right\}\right|\mathcal{A}[\hat{\phi}_A(\vec{\xi}_1,\vec{\xi}_2)\hat{\phi}_B(\vec{\xi}_4,\vec{\xi}_5)\right.$$
$$\left.\times\delta(\vec{R}-\vec{R}'')\chi_{\text{c.m.}}(\vec{R}_{\text{c.m.}})]_{ST}\right\rangle. \tag{35}$$

To solve Eq. (34), one usually uses the variational method, in which the unknown χ_{rel} is expanded by a set of well-established basis functions,

$$\chi_{rel}(\vec{R})=\sum_{i=1}^{n}c_i u(\vec{R},\vec{S}_i), \tag{36}$$

and leave c_i's to be solved. The detailed methods of calculating the binding energy of the dibaryon, the baryon-baryon interaction, and the scattering cross section can be found in the Appendix and Refs. [32,34].

[1]It should be admitted that the parameters in Table I do not exactly give the experimental values of baryon masses. However, as indicated above, these parameters can basically reproduce many existent experimental data; thus, it will be safe to use them in the single-channel calculation. Of course, if a coupled channel calculation is considered, the sets of parameters should be taken with caution so that the thresholds of opened channels can precisely be produced.

III. RESULTS

In our previous paper [26], we have investigated possible dibaryon systems with various strangeness. According to our analysis, both the symmetry structure and the chiral-field-induced interaction between quarks play important roles in forming dibaryons. The symmetry structure is characterized

TABLE III. Binding energy E_b and corresponding rms radius $\mathcal{R}$ for $[\Sigma^*\Delta]_{(0,5/2)}$ and $[\Xi^*\Omega]_{(0,1/2)}$. The units for E_b and rms are in MeV and fm, respectively. The lower indices give the total spin and the isospin of the dibaryon.

$\mathcal{S}$	System	$\langle\mathcal{A}^{sfc}\rangle$	Set I $E_b//\mathcal{R}$	set II $E_b//\mathcal{R}$
-1	$[\Sigma^*\Delta]_{(0,5/2)}$	2	24.6//0.99	19.0//1.04
-5	$[\Xi^*\Omega]_{(0,1/2)}$	2	92.6//0.71	76.5//0.72

by the expectation value of the antisymmetrizer of the baryon-baryon system in the spin-flavor-color space $\langle\mathcal{A}^{sfc}\rangle$: Systems with $\langle\mathcal{A}^{sfc}\rangle\sim 0$ are forbidden states because of the Pauli blocking effect. Systems with $\langle\mathcal{A}^{sfc}\rangle\sim 1$ are deuteron-like states. The characters of these states can also be described by the models in the baryon-meson degrees of freedom, because quark-exchange effects in these states are of minor importance and the baryons are relatively independent to each other. Systems with $\langle\mathcal{A}^{sfc}\rangle\sim 2$ have a symmetry structure which favors binding. In these systems, the kinetic energy at some interbaryon separations might be smaller than that at an infinite separation because of large quark-exchange effects. If the chiral-field-induced interaction provides a not very strong attraction, loosely bound $\Delta\Delta$-like states might be formed, but if this attraction is substantial, deeply bound $\Omega\Omega$-like states might exist. States of this kind are of special interest because of their special symmetry structures. Further investigations of these states might provide us with more information regarding the roles of quarks in the baryon-baryon interactions.

We choose $[\Sigma^*\Delta]_{(0,5/2)}$ and $[\Xi^*\Omega]_{(0,1/2)}$, where the subscripts denote the spin and isospin of the systems, respectively, as the samples of $\Delta\Delta$-like and $\Omega\Omega$-like states, respectively.

A. Binding behaviors

For the convenience of the reader, we outline the binding behaviors of these two systems. The binding energies and the corresponding root-mean-square (rms) radii of $[\Sigma^*\Delta]_{(0,5/2)}$ and $[\Xi^*\Omega]_{(0,1/2)}$ with parameters sets I and II are tabulated in Table III. It can be seen that $[\Sigma^*\Delta]_{(0,5/2)}$ has a binding energy of 19.0–24.6 MeV and $[\Xi^*\Omega]_{(0,1/2)}$ of 76.5–92.6 MeV. A variation of the model parameters does not change their binding behaviors qualitatively. The binding energy of $[\Xi^*\Omega]_{(0,1/2)}$ is much larger than that of $[\Sigma^*\Delta]_{(0,5/2)}$, and the energy level of the former dibaryon is below the threshold of the strong-decay mode $\Xi\Omega\pi$ in either set of parameters. The relative wave functions of these systems with set I are presented in Fig. 1. It can be seen that these wave functions are qualitatively consistent with corresponding binding energies; namely, the larger the binding energy is, the narrower the distribution of the relative wave function becomes and the more inward the distribution moves.

The reason that $[\Xi^*\Omega]_{(0,1/2)}$ is much more bound than $[\Sigma^*\Delta]_{(0,5/2)}$ can be attributed to the roles of the quark-quark interactions, especially the chiral-field-induced interactions. This can quantitatively be shown by decomposing the contributions of various interaction terms in the total binding energies. The results are tabulated in Table IV. In the table, all the energies are relative to that of the two baryon clusters which are well separated. The negative sign means an repulsive contribution to the total binding energies. It can be seen that, in both systems, the kinetic energy is smaller than that of two freely moving baryons; namely, its contribution is relatively attractive. This fact is due to the large quark-exchange effect characterized by $\langle\mathcal{A}^{sfc}\rangle\sim 2$ [26]. The main attractive contributions are provided by scalar mesons, especially by the σ meson. Its contribution is more substantial in $[\Xi^*\Omega]_{(0,1/2)}$ than in $[\Sigma^*\Delta]_{(0,5/2)}$ and leads to a larger binding energy for $[\Xi^*\Omega]_{(0,1/2)}$. It should be mentioned that the importance of σ is also seen in our previous investigations [26]. Namely, the σ meson is always an important source of intermediate attractions in various baryon-baryon systems such as scattered N-N and N-Y systems, deuteron, H dibaryon, etc. The existence and the structure of the σ meson should surely be studied continuously.

B. Baryon-baryon effective interaction

As mentioned above, baryon-baryon bound states with $\langle\mathcal{A}^{sfc}\rangle\sim 2$ are QCD motivated. Their structures cannot be described in terms of the models in the baryon-meson degrees of freedom. The investigation of the baryon-baryon interactions induced by the quark-quark interactions is of interest, because it not only reflects the roles of quarks and

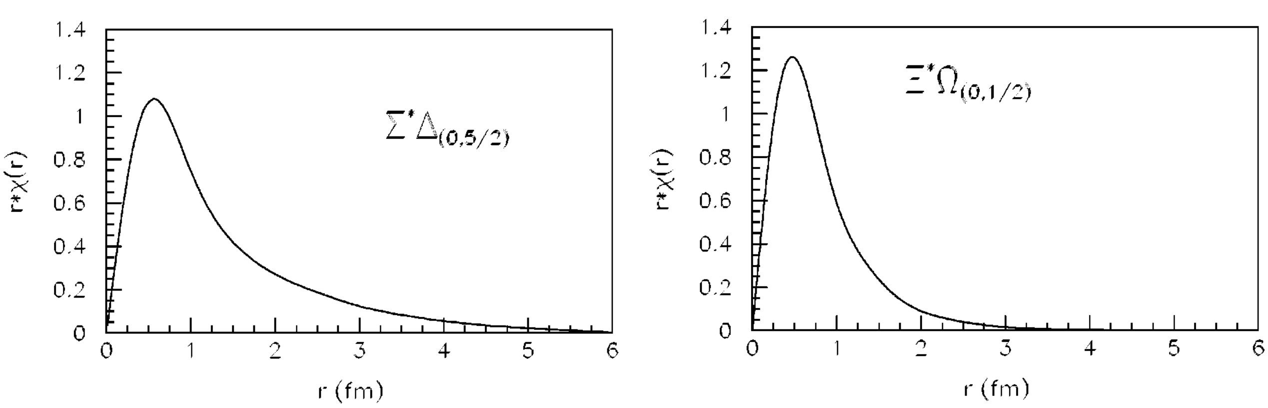

FIG. 1. S-wave relative wave functions of $[\Sigma^*\Delta]_{(0,5/2)}$ and $[\Xi^*\Omega]_{(0,1/2)}$ with set I.

TABLE IV. Contributions of various interaction terms in the binding energies of $[\Sigma^*\Delta]_{(0,5/2)}$ and $[\Xi^*\Omega]_{(0,1/2)}$. The energy is given in MeV.

	Kine.	OGE	π	K	η	η'	σ	σ'	κ	ϵ
$[\Sigma^*\Delta]_{(0,5/2)}$	10.1	−35.5	−14.9	−2.5	−2.6	−1.0	44.0	13.5	13.3	0.1
$[\Xi^*\Omega]_{(0,1/2)}$	9.1	−73.7	0.0	−13.7	0.1	−9.9	106.8	0.0	36.0	37.8

gluons in baryon-baryon interactions, it also provides us with a better understanding of the appearance of the large binding energies in such systems.

We use the equivalent potential $V(r)$, where r is the relative coordinate between the clusters, to describe the effective interaction between two clusters. We also employ a so-called ''adiabatic potential'' $V(S)$ to approximately present the general feature of the effective potential, although $V(S)$ is not a potential because S is not a coordinate of the relative motion between clusters.

To extract effective baryon-baryon interactions from quark-quark interactions, one should first transform Eq. (34) into a form in which the normalization kernel does not appear: namely, a normal Schrödinger equation. This can be achieved by using the orthogonal condition method (OCM) [32]. The obtained Schrödinger equation can be written as

$$\int \tilde{\mathcal{H}}(\vec{R},\vec{R}\,')\tilde{\chi}(\vec{R}\,')d^3\vec{R}\,' = E\tilde{\chi}(\vec{R}), \tag{37}$$

where

$$\tilde{\mathcal{H}}(\vec{R},\vec{R}\,') = \int \mathcal{N}^{-1/2}(\vec{R},\vec{R}\,'')\mathcal{H}(\vec{R}\,'',\vec{R}\,''') \times \mathcal{N}^{-1/2}(\vec{R}\,''',\vec{R}\,')d^3\vec{R}\,''d^3\vec{R}\,''', \tag{38}$$

$$\tilde{\chi}(\vec{R}\,') = \int \mathcal{N}^{1/2}(\vec{R}\,',\vec{R})\chi_{rel}(\vec{R})d^3\vec{R}. \tag{39}$$

The nonlocal baryon-baryon interaction $\mathcal{V}(\vec{R},\vec{R}\,')$ is then derived by substituting the contribution of the relative kinetic energy term from the OCM Hamiltonian kernel:

$$\mathcal{V}(\vec{R},\vec{R}\,') = \tilde{\mathcal{H}}(\vec{R},\vec{R}\,') - \mathcal{T}(\vec{R},\vec{R}\,'), \tag{40}$$

with

$$\mathcal{T}(\vec{R},\vec{R}\,') = \frac{\vec{P}^2_{\vec{R}}}{2\mu}\delta^3(\vec{R},\vec{R}\,'), \tag{41}$$

where μ is the reduced mass of interacting baryons and $\vec{P}_{\vec{R}}$ is the momentum operator of the relative motion of the system.

Consequently, we get the equivalent local baryon-baryon potential

$$V(\vec{R}) = \frac{\int \mathcal{V}(\vec{R},\vec{R}\,')\chi(\vec{R}\,')d\vec{R}\,'}{\chi(\vec{R})}. \tag{42}$$

On the other hand, the ''adiabatic potential'' expressed with the generating coordinate S is defined as

$$V(S_\alpha) = H_{\alpha\alpha} - \frac{(H_{nn} - T_{Dnn})}{N_{nn}}N_{\alpha\alpha} - T_{D\alpha\alpha}, \tag{43}$$

where generate coordinate method (GCM) matrices can be calculated by

$$\begin{Bmatrix} H_{ij} \\ N_{ij} \\ T_{Dij} \end{Bmatrix} = \langle \chi(\vec{R},\vec{S}_i)| \begin{Bmatrix} \mathcal{H}(\vec{R},\vec{R}\,') \\ \mathcal{N}(\vec{R},\vec{R}\,') \\ -\dfrac{\hbar^2}{2\mu}\delta^3(\vec{R}-\vec{R}\,')\nabla^2_{\vec{R}} \end{Bmatrix} |\chi(\vec{R}\,',\vec{S}_j)\rangle. \tag{44}$$

It should be noted that $V(S_\alpha)$ goes to zero when α approaches to the last point n where the interacting baryons are well separated.

In terms of set I, we plot the ''adiabatic potential'' $V(S)$ and the equivalent local potential $V(r)$ for the $\Xi^* - \Omega$ interaction in the $(S,T) = (0,1/2)$ state in Fig. 2. It shows that the shapes of $V(S)$ and $V(r)$ are very similar. Both of them have a depth of about 100 MeV and a minimum located around at $S \approx 1$ fm for $V(S)$ and $r \approx 0.8$ fm for $V(r)$, respectively. This similarity reflects that both $V(S)$ and $V(r)$ can qualitatively give the basic character of the baryon-baryon interaction at large distances. However, the characters of potentials at small values of variables are quite different: $V(S)$ has a finite value when $S \to 0$, while $V(r)$ is infinite when $r \to 0$. It should be mentioned that the repulsive core originates from the zero value of the S-wave relative wave function. In fact, the baryon-baryon interaction is nonlocal, because the baryon has an internal quark structure and, consequently, a finite size. This definitely affects the potential, especially when two baryons are very close to each other. Therefore, the baryon-baryon interaction can correctly be described only when the nonlocal effect is considered—for instance, the results in our previous investigations of Λ hypernuclei [27].

We present $V(S)$ in the $[\Sigma^*\Delta]_{(0,5/2)}$ system in Fig. 3. The basic character of $V(r)$ is similar to that of $V(S)$. It shows that a minimum of about -50 MeV is located at $S \approx 1.3$ fm, and the repulsion at short distance is stronger than that in the $[\Xi^*\Omega]_{(0,1/2)}$ case. This is consistent with the result that the binding energy of $[\Sigma^*\Delta]_{(0,5/2)}$ is much smaller than that of $[\Xi^*\Omega]_{(0,1/2)}$.

It should be emphasized again that although $V(S)$ depicts the basic feature of the baryon-baryon interaction, it cannot be regarded as a *true* potential in a rigorous sense, because the generator coordinate S is not a relative coordinate. It

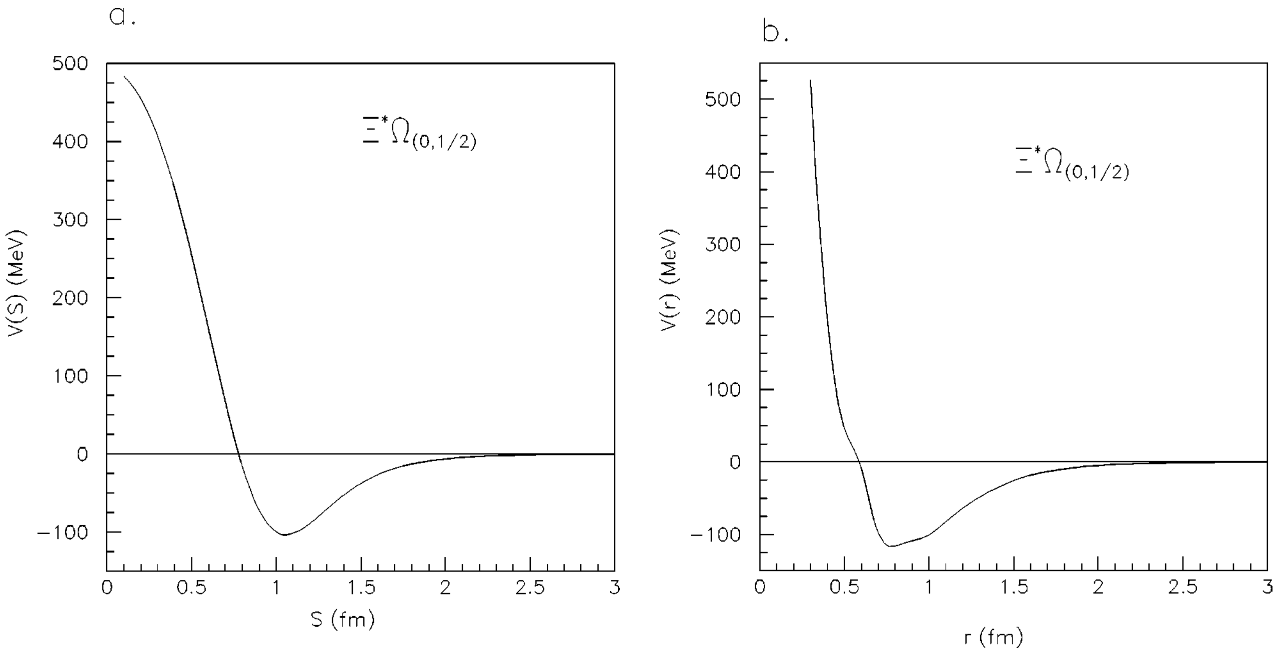

FIG. 2. Baryon-baryon interaction in $[\Xi^*\Omega]_{(0,1/2)}$ with set I: (a) "Adiabatic potential" in GCM; S is the generator coordinate. (b) The equivalent local potential.

actually reflects the dependence of the baryon-baryon potential on the distribution of the Gaussian basis set by which the trial wave function is expanded. On the other hand, $V(r)$ qualitatively reflects the dependence of the baryon-baryon interactions on the relative coordinate except the fact that the important nonlocal effect is missing due to the localization operation in calculating $V(r)$. The more rigorous baryon-baryon potential should be a nonlocal one and can explicitly be derived in the RGM. Anyway, $V(S)$ can be used to discuss the general feature of the baryon-baryon potential, especially at large distances.

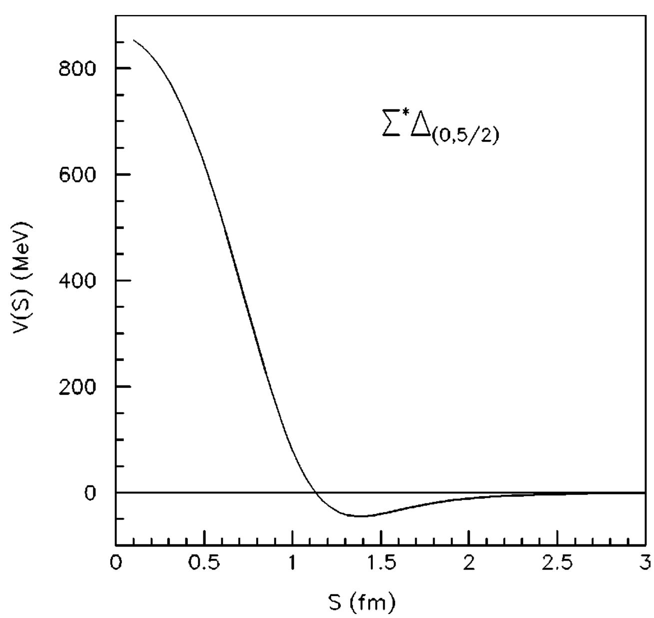

FIG. 3. "Adiabatic potential" in $[\Sigma^*\Delta]_{(0,5/2)}$ system with set I. S is the generator coordinate.

C. Total scattering cross sections

The scattering amplitude is calculated by summing over partial waves up to orbital angular momentum $L=2$. Inclusion of higher partial waves is necessary if a detailed comparison with data were to be made in the future. Moreover, the pure Coulomb contribution and interference with the nuclear amplitude are considered if both baryons are charged.

We first study $\Xi^*\Omega$ scattering in the $T=\frac{1}{2}$ channel which is closely related to the most interesting candidate of dibaryons $[\Xi^*\Omega]_{(0,1/2)}$. In isospin space, we have

$$|\Xi^{*0}\Omega\rangle=|(\Xi^*\Omega),T=1/2,T_3=1/2\rangle,$$

$$|\Xi^{*-}\Omega\rangle=|(\Xi^*\Omega),T=1/2,T_3=-1/2\rangle. \qquad (45)$$

Then, the scattering cross sections caused by the nuclear forces in the $\Xi^{*0}\Omega\to\Xi^{*0}\Omega$ and $\Xi^{*-}\Omega\to\Xi^{*-}\Omega$ processes are the same. The only difference between two entire cross sections comes from the Coulomb force and does not significantly affect the overall scattering results. The resultant total cross sections are plotted in Fig. 4.

Following, we study $\Sigma^*\Delta$ scattering in the $T=\frac{5}{2}$ channel where more information related to the possible dibaryon $[\Sigma^*\Delta]_{(0,5/2)}$ may be found. In the isospin basis

$$|\Sigma^{*+}\Delta^{++}\rangle=|(\Sigma^*\Delta),T=5/2,T_3=5/2\rangle,$$

$$|\Sigma^{*-}\Delta^{-}\rangle=|(\Sigma^*\Delta),T=5/2,T_3=-5/2\rangle. \qquad (46)$$

Similarly, the scattering cross sections in the $\Sigma^{*+}\Delta^{++}\to\Sigma^{*+}\Delta^{++}$ and the $\Sigma^{*-}\Delta^{-}\to\Sigma^{*-}\Delta^{-}$ processes should also be the same when Coulomb effects are disregarded. Again, the Coulomb effect is of minor importance for the total cross sections. The predicted total cross sections are

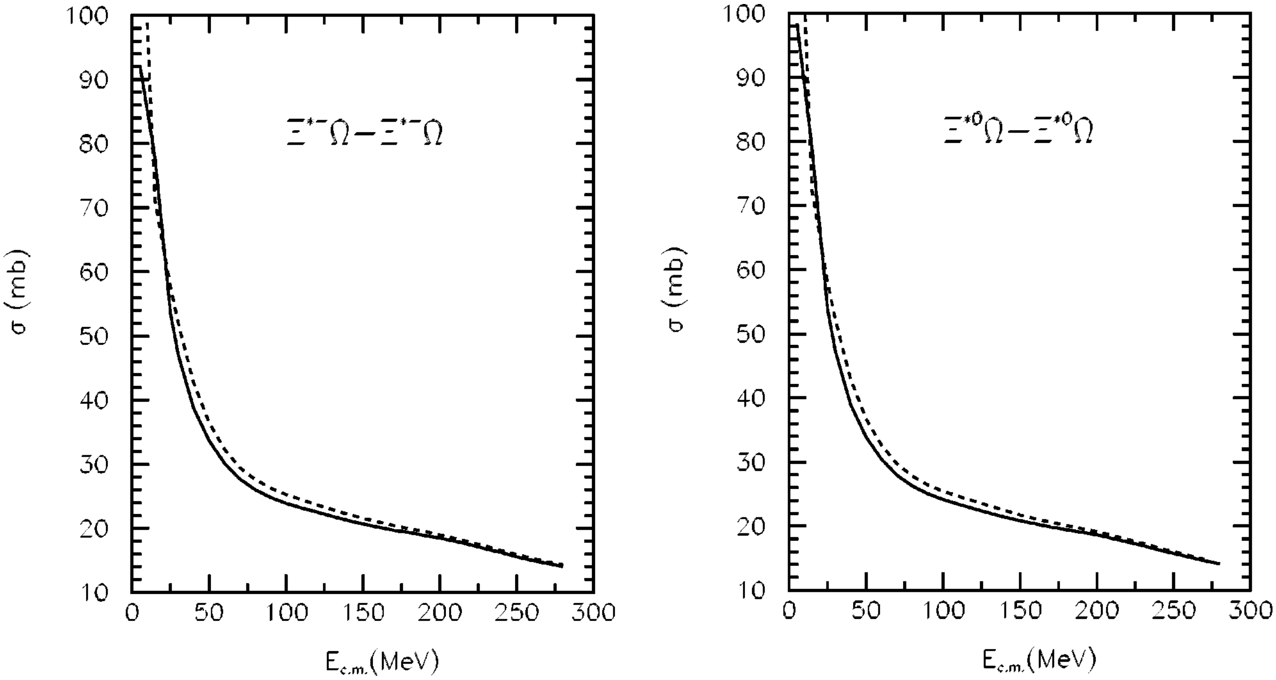

FIG. 4. Scattering cross section for $\Xi^*\Omega\rightarrow\Xi^*\Omega$. The solid and dashed curves denote the results with set I and set II, respectively.

plotted in Fig. 5. From Fig. 5, it can be seen that the parameter dependence of the total cross sections is also very small.

It should be mentioned that the presented cross sections are not accurate enough at higher energies due to the lack of contributions from higher partial waves and disregarding decay channels. However, as a preliminary investigation, a single-channel calculation with $L\leqslant 2$ waves only would be enlightening enough, especially at lower energies before data become available. Although it is rather difficult to measure decuplet baryon-baryon scatterings at this moment, the quark model may provide a general picture of all possible decuplet baryon-baryon scatterings.

IV. CONCLUSIONS

In terms of the chiral SU(3) quark model, we investigate the binding energies, the effective baryon-baryon interactions, and the total scattering cross sections of two-decuplet-

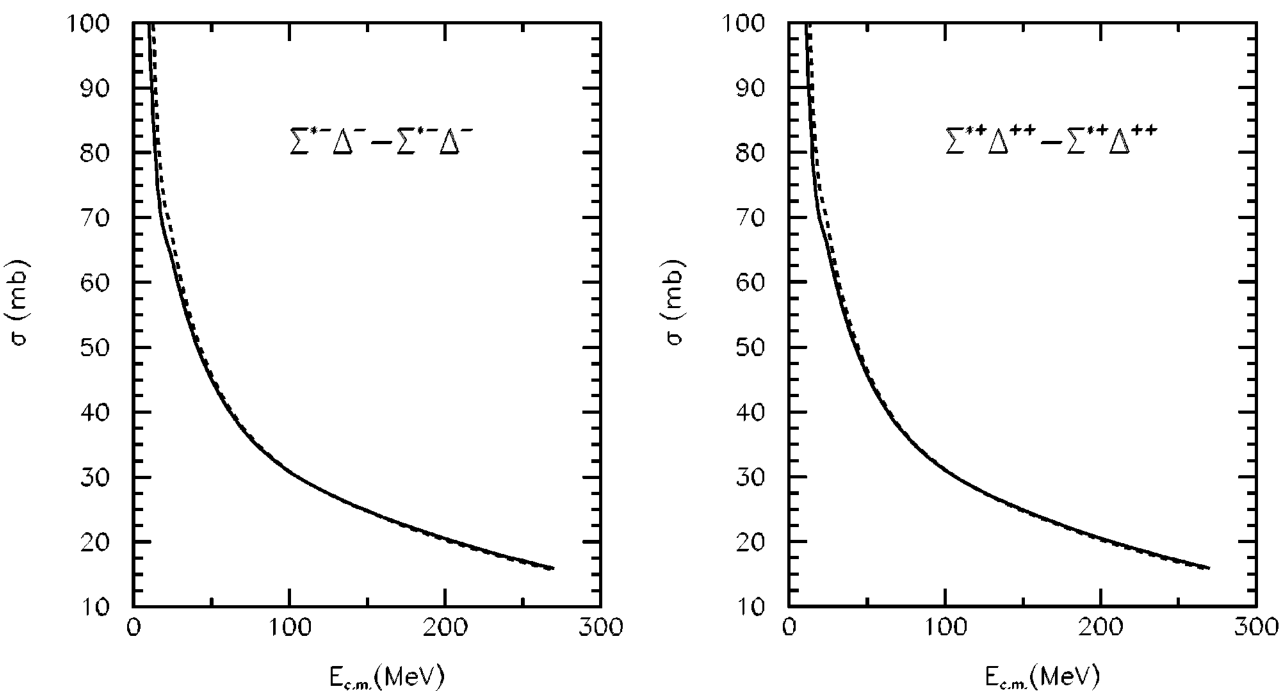

FIG. 5. Scattering cross section of $\Sigma^*\Delta\rightarrow\Sigma^*\Delta$. The solid and dashed curves denote the results with set I and set II, respectively.

baryon systems $\Sigma^*\Delta(s=-1)$ and $\Xi^*\Omega(s=-5)$. The reason for choosing these two systems is to study the strong effect of antisymmetrization which might lead to an important quark exchange effect in the baryon-baryon interactions.

In our assertion [26], two-baryon bound states with $\langle \mathcal{A}^{sfc}\rangle \sim 2$, in terms of their interaction characters, can be divided into two classes: the $\Delta\Delta$-like state such as $[\Sigma^*\Delta]_{(0,5/2)}$ and the $\Omega\Omega$-like state such as $[\Xi^*\Omega]_{(0,1/2)}$. The main characters of these states are in the following: (1) The binding energy is 19.0–24.6 MeV for $[\Sigma^*\Delta]_{(0,5/2)}$ and 76.5–92.4 MeV for $[\Xi^*\Omega]_{(0,1/2)}$. (2) The kinetic energy term contributes to attraction if we compare with two free baryons at infinite separation. (3) Scalar chiral fields—for instance, the σ field—provide a larger attraction in $[\Xi^*\Omega]_{(0,1/2)}$ than in $[\Sigma^*\Delta]_{(0,5/2)}$. (4) $[\Xi^*\Omega]_{(0,1/2)}$ might have a narrow width, because its binding energy is so large that the system is below the threshold for the strong-decay mode $\Xi\Omega\pi$, while $[\Sigma^*\Delta]_{(0,5/2)}$ might be a resonance with a broad width. Moreover, the *S*-wave relative wave functions which qualitatively reflect the binding energies of the corresponding systems are also given.

The effective baryon-baryon interaction are calculated. The "adiabatic potential" $V(S)$ is quite similar to the equivalent local potential $V(r)$ except for the repulsive core at small distances. Both $V(S)$ and $V(r)$ can give the general picture of the baryon-baryon interactions. The basic feature of $V(S)$ reflects the fact that the binding energy of the $[\Sigma^*\Delta]_{(0,5/2)}$ system is much smaller than that of the $[\Xi^*\Omega]_{(0,1/2)}$ system.

In addition, total scattering cross sections of these systems are calculated. It can be seen that the model parameter dependence is weak and the Coulomb interaction only affects the total cross sections slightly. These results might give us a preliminary impression of the scattering processes.

It should be mentioned that as a result of the great importance of the quark structure in the considered systems, the corresponding baryon-baryon interactions cannot be derived in the baryon-meson degrees of freedom. The resultant baryon-baryon interactions must reflect the effects from the quark degree of freedom. Of course, many factors may affect the results of the baryon-baryon interactions: for instance, higher partial waves and the coupled channel effects have not been taken into account. However, because the chiral SU(3) quark model is quite successful in reproducing most of existent data, especially in the strange sector, such as the *N*-Y scattering [27] and binding energy of the H particle [13], we believe that the results for the systems with strangeness $s=-1$ and -5 are reliable. We hope our work will initiate experimental work to investigate the two-baryon systems considered here.

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APPENDIX: KOHN-HULTHÉN-KATO VARIATIONAL METHOD

The RGM equation (34) can be solved by using the Kohn-Hulthén-Kato variational method [32,34]. In this method, the unknown χ_{rel} is expanded by a set of well-established basis functions

$$\chi_{rel}(\vec{R})=\sum_{i=1}^{n} c_i u(\vec{R},\vec{S}_i) \tag{A1}$$

and leave c_i's to be solved. In this paper, $u(\vec{R},\vec{S}_i)$ is chosen in the following form: For a bound-state problem, a local peaked Gaussian function is employed:

$$\begin{aligned} u_b(\vec{R},\vec{S}_i) &\equiv u(\vec{R},\vec{S}_i)\\ &=\left(\frac{\omega\mu_{AB}}{\pi}\right)^{3/4}\exp\left[-\frac{\omega\mu_{AB}}{2}(\vec{R}-\vec{S}_i)^2\right], \end{aligned}\tag{A2}$$

where μ_{AB} is the reduced mass of the two-baryon system and $\vec{S}_i$ is the generate coordinate. This function can also be expanded by partial waves:

$$u(\vec{R},\vec{S}_i)=\sum_L \frac{1}{R}u^L(R,S_i)Y_{LM}(\hat{R}), \tag{A3}$$

with the Lth partial wave function

$$\begin{aligned} u^L(R,S_i)&=4\pi R\left(\frac{\omega\mu_{AB}}{\pi}\right)^{3/4}\exp\left[-\frac{\omega\mu_{AB}}{2}(R^2+S_i^2)\right]\\ &\quad\times i_L(\omega\mu_{AB}RS_i) \end{aligned}\tag{A4}$$

and the modified spherical Bessel function i_L. For a scattering problem, the function $u(\vec{R},\vec{S}_i)$ is somehow complicated. The Lth partial wave function can be written as

$$u^L_{scat}(R,S_i)=\begin{cases} u^L(R,S_i), & R\leqslant R_c,\\ [h_L^-(k_{AB}R)-s_i h_L^+(k_{AB}R)], & R>R_c, \end{cases}\tag{A5}$$

with $h_L^{\pm}$ being Lth spherical Hankel functions, $k_{AB}=\sqrt{2\mu_{AB}E_{rel}}$ the momentum of relative motion, s_i the *S*-matrix elements and R_c is a cutoff radius beyond which all the strong interactions can be disregarded. Performing variational procedure, one can deduce a Lth partial-wave equation for the bound-state problem,

$$\sum_{i=j}^{n} \mathcal{L}^L_{ij}c_j=0, \quad (i=1,\ldots,n), \tag{A6}$$

and a Lth partial-wave equation for the scattering problem,

$$\sum_{j=1}^{n} \tilde{\mathcal{L}}^L_{ij}c_j=\mathcal{M}_i \quad (i=1,\ldots,n), \tag{A7}$$

with

$$\tilde{\mathcal{L}}^L_{ij}=\tilde{\mathcal{K}}^L_{ij}-\tilde{\mathcal{K}}^L_{i0}-\tilde{\mathcal{K}}^L_{0j}+\tilde{\mathcal{K}}^L_{00}, \tag{A8}$$

$$\tilde{\mathcal{M}}^L_i=\tilde{\mathcal{K}}^L_{00}-\tilde{\mathcal{K}}^L_{i0}, \tag{A9}$$

$$\tilde{\mathcal{K}}^L_{ij}=\mathcal{L}^L_{ij}-\mathcal{K}^{L(ex)}_{ij}, \tag{A10}$$

$$\mathcal{L}^{L}_{ij}=\int u^{L}(R',S_i)[\mathcal{H}(R',R)-E\mathcal{N}(R',R)]\times u^{L}(R,S_j)R'RdR'dR, \tag{A11}$$

$$\mathcal{K}^{L(ex)}_{ij}=\int_{R_c}^{\infty}u^{L}(R',S_i)\times\left(-\frac{\hbar^2}{2\mu_{AB}}\frac{d^2}{dR^2}+\frac{\hbar^2}{2\mu_{AB}}\frac{L(L+1)}{R^2}-E_{rel}\right)\times u^{L}(R,S_j)R'RdR'dR, \tag{A12}$$

$$\begin{Bmatrix}\mathcal{H}(R',R)\\ \mathcal{N}(R',R)\end{Bmatrix}=\int Y^{*}_{LM}(\hat{R}')\begin{Bmatrix}\mathcal{H}(\vec{R}\,',\vec{R})\\ \mathcal{N}(\vec{R}\,',\vec{R})\end{Bmatrix}\times Y_{LM}(\hat{R})d\hat{R}\,'d\hat{R}. \tag{A13}$$

Solving Eq. (A6), one can obtain the binding energy and the wave function of the baryon-baryon system. While solving Eq. (A7), one gets the wave function and the *S*-matrix elements of the scattered system and, consequently, phase shifts and scattering cross sections.

[1] S. Godfrey and N. Isgur, Phys. Rev. D **32**, 189 (1985); S. Capstick and N. Isgur, *ibid.* **34**, 2809 (1986); W. Lucha, F. F. Schoberl, and D. Gromes, Phys. Rep. **200**, 127 (1991); J. M. Richard, *ibid.* **212**, 1 (1992).

[2] E. Witten, Nucl. Phys. **B223**, 422 (1983); **B223**, 433 (1983); I. Zahed and G. E. Brown, Phys. Rep. **142**, 1 (1986).

[3] R. L. Jaffe, Phys. Rev. Lett. **38**, 195 (1977).

[4] P. J. G. Mulders *et al.*, Phys. Rev. Lett. **40**, 1543 (1978); Phys. Rev. D **21**, 2653 (1980).

[5] P. Gonzales, P. LaFrance, and E. Lomon, Phys. Rev. D **35**, 2142 (1987).

[6] Yu. S. Kalashnikova *et al.*, Sov. J. Nucl. Phys. **46**, 689 (1987).

[7] B. Sivestre-Brac and J. Leandri, Phys. Rev. D **45**, 4221 (1992); **47**, 5083 (1993), and references therein.

[8] K. Yazaki, Suppl. Prog. Theor. Phys. **91**, 146 (1987).

[9] F. J. Dyson and N. H. Xuong, Phys. Rev. Lett. **13**, 815 (1964); Y. Kamaen and T. Fujita, *ibid.* **38**, 471 (1977); M. Cvetič *et al.*, Phys. Lett. **93B**, 489 (1980); N. R. Walet, Phys. Rev. C **48**, 2222 (1993); F. Wang, J. L. Ping, G. H. Wu, L. J. Teng, and T. Goldman, *ibid.* **51**, 3411 (1995).

[10] X. Q. Yuan, Z. Y. Zhang, Y. W. Yu, and P. N. Shen, Phys. Rev. C **60**, 045203 (1999).

[11] R. Bilger *et al.*, Phys. Lett. **26B**, 247 (1991); R. Bilger and H. A. Clement, Phys. Rev. Lett. **71**, 42 (1993); George Wagner, L. Ya. Glozman, A. J. Buchmann, and Amand Faessler, Nucl. Phys. **A594**, 263 (1995); A. J. Buchmann, G. Wagner, and Amand Faessler, Phys. Rev. C **57**, 3340 (1998).

[12] A. P. Balachandran *et al.*, Phys. Rev. Lett. **52**, 887 (1984); U. Straub, *et al.*, Phys. Lett. B **200**, 241 (1988); M. Oka, K. Shimizu, and K. Yazaki, Phys. Lett. **130B**, 365 (1983); N. Kodama, M. Oka, and T. Hatsuda, Nucl. Phys. **A580**, 445 (1994); P. H. Pile *et al.*, Nucl. Instrum. Methods Phys. Res. A **321**, 48 (1992).

[13] P. N. Shen, Z. Y. Zhang, Y. W. Yu , X. Q. Yuan, and S. Yang, J. Phys. G **25**, 1807 (1999); Chin. Phys. Lett. **17**, 7 (2000).

[14] K. Yazaki, *Hadrons and Nuclei with Strangeness*, Proceedings of the First Sino-Japan Symposium on Strangeness Physics (World Scientific, Singapore, 1998).

[15] K. Imai, in *Hadrons and Nuclei with Strangeness* [14].

[16] C. W. Wong, Phys. Rev. C **57**, 1962 (1998), and references therein.

[17] J. Schaffner-Bielich, R. Mattiello, and H. Sorge, nucl-th/9908043.

[18] V. G. J. Stoks and T. A. Rijken, Phys. Rev. C **59**, 3009 (1999).

[19] P. N. Shen *et al.*, *Proceedings of Hadron'99*, Beijing, China, 1999 [Nucl. Phys. **A675**, 234c (2000)].

[20] See, e.g., C. B. Dover and G. E. Walker, Phys. Rep. **89**, 1 (1982), and references therein.

[21] See, e.g., B. Shiva Humar, Nucl. Phys. **A590**, 29c (1995), and references therein.

[22] See, e.g., G. Baym, Nucl. Phys. **A590**, 233c (1995), and references therein.

[23] *Proceedings of the International Symposium on Strangeness in Quark Matter*, 1998 [J. Phys. G **251**, 1 (1998)].

[24] T. Goldman *et al.*, Phys. Rev. Lett. **59**, 627 (1987).

[25] Y. W. Yu, Z. Y. Zhang, and X. Q. Yuan, Commun. Theor. Phys. **31**, 1 (1999); High Energy Phys. Nucl. Phys. **23**, 859 (1999).

[26] Q. B. Li, P. N. Shen, Z. Y. Zhang, and Y. W. Yu, Nucl. Phys. **A683**, 487 (2001).

[27] M. J. Zhang, P. N. Shen, X. H. Lin and F. X. Kong, High Energy. Phys. Nucl. Phys. **25**, 743 (2001).

[28] P. N. Shen, Y. B. Dong, Z. Y. Zhang, Y. W. Yu, and T.-S. H. Lee, Phys. Rev. C **55**, 2024 (1997); H. Chen and Z. Y. Zhang, High Energy Phys. Nucl. Phys. **20**, 379 (1996).

[29] Z. Y. Zhang, Y. W. Yu, P. N. Shen, L. R. Dai, Amand Faessler, and U. Straub, Nucl. Phys. **A625**, 59 (1997); S. Yang, P. N. Shen, Z. Y. Zhang, and Y. W. Yu, *ibid.* **A635**, 146 (1998).

[30] H. Georgi, *Weak Interactions and Modern Particle Theory* (Benjamin/Cummings, New York, 1984).

[31] I. T. Obukhovsky and V. G. Kusainov, Phys. Lett. B **238**, 142 (1990); A. Buchmann, E. Fernandez, and K. Yazaki, *ibid.* **269**, 35 (1991); E. M. Henley and G. A. Miller, *ibid.* **251**, 453 (1991).

[32] K. Wildermuth and Y. C. Tang, *A Unified Theory of the Nucleus* (Academic, New York, 1977); U. Straub, J. Nieves, Amand Faessler, and E. Oset, Nucl. Phys. **A556**, 531 (1993); John M. Blatt and L. C. Biedenharn, Rev. Mod. Phys. **24**, 259 (1952).

[33] Tadashi Fujita, Yoshikazu Fujiwara, Choki Nakamoto, and Yasuyuki Suzuki, Prog. Theor. Phys. **100**, 931 (1998); Yoshikazu Fujiwara, Tadashi Fujita, Choki Nakamoto, and Yasuyuki Suzuki, *ibid.* **100**, 957 (1998); Yoshikazu Fujiwara, Michio Kohno, Tadashi Fujita, Choki Nakamoto, and Yasuyuki Suzuki, *ibid.* **103**, 755 (2000).

[34] M. Kamimura, Suppl. Prog. Theor. Phys. **62**, 236 (1977).