Flavor-octet dibaryons in the quark model

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(Received 11 February 1988)

The possibility of flavor-SU(3)-octet dibaryon resonances is studied in the quark model. Because the color-magnetic gluon-exchange interaction is attractive for the flavor octets as well as the singlet, bound and/or resonance states are expected. The importance of two-baryon channel couplings and quark antisymmetrization is stressed for low-lying bound/resonance states. The dibaryon spectrum is studied by applying the quark cluster model to the $J=1^+$ and 2^+ six-quark states. Several $J=2^+$ dibaryon resonances are predicted.

I. INTRODUCTION

Quantum chromodynamics (QCD) has not completely been understood in the low-energy region, although it is established as a part of the standard model for particles. It is hard to solve QCD reliably and thus many effective low-energy models have been proposed. The quark model is one of them, where the gluon field is replaced by a potential,[1] or a bag boundary condition,[2] etc. There are two main effects of the gluon field: It is believed to provide a quark confinement. A bag boundary or a confinement potential represents this effect in the effective models. There is also a perturbative gluon-exchange effect, which explains fine structures of the hadron spectrum. Among various terms of the gluon-exchange interaction, the color-magnetic interaction (CMI) plays the most important role. It gives, for instance, in the baryon spectrum, the octet-decuplet mass difference, $\Lambda-\Sigma$ mass difference, nonzero charge radius of the neutron by breaking the SU(6) symmetry, and D state mixing in some baryon excited states, etc.[3] The magnitude of CMI is typically 300 MeV, which is the mass difference of the nucleon (N) and the delta (Δ).

The origin of the hadron-hadron interaction is of great interest. The role of the confinement, quark antisymmetrization and the perturbative gluon exchange have been studied by several different approaches.[4,5] The potential quark model has been applied to the nuclear force, on which we have the most information. It was shown that CMI with quark antisymmetrization explains the strong short-range repulsion.[5] The nuclear force, however, has unfortunately a disadvantage in this respect. It consists of two major pieces canceling with each other, i.e., strong short-distance repulsion and a medium-range attraction. The former is likely to come from the substructure of the nucleon, while the latter is predominantly a meson-exchange effect. Quark model descriptions of the nuclear force are thus obscured by the ambiguity in describing meson-exchange effects in this language.

CMI, however, is not always repulsive, but in some $B=2$ states, gives an attractive effective interaction.[6] $\Delta\Delta$ with the spin $J=3$ and isospin $I=0$, for example, may have an S-wave bound state. Because of the absence of the short-range repulsion, such a bound state will be deeply bound, compact "six-quark"-like object, where the baryons Δ lose their identity. It is totally different from the deuteron, the only observed "dibaryon" bound state. Another possibility is dibaryons with strangeness. In 1977, Jaffe pointed out that CMI is strongly attractive in the flavor-SU(3)-singlet six-quark system.[7] The state couples the two baryon systems: $\Lambda\Lambda-N\Xi-\Sigma\Sigma$ with $I=0$ and $J=0$. The dibaryon states are of great importance in confirming the role of CMI and quark antisymmetrization in multihadron systems. The observation of the dibaryon state would also distinguish different low-energy models, which mainly differ in quark confinement mechanism but give a similar prediction for the hadron spectrum.

The Skyrme model of the baryon[8] is another low-energy model of QCD. There confinement is replaced by a topological stability of the classical soliton solution and the color-magnetic gluon exchange is equivalent to the semiclassically quantized rotation of the soliton. The baryon spectrum is fairly well described, as are many static properties of the nucleon.[9] It is quite interesting to see that the dibaryon spectrum is also similar to the quark model.[10]

In this paper we examine flavor-SU(3)-octet six-quark systems, where CMI is attractive as in the singlet system. In order to study resonance states, it is important to know the two-baryon configuration of the flavor-octet states. In Sec. II, possible quantum numbers and the two-baryon configuration for each flavor-octet state are presented. It is shown that couplings of several two-baryon channels and quark antisymmetrization are essential. We argue that, in general, $J=2^+$ resonances are more likely than $J=1^+$ ones.

In order to study the dibaryon spectrum, we employ the quark-cluster model, which takes into account the full quark antisymmetrization and the coupling of two-baryon channels. In Sec. III we present results of numerical calculations of baryon-baryon scattering phase shifts after a brief review of the quark-cluster model. We show that all of the $J=2^+$ states exhibit a quasibound state or a resonance. Conclusions are given in Sec. V.

II. FLAVOR-OCTET SIX-QUARK STATES

The effective baryon-baryon interaction is sensitive to the color-magnetic interaction, which is proportional to $-(\lambda_i\cdot\lambda_j)(\sigma_i\cdot\sigma_j)$. Suppose that all the six (valence) quarks are in the ground state in a bag or a potential well.

Then the orbital wave function is totally symmetric and therefore the spin-color symmetry and the flavor symmetry are conjugate with each other in order to make a totally antisymmetric state. Eigenvalues of the operator

$$\Gamma_{\rm CM} \equiv \sum_{i<j} -(\lambda_i \cdot \lambda_j)(\sigma_i \cdot \sigma_j)$$

are determined by the spin-color symmetry and the total spin of the state:

$$\langle \Gamma_{\rm CM} \rangle = 48 - \tfrac{1}{2} C[\mathrm{SU}(6)] + \tfrac{4}{3} J(J+1) , \quad (1)$$

where $C[\mathrm{SU}(6)]$ is the spin-color-SU(6) Casimir operator and J is the total spin. CMI is most attractive in the flavor-singlet, [222], symmetric, state. The corresponding spin-color symmetry is [33], where only $J=0$ is allowed for the color-singlet state, and $\langle \Gamma_{\rm CM} \rangle = -24$. The flavor-singlet $J=0$ state is known as the H dihyperon with strangeness $\mathcal{S}=-2$ and $I=0$ (Ref. 7). In the two-baryon (BB) configuration, H corresponds to a combination of $\Lambda\Lambda - N\Xi - \Sigma\Sigma$:

$$| H \rangle = \frac{1}{\sqrt{8}} (| \Lambda\Lambda \rangle + \sqrt{4} | N\Xi \rangle - \sqrt{3} | \Sigma\Sigma \rangle) . \quad (2)$$

Several different calculations[11,12] suggest either a bound state or a low-lying resonance in the Λ-Λ system and experimental searches are going on.[13]

The next low-lying state is labeled by the flavor symmetry [321], the octet. The spins $J=1$ and 2 are allowed. The flavor octet consists of strangeness $\mathcal{S}=-1$, $I=\frac{1}{2}$; $\mathcal{S}=-2$, $I=0$; $\mathcal{S}=-2$, $I=1$; and $\mathcal{S}=-3$, $I=\frac{1}{2}$. Tables I and II summarize the two-baryon contents of these flavor-octet states with $J=1^+$ and 2^+, respectively. They are obtained by diagonalizing the $\mathrm{SU}(3)_f$ Casimir operator within the given spin-isospin-strangeness space. Note that the $J=2$ states are given only by combinations of an octet and a decuplet baryon, while the $J=1$ states are also made of two octet baryons. The eigenvalue of the color-magnetic operator $\Gamma_{\rm CM}$ is $-12+\frac{4}{3}J(J+1) = -\frac{28}{3}$ for $J=1$ and -4 for $J=2$. All the other $\mathrm{SU}(3)_f$ eigenstates have non-negative eigenvalues of $\Gamma_{\rm CM}$, and therefore CMI tends to be repulsive.

The two baryon systems which participate in the flavor-octet states may exhibit a bound or a resonance state. For instance, $N\Omega$ with $\mathcal{S}=-3$, $I=\frac{1}{2}$, and $J=2$ has been discussed in this context.[14] Two general comments are in order. The flavor-octet states are given as particular combinations (given in Tables I and II) of the BB systems. The BB channel coupling should be substantial in the resonances. The threshold energies of those channels, however, are different due to the flavor-

TABLE I. Flavor-octet dibaryon states with $J=1^+$.

$\mathcal{S}$	I	BB
-3	$\frac{1}{2}$	$\frac{1}{3}(\sqrt{2}\,|\Lambda\Sigma\rangle - \sqrt{2}\,|\Sigma\Xi\rangle$
		$-\sqrt{2}\,|N\Omega\rangle + |\Lambda\Xi^*\rangle + |\Sigma\Xi^*\rangle + |\Sigma^*\Xi\rangle)$
-2	0	$\frac{1}{3}(\sqrt{4}\,|N\Xi\rangle + \sqrt{2}\,|N\Xi^*\rangle - \sqrt{3}\,|\Sigma\Sigma^*\rangle)$
-2	1	$(1/\sqrt{27})(-\sqrt{4}\,|N\Xi\rangle + \sqrt{8}\,|\Sigma\Sigma\rangle$
		$+\sqrt{2}\,|N\Xi^*\rangle - \sqrt{3}\,|\Lambda\Sigma^*\rangle - \sqrt{2}\,|\Sigma\Sigma^*\rangle - \sqrt{8}\,|\Delta\Xi\rangle)$
-1	$\frac{1}{2}$	$\frac{1}{3}(-\sqrt{2}\,|N\Lambda\rangle + \sqrt{2}\,|N\Sigma\rangle + |N\Sigma^*\rangle - \sqrt{4}\,|\Delta\Sigma\rangle)$

symmetry breaking and also the octet-decuplet mass difference. If their thresholds are distributed in a wide energy range, a bound or a low-lying resonance state around the threshold of the lowest BB state is unlikely. Thus, we expect that $J=1^+$ dibaryon resonances are less likely than $J=2^+$ ones because of the large octet-decuplet splitting.

The lowest BB state of $J=2^+$ dibaryons is not stable. Except a deeply bound state, they will decay by emitting a pion and/or by converting into an octet-octet system by the tensor force. The former is expected to be similar to the decay of the decuplet baryon, Σ^* and Ξ^*. The widths of those baryons are 36 and 10 MeV, respectively. The latter may give a similar width, although the tensor force between quarks is weak.

III. QUARK-CLUSTER MODEL

In order to estimate the masses of the dibaryon resonances, we employ the quark-cluster model approach.[5] Advantages of this approach are that it takes into account the full antisymmetrization of quarks and that connection of the six-quark state with the BB scattering states is unambiguous. The six-quark state is represented by a cluster wave function, defined by an antisymmetrized product of the baryon internal wave function ϕ_B of the two baryons (clusters) and the relative wave function χ:

$$\Phi(1,2,3,4,5,6) = \mathcal{A}[\phi_B(1,2,3)\phi_B(4,5,6)\chi(\mathbf{R}_{123-456})] . \quad (3)$$

For a given quark Hamiltonian and a fixed ϕ_B, the Schrödinger equation can be written as a nonlocal equation for χ, called the resonating-group-method equation:[15]

$$\int d\mathbf{R}'[H(\mathbf{R},\mathbf{R}') - EN(\mathbf{R},\mathbf{R}')]\chi(\mathbf{R}') = 0 , \quad (4)$$

TABLE II. Flavor-octet dibaryon states with $J=2^+$.

$\mathcal{S}$	I	BB
-3	$\frac{1}{2}$	$(1/\sqrt{5})(\sqrt{2}\,|N\Omega\rangle - |\Lambda\Xi^*\rangle - |\Sigma\Xi^*\rangle + |\Sigma^*\Xi\rangle)$
-2	0	$(1/\sqrt{5})(\sqrt{2}\,|N\Xi^*\rangle - \sqrt{3}\,|\Sigma\Sigma^*\rangle)$
-2	1	$(1/\sqrt{15})(\sqrt{2}\,|N\Xi^*\rangle - \sqrt{3}\,|\Lambda\Sigma^*\rangle - \sqrt{2}\,|\Sigma\Sigma^*\rangle + \sqrt{8}\,|\Delta\Xi\rangle)$
-1	$\frac{1}{2}$	$(1/\sqrt{5})(|N\Sigma^*\rangle + \sqrt{4}\,|\Delta\Sigma\rangle)$

where $H(N)$ is the Hamiltonian (normalization) integral kernel. The effective baryon-baryon interaction is nonlocal due to the antisymmetrization. The coupled-channel formalism gives a similar equation, where the Hamiltonian kernel as well as the normalization one has off-diagonal components. By solving Eq. (4) for a relevant boundary condition, we obtain bound states and scattering phase shifts of a two-baryon system.

For the choice of the Hamiltonian, we follow the previous work on the hyperon-nucleon and hyperon-hyperon interaction.[11] The Hamiltonian consists of the kinetic- and potential-energy terms. The kinetic energy is the nonrelativistic one with an effective quark mass m_q. The potential energy contains a two-body linear confinement force and a static one-gluon-exchange potential of the Fermi-Breit form. Breaking of the flavor-$SU(3)_f$ symmetry is taken into account in the contact two-body interaction in the gluon-exchange potential:

$$V_c(i,j)=\xi_{ij}(1+\tfrac{2}{3}\sigma_i\cdot\sigma_j)\delta(\mathbf{r}_{ij}) \, , \tag{5}$$

where $\xi_{ij}=1$ when both i and j are u or d, $\xi_{ij}=\xi_1$ when either i or j is s, and $\xi_{ij}=\xi_2$ when both i and j are s quarks. The parameters ξ_1 and ξ_2 are determined so as to reproduce the baryon spectrum and therefore the correct threshold for various BB channels. Mass differences of u,d, and s quarks are not considered here. They appear in two places in the Hamiltonian: the mass term and the kinetic-energy term $\mathbf{p}^2/2m_q$. The former only gives a constant shift of the total energy and does not contribute to the two-baryon interaction. The latter contributes mainly to the internal energy of each baryon giving again a constant shift for a given strangeness and may not be so important for the two-baryon interaction. Since the introduction of different masses makes the quark-cluster-model calculation complicated, we will neglect symmetry breaking in the kinetic-energy term.

TABLE III. Masses of the octet and decuplet baryons. A is a matrix element given for the present parameters by $A=73$ MeV. $M_0=866$ MeV and $\Delta_s=227$ MeV are chosen to fit the N and Λ masses. The numbers with an asterisk are fitted and those in parentheses are observed values.

		(MeV)	(MeV)
N	M_0+A	939*	(939)
Λ	$M_0+\Delta_s+(-1+2\xi_1)A$	1116*	(1116)
Σ	$M_0+\Delta_s+(\frac{5}{3}-\frac{2}{3}\xi_1)A$	1183	(1193)
Ξ	$M_0+2\Delta_s+(-\frac{2}{3}\xi_1+\frac{5}{3}\xi_2)A$	1325	(1318)
Δ	M_0+5A	1232*	(1232)
Σ^*	$M_0+\Delta_s+(\frac{5}{3}+\frac{10}{3}\xi_1)A$	1375	(1385)
Ξ^*	$M_0+2\Delta_s+(\frac{10}{3}\xi_1+\frac{5}{3}\xi_2)A$	1518	(1530)
Ω	$M_0+3\Delta_s+(5\xi_2)A$	1660	(1672)

The parameters for the present calculation are chosen as follows. The quark mass, which is now assumed to be the same for all quarks, is set to be $\frac{1}{3}$ of the average octet-baryon mass, i.e., $m_q=383.7$ MeV. The size parameter b of the baryon wave function is arbitrarily chosen as $b=0.5$ fm. The quark-gluon coupling constant α_s is determined from the $N-\Delta$ mass difference, giving $\alpha_s=1.319$. The strength of the linear confinement potential is $a=95.1$ MeV/fm, which satisfies the stability condition for the average octet-baryon mass at $b=0.5$ fm. Table III shows masses of the octet and decuplet baryons, where we chose $\xi_1=0.66$ and $\xi_2=0.31$ by fitting the $\Lambda\Xi^*-N\Omega$ and $\Sigma\Xi^*-N\Omega$ threshold differences. One sees that the other mass differences are also well reproduced.

The same model has been applied to the flavor-singlet H dihyperon state in a previous paper.[11] The model predicts no $\Lambda\Lambda$ bound state but a narrow resonance below the $N\Xi$ threshold. Coupling of $N\Xi$ and $\Sigma\Sigma$ channels is

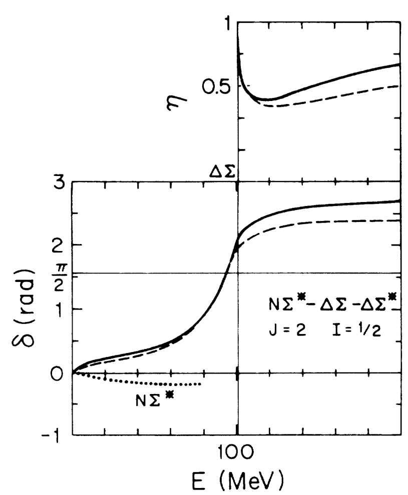

FIG. 1. Phase shift δ and elasticity η of $N-\Sigma^*$ with $J=2^+$ and $I=\frac{1}{2}$ S-wave scattering. The solid (dashed) curve shows the $N\Sigma^*-\Delta\Sigma-\Delta\Sigma^*(N\Sigma^*-\Delta\Sigma)$ coupled-channel calculation. The dotted curve is for the single-channel calculation.

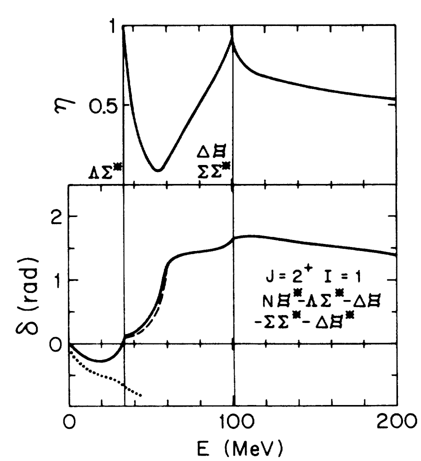

FIG. 2. Phase shift δ and elasticity η of $N-\Xi^*$ with $J=2^+$ and $I=1$ S-wave scattering. The solid (dashed) curve shows the $N\Xi^*-\Lambda\Sigma^*-\Delta\Xi-\Sigma\Sigma^*-\Delta\Xi^*(N\Xi^*-\Lambda\Sigma^*-\Delta\Xi-\Sigma\Sigma^*)$ coupled-channel calculation. The dotted curve is for the single-channel calculation.

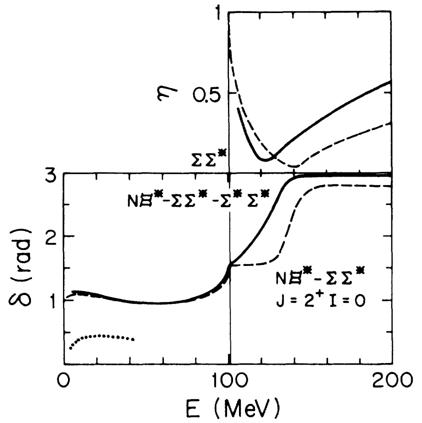

FIG. 3. Phase shift δ and elasticity η of $N-\Xi^*$ with $J=2^+$ and $I=0$ S-wave scattering. The solid (dashed) curve shows the $N\Xi^*-\Sigma\Sigma^*-E^*E^*(N\Xi^*-\Sigma\Sigma^*)$ coupled-channel calculation. The dotted curve is for the single-channel calculation.

essential to make this resonance.

In Figs. 1–8, we show the calculated S-wave phase shift, δ, and the elasticity, η, of the lowest-lying *BB* channel for each $(\mathcal{S},I,J)$ state. The solid and dashed (or dotted) curves are obtained by the full and a partial channel-coupling calculations, respectively. One sees, in general, the importance of channel couplings within the *BB* configurations which appear in the flavor-octet state. Coupling of the decuplet-decuplet channels, which do not participate in the flavor-octet states, is not so crucial. There is no bound state obtained in any of the calculated channels. In Figs. 1–4, the $J=2^+$ channels show rich structure, a quasibound state for ($\mathcal{S}=-2$, $I=0$, $J=2^+$) and ($\mathcal{S}=-3$, $I=\frac{1}{2}$, $J=2^+$), and low-energy resonances ($E<100$ MeV) for ($\mathcal{S}=-1$, $I=\frac{1}{2}$, $J=2^+$) and ($\mathcal{S}=-2$, $I=1$, $J=2^+$). At the resonance energies, the scattering wave function shows a superposition of the *BB* channels similar to the flavor-octet combination given in Table II.

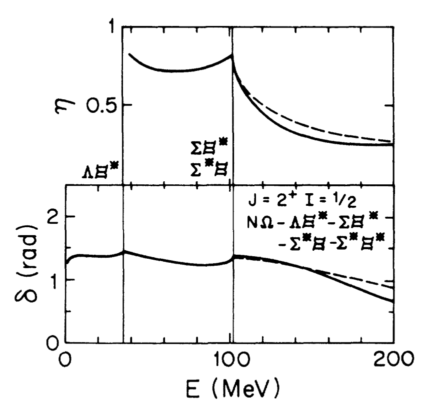

FIG. 4. Phase shift δ and elasticity η of $N-\Omega$ with $J=2^+$ and $I=\frac{1}{2}$ S-wave scattering. The solid (dashed) curve shows the $N\Omega-\Lambda\Xi^*-\Sigma\Xi^*-\Sigma^*\Xi-\Sigma^*\Xi^*(N\Omega-\Lambda\Xi^*-\Sigma\Xi^*-\Sigma^*\Xi)$ coupled-channel calculation. The $N\Omega$ single-channel calculation gives $\delta\equiv0$ due to the absence of the exchange interaction.

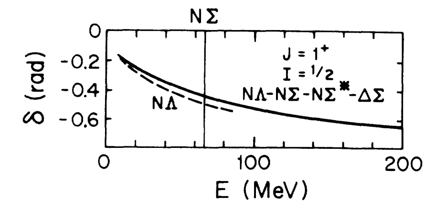

FIG. 5. Phase shift δ of $N-\Lambda$ with $J=1^+$ and $I=\frac{1}{2}$ S-wave scattering. The solid curve shows the $N\Lambda-N\Sigma-N\Sigma^*-\Delta\Sigma$ coupled-channel calculation. The elasticity η is almost 1 up to $E=200$ MeV. The dashed curve is for the single-channel calculation.

By contrast, the $J=1^+$ channels (Figs. 5–8) do not show low-energy resonances, although effect of the channel coupling is sometimes significant. The difference between the $J=1^+$ and 2^+ states comes from the fact that $J=1^+$ flavor-octet states are distributed over many *BB* channels, the octet octet as well as the octet decuplet. They have a wide range of the threshold energies due to the octet-decuplet mass difference and the flavor-SU(3) breaking. Some of the $J=1^+$ channels have been studied in a previous paper[11] in the context of the hyperon-nucleon and hyperon-hyperon interactions. Although the coupling of the octet-decuplet channels were not considered then, the low-energy octet-octet scattering phase shifts are qualitatively the same as the present calculation.

IV. CONCLUSION

The flavor-octet six-quark systems seem to have rich structure. We have studied the roles of the color-magnetic interaction and the quark-exchange processes by using the quark-cluster-model approach. We showed that all the $J=2^+$ channels have a resonance and/or a quasibound state. Couplings among the octet-decuplet two-baryon channels are important because the flavor-octet combinations are favored by CMI. It should be stressed that CMI plays a major role in the baryon-baryon interactions. The magnitude of CMI is determined by the $N-\Delta$ mass difference and therefore is model independent. The quark-exchange effect, i.e., the antisymmetrization among the six quarks, is also crucial. In fact, without the antisymmetrization, CMI does not induce any interbaryonic interaction. For example, the

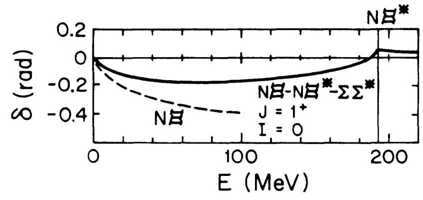

FIG. 6. Phase shift δ of $N-\Xi$ with $J=1^+$ and $I=0$ S-wave scattering. The solid curve shows the $N\Xi-N\Xi^*-\Sigma\Sigma^*$ coupled-channel calculation. The dashed curve is for the single-channel calculation.

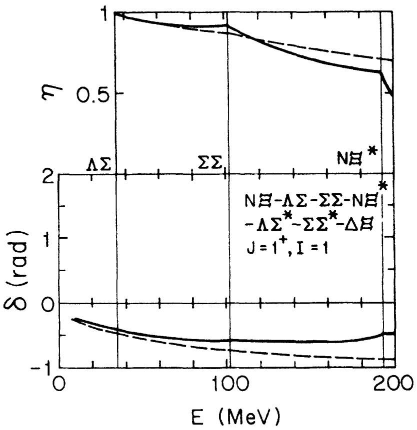

FIG. 7. Phase shift δ and elasticity η of $N-\Xi$ with $J=1^*$ and $I=1$ S-wave scattering. The solid (dashed) curve shows the $N\Xi-\Lambda\Sigma-\Sigma\Sigma-N\Xi^*-\Lambda\Sigma^*-\Sigma\Sigma^*-\Delta\Xi(N\Xi-\Lambda\Sigma-\Sigma\Sigma)$ coupled-channel calculation.

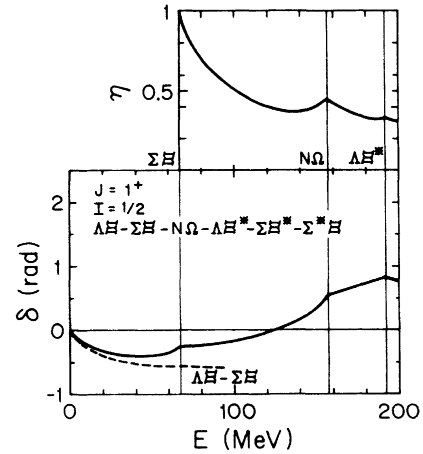

FIG. 8. Phase shift δ and elasticity η of $\Lambda-\Xi$ with $J=1^+$ and $I=\frac{1}{2}$ S-wave scattering. The solid (dashed) curve shows the $\Lambda\Xi-\Sigma\Xi-N\Omega-\Lambda\Xi^*-\Sigma\Xi^*-\Sigma^*\Xi(\Lambda\Xi-\Sigma\Xi)$ coupled-channel calculation.

$N-\Omega$ system does not have a quark-exchange contribution unless it is coupled with other channels such as Λ-Ξ^*, Σ-Ξ^*, or Σ^*-Ξ. The quark-cluster-model approach shows no $N-\Omega$ interaction in the single-channel calculation. Coupling of other channels is crucial.

No $J=1^+$ resonance is predicted up to 200 MeV above the lowest two-baryon threshold. There the flavor-octet state is distributed to many two-baryon channels, where thresholds are split by the octet-decuplet mass difference as well as the flavor-SU(3) breaking. Resonances do not appear in the low-energy region.

The present calculation is limited in the following two ways. First, the flavor-SU(3)-symmetry-breaking effect is not taken into account in the kinetic-energy term of the Hamiltonian. The reason for this shortcoming is purely technical. We might cure the problem in a different approach, for example, the quark shell model of the six-quark system. There the quark mass difference is easily handled, although the shell model would give the resonances as discrete states and therefore the connection to the two-baryon asymptotic channels become obscure. Second, because of the nonrelativistic nature of the model, meson ($q\bar{q}$) emission from the baryon is not incorporated. Thus we have to neglect the widths of the decuplet baryons and the long-range meson-exchange interactions are not considered. We have also neglected couplings of radial excited baryon states, which effectively correspond to deformations of the cluster wave functions. Considering these caveats, the energy of the resonances should not be taken too seriously. Instead we stress that the present study provides a systematic view of the dibaryon resonances and predicts the quantum numbers of state in which the resonances should be sought.

ACKNOWLEDGMENT

This work was supported in part by a grant from the National Science Foundation.

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