Quark delocalization, color screening, and dibaryons

Fan Wang
Center for Theoretical Physics, Nanjing University, Nanjing, 210008, China

Jia-lun Ping
Department of Physics, Nanjing Normal University, Nanjing, 210024, China

Guanh-han Wu and Li-jiang Teng
Institute of Nuclear Science and Technology, Sichuan University, Chengdu, 610064, China

T. Goldman
Theoretical Division, Los Alamos National Laboratory, Los Alamos, New Mexico 87545
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The quark delocalization and color screening model, a quark potential model, is used for a systematic search of dibaryon candidates in the u, d, and s three flavor world. Color screening, which appears in unquenched lattice gauge calculations, and quark delocalization (which is similar to electron delocalization in molecular physics) are both included. Flavor symmetry breaking and channel coupling effects are studied. The model is constrained not only by baryon ground state properties but also by the N-N scattering phase shifts. The deuteron and zero energy dinucleon resonance are both reproduced qualitatively. The model predicts two extreme types of dibaryonic systems: "molecular" like the deuteron, and highly delocalized six-quark systems among which only a few narrow dibaryon resonances occur in the u, d, and s three flavor world. Possible high spin dibaryon resonances are emphasized.

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I. INTRODUCTION

Quantum chromodynamics (QCD) is believed to be the fundamental theory of the strong interaction. High energy processes are calculable due to the asymptotic freedom property of QCD. The majority of low energy processes are uncalculable due to infrared confinement. Lattice gauge calculations may suffice in the confinement regime, but will still suffer from large numerical uncertainties for the prediction of many hadron properties in the near future. This leads to a reliance on QCD-inspired models to explore hadron physics for the time being and perhaps even in the future, due to the complexity of QCD. The existing models (potential, bag, soliton, etc.) are quite successful in understanding hadron (meson and baryon) properties, but have not been very successful for hadron interactions. Only recently have there been positive indications for obtaining the whole N-N interaction from QCD models [1, 2].

An outstanding problem is the fact that all of these models, including lattice gauge calculations, predict that there should be multiquark systems $(q\bar{q})^2, q^4\bar{q}, q^6$; quark gluon hybrids $q\bar{q}g, q^3g$; and glueballs, in addition to the $q\bar{q}$ mesons and q^3 baryons. Experimentally there are no well-established candidates for these exotics. In a relativistic theory, since quark and gluon number is not conserved, any meson state can be a mixture of $q\bar{q}$, $(q\bar{q})^2$, g^2, and $q\bar{q}g$; any baryon state can be a mixture of $q^3, q^4\bar{q}$, and q^3g. It is quite possible that these exotics, $(q\bar{q})^2, q\bar{q}g, g^2, q^4\bar{q}, q^3g$, exist in the normal meson and baryon states [3]. Polarized lepton nucleon scattering measurements have aroused a new round of hadron structure studies, wherein these exotic components are explored in connection with the normal q^3 and $q\bar{q}$ components [4]. However, q^6 is really a new quark system sector, different from that of mesons and baryons. We call a baryon number $B = 2$ state, which is quasistable, a dibaryon. Its minimum configuration is q^6. Since Jaffe predicted the first dibaryon, the H particle [5], a large number of dibaryon calculations with all the above-mentioned QCD models have been carried out and almost all support the existence of dibaryons [6]. If the present absence of an experimental dibaryon signal continues, then all these QCD models (and even QCD itself) should be questioned. Therefore, the dibaryon is a good place to test QCD and its models.

Recently, Silvestre-Brac *et al.* reported a new systematic dibaryon calculation based on the chromomagnetic model [6]. As pointed out by Lichtenberg and Roncaglia, the chromomagnetic model Hamiltonian is oversimplified [7]. The chromomagnetic interaction can give only the N-N short range repulsion but not any N-N attraction. Many dibaryon model calculations have the same deficiencies unless a phenomenological meson exchange is invoked. To study the dibaryon, it is better to have a model Hamiltonian which can at least fit the N-N interaction qualitatively. Then we can expect that such a model prediction may be relevant to real dibaryon states. Another deficiency of many prevailing model calculations is that the model Hilbert space is rather restricted. In

some model calculations [5, 6, 8], the six quarks are assumed to be completely merged into a single confinement region (which we term a "fully deconfined" model). In other model calculations, the quarks are assumed to be always confined separately in two distinct baryons (which we term a "fully confined" model) [9]. The real situation is quite possibly neither completely merged nor always separately confined, but, rather, in between, i.e., partially deconfined due to the interaction of two baryons. A more realistic model calculation allows the six-quark system to choose the preferred configuration by its own dynamics.

To remedy these model deficiencies, we developed a model which we will call the quark delocalization, color screening model (QDCSM) [2]. The model Hilbert space is enlarged to include the fully confined and fully deconfined models as two extremes and the real configuration is determined variationally by the dynamics of the six-quark system. In this way the system is allowed to develop its own preferred distortion. The model Hamiltonian is sufficiently realistic to produce a qualitatively correct N-N phase shift. We especially take into account the possible difference of the q-q interaction inside a hadron and between two colorless hadrons due to the nonlinearity of QCD (see Sec. II). We use this model to do a systematic search within the u, d, s-quark three-flavor world, expecting it to improve the reliability of estimates on promising dibaryon candidates. This expectation is realized in the N-N channels (see Sec. IV). Of course we cannot expect that the model estimate is quantitatively correct, because both the model Hamiltonian and Hilbert space are restricted to be simple enough to do a systematic search. As emphasized by Silvestre-Brac [10] this kind of systematic search serves the purpose of delimiting, among the thousands of multiquark states, the most promising candidates. Our intent is to assist experimental efforts to explore a challenging question in hadron physics by providing more reliable theoretical estimates.

This paper is organized as follows. In Sec. II, the model Hamiltonian and Hilbert space are described. Section III is devoted to a sketch of the calculation method. (A more complete description will be reported separately.) The results are given in Sec. IV and a conclusion in Sec. V.

II. QUARK DELOCALIZATION, COLOR SCREENING MODEL

Quark potential models are quite successful in describing single hadrons; therefore we adopt the usual potential model Hamiltonian to describe a single baryon:

$$H(3)=\sum_{i=1}^{3}(m_i+\frac{p_i^2}{2m_i})+\sum_{i<j=1}^{3}V_{ij}-T_c, \tag{1}$$

$$V_{ij}=V_{ij}^c+V_{ij}^G,$$

$$V_{ij}^c=-\vec{\lambda}_i\cdot\vec{\lambda}_j a r_{ij}^2,$$

$$V_{ij}^G=\alpha_s\frac{\vec{\lambda}_i\cdot\vec{\lambda}_j}{4}\left[\frac{1}{r_{ij}}-\frac{\pi}{2}\left(\frac{1}{m_i^2}+\frac{1}{m_j^2}+\frac{4}{3m_im_j}\vec{\sigma}_i\cdot\vec{\sigma}_j\right)\delta(\vec{r}_{ij})+\cdots\right].$$

The symbols in Eq. (1) have their usual meaning. For the confinement potential V^c, we assume a quadratic form to simplify the calculation. A possible constant part is omitted to reduce the number of parameters. In the effective one-gluon-exchange potential V^G, only the color Coulomb and color magnetic terms are retained, because we are only interested in the ground state. The effect of the momentum-dependent Darwin term has been checked and found not to be critical; hence it is also omitted to reduce the calculational burden.

A variational three-quark wave function (WF) of the form

$$\psi(123)=\chi(123)\eta_{SIJ}(123)\phi(123) \tag{2}$$

is assumed to describe the ground state baryons. Here $\chi(123)$ is the color singlet WF, $\eta_{SIJ}(123)$ is the symmetric spin-flavor $\mathrm{SU}_{2\times f}^{f\sigma}\supset\mathrm{SU}_f\times\mathrm{SU}_2^\sigma$ WF (S = strangeness, I = isospin, J = spin),

$$\phi(123)=\phi_{1s}(\vec{r}_1)\phi_{1s}(\vec{r}_2)\phi_{1s}(\vec{r}_3), \tag{3a}$$

$$\phi_{1s}(\vec{r})=\left(\frac{1}{\pi b^2}\right)^{\frac{3}{4}}e^{-\frac{(\vec{r}-\vec{s})^2}{2b^2}}, \tag{3b}$$

$\vec{s}$ is a reference center, and b is a baryon size parameter to be determined by the stability condition

$$\frac{\partial M(123)}{\partial b}=0, \tag{4}$$

where $M(123)=\langle\psi(123)|H(123)|\psi(123)\rangle$.

The other model parameters are fixed as follows: The u, d-quark mass difference is neglected and $m=m_u=m_d$ is assumed to be exactly $\frac{1}{3}$ of the nucleon mass M, i.e., $m=313\,\mathrm{MeV}$. The quark gluon coupling constant α_s is determined by the N-Δ mass difference. The confinement potential strength a is determined by the zero nucleon binding. The strange quark mass m_s is determined by an overall fit to the strange baryon masses under the assumption that all the flavor octet and decuplet baryons have the same rms radius b. (Choosing different values of b for different baryons will make the calculation more elaborate and is left as a future refinement.) The fitted parameters are $m=313\,\mathrm{MeV}$, $m_s=634\,\mathrm{MeV}$, $b=0.603\,\mathrm{fm}$, α_s=1.54, and $a=25.13\,\mathrm{MeV\,fm^{-2}}$. The theoretical baryon masses are compared with experimental values in Table I.

The strange baryon masses do not agree perfectly due to our simple model assumptions. The largest deviation is 41 MeV. The lower bounds on the constituent quark mass differences derived from the Feynman-Hellman theorem are fulfilled for our fitted quark masses, which are quite close to those of Lichtenberg [11]. The other parameters are very similar in value to the usual baryon spectroscopic results [12].

The direct extension of the single-baryon Hamiltonian (1) to q^6 is neither reasonable nor successful. First, the two-body confinement interaction will give rise to a spurious color van der Waals force. We take this as an indication that the q-q confinement interaction between two color singlet hadrons is modified due to the nonlinearity of QCD [13]. The nonperturbative and lattice gauge approaches both give rise to a string structure instead of a

TABLE I. Baryon masses (MeV).

k N	Δ	Λ	Σ	Ξ	Σ^*	Ξ^*	Ω	
Expt.	939	1232	1115	1193	1318	1385	1533	1672
Theor.	939	1232	1118	1217	1359	1361	1504	1658

two-body q-q interaction [14]. The string structure and two-body confinement give rise to similar spectroscopic results for simple quark system but are not identical [15, 16]. Lattice gauge calculations, after taking the (light) $q\bar{q}$ excitations into account, show that the (heavy) Q-$\bar{Q}$ confinement interaction is screened. Numerical results can be fitted by the following color screening interaction [17]

$$V(r) = \left(-\frac{\alpha_s}{r} + \sigma r\right)\left(\frac{1-e^{-\mu r}}{\mu r}\right), \tag{5}$$
$$\alpha_s = 0.21 \pm 0.01, \quad \sqrt{\sigma} = 400\,\text{MeV},$$
$$\mu^{-1} = 0.90 \pm 0.20\ \text{fm}.$$

Based on these results, we model the Hamiltonian of q^6 as an otherwise direct extension of Eq. (1), but modify the confinement part as

$$V_{ij}^c = \begin{cases} -\vec{\lambda}_i \cdot \vec{\lambda}_j a r_{ij}^2 & \text{if } i,j \text{ occur in the same baryon orbit,} \\ -\vec{\lambda}_i \cdot \vec{\lambda}_j \frac{a}{\mu}(1-e^{-\mu r_{ij}^2}) & \text{if } i,j \text{ occur in different baryon orbits.} \end{cases} \tag{6a}$$

Here the exponential $e^{-\mu r}$ appearing in Eq. (5) has been replaced by a Gaussian $e^{-\mu r^2}$, solely to simplify the numerical calculations. Another reason is this form will automatically match the quadratic confinement in the short distance ($\mu r^2 \ll 1$) region. Keeping the same form of confinement as that of a single baryon when the interacting pair of quarks occur in the same baryon guarantees that when the two baryons are separated to large distances, the energy of the q^6 system evolves into the two-baryon internal energy calculated by the Hamiltonian Eq. (1).

Svetitsky [18] gave a qualitative description of the Q-$\bar{Q}$ potential: The short range Coulomb potential evolves into a linear part at larger distance, but at still larger distances, the linear part evolves into an exponentially decaying Yukawa potential due to light meson exchange. To take this possibility into account, we assume another screening confinement potential (as described above)

$$V_{ij}^c = \begin{cases} -\vec{\lambda}_i \cdot \vec{\lambda}_j a r_{ij}^2 & \text{if } i,j \text{ occur in the same baryon orbit,} \\ -\vec{\lambda}_i \cdot \vec{\lambda}_j a r_{ij}^2 e^{-\nu r_{ij}^2} & \text{if } i,j \text{ occur in different baryon orbits.} \end{cases} \tag{6b}$$

Although Eq. (6a) would be favored as better representing the Yukawa falloff to a constant potential at large distances, that constant is not the zero value expected as between colorless hadrons. Thus, Eq. (6b) both tests the sensitivity (which turns out to be very mild) of our model to the long distance behavior of the potential and provides a model with the expected zero value of the hadronic potential at infinity.

The screening parameter $\mu(\nu)$ is determined by fitting our model to the N-N scattering phase shifts and they are [2]

$$\nu = 0.40\,\text{fm}^{-2}, \quad \mu = 1.0\,\text{fm}^{-2}.$$

Note particularly that the latter is consistent with the lattice result [17].

For the q^6 model space, we extend the quark cluster model space by introducing a delocalized single-quark orbit. In the usual quark cluster model approach, two single-quark orbits are assumed:

$$\phi_L(\vec{r}) = \left(\tfrac{1}{\pi b^2}\right)^{\frac{3}{4}} e^{-\frac{(\vec{r}+\vec{s}/2)^2}{2b^2}} \quad \text{(left centered orbit)},$$
$$\phi_R(\vec{r}) = \left(\tfrac{1}{\pi b^2}\right)^{\frac{3}{4}} e^{-\frac{(\vec{r}-\vec{s}/2)^2}{2b^2}} \quad \text{(right centered orbit)}. \tag{7}$$

Here $s = \|\vec{s}\|$ is the separation of the centers of two q^3 clusters. We introduce the delocalized quark orbit

$$\psi_L(\vec{r}) = [\phi_L(\vec{r}) + \epsilon(s)\phi_R(\vec{r})]\,/N(s),$$
$$\psi_R(\vec{r}) = [\phi_R(\vec{r}) + \epsilon(s)\phi_L(\vec{r})]\,/N(s), \tag{8}$$
$$N^2(s) = 1 + \epsilon^2(s) + 2\epsilon(s)\langle\phi_L|\phi_R\rangle.$$

The delocalization parameter $\epsilon(s)$ is determined variationally for every $\vec{s}$ by the dynamics of the q^6 system (see Sec. III). This orbit is a generalization of the quark molecular orbit introduced by Stancu and Wilets [19], The six-quark space is restricted to be the space spanned by the following quark cluster bases:

$$\Psi^{\alpha}_{\alpha_1 F_1, \alpha_2 F_2}(1\cdots 6) = \mathcal{A}[\psi_{\alpha_1 F_1}(123)\psi_{\alpha_2 F_2}(456)]_\alpha,$$
$$\psi_{\alpha_1 F_1}(123) = \chi(123)\eta_{S_1 I_1 J_1 F_1}(123)\psi_L(1)\psi_L(2)\psi_L(3)\ , \tag{9}$$
$$\psi_{\alpha_2 F_2}(456) = \chi(456)\eta_{S_2 I_2 J_2 F_2}(456)\psi_R(4)\psi_R(5)\psi_R(6).$$

Here $\alpha = (SIJ)$ describes the strong interaction conserved quantum numbers of strangeness, isospin, and spin.[1] The q^3 cluster WF is almost the same as given in Eq. (2), but the single cluster Gaussian WF (3b) is replaced by the delocalized orbital WF (8), and a flavor symmetry quantum number F is shown explicitly. $[\ ]_\alpha$ refers to isospin and spin coupling by means of the SU_2 Clebsch-Gordan coefficients. The SU_3 color coupling is trivial because only color singlet hadron states are used in our calculation. $\mathcal{A}$ is the normalized antisymmetry operator

$$\mathcal{A} = \frac{1}{\sqrt{20}} \sum (-)^{\delta_p} p \ ,$$

where p is the two-quark permutation operator. Equation (9) is termed the physical basis by Harvey [20]. Even though only totally symmetric q^3 orbital configurations are included in Eq. (9), q^3 orbital excitation configurations are included due to the delocalized single-quark orbit (8) used in Eq. (9). Hidden color channels are *not* included in Eq. (9). The reason is that the colorless channels are already complete when orbital-spin-isospin excitation configurations are included [15, 21]. Furthermore, color states have not been well constrained in QCD models. It should be clearly kept in mind that these physical bases are dependent on the separation $\vec{s}$ of the two q^3 cluster centers and the delocalization parameter $\epsilon(s)$.

A check had been done in the N-N scattering dynamical calculation, that if we start with the Stancu-Wilets molecular orbits instead of the left and right centered orbits ϕ_L and ϕ_R, exactly the same results are obtained [2].

III. CALCULATION METHOD

A dynamical calculation of single-channel N-N scattering has been done first to fix the screening parameter $\mu(\nu)$ by fitting the model N-N phase shifts to the experimental ones [2].

The q^6 states of a given set of quantum numbers $\alpha = (SIJ)$ are expressed as a multiple physical channel coupling WF

$$\Psi_\alpha(1\cdots 6) = \sum_{\alpha_1 F_1, \alpha_2 F_2} C^\alpha_{\alpha_1 F_1, \alpha_2 F_2} \Psi^\alpha_{\alpha_1 F_1, \alpha_2 F_2} . \quad (10)$$

The channel coupling coefficients $C^\alpha_{\alpha_1 F_1, \alpha_2 F_2}$ are determined by the diagonalization of the q^6 Hamiltonian. The maximum number of channels coupled is 16. In the diagonalization, the nonorthogonality property of the physical bases is properly accounted for.

The six-quark Hamiltonian matrix elements

$$\langle \Psi^\alpha_{\alpha'_1 F'_1, \alpha'_2 F'_2} | H(1\cdots 6) | \Psi^\alpha_{\alpha_1 F_1, \alpha_2 F_2} \rangle$$

are calculated by the group theory method developed by Harvey and by Chen and ourselves [20]:

(1) The physical bases $\Psi^\alpha_{\alpha_1 F_1, \alpha_2 F_2}$ are expressed in terms of the symmetry bases (group chain classification bases) by the $6 \to 3\times 3$, $SU_{mn} \supset SU_m \times SU_n$, and $SU_3^f \supset SU_2^\tau \times U_1^S$ isoscalar factors calculated by Chen *et al.* [22].

(2) The six-quark Hamiltonian matrix elements in the symmetry bases are reduced to a two-body matrix element and a four-quark overlap (due to the nonorthogonality of the delocalized orbit) by the traditional parentage expansion ($6 \to 4 \times 2$, $SU_{mn} \supset SU_m \times SU_n$, and $SU_3^f \supset SU_2^\tau \times U_1^S$ isoscalar factor [22]).

(3) The four-quark overlap is reduced to a one-body overlap by the permutation symmetry property of the four-quark state [23].

(4) Two-body confinement interaction matrix elements are calculated as follows: If the interacting quark pair occurs in the same left or right orbit, i.e., $\langle LL|V|LL\rangle$, $\langle RR|V|RR\rangle$, $\langle LL|V|RR\rangle$, and $\langle RR|V|LL\rangle$, then the usual quadratic confinement form is used. We do so since the two quarks involved are then always in the same "baryon" and hence in a relative color antitriplet state. For all the other two-body confinement interaction matrix elements, such as $\langle LR|V|LR\rangle$, $\langle LL|V|LR\rangle$, ..., the color screening confinement form is used. Again, we do so since the confining interaction will cancel over all the quark pairs identifiable as originating in different color singlet hadrons; we effectively remove it in advance from each pairwise interaction, leaving an interaction which mimics color singlet (mesonic) exchanges. Here $\langle LL|V|RR\rangle$ means $\langle \phi_L(1)\phi_L(2)|V_{12}|\phi_R(1)\phi_R(2)\rangle$.

The eigenenergies obtained in this way are dependent on the separation $\vec{s}$ and the delocalization parameter $\epsilon(s)$. We repeat the calculation for each $\vec{s}$ by varying $\epsilon(s)$ from 0.1 to 1.0 with step size 0.1 to get a minimum of the eigenenergy, which thus also determines the delocalization parameter $\epsilon(s)$. The difference of the minimum eigenenergy at separation $\vec{s}$ and the minimum eigenenergy at infinite separation is taken to be the baryon-baryon potential energy $V_\alpha(s)$ (an adiabatic approximation). Numerically the asymptotic values of the eigenenergy are indistinguishable from the calculation at $s = 3$ fm. It is equal to the threshold sum, i.e., the sum of theoretical masses of the corresponding channel baryon pair (single-channel case) or of the lightest baryon pair (channel coupling case), and the model relative kinetic energy of this baryon pair, which is equal to 1/6 of the total kinetic energy of the q^6 system due to our model WF assumption, Eq. (9). (This is one of the checks of our numerical calculation; another check is that all the channel mixing matrix elements are vanishingly small at $s = 3$ fm.) A zero-point harmonic oscillation energy $\frac{3\hbar^2}{4\mu s_0^2}$ is added to the minimum potential energy $V_\alpha(s_0)$ to obtain the binding energy B_α of a q^6 state with the quantum number α:

$$B_\alpha = V_\alpha(s_0) + \frac{3\hbar^2}{4\mu_\alpha s_0^2} \ . \quad (11)$$

Here μ_α is taken simply to be the reduced mass of the corresponding channel baryon pair (single-channel case)

[1] Orbital angular momentum is assumed to be zero for the lowest states. An angular momentum projection which should be done is left for future work. Preliminary estimates indicate that this correction to the state energy is small.

or the lightest baryon pair within the quantum number α set (channel coupling case). In principle, we should do a multichannel coupling dynamical calculation; this program is being pursued only for the few most promising dibaryon candidates.

Finally the experimental channel baryon pair mass (single-channel case) or the lightest baryon pair mass $(M_1 + M_2)_\alpha$ (channel coupling case) is added to B_α to obtain the lowest q^6 mass of each quantum number α set:

$$M_\alpha(q^6) = (M_1 + M_2)_\alpha + B_\alpha. \tag{12}$$

The mass $M_\alpha(q^6)$ is compared not only to the two-body decay threshold, the experimental lightest baryon pair mass $(M_1 + M_2)_\alpha$, but also to the possible multiparticle final states allowed by strong interaction to determine if there is a strong interaction quasistable dibaryon state.

In order to show the effects of channel coupling, flavor-symmetry breaking, and the different forms of color screening, the following eight sets of calculations have been done: (1) single-channel, flavor symmetry (scs), (2) multichannel, flavor symmetry (ccs), (3) single-channel, flavor-symmetry breaking (scb), and (4) multichannel, flavor-symmetry breaking (ccb), where (1)–(4) have been calculated using the color screening form (6a) and (5)–(8) are (1)–(14) recalculated with the color screening form (6b).

To indicate the level of uncertainty due to the choice of the color screening parameter value, we have calculated the results corresponding to $\mu = 1.6$ fm^{-2} and $\nu = 0.46, 0.60$ fm^{-2} in addition to the best fit values $\mu = 1.0$ fm^{-2} and $\nu = 0.40$ fm^{-2}. (The reasoning behind the parameter values chosen for these additional cases is described in the Appendix.) Only dibaryon candidates based on both the channel coupling and symmetry breaking results should be considered serious possibilities for experimental searches.

A computer program package which incorporates all the needed group theory results has been written to automate the numerical calculation. It can be used for other model calculations by simply replacing the one- and two-body matrix elements. In particular, it is may also be used for a relativistic quark cluster model calculation. As a cross-check on the program, the $S = 0, (IJ) = (01), (10)$, and (03) channels have been done both by the method described above and by direct diagonalization with the physical bases.

IV. RESULTS

All possible sets of $\alpha = (SIJ)$ within the u, d, and s three-flavor world have been calculated. Only a few states are strong interaction quasistable or narrow resonances, and they are listed in Table II.

The strong interaction unstable states have been omitted to simplify the presentation. However, some general features are listed here.

(1) The two color screening forms give quite similar results: The form (6a) gives slightly higher six-quark masses than the form (6b) but the largest difference is only about 10 MeV for the $(SIJ) = (003)$ case. This is consistent with the findings for N-N scattering where the forms (6a) and (6b) give similar N-N phase shifts while form (6b) yields a little stronger attraction in the $^3S_1$ channel [2].

(2) There are two extreme kinds of dibaryon candidates. One kind is a loosely bound ("molecular") two-baryon state. The binding energy B_α is small (usually around zero), $J \leq 1$; the delocalization $\epsilon(s_0)$ is also small (usually $\lesssim 0.2$). (The deuteron is a typical example while the H particle is an exception.) Their masses are close to the lowest two-body decay threshold and they might be stable with respect to the strong interactions, but their stability is very sensitive to model details. The other kind is a tightly bound state. The binding energy is large ($\gtrsim 100$ MeV), $J \geq 2$; the delocalization is also large ($\gtrsim 0.8$). Their masses are larger than the lowest two-body threshold and they are therefore unstable. However, their masses are lower than the favorable three- or four-body decay threshold. Their two- or three-body decay is hindered due to large angular momentum and so they might nonetheless appear as narrow resonances. A typical example is the di-$\Delta(003)$ state. This quasistability property is *not* sensitive to the model details, unless radically different model assumptions are made.

(3) Flavor-symmetry breaking effects are channel dependent as reported by Maltman [24], decreasing the binding by an amount ranging from negligible to 70 MeV. For some multichannel coupling cases, the flavor-symmetry breaking effect changes which is the lowest channel. In those cases, there is large apparent flavor-symmetry breaking effect ($\Delta B_\alpha \sim 120$–180 MeV). The mass of the q^6 state is increased correspondingly and this produces a large difference in the stability of the q^6 state with respect to the multiparticle decay channels, due to the flavor-symmetry breaking.

(4) The channel coupling effect is small for most cases, even after taking into account quark delocalization in the extended u, d, s world. For the H particle case, channel crossing occurs: The dominant channel for the minimum potential ($\Sigma\Sigma$) and the asymptotic channel ($\Lambda\Lambda$) are different. This channel coupling decreases the q^6 mass by about 100 MeV and makes the H particle stable. A similar channel coupling effect occurs in the (SIJ)=(-202) state.

We emphasize that all the QDCSM parameters are fixed by the ground state baryon properties and the N-N scattering phase shifts. Therefore the dibaryon candidates noted here are a relatively robust theoretical prediction. We wish to emphasize a few additional points.

(1) In the $S = 0, (IJ) = (01)$ channel (the deuteron channel), the model predicts that there is a state with $M(001) = 1880(\pm 14)$ MeV. It is a dinucleon state because the delocalization $\epsilon = 0.2$ is small and 4 MeV away from the deuteron energy. We take this as a measure of the predictive power of the QDCSM for the dibaryon state. Due to this uncertainty, we have included those states which are close to the lowest two-body decay threshold in Table II.

(2) In the $S = 0, (IJ) = (10)$ channel, the model predicts that there is a dinucleon [$\epsilon(s_0) = 0.1$] resonance state at $1889(\pm 4)$ MeV. It is about 10 MeV away from a possible zero binding dinucleon resonance. This is an-

TABLE II. Masses (M_α in MeV), potentials (V_α in MeV) at equilibrium separations (s_0 in fm), binding energies (B_α in MeV) relative to an asymptotic two-baryon system with channel quantum numbers and delocalization parameter values (ϵ) for all six quark three-flavor systems without orbital angular momentum contributions to total spin.

(a) $\mu = 1.0$ fm^{-2}, $\nu = 0.4$ fm^{-2}

		(6b)					(6a)					
SIJ		M_α	V_α	B_α	ϵ	s_0	M_α	V_α	B_α	ϵ	s_0	Threshold
001	scs	1885	-21	7	0.1	1.5	1885	-20	7	0.1	1.5	1878(NN)
	ccs	1886	-23	8	0.2	1.4	1894	-21	16	0.2	1.3	
010	scs	1891	-11	13	0.1	1.6	1893	-10	15	0.1	1.6	1878(NN)
	ccs	1891	-11	13	0.1	1.6	1892	-9	14	0.1	1.6	
003	scs	2134	-363	-330	1.0	1.2	2144	-359	-320	1.0	1.1	2464($\Delta\Delta$)
												2158($NN\pi\pi$)
-1$\frac{1}{2}$3	scs	2285	-363	-332	1.0	1.2	2294	-359	-322	1.0	1.1	2617($\Delta\Sigma^*$)
	ccs	2285	-363	-363	1.0	1.2	2294	-359	-322	1.0	1.1	2335($N\Lambda\pi\pi$)
	scb	2343	-311	-274	1.0	1.1	2346	-308	-271	1.0	1.1	
	ccb	2343	-311	-274	1.0	1.1	2346	-308	-271	1.0	1.1	
1$\frac{3}{2}$0	scs	2133	-24	1	0.1	1.5	2134	-23	2	0.1	1.5	2132($N\Sigma$)
	ccs	2133	-24	1	0.1	1.5	2134	-23	2	0.1	1.5	
	scb	2137	-19	5	0.1	1.5	2138	-18	6	0.1	1.5	
	ccb	2137	-19	5	0.1	1.5	2138	-18	6	0.1	1.5	
-200	scs	2145	-195	-112	0.5	0.8	2143	-198	-114	0.5	0.8	2231($\Lambda\Lambda$)
	ccs	2059	-318	-172	1.0	0.6	2055	-321	-176	1.0	0.6	
	scb	2328	-194	-58	1.0	0.6	2322	-198	-62	1.0	0.6	
	ccb	2222	-155	-9	1.0	0.6	2218	-159	-14	1.0	0.6	
-202	scs	2297	-217	-175	0.6	1.1	2307	-215	-165	0.6	1.0	2472($N\Xi^*$)
	ccs	2205	-318	-268	1.0	1.0	2216	-319	-257	1.0	0.9	2397($N\Xi\pi$)
	scb	2478	-171	-99	0.7	0.8	2476	-173	-102	0.7	0.8	
	ccb	2369	-182	-103	1.0	1.8	2367	-184	-106	1.0	0.8	
-213	scs	2432	-363	-333	1.0	1.2	2442	-359	-324	1.0	1.1	2765($\Delta\Xi^*$)
	ccs	2432	-363	-333	1.0	1.2	2442	-359	-324	1.0	1.1	2690($\Delta\Xi\pi$)
	scb	2559	-252	-210	1.0	1.0	2560	-251	-209	1.0	1.0	2511($\Lambda\Lambda\pi\pi$)
	ccb	2556	-252	-209	1.0	1.0	2557	-251	-209	1.0	1.0	
-220	scs	2394	-11	8	0.1	1.6	2395	-10	9	0.1	1.6	2386($\Sigma\Sigma$)
	ccs	2394	-11	8	0.1	1.6	2395	-10	9	0.1	1.6	
	scb	2393	-12	7	0.1	1.6	2397	-11	11	0.1	1.5	
	ccb	2393	-12	7	0.1	1.6	2397	-11	11	0.1	1.5	
-3$\frac{3}{2}$3	scs	2570	-363	-335	1.0	1.2	2579	-359	-325	1.0	1.1	2904($\Delta\Omega$)
	ccs	2570	-363	-335	1.0	1.2	2579	-359	-325	1.0	1.1	2788($\Lambda\Xi^*\pi$)
	scb	2767	-201	-151	1.0	0.9	2766	-202	-152	1.0	0.9	2714($\Lambda\Xi\pi\pi$)
	ccb	2754	-201	-150	1.0	0.9	2754	-201	-150	1.0	0.9	
-3$\frac{3}{2}$1	scs	2510	-22	-1	0.1	1.5	2511	-20	0	0.1	1.5	2511($\Sigma\Xi$)
	ccs	2512	-23	-0	0.2	1.4	2518	-21	7	0.2	1.3	
	scb	2525	-25	14	0.2	1.1	2525	-25	14	0.2	1.1	
	ccb	2525	-25	14	0.2	1.1	2525	-25	14	0.2	1.1	
-3$\frac{1}{2}$2	scs	2394	-266	-218	0.8	1.0	2394	-266	-217	0.8	1.0	2611($N\Omega$)
	ccs	2342	-318	-269	1.0	1.0	2353	-319	-259	1.0	0.9	2574($\Lambda\Xi\pi$)
	scb	2556	-267	-147	1.0	0.6	2552	-271	-151	1.0	0.6	
	ccb	2552	-195	-60	1.0	0.6	2548	-198	-64	1.0	0.6	
-3$\frac{1}{2}$1	scs	2394	-266	-218	0.8	1.0	2394	-266	-217	0.8	1.0	2434($\Lambda\Xi$)
	ccs	2575	-107	-59	0.4	1.0	2385	-108	-48	0.5	0.9	
	scb	2543	-84	109	1.0	0.5	2540	-87	106	1.0	0.5	
	ccb	2527	-100	93	1.0	0.5	2523	-104	89	1.0	0.5	
-400	scs	2632	-24	-4	0.1	1.5	2633	-23	-3	0.1	1.5	2636($\Xi\Xi$)
	ccs	2632	-24	-4	0.1	1.5	2633	-23	-3	0.1	1.5	
	scb	2643	-24	7	0.2	1.2	2643	-24	7	0.1	1.2	
	ccb	2643	-24	7	0.2	1.2	2643	-24	7	0.1	1.2	
-600	scs	3287	-74	-58	0.2	1.5	3291	-69	-54	0.2	1.5	3345($\Omega\Omega$)
	scb	3351	-29	-6	0.1	1.0	3350	-30	5	0.1	1.0	

(b) $\nu = 0.46$ fm^{-2}

		(6b)					Threshold
001	scs	1885	-30	7	0.2	1.3	1878(NN)
	ccs	1878	-31	0	0.2	1.4	

TABLE II *(Continued).*

		(6b)					Threshold
010	scs	1888	-14	10	0.1	1.6	1878(NN)
	ccs	1888	-14	10	0.1	1.6	
003	scs	2112	-385	-352	1.0	1.2	2464($\Delta\Delta$)
							2158($NN\pi\pi$)
-1$\frac{1}{2}$3	scs	2263	-385	-353	1.0	1.2	2617($\Delta\Sigma^*$)
	ccs	2263	-385	-353	1.0	1.2	2335($N\Lambda\pi\pi$)
	scb	2322	-331	-294	1.0	1.1	
	ccb	2322	-331	-294	1.0	1.1	
1$\frac{3}{2}$0	scs	2132	-29	0	0.2	1.4	2132($N\Sigma$)
	ccs	2132	-29	0	0.2	1.4	
	scb	2134	-23	2	0.1	1.5	
	ccb	2134	-23	2	0.1	1.5	
-200	scs	2128	-212	-129	0.6	0.8	2231($\Lambda\Lambda$)
	ccs	2043	-333	-188	1.0	0.6	
	scb	2312	-210	-74	1.0	0.6	
	ccb	2206	-170	-25	1.0	0.6	
-202	scs	2278	-236	-195	0.6	1.1	2472($N\Xi^*$)
	ccs	2185	-337	-287	1.0	1.0	2397($N\Xi\pi$)
	scb	2461	-188	-117	1.0	0.8	
	ccb	2352	-199	-121	1.0	0.8	
-213	scs	2410	-385	-355	1.0	1.2	2765($\Delta\Xi^*$)
	ccs	2410	-385	-355	1.0	1.2	2690($\Delta\Xi\pi$)
	scb	2540	-271	-229	1.0	1.0	2511($\Lambda\Lambda\pi\pi$)
	ccb	2537	-271	-229	1.0	1.0	
-220	scs	2391	-14	5	0.1	1.6	2386($\Sigma\Sigma$)
	ccs	2391	-14	5	0.1	1.6	
	scb	2393	-18	7	0.2	1.4	
	ccb	2393	-18	7	0.2	1.4	
-3$\frac{3}{2}$3	scs	2548	-385	-356	1.0	1.2	2904($\Delta\Omega$)
	ccs	2548	-385	-356	1.0	1.2	2788($\Lambda\Xi^*\pi$)
	scb	2749	-219	-169	1.0	0.9	2714($\Lambda\Xi\pi\pi$)
	ccb	2736	-219	-168	1.0	0.9	
-3$\frac{3}{2}$1	scs	2509	-30	-2	0.2	1.3	2511($\Sigma\Xi$)
	ccs	2504	-31	-8	0.2	1.4	
	scb	2515	-35	3	0.3	1.1	
	ccb	2514	-34	2	0.3	1.1	
-3$\frac{1}{2}$2	scs	2375	-285	-237	0.9	1.0	2611($N\Omega$)
	ccs	2323	-337	-289	1.0	1.0	2574($\Lambda\Xi\pi$)
	scb	2540	-283	-163	1.0	0.6	
	ccb	2536	-210	-76	1.0	0.6	
-3$\frac{1}{2}$1	scs	2375	-285	-237	0.9	1.0	2434($\Lambda\Xi$)
	ccs	2358	-124	-75	0.5	1.0	
	scb	2469	-99	36	1.0	0.6	
	ccb	2452	-116	19	1.0	0.6	
-400	scs	2630	-29	-6	0.2	1.4	2636($\Xi\Xi$)
	ccs	2630	-29	-6	0.2	1.4	
	scb	2642	-31	6	0.2	1.1	
	ccb	2641	-32	5	0.2	1.1	
-600	scs	3274	-86	-71	0.3	1.5	3345($\Omega\Omega$)
	scb	3341	-38	-4	0.2	1.0	

(c) $\mu = 1.6$ fm^{-2}, $\nu = 0.6$ fm^{-2}

		(6b)					(6a)					
SIJ		M_α	V_α	B_α	ϵ	s_0	M_α	V_α	B_α	ϵ	s_0	Threshold
001	scs	1869	-46	9	0.2	1.3	1873	-42	5	0.2	1.3	1878(NN)
	ccs	1867	-48	11	0.3	1.3	1873	-42	5	0.2	1.3	
010	scs	1890	-20	12	0.2	1.4	1885	-18	7	0.1	1.6	1878(NN)
	ccs	1888	-22	10	0.2	1.4	1885	-18	7	0.1	1.6	
003	scs	2074	-422	-390	1.0	1.2	2084	-413	-380	1.0	1.2	2464($\Delta\Delta$)
												2158($NN\pi\pi$)

TABLE II *(Continued).*

		(6b)					(6a)					
$-1\frac{1}{2}3$	scs	2225	-422	-390	1.0	1.2	2235	-413	-382	1.0	1.2	$2617(\Delta\Sigma^*)$
	ccs	2225	-422	-390	1.0	1.2	2235	-413	-382	1.0	1.2	$2335(N\Lambda\pi\pi)$
	scb	2285	-369	-332	1.0	1.1	2292	-361	-324	1.0	1.1	
	ccb	2285	-369	-332	1.0	1.1	2292	-361	-324	1.0	1.1	
$1\frac{3}{2}0$	scs	2119	-41	13	0.2	1.4	2123	-37	9	0.2	1.4	$2132(N\Sigma)$
	ccs	2119	-41	13	0.2	1.4	2123	-37	9	0.2	1.4	
	scb	2126	-34	6	0.2	1.4	2130	-30	2	0.2	1.4	
	ccb	2126	-34	6	0.2	1.4	2130	-30	2	0.2	1.4	
-200	scs	2097	-244	-161	1.0	0.8	2097	-244	-161	0.9	0.8	$2231(\Lambda\Lambda)$
	ccs	1974	-364	-257	1.0	0.7	2012	-365	-220	1.0	0.6	
	scb	2282	-240	-104	1.0	0.6	2281	-241	-105	1.0	0.6	
	ccb	2176	-200	-55	1.0	0.6	2175	-202	-56	1.0	0.6	
-202	scs	2242	-272	-231	0.8	1.1	2249	-265	-223	0.7	1.1	$2472(N\Xi^*)$
	ccs	2150	-373	-323	1.0	1.0	2255	-368	-318	1.0	1.0	$2397(N\Xi\pi)$
	scb	2413	-221	-164	0.9	0.9	2429	-219	-148	1.0	0.8	
	ccb	2319	-231	-153	1.0	0.8	2320	-230	-152	1.0	0.8	
-213	scs	2372	-423	-393	1.0	1.2	2382	-413	-383	1.0	1.2	$2765(\Delta\Xi^*)$
	ccs	2372	-423	-393	1.0	1.2	2382	-413	-383	1.0	1.2	$2690(\Delta\Xi\pi)$
	scb	2496	-308	-273	1.0	1.1	2509	-302	-260	1.0	1.0	$2511(\Lambda\Lambda\pi\pi)$
	ccb	2493	-308	-272	1.0	1.1	2506	-302	-259	1.0	1.0	
-220	scs	2391	-20	5	0.2	1.4	2388	-18	2	0.1	1.6	$2386(\Sigma\Sigma)$
	ccs	2390	-21	4	0.2	1.4	2388	-18	2	0.1	1.6	
	scb	2382	-33	-4	0.2	1.3	2386	-29	0	0.2	1.3	
	ccb	2381	-34	-5	0.2	1.3	2385	-30	-1	0.2	1.3	
$-3\frac{3}{2}3$	scs	2510	-423	-394	1.0	1.2	2520	-413	-384	1.0	1.2	$2904(\Delta\Omega)$
	ccs	2510	-423	-394	1.0	1.2	2520	-413	-384	1.0	1.2	$2788(\Lambda\Xi^*\pi)$
	scb	2704	-254	-214	1.0	1.0	2717	-250	-201	1.0	0.9	$2714(\Lambda\Xi\pi\pi)$
	ccb	2691	-255	-213	1.0	1.0	2705	-250	-200	1.0	0.9	
$-3\frac{3}{2}1$	scs	2493	-46	-18	0.2	1.3	2497	-42	14	0.2	1.3	$2511(\Sigma\Xi)$
	ccs	2491	-48	-20	0.3	1.3	2497	-42	14	0.2	1.3	
	scb	2498	-60	13	0.4	1.0	2501	-57	10	0.3	1.0	
	ccb	2497	-61	14	0.4	1.0	2500	-58	11	0.4	1.0	
$-3\frac{1}{2}2$	scs	2339	-321	-272	1.0	1.0	2344	-316	-268	1.0	1.0	$2611(N\Omega)$
	ccs	2287	-373	-324	1.0	1.0	2292	-368	-319	1.0	1.0	$2574(\Lambda\Xi\pi)$
	scb	2510	-313	-193	1.0	0.6	2508	-304	-194	1.0	0.6	
	ccb	2506	-240	-106	1.0	0.6	2504	-242	-107	1.0	0.6	
$-3\frac{1}{2}1$	scs	2339	-321	-272	1.0	1.0	2344	-316	-268	1.0	1.0	$2434(\Lambda\Xi)$
	ccs	2337	-156	-97	1.0	0.9	2340	-154	-94	1.0	0.9	
	scb	2474	-167	-38	1.0	0.6	2437	-131	3	1.0	0.6	
	ccb	2422	-146	11	1.0	0.6	2421	-147	13	1.0	0.6	
-400	scs	2617	-41	-19	0.2	1.4	2622	-37	-15	0.2	1.4	$2636(\Xi\Xi)$
	ccs	2617	-41	-19	0.2	1.4	2622	-37	-15	0.2	1.4	
	scb	2628	-45	-8	0.3	1.1	2630	-43	-6	0.2	1.1	
	ccb	2627	-46	-9	0.3	1.1	2630	-43	-6	0.2	1.1	
-600	scs	3251	-112	-94	0.3	1.4	3260	-102	-84	0.3	1.4	$3345(\Omega\Omega)$
	scb	3376	-65	32	1.0	0.6	3375	-67	30	1.0	0.6	

other example that shows the predictive power of the QDCSM and appears to limit the model uncertainty.

(3) In the nonstrange sector ($S = 0$), the model predicts an $\alpha = (003)$ state with mass $M(003) = 2110(\pm 36)$ MeV. It has the largest binding (320–390 MeV) within the u, d, s three-flavor world. Although its mass is above $NN\pi$ threshold, the transition to $NN\pi$ and NN is hindered by the large angular momentum, so that it is still possible for this to be a narrow resonance. The large delocalization $\epsilon = 1.0$ means this is a true six-quark state. All these results are consistent with our earlier simple relativistic model result [25]. Although the Skyrmion model calculation of Walet does not obtain a binding as large as QDCSM result [26], we note that this model does not obtain sufficient attraction in the N-N channel, either. We know of no reason for the predictive power of the QDCSM shown in the N-N channel to be totally lost in the Δ-Δ channel since both channel results are stable under the same reasonable variations in the values of the model parameters. Therefore, we con-

tinue to recommend this state highly as a good candidate for discovery of a dibaryon resonance.

There is no other interesting channel in the u,d two-flavor world as found earlier by Maltman [9].

(4) Jaffe's H particle remains as the unique strong interaction stable dibaryon in the QDCSM. $M(H) = 2199(\pm 24)$ MeV is 32 MeV lower than the $\Lambda\Lambda$ threshold. The delocalization $\epsilon(s_0) = 1.0$ is also large and the adiabatic channel coupling WF is quite close to Jaffe's pure symmetric flavor singlet basis. However, because of the sensitivity to details of the model, it is not possible to claim that it is indeed a strong interaction stable state.

(5) Another interesting state is the $\alpha = (-3\frac{1}{2}2)$ state [27]. $M(-3\frac{1}{2}2) = 2529(\pm 25)$ MeV is 45 MeV lower than the favorable $(\Lambda\Xi\pi)$ three-body decay threshold. Also because of its large angular momentum, its decay into $\Lambda\Xi$ should be inhibited and so it too might show up as a dibaryon resonance. Another good point about this state is that all the other states with the same quantum number set are about 100 MeV higher than it (not listed in Table II). This might make it a cleaner resonance to observe.

(6) The states $M(-1\frac{3}{2}0) = 2133(\pm 7)$ MeV, $M(-220) = 2390(\pm 9)$ MeV, $M(-3\frac{3}{2}1) = 2512(\pm 15)$ MeV, $M(-3\frac{1}{2}1) = 2469(\pm 58)$ MeV, $M(-400) = 2637(\pm 10)$ MeV, and $M(-600) = 3359(\pm 18)$ MeV all have their masses close to the corresponding thresholds: $N\Sigma$, $\Sigma\Sigma$, $\Sigma\Xi$, $\Lambda\Xi$, $\Xi\Xi$, and $\Omega\Omega$. Other states $M(-1\frac{1}{2}3) = 2318(\pm 33)$ MeV, $M(-213) = 2530(\pm 37)$ MeV, $M(-202) = 2345(\pm 26)$ MeV, and $M(-3\frac{3}{2}3) = 2728(\pm 37)$ MeV have a very large binding and their masses are all less than the favorable multibody channel. These states bear further study. The high-spin dibaryon resonances seem to be especially worth experimental searches, in addition to the spin-zero H particle.

V. CONCLUSION

Since Jaffe's first prediction of the H particle [5], there have been many efforts both theoretically and experimentally to search for dibaryons. Whether all these QCD-inspired models miss some physics when they are extended from the single-hadron to the multihadron case, so that their predictions are not reliable, remains a question. The QDCSM is a more realistic model by taking into account the possible difference of the q-q interaction inside a single baryon and between two color singlet baryons, by allowing each system to choose its favorable configuration in a larger Hilbert space, and by having the model constrained not only by qualitatively fitting hadron spectroscopy but also N-N scattering. This model approach has some moderate success to support it: It predicts two dinucleon states not too far from the experimentally known deuteron and quasideuteron states. If we take this as a measure of the predictive accuracy of the QDCSM, then there are few promising dibaryon resonance candidates within the u,d, and s three-flavor world, as listed in Table II. Because of the simplicity of the model assumptions, the quantitative predictions of the dibaryon masses are uncertain ($\sim$ 10 MeV for nonstrange states and even larger uncertainty for strange states). To get more reliable estimates, especially to be able to determine whether or not the candidate states are strong interaction stable requires improvement of many aspects of the QDCSM:

(1) The QDCSM does not fit the baryon octet and decuplet perfectly, the largest deviation being 41 MeV for the Ξ. Although the adiabatic potential is obtained through a subtraction procedure which suggests that cancellation of errors is possible, there is no guarantee that the uncertainty of the strange baryon mass cancels very accurately.

(2) An adiabatic approximation has been used in this calculation which should be replaced by a dynamical channel coupling calculation.

(3) This calculation is nonrelativistic; a relativistic calculation is underway to estimate relativistic corrections. The preliminary result is that the relativistic and nonrelativistic versions give very similar mass values, especially for the nonstrange states.

(4) Only N-N scattering has been used to constrain the QDCSM. Although data are sparser, Λ-p, and Σ-p, scattering should be used as well. We have begun such an analysis.

(5) The effects of $q\bar{q}$ excitations or quark-meson couplings, which have not been included, may well be important, especially for those states which have a mass close to the lowest two-body threshold.

(6) It would be interesting to include c and b quarks with a view towards making contact with heavy quark effective theory. However, the large quark WF difference between u,d and b,c would need to be treated first. This is unlike the s-quark case where the single-quark WF distortion is not large.

We believe the QDCSM results support the value of investing additional effort, both theoretical and experimental, using more sophisticated approaches, in order to concentrate on a few promising dibaryon candidates.

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APPENDIX

In the dynamical N-N scattering calculation [2], we first calculate the interaction kernel

$$K(\vec{S},\vec{S}) = \frac{\langle\Psi(\vec{S})|H|\Psi(\vec{S})\rangle}{\langle\Psi(\vec{S})|\Psi(\vec{S})\rangle}\ , \tag{A1}$$

where H is the six-quark Hamiltonian and $\Psi(\vec{S})$ is the N-N channel WF (9). Then we do a partial wave decomposition

$$k_l(S,S') = \int d\Omega P_l(\theta_{ss'})K(\vec{S},\vec{S}'), \tag{A2}$$

where $\theta_{ss'}$ is the angle between $\vec{S}$ and $\vec{S}'$. We then as-

sume the diagonal matrix elements $k_l(s,s)$ as the effective interaction of the l partial wave between two nucleons,

$$V_l(s) = k_l(s,s) \ . \tag{A3}$$

This effective interaction is dependent on the delocalization parameter $\epsilon(s)$. We vary the value $\epsilon(s)$ to get a minimum for each separation s. The $\epsilon(s)$ so determined is l dependent; i.e., we have $\epsilon_l(s)$. Next we substitute the values $\epsilon_l(s)$ and $\epsilon_l(s')$ back into Eq. (A2) to get the final $k_l(s,s')$ for the N-N scattering calculation. Finally, we adjust the screening parameter μ or ν of Eqs. (6a) and (6b) to get the best fit to the $^1S_0$ and $^3S_1$ N-N phase shifts. This determines the best values, $\mu = 1.0$ fm^{-2}, $\nu = 0.4$ fm^{-2}, because it is a more complete and consistent calculation. However, it takes much more computer time than the next calculation we describe.

To minimize computer time for systematic dibaryon search, we tried another approximation, assuming $K(\vec{s},\vec{s})$ as the effective interaction between two nucleons,

$$V(s) = K(\vec{s},\vec{s}) \ . \tag{A4}$$

Then, we varied $\epsilon(s)$ to minimize $V(s)$ and so determine the $\epsilon(s)$ as well. This $\epsilon(s)$ is partial wave independent. Substituting the $\epsilon(s)$ and $\epsilon(s')$ values so determined back to Eq. (A1), we obtain the final $K(\vec{s},\vec{s}')$ and then do a partial wave decomposition to calculate the phase shifts. Adjusting the screening parameters again to obtain the best fit, the $^1S_0$ and $^3S_1$ channels, we obtain a second set of values: $\mu = 1.6$ fm^{-2}, $\nu = 0.6$ fm^{-2}. The approximation is not as good as the first one, but also gives a qualitatively good fit to the N-N phase shifts.

In the dibaryon calculation, we use the second variation method (variation before partial wave decomposition) to obtain the effective interaction between two baryons. The results are shown in Tables II(a) and II(c). As a check on the range of variation, we also used the intermediate value $\nu = 0.46$ fm^{-2}, and those results are shown in Table II(b).

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