The $d'$ dibaryon in the quark-delocalization, color-screening model

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Abstract

We study the questions of the existence and mass of the proposed $d'(IJ^P = 00^-)$ dibaryon in the quark-delocalization, color-screening model (QDCSM). The transformation between physical and symmetry bases has been extended to the cases beyond the SU(2) orbital symmetry. Using parameters fixed by baryon properties and $NN$ scattering, we find a mild attraction in the $IJ^P = 00^-$ channel, but it is not strong enough to form a deeply bound state as proposed for the $d'$ state. Nor does the (isospin) $I=2$ NΔ configuration have a deeply bound state. These results show that if a narrow dibaryon $d'$ state does exist, it must have a more complicated structure.

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Quantum chromodynamics (QCD) has been accepted as the fundamental theory of the strong interaction. Understanding the low energy behavior of QCD and the nature of the strong interactions of matter, however, remains a challenge. Lattice QCD has provided numerical results describing quark confinement between two static colorful quarks, a preliminary picture of the QCD vacuum and the internal structure of hadrons in addition to a phase transition of strongly interacting matter. Phenomenological quark model analyses of hadron spectroscopy have also provided useful physical information. However, the color structures available in $q\bar{q}$ and $q^3$ systems are limited. Multiquark systems involve more complicated color structures which can not be studied directly in meson and baryon systems. A simple example is given by three gluon exchange [1] and the three body instanton interaction [2] both of which do not contribute within a colorless meson or baryon but do contribute to a multiquark system. Therefore multiquark systems are indispensable for the full study of the low energy behavior of QCD and the structure of strongly interacting matter.

This report is limited to the dibaryon or $q^6$ systems. Since Jaffe predicted the H particle [3], the study of dibaryons has been waxed and waned. At the end of the 1970’s and beginning of the 1980’s, many dibaryon states were predicted based on the MIT bag model and some of them were even claimed to have been observed experimentally. However further measurement has almost completely dismissed all of them. All quark models insist that there should be dibaryon states but experimental signals are scarce, at best. The search for the H particle has been continued for more than twenty years with no indication of its existence. Lomon predicted a high mass NN $I^=1^ 1S_0$ resonance at around 2.7 GeV [4] which does seem to be supported by SATURNE pp scattering data [5]. The Moscow-Tübingen-Warsaw-Uppsala collaboration [6] has claimed a narrow dibaryon resonance, $d'$, centered at 2.06 GeV with a small width of 0.5 MeV; the preferred quantum numbers are $IJ^P = 00^-$. In contrast to all the other cases, the $d'$ dibaryon, with a mass as small as 2.06 GeV and $IJ^P = 00^-$, is hard to accommodate by the available quark models [7].
A new quark model, the quark-delocalization, color-screening model (QDCSM), has been
developed with the aimed of understanding the well known similarities between nuclear and
molecular forces despite the obvious energy and length scale differences [8]. This model
predicts a small mass narrow dibaryon resonance $d^*$ with $IJ^P = 0^+\,$, $M\sim 2.1$ GeV, $\Gamma(NN) \sim 1$ MeV [9]. This paper reports our search for the $d'$ dibaryon using the QDCSM. Sect.II
gives a brief description of the Hamiltonian and wave function in QDCSM; Sect.III describes
the calculation method. For this, a very useful extension of the transformation between the
physical bases and symmetry bases beyond the SU(2) orbital symmetry is presented. The
results are presented in Sect.IV. The final section provides a summary.

II. MODEL HAMILTONIAN AND WAVE FUNCTION

The details of the QDCSM can be found in Ref. [8,9]. Here we present only the model
Hamiltonian and wave functions used in the calculation.

The Hamiltonian for the 3-quark system is the same as the usual potential model. For
the six-quark system, it is assumed to be

$$
H_6 = \sum_{i=1}^{6} (m_i + \frac{p_i^2}{2m_i}) - T_{CM} + \sum_{i<j=1}^{6} \left( V_{ij}^c + V_{ij}^G \right),
$$

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

$$
V_{ij}^c = -a_c \bar{\chi}_i \cdot \bar{\chi}_j \left\{ \begin{array}{ll}
\frac{r_{ij}}{1-e^{-\mu r_{ij}}} & \text{if } i, j \text{ occur in the same baryon orbit}, \\
\frac{1-e^{-\mu r_{ij}}}{\mu} & \text{if } i, j \text{ occur in different baryon orbits},
\end{array} \right.
$$

where all the symbols have their usual meaning except the confinement potential $V_{ij}^C$ which
will be explained below.

The wave function of the six-quark system is written as

$$
|\Psi_{6q}\rangle = A[\Psi_{B_1} \Psi_{B_2}]_{WcM_iM_j}^{[\sigma]IJ},
$$

where $\Psi_{B_i}$ is the 3-quark system wave function,

$$
\Psi_{B_1} = [[\psi_L(1)\psi_L(2)\psi_L(3)]^i]_{l_1} \eta_{i_1} [j^i ]_{m_{ij}} x_c(123)
$$
and

$$\Psi_{B_2} = [[\psi_R(4)\psi_R(5)\psi_R(6)]^2J_2\eta_2S_2J_{m_2}^2\chi_c(456)].$$

(4)

The single particle orbital wave functions are delocalized, as

$$\psi_L = (\phi_L + \epsilon_s\phi_R)/N_s,$$

$$\psi_R = (\phi_R + \epsilon_s\phi_L)/N_s,$$

$$\phi_L = \left(\frac{1}{\pi b^2}\right)^{3/4}e^{-\frac{1}{2b^2}(\vec{r}-\vec{s}_0/2)^2},$$

$$\phi_R = \left(\frac{1}{\pi b^2}\right)^{3/4}e^{-\frac{1}{2b^2}(\vec{r}+\vec{s}_0/2)^2},$$

$$N_s = [1 + \epsilon_s^2 + 2\epsilon_s\epsilon_p^2/4b^2]^{1/2}$$

if the particle is in s-wave and as

$$\psi_L = (\phi_L + \epsilon_p\phi_R)/N_p,$$

$$\psi_R = (\phi_R + \epsilon_p\phi_L)/N_p,$$

$$\phi_L = \left(\frac{1}{\pi b^2}\right)^{3/4}\frac{\sqrt{2}}{b}|\vec{r}+\vec{s}_0/2|Y_{1m}e^{-\frac{1}{2b^2}(\vec{r}+\vec{s}_0/2)^2},$$

$$\phi_R = \left(\frac{1}{\pi b^2}\right)^{3/4}\frac{\sqrt{2}}{b}|\vec{r}−\vec{s}_0/2|Y_{1m}e^{-\frac{1}{2b^2}(\vec{r}−\vec{s}_0/2)^2},$$

$$N_s = [1 + \epsilon_p^2 + 2\epsilon_p(1 − \frac{\vec{s}_0^2}{2b^2})\epsilon_p^2/4b^2]^{1/2}$$

if the particle is in p-wave, where $\vec{s}_0$ is the separation of the two clusters.

Although we use a potential model language in the description of the confinement potential $V_{ij}^C$ in Eq.(1), our calculations employ an extended effective matrix element method. Even though understanding of quark confinement is limited, the use of a two body interaction to describe the quark confinement may well be highly oversimplified, especially for multiquark systems. For example, the three-gluon and three-body instanton interactions mentioned above [1,2] cannot be expressed in terms of two-body interactions and the full nonperturbative QCD interaction contains additional varieties, such as condensates, which can not be expressed in terms of two body interactions either. Our confinement “potential”
as defined in Eq. (1) has the meaning of the usual interaction potential only in the asymptotic region where the overlap of the $\phi_L$ and $\phi_R$ wave function orbitals is negligible. In the interaction region, it defines a recipe for determining the effective matrix elements of the Hamiltonian of a six quark system. (More details will be provided in the next section.) Our goal is to include the major portion of nonperturbative many-body quark interactions by this method.

We refer to the six-quark wave function defined in Eq. (2) as the physical basis because it has an obvious physical meaning in terms of a pair of baryons. In our channel coupling calculation, all of the colorless p-wave excitation baryon-baryon channels compatible with the $d'$ quantum numbers $IJ^P = 00^-$ are included. Although the six-quark wave function, Eq. (2), appears to have the form of $q^3 - q^3$ clusters, in fact the other clusters such as $q^6, q^5q$ and $q^4q^2$ have been included because of the quark delocalization. Hidden color channels have been omitted because we do not know anything about the mass and interaction of colorful baryons and because a hidden color channel can be expressed in terms of a sum of colorless channels [10].

III. CALCULATION METHOD

The physical basis used in Eq. (2) is not convenient for the calculation of matrix elements of the six-quark Hamiltonian. In order to simplify the calculation, the physical basis (cluster basis) is expanded in terms of symmetry bases first. Next, we use the powerful fractional parentage (fp) expansion method to calculate the matrix elements between symmetry basis components of six-quark system. Finally, the matrix elements between physical basis components are obtained by reversing the transformation between the physical and symmetry bases.

The transformation between the physical basis and symmetry basis has been studied by Harvey [11], Chen [12], and ourselves [13] in the case where the orbital symmetry of the three-quark cluster was limited to $[3]$ and $SU_x(2)$ orbital symmetry was assumed. In the
case of interest here, there is a p-wave quark with respect to the right- or left-center, in addition to right- and left-centered s-wave quarks. The orbital symmetry group needs to be extended to $SU^x(4)$. (Only one state among the three p-wave states is included in the calculation of the transformation coefficients, since the spin-orbit coupling is a trivial one.) The orbital symmetry of individual baryons should include both $[3]$ and $[21]$ configurations. This requires and extension of the transformation method studied previously. Here we develop two new methods to calculate the transformation coefficients.

In the symmetry basis, the group-chain classified basis is denoted by

$$|\Phi_S\rangle = \begin{pmatrix} [\nu]W_x \\ [\sigma]W_c[\mu]IMJ\, SLMJ \end{pmatrix},$$

where $[\nu], [\sigma]$ and $[\mu]$ represent the symmetries of orbital, color and spin-isospin features, and $W_x, W_c$ are Weyl tableaux for $[\nu]$ and $[\sigma]$. Other symbols have their usual meanings. The group chain we use here is

$$SU^{x\tau\sigma}(36) \supset SU^x(3) \times SU^{c\tau\sigma}(12) \supset SU^c(3) \times SU^{\tau\sigma}(4) \supset SU^\tau(2) \times SU^\sigma(2).$$

The physical basis is constructed from two 3q states,

$$|\Psi_P\rangle = A[B_1B_2[\sigma]^{ISLJ}W_cM_J,]$$

where $B_1$ and $B_2$ represent the wave functions of two three-quark clusters. The physical basis defined here is a little different from the $\Psi_{6q}$ defined in Eq.(2). The two bases are related by Racah coefficients. Fortunately, for the quantum numbers of $d'$, all of the coefficients equal unity. From now on, we replace $\Psi_{6q}$ by $\Psi_P$.

The physical basis can be expanded in terms of the symmetry basis as

$$|\Psi_P\rangle = \sum_{\nu W_x \mu} C_{P-s} |\Phi_S\rangle.$$ 

By using the unitary condition of the symmetry basis, the expansion coefficients can be expressed as,
\[ C_{\mu-s} = \langle \Phi_S | \Psi_P \rangle. \]  

(11)

On the other hand, the six-quark symmetry basis can be expanded into two three-quark clusters by application of the fractional parentage technique:

\[
\begin{align*}
|\Phi_S\rangle &= \sum_{1,2} C^{[1\sigma],[\nu][\bar{\nu}]}_{[1^3][\bar{\nu}_1][\bar{\nu}_2][\bar{\nu}_2]} C^{[\bar{\nu}],[\bar{\nu}][\sigma],[\bar{\nu}]}_{[1^3][\bar{\nu}_1][\bar{\nu}_1][\bar{\nu}_1]} C^{[\bar{\nu}],[\sigma],[\mu]}_{[1^3][\bar{\nu}_1][\bar{\nu}_1][\bar{\nu}_1]} C^{[\mu],IS}_{[\mu_1]I_1S_1,[\mu_2]I_2S_2} C^{[\nu]W_{x_1}}_{[\nu_1]W_{x_1},[\nu_1]} W_{x_2} \\
&\quad \times C^{[\sigma]W_{x_2}}_{[\sigma_1]W_{x_2},[\sigma_2]} W_{x_2} C^{L_{M_1L_1}}_{L_{M_1L_1}} C^{I_{M_1}}_{I_{M_1}} C^{I_{M_2}}_{I_{M_2}} C^{S_{M_1}S_{M_2}}_{S_{M_1}S_{M_2}} C^{[\nu_1]W_{x_2}}_{[\nu_2]W_{x_2}} \\
&\quad \times [B_1B_2]^{[\sigma]ISLM}_{W_{x1}W_{x2},M_1M_2}.
\end{align*}
\]

(12)

where 1, 2 stand for \( \nu_1, \nu_2, W_{x_1}, W_{x_2}, \mu_1, \mu_2, \ldots \), etc. The last line is just the wave functions of two three-quark clusters. Combining with the relevant Clebsch-Gordan (CG) coefficients, we have

\[
|\Phi_S\rangle = \sum_{1,2} C^{[1\sigma],[\nu][\bar{\nu}]}_{[1^3][\bar{\nu}_1][\bar{\nu}_2][\bar{\nu}_2]} C^{[\bar{\nu}],[\bar{\nu}][\sigma],[\bar{\nu}]}_{[1^3][\bar{\nu}_1][\bar{\nu}_1][\bar{\nu}_1]} C^{[\bar{\nu}],[\sigma],[\mu]}_{[1^3][\bar{\nu}_1][\bar{\nu}_1][\bar{\nu}_1]} C^{[\mu],IS}_{[\mu_1]I_1S_1,[\mu_2]I_2S_2} C^{[\nu]W_{x_2}}_{[\nu_1]W_{x_2},[\nu_1]} W_{x_2} \\
&\quad \times [B_1B_2]^{[\sigma]ISLM}_{W_{x1}W_{x2},M_1M_2}.
\]

(13)

We next apply the inter-cluster antisymmetrization operator

\[
\mathcal{A} = \sqrt{\frac{1}{20}} \sum_{P:\text{inter-cluster permutations}} (-)^P P
\]

to Eq.(13). Because of the total antisymmetrization of the six-quark symmetry basis, when the operator \( \mathcal{A} \) is applied to the left-hand side of Eq.(13), it produces 20 copies. In this way, we obtain

\[
\sqrt{20} \left| [\nu]W_{x_2} \right| [\sigma]W_{c}[\mu]IMSL_{M_1M_2} = \sum_{1,2} C^{[1\bar{\nu}],[\bar{\nu}][\bar{\nu}][\bar{\nu}]}_{[1^3][\bar{\nu}_1][\bar{\nu}_1][\bar{\nu}_1]} C^{[\bar{\nu}],[\bar{\nu}][\sigma],[\bar{\nu}]}_{[1^3][\bar{\nu}_1][\bar{\nu}_1][\bar{\nu}_1]} C^{[\bar{\nu}],[\sigma],[\mu]}_{[1^3][\bar{\nu}_1][\bar{\nu}_1][\bar{\nu}_1]} C^{[\mu],IS}_{[\mu_1]I_1S_1,[\mu_2]I_2S_2} C^{[\nu]W_{x_2}}_{[\nu_1]W_{x_2},[\nu_1]} W_{x_2} \mathcal{A}[B_1B_2]^{[\sigma]ISLM}_{W_{x1}W_{x2},M_1M_2}.
\]

(14)

Substituting Eq.(14) into Eq.(11), and making use of the orthonormal property of the physical bases, we obtain the expression for the transformation coefficients,

\[
C_{\mu-s} = \sqrt{20} C^{[1\sigma],[\nu][\bar{\nu}]}_{[1^3][\bar{\nu}_1][\bar{\nu}_1][\bar{\nu}_1]} C^{[\bar{\nu}],[\bar{\nu}][\sigma],[\bar{\nu}]}_{[1^3][\bar{\nu}_1][\bar{\nu}_1][\bar{\nu}_1]} C^{[\bar{\nu}],[\sigma],[\mu]}_{[1^3][\bar{\nu}_1][\bar{\nu}_1][\bar{\nu}_1]} C^{[\mu],IS}_{[\mu_1]I_1S_1,[\mu_2]I_2S_2} C^{[\nu]W_{x_1}}_{[\nu_1]W_{x_1},[\nu_1]} W_{x_2}
\]

(15)
This expression differs from the transformation coefficients obtained before [12,13]. It has two "new" factors: $C^{[16],[\nu][\bar{\nu}]}_{[1^3] [\nu_1][\bar{\nu}_1],[1^3] [\nu_2][\bar{\nu}_2]}$ and $C^{[\nu]W_x}_{[\nu_1]W_{x_1},[\nu_2]W_{x_2}}$. In the limit of $SU^x(2)$ with $[\nu_1] = [\nu_2] = [3]$, the product of these two factors is a constant with the value, $\sqrt{\frac{1}{20}}$, which returns us to the expression of Ref. [12,13].

Recently Chen [14] also proposed a method to generalize the transformation between physical and symmetry bases. Our new method is based on his expressions in Eqs.(9-20a) of Ref. [15],

$$
\mathcal{A} \begin{bmatrix} [\nu_1]a^f_1 \\
\sigma_1 \mu_1 I_1 S_1 \end{bmatrix}_{\sigma IS} \begin{bmatrix} [\nu_2]a^f_2 \\
\sigma_2 \mu_2 I_2 S_2 \end{bmatrix}_{\sigma IS} = \sum_{[\nu]} C^{[\nu],[\sigma][\mu]}_{[\nu_1][\sigma_1][\mu_1],[\nu_2][\sigma_2][\mu_2]} C^{[\nu],IS}_{[\mu]I_1 S_1,[\mu]I_2 S_2} \begin{bmatrix} [\nu]a^f_1 b^f_2 \\
\sigma_\mu IS \end{bmatrix},
$$

To extend this to the case of interest here, first set $a^f_1$ to $W_{x_1}$, $b^f_2$ to $W_{x_2}$. The symmetry basis used in Eq.(16) is not the one which we defined previously. Rather, it is an $SU_4 \supset SU_2 \times SU_2$ irreducible basis, which is a non-standard $SU_4$ basis, with respect to the orbital symmetry and must be expanded in terms of the standard Gel’fand bases of $SU_4$ and the Yamanouchi bases of $S_6$. It is only in the special $SU^x(2)$ case that the Weyl tableau, $W_x$, is automatically fixed by the orbital symmetry $[\nu]$. In general a further transformation from non-standard to standard $SU_n$ bases is needed.

$$
\begin{bmatrix} [\nu] & [\nu_1][\nu_2] \\
m & W_{x_1} W_{x_2} \end{bmatrix} = \sum_{W_x} \begin{bmatrix} [\nu] & [\nu_1][\nu_2] \\
m & W_{x_1} W_{x_2} \end{bmatrix} \begin{bmatrix} [\nu] & [\nu_1][\nu_2] \\
W_x & W_{x_1} W_{x_2} \end{bmatrix},
$$

where the transformation coefficients, termed subduction coefficients (SDC) of $SU_4$, are independent of $m$. From the SDC and isoscalar factor tables of $SU_4$ and CG coefficient tables of $SU_2$, it is easy to verify that the SDC used here is just the product of $\sqrt{20}$, $C^{[16],[\nu][\bar{\nu}]}_{[1^3] [\nu_1][\bar{\nu}_1],[1^3] [\nu_2][\bar{\nu}_2]}$ and $C^{[\nu]W_x}_{[\nu_1]W_{x_1},[\nu_2]W_{x_2}}$. So the two methods give the same transformation coefficients as they should.

For the quantum numbers of $d'$, there are six physical channels, (taking only color singlet baryons into account,) which are given in the first column of Table I, where: $N$ denotes the
state of nucleon with quantum numbers \([\nu_i] = [3], I_i = \frac{1}{2}, S_i = \frac{1}{2}, l_i = 0, j_i = \frac{1}{2}; \Delta\) denotes \([\nu_i] = [3], I_i = \frac{3}{2}, S_i = \frac{3}{2}, l_i = 0, j_i = \frac{3}{2}; N_1^*\) denotes \([\nu_i] = [21], I_i = \frac{1}{2}, S_i = \frac{1}{2}, l_i = 1, j_i = \frac{1}{2}; N_2^*\) denotes \([\nu_i] = [21], I_i = \frac{1}{2}, S_i = \frac{3}{2}, l_i = 1, j_i = \frac{1}{2}; \Delta^*\) denotes \([\nu_i] = [21], I_i = \frac{3}{2}, S_i = \frac{1}{2}, l_i = 1, j_i = \frac{3}{2}. \) The corresponding symmetry bases are given in the first row of Table I. There are thirteen bases. The reason we have a \(6 \times 13\) table is that the hidden color channels have been omitted. The \(d'\) is assumed to be a linear combination of these six physical channels. The combination coefficients are determined by diagonalizing the Hamiltonian in the 6-dimensional space.

To calculate the matrix elements of the six-quark Hamiltonian between symmetry bases, we use a fractional parentage expansion. Then only the two-body matrix elements and four-body overlaps are required. Details can be found in Ref. [13]. By using the transformation coefficients derived above, the matrix elements of the six-quark Hamiltonian between physical bases can be obtained straightforwardly.

In the calculation of the two body matrix elements of the confinement interaction, we assume that the normal confinement interaction should be used for \(\langle LL|V_C|LL\rangle, \langle RR|V_C|RR\rangle, \langle LR|V_C|LR\rangle\), while the color-screened confining interaction should be used for \(\langle LR|V_C|RL\rangle\) and all others. Here

\[
\langle LL|V_C|LL\rangle = \langle \phi_L(\vec{r}_1)\phi_L(\vec{r}_2)|V_{1,1}^C|\phi_L(\vec{r}_1)\phi_L(\vec{r}_2)\rangle.
\]

This recipe fixes the matrix elements of the confinement interaction and distinguishes our extended effective matrix element approach from a two-body potential-model approach. It reduces to the usual two-body potential-model for a single hadron and therefore maintains the successes of the constituent quark model for hadron spectroscopy. It also reproduces qualitatively correct phase shifts in ten channels: \(NN\) \(ST = 10, 01, 00, 11, N\Lambda\) \(ST = 0\frac{1}{2}, 1\frac{1}{2}, N\Sigma\) \(ST = 0\frac{1}{2}, 1\frac{1}{2}, 0\frac{3}{2}, 1\frac{3}{2}, 3\frac{3}{2}\), with only one additional adjustable parameter, the color screening constant \(\mu\) [8,16]. It is the only model which demonstrates a similarity between nuclear and molecular forces. We take these successes as an indication that this effective matrix element approach includes significant components of correct physics.
For a preliminary study, we use the adiabatic approximation. That is, for each separation $s_0$, we determine the energy of the six-quark system by the variational condition

$$\frac{\partial E_6}{\partial \epsilon_l} = 0, \quad l = s, p$$

(18)

The effective potential between two baryons is obtained by a subtraction [17],

$$V_e(s_0) = E_6(s_0) - E_6(\infty).$$

(19)

and the mass of the six-quark system is estimated by the formula,

$$M_6 = m_1 + m_2 + V_e + E_0$$

(20)

where $m_1, m_2$ are the masses of the two baryons in isolation and $E_0$ is the zero-point energy of the pair, $E_0 = \frac{4k^2}{3m^2}$, with $m$ the reduced mass of the two baryons.

**IV. RESULTS AND DISCUSSION**

The model parameters, which are fixed by the baryon spectrum and $NN$ scattering, are:

$$m_u = m_d = 313 \text{ MeV}, \quad b = 0.603 \text{ fm}, \quad a = 25.13 \text{ MeV/fm}^2, \quad \alpha_s = 1.54, \quad \mu = 1.6 \text{ fm}^{-2}.$$ 

Both single channel and channel coupling calculations were carried out. The results are shown in Table II. From these results, we observe a mildly attractive interaction between $N$ and $N^*$, $\sim 100$ MeV. Since the delocalization is approximately 1, this is a six-quark state, not a two-baryon state. However, this attraction is not enough to form a deeply bound state such as $d'$. Channel coupling adds a little more attraction ($\sim 10$ MeV), does not change this conclusion.

These results are expected in the QDCSM. From our previous study of the effective potential between baryons [17], we found that as a general trend there are strong attractions in decuplet-decuplet channels, but the attraction in octet-octet channels is weak. The excited $N^*$ state of the nucleon is still in a flavor octet, so the attraction between $N$ and $N^*$ would not be expected to be strong, although the existence of the $p$-wave quark adds a little more
attraction. The mass of $\Delta \Delta^*$ channel is much larger than that of the $NN^*$ channels, so the effect of channel-coupling to it is small.

The effect of the harmonic oscillator parameter, $b$, of the quark orbital wave function has also been studied. G. Wagner et al. [18] concluded that it is impossible to describe the dibaryon, $d'$, with the same parameters as those used for single baryons. In particular, they found it impossible to get a $d'$ mass as low as 2.06 GeV if the same $b$ is used for both 3-quark and six-quark systems. To check their statement, the results obtained with a larger value of $b$ in the six-quark calculation are also shown in Table II. By varying $b$, the attraction between $N$ and $N^*$ could be increased. However, if a different $b$ is used in the 3-quark and six-quark calculations, the subtraction procedure, shown in Eq.(19) to obtain the effective potential, is no longer reliable. A dynamic calculation is needed. It is possible that a deeply bound $d'$ state might be obtained in the QDCSM if a larger parameter $b$ is used for the six quark system. However, the physical meaning of such a model calculation is obscure.

The possibility that the $d'$ has isospin $I = 2$ has been discussed in the literature. We have carried out both adiabatic and dynamical $IJ^P = 20^-$ $N\Delta$ single channel calculations. With the same model parameters, the adiabatic calculation produces a mass of 2173 MeV for the $d'$. A dynamical calculation with eight Gaussians (spanning cluster separation coordinate values from 0.6 to 4.8 fm produces a mass of 2191 MeV and one with fifteen Gaussian (spanning cluster separation coordinate values from 0.6 to 9.0 fm produces a mass of 2178 MeV. These trends strongly suggest that this is a scattering state, not a bound state. So we conclude that no state with a mass as low as 2.06 GeV can be obtained in the $N\Delta$ channel in the QDCSM.

V. SUMMARY

We have studied the proposed $d'$ state in the QDCSM using the physical bases $NN^*$ and $\Delta\Delta^*$. Our results show that there a mild attraction does develop in the $IJ^P = 00^-$ state, but this attraction is not strong enough to form a deeply bound state with a mass as low as
2.06 GeV. All of the p-wave excitation baryon-baryon channels with $d'$-compatible quantum numbers have been included in our channel coupling calculations. Because the QDCSM includes quark delocalization, all of the configurations, $q^6, q^5q, q^4q^2$ and $q^3q^3$, are included in our model space.

Nor do we find such a low mass state in the $I = 2 N\Delta$ channel in our model approach.

Including the $d'$ results of the Faessler group [7,18], we conclude that, if the $d'$ is experimentally verified as a real dibaryon resonance, it must have a more complicated structure than studied here, such as including $q^7\bar{q}$ components.

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REFERENCES


### TABLE I. The transformation coefficients between physical and symmetry bases.

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<th>[51]_{1}[321]</th>
<th>[51]_{2}[321]</th>
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<td>(NN_{2}^{*})</td>
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<td>(\sqrt{\frac{1}{180}})</td>
<td>(-\sqrt{\frac{4}{81}})</td>
<td>(-\sqrt{\frac{16}{405}})</td>
<td>(-\sqrt{\frac{8}{81}})</td>
<td>(\sqrt{\frac{32}{405}})</td>
<td>(\sqrt{\frac{5}{324}})</td>
</tr>
<tr>
<td>(\Delta\Delta^{*})</td>
<td>(\sqrt{\frac{1}{45}})</td>
<td>(-\sqrt{\frac{1}{180}})</td>
<td>(-\sqrt{\frac{1}{81}})</td>
<td>(\sqrt{\frac{4}{405}})</td>
<td>(-\sqrt{\frac{8}{81}})</td>
<td>(\sqrt{\frac{32}{405}})</td>
<td>(-\sqrt{\frac{5}{324}})</td>
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<tr>
<td>(N_{1}^{*}N)</td>
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</table>

### TABLE II. Effective potentials and masses of six-quark systems.

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<tr>
<th></th>
<th>(b)(fm)</th>
<th>(s_0)(fm)</th>
<th>(\epsilon_{s})</th>
<th>(\epsilon_{p})</th>
<th>(E_0(s_0)(\text{MeV}))</th>
<th>(E_0(\infty)(\text{MeV}))</th>
<th>(V_{c}(\text{MeV}))</th>
<th>(M_{6}(\text{MeV}))</th>
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</thead>
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<td>s.c.</td>
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<td>2364</td>
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<td>1.0</td>
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<td>2473</td>
<td>-109</td>
<td>2454</td>
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<tr>
<td>c.c.</td>
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<td>0.9</td>
<td>1.0</td>
<td>1.0</td>
<td>2123</td>
<td>2473</td>
<td>-109</td>
<td>2454</td>
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