Strangeness −3 Dibaryons

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(Received 12 May 1987)

We demonstrate, using two different quark models of hadrons, that there should be isodoublets of dibaryons with strangeness −3 and $J=1,2$, which are stable with respect to strong decay.

PACS numbers: 14.20.Pt, 12.40.Aa

In any quark model of hadrons, the prediction of the energy of most dibaryons, whether bound or not, depends in a critical way upon the contribution of the color-magnetic interaction. In this Letter, we wish to point out that there should be bound states with strangeness −3, which, because of their particular structure, have very little contribution from the color-magnetic interaction. Because of this, the prediction of these bound states occurs in very different models and one may expect the energy estimates to be relatively reliable. We shall pursue the general argument and then estimate the energy in two different models. The first, which we believe to be more reliable, is a potential model,[1-4] and the second is the Massachusetts Institute of Technology (MIT) bag model[5-8] in the spherical-bag limit.

In either model, the one-body quark energy will go down as a result of the expansion of the wave function into a larger volume and this effect provides binding energy. For a given pair of baryons (or a given dibaryon) this may be augmented or counteracted by the two-body color-magnetic interaction (the color-electric interaction is modeled by the potential or the bag).

The general argument follows from the behavior of the color-magnetic spin-splitting (CMSS) matrix elements under symmetry transformations. The form of this interaction is

$$H_I = -(g^2/4\pi)\boldsymbol{\alpha}_i\cdot\boldsymbol{\alpha}_j\lambda_i^a\lambda_j^a\Delta(r_i,r_j), \qquad (1)$$

where $g^2/4\pi$ $(=\alpha_s)$ is an effective strength parameter (related to the average of the QCD coupling constant at intermediate distance scales), $\boldsymbol{\alpha}_i$ are the Dirac α matrices for the ith quark, λ_i^a are the SU(3) color matrices, and $\Delta(r_i,r_j)$ is a gluon propagator from $\mathbf{r}_i$ to $\mathbf{r}_j$ in the background field. To understand the symmetry properties of H_I, one may treat $\boldsymbol{\alpha}$ as if it were a Pauli spin matrix $\boldsymbol{\sigma}$. Then, for quarks in the same spatial wave function (or in the symmetric spatial wave function), the sign of H_I is given in Table I. It is clear that this interaction is always repulsive between two quarks of the same flavor in this spatial configuration.

The CMSS is mainly responsible for the splitting between the octet and the decuplet for low-mass baryons. In these baryons, which are color singlets, any pair of quarks is in a $\mathbf{3}^*$ color state. For the decuplet, the states are symmetric under flavor and are, therefore, raised in energy. The octet dibaryons have both flavor symmetries and are net attractive.

If one attempts to make a dibaryon from two members of the octet, two effects occur for the CMSS, both usually repulsive. First, because of the short effective range of the gluon propagator in the nucleon, the size of the matrix elements between quarks originally in one hadron will decrease as the quark wave functions are modified to occupy a larger volume. (Note that this makes no reference to the detailed shape of the volume.) This decreases an attractive interaction and works against binding. Second, there is an interaction between quarks initially in different hadrons. Since the gluon transforms as a color octet, this latter interaction must be accompanied by quark exchange between the two color-singlet subsystems to maintain their individual neutrality. The effect of this interaction depends in detail on the system, pro-

TABLE I. The sign of the color-magnetic spin-splitting matrix element for symmetric space states. The sign correlation flips for antisymmetric space states.

Color state	Spin	Flavor symmetry	Sign (− for attraction)
$\mathbf{3}^*$	0	A	−
$\mathbf{3}^*$	1	S	+
$\mathbf{6}$	0	S	+
$\mathbf{6}$	1	A	−

TABLE II. Coefficients for the color-magnetic spin-splitting spatial matrix element in the lowest SU(6) configuration for the given angular momentum J and strangeness S. In this convention, the corresponding weight for a Δ is +6.

J	S	Wt
2	−3	−3
1	−3	−7
0	−2	−18

ducing repulsion in the deuteron channel and additional attraction in the H dibaryon,[9] which is a singlet in spin and flavor.

Consider now the special case of a nucleon and an Ω^-. The N-Ω^- is a combination of an octet member and a decuplet member, so that the first effect of the CMSS (referred to above) cancels to a good approximation. Furthermore, since there are no quarks of like flavor in the two hadrons, diagonal quark exchange is not allowed, and so the second effect does not contribute. Therefore, to the extent that the physical system is dominated by the N-Ω^- channel (which is the lowest $S=-3$, $J=2$ channel), there is no color-magnetic effect and the energetics are dominated by modifications to the single-quark wave function.[1]

The approximation that the N-Ω^- channel dominates will be better in the $J=2$ state, since it has no s-wave components from the octet-octet sector. For $J=1$, with more components available, the CMSS should have a greater effect. This is borne out by a SU(6) analysis, from which we determine the coefficients of the CMSS spatial matrix elements in Table II.

To estimate the binding energy of these dibaryons, we use a model which has been developed to study nuclei.[3] In this model, relativistic quarks are confined in color-singlet hadrons by means of a linear scalar potential which is an approximation to the potentials of Buchmüller and Tye.[10] The wave functions for these quarks are solutions of the Dirac equation with this potential.

The potential from two separated hadron centers goes to the single-hadron value at large distance and is truncated at the value from either hadron at the midplane. (It is W shaped along the line determined by the centers.) It therefore reflects the mean-field concept of screening, in that the potential acting on a quark is always the single-hadron potential from the nearest center. Since the potential in between the centers is finite, tunneling of the quark wave function from the region of one center to the other is allowed. The extension to $A=3$ and 4, allowing for a space-spin-color clustering to minimize the color-magnetic interaction, gives very encouraging results.[3]

For the case at hand, we vary the distance between centers and calculate the single-quark energy using a trial function which is the normalized symmetrized sum of the lowest eigenfunctions in each well. Since the trial function is not an eigenfunction, some negative-energy components may be admixed. To ensure a valid variational estimate, we actually calculate the expectation value of H^2, which is positive semidefinite.

Two calculation methods are used. The first uses Monte Carlo integration for both one- and two-body matrix elements and a Yukawa propagator for the gluons. In the second, Gaussian fits to the wave function (to upper and lower components independently) and a Gaussian gluon propagator are used for analytic calculation of the two-body matrix elements.[11] The one-body matrix elements of H^2 still require some numerical integration, which is done by Gauss-Laguerre and Gauss-Hermite techniques. While we argue that the effect of the CMSS is small, we include it for completeness and show calculations for different propagators with the coupling adjusted to give 49 MeV for two light quarks, appropriate to N-Δ splitting. Changing only the mass parameter for strange quarks ($m_s=300$ MeV) gives the matrix element as 36 MeV, similar to bag-model results.[8] To bracket the true problem, we use either light-light or strange-strange values for all pairs.

The results are shown in Table III. These show that

TABLE III. Downward energy shifts for various dibaryons in megaelectronvolts. For $S=-3$, the reference is the N-Ω^- threshold; for the H it is the Λ-Λ threshold. The calculations are performed in the potential model. MC means Monte Carlo integration; fit refers to the sum-of-Gaussians approximation; l-l refers to a pair of massless quarks (either up or down); and s-s to a pair of strange quarks. In each case, fits were made to the full set of octet and decuplet masses. This is the source of the extremely small variations in the values quoted for the N-Ω^-. Note that the physical Λ-Ξ^- threshold is only 174 MeV below the physical N-Ω^- threshold, so that the $J=1$ state is well bound.

Method	Propagator	Quark pair	N-Ω^-	$S=-3$, $J=2$	$J=1$	H
MC	Yukawa 0.4 GeV	l-l	244	295	370	281
MC	Yukawa 0.4 GeV	s-s	245	275	328	104
Fit	Gaussian $(0.2\text{ GeV})^2$	l-l	246	305	388	321
Fit	Gaussian $(0.2\text{ GeV})^2$	s-s	246	289	347	285
Fit	Gaussian $(0.5\text{ GeV})^2$	l-l	244	298	371	148
Fit	Gaussian $(0.5\text{ GeV})^2$	s-s	242	281	327	107

the N-Ω^- system is independent of the CMSS, while the full $J=2$, $S=-3$ state is very weakly dependent, and the $J=1$, $S=-3$ is only slightly more dependent. For comparison, we show the H calculated the same way, displaying the great sensitivity of that prediction. For this reason we believe that our prediction of $S=-3$ dibaryons is quite reliable.

When we examine the density profiles of the quark wave functions for these dibaryons, we find a significant value on the midplane between centers. In a strict sense, this violates the assumption which underlies a truncated two-centered potential,[3] calling the energy estimate into some doubt. Such a distribution naturally leads to consideration of the possibility of a distorted bag.[12] We leave that for future work, but here consider only the predictions in a simple spherical bag.[5-8]

If we ignore the color-magnetic spin splitting, the N-Ω^- system has binding energy $E_B=E_N+E_\Omega-E_{N,\Omega}$, where

$$E_N=\frac{3\omega_l-Z_0}{R_N}+\frac{4\pi B}{3}R_N^3, \quad (2)$$

$$E_\Omega=\frac{3\omega_s-Z_0}{R_\Omega}+\frac{4\pi B}{3}R_\Omega^3, \quad (3)$$

and

$$E_{N-\Omega}=\frac{3\omega_l+3\omega_s-Z_0}{R}+\frac{4\pi B}{3}R^3. \quad (4)$$

Using parameters from Ref. 7 (see Table III of that reference), $\omega_l=2.043$, $\omega_s=2.923$, $Z_0=1.84$, and $B^{1/4}=145$ MeV, we find

$$E_B=140 \text{ MeV}. \quad (5)$$

While this is less than the corresponding 250 MeV in the potential model, it is still very similar. Furthermore, we expect that including distortion will improve the agreement. Details may differ, but the color-magnetic interactions should have the same qualitative effects as those in Table III. Note that any corrections, such as center-of-mass motion, applicable to the potential model also apply to the bag model and should be similar in size.

If we argue that the binding energies are calculated more reliably than actual masses would be, then we can estimate the masses using physical thresholds. If we do that, the $J=2$, $S=-3$ dibaryon is, in all cases shown in Table III, stable against strong decay. This state presumably will radiatively decay to the $J=1$, $S=-3$ dibaryon, emitting a γ ray with energy between 50 and 80 MeV. The $J=1$ state can β decay to an H or to two Λ's. We leave an estimate of these rates and additional details to a more complete paper.[4]

However, we may make a crude estimate of the production rate in p-nucleus reactions from the known $p+A\to\Omega^-$ inclusive invariant differential cross sections which is about 500 nb/GeV^2 per nucleon at 200 GeV laboratory energy.[13] At such high energies, the dibaryon will most likely be formed in the fireball in analogy with the formation of antideuterons. Since the formation of a proton or a neutron should be the same as for the antineutron in the $\bar{d}$, the ratio of production cross sections should be about the observed ratio of Ω^- to $\bar{p}$ formation,[13] approximately 10^{-3}. This leads to an estimate on the order of 10 pb/GeV^2 per nucleon at 200 GeV, for dibaryons with momentum around 100 GeV/c. At lower dibaryon momenta, other production mechanisms, utilizing target nucleons, could be more important.

In summary, we have argued that, for the $S=-3$ sector, quark models unambiguously predict deeply bound dibaryons, stable under the strong interactions. We expect that the N-Ω^- physical channel should dominate the wave function and that because of this, production cross sections should be sufficient for experimental observation of these states.

This work has been supported in part by the U.S. Department of Energy and in part by the National Science Foundation. We are pleased to acknowledge valuable conversations with Dan Strottman and Jen-Chieh Peng.

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