"Inevitable" nonstrange dibaryon

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Certain basic features, common to all phenomenological models of hadron structure based on the picture of confinement at large distances and effective one-gluon exchange within the confinement region, necessarily lead to the prediction of the existence of a nonstrange dibaryon resonance with quantum numbers $IJ^P=03^+$, the d^*, independent of more detailed features of the dynamics of any of the models. We discuss the qualitative physics underlying this claim, comment on the probable mass and decay properties of the resulting state, and provide estimates of the expected production cross sections in $np \to d^*$ and $\pi^\pm d \to \pi^\pm d^*$.

I. INTRODUCTION

According to our current understanding, quantum chromodynamics (QCD) at low temperature and density is absolutely color confining. The requirement that all states have zero net color uniquely specifies the valence color structure of ordinary mesons and baryons; the same cannot be said for "multiquark" configurations (those with zero net color and ≥ 4 quarks and/or antiquarks) for which color singlet subclusterings exist. Whether or not a given multiquark configuration is resonant with respect to some ordinary multihadron subchannel depends on the nature of the underlying dynamics as applied to the physical breakup channels as well as the state in question. To date, despite considerable effort in identifying those channels most favorable to their existence, exotic resonances (as opposed to multiquark states which can be described as weakly bound states of ordinary hadrons, such as nuclei) have proven rather elusive.[1] In this paper we wish to point out that the most strongly favored nonstrange, positive parity dibaryon channel is that with $IJ^P=03^+$. The possible existence of such a state has been noted before, but we wish to stress that *any* QCD model having certain rather general features, which we discuss below, will necessarily predict a resonance in this channel. We then speculate on its likely mass, decay properties, and possible methods of production.

II. WHY THE $IJ^P=03^+$ CHANNEL?

In what follows we restrict ourselves to six-quark, orbitally nonexcited states. By the lowest-lying asymptotic two hadron (LLATH) state of a given channel we will mean the lowest-lying two-baryon state, in a relative s wave, having the given channel quantum numbers. The sum of the masses of the two baryons in the LLATH state then represents the threshold for fall-apart "decay" in the channel in question. A six-quark configuration lying below this threshold is necessarily resonant since it can be connected to a lower-lying two-baryon state (by definition of the threshold, in a relative d or higher wave) only by the action of some operator with nontrivial spin and orbital transformation properties. Moreover, since, as we see from the hadron spectrum, quark tensor forces are quite weak,[2] such states will have small tensor decay widths[3] to any such lower-lying two-baryon states and, in consequence, be very inelastic with respect to such channels.

Two features common to all phenomenological QCD-inspired models[4-9] of hadron structure are responsible for singling out the $IJ^P=03^+$ channel from among all available positive parity, nonstrange dibaryon channels: (1) the large internal kinetic energy which results from the confinement of (massless) quarks inside a baryon and (2) the specifics of the color-spin structure of the quark color magnetic (hyperfine) interaction.

The argument follows. In both bag[4-6] and potential (nonrelativistic[7,8] and relativistic[9]) models the effective confinement volume for a quark is larger in a six-quark than in a three-quark system. This delocalization alleviates the large internal kinetic energy and provides an effective binding relative to the LLATH state of the channel in question. In most channels, however, this effect is more than offset by the hyperfine interaction, the color-spin structure of which is

$$-\left(\tfrac{3}{4}\right)\sum_{i<j}\boldsymbol{\sigma}_i\cdot\boldsymbol{\sigma}_j\lambda_i^a\lambda_j^a \ . \qquad (1)$$

The structure of (1) is such as to favor color-spin symmetric pairs and hence, if we restrict ourselves to totally symmetric spatial configurations, lower flavor symmetries. The minimum value of the expectation of (1), as is well known, occurs in the H channel.[4]

Two points should be made with regard to this type of "counting." First, in channels favored by the hyperfine interaction, one, in fact, expects potentially significant admixtures of the spatial (42) symmetry,[10] since a lower spatial symmetry allows a higher color-spin symmetry. (The deuteron is presumably an extreme limiting case of this effect.) Second, since the hyperfine interaction is a short-distance effect, the increase in the effective confinement scale in the six-quark sector dilutes the strength of the spatial matrix element which accompanies the expectation of (1) relative to its value in an isolated baryon. If the LLATH state is flavor octet-flavor octet this represents a repulsive effect; for octet-decuplet the effect roughly cancels,[11] and for decuplet-decuplet it produces effective attraction. In general, however, the totally symmetric spatial configuration has an expectation of (1) different from that in the LLATH state; for most channels the overall hyperfine effect is repulsive.

Let us now examine the expectation of (1) in the spatially symmetric and LLATH configurations for the channels we are considering. These values are given in Table I, where $\langle(1)\rangle$ represents the expectation value of (1) and is to be multiplied by the corresponding spatial expectation in obtaining the total hyperfine energy.

One immediately sees that for all channels but $IJ^P=03^+$ the hyperfine expectation is far more repulsive in the symmetric configuration than in the corresponding LLATH channel. Furthermore, the dilution effect discussed above makes it clear that the $IJ^P=03^+$ state actually experiences an effective hyperfine attraction relative to the relevant two hadron ($\Delta\Delta$) threshold. Since the dilution factor typically lies between 1 and 2, it is also clear that this is the only channel for which this is the case. Moreover, we see that *any* model incorporating a hyperfine interaction with the one-gluon exchange color-spin structure (1) and a delocalization mechanism in its dynamics must necessarily produce a dibaryon resonance in this channel. For ease of reference in what follows, we will refer to the resulting state, whose quantum numbers are those of a spin excitation of the deuteron, as the d^*.

TABLE I. For the various possible spatially symmetric dibaryons, labeled by isospin and spin, we display the two hadron configuration of the LLATH (column 3), and the expectation of the discrete part of the color magnetic spin operator for the symmetric state (column 4) and for the separated hadrons (column 5).

I	J	LLATH	$\langle(1)\rangle_{\text{symmetric}}$	$\langle(1)\rangle_{\text{LLATH}}$
0	1	NN	2	−12
1	0	NN	6	−12
1	2	ΔN	12	0
2	1	ΔN	20	0
0	3	$\Delta\Delta$	12	12
3	0	$\Delta\Delta$	36	12

It is illuminating, at this point, to compare the qualitative features of the hyperfine interaction in the d^* and H channels.[12] In the latter the expectation of (1) is −18 for the H and −12 for the corresponding LLATH state, $\Lambda\Lambda$. Denoting by $\langle jj\rangle_{3q}$ and $\langle jj\rangle_{6q}$ the current-current spatial matrix elements accompanying the expectations of (1) in the (spatially symmetric) three- and six-quark systems, the change in hyperfine energy in going from $\Lambda\Lambda$ to the H is, therefore,

$$-18\langle jj\rangle_{6q}-2(-6)\langle jj\rangle_{3q} =12\langle jj\rangle_{3q}(1-1.5\langle jj\rangle_{6q}/\langle jj\rangle_{3q})\ . \quad (2)$$

If $\langle jj\rangle_{6q}/\langle jj\rangle_{3q}<\frac{2}{3}$ this is actually positive and so repulsive. The question of whether or not the hyperfine interaction is attractive (or, if repulsive, whether it overcomes the attraction due to the kinetic energy effect) in the H, relative to $\Lambda\Lambda$, is therefore one whose answer depends on the detailed dynamics of the model under consideration. In contrast, for the d^*, the hyperfine energy, relative to that in the corresponding LLATH state, $\Delta\Delta$, is

$$12\langle jj\rangle_{6q}-2(6)\langle jj\rangle_{3q} =-12\langle jj\rangle_{3q}[1-(\langle jj\rangle_{6q}/\langle jj\rangle_{3q})]\ . \quad (3)$$

Since $\langle jj\rangle_{6q}<\langle jj\rangle_{3q}$, this quantity is always negative and hence always favors resonance formation, independent of the detailed structure of the dynamics.

One further comparison bears mentioning, this from existing resonating group and resonating group-like calculations of the channels in question in the context of the nonrelativistic quark model (NRQM).[7,8] The NRQM provides a convenient framework within which to study the qualitative features of the quark-induced interactions between ordinary baryons which might be considered to form plausible entry channels for the formation of localized six-quark resonances. In the d^* channel one discovers that the induced residual $\Delta\Delta$ interaction is attractive,[7,8] while in the H channel, the residual $\Lambda\Lambda$ interaction is repulsive,[13] a fact which may have some bearing on production estimates in double-Λ hypernuclei.[14] Resonance formation by fusion of the corresponding LLATH state is thus unhindered for the d^* but hindered for the H.[15]

In Table II we list the binding energies for the d^*, relative to $\Delta\Delta$ threshold, obtained in a number of different models. The Massachusetts Institute of Technology (MIT) bag[4,5] and cloudy bag[6] and nonrelativistic quark[7,8] model calculations are described in the literature. We next briefly describe the Los Alamos potential model[9] (LAMP) calculation.

III. THE LAMP CALCULATION

The LAMP is a model[9] which represents the confining structure of the QCD vacuum together with the effect of additional quarks and/or antiquarks in a hadron by a linearly rising one-body Lorentz-scalar confining potential. Single quark wave functions are obtained by solving the Dirac equation for this potential. The underlying pic-

TABLE II. Binding energies of the d^* in various models described in the text. The last column indicates whether the calculation requires spatially symmetric states or not.

Model	Binding energy (MeV)	Symmetric spatial approximation
MIT bag (Ref. 5)	115	yes
cloudy bag (Ref. 6)	85	yes
NRQM (Ref. 8)	260	no
LAMP (Ref. 9)	350	yes

ture is baglike, but with a gradual restoration of the full nonperturbative vacuum surrounding the hadron in question over the region in which the quark density becomes small. Since the structure of QCD is believed to be such that gluonic configurations characterizing the physical vacuum are suppressed in the presence of quark density, this means that, for systems of low baryon density, the appropriate (mean-field) potential is obtained by truncation of the potentials from the individual baryon wells. In applying the model to the dibaryon sector, we will extend this ansatz to the regime of intermediate well separations where it is less well justified. Uncertainties associated with this procedure are discussed below. Note that the truncation ansatz also incorporates the notion that, on the confinement scale (~ 1 fm), QCD prefers to produce local color singlets. In the LAMP, therefore, the potential seen by a single quark in a dibaryon system along the line joining the two wells is W shaped, the outer arms of the W continuing toward infinity. Quark tunneling occurs through the barrier between the two wells and is dynamically determined. Note that, in using the Dirac equation and massless light quarks for our baryon, we include, except for retardation, all relativistic effects.

In the six-quark sector, rather than attempt to solve for the ground state of the truncated potential, we employ the following single-body trial wave functions

$$\psi_{i\epsilon}(x)=[\psi(x-x_i)+\epsilon\psi(x-x_j)]/N(\epsilon) ,$$

where x_i is the center of the ith potential well, $N(\epsilon)$ is a normalization factor, ψ is the $1S_{1/2}$ wave function in the single baryon well and, for $i=1,2$, $j=2,1$. Although we restrict ourselves here to the case $\epsilon=1$, such trial functions are in general a practical necessity since the repulsive character of the color hyperfine interaction, discussed below, in most channels tends to produce localized quark structure, the wave functions of which may be represented only by a superposition of many single-particle levels of the truncated two-centered potential.

In the limit $\epsilon=1$ the spatial d^* wave function is symmetric and the color-spin-isospin structure is uniquely determined. We restrict ourselves here to this limit, though it should be pointed out that the use of trial wave functions with $\epsilon\neq 1$ would allow us to construct states having the (42) spatial symmetry. Since ϵ and R ($=|\mathbf{x}_1-\mathbf{x}_2|$) are in general determined variationally, however, deviations of ϵ from 1 would serve only to lower the predicted d^* energy. We note in passing that this more general formulation of the model, in which ϵ as well as R is determined variationally, is free from the problem of artificial confinement since the LLATH state ($\epsilon=0$, $R=\infty$) is included in the allowed variational parameter space.

We thus consider a six-quark state of the form $(\psi_{\epsilon=1})^6$, the discrete couplings required to produce the d^* channel quantum numbers $IJ^P=03^+$, $C=0$ being suppressed. The required one-body energies are then obtained straightforwardly by numerical integration. To evaluate the color magnetic energy of the system one requires the quark color currents constructed from the $\epsilon=1$ trial functions and a local effective gluon propagator in the background, confining mean field. The latter, owing to confinement, has a restricted range; in order to facilitate calculation, we take the form to be Gaussian. The plausibility of such a form, together with some of its shortcomings, is discussed in Ref. 9. All model parameters are fixed by fitting to the N-Δ spectrum. Since we have restricted ourselves to $\epsilon=1$, the only dynamical parameter remaining is R, which is determined variationally. The resulting d^* energy in the model is quoted in Table II, and corresponds to $R\simeq 1.4$ fm.

As discussed above, the truncation ansatz for the mean-field potential is certainly not completely justifiable in the dibaryon regime, but is forced on us at present by the absence of a dynamical relationship between the mean field and the quark density in the model. A proper treatment would presumably somewhat raise the d^* energy since the nonperturbative vacuum between the well centers must, according to standard lore, be suppressed in order to accommodate the increased quark density.

We may make a crude estimate of this effect by using the MIT bag constant for the vacuum energy and imagining that we must create a cylinder of full perturbative vacuum of hadronic radius between the two potential centers.[16] The upper bound on the increase of the d^* mass so obtained is 80 MeV. We consider this to be a strong upper bound on the effect since the quark density on the midplane between wells, at the optimized well separation, is lower than its central value.[17] The resulting difference in binding energies between the LAMP and the bag results thus appears to indicate the efficiency of nonspherical distortions in lowering the hyperfine energies in this channel, without significantly increasing the kinetic energies. Note that both the bag and LAMP results suffer from the absence of configurations having (42) spatial symmetry. The inclusion of such configurations would further lower the d^* mass.

This (42) effect is included in the NRQM result. There, however, confinement is incorporated via effective two-body confining potentials whose connection to an underlying QCD-like picture is obscure, at best. The structure of confinement in the multiquark sector may, however, be more reasonable than one might at first expect. For example, the two-body confinement forces in the $QQ\bar{Q}\bar{Q}$ sector produce a lowest-lying adiabatic potential surface qualitatively similar to that of a string confinement picture.[18] Moreover, if one compares two-body confining potentials and the many-body confinement potential derived from the adiabatic approximation to the MIT bag,[19] one again finds the lowest lying adiabatic potential surfaces to be in striking agreement for the $QQ\bar{Q}\bar{Q}$ system, and to lie within the (albeit very crude) error bars of the quenched lattice calculation of Ref. 20. (One should also note that hidden color degrees of freedom are not admitted in the calculation of Ref. 7 in full generality, a situation which, however, appears remediable via Monte Carlo Green's function techniques.[21]) Thus, the NRQM may be a better guide than might be supposed.

It should be noted that, by construction, NRQM calculations remove all center-of-mass (c.m.) motion effects. The MIT and cloudy bag results can also claim to include c.m. corrections in the phenomenological fit, through the contribution of the $-Z/R$ term found in them. (One should be concerned, however, about how well this functional form includes c.m. corrections for nonspherical systems, or those with more than three quarks.) The LAMP has neither of these advantages and must, in principle, have binding energy results, such as those in Table II, corrected for c.m. effects. The manner in which this should be done is described in detail in Ref. 9. We only reiterate here that it is really the *difference* in c.m. effects (and others, such as rotational and vibrational effects) between the $\epsilon=1$ wave functions and noninteracting, $\epsilon=0$ wave functions at the same separation of well centers that affects the binding energy. For six light quarks in these two configurations, we find only an approximately 25-MeV correction to (reduction of) the binding energy. (Note that we cannot make c.m. corrections separately for the two noninteracting, separated sets of three quarks as this also removes rotational and vibrational effects from the "two noninteracting baryon," $\epsilon=0$ configuration that are *not* correspondingly removed by the c.m. correction procedure from the six-quark, $\epsilon=1$ configuration.) We have chosen not to include any c.m. correction in Table II as it is small and as there *are* rotational and vibrational collective-mode corrections of comparable size which should also be made, and which could even reverse the sign of the c.m. correction.

We conclude that, while explicit binding energies obtained must be treated with considerable caution, a binding energy of 200 MeV or more relative to $\Delta\Delta$ threshold represents a conservative estimate. This would put the resulting state below $N\Delta\pi$ threshold, leaving only NN, $NN\pi$, and $NN\pi\pi$ as available decay modes. The first is almost certainly[3] suppressed, owing to the weakness of quark tensor forces; the second as well, as is evident by considering the all spins up ($\uparrow^6$) substate of the d^* and any model for the pair creation mechanism. The $NN\pi\pi$ decay mode, if open with sufficient phase space, is, therefore, expected to dominate. As we will see below, however, for values of the d^* binding $\gtrsim 200$ MeV sufficient phase space is not available and the resulting 2π decay width is expected to be small.[22] The $NN\pi\pi$ coupling of the d^* will, nonetheless, represent the entry channel with greatest intrinsic strength for d^* production. This remains the case even if the d^* binding is greater than 300 MeV, in which case the $NN\pi\pi$ decay channel is closed. Note that, in the NRQM (Ref. 7), the NN one-gluon–exchange-mediated tensor decay width has been found to be ~ 5 MeV, in keeping with the qualitative arguments given above. Since the pionic decay widths are also expected to be relatively small (unless the d^* binding is considerably less than 200 MeV in which case the $NN\pi\pi$ width may become moderate) the total width should be of order a few 10's of MeV. We do not see any sufficiently reliable means of providing a more quantitative estimate.

IV. PRODUCTION OF THE d^*

As mentioned above, the largest intrinsic coupling of the d^* will certainly be to $NN\pi\pi$. We have argued, however, that the binding is likely to be 200 MeV or greater, relative to $\Delta\Delta$ threshold. As such, the $NN\pi\pi$ decay has less than 100 MeV of phase space and estimates presented below suggest that the resulting decay width is highly phase-space suppressed. As a result, for such d^* masses, the two-body np tensor decay mode may represent a reasonable fraction of the total decay width. One might then entertain the possibility of d^* production in np scattering. The $np\to d^*\to X$ on-resonance production cross section as a function of d^* mass is presented in Fig. 1. Note that the incoming and outgoing branching fraction have been scaled out. We may reasonably expect elastic branching fractions of order $\frac{1}{3}$ or so in this mass regime and widths of order of few 10's of MeV or less. However, even if the d^* were to lie below ~ 2200 MeV in mass, it would still have been difficult to see in np total cross section measurements at LAMPF (Ref. 23), since the cross sections in elastic np scattering are expected to be of order 1 mb. Nonetheless, detection of the d^* in np elastic scattering (after a partial-wave analysis) in the $^3D_3$ wave, appears feasible, there or at SATURNE. It should be pointed out that the above scenario can be somewhat altered if the QCD confinement and chiral symmetry breaking scales are significantly different, allowing the pion to penetrate the perturbative region, since, as discussed above, in that case effective one-pion-exchange quark-quark tensor forces may increase the total width, making detection more difficult. Note, however, that

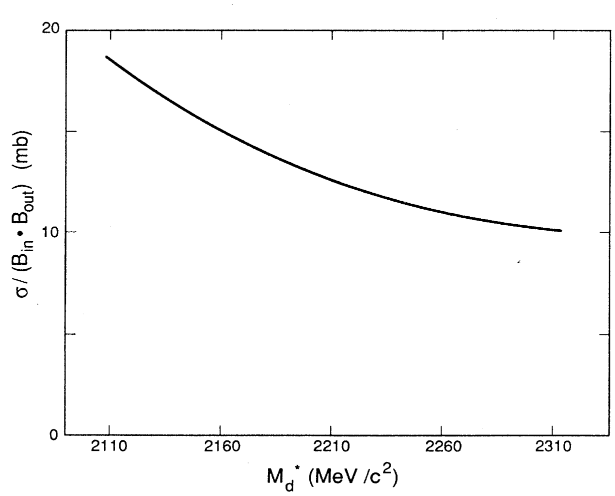

FIG. 1. Cross section for $np\to d^*\to X$ at resonance as a function of d^* mass, scaled by the product of branching fractions.

such pionic effects are not expected to produce large shifts in the d^* mass: In the cloudy bag model, the resonance is shifted[6] up 25 MeV, which is almost certainly an upper bound on the effect, given the pointlike nature of the effective πqq coupling in the model.

Owing to the expected large intrinsic $d^* \to NN\pi\pi$ width, it is of interest to consider the $\pi^{\pm} d \to \pi^{\pm} d^*$ reaction, since the phase-space suppression of the coupling in the d^* decay is not relevant. The $I=0$ character of the deuteron and d^* means that one can tag on the same sign outgoing pion. The background from quasifree Δ production may require eliminating all events with at least one soft (spectator) nucleon. One may make a rough estimate of the production cross section using a model with an effective πqq coupling whose strength is adjusted to generate the correct asymptotic NN one-pion-exchange-potential strength. The basic production process is then as shown in Fig. 2, where the fusion vertex, M_{fi}, depends on $\mathbf{q}_1, \mathbf{q}_2$ and the d, d^* wave functions. In making this estimate we take for M_{fi} the simple overlap between the four spectator and two scattered quarks of the deuteron and the six-quark d^* wave function. As such we ignore the (not inconsiderable!) soft QCD effects certainly present which may enhance the fusion vertex. The differential cross section is

$$\left[\frac{d\sigma}{d\Omega}\right] = q_1 q_2^3 \left[\frac{2^7 7}{5^5}\right] f_N^2 \frac{|M_{fi}|^2 (3+\cos^2\theta_{q_1 q_2})}{m_\pi^4 |(E - E_{in})|^2} , \tag{4}$$

where $f_N = 0.08$ and

$$M_{fi} = \frac{(\frac{2}{3})^{3/4} 2^6 6^{3/2} \beta^{15/2} \alpha^6}{\left[\sum_{k,j} \xi_k \xi_j \left[\frac{2}{\gamma_k^2 + \gamma_j^2}\right]^{3/2}\right]^{1/2} (\alpha^2+\beta^2)^6} \sum_k \frac{\xi_k}{(3\beta^2 + 2\gamma_k^2)^{3/2}} \exp\left[-\frac{(q_1^2+q_2^2)}{3(\alpha^2+\beta^2)} - \frac{(\mathbf{q}_1 - \mathbf{q}_2)^2}{4(3\beta^2+2\gamma_k^2)} \right] , \tag{5}$$

where ξ_k, γ_k $(k=1,2)$ are "fit" to make a two Gaussian (strength and width, respectively) description of the deuteron wave function, α is the width parameter for a single Gaussian description of the quark distributions in a nucleon, and β is a similar single Gaussian representation of the d^* quark wave function.

Because of the diffuse nature of the deuteron wave function, the fusion factors $|M_{fi}|^2$ entering the cross section are rather small, 10^{-5} to 10^{-4}. These may be enhanced by an order of magnitude by choosing a more compact nucleus like ^{3}He, but at the cost of more complicated detection problems in the final state. Results of this (conservative) estimate for a typical d^* mass, $m = 2200$ MeV, and a laboratory π momentum, $q_\pi = 700$ MeV/c, are shown in Fig. 3 as a function of the outgoing π scattering angle. The squared fusion vertex, $|M_{fi}|^2$, accounts for a factor of about 3.2 in the forward to backward peaking. The total cross section is 430 nb. These results are typical: Cross sections are of the order of a few hundred nb for π momenta of order 100 MeV/c above threshold.

As discussed above, the estimates just presented represent a rather conservative treatment of the fusion process. In order to investigate by how much these considerations of substructure might modify the estimates we can take a (considerably) overoptimistic view and ignore the d, d^* substructure completely. One may then consider an effective $dd^*\pi\pi$ vertex of the form

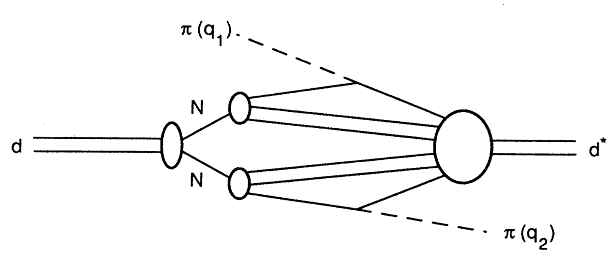

FIG. 2. Graph for $\pi d \to \pi d^*$ in a nonrelativistic quark model where the pions couple to the quarks from each nucleon.

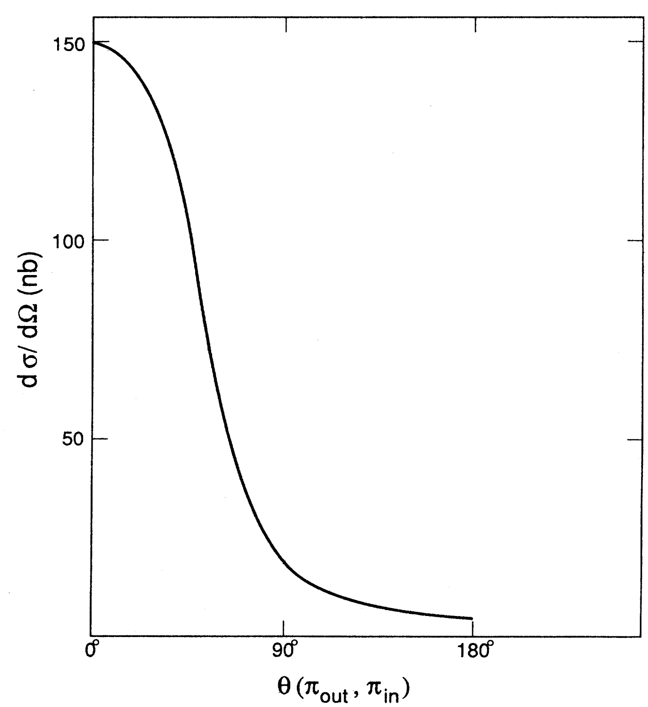

FIG. 3. Differential cross section for the process in Fig. 2 at $q_\pi = 700$ MeV/c.

$$\frac{g}{m_\pi^2} d^*_{\{\mu\nu\lambda\}} d^\lambda \partial^\mu \pi \partial^\nu \pi \ , \tag{6}$$

where $\{\mu\nu\lambda\}$ are symmetrized indices. This has the further advantage that, to the extent the low energy NN wave function in the $d^* \to NN\pi\pi$ decay can be approximated by the deuteron wave function, one can use the same effective vertex to estimate the two-pion decay width of the d^*, thus correlating the $\pi d \to \pi d^*$ cross section with the $d^* \to NN\pi\pi$ width.

Before beginning this calculation, we consider the question of the expected size of the dimensionless coupling constant g in (6), in the absence of direct data. A hadrodynamic model for this vertex is indicated in Fig. 4. From this, we see that

$$g \simeq g^2_{\pi N \Delta} \left[\frac{m_\pi}{M_N} \right]^2 f(d^* \to \Delta\Delta) f(d \to NN) \ . \tag{7}$$

Ignoring momentum dependence (beyond the factored out $\partial_\mu \pi$'s), we take $f(d \to NN) \sim 1$ and $f(d^* \to \Delta\Delta) \sim 2^{-3}$, where the latter vertex overlap represents only the difficulty of finding three quarks in a Δ-sized region and the three others in a similar, but modestly displaced region. We have in mind here the d^* quark wave functions of the LAMP which are spatially symmetric between two confining wells separated by ~ 1.4 fm. Using conventional values, we find

$$g \sim 1.5 \ , \tag{8}$$

although, as indicated above, momentum dependencies can be expected to reduce this value.

Applying the vertex (6) above to d^* decay, then, we find, after a straightforward calculation using nonrelativistic approximations for the pions and assuming that the d carries negligible energy in the final state

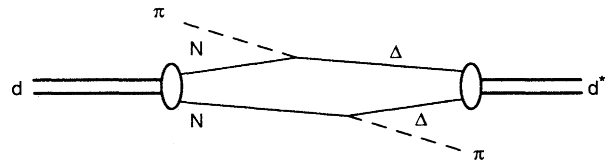

FIG. 4. Basic amplitude for the process $\pi d \to \pi d^*$ via a $\Delta\Delta$ intermediate state.

$$\Gamma(d^* \to d\pi\pi) = \frac{g^2(a + b/3)}{28\pi^3 M M^*} \delta^3 I(\delta) \ , \tag{9}$$

where $a = \frac{21}{45}$ and $b = \frac{7}{45}$ arise from polarization sums, $M^* = M + 2m_\pi + \delta$ is the d^* mass, M is the deuteron mass, and $I(\delta)$ is given by

$$I(\delta) = \int_0^1 dy \frac{[y(1-y)]^{3/2}}{\left[\left[y + \frac{m_\pi}{2\delta} \right] \left[1 + \frac{m_\pi}{2\delta} - y \right] \right]^{1/2}} \ . \tag{10}$$

For $\delta \sim 100$ MeV, $I(100 \text{ MeV}) \simeq 0.0623$ and then $\Gamma \sim 20$ eV using the value of g above. This confirms our contentions regarding phase space suppressions; in fact, the true result should be even smaller.

Now, similarly, using (6) to calculate the d^* production cross section in $\pi^\pm d \to \pi^\pm d^*$, we find ($s = W^2$ where W is the total center of mass energy)

$$\sigma = \frac{g^2(\frac{14}{27})}{768\pi m_\pi^4} \frac{[s-(M+m_\pi)^2]^{1/2}[s-(M-m_\pi)^2]^{1/2}[(s-M^{*2}+m_\pi^2)^2 - 4sm_\pi^2]^{3/2}}{s^3} \ . \tag{11}$$

For δ and g as above, and incident pion momentum of 580 MeV/c (which is then barely above threshold),

$$\sigma \sim 0.1 \ \mu\text{b} \ , \tag{12}$$

and the result is more than an order of magnitude larger at 700 MeV/c.

As stated earlier, this result is likely to be overly optimistic. As the f's are reduced to include the effects of momentum dependence, σ is easily reduced to the level of the previous estimate. The point to be noted, however, is that both estimates lie in the same ballpark, and suggest that experiments to search for this dibaryon are indeed feasible.

V. CONCLUSION

We have shown that the d^* is bound relative to Δ-Δ threshold in any QCD model that incorporates confinement and the hyperfine color-magnetic interaction. This is the only nonstrange dibaryon for which this statement can be made. The result is due to the fact that the six quarks occupy a larger volume in the d^* than they do in an individual baryon and that, for this channel ($IJ^P = 03^+$), such a delocalization also induces a net attraction from the hyperfine interaction. Our estimates of the production cross sections suggest that, over the range of expected binding energies for the d^*, the state should be experimentally observable in π-d and n-p scattering. We conclude that if a concerted experimental effort fails to confirm the existence of *this* dibaryon state, then there must be some serious flaw in the most basic features of all current models which purport to represent QCD in the strongly interacting regime, and that the determination of the existence or nonexistence of the d^* is therefore of considerable importance.

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[3]In making this statement we have tacitly assumed that the pion field is suppressed in the perturbative region. If the pions leak into the confinement volume, however, (corresponding to a scenario in which chiral symmetry breaking in QCD occurs at a significantly higher energy scale than does confinement) tensor forces mediated by pion exchange between quarks may be present and enhance tensor decay widths. [For an example of this approach in the context of R-matrix descriptions of dibaryon resonances, see P. Gonzales, P. La France, and E. L. Lomon, Phys. Rev. D **35**, 2142 (1987), and references therein.] Given the absence of evidence for strong tensor forces of the sort expected from effective *pointlike* pion exchange in the excited baryon spectrum, however, we do not expect such effects, even if present, to lead to large tensor couplings.

[4]R. L. Jaffe, Phys. Rev. Lett. **38**, 195 (1977); **38**, 617(E) (1977).

[5]A. Th. M. Aerts, P. J. G. Mulders, and J. J. deSwart, Phys. Rev. D **17**, 260 (1978).

[6]P. J. G. Mulders and A. W. Thomas, J. Phys. G **9**, 1159 (1983); K. Saito, Prog. Theor. Phys. **72**, 674 (1984).

[7]K. Maltman, Nucl. Phys. **A438**, 669 (1985).

[8]M. Oka and K. Yazaki, Phys. Lett. **90B**, 41 (1980).

[9]T. Goldman, K. Maltman, G. J. Stephenson, Jr., K. E. Schmidt, and Fan Wang, Phys. Rev. Lett. **59**, 627 (1987); T. Goldman, K. Maltman, G. J. Stephenson, Jr., and K. E. Schmidt, Nucl. Phys. **A481**, 621 (1988).

[10]Yu. F. Smirnov *et al.*, Yad. Fiz. **27**, 860 (1978) [Sov. J. Nucl. Phys. **27**, 456 (1978)]; I. T. Obukhovsky *et al.*, *ibid.* **31**, 516 (1980) [*ibid.* **31**, 269 (1980)].

[11]This, for example, is the reason why the $N\Omega$ system is expected to produce a resonance dominated by delocalization effects and is essentially insensitive to the structure of the hyperfine interaction. See first paper in Ref. 9.

[12]It should be pointed out that the d^* is the strangeness 0 member of a spin 3 flavor $\overline{10}$ multiplet whose isospin-strangeness content (I,S) is $(0,0)\oplus(\frac{1}{2},-1)\oplus(1,-2)\oplus(\frac{3}{2},-3)$. For all members of this multiplet the expectation of (1) in the spatially symmetric six-quark configuration is $+12$, the same as for the corresponding LLATH state. One thus expects resonances in all these channels; we concentrate here on the d^* channel as representing the channel most accessible to experimental investigation at the present time. Other $SU(3)_F$ multiplets with potentially attractive hyperfine expectations in the spatially symmetric configuration are the spin-zero flavor singlet H multiplet $(I,S)=(0,-2)$, and the spin-2 flavor octet-multiplet [containing the so-called "omegon" (Ref. 9)] and having isospin strangeness content $(I,S)=(\frac{1}{2},-1)\oplus(0,-2)\oplus(1,-2)\oplus(\frac{1}{2},-3)$. These do *not* contain nonstrange members.

[13]M. Oka, K. Shimizu, and K. Yazaki, Nucl. Phys. **A464**, 700 (1987).

[14]B. O. Kerbikov, Yad. Fiz. **39**, 816 (1984) [Sov. J. Nucl. Phys. **39**, 516 (1984)].

[15]The $N\Xi$ subchannel is, however, not so hindered (Ref. 13), and, in fact, has a larger discrete overlap with the H state.

[16]Here one must bear in mind that the short distance, attractive portion of the potential represents the perturbative regime.

[17]This is in contrast to the strangeness -3 state in this model, which is more tightly bound and more extensively delocalized.

[18]N. Isgur and J. Paton, Phys. Lett. **124B**, 247 (1983).

[19]K. Maltman (private communication).

[20]S. Ohta, M. Fukugita, and H. Ukawa, Phys. Lett. B **173**, 15 (1986).

[21]J. Carlson (private communication).

[22]This expectation may seem surprising if one were to view the d^* as a $\Delta\Delta$ bound state and to consider the effect of the intrinsic Δ width. However, although the coupling of this six-quark state to two Δ's is large, there is a factor of approximately 2^6 suppression in rate, due to the difference in the quark spatial wave function in the d^* vs in two separated Δ's and an additional suppression of the offshell Δ effective width, due to the cubic momentum dependence of that width. When both Δ's are about 100 MeV below their nominal mass shell value, the combination of only these two effects reduces the contribution of the Δ width (to the total effective width) to less than 1 MeV.

[23]R. R. Silbar (private communication). For the basic data, see P. Lisowski *et al.*, Phys. Rev. Lett. **49**, 255 (1982).