Multiquark states. III. Q^6 dibaryon resonances

P. J. Mulders, A. T. Aerts,* and J. J. de Swart
Institute for Theoretical Physics, University of Nijmegen, The Netherlands
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The mass spectrum for orbitally excited dibaryon resonances is predicted under the assumption of two clusters of quarks in a stretched bag. Decay mechanisms, stability, and experimental candidates in the $Y=2$, 1, and 0 channels are discussed. Natural explanations are found for, e.g., the $^3F_3$ and $^1D_2$ pp resonances, for the shoulder at 2.14 GeV, and the enhancement at 2.25 GeV in the Λp invariant-mass spectrum.

I. INTRODUCTION

Although during the last twenty years several candidates for dibaryon resonances have been found, they have not been considered very seriously until recently. One tried to explain nonstrange dinucleon resonances as strong $N\Delta$ or $\Delta\Delta$ interactions[1,2] or in three-body treatments as $NN\pi$ resonances.[3,4] In this paper, which is one of a series[5,6] in which we study multiquark states, most dibaryon resonances will be explained as six-quark states in one bag. The first reason for this study is that experimental evidence has accumulated for more and more of these dibaryon resonances.[7,8] Secondly, with the development of the MIT bag model[9] and its phenomenological extensions,[10,11] it has become possible to perform probably reliable calculations of the masses of multiquark states. As examples we mention here the calculations of the masses of the lowest $Q^2\overline{Q}^2$ mesons,[12,13] the $Q^4\overline{Q}$[13-16] and the Q^6 dibaryons.[17-19] For the orbital excitations there are the different calculations of the baryonium states,[6,20-22] the excited Q^3 baryons,[5,15,23] the orbital excitations of the $Q^4\overline{Q}$ baryons,[15,24,25] and the orbital excitations of the Q^6 dibaryons.[15,24,26,27]

We will denote experimentally observed dibaryon resonances as $B^2(Y,I,J^P;\text{mass})$ and the predicted six-quark states as $D(Y,I,J^P;\text{mass})$, where Y, I, J, and P are the hypercharge, isospin, spin, and parity of these resonances. In the Secs. IV, V, and VI where we discuss the $Y=2$, 1, and 0 channels, respectively, Y is omitted. The mass is quoted in GeV.

Except for the deuteron $B^2(2,0,1^+;1.875)$, which is a bound state in the $^3S_1+{}^3D_1$ NN wave, the longest-known dibaryon resonance is the $B^2(2,1,2^+;2.17)$ which first showed up as an enhancement at the $N\Delta$ threshold in the cross section of the photodisintegration of the deuteron[28]. Later some more NN resonances with masses above 2.6 GeV were found in the reactions $pp\to\pi^+X^+$ (Ref. 29), $pp\to\pi^+d$ (Refs. 30 and 31), $K^-d\to K^-\pi^+\pi^-d$ (Ref. 32), and $dp\to ppn$ (Ref. 33). Recent measurements[34] of the proton polarization in $\gamma d\to pn$ revealed a structure around 2.38 GeV that can be interpreted as a dibaryon resonance $B^2(2,0,3^+;2.38)$. Further recent evidence for dibaryon resonances comes from the Argonne experiments[35] with polarized beams and targets. These experiments indicate several resonances in the energy region below 2.5 GeV, the clearest one being $B^2(2,1,3^-;2.26)$, a $^3F_3$ NN resonance.[36,37] Also the already mentioned resonance $B^2(2,1,2^+;2.17)$ is seen in these experiments.[38] In a recent πd elastic-scattering experiment,[39] the resonances in the $^3F_3$ and probably also the $^1S_0$ wave have been seen. These resonances are naturally explained as six-quark states.

The experimental information for the strange dibaryon resonances mainly comes from invariant-mass plots for the different channels. The best established resonance is the $B^2(1,\frac{1}{2},1^+;2.129)$ which has been seen in many different experiments.[40-50] The spin and parity of this state comes from models[8] where this state is satisfactorily explained as the companion of the deuteron in the flavor-SU(3) irreducible representation (irrep) $\underline{10}^*$. A long-standing difficulty has been the shoulder[50] in the Λp invariant-mass spectrum around 2.14 GeV. In this paper, a quite natural explanation of this shoulder as a $J^P=1^-$ six-quark state will be given. In recent experiments[47] Shahbazian has clearly seen a Λp resonance at 2.256 GeV and found evidence for more strange dibaryon resonances.

If the explanation for the dibaryon resonances as six-quark states is correct, then it is quite easy to understand why the lowest *strange* dibaryon resonances are much closer to the corresponding two-body thresholds than the lowest dinucleon resonances. The color-magnetic interactions split dibaryons which have different flavor and spin structure. The more antisymmetric flavor irreps [with lower values for the flavor-SU(3) quadratic Casimir operator] have a lower energy. These flavor irreps, e.g., $\underline{f}=\underline{8}$ or $\underline{f}=\underline{1}$, only contain members with at least one or two strange quarks;

they do not contain members with only nonstrange quarks. As a consequence, we expect the study of these dibaryon resonances to be easier in ΛN final-state interactions, despite the fact that the statistics and variety in pp scattering experiments is of course much larger.

We would like to urge experimentalists to plan high-statistics experiments looking for these strange dibaryons at sufficiently high energy.

II. MASSES OF THE SIX-QUARK STATES

In order to calculate the masses of the Q^6 states we use the mass formula given in Ref. 5. Because only a few candidates exist we cannot include the phenomenological contributions to the mass discussed in Ref. 5. Moreover, we think they are less important here. The mass formula then becomes

$$M = M_l + M_m . \qquad (1)$$

When we neglect the contribution M_m, due to the color-magnetic interaction, the states lie on linear trajectories, e.g., for the leading trajectory

$$M_l^2 = M_0^2 + (1/\alpha')l . \qquad (2)$$

We assume that the intercept M_0 of the leading trajectory can be calculated in the spherical-cavity approximation of the MIT bag model.[10] Thus

$$M_0 = \frac{4\pi}{3} BR^3 - \frac{Z_0}{R} + \sum_i N_i \frac{\alpha_i(R)}{R} , \qquad (3)$$

which contains a volume term ($B^{1/4} = 0.146$ GeV), a contribution from the zero-point energy and the center of mass motion ($Z_0 = 1.84$), and the quark energies. The sum runs over all flavors in the bag, and N_i counts the quarks with a specific flavor i (n for nonstrange, s for strange). Our prescription[5] for the bag radius of the N-quark ground state gives

$$R = r_0 N^{1/3} \qquad (4)$$

with $r_0 = 3.63\ \text{GeV}^{-1}$. The values of the functions $\alpha_i(R)$, which measure the energy of a quark in a spherical bag with radius R can be found in Ref. 10. For a nonstrange, massless quark $\alpha_n = 2.043$, independent of the radius.

When orbitally excited, hadrons consist by assumption of two clusters of quarks at the ends of a stretched and fast-rotating stringlike bag. The quarks inside each cluster couple to a nonzero color charge, such that the whole hadron is colorless. By $(q^N)_c$ we indicate an N-quark cluster coupling to the color-SU(3) irrep $\underline{c}$, which is identified by its dimension c. For dibaryons we have the following possibilities:

$$(Q^5)_{3*} - (Q)_3; \quad (Q^4)_3 - (Q^2)_{3*}; \quad (Q^4)_{6*} - (Q^2)_6; \quad (Q^3)_8 - (Q^3)_8 .$$

For high l, the energy and the orbital angular momentum of the rotating bag mainly come from the color fields in the bag. They depend on the color charge of the quark clusters at the bag ends. This shows up in the slope of the trajectories,

$$(1/\alpha') = (1.1\ \text{GeV}^2)(\tfrac{3}{4} f_c^2)^{1/2} , \qquad (5)$$

where f_c^2 is the eigenvalue of the quadratic Casimir operator F_c^2 in the color-SU(3) irrep $\underline{c}$ to which the quarks in a cluster couple. For $N = 6$, the intercept M_0 and the multiplet masses M_l are given in Table I for several l values.

The color dependence enters the multiplet masses M_l through the slope of the trajectories. The intercepts are assumed to be independent of the color structure of the trajectories. The flavor dependence of the masses M_l is completely determined by M_0 i.e., by the quark energies in Eq. (3).

An alternate way to calculate M_l is the semiclassical approach to the relativistic string, in which a relation $M_l = f(l)$ between the mass M and the orbital angular momentum l is established classically. To account for quantum-mechanical effects, a quantum defect l_0 is introduced and M_l is assumed to be given by

TABLE I. The masses of orbitally excited six-quark states in MeV, neglecting the color-magnetic interactions. The irreps ($\underline{c}$-$\underline{c}$*) indicate the color structure.

$M_l(\underline{3}$-$\underline{3}^*)$

Y \ l	0	1	2	3	4	5
2	2125	2370	2591	2796	2986	3165
1	2295	2523	2733	2927	3109	3281
0	2464	2678	2876	3061	3236	3402
−1	2634	2835	3023	3200	3367	3527
−2	2803	2993	3171	3340	3501	3655
−3	2973	3153	3322	3153	3322	3484
−4	3142	3312	3475	3629	3778	3921

$M_l(\underline{6}$-$\underline{6}^*)$

Y \ l	0	1	2	3	4
2	2125	2501	2827	3120	3387
1	2295	2647	2957	3238	3496
0	2464	2795	3090	3360	3609

$M_l(\underline{8}$-$\underline{8})$

Y \ l	0	1	2	3	4
2	2125	2483	2796	3077	3334
1	2295	2630	2927	3196	3445
0	2464	2779	3061	3320	3560

$$M_l = f(l + l_0)\,.$$

For mesons[51] and baryons, massless nonstrange quarks and light ($m_s \simeq 0.29$ GeV) strange quarks can account for the flavor dependence of the intercepts and the orbitally excited states, when we assume a flavor-independent quantum defect. The same assumption for dibaryons enables us to determine l_0 from the nonstrange intercept mass $M_0 = 2.125$ GeV. We find $l_0 \simeq 4.1$ and an intercept mass $M_0 = 2.238$ for $Y=1$ states. This is about 57 MeV lower than the result in Table I. For $Y=0$ states, two different intercepts appear: $M_0 = 2.349$ GeV when the strange quarks are at different bag ends and $M_0 = 2.423$ GeV when the strange quarks are at the same bag end.

The color-magnetic interaction energy is given by

$$M_m = m_1\Delta_1 + m_2\Delta_2\,, \tag{6}$$

where m_1 and m_2 measure the strength of the color-magnetic interaction in the clusters. This strength m depends [after elimination of R via Eq. (4)] on the number and flavor of the quarks (see Ref. 5 and Table II) and can be approximated by

$$m = aN^{-1/3} - bN_s N^{-1}\,,$$

with $a = 107$ MeV and $b = 28$ MeV. The factors Δ_1 and Δ_2 give the dependence on the spin and color-spin of the quarks. For an N-quark cluster (no antiquarks)

$$\Delta = -\tfrac{1}{4}N(10-N) + \tfrac{1}{3}S^2 + F_f^{\,2} + \tfrac{1}{2}F_c^{\,2}\,, \tag{7}$$

where $F_f^{\,2}$ is the quadratic Casimir operator for the flavor SU(3) and S is the spin of the cluster. For the relevant clusters the color, flavor, and spin content and the eigenvalues of Δ have been summarized in Table III.

The aforementioned prescription enables us to compute the masses of the six-quark states (only one cluster present) and their orbital excitations (two clusters present). Several effects which can influence the actual value of the masses have been omitted. First of all, we have neglected the influence of decay channels and final-state interactions. Secondly, we did not include splittings arising from the spin-orbit and tensor forces, which presumably are rather small.[5] Apart from these contributions, small mass shifts may arise from mixing but they are neglected. We mention the color-magnetic interactions between the two clusters and the exchange contributions, due to the overlap of the quark wave functions. They cause mixing between $c=3$ and $c=6$ Q^4-Q^2 states and also between $c=8$ and $c=1$ Q^3-Q^3 configurations (see also Sec. III). We also want to mention the effect of tunneling of a quark from one end of the bag to the other end. For baryons this causes an exchange-

TABLE II. The strength of the color-magnetic interaction $m(N, N_s)$ in MeV for an N-quark cluster with N_s strange quarks.

N \ N_s	0	1	2	3	4	5	6
2	85.1	70.2	58.2				
3	74.4	64.3	55.5	47.9			
4	67.6	60.1	53.3	47.0	41.2		
5	62.7	56.8	51.2	46.0	41.3	36.9	
6	59.0	54.1	49.5	45.1	40.9	37.0	33.4

TABLE III. N-quark clusters. c = color, f = flavor, s = spin.

$(Q^N)_c$	f	s	Δ	$(Q^N)_c$	f	s	Δ
$(Q)_3$	$\underline{3}$	$\frac{1}{2}$	0	$(Q^4)_{6^*}$	$\underline{3}$	1	$-\frac{7}{3}$
					$\underline{6}^*$	0	-1
$(Q^2)_{3^*}$	$\underline{3}^*$	0	-2		$\underline{6}^*$	2	1
	$\underline{6}$	1	$\frac{2}{3}$		$\underline{15}$	1	$\frac{5}{3}$
					$\underline{15}_s$	0	5
$(Q^2)_6$	$\underline{3}^*$	1	$-\frac{1}{3}$				
	$\underline{6}$	0	1				
$(Q^3)_1$	$\underline{8}$	$\frac{1}{2}$	-2	$(Q^5)_{3^*}$	$\underline{3}^*$	$\frac{1}{2}$	-4
	$\underline{10}$	$\frac{3}{2}$	2		$\underline{3}^*$	$\frac{3}{2}$	-3
					$\underline{6}$	$\frac{1}{2}$	-2
$(Q^3)_8$	$\underline{1}$	$\frac{1}{2}$	$-\frac{7}{2}$		$\underline{6}$	$\frac{3}{2}$	-1
	$\underline{8}$	$\frac{1}{2}$	$-\frac{1}{2}$		$\underline{15}^*$	$\frac{1}{2}$	0
	$\underline{8}$	$\frac{3}{2}$	$\frac{1}{2}$		$\underline{15}^*$	$\frac{3}{2}$	1
	$\underline{10}$	$\frac{1}{2}$	$\frac{5}{2}$		$\underline{15}^*$	$\frac{5}{2}$	$\frac{8}{3}$
					$\underline{24}$	$\frac{1}{2}$	3
					$\underline{24}$	$\frac{3}{2}$	4
					$\underline{21}$	$\frac{1}{2}$	8
$(Q^4)_3$	$\underline{3}$	0	-4	$(Q^6)_1$	$\underline{1}$	0	-6
	$\underline{3}$	1	$-\frac{10}{3}$		$\underline{8}$	1	$-\frac{7}{3}$
	$\underline{6}^*$	1	$-\frac{4}{3}$		$\underline{8}$	2	-1
	$\underline{15}$	0	0		$\underline{10}$	1	$\frac{2}{3}$
	$\underline{15}$	1	$\frac{2}{3}$		$\underline{10}^*$	1	$\frac{2}{3}$
	$\underline{15}$	2	2		$\underline{27}$	0	2
	$\underline{15}_s$	1	$\frac{14}{3}$		$\underline{27}$	2	4
					$\underline{10}^*$	3	4
					$\underline{35}$	1	$\frac{20}{3}$
					$\underline{28}$	0	12

like contribution to the mass, as the tunneling of a quark is equivalent to the exchange of the other two quarks. The Q^2-Q structure is preserved. For dibaryons, tunneling causes transitions between the various structures

$$Q^5\text{-}Q \longleftrightarrow Q^4\text{-}Q^2 \longleftrightarrow Q^3\text{-}Q^3 .$$

In Ref. 1, the tunneling effects proved to be at most 100 MeV for $l=1$ and to disappear rapidly for higher l values. The mixing due to tunneling is strongest between those dibaryons which have approximately the same masses, e.g., for $l=1$ between the $(Q^5)_{3*}$-$(Q)_3$ and $(Q^4)_3$-$(Q^2)_{3*}$, or between the $(Q^4)_{6*}$-$(Q^2)_6$ and $(Q^3)_8$-$(Q^3)_8$ states. Tunneling also may lead to the decay mode $(Q^4)_3$-$(Q^2)_{3*} \to (Q^3)_1$-$(Q^3)_1$, which will be considered in Sec. III.

In order to illustrate the mass formula we consider two examples. We calculate the spectrum of the nonstrange $(Q^6)_1$ and $(Q^4)_3$-$(Q^2)_{3*}$ dibaryons. For the $(Q^6)_1$ dibaryons, the color-magnetic contribution, calculated using Tables II and III, is added to the multiplet mass M_0 which is found in Table I. This is illustrated for dinucleons in Fig. 1. For the $(Q^4)_3$-$(Q^2)_{3*}$ dibaryons, the color-magnetic interactions contain two contributions, one from the Q^2 system and the other from the Q^4 system. In the final spectrum this leads to degenerate levels, as illustrated for dinucleons in Fig. 2. In this figure, the masses for the $l=1$ levels have been given. The other orbitally excited $(Q^4)_3$-$(Q^2)_{3*}$

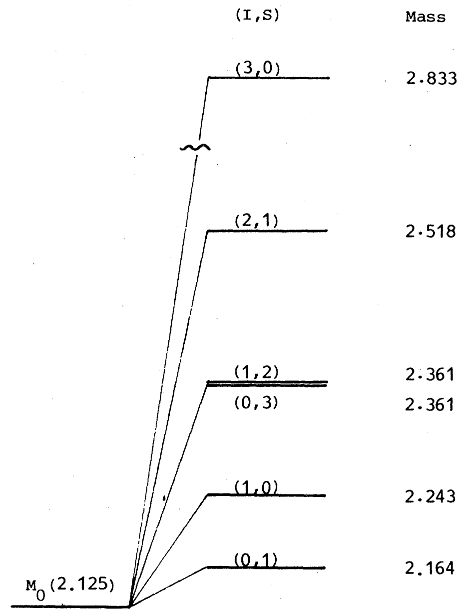

FIG. 1. The color-magnetic splitting for the $l=0$ Q^6 dibaryons. All masses are in GeV.

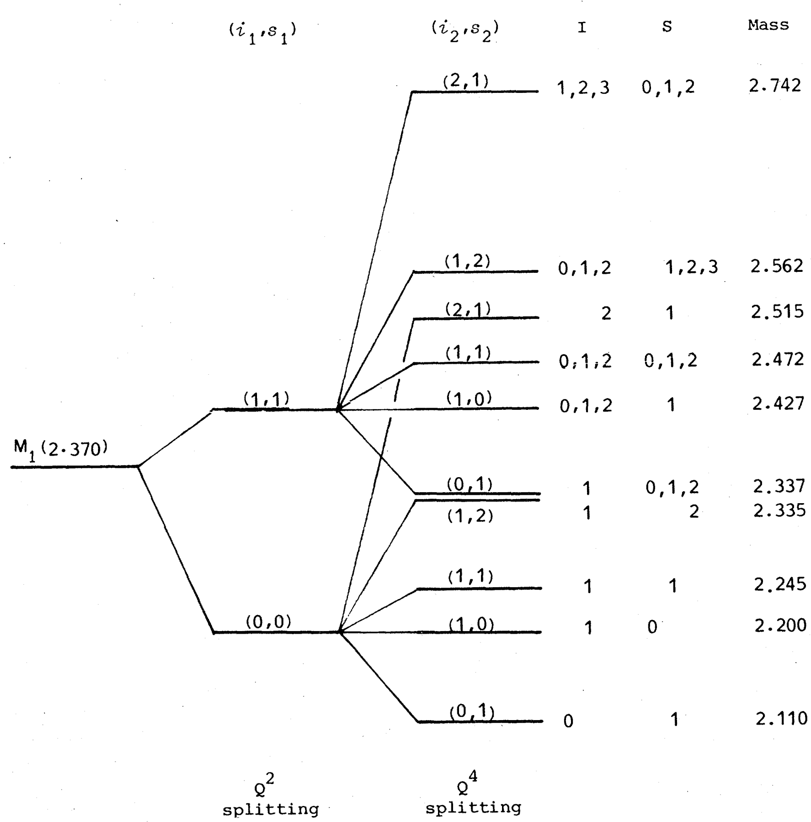

FIG. 2. The color-magnetic splitting for the $l=1$ color-triplet Q^4-Q^2 dibaryons. All masses are in GeV.

levels show the same color-magnetic splitting, only the multiplet mass is different: e.g., for $l=2$, $M_2=2.591$ GeV.

III. STABILITY AND DECAY OF DIBARYONS

One of the main decay modes of the s-wave Q^6 states is fission. If it is energetically favorable, a Q^6 state will decay into two colorless baryons. The change in energy, neglecting the color-magnetic interaction

$$\delta M_0 = M_0(Q^6) - 2M_0(Q^3),$$

is not very large ($\delta M_0 \simeq -50$ MeV). To determine whether fission into two colorless parts is energetically favorable, one has to look at the change in the color-magnetic interaction energy

$$\delta M_m = M_m(Q^6) - 2M_m(Q^3).$$

We will approximate this by

$$\delta M \simeq m\Delta_{12}(Q^6),$$

with

$$\Delta_{12}(Q^6) = \Delta(Q^6) - \Delta_1(Q^3) - \Delta_2(Q^3).$$

Δ_{12} is a measure for the color-magnetic attraction in the Q^6 cluster between the two $(Q^3)_1$ subsets. The stability increases when Δ_{12} decreases. We will assume that the color-magnetic interaction energy also determines the stability of clusters other than $(Q^6)_1$. For example, a $(Q^4)_3$ cluster can fission into a colorless baryon and a quark

$$(Q^4)_3 \to (Q^3)_1(Q)_3.$$

The stability of this cluster is then measured by

$$\Delta_{12}(Q^4) = \Delta(Q^4) - \Delta(Q^3).$$

In the following, we will discuss briefly the various possibilities for the decay of dibaryon resonances.

The first possibility is the fission of an s-wave Q^6 dibaryon into two baryons (BB). Because of parity conservation, the final-state baryons are in an even L wave. The decay in S waves is expected to have a very large width. The decay in D or higher waves is suppressed due to the angular momentum barrier and due to the spin flip of the quarks that is required in order to conserve the total angular momentum. The vectorial change of the angular momenta, such that $\Delta J^P = 0^+$, is given by $\Delta L = \Delta S = 0$ for the decay in S waves and by $\Delta L = \Delta S = 2$ for the decay in D waves.

Fission is also possible for orbitally excited dibaryons but it will not be important because other decay modes will dominate. The reason is that the orbitally excited baryon (Q-Q^2) which is formed in the decay

$$((Q^4)_3\text{-}(Q^2)_{3*}) \to ((Q^3)_1(Q_3\text{-}(Q^2)_3) \to (Q^3)_1(Q\text{-}Q^2)_1$$

is usually quite heavy.

Very important for the coupling of excited dibaryons to the BB channels is the tunneling mode. In excited multiquark states, the quarks reside at the ends of the rotating stringlike bag. Nevertheless, it is possible that a quark tunnels from one end of the bag through the angular momentum barrier to the other end. This gives the recoupling

$$((Q^4)_3 \overset{l}{—} (Q^2)_{3*}) \to ((Q_3)_1 \overset{L}{—} (Q^3)_1).$$

This tunneling will be easy when the tunneling quark is not "bound" to the Q^4 cluster $[\Delta_{12}(Q^4) > 0]$ and when $L = l$. This coupling to the BB channel will not be so strong when $l \neq L$, because then the process has to be accompanied by a spin flip in order to conserve the total angular momentum.

The final decay mode which we want to discuss proceeds via quark-antiquark creation. The s-wave and excited color-triplet dibaryons will decay via $Q\bar{Q}$ creation into BBM channels,

$$Q^6 \to Q^7\bar{Q} \to (Q^3)(Q^3)(Q\bar{Q}),$$
$$(Q^4\text{-}Q^2) \to (Q^4(\bar{Q}Q)Q^2) \to (Q^3)(Q\bar{Q})(Q^3),$$
$$(Q^5\text{-}Q) \to (Q^5(Q\bar{Q})Q) \to (Q^3)(Q^3)(\bar{Q}Q).$$

In order to conserve angular momentum and parity, $\Delta J^P = 0^+$, the $Q\bar{Q}$ pair is created in a $^3P_0$ wave; i.e., $\Delta L = \Delta S = 1$ for this decay. The total orbital angular momentum in the final state is $L = l \pm 1$.

Color-sextet and color-octet orbital excitations with $l = 1$ can also decay easily via $Q\bar{Q}$ creation. The reason is that the pair creation can take away one unit of orbital angular momentum. This can then leave all quarks in relative s waves and the color is easily annihilated by recoupling, e.g.,

$$((Q^4)_{6*}\text{-}(Q^2)_6) \to ((Q^4)_{6*}(Q)_3(\bar{Q})_{3*}(Q^2)_6) \to (Q^3)_1(Q^3)_1(Q\bar{Q})_1 .$$

IV. NONSTRANGE ($Y = 2$) DIBARYON RESONANCES

In this section we discuss the nonstrange dibaryon resonances. The predicted mass spectrum will be compared with the experimentally known resonances, which have been listed in Table IV. The lowest experimentally observed resonances are good candidates for six-quark states. They are $B^2(1, 2^+; 2.17)$, $B^2(1, 3^-; 2.2\text{–}2.3)$, $B^2(0, 3^+; 2.38)$, and $B^2(1, 0^+ \text{ or } 4^+; 2.4\text{–}2.5)$, although the (I, J^P) assignments of the higher two are less definite.

The predicted mass spectrum for $(Q^6)_1$, $(Q^5)_{3*}$-$(Q)_3$, and $(Q^4)_3$-$(Q^2)_{3*}$ nonstrange dibaryons for $I = 0$ and $I = 1$ is shown in Figs. 3 and 4. The flavor

TABLE IV. Candidates for nonstrange ($Y=2$) dibaryon resonances. PWA = partial-wave analysis, DA = dispersion analysis, LCM = Legendre-coefficient method, P = polarization.

Mass (GeV)	Width (MeV) elasticity	I	J^P (NN wave)	Remarks	Refs.
1.875		0	1^+ ($^3S_1+{}^3D_1$)	deuteron	
2.17	$\Gamma\simeq$ 35–100	1	2^+	$\gamma d \to pn$	28
	$x\simeq 0.1$		($^1D_2$)	$K^-d\to K^-\pi^-\pi^+d$	32
				$dp \to ppn$	33
				$pp \to pp$ (PWA)	38
2.2–2.3	$\Gamma\simeq$ 100–300	1	3^-	$pp \to pp$ (PWA, DA, P, LCM)	35–37
	$x\simeq 0.2$		($^3F_3$)	$\pi d \to \pi d$	39
2.38		0?	3^+? ($^3D_3+{}^3G_3$)	$\gamma d \to pn$ (P)	34
2.4–2.5	$\Gamma\simeq$ 100–200	1	0^+ or 4^+?	$pp \to pp$ (PWA, DA)	36, 37
			($^1S_0$ or $^1G_4$)	$\pi d \to \pi d$	39
~2.6				$pp \to \pi^+X^+$	29
~2.9				$pp \to \pi^+X^+$	29
				$pp \to \pi^+d$	30, 31
~3.6				$pp \to \pi^+d$	31
				$pp \to \pi^+X^+$	29
~3.9				$pp \to \pi^+X^+$	29

and spin structure of the $Y=2$ dibaryons is given in Table V. This structure determines the magnitude of the color-magnetic energy M_m. It also determines to which baryon-baryon (BB) channels the dibaryon can couple. For the $(Q^6)_1$ states, which fission into two colorless baryons in S waves, and for the $(Q^4)_3$-$(Q^2)_{3*}$ states, which decay via quark tunneling into two colorless baryons in L waves with $L=l$, the BB channels have been given in Table V. The flavor and spin structure is independent of the orbital angular momentum l of the quark bag. The mass of every dibaryon resonance is found by adding the color-magnetic interaction energy M_m to the multiplet mass M_l (Table I). For the $(Q^3)_8$-$(Q^3)_8$ dibaryons we have to take into account the fact that we are combining two identical fermion systems.

The predicted masses are listed with greater accuracy than warranted by the model in order to distinguish between the resonances. We start the discussion with the s-wave Q^6 states, which can fission into baryon-baryon channels in even-L

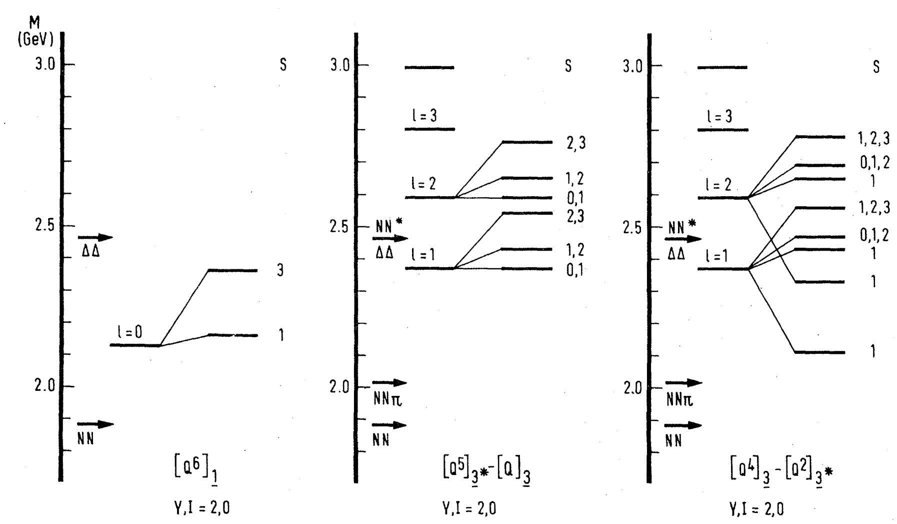

FIG. 3. The predicted mass spectrum for S wave and orbitally excited $(Y,I)=(2,0)$ color-triplet dibaryon resonances. The spin s and l must be added to find the total angular momentum J.

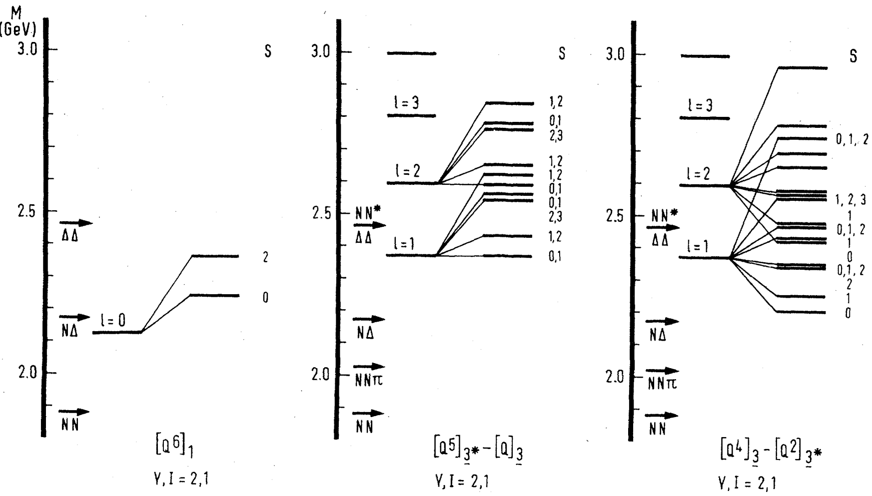

FIG. 4. The predicted mass spectrum for S wave and orbitally excited $(Y,I)=(2,1)$ color-triplet dibaryon resonances. The spin s and l must be added to find the total angular momentum J.

waves.

$D(0,1^+;2.16)$ and $D(1,0^+;2.24)$ are the lowest predicted $(Q^6)_1$ nonstrange dibaryons. They fission into S waves, the $^3S_1$ and $^1S_0$ NN waves, respectively. These dibaryons are very unstable. The change in color-magnetic energy, measured by Δ_{12}, is very large; $\Delta_{12}=\frac{14}{3}$ and 6, respectively. Such states probably do not show up as pronounced resonances in a BB channel; rather they only give a background contribution. This can be compared with the $\epsilon(760)\to\pi\pi$, which can be interpreted as the fission of a $Q^2\overline{Q}^2$ bag into two mesons.[12,52]

$D(0,3^+;2.36)$ cannot fission in S waves. It couples to the $^7S_3$ $\Delta\Delta$ wave, but its mass is below the $\Delta\Delta$ threshold at 2.47 GeV. It can, however, fission into BB channels in an (even) L wave with $L\neq 0$, accompanied by a spin flip, or it can decay into BBM channels through quark-pair creation. Fission into a $^3D_3$ NN wave is possible. As the coupling to NN is suppressed due to the angular-momentum barrier and the necessary spin flip, a reasonable width may emerge. We think that this state is responsible for the experimental resonance structure found around 2.38 GeV in deuteron photodisintegration.[34] Therefore we prefer to make the identification $B^2(2.38)\equiv D(2,0,3^+;2.36)$.

$D(1,2^+;2.36)$ is an intermediate case. It couples to the $^1D_2$ NN wave through fission in a D wave, but it can also couple to the $^5S_2$ $N\Delta$ wave through fission in S waves. While for the latter decay possibility, the coupling is larger (S-wave vs D-wave fission), for the former more phase space is available. Experimentally an $(I,J^P)=(1,2^+)$ state shows up in the $^1D_2$ NN wave and in $NN\pi$ just at the $N\Delta$ threshold (Table IV). We therefore prefer to make the identification $B^2(2,1,2^+;2.17)\equiv D(2,1,2^+;2.36)$. This situation can be understood in a potential model[53] where the bag is coupled to a $^1D_2$ NN and a $^5S_2$ $N\Delta$ channel. In this model, the pole positions are followed in the complex energy plane. When the bag is weakly coupled to the NN and $N\Delta$ channels, the $D(1,2^+;2.36)$ is represented by one set of conjugate poles on the third Riemann sheet ($\mathrm{Im}k_{NN}<0$, $\mathrm{Im}k_{N\Delta}<0$), connected with the physical sheet above the $N\Delta$ threshold and another set of conjugate poles on the fourth sheet ($\mathrm{Im}k_{NN}>0$, $\mathrm{Im}k_{N\Delta}<0$), connected with the second sheet above the $N\Delta$ threshold. The second sheet ($\mathrm{Im}k_{NN}<0$, $\mathrm{Im}k_{N\Delta}>0$) is reached from the physical sheet by passing the unitarity cut between the NN and $N\Delta$ thresholds, the poles all lie near $E\simeq 2.36$ GeV. When the coupling of the bag and the NN and $N\Delta$ channels is increased, the poles move. The poles on the third sheet move away from the unitarity cut and the resonance structure around $E=2.36$ GeV weakens. Increasing the coupling strength, the poles on the fourth sheet move towards the $N\Delta$ thresholds. Still above the threshold, they cross the unitarity cut into the second sheet and finally show up as poles in the second sheet quite close to the $N\Delta$ threshold. This means a NN resonance near the $N\Delta$ threshold. An analogous case is the coupled $\pi\pi$ and $K\overline{K}$ system. Here the $S^*(0.98$ GeV) is predicted in the bag model as a $Q^2\overline{Q}^2$ state at 1.15 GeV and shows up as a $\pi\pi$ resonance near the $K\overline{K}$ threshold.[52,54]

The higher $(Q^6)_1$ dibaryons $D(2,1^+;2.52)$ and $D(3,0^+;2.83)$ do not couple to NN, but only to $N\Delta$ and $\Delta\Delta$.

Other dibaryons which we expect to couple strongly to BB channels are the $(Q^4)_3$-$(Q^2)_{3^*}$ states. Through tunneling they decay into BB channels. The stability depends on the color-magnetic energy gained when the Q^4 cluster fissions into a baryon and a quark. This is measured by $\Delta_{12}(Q^4)$. Of all nonstrange Q^4 clusters, the cluster with $(\underline{f},s;\Delta)=(\underline{15},2;2)$ has the smallest value Δ_{12}, namely $\Delta_{12}=0$. The cluster with $(\underline{f},s;\Delta)=(6,1;-\frac{4}{3})$ has $\Delta_{12}=\frac{2}{3}$; the others have $\Delta_{12}\geq 2$.

The clusters with $\Delta_{12}>0$ are unstable. They easily fission and the quark recombines after tunneling with the other diquark to a baryon or the

TABLE V. $Y=2$ dibaryon resonances. All masses are quoted in MeV.

$(Q^6)_1$

$f(y,i)s$	$NN(1.88)$	$N\Delta(2.17)$	$\Delta\Delta(2.46)$	M_m	M_0
$\underline{10}^*(2,0)1$	×		×	39	2164
$\underline{27}$ (2,1)0	×		×	118	2243
$\underline{27}$ (2,1)2		×	×	236	2361
$\underline{10}^*(2,0)3$			×	236	2361
$\underline{35}$ (2,2)1		×	×	393	2518
$\underline{28}$ (2,3)0			×	708	2833

$(Q^4)_3$-$(Q^2)_{3^*}$

$f(y,i)s$	$f(y,i)s$	s	i	$NN(1.88)$	$\Delta N(2.17)$	$\Delta\Delta(2.46)$	M_m	M_1	M_2	M_3
$\underline{6}^*(\frac{4}{3},0)1$	$\underline{3}^*(\frac{2}{3},0)0$	1	0	×			−260	2110	2331	2536
$\underline{15}$ $(\frac{4}{3},1)0$	$\underline{3}^*(\frac{2}{3},0)0$	0	1	×			−170	2200	2421	2626
$\underline{15}$ $(\frac{4}{3},1)1$	$\underline{3}^*(\frac{2}{3},0)0$	1	1	×	×		−125	2245	2466	2671
$\underline{15}$ $(\frac{4}{3},1)2$	$\underline{3}^*(\frac{2}{3},0)0$	2	1		×		−35	2335	2556	2761
$\underline{6}^*(\frac{4}{3},0)1$	$\underline{6}$ $(\frac{2}{3},1)1$	0,1,2	1	×	×		−33	2337	2558	
$\underline{15}$ $(\frac{4}{3},1)0$	$\underline{6}$ $(\frac{2}{3},1)1$	1	0,1,2	×	×		57	2427	2648	
$\underline{15}$ $(\frac{4}{3},1)1$	$\underline{6}$ $(\frac{2}{3},1)1$	0,1,2	0,1,2	×	×	×	102	2472	2693	
$\underline{15}_s(\frac{4}{3},2)1$	$\underline{3}^*(\frac{2}{3},0)0$	1	2		×	×	145	2515	2736	
$\underline{15}$ $(\frac{4}{3},1)2$	$\underline{6}$ $(\frac{2}{3},1)1$	1,2,3	0,1,2		×	×	192	2562	2783	
$\underline{15}_s(\frac{4}{3},2)1$	$\underline{6}$ $(\frac{2}{3},1)1$	0,1,2	1,2,3		×	×	372	2742	2963	

$(Q^5)_{3^*}$-$(Q)_3$

$f(y,i)s$	$f(y,i)s$	s	i	M_m	M_1	M_2
$\underline{15}^*(\frac{5}{3},\frac{1}{2})\frac{1}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	0,1	0,1	0	2370	2591
$\underline{15}^*(\frac{5}{3},\frac{1}{2})\frac{3}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	1,2	0,1	63	2433	2654
$\underline{15}^*(\frac{5}{3},\frac{1}{2})\frac{5}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	2,3	0,1	167	2537	2758
$\underline{24}$ $(\frac{5}{3},\frac{3}{2})\frac{1}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	0,1	1,2	188	2558	2779
$\underline{24}$ $(\frac{5}{3},\frac{3}{2})\frac{3}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	1,2	1,2	251	2621	2842
$\underline{21}$ $(\frac{5}{3},\frac{5}{2})\frac{1}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	0,1	2,3	502	2872	3093

$(Q^4)_{6^*}$-$(Q^2)_6$

$f(y,i)s$	$f(y,i)s$	s	i	M_m	M_1
$\underline{6}^*(\frac{4}{3},0)0$	$\underline{3}^*(\frac{2}{3},0)1$	1	0	−96	2405
$\underline{6}^*(\frac{4}{3},0)0$	$\underline{6}$ $(\frac{2}{3},1)0$	0	1	18	2519
$\underline{6}^*(\frac{4}{3},0)2$	$\underline{3}^*(\frac{2}{3},0)1$	1,2,3	0	39	2540
$\underline{15}$ $(\frac{4}{3},1)1$	$\underline{3}^*(\frac{2}{3},0)1$	0,1,2	1	84	2585
$\underline{6}^*(\frac{4}{3},0)2$	$\underline{6}$ $(\frac{2}{3},1)0$	2	1	153	2654
$\underline{15}$ $(\frac{4}{3},1)1$	$\underline{6}$ $(\frac{2}{3},1)0$	1	0,1,2	198	2699
$\underline{15}_s(\frac{4}{3},2)0$	$\underline{3}^*(\frac{2}{3},0)1$	1	2	310	2811
$\underline{15}_s(\frac{4}{3},2)0$	$\underline{6}$ $(\frac{2}{3},1)0$	0	1,2,3	423	2924

TABLE V. *(Continued.)*

$(Q^3)_8$-$(Q^3)_8$					
$f(y,i)s$	$f(y,i)s$	s	i	M_m	M_1
$\underline{8}(1,\frac{1}{2})\frac{1}{2}$	$\underline{8}(1,\frac{1}{2})\frac{1}{2}$	0, 1	0, 1 [a]	−74	2409
$\underline{8}(1,\frac{1}{2})\frac{3}{2}$	$\underline{8}(1,\frac{1}{2})\frac{1}{2}$	1, 2	0, 1	0	2483
$\underline{8}(1,\frac{1}{2})\frac{3}{2}$	$\underline{8}(1,\frac{1}{2})\frac{3}{2}$	0, 1, 2, 3	0, 1 [a]	74	2557
$\underline{10}(1,\frac{3}{2})\frac{1}{2}$	$\underline{8}(1,\frac{1}{2})\frac{1}{2}$	0, 1	1, 2	149	2632
$\underline{10}(1,\frac{3}{2})\frac{1}{2}$	$\underline{8}(1,\frac{1}{2})\frac{3}{2}$	1, 2	1, 2	223	2706
$\underline{10}(1,\frac{3}{2})\frac{1}{2}$	$\underline{10}(1,\frac{3}{2})\frac{1}{2}$	0, 1	0, 1, 2, 3 [a]	372	2859

[a] For i even, either s even, l odd or s odd, l even. For i odd, either s even, l even or s odd, l odd.

diquark and quark recouple after quark-antiquark creation to a baryon and meson. If a large enough phase space is available, those $(Q^4)_3$-$(Q^2)_{3*}$ states decay easily into $NN, N\Delta$, and $NN\pi$. Probably their widths still are large ($\Gamma \gtrsim 300$ MeV) and their elasticity is small ($x < 0.3$).

The $(Q^4)_3$ cluster with $(\underline{f}, s; \Delta) = (\underline{15}, 2; 2)$ is different. It cannot easily fission in a Δ and a quark, because $\Delta_{12} = 0$, and also the fission into a nucleon and a quark is suppressed because it cannot happen in S waves and has to be accompanied by a spin flip. This cluster therefore is less unstable and the width of the $(Q^4)_3$-$(Q^2)_{3*}$, built from this cluster, might not be too large, and the elasticity will still be small. The resonances with the highest total angular momentum ($j = l + s$), which are most easy to detect, are the resonances in the $^3F_3, {}^1G_4, {}^3H_5, \ldots$ waves. Therefore, we think that it is the $(Q^4)_3$-$(Q^2)_{3*}$ states, containing the $s = 2$ (Q^4) cluster, that have been observed experimentally.

$D(0, J^P; 2.110)$ with $J^P = 0^-$, 1^-, and 2^- are the lowest nonstrange dibaryon resonances. The $D(0, 1^-; 2.11)$ can decay in the $^1P_1$ NN wave. It has a strong coupling to NN and cannot decay into $NN\pi$ in S waves, and therefore should have a large elasticity. This $^1P_1$ resonance, predicted at $T_{1ab} \simeq 0.5$ GeV is an important test for the validity of this model for orbitally excited Q^4-Q^2 dibaryon resonances. In the neighborhood of this $D(0, 1^-; 2.11)$ resonance, no other $I = 0$ dibaryon resonances are predicted. Because its mass is so low it will probably give a clear signal in the $^1P_1$ NN wave. We would like to urge the experimentalists to look for this resonance in np scattering experiments in the mass range $2.06 < M < 2.16$ GeV, that is the laboratory momentum range $0.94 < P_{1ab} < 1.23$ GeV/c, or the laboratory kinetic-energy range $390 < T_{1ab} < 610$ MeV. The presence of this resonance in the lower part of this range is perhaps already excluded by present day experiments.

The states with $J^P = 0^-$ and 2^- are extraneous states.[26] The quantum numbers $(I, J^P) = (0, 0^-)$ and $(0, 2^-)$ are forbidden for the NN system. Thus these extraneous states cannot decay into NN; however, they can decay into $NN\pi$. They can be produced in the reaction

$$pp \to \pi^+ X^+, \quad X^+ \searrow NN\pi\,.$$

The $D(0, 2^-; 2.11)$ is especially likely to be rather narrow as it cannot decay into $NN\pi$ in S waves.

$D(1, 1^-; 2.200)$ couples to the $^3P_1$ NN wave and to $NN\pi$ in S waves. $D(1, J^P; 2.245)$ with $J^P = 0^-$, 1^-, and 2^- couples to the ^3P NN waves, to $N\Delta$, and to $NN\pi$. All these resonances are probably rather unstable and have small elasticity.

$D(1, J^P; 2.335)$ with $J^P = 1^-$, 2^-, and 3^- contains the relatively stable (Q^4) cluster with spin $s = 2$. They yield $^3P_1$, $^3P_2 + {}^3F_2$, and $^3F_3$ NN resonances whose widths are not too large, however, with a small elasticity. Experimentally, structure is seen in the ^3P and ^3F waves ($J \leq 3$) in the region 2.2–2.3 GeV (Refs. 36 and 37). Clear evidence exists for a resonance in the $^3F_3$ NN wave (Ref. 36) with a small elasticity ($x \simeq 0.2$). The dispersion analysis[37] shows that the structure comes mainly from the triplet uncoupled waves $^3P_1$ and $^3F_3$. We think that the complete structure in this region is rather complex, due to the presence of many dibaryons. Including the $D(1, J^P; 2.337)$ as well, there are predicted in the region 2.20–2.35 GeV two $^3P_0$, six $^3P_1$, four $^3P_2 + {}^3F_2$, and two $^3F_3$ NN resonances. As $J = 3$ is the highest spin and one of the $^3F_3$ resonances is somewhat stable, it is understandable why this resonance is most clearly seen. The great number of $^3P_1$ resonances might explain the effect in the triplet uncoupled channels.[37]

$D(0, J^P; 2.331)$ with $J^P = 1^+, 2^+$, and 3^+ is the lowest $l = 2$ dibaryon. It couples to NN and $NN\pi$ through tunneling and quark-antiquark creation, respectively. It probably is a rather inelastic, unstable resonance.

$D(1,J^P;2.556)$ with $J^P=0^+$, 1^+, 2^+, 3^+, and 4^+ is the $l=2$ $(Q^4)_3$-$(Q^2)_{3*}$ dibaryon which contains the $s=2$ (Q^4) cluster. The $D(1,1^+)$ and $D(1,3^+)$ are extraneous in NN. While many dibaryons (most of them unstable) with $J^P=2^+$ ($^1D_2$ NN wave) appear in this region, the $D(1,0^+)$ and $D(1,4^+)$ resonances are more isolated. Therefore, resonances in the $^1S_0$ and $^1G_4$ NN waves will show the clearest resonant behavior, and they are candidates for the experimentally observed structure in the region 2.4–2.5 GeV.

The higher recurrences of the $D(1,J^P;2.335)$ and $D(1,J^P;2.556)$ lie at 2.76, 2.95, and 3.13 GeV for $l=3$, 4, and 5.

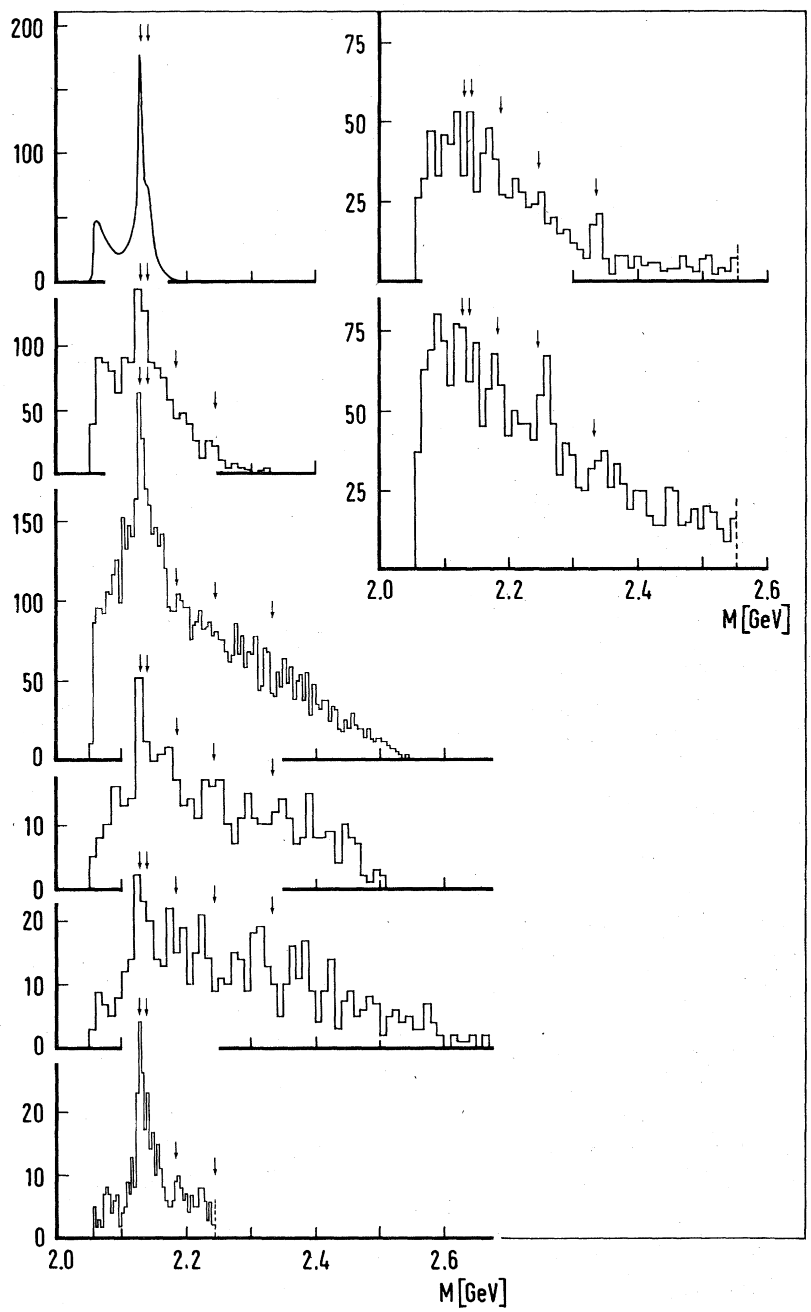

FIG. 5. Λp invariant-mass plots in the reaction $K^-d\to\Lambda p\pi^-$ (left) or K^--nucleus interactions (right). On the left from top to bottom taken from Refs. 50, 42, 40, 43, 41, and 44. On the right from Refs. 45 and 47. The plots are ordered by their (increasing) K^- incident momentum.

We think that the splittings due to the color-magnetic interaction are reliable in so far as we may neglect the final-state interactions. It appears that for the nonstrange dibaryon resonances $(Q^4)_3$-$(Q^2)_{3*}$, the observed masses are about 50–100 MeV lower than the predicted ones, but this of course strongly depends on the assignments. Moreover, it is difficult to determine an experimental mass. This depends on the method of analysis, e.g., for the $^3F_3$ in Refs. 36 and 37.

The importance of dibaryon resonances, other than Q^6 or $(Q^4)_3$-$(Q^2)_{3*}$, is at present not clear to us. The Q^5-Q resonances probably do not couple strongly to BB channels. The nonstrange $(Q^3)_8$-$(Q^3)_8$ dibaryons are probably very unstable. Through gluon exchange (electric) they easily couple to BB channels.

V. THE $Y=1$ DIBARYON RESONANCES[8]

Experimental evidence exists for several $Y=1$ dibaryon resonances $B^2(I, \text{mass})$. The evidence for the $I=\frac{1}{2}$ resonances comes from Λp invariant-mass plots.[40-50] A collection of such plots is given in Fig. 5.

The most pronounced enhancement, $B(\frac{1}{2}, 2.13)$, lies near the ΣN threshold with a mass $M=2.129$ GeV and a width $\Gamma \simeq 6$ MeV. This certainly is not a candidate for a six-quark state; rather, just like the deuteron, it is explained very well in potential theory[55,56] as being a ΣN "bound" state showing up as a ΛN resonance. This enhancement is accompanied by a shoulder[50] which can be fit by a Breit-Wigner resonance $B^2(\frac{1}{2}, 2.14)$ with $M = 2.139$ GeV and $\Gamma = 9$ MeV.

Recently, Shahbazian and coworkers at Dubna[47] determined Λp invariant-mass spectra in the reactions $n^{12}\text{C} \to \Lambda(mp)X$ and $\pi^{-12}C \to \Lambda(mp)X$. They found evidence for two more enhancements: $B^2(\frac{1}{2}, 2.18)$ and $B^2(\frac{1}{2}, 2.25)$. They try to explain the $B^2(\frac{1}{2}, 2.18)$ enhancement as an effect due to $\Sigma N \to \Lambda p$ conversion at large relative momenta. They want to explain for the $B^2(\frac{1}{2}, 2.14)$ resonance the same way. We prefer to retain the resonance explanation for the $B^2(\frac{1}{2}, 2.18)$ enhancement. We note that it is also recognizable (although of course not statistically significant) in most of the other analyses of Fig. 5.

The resonance $B^2(\frac{1}{2}, 2.25)$ with mass $M=2.256$ GeV and width $\Gamma \sim 15$ MeV is a 5-to-6-standard-deviation effect in the Dubna experiments. This state also shows up weakly in most of the other analyses. Shahbazian[47] gives arguments why the $B^2(\frac{1}{2}, 2.25)$ state is clearly visible in their $n^{12}\text{C}$ and $\pi^{-12}\text{C}$ experiments, while it is not clearly visible in the K^-d experiments. Around 2.34 GeV, an enhancement $B^2(\frac{1}{2}, 2.34)$ shows up in several analyses.[45-48] Beillière *et al.*[45] calculate a 2.8-standard-deviation significance, but conclude for no evidence for a resonance at this position.

In the $\Lambda p \pi^+$ invariant-mass plots Shahbazian *et al.*[57] found evidence for $I=\frac{3}{2}$ resonances around 2.5 and 2.99 GeV. A state at 2.5 GeV also shows up in the $\Lambda p \pi^-$ invariant-mass plots.

In Table VI we have listed the predicted $Y=1$ dibaryon resonances. What strikes us is the enormous number of predicted resonances. In order to show up in the experimental data, a resonance must have strong enough coupling to the ΛN and/or ΣN channels and its width may not be unreasonably

TABLE VI. $Y=1$ dibaryon resonances. All masses are quoted in MeV.

$f(y,i)s$	$(Q^6)_1$ ΛN (2.05)	ΣN (2.13)	$\Sigma^* N$ (2.32)	$\Lambda\Delta$ (2.35)	$\Sigma\Delta$ (2.42)	$\Sigma^*\Delta$ (2.62)	M_m	M_0
$\underline{8}$ $(1,\frac{1}{2})1$	×	×	×		×		−126	2169
$\underline{8}$ $(1,\frac{1}{2})2$			×		×		−54	2241
$\underline{10}$ $(1,\frac{3}{2})1$		×	×	×	×		36	2331
$\underline{10}^*(1,\frac{1}{2})1$	×	×				×	36	2331
$\underline{27}$ $(1,\frac{1}{2})0$	×	×				×	108	2403
$\underline{27}$ $(1,\frac{3}{2})0$		×				×	108	2403
$\underline{27}$ $(1,\frac{1}{2})2$			×		×	×	216	2511
$\underline{27}$ $(1,\frac{3}{2})2$			×	×	×	×	216	2511
$\underline{10}^*(1,\frac{1}{2})3$						×	216	2511
$\underline{35}$ $(1,\frac{3}{2})1$			×	×	×	×	361	2656
$\underline{35}$ $(1,\frac{5}{2})1$					×	×	361	2656
$\underline{28}$ $(1,\frac{5}{2})0$						×	649	2944

TABLE VI. *(Continued.)*

						$(Q^4)_3-(Q^2)_{3^*}$						
$f(y,i)\,s$	$f(y,i)s$	s	i	ΛN (2.05)	ΣN (2.13)	$\Sigma^* N$ (2.32)	$\Lambda\Delta$ (2.35)	$\Sigma\Delta$ (2.42)	$\Sigma^*\Delta$ (2.62)	M_m	M_1	M_2
$\underline{3}\ (\frac{1}{3},\frac{1}{2})0$	$\underline{3}^*(\frac{2}{3},0)\,0$	0	$\frac{1}{2}$	×	×					−411	2112	2322
$\underline{3}\ (\frac{1}{3},\frac{1}{2})1$	$\underline{3}^*(\frac{2}{3},0)0$	1	$\frac{1}{2}$	×	×					−371	2152	2362
$\underline{6}^*(\frac{1}{3},\frac{1}{2})1$	$\underline{3}^*(\frac{2}{3},0)0$	1	$\frac{1}{2}$	×	×					−250	2273	2483
$\underline{6}^*(\frac{4}{3},0)1$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})\,0$	1	$\frac{1}{2}$	×	×					−231	2292	2502
$\underline{3}\ (\frac{1}{3},\frac{1}{2})0$	$\underline{6}\ (\frac{2}{3},1)1$	1	$\frac{1}{2},\frac{3}{2}$	×	×	×	×	×		−184	2339	2549
$\underline{15}\ (\frac{1}{3},\frac{1}{2})0$	$\underline{3}^*(\frac{2}{3},0)0$	0	$\frac{1}{2}$	×	×					−170	2353	2563
$\underline{15}\ (\frac{1}{3},\frac{3}{2})0$	$\underline{3}^*(\frac{2}{3},0)0$	0	$\frac{3}{2}$		×					−170	2353	2563
$\underline{3}\ (\frac{1}{3},\frac{1}{2})1$	$\underline{6}\ (\frac{2}{3},1)1$	0, 1, 2	$\frac{1}{2},\frac{3}{2}$	×	×	×	×	×		−144	2379	2589
$\underline{15}\ (\frac{4}{3},1)0$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})0$	0	$\frac{1}{2},\frac{3}{2}$	×	×					−140	2383	2593
$\underline{15}\ (\frac{1}{3},\frac{1}{2})1$	$\underline{3}^*(\frac{2}{3},0)0$	1	$\frac{1}{2}$	×	×	×				−130	2393	2603
$\underline{15}\ (\frac{1}{3},\frac{3}{2})1$	$\underline{3}^*(\frac{2}{3},0)0$	1	$\frac{3}{2}$		×	×	×			−130	2393	2603
$\underline{15}\ (\frac{4}{3},1)1$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})0$	1	$\frac{1}{2},\frac{3}{2}$	×	×		×	×		−95	2428	2638
$\underline{15}\ (\frac{1}{3},\frac{1}{2})2$	$\underline{3}^*(\frac{2}{3},0)0$	2	$\frac{1}{2}$			×				−50	2473	2683
$\underline{15}\ (\frac{1}{3},\frac{3}{2})2$	$\underline{3}^*(\frac{2}{3},0)0$	2	$\frac{3}{2}$			×	×			−50	2473	2683
$\underline{6}^*(\frac{4}{3},0)1$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})1$	0, 1, 2	$\frac{1}{2}$	×	×	×				−43	2480	2690
$\underline{6}^*(\frac{1}{3},\frac{1}{2})1$	$\underline{6}\ (\frac{2}{3},1)1$	0, 1, 2	$\frac{1}{2},\frac{3}{2}$	×	×	×	×	×		−23	2500	2710
$\underline{15}\ (\frac{4}{3},1)2$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})0$	2	$\frac{1}{2},\frac{3}{2}$				×	×		−5	2518	2728
$\underline{15}\ (\frac{4}{3},1)0$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})1$	1	$\frac{1}{2},\frac{3}{2}$	×	×	×				47	2570	
$\underline{15}\ (\frac{1}{3},\frac{1}{2})0$	$\underline{6}\ (\frac{2}{3},1)1$	1	$\frac{1}{2},\frac{3}{2}$	×	×	×	×	×		57	2580	
$\underline{15}\ (\frac{1}{3},\frac{3}{2})0$	$\underline{6}\ (\frac{2}{3},1)1$	1	$\frac{1}{2},\frac{3}{2},\frac{5}{2}$		×			×		57	2580	
$\underline{15}\ (\frac{4}{3},1)1$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})1$	0, 1, 2	$\frac{1}{2},\frac{3}{2}$	×	×	×	×	×	×	92	2615	
$\underline{15}\ (\frac{1}{3},\frac{1}{2})1$	$\underline{6}\ (\frac{2}{3},1)1$	0, 1, 2	$\frac{1}{2},\frac{3}{2}$	×	×	×	×	×	×	97	2620	
$\underline{15}\ (\frac{1}{3},\frac{3}{2})1$	$\underline{6}\ (\frac{2}{3},1)1$	0, 1, 2	$\frac{1}{2},\frac{3}{2},\frac{5}{2}$		×	×		×	×	97	2620	
$\underline{15}_s(\frac{1}{3},\frac{3}{2})1$	$\underline{3}^*(\frac{2}{3},0)0$	1	$\frac{3}{2}$			×				110	2633	
$\underline{15}_s(\frac{4}{3},2)1$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})0$	1	$\frac{3}{2},\frac{5}{2}$				×	×		175	2698	
$\underline{15}\ (\frac{1}{3},\frac{1}{2})2$	$\underline{6}\ (\frac{2}{3},1)1$	1, 2, 3	$\frac{1}{2},\frac{3}{2}$			×			×	177	2700	
$\underline{15}\ (\frac{1}{3},\frac{3}{2})2$	$\underline{6}\ (\frac{2}{3},1)1$	1, 2, 3	$\frac{1}{2},\frac{3}{2},\frac{5}{2}$			×		×	×	177	2700	
$\underline{15}\ (\frac{4}{3},1)2$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})1$	1, 2, 3	$\frac{1}{2},\frac{3}{2}$				×	×	×	182	2705	
$\underline{15}_s(\frac{1}{3},\frac{3}{2})1$	$\underline{6}\ (\frac{2}{3},1)1$	0, 1, 2	$\frac{1}{2},\frac{3}{2},\frac{5}{2}$			×		×	×	337	2860	
$\underline{15}_s(\frac{4}{3},2)1$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})1$	0, 1, 2	$\frac{3}{2},\frac{5}{2}$				×	×	×	362	2885	

TABLE VI. *(Continued.)*

$f(y,i)$	$(Q^5)_{3^*}$–$(Q)_3$ $f(y,i)$	s	i	M_m	M_1	M_2
$\underline{3}^*(\frac{2}{3},0)\frac{1}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	0,1	$\frac{1}{2}$	−227	2296	2506
$\underline{3}^*(\frac{2}{3},0)\frac{3}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	1,2	$\frac{1}{2}$	−170	2353	2563
$\underline{6}\ (\frac{2}{3},1)\frac{1}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	0,1	$\frac{1}{2},\frac{3}{2}$	−114	2409	2619
$\underline{6}\ (\frac{2}{3},1)\frac{3}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	1,2	$\frac{1}{2},\frac{3}{2}$	−57	2466	2676
$\underline{15}^*(\frac{2}{3},0)\frac{1}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	0,1	$\frac{1}{2}$	0	2523	2733
$\underline{15}^*(\frac{2}{3},1)\frac{1}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	0,1	$\frac{1}{2},\frac{3}{2}$	0	2523	2733
$\underline{15}^*(\frac{5}{3},\frac{1}{2})\frac{1}{2}$	$\underline{3}(-\frac{2}{3},0)\frac{1}{2}$	0,1	$\frac{1}{2}$	0	2523	2733
$\underline{15}^*(\frac{2}{3},0)\frac{3}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	1,2	$\frac{1}{2}$	57	2580	
$\underline{15}^*(\frac{2}{3},1)\frac{3}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	1,2	$\frac{1}{2},\frac{3}{2}$	57	2580	
$\underline{15}^*(\frac{5}{3},\frac{1}{2})\frac{3}{2}$	$\underline{3}(-\frac{2}{3},0)\frac{1}{2}$	1,2	$\frac{1}{2}$	63	2586	
$\underline{15}^*(\frac{2}{3},0)\frac{5}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	2,3	$\frac{1}{2}$	152	2675	
$\underline{15}^*(\frac{2}{3},1)\frac{5}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	2,3	$\frac{1}{2},\frac{3}{2}$	152	2675	
$\underline{15}^*(\frac{5}{3},\frac{1}{2})\frac{5}{2}$	$\underline{3}(-\frac{2}{3},0)\frac{1}{2}$	2,3	$\frac{1}{2}$	167	2690	
$\underline{24}\ (\frac{2}{3},1)\frac{1}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	0,1	$\frac{1}{2},\frac{3}{2}$	170	2693	
$\underline{24}\ (\frac{2}{3},2)\frac{1}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	0,1	$\frac{3}{2},\frac{5}{2}$	170	2693	
$\underline{24}\ (\frac{5}{3},\frac{3}{2})\frac{1}{2}$	$\underline{3}(-\frac{2}{3},0)\frac{1}{2}$	0,1	$\frac{3}{2}$	188	2711	
$\underline{24}\ (\frac{2}{3},1)\frac{3}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	1,2	$\frac{1}{2},\frac{3}{2}$	227	2750	
$\underline{24}\ (\frac{2}{3},2)\frac{3}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	1,2	$\frac{3}{2},\frac{5}{2}$	227	2750	
$\underline{24}\ (\frac{5}{3},\frac{3}{2})\frac{3}{2}$	$\underline{3}(-\frac{2}{3},0)\frac{1}{2}$	1,2	$\frac{3}{2}$	251	2774	
$\underline{21}\ (\frac{2}{3},2)\frac{1}{2}$	$\underline{3}(\frac{1}{3},\frac{1}{2})\frac{1}{2}$	0,1	$\frac{3}{2},\frac{5}{2}$	454	2977	
$\underline{21}\ (\frac{5}{3},\frac{5}{2})\frac{1}{2}$	$\underline{3}(-\frac{2}{3},0)\frac{1}{2}$	0,1	$\frac{5}{2}$	502	3025	

large.

The Q^6 states with $J^P = 0^+$ in the flavor irreps $\underline{8}$ and $\underline{27}$, and with $J^P = 1^+$ in the irreps $\underline{8}$, $\underline{10}$, and $\underline{10}^*$, can decay spontaneously in the S-wave ΛN or ΣN channels. As in the NN case, we expect these states to have a very large width and therefore not to be visible in invariant-mass plots.

$D(\frac{1}{2}, 2^+; 2.24)$ is the lowest Q^6 state which could be visible. It belongs to an octet and couples to the ${}^1D_2$ and ${}^3D_2$ ΛN and ΣN channels. As can be seen in Table VI, it also couples to the S-wave $\Sigma^*(1385)N$ and $\Sigma\Delta$ channels, but its mass is below the corresponding thresholds. We would like to make the assignment $D(\frac{1}{2}, 2^+; 2.24) \equiv B^2(\frac{1}{2}; 2.25)$. Because this state is above the $\Lambda N\pi$ threshold ($E_{\text{th}} = 2.19$ GeV), this state could also decay via $Q\bar{Q}$-pair creation. The final state must then also contain an angular-momentum barrier. The observed small width is perhaps not in contradiction with this assignment.

$D(\frac{1}{2}, 2^+; 2.51)$ and $D(\frac{3}{2}, 2^+; 251)$ are companions of the $D(2, 1, 2^+; 2.36)$ dinucleon resonance in the irrep $\underline{27}$. They couple not only to the ${}^1D_2$ ΛN and $I = \frac{1}{2}$ ΣN channel and to the ${}^1D_2$ $I = \frac{3}{2}$ ΣN channel, but also to the ${}^5S_2$ $\Sigma^* N$ and $\Sigma\Delta$ channels. Because their mass is above the thresholds for these latter channels we expect, as observed in the NN case, the resonance poles to shift (due to the final-state interactions) to the neighborhood of these thresholds. These states would therefore to be expected to have an experimental mass of about 2.32 GeV.

$D(\frac{1}{2}, 3^+; 2.51)$ is a companion of the ${}^3D_3$ NN resonance $D(2, 0, 3^+; 2.36)$ in the irrep $\underline{10}^*$. It couples to the ${}^3D_3$ ΛN and ΣN channels and to the ${}^7S_3$ $\Sigma^*\Delta$ channel. It is below the threshold for the latter channel and it is also coupled to $\Lambda N\pi$ via $Q\bar{Q}$-pair creation. We expect, therefore, a width for this state of the order of 100 MeV.

Having discussed the relevant Q^6 states, we will turn now our attention to the orbitally excited states. Of these we will discuss only the states with the lowest masses.

TABLE VI. *(Continued.)*

$f(y,i)s$	$(Q^4)_{6^*}$-$(Q^2)_6$ $f(y,i)s$	s	i	M_m	M_1
$\underline{3}\ (\frac{1}{3},\frac{1}{2})1$	$\underline{3}^*(\frac{2}{3},0)1$	0,1,2	$\frac{1}{2}$	−169	2478
$\underline{6}^*(\frac{4}{3},0)0$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})1$	1	$\frac{1}{2}$	−91	2556
$\underline{6}^*(\frac{1}{3},\frac{1}{2})0$	$\underline{3}^*(\frac{2}{3},0)1$	1	$\frac{1}{2}$	−89	2558
$\underline{3}\ (\frac{1}{3},\frac{1}{2})1$	$\underline{6}\ (\frac{2}{3},1)0$	1	$\frac{1}{2},\frac{3}{2}$	−55	2592
$\underline{6}^*(\frac{4}{3},0)0$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})0$	0	$\frac{1}{2}$	3	2650
$\underline{6}^*(\frac{1}{3},\frac{1}{2})0$	$\underline{6}\ (\frac{2}{3},1)0$	0	$\frac{1}{2},\frac{3}{2}$	25	2672
$\underline{6}^*(\frac{1}{3},\frac{1}{2})2$	$\underline{3}^*(\frac{2}{3},0)1$	1,2,3	$\frac{1}{2}$	32	2679
$\underline{6}^*(\frac{4}{3},0)2$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})1$	1,2,3	$\frac{1}{2}$	44	2691
$\underline{15}\ (\frac{1}{3},\frac{1}{2})1$	$\underline{3}^*(\frac{2}{3},0)1$	0,1,2	$\frac{1}{2}$	72	2719
$\underline{15}\ (\frac{1}{3},\frac{3}{2})1$	$\underline{3}^*(\frac{2}{3},0)1$	0,1,2	$\frac{3}{2}$	72	2719
$\underline{15}\ (\frac{4}{3},1)1$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})1$	0,1,2	$\frac{1}{2},\frac{3}{2}$	89	2736
$\underline{6}^*(\frac{4}{3},0)2$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})0$	2	$\frac{1}{2}$	138	2785
$\underline{6}^*(\frac{1}{3},\frac{1}{2})2$	$\underline{6}\ (\frac{2}{3},1)0$	2	$\frac{1}{2},\frac{3}{2}$	145	2792
$\underline{15}\ (\frac{4}{3},1)1$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})0$	1	$\frac{1}{2},\frac{3}{2}$	183	2830
$\underline{15}\ (\frac{1}{3},\frac{1}{2})1$	$\underline{6}\ (\frac{2}{3},1)0$	1	$\frac{1}{2},\frac{3}{2}$	185	2832
$\underline{15}\ (\frac{1}{3},\frac{3}{2})1$	$\underline{6}\ (\frac{2}{3},1)0$	1	$\frac{1}{2},\frac{3}{2},\frac{5}{2}$	185	2832
$\underline{15}_S(\frac{1}{3},\frac{3}{2})0$	$\underline{3}^*(\frac{2}{3},0)1$	1	$\frac{3}{2}$	272	2919
$\underline{15}_S(\frac{4}{3},2)0$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})1$	1	$\frac{3}{2},\frac{5}{2}$	315	2962
$\underline{15}_S(\frac{1}{3},\frac{3}{2})0$	$\underline{6}\ (\frac{2}{3},1)0$	0	$\frac{1}{2},\frac{3}{2},\frac{5}{2}$	386	3033
$\underline{15}_S(\frac{4}{3},2)0$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})0$	0	$\frac{3}{2},\frac{5}{2}$	408	3055

$D(\frac{1}{2},1^-;2.11)$ belongs to a nonet in the configuration $(Q^4)_3$-$(Q^2)_{3^*}$ with $S=0$, $l^P=1^-$, and, therefore $J^P=1^-$. We would like to assign this state to the shoulder $B^2(\frac{1}{2};2.14)$. This state is then coupled to the $^1P_1$ and $^3P_1$ ΛN and ΣN channels. This state decays via the tunneling of a nonstrange quark into ΛN or ΣN, or via the tunneling of a strange quark into ΛN. No change of orbital angular momentum and therefore no spin flip is required in this tunneling. The observed small width, $\Gamma\approx 9$ MeV, does not seem unreasonable.

$D(\frac{1}{2},J^P;2.15)$ belongs again to a nonet in the configuration $(Q^4)_3$-$(Q^2)_{3^*}$ but now with $S=1$, $l^P=1^-$ and therefore $J^P=0^-$, 1^-, and 2^-. These states are coupled to the $^3P_0$, $^1P_1+{}^3P_1$, and the $^3P_2+{}^3F_2$ waves of the ΛN and ΣN channels. We would like to assign these states to the $B^2(\frac{1}{2};2.18)$ enhancement. The decay via tunneling goes exactly the same way as for the $D(\frac{1}{2},1^-;2.11)$ state. The assignments of $D(2.11)$ and $D(2.15)$ to the states $B^2(2.14)$ and $B^2(2.18)$ is supported by the fact that their mass difference is only 40 MeV. This mass difference is of color-magnetic origin. We believe that the mass differences between states are much more accurately known in this model (neglecting final-state interactions) than their total mass. It is even surprising that the total masses seem to be only 30 MeV off.

It is noteworthy that extraneous states can also occur in the $Y=1$ channel. For baryons with spin $s=\frac{1}{2}$ belonging to the flavor octet B_8, the flavor representation in the baryon-baryon system B_8B_8 is given by

$$\underline{8}\otimes\underline{8}=(\underline{1}\oplus\underline{8}_S\oplus\underline{27})\oplus(\underline{8}_A\oplus\underline{10}\oplus\underline{10}^*)\,.$$

The flavor part of the wave function is symmetric for the flavor irreps $\underline{1}$, $\underline{8}_S$, and $\underline{27}$, while it is antisymmetric for the irreps $\underline{8}_A$, $\underline{10}$, and $\underline{10}^*$. According to the generalized Pauli principle, the symmetric flavor wave functions are allowed only in the ^1S, ^3P, ^1D, ^3F, etc. waves and the antisymmetric flavor wave functions are allowed only in the ^3S, ^1P, ^3D, ^1F, etc. waves of the B_8B_8 system. Dibaryon resonances belonging to the flavor irrep $\underline{1}$ or $\underline{27}$ with $J^P=1^+$, 3^+, 5^+, etc., and be-

TABLE VI. *(Continued.)*

$f(y,i)s$	$f(y,i)s$	$(Q^3)_8$-$(Q^3)_8$ s	i	M_m	M_1
$\underline{1}(0,0)\frac{1}{2}$	$\underline{8}(1,\frac{1}{2})\frac{1}{2}$	0,1	$\frac{1}{2}$	−262	2368
$\underline{1}(0,0)\frac{1}{2}$	$\underline{8}(1,\frac{1}{2})\frac{3}{2}$	1,2	$\frac{1}{2}$	−188	2442
$\underline{8}(0,0)\frac{1}{2}$	$\underline{8}(1,\frac{1}{2})\frac{1}{2}$	0,1	$\frac{1}{2}$	−69	2561
$\underline{8}(0,1)\frac{1}{2}$	$\underline{8}(1,\frac{1}{2})\frac{1}{2}$	0,1	$\frac{1}{2},\frac{3}{2}$	−69	2561
$\underline{1}(0,0)\frac{1}{2}$	$\underline{10}(1,\frac{3}{2})\frac{1}{2}$	0,1	$\frac{3}{2}$	−39	2591
$\underline{8}(0,0)\frac{3}{2}$	$\underline{8}(1,\frac{1}{2})\frac{1}{2}$	1,2	$\frac{1}{2}$	−5	2625
$\underline{8}(0,1)\frac{3}{2}$	$\underline{8}(1,\frac{1}{2})\frac{1}{2}$	1,2	$\frac{1}{2},\frac{3}{2}$	−5	2625
$\underline{8}(0,0)\frac{1}{2}$	$\underline{8}(1,\frac{1}{2})\frac{3}{2}$	1,2	$\frac{1}{2}$	5	2635
$\underline{8}(0,1)\frac{1}{2}$	$\underline{8}(1,\frac{1}{2})\frac{3}{2}$	1,2	$\frac{1}{2},\frac{3}{2}$	5	2635
$\underline{8}(0,0)\frac{3}{2}$	$\underline{8}(1,\frac{1}{2})\frac{3}{2}$	0,1,2,3	$\frac{1}{2}$	69	2699
$\underline{8}(0,1)\frac{3}{2}$	$\underline{8}(1,\frac{1}{2})\frac{3}{2}$	0,1,2,3	$\frac{1}{2},\frac{3}{2}$	69	2699
$\underline{10}(0,1)\frac{1}{2}$	$\underline{8}(1,\frac{1}{2})\frac{1}{2}$	0,1	$\frac{1}{2},\frac{3}{2}$	124	2754
$\underline{8}(0,0)\frac{1}{2}$	$\underline{10}(1,\frac{3}{2})\frac{1}{2}$	0,1	$\frac{3}{2}$	154	2784
$\underline{8}(0,1)\frac{1}{2}$	$\underline{10}(1,\frac{3}{2})\frac{1}{2}$	0,1	$\frac{1}{2},\frac{3}{2},\frac{5}{2}$	154	2784
$\underline{10}(0,1)\frac{1}{2}$	$\underline{8}(1,\frac{1}{2})\frac{3}{2}$	1,2	$\frac{1}{2},\frac{3}{2}$	198	2828
$\underline{8}(0,0)\frac{3}{2}$	$\underline{10}(1,\frac{3}{2})\frac{1}{2}$	1,2	$\frac{3}{2}$	218	2848
$\underline{8}(0,1)\frac{3}{2}$	$\underline{10}(1,\frac{3}{2})\frac{1}{2}$	1,2	$\frac{1}{2},\frac{3}{2},\frac{5}{2}$	218	2848
$\underline{10}(0,1)\frac{1}{2}$	$\underline{10}(1,\frac{3}{2})\frac{1}{2}$	0,1	$\frac{1}{2},\frac{3}{2},\frac{5}{2}$	347	2977

longing to the irreps $\underline{10}$ and $\underline{10}^*$ with $J^P=0^+;0^-,2^-,4^-$, etc., therefore cannot decay into B_8B_8 and are called extraneous to B_8B_8 (see Table VII). As an example, consider the $D(\frac{3}{2},(0^-,1^-,2^-);2.339)$ states. The flavor representation f is found from Table VI; $\underline{f}=\underline{3}\otimes\underline{6}=\underline{8}\oplus\underline{10}$. As $(Y,\bar{I})=(1,\frac{3}{2})$, these states belong to the irrep $\underline{10}$. The states $D(\frac{3}{2},0^-;2.339)$ and $D(\frac{3}{2},2^-;2.339)$ are thus extraneous. Therefore, the decay $D(1,\frac{3}{2},(0^-,2^-);2.339)\to\Sigma N$ is forbidden. The $I=\frac{1}{2}$ analogues of these states,

TABLE VII. Extraneous dibaryon states.

Baryon×baryon	Flavor (isospin)	Allowed BB waves	Extraneous dibaryon states
B_8B_8	$\underline{1}+\underline{8}_D+\underline{27}$	${}^1S_0$ ${}^1D_2\cdots$	$1^+,3^+,5^+,\ldots$
(NN)	$(I=1)$	${}^3P_{0,1,2}\cdots$	
	$\underline{8}_F+\underline{10}+\underline{10}^*$	${}^3S_1$ ${}^3D_{1,2,3}\cdots$	0^+
	$(I=0)$	${}^1P_1\cdots$	$0^-,2^-,4^-,\ldots$
$B_{10}B_{10}$	$\underline{27}+\underline{28}$	${}^1S,{}^5S$ ${}^1D,{}^5D\cdots$	
$(\Delta\Delta)$	$(I=1,3)$	${}^3P,{}^7P\cdots$	
	$\underline{35}+\underline{10}^*$	${}^3S,{}^7S$ ${}^3D,{}^7D\cdots$	0^+
	$(I=0,2)$	${}^1P,{}^5P\cdots$	0^-

TABLE VIII. Y = 0 dibaryon resonances. All masses are quoted in MeV.

$f(y,i)s$	ΛΛ(2.23)	NΞ(2.25)	ΛΣ(2.31)	ΣΣ(2.38)	NΞ*(2.47) $(Q^6)_1$	ΛΣ*(2.50)	ΔΞ(2.55)	ΣΣ*(2.58)	ΔΞ*(2.76)	Ξ*Ξ*(2.77)	M_m	M_0
1 (0,0)0	×	×		×							−297	2164
8 (0,0)1	×	×		×	×			×			−115	2349
8 (0,1)1		×	×	×	×	×	×	×			−115	2349
8 (0,0)2					×			×			−50	2414
8 (0,1)2					×	×	×	×			−50	2414
10 (0,1)1		×	×	×	×	×	×	×			33	2497
10* (0,1)1		×	×	×					×	×	33	2497
27 (0,0)0	×	×		×						×	99	2563
27 (0,1)0		×	×	×					×	×	99	2563
27 (0,2)0				×					×	×	99	2563
27 (0,0)2					×			×		×	198	2662
27 (0,1)2					×	×	×	×	×	×	198	2662
27 (0,2)2							×	×	×	×	198	2662
10* (0,1)3									×	×	198	2662
35 (0,1)1					×	×	×	×	×	×	330	2794
35 (0,2)1							×	×	×	×	330	2794
28 (0,2)0									×	×	594	3058

$D(1,\frac{1}{2},(0^-,1^-,2^-);2.339)$, are not extraneous. As $(Y,I)=(1,\frac{1}{2})$ these states belong to the irrep $\underline{8}$. This irrep, however, generally couples to both the symmetric and antisymmetric octet in the B_8B_8 system. Another instructive example is formed by the Q^4-Q^2 states $D(1,\frac{3}{2},(0^-,1^-,2^-);2.393)$ and $D(1,\frac{3}{2},(0^-,1^-,2^-);2.428)$. The flavor representation is $\underline{15}\otimes\underline{3}=\underline{8}\oplus\underline{10}\oplus\underline{27}$. As the flavor symmetry is broken, the states with $(Y,I)=(1,\frac{3}{2})$ do not belong to either the irrep $\underline{10}$ or $\underline{27}$, but rather are mixtures. The 0^- and 2^- states then decay via the $\underline{27}$ component, which does couple to the B_8B_8 system. The states at 2.393 GeV have the structure $(n^3s)_3$-$(n^2)_{3*}$, and the states at 2.428 GeV have the structure $(n^4)_3$-$(ns)_{3*}$. The energy difference is due to the different color-magnetic interactions of the nonstrange quarks n and strange quark s.

The lowest ΛN resonances are predicted much closer to the ΛN threshold than the lowest NN resonances to the NN threshold. Therefore these ΛN resonances are more pronounced than the NN resonances and it is advisable to plan high-statistics experiments to reconfirm these ΛN resonances. We think here of $K^-d\to\Lambda p\pi^-$ or $K^{-\,4}$He $\to(\Lambda N)X$ at sufficiently high-incident K^- momenta.

VI. $Y=0$ DIBARYON RESONANCES

The calculated masses for the $Y=0$ states are presented in Table VIII. In the $I=0$ channel, the lowest state is $D(0,0^+;2.164)$, which is a $\Lambda\Lambda$ bound state.[17] Only via weak interactions can it decay into ΛN. The states $D(0,1^-;2.295)$ and $D(0,1^-;2.297)$ are predicted not far above the thresholds of the $\Lambda\Lambda$ and ΞN thresholds to which they couple strongly after tunneling. A probably narrow state is $D(0,2^+;2.414)$, which requires a spin flip to decay into ΞN and $\Lambda\Lambda$.

Experimental evidence for an $I=0$ dibaryon resonance is seen in the enhancement in the $\Lambda\Lambda$ invariant mass plots at 2.365 MeV with $\Gamma\simeq 50$ MeV (Refs. 58 and 59). This is a candidate for the $D(0,2^+;2.414)$.

In the $I=1$ channel, the lowest state is $D(1,1^-;2.297)$, decaying into ΞN. $D(1,2^+;2.414)$ and $D(1,3^+;2.662)$ are narrow $(Q^6)_1$ states decaying into $\Sigma\Lambda$ and ΞN after a spin flip.

In the $I=2$ channel, we mention $D(2,1^-;2.538)$ decaying into $\Sigma\Sigma$, and $D(2,(1^-,2^-,3^-);2.658)$ decaying into $\Delta\Xi$ and $\Sigma\Sigma^*$. In both cases these states are the lowest above the thresholds of the channels mentioned.

TABLE VIII. *(Continued.)*

				$(Q^4)_3$–$Q^2)_{3^*}$											
$f(y,i)s$	$f(y,i)s$	s	i	$\Lambda\Lambda(2.23)$	$N\Xi(2.25)$	$\Lambda\Sigma(2.31)$	$\Sigma\Sigma(2.38)$	$N\Xi^*(2.47)$	$\Lambda\Sigma^*(2.50)$	$\Delta\Xi(2.55)$	$\Sigma\Sigma^*(2.58)$	$\Delta\Xi^*(2.76)$	$\Sigma^*\Sigma^*(2.77)$	M_m	M_1
$\underline{3}\ (-\frac{2}{3},0)0$	$\underline{3}^*(\frac{2}{3},0)0$	0	0	×	×									−383	2295
$\underline{3}\ (\frac{1}{3},\frac{1}{2})0$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})0$	0	0, 1	×	×	×	×							−381	2297
$\underline{3}\ (-\frac{2}{3},0)1$	$\underline{3}^*(\frac{2}{3},0)0$	1	0	×	×									−348	2330
$\underline{3}\ (\frac{1}{3},\frac{1}{2})1$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})0$	0	0, 1	×	×	×	×							−341	2337
$\underline{6}^*(-\frac{2}{3},1)1$	$\underline{3}^*(\frac{2}{3},0)0$	1	1		×	×								−241	2437
$\underline{6}^*(\frac{1}{3},\frac{1}{2})1$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})0$	1	0, 1	×	×	×	×							−221	2457
$\underline{3}\ (\frac{1}{3},\frac{1}{2})0$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})1$	1	0, 1	×	×	×	×	×	×		×			−194	2484
$\underline{15}\ (-\frac{2}{3},0)0$	$\underline{3}^*(\frac{2}{3},0)0$	0	0	×	×									−170	2508
$\underline{15}\ (-\frac{2}{3},1)0$	$\underline{3}^*(\frac{2}{3},0)0$	0	1		×	×								−170	2508
$\underline{3}\ (-\frac{2}{3},0)0$	$\underline{6}\ (\frac{2}{3},1)1$	1	1		×	×			×	×				−157	2521
$\underline{3}\ (\frac{1}{3},\frac{1}{2})1$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})1$	0, 1, 2	0, 1	×	×	×	×	×	×		×			−154	2524
$\underline{15}\ (\frac{1}{3},\frac{1}{2})0$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})0$	0	0, 1	×	×	×	×							−140	2538
$\underline{15}\ (\frac{1}{3},\frac{3}{2})0$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})0$	0	1, 2			×	×							−140	2538
$\underline{15}\ (-\frac{2}{3},0)1$	$\underline{3}^*(\frac{2}{3},0)0$	1	0	×	×			×						−135	2543
$\underline{15}\ (-\frac{2}{3},1)1$	$\underline{3}^*(\frac{2}{3},0)0$	1	1		×	×		×	×					−135	2543
$\underline{3}\ (-\frac{2}{3},0)1$	$\underline{6}\ (\frac{2}{3},1)1$	0, 1, 2	1		×	×			×	×				−121	2557
$\underline{15}\ (\frac{1}{3},\frac{1}{2})1$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})0$	1	0, 1	×	×	×	×		×		×			−100	2578
$\underline{15}\ (\frac{1}{3},\frac{3}{2})1$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})0$	1	1, 2			×	×		×	×	×			−100	2578
$\underline{15}\ (-\frac{2}{3},0)2$	$\underline{3}^*(\frac{2}{3},0)0$	2	0					×						−64	2614
$\underline{15}\ (-\frac{2}{3},1)2$	$\underline{3}^*(\frac{2}{3},0)0$	2	1					×	×					−64	2614
$\underline{6}^*(\frac{4}{3},0)1$	$\underline{6}\ (-\frac{4}{3},0)1$	0, 1, 2	0		×			×						−52	2626
$\underline{6}^*(\frac{1}{3},\frac{1}{2})1$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})1$	0, 1, 2	0, 1	×	×	×	×	×	×		×			−33	2645
$\underline{15}\ (\frac{1}{3},\frac{1}{2})2$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})0$	2	0, 1						×		×			−20	2658
$\underline{15}\ (\frac{1}{3},\frac{3}{2})2$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})0$	2	1, 2						×	×	×			−20	2658
$\underline{6}^*(-\frac{2}{3},1)1$	$\underline{6}\ (\frac{2}{3},1)1$	0, 1, 2	0, 1, 2		×		×			×	×			−14	2664
$\underline{15}\ (\frac{4}{3},1)0$	$\underline{6}\ (-\frac{4}{3},0)1$	1	1		×			×						38	2716
$\underline{15}\ (\frac{1}{3},\frac{1}{2})0$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})1$	1	0, 1	×	×	×	×	×	×		×			47	2725
$\underline{15}\ (\frac{1}{3},\frac{3}{2})0$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})1$	1	1, 2			×	×				×			47	2725
$\underline{15}\ (-\frac{2}{3},0)0$	$\underline{6}\ (\frac{2}{3},1)1$	1	1		×	×			×	×				57	2735
$\underline{15}\ (-\frac{2}{3},1)0$	$\underline{6}\ (\frac{2}{3},1)1$	1	0, 1, 2		×		×			×	×			57	2735
$\underline{15}_s(-\frac{2}{3},1)1$	$\underline{3}^*(\frac{2}{3},0)0$	1	1					×	×					79	2757
$\underline{15}\ (\frac{4}{3},1)1$	$\underline{6}\ (-\frac{4}{3},0)1$	0, 1, 2	1		×			×		×		×		83	2761
$\underline{15}\ (\frac{1}{3},\frac{1}{2})1$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})1$	0, 1, 2	0, 1	×	×	×	×	×	×		×		×	87	2765
$\underline{15}\ (\frac{1}{3},\frac{3}{2})1$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})1$	0, 1, 2	1, 2			×	×		×	×	×	×	×	87	2765
$\underline{15}\ (-\frac{2}{3},0)1$	$\underline{6}\ (\frac{2}{3},1)1$	0, 1, 2	1		×	×		×	×	×		×		92	2770
$\underline{15}\ (-\frac{2}{3},1)1$	$\underline{6}\ (\frac{2}{3},1)1$	0, 1, 2	0, 1, 2		×		×	×		×	×	×	×	92	2770
$\underline{15}_s(\frac{1}{3},\frac{3}{2})1$	$\underline{3}^*(-\frac{1}{3},\frac{1}{2})0$	1	1, 2						×	×	×			140	2818

TABLE VIII. *(Continued.)*

$f(y,i)s$	$f(y,i)s$	s	i	$(Q^4)_3$–$(Q^2)_{3^*}$ $\Lambda\Lambda(2.23)$	$N\Xi(2.25)$	$\Lambda\Sigma(2.31)$	$\Sigma\Sigma(2.38)$	$N\Xi^*(2.47)$	$\Lambda\Sigma^*(2.50)$	$\Lambda\Xi(2.55)$	$\Sigma\Sigma^*(2.58)$	$\Lambda\Xi^*(2.76)$	$\Sigma^*\Sigma^*(2.77)$	M_m	M_1
$\underline{15}\ (-\frac{2}{3},0)2$	$\underline{6}\ (\frac{2}{3},1)1$	1, 2, 3	1					×				×		164	2842
$\underline{15}\ (-\frac{2}{3},1)2$	$\underline{6}\ (\frac{2}{3},1)1$	1, 2, 3	0, 1, 2					×			×	×	×	164	2842
$\underline{15}\ (\frac{1}{3},\frac{1}{2})2$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})1$	1, 2, 3	0, 1						×		×		×	167	2845
$\underline{15}\ (\frac{1}{3},\frac{3}{2})2$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})1$	1, 2, 3	1, 2						×	×	×	×	×	167	2845
$\underline{15}\ (\frac{4}{3},1)2$	$\underline{6}\ (-\frac{4}{3},0)1$	1, 2, 3	1							×		×		174	2852
$\underline{15}_s(-\frac{2}{3},1)1$	$\underline{6}\ (\frac{2}{3},1)1$	0, 1, 2	0, 1, 2					×			×	×	×	306	2984
$\underline{15}_s(\frac{1}{3},\frac{3}{2})1$	$\underline{6}\ (-\frac{1}{3},\frac{1}{2})1$	0, 1, 2	1, 2						×	×	×	×	×	327	3005
$\underline{15}_s(\frac{4}{3},2)1$	$\underline{6}\ (-\frac{4}{3},0)1$	0, 1, 2	2							×		×		354	3032

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*Present address: Los Alamos Scientific Laboratory, Los Alamos, N.M. 87545.

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