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Baryon-antibaryon decays of four-quark states

W. Roberts,* B. Silvestre-Brac, and C. Gignoux
Institut des Sciences Nucléaires, 53, Avenue des Martyrs, 38026 Grenoble CEDEX, France
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We classify all the SU(3) multiplets of T-diquonia consistent with the Pauli principle, and estimate their masses using a potential model. Within the framework of the $^3P_0$ model, we calculate the total and partial decay widths of these states into baryon-antibaryon pairs. We find that both the total widths and the partial widths range from a few MeV to several hundred MeV. We briefly discuss implications for experimental detection of diquonia.

I. INTRODUCTION AND MOTIVATION

Exotic hadrons have captured the imagination of physicists ever since the earliest days of the quark model. Objects with fractional electric charge were the targets of some of the first experimental searches for exotic hadrons.[1] Apart from a few isolated events, it appears safe to say that, within acceptable experimental limits, none have been found.[2]

The advent of quantum chromodynamics[3] (QCD) changed the concept of an exotic hadron. First of all, the color charge of quarks allows "hadrons" to have a large variety of different color charges. For instance, a "meson" ($q\bar{q}$) may belong to a color singlet or octet, while a "baryon" (q^3) may belong to a color singlet, one of two octets, or a decuplet. For many reasons, the present prejudice is that physically observable states must be color singlets, at least at currently accessible energies.

The color-singlet restriction, however, still leaves open the possibility for a rich assortment of hadrons composed of quarks. In addition, the non-Abelian nature of QCD allows the gauge fields of the theory, the gluons, to couple to each other. This allows an even more intriguing spectrum of exotics, as quarkless hadrons or glueballs,[4] as well as hadrons composed of both quarks and gluons[5] (hybrids), may now populate the particle universe.

In the pure quark sector, the lightest of the exotics satisfying the color-singlet condition are the four-quark states ($q^2\bar{q}^2$), which are commonly called diquonia or baryonia.[6] Perhaps the earliest candidates for such states are mesons such as the $a_0(980)$ and $f_0(975)$. Their decay characteristics are not easily accommodated by $q\bar{q}$ structure, and as early as 1977, Jaffe[7] suggested that they may be four-quark states. Today, there is growing support for this hypothesis.[8]

Apart from other mesons like the two mentioned above, many diquonia candidates have been observed in baryon-antibaryon systems (hence the name baryonia), where they may be produced via mechanisms such as that illustrated in Fig. 1. Narrow structures have been observed in $p\bar{p}$ and $p\bar{d}$ (Ref. 9) systems at about 1930 MeV, with widths of the order of 10 MeV. Broader enhancements have also been detected near 2190 and 2400 MeV ($\Gamma \approx 80$–250 MeV). Candidates with strangeness include one narrow structure near 2460 MeV ($\Gamma \approx 20$ MeV) and at least one broad structure in the Λp system produced in the reactions $K^+p \to \bar{\Lambda}pp$ and $K^-p \to \Lambda\bar{p}p$. Perhaps the most well known of such states are the three seen near 3 GeV, commonly called the U(3.1) (Ref. 10). A more exhaustive catalogue of diquonia candidates is given in Ref. 9.

Although most of these diquonia candidates are not firmly established, and in fact some of them have disappeared, they have inspired a number of theoretical papers discussing one aspect or another of their phenomenology.[11–19] Not surprisingly, much of the discussion has centered around their baryon-antibaryon decays. Most of the papers have focused on baryonia consisting solely of light quarks (u,d). This is understandable for two reasons. The first is that baryonia produced by the mechanism of Fig. 1 at $N\bar{N}$ colliders will not contain strange quarks. The second reason is that the number of states possible with strange quarks included is perhaps prohibitive.

The only published paper of which we are aware that treats the baryon-antibaryon decays of diquonia with strangeness is that of Ref. 11. In that paper, a few baryon-antibaryon channels are studied, in proposing that the U(3.1) is a diquonium. We point out, however, that the discussion in that paper is not nearly as exhaustive as that of Jaffe[12] or of Fukugita and Hansson,[13] who catalogue all states consisting of light quarks, and estimate widths for decay into baryon pairs, via quark pair creation, and into meson pairs, via quark rearrangement, for many of them. The work of Ader, Bonnier, and Sood[14] and of Barbour and Gilchrist[15] also treat a number of states. As yet, there have been no attempts at carrying out a similarly systematic analysis of the $B\bar{B}$ decays of diquonia that include strange quarks.

In this paper we propose to partially fill this strange void. Assuming that $SU(3)_f$ symmetry [the f denotes flavor SU(3)] is valid, we estimate the masses of diquonia in a potential model, and evaluate their total and partial decay widths into $B\bar{B}$ channels using the $^3P_0$ model. In this model, a color-singlet quark pair is produced from the QCD vacuum and is combined with the quarks of the decaying diquonia to give a baryon-antibaryon pair. In the version of the model that we use, the diquarks of the diquonia are treated as spectators. This is significant in that it allows only diquonia with a specific color composi-

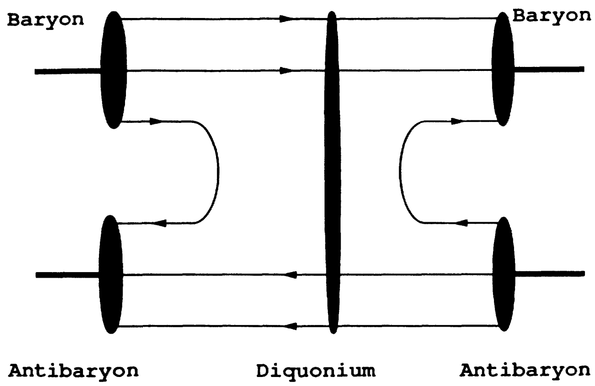

FIG. 1. Mechanism for diquonia or baryonia formation and decay in baryon-antibaryon systems.

tion to decay into baryons. This is discussed further in the next section.

The assumption of SU(3) is essential in lessening the work required. A study of diquonia states with broken SU(3) would be quite tedious since the number of states is very large. We also differ from previous work in the literature with respect to the decay channels, since not only are the baryons allowed to have strangeness, as is necessary [the SU(3) multiplets are used], but we also allow orbital excitations. This means that we can calculate the partial widths for decays into channels containing singlet baryons such as the Λ. Such channels are of some experimental interest, but have not been investigated theoretically before.

The rest of this paper is set out as follows. Section II lists the diquonium and baryon states in which we are interested. In Sec. III, we estimate the masses of the diquonia, and obtain the other parameters needed for calculation of the decay widths. Section IV presents a few salient points of the model used for describing the decays, while we tabulate and discuss the total and partial widths obtained in Sec. V. Section VI is used to present our conclusions and a brief discussion of possible future work.

II. CLASSIFICATION OF STATES

In this section we list the diquonia states that we study. These states are assumed to consist of an S-wave diquark and an S-wave antidiquark with some orbital angular momentum between them. Let us point out that the diquark-antidiquark basis we use is motivated in part by the need for simplicity in treating the decays into baryon-antibaryon pairs. Any other basis would remove the attractive simplifying assumption of diquarks that are spectators in the decay process. Note that the only other basis that has been used for the description of four-quark states is the $(q\bar{q})(q\bar{q})$ basis, where each quark-antiquark pair is in a color octet or singlet, coupled to give a color singlet.

For the basis used here, we note that in color space the diquark may belong to an antitriplet or to a sextet, while the antidiquark may belong to a triplet or an antisextet. To form a color-singlet object, a diquark in an antitriplet must be combined with an antidiquark in a triplet, or a sextet diquark must be combined with an antisextet antidiquark.

The diquonia formed from the latter combination $(\mathbf{6},\bar{\mathbf{6}})$ are the so-called "mock" or M-diquonia, and are of no interest to us here, as their decays into baryon-antibaryon pairs are forbidden without color mixing. This is because the diquark must be combined with a single quark to create a color-singlet baryon. This is only possible if the diquark belongs to the color antitriplet.

The other kind of diquonia $(\bar{\mathbf{3}},\mathbf{3})$, the "true" or T-diquonia, can decay into $B\bar{B}$ pairs within the framework of the $^3P_0$ model, with no need for color mixing. This model is illustrated in Fig. 2 and is briefly discussed in Sec. IV. We thus confine our discussion to states of T-diquonia, and point out that the widths we calculate will correspond to upper limits, since physical states, if they exist, may contain some admixture of M-diquonia. This color restriction is the only one that we place on the states that we study. We discuss all states compatible with the Pauli principle.

The states we consider are classified by the flavor content of the diquarks of which they are made, the flavor multiplet of the diquonia themselves, and the total spin and orbital angular momentum of the diquonia. To begin, we point out that the notation used in discussion of diquonia within $SU(2)_f$ finds a natural extension in $SU(3)_f$ (Ref. 12). A diquark belonging to the antitriplet of $SU(3)_f$ is denoted β, while one belonging to the sextet is denoted δ. Overall antisymmetry of the diquonium wave function under exchange of quarks or antiquarks constrains the spin of the β-type diquark to be zero and that of the δ type to be one. The diquark states are therefore $\beta(\bar{\mathbf{3}},0)$ and $\delta(\mathbf{6},1)$, where the numbers in parentheses are the flavor multiplet and the total spin, respectively.

The 15 diquonia multiplets that can be formed from these diquarks and antidiquarks are shown in Table I. Note that these are not, in general, states of definite G parity, but are, instead, states of definite flavor. The notation used in the table is the same as that used for SU(2): A denotes a diquonium with diquark content $\beta\bar{\beta}$, B corre-

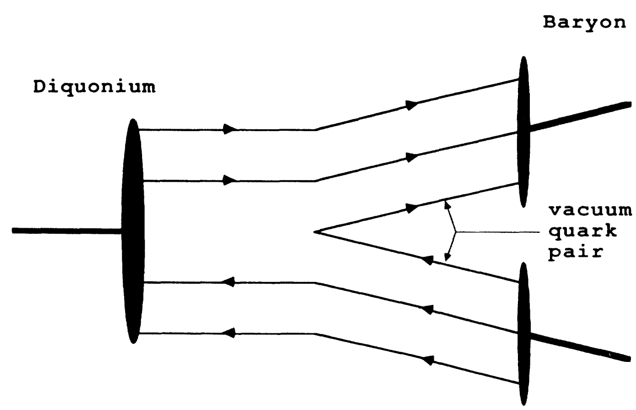

FIG. 2. Diquonium decay into baryon-antibaryon pair via pair creation model.

TABLE I. Baryonia multiplets. The notation in columns 1 and 2 are the same as in Ref. 11. The numbers in parentheses in column 1 refer to the multiplet and the orbital angular momentum, respectively. In the case of the C states, the three numbers are the multiplet, the total spin, and the orbital angular momentum, respectively.

Diquonium states State	Flavor-spin wave function
$A(1,L_D)$	$\beta\bar{\beta}$
$A(8,L_D)$	$\beta\bar{\beta}$
$B(8,L_D)$	$\beta\bar{\delta}$
$B(10,L_D)$	$\beta\bar{\delta}$
$B(8,L_D)$	$\delta\bar{\beta}$
$B(\overline{10},L_D)$	$\delta\bar{\beta}$
$C(1,0,L_D)$	$\delta\bar{\delta}$
$C(8,0,L_D)$	$\delta\bar{\delta}$
$C(27,0,L_D)$	$\delta\bar{\delta}$
$C(1,1,L_D)$	$\delta\bar{\delta}$
$C(8,1,L_D)$	$\delta\bar{\delta}$
$C(27,1,L_D)$	$\delta\bar{\delta}$
$C(1,2,L_D)$	$\delta\bar{\delta}$
$C(8,2,L_D)$	$\delta\bar{\delta}$
$C(27,2,L_D)$	$\delta\bar{\delta}$

sponds to $\beta\bar{\delta}$ and $\delta\bar{\beta}$, while C denotes diquark content $\delta\bar{\delta}$. In the table, the numbers in parentheses indicate the flavor multiplet and orbital angular momentum of each state, respectively. In the case of the C states, the second number in parentheses is the total spin of the state while the third number is the orbital angular momentum of the state. For states A, the total spin is always zero, while for states B it is always one.

In the case of the baryons, we are interested in those baryons in which there is some component of the total

TABLE II. Baryon multiplets. The symbols denote the nonstrange members of each multiplet, where possible, while the numbers in parentheses are twice the total spin and the orbital angular momentum of the baryon, respectively. The fourth column indicates the component of the total flavor-spin wave function that is of interest here.

State	Multiplet	S_B	L_B	Component
$N(1,0)$	8	$\frac{1}{2}$	0	$(\beta+\delta)q/\sqrt{2}$
$\Delta(3,0)$	10	$\frac{3}{2}$	0	δq
$\Lambda(1,1)$	1	$\frac{1}{2}$	1	$\beta q/\sqrt{2}$
$\Delta(1,1)$	10	$\frac{1}{2}$	1	$-\delta q/\sqrt{2}$
$N(3,1)$	8	$\frac{3}{2}$	1	$-\delta q/\sqrt{2}$
$N(1,1)$	8	$\frac{1}{2}$	1	$(\beta-\delta)q/2$
$\Lambda(1,2)$	1	$\frac{1}{2}$	2	$-\beta q/2$
$\Delta(3,2)$	10	$\frac{3}{2}$	2	$\delta q/\sqrt{2}$
$\Delta(1,2)$	10	$\frac{1}{2}$	2	$\delta q/2$
$N(3,2)$	8	$\frac{3}{2}$	2	$\delta q/2$
$N(1,2)$	8	$\frac{1}{2}$	2	$(\beta+\delta)q/2$
$N'(1,2)$	8	$\frac{1}{2}$	2	$(\beta-\delta)q/\sqrt{8}$

wave function that corresponds to an S-wave diquark combined with the third quark. Other components of the wave function do not contribute to the calculation of the decay width, since the decaying diquonia are assumed to consist of S-wave diquarks and antidiquarks only, and these are assumed to be spectators in the decay process.

The lowest-lying baryon states that satisfy this condition are given in Table II. The component of the flavor-spin wave function containing the S-wave diquark is given in column 4. Thus, for example, the flavor-spin wave function of the members of the lowest-lying octet may be written as $(\beta+\delta)q/\sqrt{2}$. Note that the symbol used for the multiplet is that of the nonstrange members (where possible), and the numbers in parentheses are twice the total spin of each baryon in the multiplet, and the orbital angular momentum of the baryon, respectively.

III. MASS CALCULATIONS

Let us now focus on the problem of calculating the masses of the diquonia states of the previous section or, more precisely, estimating the positions of the centroids of the $SU(3)_f$ multiplets. To do this, we use an additive potential similar to that used by Bhaduri,[20] consisting of a linear confining term with Coulomb and short-range spin-spin terms. The form used is

$$V=-\tfrac{3}{16}\sum_{i,j}\boldsymbol{\lambda}_i\cdot\boldsymbol{\lambda}_j[-K/r_{ij}+\lambda r_{ij}-D+\boldsymbol{\sigma}_i\cdot\boldsymbol{\sigma}_j Ce^{-\mu r_{ij}}/(m_i m_j r_{ij})] , \qquad (3.1)$$

where $K=0.52$, $\lambda=0.186\ \mathrm{GeV}^2$, $D=0.914$ GeV, $C=0.375\ \mathrm{GeV}^2$, and $\mu=0.434\ \mathrm{GeV}^{-1}$ in Ref. 20.

The differences between the parameters we use and those above arise because we choose to treat the baryons and diquonia in the diquark approximation: the baryon consists of an S-wave diquark and a quark with relative orbital angular momentum L_B, and the diquonium consists of an S-wave diquark and an S-wave antidiquark with orbital angular momentum L_D. While the dynamical motivation for this approximation may be poor, especially in the case of baryons and diquonia with low orbital angular momentum, it is a convenient one that allows a simple treatment of both the $B\bar{B}$ decays of the diquonia, and calculation of their masses.

We emphasize here that the full symmetrized baryon wave function contains components that correspond to a diquark with nonzero orbital angular momentum. These do not contribute to the decay in the $^3P_0$ model, and are thus ignored. We therefore concentrate on those components of the wave function that correspond to S-wave diquarks only, and use these in calculating the masses of the baryons.

For the baryon wave functions, the spatial parts that are of interest to us here are of the form

$$\Psi_B(\mathbf{r}_1,\mathbf{r}_2,\mathbf{r}_3)=\{8/[\alpha^{2L_B+3}\beta^3\Gamma(L_B+\tfrac{3}{2})\sqrt{\pi}]\}^{1/2}\times e^{-v^2/2\alpha^2}e^{-w^2/2\beta^2}\mathcal{Y}_{L_B}(\mathbf{v})\mathcal{Y}_0(\mathbf{w}) , \qquad (3.2)$$

where the $\mathcal{Y}_L$ are the solid harmonics,

$$\mathbf{v}=(\mathbf{r}_1+\mathbf{r}_2-2\mathbf{r}_3)/2$$

is the quark-diquark separation,

$$\mathbf{w}=\mathbf{r}_1-\mathbf{r}_2 \tag{3.3}$$

is the interquark separation in the diquark, and $\alpha^2=3\beta^2/4$. This relationship between α and β ensures the overall symmetry properties of the wave function. The radial parts of the diquonia wave functions used have a similar form:

$$\Psi_D(\mathbf{r}_1,\mathbf{r}_2,\mathbf{r}_3,\mathbf{r}_4)=\{32/[\gamma^{2L_D+3}\delta^6\Gamma(L_D+\tfrac{3}{2})\pi]\}^{1/2}e^{-x^2/2\gamma^2}e^{-(y^2+z^2)/2\delta^2}\mathcal{Y}_{L_D}(\mathbf{x})\mathcal{Y}_0(\mathbf{y})\mathcal{Y}_0(\mathbf{z})\ , \tag{3.4}$$

where $\mathbf{r}_1$ and $\mathbf{r}_2$ are quark coordinates, $\mathbf{r}_3$ and $\mathbf{r}_4$ are antiquark coordinates,

$$\mathbf{x}=(\mathbf{r}_1+\mathbf{r}_2-\mathbf{r}_3-\mathbf{r}_4)/2,\quad \mathbf{y}=\mathbf{r}_1-\mathbf{r}_2$$

is the internal coordinate of the diquark, and

$$\mathbf{z}=\mathbf{r}_3-\mathbf{r}_4 \tag{3.5}$$

is that of the antiquark. The momentum representations of these wave functions have similar forms.

The diquark approximation is implemented in evaluating

$$\langle E\rangle=\left\langle V+\sum[p_i^2/(2m_i)+m_i]\right\rangle$$

by expanding the various position-dependent terms in the potential as series in the coordinate w in the case of the baryons, for example, and truncating at quadratic terms. The assumption of S-wave diquarks simplifies much of the work. For example, in the baryon, $1/r_{13}$ is expanded in terms of spherical harmonics is

$$1/r_{13}=|\mathbf{w}/2+\mathbf{v}|^{-1}$$
$$=\frac{1}{v}\sum_{n=0}^{\infty}\left[\frac{w}{2v}\right]^n\sum_{m=-n}^{n}Y_n^m(\hat{\mathbf{w}})Y_n^{m*}(\hat{\mathbf{v}})\ . \tag{3.6}$$

With S-wave diquarks, only the first term of the series contributes when expectation values are calculated, so that

$$\langle 1/r_{13}\rangle=\langle 1/v\rangle\ . \tag{3.7}$$

Inherent in the above expansion is the assumption that the internal dimension of the diquark is smaller than the interdiquark separation in the diquonia, or the distance between the diquark and the third quark in the baryon. This is the essence of the diquark approximation. However, this assumption is not necessarily a good one for low L_D.

When the expectation value of the total energy is minimized as a function of the Gaussian parameters of the wave functions, some of the masses obtained are very small. This occurs only for states with $L=0$, 1, and is a consequence of the short-range spin-spin term of the potential. Note, however, that for the baryons, the minimization procedure gives masses that are in reasonable agreement with a more complete theoretical calculation[21] for baryons with $L_B=2$.

To escape the problem of very light states, we fit the theoretical baryon masses to the physical values, allowing not just the Gaussian parameters of the wave functions to vary, but also the parameters of the spin-spin term, μ and C. The results of this fit are shown in Table III. Note that the subscript added to the notation for many of the octets refers to the spin of the diquark in the baryon. Note, too, that the masses in column 4 are the averages of the masses of all the multiplets with given L_B and S_B but different J_B. The results in Table III are obtained with μ unchanged from its original value, but with a somewhat smaller value for C of 0.175.

Note that in fitting the baryon masses we do not make the two components of each octet degenerate, but instead we fit the average mass of the two components of each octet to the "experimental" mass. In addition, when we evaluate the decay widths, we use only the "experimental" masses of the baryons: the above exercise in the case of the baryons serves mainly to give us a new value for C, and to provide values of the Gaussian parameters α and β used in calculating the decay widths.

The new value of C is used in obtaining diquonia masses. Here, however, we are unable to perform a fit *per se*, so we must again minimize the total energy as a func-

TABLE III. Baryon masses and Gaussian widths. Only the width α is shown, and $\alpha^2=3\beta^2/4$. The theoretical masses obtained in this work are shown in column 2, while column 4 shows the experimental masses. A single asterisk indicates that the masses of some members of the multiplet are taken from experiment while the others are taken from the calculation of Ref. 21. Two asterisks indicate that all the masses in the multiplet are from Ref. 21. For the multiplets $N_1(1,L)$ and $N_0(1,L)$, the masses shown in column 2 are the calculated average masses of these two components of the octets.

State	Fitted mass (GeV)	α(GeV^{-1})	Mass
$N_0(1,0)$	1.150	3.45	1.150
$N_1(1,0)$		3.80	
$\Delta(3,0)$	1.482	3.29	1.381
$\Lambda(1,1)$	1.517	2.91	1.480
$\Delta(1,1)$	1.760	3.58	1.760*
$N(3,1)$	1.790	3.66	1.790*
$N_1(1,1)$	1.665	3.40	1.665*
$N_0(1,1)$		3.54	
$\Lambda(1,2)$	1.912	3.12	1.850**
$\Delta(3,2)$	2.126	3.40	2.115**
$\Delta(1,2)$	2.125	3.40	2.085**
$N(3,2)$	2.126	3.41	2.105**
$N_1(1,2)$	2.018	3.40	1.915**
$N_0(1,2)$		3.12	
$N_1'(1,2)$	2.100	2.74	2.100**
$N_0'(1,2)$		3.56	

tion of the Gaussian widths. In doing this, we again run into very light states. We remedy this by taking the masses of the diquonia with large L_D ($2 \le L_D \le 8$), where the diquark approximation is expected to have a greater degree of validity, performing a Regge fit to these masses, and using this Regge fit to extract the masses of the states at lower L_D ($L_D < 2$).

The results from this procedure are most conveniently presented as Regge trajectories, which are

$$\begin{aligned} M_A^2 &= 1.29L_D + 1.55 \ , \\ M_B^2 &= 1.44L_D + 2.56 \ , \\ M_{C_0}^2 &= 1.65L_D + 2.68 \ , \\ M_{C_1}^2 &= 1.61L_D + 2.97 \ , \\ M_{C_2}^2 &= 1.43L_D + 4.06 \ . \end{aligned} \tag{3.8}$$

In the above, C_0, C_1, and C_2 refer to the states $C(m,0,L_D)$, $C(m,1,L_D)$, and $C(m,2,L_D)$, respectively. Note that the different multiplets of a given kind of diquonium, A, B, or C, are degenerate.

For the convenience of the reader who may wish to reproduce any part of our calculation, we have also parametrized the Gaussian widths of the diquonia as functions of L_D (the Gaussian widths of the baryons are given in Table III). For the widths δ (not to be confused with the notation for the sextet diquark), a form quadratic in L_D is sufficient, and the results are

$$\begin{aligned} \delta_A &= -4.77\times 10^{-3}L_D^2 + 6.12\times 10^{-2}L_D + 2.71 \ , \\ \delta_B &= -7.14\times 10^{-3}L_D^2 + 9.28\times 10^{-2}L_D + 3.22 \ , \\ \delta_{C_0} &= -3.47\times 10^{-2}L_D^2 + 0.40L_D + 2.98 \ , \\ \delta_{C_1} &= -2.65\times 10^{-2}L_D^2 + 0.31L_D + 3.22 \ , \\ \delta_{C_2} &= -7.77\times 10^{-3}L_D^2 + 0.11L_D + 3.66 \ . \end{aligned} \tag{3.9}$$

In the case of the widths γ, it is found that a quartic form gives a satisfactory fit:

$$\begin{aligned} \gamma_A &= -3.76\times 10^{-4}L_D^4 + 8.02\times 10^{-3}L_D^3 - 6.60\times 10^{-2}L_D^2 + 0.33L_D + 2.52 \ , \\ \gamma_B &= -2.40\times 10^{-4}L_D^4 + 5.29\times 10^{-3}L_D^3 - 4.65\times 10^{-2}L_D^2 + 0.27L_D + 2.62 \ , \\ \gamma_{C_0} &= 3.57\times 10^{-3}L_D^4 - 6.61\times 10^{-2}L_D^3 + 0.40L_D^2 - 0.73L_D + 3.06 \ , \\ \gamma_{C_1} &= 2.81\times 10^{-3}L_D^4 - 4.95\times 10^{-2}L_D^3 + 0.27L_D^2 - 0.37L_D + 2.79 \ , \\ \gamma_{C_2} &= 1.06\times 10^{-3}L_D^4 - 2.09\times 10^{-2}L_D^3 + 0.14L_D^2 - 0.32L_D + 3.37 \ . \end{aligned} \tag{3.10}$$

Note that both parametrizations give Gaussian widths that are within 1% of those used in calculating the decay widths. However, we emphasize that these fits have only been verified up to $L_D = 8$, and are simply a convenient way of presenting the multitude of Gaussian parameters that we have used: profound physical content is neither implied nor expected.

Some comments on the masses obtained by the method outlined above are in order. First, it is worth pointing out that most of the masses obtained with the equations above are similar to those obtained using Jaffe's corresponding equations for SU(2) (Ref. 12), and adding 300 MeV to estimate the masses for the SU(3) multiplets. They are also similar to the masses estimated in Ref. 17, where strange quarks have been included. In the case of the A states, however, we find masses that are smaller. We expect that the masses of our A states are consistently too small by about 150 MeV. Note that the masses that we have obtained take into account the important spin-spin interaction, while those of Refs. 12 and 17 do not.

Let us also point out that although the diquark approximation is questionable, and the minimization procedure leads to problems with very light states, perhaps because of the diquark approximation, the results for $L_B = 2$ in the baryon sector are reasonable. We therefore expect that the masses obtained for the diquonia with $L_D \ge 2$ are trustworthy and, subsequently, that the masses obtained for $L_D = 0,1$, based on the Regge fit to states with higher L_D, are also dependable.

We find it instructive, however, to allow for some uncertainty in the masses of the states. Thus, when we discuss the $B\bar{B}$ widths of the diquonia, we present the widths corresponding not only to the masses of the equations above, but also those corresponding to masses 300 MeV heavier and lighter. This may be especially useful in the case of the A states. With this presentation, we hope to obtain some idea of the range of the $B\bar{B}$ widths of the diquonia states discussed herein.

Finally we end this section with a note of irony. After carrying out the procedure outlined above to obtain the masses of the diquonia with $L_D = 0,1$, we find that most of these states are below the lowest $B\bar{B}$ threshold, and so may hold little interest for $N\bar{N}$ experiments, for example. For $L_D = 0$, all of these states, with the exception of the $C(m,2,0)$, remain below threshold even with the addition of 300 MeV. However, some of them lie very close to threshold, and could actually show up as narrow structures in some $B\bar{B}$ channels. In the case of $L_D = 1$, most of the states become heavier than the lowest $B\bar{B}$ threshold when the extra 300 MeV is added to their masses. The exceptions here are the A states, which are still well below threshold, and the $C(m,2,1)$, which are above threshold without the extra 300 MeV. However, apart from the A states, we are being overly generous in allowing such a large variation in the masses. In addition, we point out that even though most of the diquonia with $L_D=0,1$ are below the lowest $B\bar{B}$ threshold, they are all well above $M\bar{M}$ thresholds (M denotes meson). None of these states are therefore bound, and will fall apart "easily," modulo angular momentum selection rules, into pairs of mesons.

IV. DECAY MODEL

The model used for description of the decays is the version of the ${}^3P_0$ model popularized by LeYaouanc *et al.*,[22] illustrated in Fig. 2. This model has been discussed several times in the literature,[23] so we choose not to include too many details of this model here. We will, however, reproduce the general formula for the partial widths obtained using this model. The partial width for the decay of diquonium A of mass M_A into baryon B and antibaryon C is given by

$$\Gamma=\frac{24\pi\lambda^2 E_B E_C k_B}{M_A}\frac{(2S_B+1)(2S_C+1)}{2L_A+1}\mathcal{R}\mathcal{I}\mathcal{N}\sum_{S_{BC}}(2S_{BC}+1)\begin{Bmatrix} s_d & \frac{1}{2} & S_B \\ s_{\bar{d}} & \frac{1}{2} & S_C \\ S_A & 1 & S_{BC}\end{Bmatrix}^2 \sum_{l,L,L_{BC}}(2L+1)\mathcal{E}^2(l,L,L_{BC},k_B)\,, \tag{4.1}$$

where we have summed over allowed values of J_B and J_C.

In the above, k_B is the baryon momentum in the center-of-mass frame of the baryon-antibaryon pair, $E_B=(k_B^2+M_B^2)^{1/2}$ with a similar definition for E_C. $S_{B,C}$ are the spins of baryon B,C, respectively, L_A is the orbital angular momentum in the diquonium, s_d is the spin of the diquark, and $s_{\bar{d}}$ is that of the antidiquark. The factor $\mathcal{I}$ contains all information on the overlaps between the flavor parts of the diquonia wave functions and those of the baryons, and $\mathcal{R}$ is the square of the overlap between the spatial parts of the wave functions of the diquarks in the diquonium and in the baryons. It is given by

$$\mathcal{R}=\left[\frac{4\beta_B\beta_C\delta^2}{(\delta^2+\beta_B^2)(\delta^2+\beta_C^2)}\right]^3 .$$

$\mathcal{N}$ contains the rest of the spatial wave-function normalization factors, and is given by

$$\mathcal{N}=\frac{8\gamma_A^{2L_A+3}\alpha_B^{2L_B+3}\alpha_C^{2L_C+3}}{\Gamma(L_A+\frac{3}{2})\Gamma(L_B+\frac{3}{2})\Gamma(L_C+\frac{3}{2})}\,,$$

$l=L_A\pm1$, and the term $\mathcal{E}$ is given by

$$\begin{aligned}\mathcal{E}(l,L,L_{BC},k_B)=&\frac{(-1)^{L_{BC}}3e^{-(F^2k_B^2)}}{4\sqrt{\pi}G^{L_A+L_B+L_C+4}}(2L_A+1)(2L_B+1)(2L_C+1)\\ &\times\sum_{l_1,l_2,l_3,l_4}(Gk_B)^{l_1+l_2+l_3+l_4}\left[\begin{bmatrix}2L_B\\2l_1\end{bmatrix}\begin{bmatrix}2L_C\\2l_2\end{bmatrix}\begin{bmatrix}2\\2l_3\end{bmatrix}\begin{bmatrix}2L_A\\2l_4\end{bmatrix}\right]^{1/2}\\ &\times\sqrt{(2l+1)(2L_{BC}+1)}(x-\tfrac{2}{3})^{l_1+l_2}(x-1)^{l_3}x^{l_4}\\ &\times\Gamma\left[\frac{L_A+L_B+L_C-l_1-l_2-l_3-l_4}{2}+2\right]\\ &\times\sum_{l_{12},l_5,l_6}(-1)^{l_6}(2l_{12}+1)(2l_5+1)(2l_6+1)\begin{Bmatrix}l_1 & L_B-l_1 & L_B\\ l_2 & L_C-l_2 & L_C\\ l_{12} & l_6 & L_{BC}\end{Bmatrix}\\ &\times\begin{Bmatrix}l_3 & 1-l_3 & 1\\ l_4 & L_A-l_4 & L_A\\ l_5 & l_6 & L\end{Bmatrix}\begin{Bmatrix}l & l_{12} & l_5\\ l_6 & L & L_{BC}\end{Bmatrix}\begin{pmatrix}l_1 & l_2 & l_{12}\\0&0&0\end{pmatrix}\begin{pmatrix}l & l_{12} & l_5\\0&0&0\end{pmatrix}\begin{pmatrix}l_3 & l_4 & l_5\\0&0&0\end{pmatrix}\\ &\times\begin{pmatrix}L_B-l_1 & L_C-l_2 & l_6\\0&0&0\end{pmatrix}\begin{pmatrix}1-l_3 & L_A-l_4 & l_6\\0&0&0\end{pmatrix}.\end{aligned}\tag{4.2}$$

TABLE IV. Baryon-antibaryon exclusive partial widths of diquonia states. Masses are in GeV and widths are in MeV. Only widths greater than 1 MeV are shown. Column 4 shows the widths corresponding to the masses calculated in Sec. IV, column 5 shows the widths corresponding to masses 300 MeV heavier, while column 6 shows the widths corresponding to masses 300 MeV lighter. Note that for most states, the widths in column 6 are zero. This is because states 300 MeV lighter are below the various thresholds, so that the decay channels are inaccessible.

State	Mass	$B\bar{B}$ channel	$\Gamma(M)$	$\Gamma(M+300)$	$\Gamma(M-300)$
$C(1,2,0)$	2.10	$N(1,0)N(1,0)$	0	423	0
$C(8,2,0)$	2.10	$N(1,0)N(1,0)$	0	317	0
$C(27,2,0)$	2.10	$N(1,0)N(1,0)$	0	141	0
$C(1,0,1)$	2.08	$N(1,0)N(1,0)$	0	14	0
$C(8,0,1)$	2.08	$N(1,0)N(1,0)$	0	10	0
$C(27,0,1)$	2.08	$N(1,0)N(1,0)$	0	5	0
$C(1,1,1)$	2.14	$N(1,0)N(1,0)$	0	38	0
$C(8,1,1)$	2.14	$N(1,0)N(1,0)$	0	28	0
$C(27,1,1)$	2.14	$N(1,0)N(1,0)$	0	13	0
$C(1,2,1)$	2.32	$N(1,0)N(1,0)$	84	453	0
$C(8,2,1)$	2.32	$N(1,0)N(1,0)$	63	340	0
		$N(1,0)\Delta(3,0)^*$	0	85	0
$C(27,2,1)$	2.32	$N(1,0)N(1,0)$	28	151	0
		$N(1,0)\Delta(3,0)^*$	0	227	0
$A(1,2)$	2.00	$N(1,0)N(1,0)$	0	3	0
$A(8,2)$	2.00	$N(1,0)N(1,0)$	0	2	0
$B(8,2)$	2.31	$N(1,0)N(1,0)$	5	212	0
$B(10,2)$	2.31	$N(1,0)N(1,0)$	6	282	0
$C(1,0,2)$	2.44	$N(1,0)N(1,0)$	11	63	0
$C(8,0,2)$	2.44	$N(1,0)N(1,0)$	8	47	0
		$N(1,0)\Delta(3,0)^*$	0	78	0
$C(27,0,2)$	2.44	$N(1,0)N(1,0)$	4	21	0
		$N(1,0)\Delta(3,0)^*$	0	207	0
$C(1,1,2)$	2.49	$N(1,0)N(1,0)$	33	180	0
		$\Delta(3,0)\Delta(3,0)$	0	243	0
$C(8,1,2)$	2.49	$N(1,0)N(1,0)$	25	136	0
		$N(1,0)\Delta(3,0)^*$	0	102	0
		$\Delta(3,0)\Delta(3,0)$	0	195	0
$C(27,1,2)$	2.49	$N(1,0)N(1,0)$	11	60	0
		$N(1,0)\Delta(3,0)^*$	0	273	0
		$\Delta(3,0)\Delta(3,0)$	0	114	0
$C(1,2,2)$	2.58	$N(1,0)N(1,0)$	134	554	0
		$\Delta(3,0)\Delta(3,0)$	0	1348	0
$C(8,2,2)$	2.58	$N(1,0)N(1,0)$	100	415	0
		$N(1,0)\Delta(3,0)^*$	18	142	0
		$\Delta(3,0)\Delta(3,0)$	0	1078	0
$C(27,2,2)$	2.58	$N(1,0)N(1,0)$	45	185	0
		$N(1,0)\Delta(3,0)^*$	47	379	0
		$\Delta(3,0)\Delta(3,0)$	0	629	0
$A(1,3)$	2.29	$N(1,0)N(1,0)$	0	504	0
$A(8,3)$	2.29	$N(1,0)N(1,0)$	0	441	0
$B(8,3)$	2.60	$N(1,0)N(1,0)$	82	380	0
$B(10,3)$	2.60	$N(1,0)N(1,0)$	109	507	0
$C(1,0,3)$	2.77	$N(1,0)N(1,0)$	34	120	5
		$\Delta(3,0)\Delta(3,0)$	0	952	0
$C(8,0,3)$	2.77	$N(1,0)N(1,0)$	26	90	4
		$N(1,0)\Delta(3,0)^*$	41	215	0
		$\Delta(3,0)\Delta(3,0)$	0	762	0
$C(27,0,3)$	2.77	$N(1,0)N(1,0)$	11	40	1
		$N(1,0)\Delta(3,0)^*$	108	572	0
		$\Delta(3,0)\Delta(3,0)$	0	444	0
$C(1,1,3)$	2.80	$N(1,0)N(1,0)$	83	260	13
		$\Delta(3,0)\Delta(3,0)$	17	1256	0
$C(8,1,3)$	2.80	$N(1,0)N(1,0)$	63	196	10
		$N(1,0)\Delta(3,0)^*$	45	220	0

TABLE IV. (*Continued*).

State	Mass	$B\bar{B}$ channel	$\Gamma(M)$	$\Gamma(M+300)$	$\Gamma(M-300)$
		$\Delta(3,0)\Delta(3,0)$	14	1005	0
$C(27,1,3)$	2.80	$N(1,0)N(1,0)$	28	87	5
		$N(1,0)\Delta(3,0)^*$	119	586	0
		$\Delta(3,0)\Delta(3,0)$	8	586	0
$C(1,2,3)$	2.84	$N(1,0)N(1,0)$	228	572	41
		$N(1,0)N(1,2)^*$	0	9	0
		$\Delta(3,0)\Delta(3,0)$	151	2004	0
$C(8,2,3)$	2.84	$N(1,0)N(1,0)$	171	429	31
		$N(1,0)\Delta(3,0)^*$	46	198	0
		$N(1,0)N(1,2)^*$	0	6	0
		$\Delta(3,0)\Delta(3,0)$	121	1603	0
$C(27,2,3)$	2.84	$N(1,0)N(1,0)$	76	191	14
		$N(1,0)\Delta(3,0)^*$	122	528	0
		$N(1,0)N(1,2)^*$	0	3	0
		$\Delta(3,0)\Delta(3,0)$	71	936	0
$A(1,4)$	2.56	$N(1,0)N(1,0)$	129	798	0
$A(8,4)$	2.56	$N(1,0)N(1,0)$	113	698	0
$B(8,4)$	2.86	$N(1,0)N(1,0)$	140	463	21
		$N(1,0)\Lambda(1,2)$	0	4	0
		$N(1,0)N(1,2)$	0	3	0
		$N(1,2)N(1,0)$	0	4	0
$B(10,4)$	2.86	$N(1,0)N(1,0)$	187	618	28
		$N(1,0)N(1,2)$	0	4	0
		$N(1,2)N(1,0)$	0	6	0
$C(1,0,4)$	3.05	$N(1,0)N(1,0)$	60	128	13
		$N(1,0)N(3,2)^*$	0	5	0
		$N(1,0)N(1,2)^*$	0	8	0
		$\Delta(3,0)\Delta(3,0)$	312	1642	0
$C(8,0,4)$	3.05	$N(1,0)N(1,0)$	45	96	10
		$N(1,0)\Delta(3,0)^*$	87	303	11
		$N(1,0)\Delta(3,2)^*$	0	2	0
		$N(1,0)N(3,2)^*$	0	4	0
		$N(1,0)N(1,2)^*$	0	6	0
		$\Delta(3,0)\Delta(3,0)$	249	1313	0
$C(27,0,4)$	3.05	$N(1,0)N(1,0)$	20	43	4
		$N(1,0)\Delta(3,0)^*$	231	804	29
		$N(1,0)\Delta(3,2)^*$	0	5	0
		$N(1,0)N(3,2)^*$	0	2	0
		$N(1,0)N(1,2)^*$	0	3	0
		$\Delta(3,0)\Delta(3,0)$	146	767	0
		$\Delta(3,0)N(1,2)^*$	0	2	0
$C(1,1,4)$	3.07	$N(1,0)N(1,0)$	130	259	30
		$N(1,0)N(3,2)^*$	0	7	0
		$N(1,0)N(1,2)^*$	0	18	0
		$\Delta(3,0)\Delta(3,0)$	391	1947	0
$C(8,1,4)$	3.07	$N(1,0)N(1,0)$	97	193	22
		$N(1,0)\Delta(3,0)^*$	84	277	12
		$N(1,0)\Delta(3,2)^*$	0	3	0
		$N(1,0)N(3,2)^*$	0	5	0
		$N(1,0)N(1,2)^*$	0	14	0
		$\Delta(3,0)\Delta(3,0)$	313	1558	0
		$\Delta(3,0)N(1,2)^*$	0	2	0
$C(27,1,4)$	3.07	$N(1,0)N(1,0)$	43	86	10
		$N(1,0)\Delta(3,0)^*$	223	738	30
		$N(1,0)\Delta(3,2)^*$	0	7	0
		$N(1,0)\Delta(1,2)^*$	0	2	0
		$N(1,0)N(3,2)^*$	0	2	0
		$N(1,0)N(1,2)^*$	0	6	0
		$\Delta(3,0)\Delta(3,0)$	183	909	0

TABLE IV. (*Continued*).

State	Mass	$B\bar{B}$ channel	$\Gamma(M)$	$\Gamma(M+300)$	$\Gamma(M-300)$
		$\Delta(3,0)N(1,2)^*$	0	4	0
$C(1,2,4)$	3.09	$N(1,0)N(1,0)$	289	511	74
		$N(1,0)N(3,2)^*$	0	10	0
		$N(1,0)N(1,2)^*$	0	45	0
		$N(1,0)N'(1,2)^*$	0	2	0
		$\Delta(3,0)\Delta(3,0)$	577	2610	1
		$N(1,1)N(1,1)$	0	2	0
$C(8,2,4)$	3.09	$N(1,0)N(1,0)$	217	383	56
		$N(1,0)\Delta(3,0)^*$	71	210	11
		$N(1,0)\Delta(3,2)^*$	0	4	0
		$N(1,0)\Delta(1,2)^*$	0	3	0
		$N(1,0)N(3,2)^*$	0	7	0
		$N(1,0)N(1,2)^*$	0	34	0
		$N(1,0)N'(1,2)^*$	0	1	0
		$\Delta(3,0)\Delta(3,0)$	461	2088	0
		$\Delta(3,0)N(1,2)^*$	0	3	0
		$N(1,1)N(1,1)$	0	1	0
$C(27,2,4)$	3.09	$N(1,0)N(1,0)$	96	170	24
		$N(1,0)\Delta(3,0)^*$	190	562	30
		$N(1,0)\Delta(3,2)^*$	0	11	0
		$N(1,0)\Delta(1,2)^*$	0	7	0
		$N(1,0)N(3,2)^*$	0	3	0
		$N(1,0)N(1,2)^*$	0	15	0
		$\Delta(3,0)\Delta(3,0)$	269	1218	0
		$\Delta(3,0)N(1,2)^*$	0	3	0

Note that $\mathcal{E}$ is also a function of L_A, L_B, and L_C. The { } are the usual 6- or 9-J symbols, while the ( ) are 3-J symbols. This nontrivial form simplifies when either of L_A, L_B, or L_C is zero, and if $L_B=L_C=0$, the form we use agrees with that of Ref. 14 and the corrected version of Ref. 15 (Ref. 15, erratum). In order to compare the form given above with those of Refs. 14 and 15, however, the reader must note that we have explicitly removed the spectator overlap (the factor $\mathcal{R}$), as well as the rest of the spatial wave-function normalization ($\mathcal{N}$) from the spatial amplitude, while in Refs. 14 and 15, they are included. λ is the pair-creation constant, for which we take a value of 3.4. This value is consistent with the values used elsewhere in calculations of the widths of various baryon and meson decay processes.[22] By using this value for λ, we are assuming that there is no significant change in the physics of the decay process in going from the two- and three-quark sectors to the four-quark sector.

In Eq. (4.2),

$$x=\frac{2(\alpha_B^2+\alpha_C^2)}{3(\gamma_A^2+\alpha_B^2+\alpha_C^2)}\ ,$$
$$G^2=(\gamma_A^2+\alpha_B^2+\alpha_C^2)/2\ , \tag{4.3}$$
$$F^2=[\gamma_A^2x^2+(\alpha_B^2+\alpha_C^2)(x-\tfrac{2}{3})^2]/2\ .$$

V. TOTAL AND PARTIAL WIDTHS

Figures 3–7 present the total $B\bar{B}$ decay widths for some of the diquonia states discussed in Sec. II, for $L_D=0$ to 4. The reader is reminded that at $L_D=0$, the $C(m,2,0)$ are the only ones that are above the lowest $B\bar{B}$ threshold when an extra 300 MeV is added to the masses. In these figures the widths are shown as functions of the mass of the decaying diquonium. The segment of each curve shown corresponds to the mass range $M-300$ MeV to $M+300$ MeV, where M is the mass given by Eq. (3.8). The widths for the states not shown exhibit behavior similar to one of the graphs shown: the states $A(8,L_D)$ are similar to the $A(1,L_D)$, the $B(10,L_D)$ and $B(\overline{10},L_D)$ are similar to the $B(8,L_D)$, and the $C(8,S,L_D)$ and

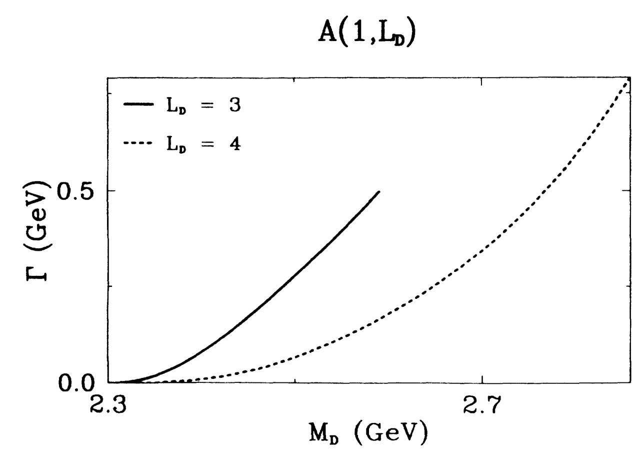

FIG. 3. Total decay widths of the states $A(1,L_D)$ as a function of the mass of the decaying state.

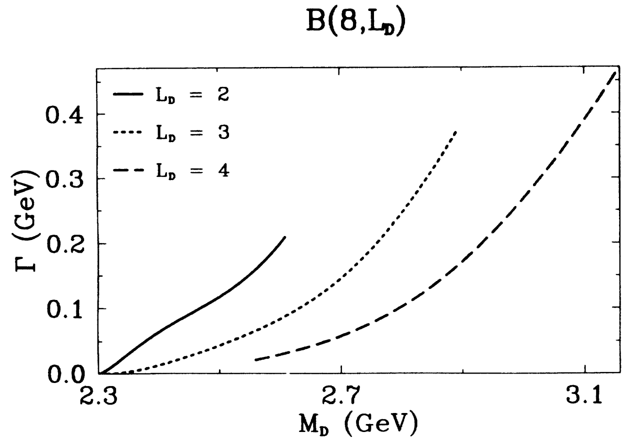

FIG. 4. Total decay widths of the states $B(8,L_D)$ as a function of the mass of the decaying state.

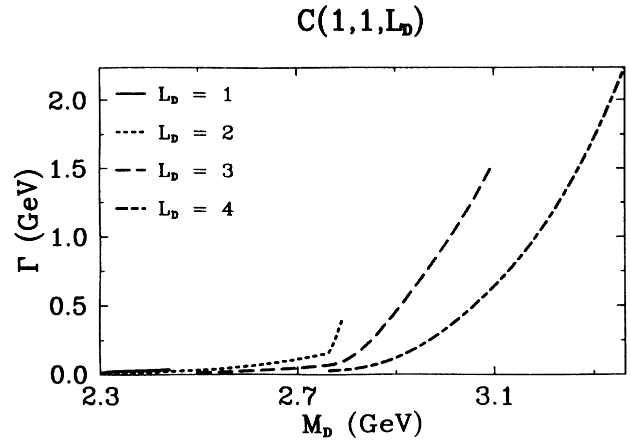

FIG. 6. Total decay widths of the states $C(1,1,L_D)$ as a function of the mass of the decaying state.

$C(27,S,L_D)$ are similar to the $C(1,S,L_D)$.

It is clear from these figures that many of the total $B\bar{B}$ widths are very large and that the corresponding states would perhaps not be very interesting. Some of the states have narrow widths which remain narrow even when the mass is increased by 300 MeV. In contrast, some states have broad widths that become immense when the mass is increased by the same amount. Such states are almost certainly without interest, unless their masses are significantly smaller than those calculated here.

The A states have small total widths corresponding to the masses calculated, but these widths increase rapidly when the mass of the state is increased. The reader is reminded that our A states are probably too light, so that the widths corresponding to more realistic masses are considerably larger. In contrast with the A states, the total widths of the B states increase less rapidly when the mass of the state is increased. The C states are perhaps the most intriguing in that their total widths remain quite small for a significant part of the mass range investigated, then become quite large as the mass is increased beyond about 2.8 GeV. This is the effect of a new threshold, and will become clearer when we discuss the partial widths. The effect of this threshold is seen most clearly in the curves for the states $C(1,1,2)$ and $C(1,2,2)$ in Figs. 6 and 7, respectively.

We must point out that the widths presented are sensitive not only to the mass, as is evident from the figures, but also to the Gaussian parameters used in the wave functions. This is especially so in the case of states with large widths. For instance, if we set $\alpha=\beta=\gamma=\delta=3.5$ GeV^{-1}, the total width of the state $C(27,2,4)$ is of the order of 800 MeV. If, however, we choose $\alpha=\beta=\gamma=\delta=2.5$ GeV^{-1}, the width of this state decreases by a factor of about 4. The partial widths show similar dependence on the Gaussian parameters.

Despite the fact that the total and partial widths are sensitive to the values chosen for the Gaussian parameters, we expect that the widths we have calculated are reliable. This is because the Gaussian parameters for the baryon wave functions are expected to be close to (or greater than) 3.0 GeV^{-1}, and those for the diquonia should be similar. Thus, while estimates of the widths

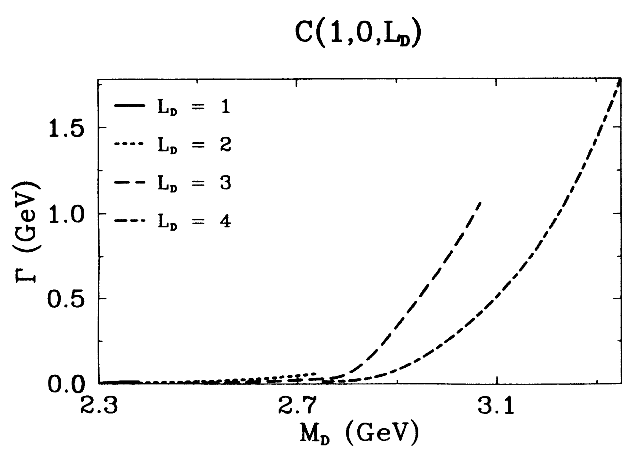

FIG. 5. Total decay widths of the states $C(1,0,L_D)$ as a function of the mass of the decaying state.

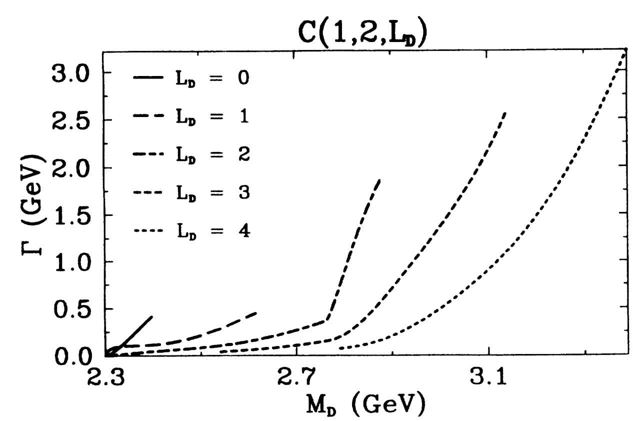

FIG. 7. Total decay widths of the states $C(1,2,L_D)$ as a function of the mass of the decaying state.

with all Gaussian parameters set to 2.5 GeV^{-1} may represent a lower limit, the widths we present should correspond more closely with physical reality.

While the total $B\bar{B}$ widths give some idea of which diquonia states may be of experimental interest, it is of some interest to look at the partial widths into specific $B\bar{B}$ channels. These are presented in Table IV, where only the partial widths greater than 1 MeV are shown. Note that we show the widths for only two of the B states: as a result of charge-conjugation symmetry, the states with diquark content $\delta\bar{\beta}$ have identical partial widths for decay into the baryon channel $B\bar{C}$ as states with diquark content $\beta\bar{\delta}$ decaying into the channel $C\bar{B}$.

We also point out that for the states A and C we have not shown all of the partial widths. For these diquonia, the width for the decay into a baryon channel $B\bar{C}$ (B and C are different) is the same as the width for decay into the charge-conjugate channel $C\bar{B}$. The partial widths missing from the table may thus be easily deduced. To aid the reader, we indicate by an asterisk all channels for which the charge-conjugate partner is missing.

In the table, column 4 shows the widths corresponding to the masses calculated in Sec. IV, column 5 shows the widths corresponding to masses 300 MeV heavier, while column 6 shows the widths corresponding to masses 300 MeV lighter. Note that for most states, the widths in column 6 are zero. This is because states 300 MeV lighter are below the various thresholds, so that the decay channels are inaccessible.

On examining the entries in Table IV, it immediately becomes clear that by far the largest partial widths correspond to decay into the channel $\Delta(3,0)\Delta(3,0)$ (we use the same symbols, without bars, to refer to the antibaryon multiplets, but we always refer to the baryon first). This is in agreement with the work of Ono and Furui[11] who find large widths for the decay of diquonia with strangeness into the $\Delta\bar{\Sigma}^*$ channel. This suggests that diquonia which decay into $\Delta\bar{\Delta}$ channels will be difficult to find since they will show up as very broad structures. On the other hand, almost all states, even those with ample phase space, have smaller widths into the decay channel $N(1,0)N(1,0)$. These widths range from tens of MeV to a few hundreds.

As with the total widths, the partial widths are sensitive to both the mass of the decaying state and the Gaussian parameters used. Note, however, that this latter sensitivity is greatest in the $\Delta(3,0)\Delta(3,0)$ channel. The partial widths into the $N(1,0)N(1,0)$ channel are also sensitive to α, β, γ, and δ, but they generally remain in the physically interesting range when either these parameters or the masses of the decaying states are changed.

Let us take a more systematic look at the diquonia decays. To begin, we note that the widths for the states $A(1,L_D)$ and $A(8,L_D)$ are small, but increase quite rapidly as their masses and consequently, the available phase space is increased. This is in overall agreement with the results of Refs. 14 and 15 where it was found that A-type states made of light quarks have the largest widths into $N\bar{N}$ channels. These will thus show up as broad structures in $N(1,0)N(1,0)$ channels ($p\bar{p}$, for example), unless the masses are very close to threshold, as is the case here. In this channel, the most interesting cases are perhaps the $A(1,3)$ and the $A(8,3)$, whose calculated masses are 7 MeV below threshold.

At $L_D=2$, the B states show widths into $N(1,0)N(1,0)$ channels that are quite small (< 10 MeV), but at the high end of the mass window these increase by more than 200 MeV. At $L_D=3$ and 4, this partial width increases to about 100 MeV, due to the increased phase space available. However, at masses larger than the ones we have calculated, these partial widths, and hence the total widths, become perhaps too large to be easily discernible. At $L_D=4$, new $B\bar{B}$ channels become accessible at the higher end of the mass window. The partial widths into these channels, which all involve orbitally excited baryons, are small, of the order of 10 MeV. However, the relatively large total widths may make these states difficult to detect experimentally.

At $L_D=1$, the states $C(m,0,1)$ are marginal in that they have narrow widths into the $N(1,0)N(1,0)$ channel, but only at the high end of the mass window. At $L_D=2$, the partial widths into this channel for the calculated masses remain small (< 10 MeV), and they are small even with masses 300 MeV heavier (< 50 MeV). In addition, at the higher mass, decays into the $N(1,0)\Delta(3,0)$ channel become kinematically allowed, but these widths increase quite rapidly. Nevertheless, within the mass window presented, they remain detectable. At $L_D=2$, therefore, the states $C(m,0,2)$ consistently show some promise.

At $L_D=3$ and 4, the partial widths of these states into $N(1,0)N(1,0)$ are still < 130 MeV, even with masses that are hundreds of MeV larger than those calculated. In addition, new channels, such as $N(1,0)\Delta(3,0)$ and $\Delta(3,0)\Delta(3,0)$, open up. For $L_D=3$, the latter channel is barely accessible for the calculated masses, so that the partial widths are small. However, as in all cases involving this channel, the width increases very rapidly with increasing mass, and these states will only be of experimental interest if their masses are very close to the threshold for this channel, as is the case here. The widths into the $N(1,0)\Delta(3,0)$ channel are also interesting, but become large at higher masses, though not as rapidly as in the $\Delta(3,0)\Delta(3,0)$ channel. For $L_D=4$, other channels become accessible, but these all involve orbitally excited baryons, and the partial widths are all very small (< 5 MeV). Nevertheless, the large total widths of these states may make them experimentally inaccessible.

All that has been said for the $C(m,0,L_D)$ is applicable to the $C(m,1,L_D)$ and $C(m,2,L_D)$ as well, but with the reminder that the $C(m,0,L_D)$ are lighter. This means that new channels become accessible at lower L_D for the $C(m,1,L_D)$ and $C(m,2,L_D)$, and the increased phase space tends to make all the partial widths larger. Total widths are all quite large, so that high-mass C states are almost certainly without experimental interest.

From the above, we may conclude that there are a few good candidates that may be fairly easily observed. Perhaps the most interesting are the low-mass C states. Their decay widths into the $N(1,0)N(1,0)$ channel are consistently small, and are almost always in a range that is quite accessible experimentally. The $A(1,L_D)$ and

$A(8,L_D)$ have very small widths into the $N(1,0)N(1,0)$ channel, but their masses are expected to be larger than those presented herein. The partial widths corresponding to masses that are 150 MeV heavier, for instance, are much larger, so that at more realistic masses, these states may be too broad to be easily detected. The B states are generally between these two extremes, but only the states with L_D small may be of real interest.

We also point out that without resorting to any admixtures of M-diquonia, we obtain widths that are very narrow, even for channels involving only the lowest baryon octet. We have also obtained widths that are quite broad, especially for the channel where both baryons belong to the lowest-lying decuplet. We expect this range of widths to persist for broken SU(3). There, however, the important question to be answered is whether the predicted states will exhibit the same narrow widths as the corresponding experimental candidates, without invoking significant admixtures of M-diquonia.

We conclude this section by pointing out that while many of the partial widths in Table IV may be small, most of the decaying diquonia have large total widths. This means that most of these states will manifest themselves as very broad structures, with resonant features that will be extremely difficult to detect.

VI. CONCLUSIONS AND OUTLOOK

We have carried out a systematic study of T-diquonia assuming that $SU(3)_f$ is valid. In this study, we have analyzed all states compatible with the Pauli principle, and have calculated their total and partial widths into all of the lowest-lying baryon channels, using the $^3P_0$ model. The C states appear to be the most interesting, especially with respect to their decays into pairs of baryons from the lightest octet.

While this study has been instructive, there is much that can still be done. Perhaps one of the more urgent needs is a systematic calculation of the masses of the diquonia. This need is greatest for states with $L_D=0,1$. We believe that the masses we have obtained for states with $L_D\geq 2$ are reliable. Note, however, that a more realistic calculation of masses, especially for $L_D=0$, may change none of our conclusions, since these states almost certainly all lie below the lowest $B\bar{B}$ threshold. For $L_D=1$, some states may be just above this threshold, and so will be interesting since the limited phase space would make their $B\bar{B}$ partial and total decay widths small.

To get a clearer idea of what to expect experimentally, a similar study with broken SU(3) should be carried out. However, this appears to be at least a very tedious task, as the number of possible diquonia states and $B\bar{B}$ channels is very large. To some extent, some of this has already been done, but involving only the SU(2) subgroup of broken SU(3).

Finally, it is essential to compare theory to experiment by estimating, for example, the contribution of diquonium formation and subsequent decay, to processes such as $p\bar{p}\rightarrow p\bar{p},\Lambda\bar{\Lambda}$. The former can be done without strange quarks, and has in fact been estimated by Barbour and Gilchrist.[15] We point out, however, that an error in this work, which has been subsequently corrected, makes the figures invalid. The $\Lambda\bar{\Lambda}$ channel has not yet been investigated.

To conclude, we mention that this study may give a hint at resolving the puzzle of the perhaps remarkable sparsity of experimental diquonia candidates, given the theoretical abundance predicted herein and elsewhere. Part of the problem is that many of the total widths are large, so that many of the theoretical states give contributions to $p\bar{p}$ scattering, for instance, that would appear as very broad, nonresonant features. In addition, the number of states present may be important. In this work, we have looked at a number of states, more than 40 of which decay into the $N(1,0)N(1,0)$ channel. The contribution of this large number of states to the elastic scattering cross section of $N(1,0)N(1,0)$, for example, could be such that very few resonant features are observed. This will be investigated in greater detail in a forthcoming paper.

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*Present address: Department of Physics, Harvard University, Cambridge, MA 02138.

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