Possibility of new dibaryons containing heavy flavors

J. Leandri*
Laboratoire de Physique Théorique, Université de Bordeaux I, rue du Solarium 33175 Gradignan CEDEX, France

B. Silvestre-Brac†
Institut des Sciences Nucléaires, Institut National de Physique Nucléaire et de Physique des Particules du Centre National de la Recherche Scientifique, Université Joseph Fourier, 53 Avenue des Martyrs 38026 Grenoble CEDEX, France

(Received 1 October 1992; revised manuscript received 9 February 1993)

In a recent paper we have shown that the possibility of including heavy flavor in the dibaryon sector can lead to some new favored configurations (relative to the baryon-baryon threshold). In this study we extend our previous work by a systematic study of all the physical Qq^5 systems in a simple chromomagnetic model. In the first part we assume that the q quarks belong to the fundamental irrep of $SU(3)_F$ and that the Q quark has infinite mass. These assumptions are subsequently relaxed by introducing two mass parameters δ and η. Once these symmetries are broken we gain access in our model to a large number of new dibaryons containing heavy flavor. Some of them could be stable against decay via strong interactions, and we indicate the most favorable cases.

PACS number(s): 14.20.Pt, 11.30.Hv, 12.40.Qq

I. INTRODUCTION

Strong interactions govern in a large part the phenomena occurring in systems with sizes of a few fermis. The up-to-date theory for strong interactions is quantum chromodynamics (QCD). In this theory the matter fields are the quarks which interact by exchanging gluons. Both quarks and gluons carry color, quarks (antiquarks) belonging to the **3** ($\overline{\mathbf{3}}$) representation and the gluons to the **8** representation of the $SU(3)_C$ gauge group. The color wave function of a complicated system composed of quarks and gluons may be indeed quite complex. However, a physical principle states that any observable system must be in a color singlet state. Despite this drastic simplification, even the simplest systems are far too intricate to be solved using first QCD principles. One must rely upon some approximations which preserve as much as possible of the fundamental features of QCD. The most popular approaches are the QCD sum rules [1], the Skyrmion theories [2], the bag models [3], and the nonrelativistic models [4]. In those theories the gluonic degrees of freedom are removed and replaced by some phenomenological ansatz. In that case the condition of a color singlet function is known as the triality rule because it imposes that the number m of quarks minus the number n of antiquarks must be a multiple of 3: $m-n=3B$ where B is the baryon number. The quark degrees of freedom are in a large part different from the bare ones occurring in the QCD Lagrangian. In particular the masses of these "constituent" quarks are larger than those corresponding to the "current" quarks. At our level the baryon number B is a good quantum number; in principle the Fock space necessary to describe a system with a given B is infinite since an arbitrary number of $q\overline{q}$ pairs can be added without modifying the B value. However, there is usually a dominant term in the expansion $q^m\overline{q}^n + q^{m+1}\overline{q}^{n+1} + \cdots$. The quarks (and antiquarks) appearing in this leading term are called "valence" quarks (and antiquarks); this denomination is purely semantic. In all the following, when we speak of quarks, we mean *valence* and *constituent* quarks. In this respect, the simplest systems (containing the smallest number of particles $m+n$) are the mesons ($q\overline{q}$) and baryons (q^3). Some of them are bound under strong interactions, while others appear as resonances. The common mechanism responsible for this binding is multigluon exchange and the resulting composite systems are very compact objects with dimensions of order 0.5 fm. There exists another mechanism which can lead to binding among quark systems; it has been known for a long time and is responsible for the cohesion of atomic nuclei. It relies on the exchange of correlated $q\overline{q}$ pairs, the most important one being the π meson. In that case the resulting systems are more diffuse objects (the size of deuteron is more than 2 fm); this mechanism can bind composite quark systems containing more than 600 particles. Immediately the question arises whether the gluon exchange mechanism is able to bind systems more complicated than mesons and baryons. We will give these objects the generic name of "multiquarks" ($q^m\overline{q}^n$ with $m-n=3B$ and $n+m\geq 4$, bound by gluon exchanges). This question is of capital importance. First, the search for multiquarks is a natural evolution in the

*Electronic address (bitnet): LEANDRI@FRCPN11
†Electronic address (bitnet): BSILVEST@FRCPN11

curiosity of mankind: everybody wants to discover new things. Second, and probably more important, is the fact that the multiquarks are a crucial key for testing the validity of QCD theory. At present, this theory predicts the existence of multiquarks. Since they are not forbidden, one may hope to find them in experimental searches. Consequently, if this search program remains unsuccessful, QCD must be seriously questioned. Either the theory is wrong and must be replaced, or it is incomplete and needs to be supplemented by some mechanism which forbids the existence of multiquarks. To be interesting a multiquark state should be bound under strong interaction (the only decay mechanisms are then electromagnetic for excited levels and weak interaction for the ground state) or it should appear as a narrow resonance in some definite channels. For hunting such particles the role of theoreticians is decisive. They must specify the flavor sector, the energy range, the width, and branching ratios. Such an extensive research program is very ambitious and must be constructed in several steps.

First it should be noted that the number of physical configurations allowed for multiquark systems increases dramatically with the number of valence quarks (antiquarks). Second, to be complete, a study of hadronic systems must cover all the flavors allowed by the standard model (even though the corresponding multiquark might be very difficult to produce). This last point enlarges even more the number of possible multiquark states.

Given this increasing complexity we feel that it is essential to build, on *qualitative grounds*, a reliable hierarchy of the most promising candidates. As has been explained with more details in Refs. [5,6,11], to achieve this aim we use the so-called pure chromomagnetic (CM) interaction

$$H = -a \sum_{i<j} \frac{(\boldsymbol{\lambda}_i \cdot \boldsymbol{\lambda}_j)\,(\boldsymbol{\sigma}_i \cdot \boldsymbol{\sigma}_j)}{m_i m_j}, \tag{1}$$

where $\boldsymbol{\lambda}$ indicates the Gell-Mann matrices and $\boldsymbol{\sigma}$ the Pauli matrices acting respectively on color and spin degrees of freedom of interacting quarks with masses m_i, m_j. The only justification for such an expression is simplicity. In fact, the hyperfine contribution due to one gluon exchange contains a term proportionnal to $\delta(\mathbf{r}_{ij})(\boldsymbol{\lambda}_i.\boldsymbol{\lambda}_j)(\boldsymbol{\sigma}_i.\boldsymbol{\sigma}_j)/m_i m_j$. This term is known to be very important since it is responsible for Δ-N splitting and it is absolutely necessary to explain the short range repulsion in the nucleon-nucleon interaction [7]. Here we will be mostly interested in ground-state multiquarks, whose spatial wave function has the highest symmetry for identical particles and consequently is in an $L = 0$ orbital angular momentum. If we make the drastic assumption that $\langle \delta(\mathbf{r}_{ij}) \rangle$ on the spatial wave function is constant for every quark pair, we recover the chromomagnetic Hamiltonian (1). Thus all spatial degrees of freedom are frozen in the coupling constant a, and we deal in this study only with color, isospin, and spin degrees of freedom. In addition to this already severe condition we neglect kinetic-energy terms as well as Coulomb + linear flavor-independent potentials. These terms are generally important but their effect is also important in the baryon-baryon threshold energy and the behavior of energy relative to threshold due to these terms should be rather smooth.

Another approximation in our framework is to treat on the same footing the multiquarks and the thresholds. That is to adopt for the CM coupling constant a the same value for all hadronic systems. This is also a drastic simplification; in fact a does depend on the confinement volume of the quark. An increase of this volume in the multiquark leads to a lower value of the coupling constant (connected to a reduction of the spatial hyperfine matrix element) in comparison with the thresholds [8]. The binding energy, in case of attraction, should be decreased. Of course such a simplified dynamical process is a crude approximation to reality. Recently, Lichtenberg and Roncaglia [9] have applied this model to some experimental situations and proved that in some cases the model is off by 600 MeV or so. Although they exhibited the worse cases for their demonstration and used a set of parameters which can be easily improved, it is true that the chromomagnetic model is unable to give correct quantitative binding energies. As explained above, because of the dependence of a with the system under consideration, quite often the binding energies are overestimated. This can lead to false predictions for the stability of multiquarks. We have some experience on this subject since the H particle ($uuddss$ with null spin and isospin) is bound by 20 MeV in the CM framework while two more refined calculations predicted it to be unbound. The situation is exactly the same for the so-called P multiquarks [10]. Thus, conclusions obtained with CM concerning the multiquark stability must be handled with much care.

Despite these apparently very dramatic shortcomings, we decided to use this model for performing systematics on a large number of samples. Since, for multiquarks, we have no experimental indications, it is urgent for theoreticians to delimit the interesting sectors for stability. We have no objective arguments for rejecting some of them *a priori*. Thus, a systematic study of all possible flavors is unavoidable. Of course, to be really interesting candidates, multiquarks must have a large binding and a flavor content which makes their production and observation a definite possibility. A complete and rigorous treatment of such a tremendous amount of data is out of question for the moment. We need to use a model which is simple enough to allow this systematics while retaining most of the physical principles. From our point of view, the chromomagnetic model is best suited for that.

Our aim is not to give exactly the binding energies of each state. It is not even to predict whether or not a multiquark is stable. Our hope is that the chromomagnetic model helps us to select, among the thousands of possibilities, those that are the most serious candidates for exhibiting stability or narrow resonance properties. In a sense it is a necessary first step towards constructing a hierarchy of interesting objects. In a second step, more rigorous and time-consuming methods could be brought into play in order to make more reliable predictions. But the present study is concerned with a much more restricted set resulting from the CM approach. The proof that the

chromomagnetic model can do this task was established very seriously in the diquonia $(\overline{q_1}\overline{q_2}q_3q_4)$ sector with total orbital momentum $L = 0$.

A systematic study of *all possible* diquonia (all total spin values $S = 0, 1, 2$ and all flavor possibilities for the quarks (u, d, s, c, b) has been carried out within the framework of the CM model in Ref. [11]. On the other hand, a similar analysis was performed in Ref. [12] using a realistic quark-quark potential and a sophisticated treatment for solving the four-body problem. Although the absolute values for the energies can differ appreciably between both approaches (the CM giving systematically more binding), the interesting conclusion is that very promising candidates remain essentially the same.

The second justification of our simplified approach appears if we consider the dibaryon sector. The particle that has retained the most attention is the so-called H particle proposed long ago by Jaffe [13] with a binding energy around 80 MeV. After a recent systematic study [5] in the CM framework with all possibilities of q belonging to an $SU(3)_F$ representation (all dibaryons formed with u, d, and quarks of only one additional flavor) seven multiquarks were found bound and, in contrast with what was expected, the most tightly bound particle was not the H but the strangeness-3 particle $uudsss$ (binding energy 33 MeV in CM). As in the diquonium sector our classification of dibaryons is in global agreement with the previous studies dealing with the same subject (see [14,15] and references quoted therein) and using much more sophisticated models. In particular the existence of the $N\Omega$ resonance had been demonstrated previously by Goldman *et al.* [16] within two different quark models. Through these examples one can check that the interesting candidates do not depend crucially upon whether one works in the simple CM model or with a more refined approach.

After having made several systematic studies it is our experience that the most favorable candidates are found whenever we "share the flavor at maximum." For example n^6 (n denotes an ordinary u or d quark) is never bound, n^5s has yet more binding, n^4ss (such as H) has more binding, and n^3sss (such as the $N\Omega$ resonance) even more. There is a qualitative explanation for this phenomenon. In the CM model it is not difficult to show that identical particles always (i.e., whatever are the pair quantum numbers) feel repulsion, while different particles may feel attraction in some channels. Thus to have a maximum binding it is judicious to maximize the number p of pairs containing different flavors. Among the 15 possible pairs in a dibaryon, $p = 0$ for n^6, $p = 5$ for n^5Q, $p = 8$ for n^4QQ, and $p = 9$ for n^3QQQ, which explains the hierarchy found in Ref. [5]. In this reference only n and *one* heavy flavor Q [$n = u, d$, and Q forming an $SU(3)_F$ representation] were allowed and $p = 9$ is the maximum possible value. But the situation is even more favorable if one allows in addition to the heavy flavor Q [entering in the $SU(3)_F$ representation] a second heavy flavor Q'. For example $p = 9$ for n^4QQ' and nQ^4Q', while $p = 11$ for n^3QQQ' and nnQ^3Q'. The aim of this paper is the extensive study of $Q'q^5$ [$q = n, Q$ is a member of a $SU(3)_F$ representation] with the hope of discovering new and *a priori* more interesting (in CM picture) candidates than in the q^6 sector. The paper is organized as follows. In Sec. II the model for $Q'q^5$ is presented. In Sec. III we focus our attention on the corresponding thresholds, while in Sec. IV all the interesting results are quoted. In Sec. V we give a short conclusion.

II. THE MODEL

Since the techniques are similar to those used in our previous work [5], we will sketch them briefly, focusing our attention on the novelties in comparison with the q^6 sector. In the systems under consideration Qq^5, Q denotes any quark with isospin 0 and q belongs to the fundamental representation 3 of $SU(3)_F$ flavor group. In a first step (A) the $SU(3)_F$ symmetry is supposed to be exact and the Q quark is assigned infinite mass. In a second step (B) the $SU(3)_F$ symmetry is broken by the introduction of a parameter δ. Finally, in a third step (C), the condition $m_Q = \infty$ is relaxed and another parameter η is defined. Let us recall here a very important point. Once the flavor symmetry is broken, the mass of the $I = 0$, $Y = -2/3$ particle can be anything. Therefore, the corresponding particle is not limited to be only the strange quark. It can be the c or b quark as well. The hypercharge Y is in that case related to the charm and beauty quantum numbers. In the case of charmed q^5 systems our definition of hypercharge does not coincide with the usual one since the c quark carries $Y = 4/3$. Thus not only multiquarks like $cuudds$, $budsss$, ... appear in our study but very exotic flavor objects like $sudccc$, $cuudbb$, $buddcc$, ... are described as well and open the door to many interesting possibilities for new resonance channels. For simplicity in the notation we will nevertheless speak of the $I = 0$, $Y = -2/3$ quark coming from $SU(3)_F$ symmetry breaking as the "strange" quark despite the fact that one can choose another flavor for some specific multiquarks. In our model Y is related to the number of $I = 0$ quarks (N_0) in the q^5 systems by the relation $Y = 5/3 - N_0$.

A. The q^5 system in $SU(3)_F$ limit

In step (A), because of its infinite mass, the quark Q does not interact with the "light" ones and the Hamiltonian (1) reduces to

$$H = -\frac{a}{m^2} \sum_{i<j} (\boldsymbol{\lambda}_i \cdot \boldsymbol{\lambda}_j)\,(\boldsymbol{\sigma}_i \cdot \boldsymbol{\sigma}_j), \tag{2}$$

where the summation is over the five light quarks q.

The factor $A = \frac{a}{m^2}$ defines the energy scale. From the experimental Δ-N mass difference we obtain $A \approx 18$ MeV. Let us recall here that we choose the same coupling constant for both the multiquarks and the threshold calculations. Thus we can express all the energies in A units. The eigenenergies of (2) for N identical particles belonging to the fundamental irrep of an $SU(n)_F$ flavor group

can be obtained easily by very simple formulas (see Ref. [5]) for $n = 2$ and $n = 3$.

Even if in step (A) the Q quark does not participate dynamically, it does fix some constraints on spin and color quantum numbers. We need to classify the q^5 states under irreducible representation of $SU(3)_c \otimes SU(3)_F \otimes SU(2)_S$. First of all, because of the Pauli principle, one must retain only the $[1^5]$ antisymmetric representation of $SU(18)_{CFS}$ which is 8568-fold. It splits into

$$(21, \bar{6}) \oplus (24, 84) \oplus (\overline{15}, 210) \oplus (6, 336) \oplus (\bar{3}, 420)$$

under the chain $SU(18)_{CFS} \supset SU(3)_C \otimes SU(6)_{FS}$. The Qq^5 system must be color singlet and since Q belongs to the 3 representation, the q^5 system must belong to the $\bar{3}$ representation; this limits our attention to the 420 representation of $SU(6)_{FS}$ which splits into

$$\begin{aligned}&(21, 1/2) \oplus (24, 3/2) \oplus (24, 1/2) \oplus (\overline{15}, 5/2) \oplus (\overline{15}, 3/2) \\ &\oplus (\overline{15}, 1/2), \oplus (6, 3/2) \oplus (6, 1/2) \oplus (\bar{3}, 3/2) \oplus (\bar{3}, 1/2)\end{aligned} \tag{3}$$

under $SU(6)_{FS} \supset SU(3)_F \otimes SU(2)_S$. These are the only representations ($C = \bar{3}$, f, S) allowed by the Pauli principle. We will consider all of them in this work.

We consider now the problem of building the corresponding wave functions, and, for that, we adopt a coupling scheme of type $\{[(1,2),3]\,(4,5)\}$ because it will be the most natural for studying the decay into two baryons. For each representation $\alpha = (\bar{3}, f, S)$, we know from group theory that there exists a function $|\psi_\alpha\rangle$ whose quantum numbers are completely antisymmetric under particle exchange. We are faced with the problem of explicitly constructing it. The natural basis to expand $|\psi_\alpha\rangle$ is given by the complete set of states $|C_i, F_j, S_k\rangle$, where the C_i, F_j, S_k are the wave function associated with each standard tableaux corresponding to the Young diagram (of the symmetric group S_5) having the symmetry of the $SU(3)_C$, $SU(3)_F$, and $SU(2)_S$ irrep involved in the decomposition (3). The definitions for the color, flavor, and spin basis states are given in the Appendix.

The algorithm necessary to compute the expansion of $|\psi_\alpha\rangle$ was explained in detail in Ref. [5]. The fact that it gives a solution only for the physically allowed states is a good check of the computer code. As a matter of indication the state $(\bar{3}, \overline{15}, 1/2)$ is expanded over 29 components. Since the corresponding functions $|\psi_\alpha\rangle$ are antisymmetric the eigenenergies are readily obtained by

$$E_\alpha = \langle\psi_\alpha| H |\psi_\alpha\rangle = 10\,\langle\psi_\alpha| H_{12} |\psi_\alpha\rangle. \tag{4}$$

Then in the following only $\langle C_i|\boldsymbol{\lambda}_1.\boldsymbol{\lambda}_2|C_{i'}\rangle$ and $\langle S_k|\boldsymbol{\sigma}_1.\boldsymbol{\sigma}_2|S_{k'}\rangle$ are needed. Because of our coupling scheme, these operators are diagonal. Owing to the properties

$$\begin{aligned}\boldsymbol{\lambda}_1.\boldsymbol{\lambda}_2|(12)_{\bar{3}}\rangle &= -\tfrac{8}{3}\ |(12)_{\bar{3}}\rangle, \\ \boldsymbol{\lambda}_1.\boldsymbol{\lambda}_2|(12)_6\rangle &= \tfrac{4}{3}\ |(12)_6\rangle, \\ \boldsymbol{\sigma}_1.\boldsymbol{\sigma}_2|(12)_0\rangle &= -3\ |(12)_0\rangle, \\ \boldsymbol{\sigma}_1.\boldsymbol{\sigma}_2|(12)_1\rangle &= \ |(12)_1\rangle,\end{aligned} \tag{5}$$

one easily derives

$$\langle C_i|\boldsymbol{\lambda}_1.\boldsymbol{\lambda}_2|C_{i'}\rangle = c_i\,\delta_{ii'} \begin{cases} c_i = -\frac{8}{3}, & i = 2, 4, 5, \\ c_i = \frac{4}{3}, & i = 1, 3, \end{cases} \tag{6}$$

$$\langle S_k|\boldsymbol{\sigma}_1.\boldsymbol{\sigma}_2|S_{k'}\rangle = s_k\,\delta_{kk'} \begin{cases} s_k = 1, & k = 1, 2, 3, 4, 6, 7, 8, \\ s_k = -3, & k = 5, 9, 10. \end{cases} \tag{7}$$

Since in step (A) the Hamiltonian (2) does not depend on flavor, the only relation we need is the orthogonality of flavor functions.

B. Breaking the $SU(3)_F$ symmetry

We proceed further, going to step (B), that is, we break $SU(3)_F$ symmetry but we maintain the condition $m_Q = \infty$. This situation should simulate the case of a heavy top quark Q, although the physics of such exotic objects would probably be more complicated than the naive view exposed here. In the Hamiltonian (1), the mass term can no longer be factorized and the formula (2) is not valid. Nevertheless we can still use the formalism described above at the cost of introducing a "flavor operator" X_{ij}, which is a function of the breaking parameter $\delta = 1 - m/m_s$. As explained in Ref. [5], the Hamiltonian (1) can be rewritten

$$H = -\frac{a}{m^2} \sum_{i<j} X_{ij}(\delta)\,(\boldsymbol{\lambda}_i \cdot \boldsymbol{\lambda}_j)\,(\boldsymbol{\sigma}_i \cdot \boldsymbol{\sigma}_j). \tag{8}$$

There is no fundamental change from the previous treatment. One can still use the same basis states, the same wave function $|\psi_\alpha\rangle$, but in this case the label f of the $SU(3)_F$ representation is no longer a good quantum number. Only hypercharge, isospin, and the total spin S of the five "light" particles are still good quantum numbers. So, each physical state is now specified by the labels $(Y, I, S) = \rho$. However, given a specific set ρ, there may exist several flavor representations f which contain those states as subset. Because of the presence of the $X_{ij}(\delta)$ operator in the Hamiltonian those $|\rho; f\rangle$ states, which were decoupled in the pure $SU(3)_F$ limit, are now coupled and give rise to a matrix which must be diagonalized to obtain the physical eigenvectors. To calculate the matrix element $\langle \rho, f^1| H |\rho, f^2\rangle$ we must introduce a flavor element

$$f^\rho_{ij}(\delta) = \langle F_i| X_{12}(\delta) |F_j\rangle \tag{9}$$

where the F_i and F_j are the flavor basis states contained in the expansion of the initial wave function $|\psi_{\alpha^p}\rangle$ with $\alpha^p = (\bar{3}, f^p, S)(p = 1, 2)$.

The first thing to do is to calculate the $f^\rho_{ij}(\delta)$ coefficients. Since we need the value of X_{12} we expand the F_j with the help of SU(3) Clebsch-Gordan coefficients in terms of our (123), (45) coupling states. Then the action of X_{12} on these states is straightforward. There are many $f^\rho_{ij}(\delta)$ coefficients to be calculated; all these manip-

ulations have been performed using the formal language MAPLE. Thus each element $f^{\rho}_{ij}(\delta)$, which is a second-order polynomial in δ, has been calculated analytically. Then the matrix elements $\langle \rho,\ f^i |\, H \,| \rho,\ f^j \rangle$ are calculated analytically, again using the MAPLE program. We then diagonalize numerically the corresponding matrices for several values of δ. The largest matrix is of order 4×4 describing the cases $(Y = -1/3, I = 1/2, S = 3/2)$ and $(Y = -1/3, I = 1/2, S = 1/2)$.

In summary, as compared to step (A) calculations, our formalism is affected essentially in two ways.

(i) Each matrix element depends on the $(Y, I, S) = \rho$ label and is a function of the breaking parameter δ. Thus there is a splitting among the various ρ subsets of a given SU(3) representation f.

(ii) For a given set of physical quantum numbers ρ, there exists a configuration mixing between n various representations f. One must diagonalize an $n \times n$ matrix, each element of which depends on δ, in order to obtain the physical energies $E_{k,\rho}(\delta)$ $(k = 1, \ldots, n)$.

C. General case

We proceed now a step further (C) and relax the last constraint $m_Q = \infty$. Since the techniques used in this section are very similar to those applied in a previous study of $\overline{Q}q^4$ systems [6], we will sketch them briefly in the following. The problem is to solve *exactly* the Schrödinger equation for the Qq^5 system with particles interacting through the CM interaction (1). In addition to the δ parameter already introduced we define another breaking parameter $\eta = m/m_Q$ which takes care of the finite Q mass. The interaction between the "heavy" particle Q and the "light" q can no longer be neglected. As in the previous case, one can deal with the general formalism by introducing another flavor operator

$$Z_i(\eta, \delta) \ = \ \eta\, P_i^n \ + \ \eta(1-\delta)\, P_i^s, \qquad (10)$$

where P_i^q, $q = (n, s)$ is a projection operator which takes the value 1 if the i particle is of type q, and 0 otherwise. The total Hamiltonian (1) contains, in addition to the Hamiltonian (8), a residual term H_R that takes care of the interaction between the "heavy" quark Q (denoted as the sixth particle) and the five "light" quarks i:

$$H_R \ = \ -\frac{a}{m^2} \sum_{i=1}^{5} Z_i(\eta, \delta)\, (\boldsymbol{\lambda}_i \cdot \boldsymbol{\lambda}_6)\, (\boldsymbol{\sigma}_i \cdot \boldsymbol{\sigma}_6). \qquad (11)$$

Let us now examine the consequences of the existence of H_R.

The spin S of the first five particles that, up to now, was a good quantum number does not preserve this property. Only the total spin J has a definite value. The Q quark has isospin 0 so that I, the isospin of q^5, is also the total isospin of the system. Finally, the hypercharge, which is closely related to the number of "strange" quarks in q^5, is not affected by the presence of another different particle. Thus, each physical state is characterized by a set of quantum numbers $\tau = (Y, I, J)$. The physical wave function $|\psi_\tau\rangle$ can be expanded on a new natural basis $|[C_i 3]_1; [F_j \rho]; [S_l 1/2]_J\rangle$. To obtain these basis states, we must couple the quantum numbers of the previous basis to those of the single particle Q in order to give the quantum numbers τ of the total system. The matrices that we must diagonalize in this general case are often larger than the ones appearing in the preceding section; the reason is that it is possible to mix states of the old basis with different S which give nevertheless the same J value. For our problem the largest matrix is 8×8. In addition to the fact that we must calculate more matrix elements than before, each element is itself more complicated, since the previous value, if any, must be supplemented by the construction of the H_R part of the Hamiltonian. Because of the already antisymmetrized basis, it can be traced back to five times the contribution due to H_{56}. The color and spin matrix elements due to H_{56} are obtained by usual Racah techniques, while the flavor matrix elements due to Z_5 are derived by direct inspection of the flavor functions $|F_j, \rho\rangle$. These latter elements as well as the energy matrix elements in the new basis are calculated analytically using the MAPLE code. They will be provided on request. The matrix diagonalizations are performed numerically for several values of η and δ parameters.

Before leaving this section, let us shortly comment on the various ways to check the calculations at each step. First, in step (A), application of (4,6,7) allows us to recover the simple values of the energies already obtained by the magic formula (4) of Ref. [5]. In step (B) the most severe test is to consider the $\delta = 1$ limit. The "strange" quark acquires an infinite mass and consequently does not interact with other ordinary quarks (denoted n). For a given hypercharge Y, the quark content of the q^5 system is $n^{Y+10/3} s^{5/3-Y}$. The $s^{5/3-Y}$ quarks have zero energy and can be classified under the group $\mathrm{SU}(3)_C \otimes \mathrm{SU}(2)_S$. They can be labeled by the representations (C_b, S_b). In the same way the $n^{Y+10/3}$ system can be classified under the group $\mathrm{SU}(3)_C \otimes \mathrm{SU}(2)_I \otimes \mathrm{SU}(2)_S$, and labeled by (C_a, I_a, S_a). The corresponding energy E_a (which is in fact the total value of the entire system) can be calculated directly by the magic formula (3) of Ref. [5]. We now seek all the representations (C_b, S_b) and (C_a, I_a, S_a) compatible with (C, I, S) states under consideration; the corresponding E_a give us an analytical value for the energy in the limit $\delta = 1$.

In the general case we have also several checks of our calculation. For instance if we consider the states with hypercharge 2/3, only one "strange" quark appears in the q^5 system so that we deal with $n^4 qQ$, $q = (s, c, b)$, dibaryons with $\eta = m/m_Q$ and $\delta = 1 - m/m_q$. Let us now consider the system with $\eta' = 1 - \delta$ and $\delta' = 1 - \eta$. Obviously one has permuted the role of q and Q and the corresponding system is $n^4 Qq$ which is identical to the original one. Thus the energies coming with the sets (η, δ) and (η', δ') must coincide.

There is another test which is much more relevant. In our study we limit ourselves to the case where the Q quark is different from the "strange" quark appearing in q^5. However, one can impose artificially $\eta = 1 - \delta$ $(m_s = m_Q)$, but, in that stage, the Q quark is not re-

ally identical to the "strange" quark because the Pauli principle does not apply. Nevertheless among all these calculated states some possess the symmetry imposed by the Pauli principle, and the energies of the corresponding states must coincide with the energies of the q^6 system studied in Ref. [5]. All these tests are satisfied by our formal and numerical results and we are therefore very confident of our final results.

III. BARYONS AND THRESHOLDS

Under strong interactions the system Qq^5 may decay into $q^3 + Qq^2$. Since there are two different types of baryon appearing in the threshold, we examine them separately, focusing in particular on the symmetry-breaking occurring from stage (A) to stage (C). By analogy with the dibaryon case the total spin of a baryon is called J and the partial spin of the diquark q^2 is called S.

A. q^3 baryons

The observable baryons are classified in the well-known 56 symmetric representation of $SU(6)_{FS}$. This is split again into the classical $(10, \frac{3}{2}) \oplus (8, \frac{1}{2})$ representations following the chain $SU(6)_{FS} \supset SU(3)_F \otimes SU(2)_S$. Application of the magic formula gives immediately the decuplet energy $E_{10} = 8$ and the octet energy $E_8 = -8$.

In step (B) the flavor multiplets are no longer degenerate; each (Y, I, J) multiplet now possesses its own energy $E_{YIJ}(\delta)$, which is a function of the breaking parameter. The resulting energies are given in Table I. Note that the third step does not exist in this case.

B. Qq^2 baryons

In step (A) the q^2 system can be classified in the antisymmetric $[1^2]$ irrep of $SU(18)_{CFS}$ which is here 153-fold. It splits into $(6, 15) \oplus (\overline{3}, 21)$ under the chain $SU(18)_{CFS} \supset SU(3)_C \otimes SU(6)_{FS}$. Here the pair q^2 must be in the $\overline{3}$ color representation, selecting the 21 representation of $SU(6)_{FS}$. The latter splits again into $(6, 1) \oplus (\overline{3}, 0)$ under $SU(6)_{FS} \supset SU(3)_F \otimes SU(2)_S$. The sextet representation has energy $E_6 = \frac{8}{3}$ and the triplet representation energy $E_3 = -8$.

Since in stage (A) $m_Q = \infty$, the Q-q interaction vanishes, thus E_6 and E_3 are the total baryon energies. As for the total spin of the baryon, it can be either 3/2 or 1/2 for the sextet, and 1/2 for the triplet.

In step (B) the 3/2,1/2 sextet irrep splits into three members (Y, I) namely $A(2/3, 1)$, $B(-1/3, 1/2)$, $C(-4/3, 0)$, with the following energies:

$$E_A = \frac{8}{3}, \quad E_B = \frac{8}{3} - \frac{8}{3}\delta, \quad E_C = \frac{8}{3} - \frac{16}{3}\delta + \frac{8}{3}\delta^2.$$

In the same way the $J = 1/2$ triplet representation splits into two members, namely $X(2/3, 0)$, $W(-1/3, 1/2)$, with energies given by

$$E_X = -8, \qquad E_W = -8 + 8\delta.$$

In step (C) each member of a multiplet is affected by a diagonal contribution depending on η. This contribution is also sensitive to the total spin J and consequently the members A, B, C of the sextet are no longer degenerate for $J = 3/2$ and $J = 1/2$ but split into two levels (indexed by 3 and 1 respectively). The states B_1 and W with the same quantum numbers are coupled by the interaction and the eigenmass states result from the diagonalization of the matrix

$$\begin{pmatrix} \frac{8}{3} - \frac{8}{3}\delta + \frac{16}{3}\eta(\delta - 2) & -\frac{8}{3}\sqrt{3}\,\eta\delta \\ -\frac{8}{3}\sqrt{3}\,\eta\delta & -8 + 8\delta \end{pmatrix}. \tag{12}$$

Let us denote by $e_1(\eta, \delta)$ the lowest eigenvalue of the matrix. Only the corresponding eigenvector will be interesting for our study, since we compare the multiquark energy to that of the lowest threshold. We will denote by B_1 the corresponding baryon although it is in fact a mixture of B_1 and W. The energies of the physical Qq^2 systems are presented in Table II.

C. Thresholds

The decay of $Qq^5 \to q^3 + Qq^2$ is governed by a number of conservation laws. In the pure $SU(3)_F$ limit, the Q quark does not really participate, and a better way to indicate the decay would be $q^5 \to q^3 + q^2$. The first threshold q^3 is characterized by the quantum numbers (f_1, J_1) and the second threshold q^2 by (f_2, S_2). The quantum number of the q^5 multiquark are in the same way (f, S). The decay is allowed if S can be formed by coupling J_1 and S_2. We have also the constraint that the f representation can be formed with f_1 and f_2. For this study we need the decompositions

$$\begin{aligned} \mathbf{10} \otimes \mathbf{6} &= \mathbf{21} \oplus \mathbf{24} \oplus \overline{\mathbf{15}}, \\ \mathbf{10} \otimes \overline{\mathbf{3}} &= \mathbf{24} \oplus \mathbf{6}, \\ \mathbf{8} \otimes \mathbf{6} &= \mathbf{24} \oplus \overline{\mathbf{15}} \oplus \mathbf{6} \oplus \overline{\mathbf{3}}, \\ \mathbf{8} \otimes \overline{\mathbf{3}} &= \overline{\mathbf{15}} \oplus \mathbf{6} \oplus \overline{\mathbf{3}}. \end{aligned} \tag{13}$$

If we are in the framework of steps (B) or (C) the decay of dibaryon $D(Y, I, J)$ into two baryons $B_a(Y_a, I_a, J_a)$ (q^3), $B_b(Y_b, I_b, J_b)$ (Qq^2) is allowed if the classical selection

TABLE I. Energies (in A units) for the various (Y, I, J) members of the q^3 baryon coming from decuplet and octet representations after flavor symmetry breaking.

Δ	(1,3/2)	8
Σ^*	(0,1)	$8 - \frac{16}{3}\delta$
Ξ^*	$(-1, 1/2)$	$8 - \frac{32}{3}\delta + \frac{8}{3}\delta^2$
Ω	$(-2, 0)$	$8 - 16\delta + 8\delta^2$
N	$(1, 1/2)$	-8
Σ	$(-1, 1)$	$-8 + \frac{32}{3}\delta$
Λ	$(0, 0)$	-8
Ξ	$(-2, 1/2)$	$-8 + \frac{16}{3}\delta + \frac{8}{3}\delta^2$

TABLE II. Energies (in A units) for the various (Y, I, J) members of the Qq^2 baryon as a function of the breaking parameters δ and η. $e_1(\eta, \delta)$ is the lowest eigenvalue of the matrix (3.1).

A_3	$(2/3, 1, 3/2)$	$\frac{8}{3} + \frac{16}{3}\eta$
A_1	$(2/3, 1, 1/2)$	$\frac{8}{3} - \frac{32}{3}\eta$
X	$(2/3, 1, 3/2)$	-8
B_3	$(-1/3, 1/2, 3/2)$	$\frac{8}{3} - \frac{8}{3}\delta + \frac{8}{3}\eta(2-\delta)$
B_1, W	$(-1/3, 1/2, 1/2)$	$e_1(\eta, \delta)$
C_3	$(-4/3, 0, 3/2)$	$\frac{8}{3} - \frac{16}{3}\delta + \frac{8}{3}\delta^2 + \frac{16}{3}\eta(1-\delta)$
C_1	$(-4/3, 0, 1/2)$	$\frac{8}{3} - \frac{16}{3}\delta + \frac{8}{3}\delta^2 - \frac{32}{3}\eta(1-\delta)$

rules on hypercharge, total isospin, and spin are satisfied. Moreover the decay is forbidden if the dibaryon energy $E_D(\eta, \delta)$ is less than the energy of an allowed threshold $E_T(\eta, \delta) = E_{B_a}(\delta) + E_{B_b}(\eta, \delta)$. Let us emphasize a point concerning the decay. The CM Hamiltonian commutes with $\mathbf{L}$ and $\mathbf{S}$ separately; the orbital and spin angular momenta are good quantum numbers. Thus, in our model, they must be conserved during the decay. Since we are interested only in $L = 0$ multiquarks, the only relevant thresholds (those with the lowest energy) have obviously $L_a = L_b = 0$ and appear in the s wave. Concerning the spin, the total value S must result from the coupling of S_a and S_b. In fact, the true total interaction conserves only the total angular momentum J, resulting from the coupling of L and S, and the parity. In that case a dibaryon with $L = 0$ and $S = 2$ is allowed to decay into two baryons (L_a=0,S_a=1/2) and (L_b=0,S_b=1/2) but in a relative D wave. Since we want to be coherent in our interaction used for the complex system on one hand and for the thresholds on the other, such a decay is not considered here because it is strictly forbidden by the chromomagnetic interaction. If the part of the interaction physically responsible for such a decay (tensor force, annihilation, etc.) is weak as compared to the rest of it, the decay would be strongly disfavored.

IV. RESULTS

We applied the program presented in Sec. II to all possible states. Although we have made a systematic study it would be unrealistic to present all the possibilities in this paper. In stage (A), among the nine possible $SU(3)_F$ irreps we found only three possibilities below or at threshold: the state $(f, S) = (\overline{3}, 3/2)$ is bound by $-20/3$ under strong interaction; on the other hand, the states $(\overline{15}, 5/2)$ and $(\overline{3}, 1/2)$ lie exactly at the threshold energy and need further calculations to decide whether or not they are bound. None of these states have been studied previously.

Let us proceed now to step (B); the breaking parameter δ is introduced. Realistic values for (δ) are $\delta = 0.4, 0.8, 0.93$ for s, c, and b quarks, respectively. Let us recall that the Q quark is taken as the top quark in this case. We present below only those situations where binding occurs. In Fig. 1 we plot the values of $\Delta E = E - E_T$ as functions of δ for the states $Y = 2/3, I = 0$. The S values are also indicated for each curve. The state with $S = 5/2$, which was precisely at threshold for $\delta = 0$, becomes weakly bound after symmetry breaking. One state $S = 3/2$, which was tightly bound in a pure $SU(3)_F$ limit, becomes less and less bound but is still below threshold up to $\delta = 0.5$.

In Fig. 2 we examine the states with $Y = -1/3, I = 1/2$. The $S = 5/2$ level which was at threshold for $\delta = 0$ acquires a very weak binding. One state with $S = 3/2$ is always bound but its binding is greater at small δ values. Finally, one state with $S = 1/2$ becomes rather strongly bound when symmetry breaking is included.

In Fig. 3 we present the states with $Y = -4/3$. The only states displaying binding have characteristics $I = 0, S = 1/2, S = 3/2$. In any case the binding is very weak and occurs only for $\delta > 0.47$. It is remarkable that the only interesting states once $SU(3)_F$ symmetry is bro-

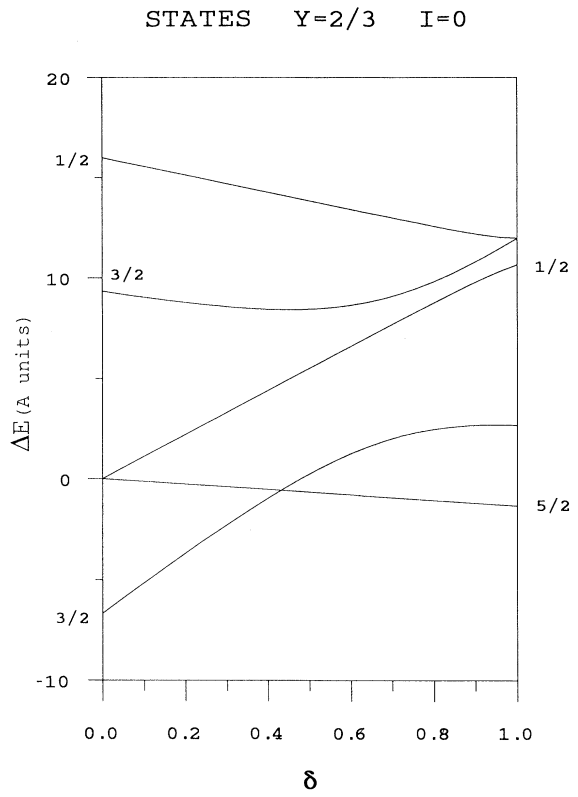

FIG. 1. Energies relative to threshold (in a/m^2 units) vs $SU(3)_F$-breaking parameter δ for q^5 states of hypercharge 2/3 and isospin 0. The total spins S are shown with the corresponding curves.

TABLE III. All possible (Y, I, J) studied in this paper with their corresponding multiplicity N.

Y	I	J	N	Y	I	J	N	Y	I	J	N
$\frac{5}{3}$	$\frac{5}{2}$	1	1	$\frac{2}{3}$	2	2	1	$-\frac{1}{3}$	$\frac{3}{2}$	3	1
		0	1			1	3			2	3
	$\frac{3}{2}$	2	1			0	2			1	5
		1	2		1	3	1			0	3
		0	1			2	4		$\frac{1}{2}$	3	1
	$\frac{1}{2}$	3	1			1	6			2	5
		2	2			0	3			1	8
		1	2		0	3	1			0	4
		0	1			2	3				
						1	4				
						0	2				
$-\frac{4}{3}$	1	3	1	$-\frac{7}{3}$	$\frac{1}{2}$	2	1	$-\frac{10}{3}$	0	1	1
		2	3			1	3			0	1
		1	5			0	2				
		0	3								
	0	2	2								
		1	4								
		0	2								

ken are those coming by continuity from the interesting candidates mentioned previously in stage (A).

We now proceed a step further and come to the general stage (C). In Table III we present all the possible dibaryons for each (Y, I, J) value with the corresponding multiplicity. Thus for each (η, δ) we obtained 100

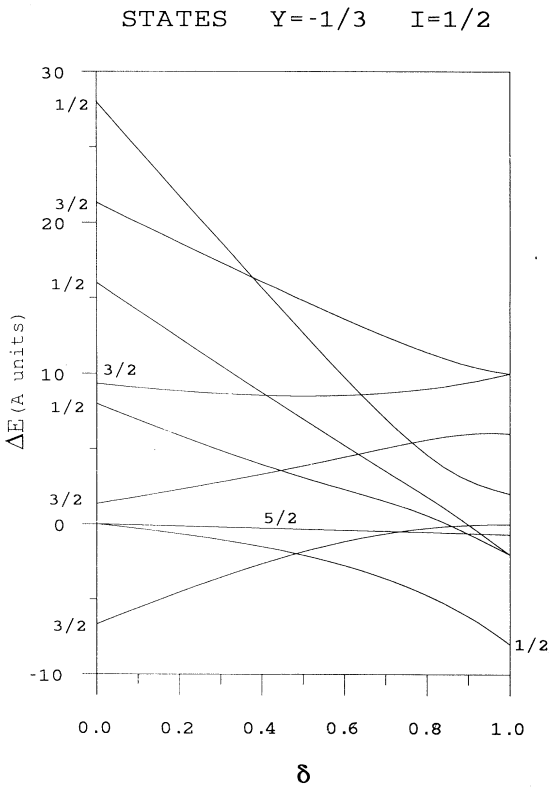

FIG. 2. Same as Fig. 1 for the hypercharge $-1/3$ and isospin 1/2 sector.

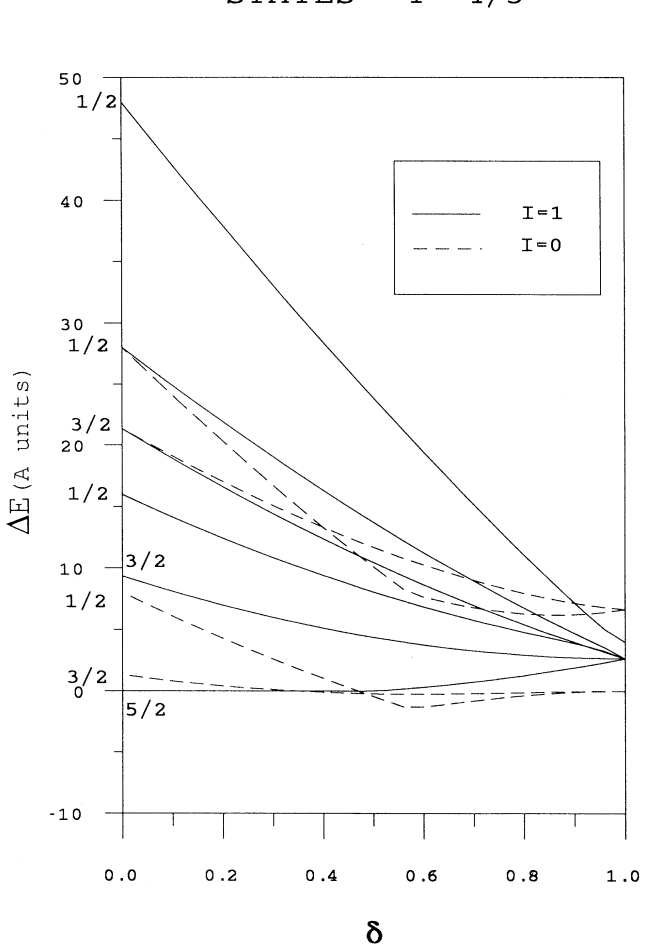

FIG. 3. Same as Fig. 1 for the hypercharge $-4/3$ sector. The solid lines correspond to isospin 1 and the dashed lines to isospin 0.

multiquarks and among them 40 are ground states. Realistic values for (η) are $\eta = 0.6, 0.2, 0.07$ for s, c, and b quarks, respectively. If we choose all possible physical values for (η, δ) we find 300 dibaryons. We have represented the results for each ground state as a two-parameter surface $\Delta E = E(\delta, \eta)$. The systematic study was made over the 40 ground states. However, following the same philosophy as earlier, we present here only a selected sample of these. Since we are interested mainly by "bound" multiquarks, it is natural to retain in our analysis only the diagrams which present domains with binding. Moreover, only some of these domains are of interest because, although the variation over our parameters represents all possibilities for quarks masses, only the realistic values for (η, δ) quoted before make sense. In addition we must impose $\eta \neq 1 - \delta$ in order that the Q quark is different from the "strange" quark. These physical constraints eliminate some diagrams (this eliminates also all the physical dibaryons already studied in Ref. [5]). Thus in the following we present only the diagrams for which a physically interesting region exists. We discuss sectors with different hypercharge separately.

A. Y=5/3 states

There is no "strange" quark present so that the parameter δ is irrelevant. The only parameter of interest is η. No interesting dibaryons emerge from that study; the multiplet $I = 1/2$, $J = 2$ acquires binding but for unrealistic high values of the parameter η while the multiplet $I = 1/2$, $J = 3$ remains always at threshold energy regardless of the nature of the Q quark.

B. Y=2/3 states

The $I = 2$ states lie very high in energy. The situation is more favorable for the $I = 1$ sector but the rare case where binding occurs concerns unphysical values of the parameters. The $I = 0$ states are much more interesting; some of them are represented in Fig. 4. In part (a) the $J = 3$ states are displayed; binding always exists althought it is quite weak. In part (b) we plot the state $J = 2$ which exhibits more binding; the most interesting candidate requires low values of δ and η ($uuddsb$). For one state with $J = 1$ the binding is very impressive but only with nonphysical values of the parameters. The situation is identical for the $J = 0$ state. The part (a) comes by continuity from the $(\overline{15}, 5/2)$ representation of stage (A), while part (b) comes from the corresponding $(\overline{3}, 3/2)$.

C. Y=−1/3 states

All states with $I = 3/2$ lie above threshold. Some $I = 1/2$ states are displayed in Fig. 5. In (a) the $J = 3$ states are represented; the region of binding is very small. In part (b) we represent one $J = 2$ state. It is quite interesting because binding exists everywhere in the (δ, η) plane. In part (c) we show one state with $J = 1$; the region where binding occurs is quite extended but the deepest binding occurs only for unphysical values of the parameters. For one state with $J = 0$ the most strongly bound systems occur for unrealistic values of η and δ. Here again parts (a) and (b) follow by continuity from the $(\overline{15}, 5/2)$, $(\overline{3}, 3/2)$ representations of stage (A). Because

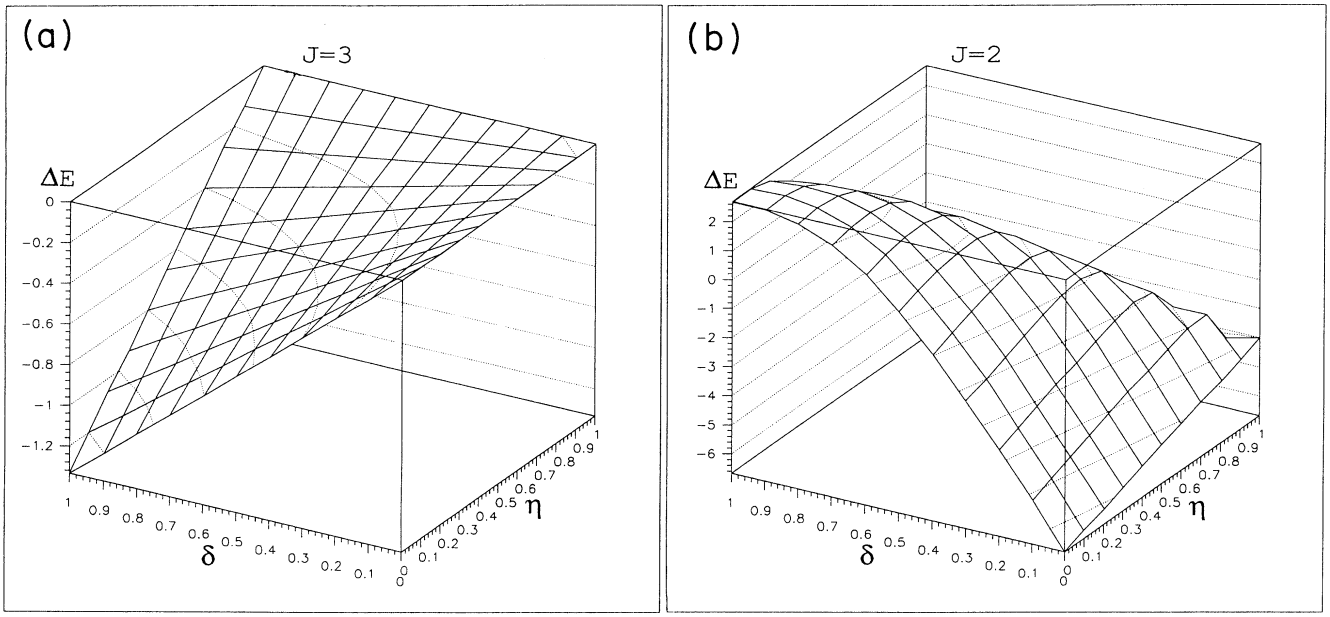

FIG. 4. Energies relative to threshold (in a/m^2 units) vs symmetry-breaking parameters η and δ for Qq^5 states of hypercharge 2/3 and isospin 0. The total spins J are shown with the corresponding surfaces. Part (a) for $J = 3$, part (b) for $J = 2$.

of sudden changes in the thresholds between $\eta = 0$ and $\eta \neq 0$, part (c) is much more difficult to interpret.

D. $Y=-4/3$ states

All $I = 1$ states are unbound. Some $I = 0$ states are displayed in Fig. 6. In part (a) for total spin 2, we show states which exhibit binding throughout a large domain. For all physically acceptable choices of (η, δ) dibaryons are found below threshold. In part (b) for $J = 1$, some physical multiquarks are found to be bound. In part (c) for $J = 0$ states, we see a large valley of strong binding. Fortunately, almost all physical multiquarks lie inside the valley. It is worthwhile noting that the states appearing in Fig. 6 are those of Fig. 3 which gave a tiny binding. In these special cases, reducing the mass of the Q quark helps a lot in getting favorable binding.

E. $Y=-7/3, Y=-10/3$ states

All the states with these quantum numbers lie above threshold.

V. CONCLUSION AND OVERVIEW

We have made a somewhat qualitative overview of the possibly interesting candidates. To have a better quantitative grasp of our whole study we present in Table IV all the dibaryons bound by more than 5 MeV with the chromomagnetic interaction. The binding energies $B = -\Delta E$ are calculated in MeV using the energy scale $A = 18$ MeV. The states are ordered by decreasing binding. There exist 32 stable dibaryons of type Qq^5 within the CM framework (this must be compared to 300, the total number of multiquarks studied in this paper). This

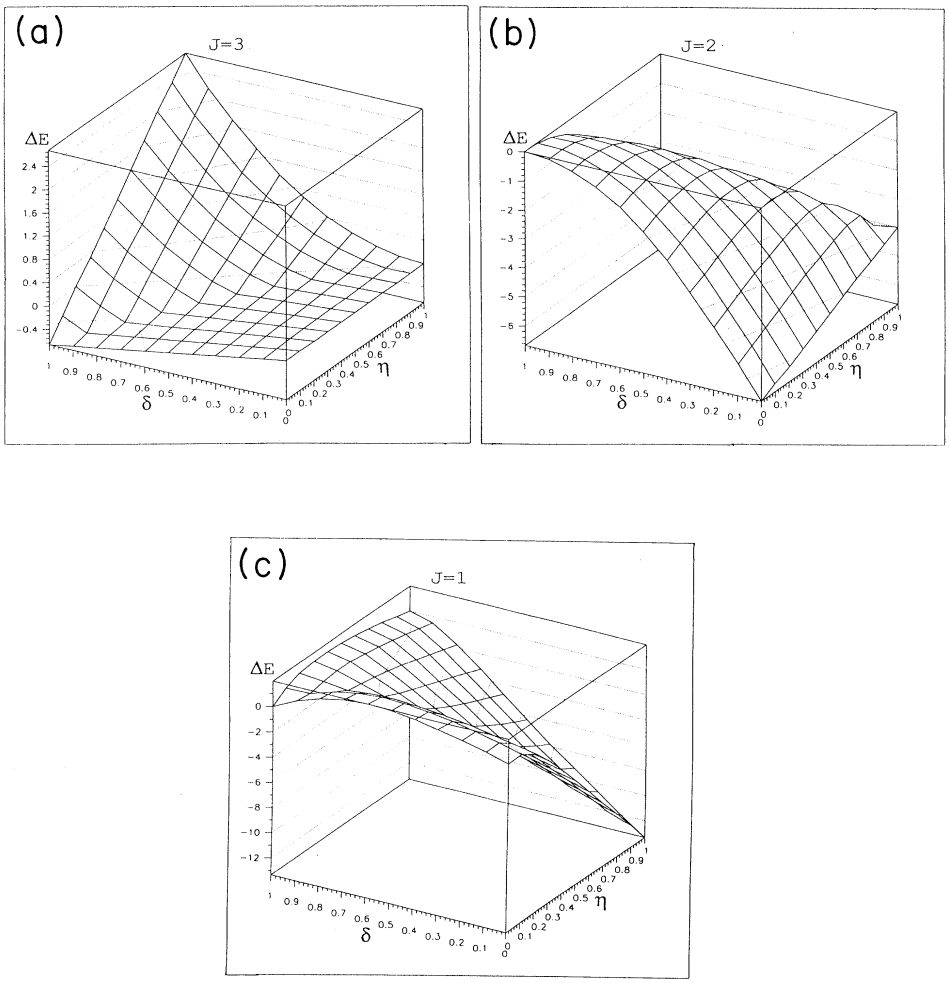

FIG. 5. Energies relative to threshold (in a/m^2 units) vs symmetry-breaking parameters η and δ for Qq^5 states of hypercharge $-1/3$ and isospin $1/2$. The total spins J are shown with the corresponding surfaces. Parts (a),(b),(c) refer to $J = 3, 2, 1$, respectively.

TABLE IV. Dibaryon systems bound within the chromomagnetic model ($B \geq 5$ MeV). The entries are the isospin I, spin J, and the binding energy B in MeV. We adopt $A = 18$ MeV and $a(Qq^5) = a(q^3) = m^2 A$ [17].

I	J		B (MeV)	I	J		B (MeV)
0	0	*sudccc*	76.7	0	2	*cudsss*	12.6
0	0	*sudbbb*	57.1	0	0	*budccc*	12.3
$\frac{1}{2}$	2	*buudss*	43.9	0	0	*cudbbb*	10.8
$\frac{1}{2}$	2	*cuudss*	38.2	0	3	*suuddb*	8.9
$\frac{1}{2}$	2	*suudbb*	30.8	$\frac{1}{2}$	3	*buudcc*	8.9
0	1	*sudbbb*	29.7	0	3	*suuddc*	7.7
$\frac{1}{2}$	2	*suudcc*	27.9	0	2	*suuddc*	7.5
0	0	*cudsss*	27.8	$\frac{1}{2}$	1	*suudcc*	6.5
0	3	*cuuddb*	17.9	0	2	*budsss*	5.6
0	2	*suuddb*	13.8	0	2	*sudccc*	5.0

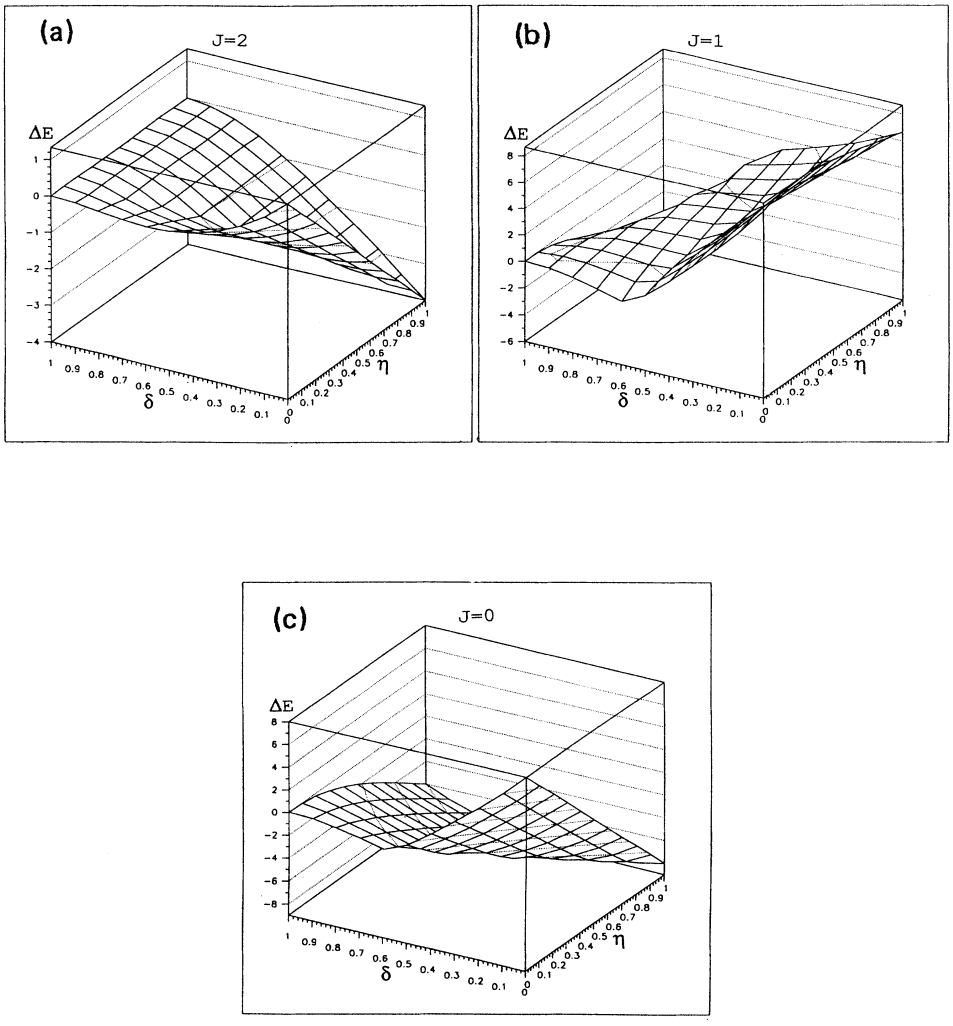

FIG. 6. Energies relative to threshold (in a/m^2 units) vs symmetry-breaking parameters η and δ for Qq^5 states of hypercharge $-4/3$ and isospin 0. The total spins J are shown with the corresponding surfaces. Parts (a),(b),(c) refer to $J = 2, 1, 0$, respectively.

is only a small fraction, indicating the usefulness of systematics made in the CM framework to drastically reduce the number of interesting candidates which should be further studied by more refined methods. Moreover, most of these states are weakly bound so that they probably would become unstable once a proper dynamical treatment is done. Several remarks are in order. All possible values of J appear in the table; thus there is not really a most favorable spin value. For a given flavor quark content, it is *always* the states with lowest isospin that are most deeply bound. This is because in such cases the number of identical pairs is minimized. In the same spirit it is always that flavor configuration which minimizes this latter number that is favored by the CM force. In the Qq^5 system studied in this paper, there exist most possibilities of finding different pairs than in the q^6 system considered in Ref. [5]. As a matter of comparison we found only seven stable dibaryons in the q^6 case, compared to 32 in the Qq^5 system.

It has often been argued in the past that the H particle was the best possible candidate for stability among the dibaryons, whereas, using only the CM interaction, we have already confirmed the existence of a better candidate with strangeness -3 in the q^6 sector [16]. If we extend our analysis to the Qq^5 system, eight new possibilities emerge (the most strongly bound candidates of Table IV). All of these possibilities were, to the best of our knowledge, unknown before our study and can be considered as new predictions of promising multiquarks. However, it will be very hard to produce them experimentally because one has to create three or four units of heavy flavor (c, s, b). Moreover, the values of binding given by the CM potential are probably overestimated. Nevertheless, the new proposed candidates such as *cuudss* $(I = 1/2, J = 2)$, and *suudcc* $(I = 1/2, J = 2)$, as well as the previously studied *uudsss* $(I = 1/2, J = 2)$ and *uuddss* $(I = 0, J = 0)$, even if they are not physically bound, may well appear as resonances in a spectrum based on invariant masses of two baryons already discovered experimentally. The multiquark *sudccc* $(I = 0, J = 0)$, because of its large binding, may become a very interesting challenge in the near future. Recent calculation based on the relativistic coalescence model [18] show that at the CERN Super Proton Synchrotron (SPS) and Brookhaven Alternate Gradient Synchrotron (AGS) energies ultrarelativistic heavy ion collision experiments should be able to discover H and P multiquarks. To produce multiquarks containing more than one unit of charm we can follow the same path but at higher energy and thus we have to wait for the next generation of colliders [BNL Relativistic Heavy Ion Collider (RHIC), CERN Large Hadron Collider (LHC)] (nucleon-nucleon collisions produce mostly D mesons and baryons with one unit of charm). A significant production rate is predicted for certain multiheavy baryon (e.g., Ω_{ccc}) at these high energies [19].

Let us note in passing that the quantum numbers favorable for strong binding seem to be $(I = 0, J = 0)$ for n^2q^3Q multiquarks or $(I = 1/2, J = 2)$ for n^3q^2Q multiquarks. We have no intuitive explanation for such a phenomenon.

How many stable dibaryons will survive after a proper treatment, including kinetic-energy terms and a more elaborate quark-quark potential, has been performed? It is very hard to answer without doing such calculations, because very subtle mechanisms enter both in the dibaryon and in the thresholds. Our guess is that there would remain very few, if any, stable multiquarks in this sector. But, here too, one needs much more refined calculations before presenting any reliable quantitative conclusions. The principal merit of this study is to make a very drastic selection, among the several hundred possibilities, of ten or twenty most favorable candidates that could be detected experimentally. This conclusion should remain valid, even after performing more complete calculations.

ACKNOWLEDGMENTS

We thank Dr. J. T. Donohue for reading the manuscript. Laboratoire de Physique Théorique is Unité associée au CNRS No. 764.

APPENDIX: COLOR, FLAVOR, AND SPIN BASIS STATES

We give in this section all the possible color C_i, flavor F_j, and spin S_k functions corresponding to the symmetry of each allowed Young diagram of the symmetric group S_5.

	$IR\,[2^21]\ n=5$		
$\overline{3}$			
$C_1 = F_1$	$\{[(12)_6\,3]_8\,(45)_6\}_{\overline{3}}$		
$C_2 = F_2$	$\{[(12)_{\overline{3}}\,3]_8\,(45)_6\}_{\overline{3}}$		
$C_3 = F_3$	$\{[(12)_6\,3]_8\,(45)_{\overline{3}}\}_{\overline{3}}$		
$C_4 = F_4$	$\{[(12)_{\overline{3}}\,3]_8\,(45)_{\overline{3}}\}_{\overline{3}}$		
$C_5 = F_5$	$\{[(12)_{\overline{3}}\,3]_1\,(45)_{\overline{3}}\}_{\overline{3}}$		
	$IR\,[5]\ n=1$		
21		5/2	
F_6	$\{[(12)_6 3]_{10}\,(45)_6\}_6$	S_1	$\{[(12)_1 3]_{3/2}\,(45)_1\}_{5/2}$
	$IR\,[4,1]\ n=4$		
24		3/2	
F_7	$\{[(12)_6 3]_{10}\,(45)_6\}_{24}$	S_2	$\{[(12)_1 3]_{3/2}\,(45)_1\}_{3/2}$
F_8	$\{[(12)_6 3]_8\,(45)_6\}_{24}$	S_3	$\{[(12)_1 3]_{3/2}\,(45)_0\}_{3/2}$
F_9	$\{[(12)_{\overline{3}} 3]_8\,(45)_6\}_{24}$	S_4	$\{[(12)_1 3]_{1/2}\,(45)_1\}_{3/2}$
F_{10}	$\{[(12)_6 3]_{10}\,(45)_{\overline{3}}\}_{24}$	S_5	$\{[(12)_0 3]_{1/2}\,(45)_1\}_{3/2}$
	$IR\,[3,2]\ n=5$		
$\overline{15}$		1/2	
F_{11}	$\{[(12)_6 3]_{10}\,(45)_6\}_{\overline{15}}$	S_6	$\{[(12)_1 3]_{3/2}\,(45)_1\}_{1/2}$
F_{12}	$\{[(12)_6 3]_8\,(45)_6\}_{\overline{15}}$	S_7	$\{[(12)_1 3]_{1/2}\,(45)_1\}_{1/2}$
F_{13}	$\{[(12)_{\overline{3}} 3]_8\,(45)_6\}_{\overline{15}}$	S_8	$\{[(12)_1 3]_{1/2}\,(45)_0\}_{1/2}$
F_{14}	$\{[(12)_6 3]_8\,(45)_{\overline{3}}\}_{\overline{15}}$	S_9	$\{[(12)_0 3]_{1/2}\,(45)_1\}_{1/2}$
F_{15}	$\{[(12)_{\overline{3}} 3]_8\,(45)_{\overline{3}}\}_{\overline{15}}$	S_{10}	$\{[(12)_0 3]_{1/2}\,(45)_0\}_{1/2}$
	$IR\,[3,1^2]\ n=6$		
6			
F_{16}	$\{[(12)_6 3]_8\,(45)_6\}_6$		
F_{17}	$\{[(12)_{\overline{3}} 3]_8\,(45)_6\}_6$		
F_{18}	$\{[(12)_{\overline{3}} 3]_1\,(45)_6\}_6$		
F_{19}	$\{[(12)_6 3]_{10}\,(45)_{\overline{3}}\}_6$		
F_{20}	$\{[(12)_6 3]_8\,(45)_{\overline{3}}\}_6$		
F_{21}	$\{[(12)_{\overline{3}} 3]_8\,(45)_{\overline{3}}\}_6$		

[1] D. B. Leinweber, Ann. Phys. (N.Y.) **198**, 203 (1990), and references quoted therein.

[2] A. P. Balachandran *et al.*, Phys. Rev. Lett. **52**, 887 (1984); S. A. Yost and C. R. Nappi, Phys. Rev. D **32**, 816 (1985); R. L. Jaffe and C. L. Korpa, Nucl. Phys. **B258**, 468 (1985).

[3] A. T. M. Aerts and J. Rafelski, Phys. Lett. **148B**, 337 (1984); K. F. Liu and C. W. Wong, *ibid.* **113B**, 1 (1982); P. J. G. Mulders and A. W. Thomas, J. Phys. G **9**, 1159 (1983); P. J. G. Mulders, A. T. M. Aerts, and J. J. de Swart, Phys. Rev. D **17**, 260 (1978).

[4] M. Oka, K. Shimizu, and K. Yazaki, Phys. Lett. **130B**, 365 (1983); Nucl. Phys. **A464**, 700 (1987); B. Silvestre-Brac, J. Carbonell, and C. Gignoux, Phys. Rev. D **36**, 2083 (1987); U. Straub, Z. Y. Zhang, K. Bräuer, A. Faessler, and S. B. Khadkikar, Phys. Lett. B **200**, 241 (1988).

[5] B. Silvestre-Brac and J. Leandri, Phys. Rev. D **45**, 4221 (1992).

[6] J. Leandri and B. Silvestre-Brac, Phys. Rev. D **40**, 2340 (1989).

[7] M. Oka and K. Yazaki, Prog. Theor. Phys. **66**, 556 (1981); A. Faessler, F. Fernandez, G. Lübeck, and K. Shimizu, Phys. Lett. **112B**, 201 (1982); Y. Suzuki and K. T. Hecht, Nucl. Phys. **A420**, 525 (1984); B. Silvestre-Brac, J. Carbonell, and C. Gignoux, Phys. Lett. B **179**, 9 (1986).

[8] T. Goldman, in *Proceedings of the European Workshop on Hadronic Physics with Electrons Beyond 10 GeV*, Dourdan, France, 1990, edited by B. Frois and J.-F. Matthiot [Nucl. Phys. **A532**, 389c (1991)].

[9] D.B Lichtenberg and R. Roncaglia, Report No. IUHET 230, 1992 (unpublished).

[10] S. Fleck, C. Gignoux, J. M. Richard, and B. Silvestre-Brac, Phys. Lett. B **220**, 616 (1989).

[11] B. Silvestre-Brac, Phys. Rev. D **46**, 2179 (1992).

[12] B. Silvestre-Brac and C. Semay, Z. Phys. C (to be published).

[13] R. L. Jaffe, Phys. Rev. Lett. **38**, 195 (1977).

[14] K. Maltman, Nucl. Phys. **A501**, 843 (1989).

[15] K. Maltman, T. Goldman, G. J. Stephenson, Jr., K. E. Schmidt, and F. Wang, in *Glueballs, Hybrids and Exotic Hadrons*, Proceedings of the Workshop, Upton, New York, 1988, edited by S. U. Chung, AIP Conf. Proc. No. 185 (AIP, New York, 1989).

[16] T. Goldman, K. Maltman, G. J. Stephenson, Jr., K. E. Schmidt, and F. Wang, Phys. Rev. Lett. **59**, 627 (1987).

[17] In the simplified CM model we have no decisive and reliable criteria to estimate the value of the energy scale A in Qq^5 states. On qualitative grounds, one expects that an increase of the confinement volume of the quark in the multiquark relative to that in a baryon leads to a reduction of the corresponding spatial hyperfine matrix element, and that would suggest that the energy scale A in Qq^5 states is smaller than in q^3 systems. As already said, this effect is potentially important and could change the hyperfine contribution to the binding energy. To simulate this effect and give an idea on what could be the hierarchy of the most bound Qq^5 multiquarks if we include a reduction of $A(Qq^5)$ versus $A(q^3)$, $A(Qq^5)$ is chosen to fix the binding energy of the H particle at the value zero. With this new value for $A(Qq^5)$ (16.8 MeV) the hierarchy is not affected. For example, the two most interesting charmed multiquarks (*sudccc*, *cuudss*) lie now at 60.9 and 28.1 MeV below their corresponding thresholds.

[18] C. B. Dover, U. Heinz, E. Schnedermann, and J. Zimányi, Phys. Rev. C **44**, 1636 (1991).

[19] P. Lévai and J. Zimányi, KFKI Report No. 31/A, 1992 (unpublished).