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BARYONIUM STATES IN MULTIQUEARK SPECTROSCOPY

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ABSTRACT

The spectrum of $qq\bar{q}\bar{q}$ states, examined in a quark-gluon model combined with dual unitarization, yields two types of baryoniums both of which have narrow widths into mesons. One type has normal hadronic widths ~ 100 MeV into $B\bar{B}$. The other is narrow ~ 10 MeV and goes into $B\bar{B}$ only reluctantly, preferring, if possible, to decay by cascade, and, being a consequence only of hidden colour, is an important object to verify experimentally.

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There have been discovered recently in experiments a number of new states in the $N\bar{N}$ spectrum with masses ranging from below thresholds to about 3 GeV ¹⁾. Some of the states are wide with typical hadronic widths of around 100 MeV, while others are narrow with widths consistent with the experimental resolution of about 10 MeV, even at high masses. Most of these states are characterized by an unexpected reluctance for decaying into meson modes, as inferred either from the measured branching ratios, or from their small total widths. They are therefore popularly called "baryoniums", and are generally regarded as candidates for the long-sought multiquark bound states, probably $qq\bar{q}\bar{q}$.

We examine in this paper the spectrum for the $qq\bar{q}\bar{q}$ system using the quark gluon model combined with some results from dual unitarization (or topological expansion). We found two types of states which may be associated with 'baryonium', both being weakly coupled to mesons. The first type, which we shall call true baryonium (T), has normal hadronic widths into baryon-antibaryons. Whereas, the second type is not genuine $B\bar{B}$ states at all and decays into $B\bar{B}$ pairs only by default of the meson modes. They are expected to have narrow widths even at high masses and prefer whenever possible, to decay by emitting a pion and cascading down to a resonance of the same type. Since these latter mock baryonium states (M) exist only by virtue of their hidden colour, it will be of great interest either to establish or to disprove their existence experimentally.

The experience gained from ordinary meson and baryon spectroscopy allows one to make quite detailed predictions for the spectrum and decay pattern of baryonium states on the leading Regge trajectories, and also to give some estimates for their widths. For instance, we found a M baryonium state with $I^G_J = 1^+_{4^-}$ and mass ~ 2.95 GeV which decays mainly into a pion plus another narrow M baryonium state. The product state can be either $I^G_J = 0^-_{3^+}$ with mass ~ 2.56 GeV, or $I^G_J = 0^-_{3^-}$ with mass ~ 2.15 GeV, and decays prominently into a $N\bar{N}$ pair. This is strikingly similar to the following observation by an Ω spectrometer experiment ²⁾:

$$\begin{aligned} (\bar{p} p \pi^-)_{2.96} &\longrightarrow (p \bar{p})_{2.62} + \pi^- \\ &\longrightarrow (p \bar{p})_{2.10} + \pi^- \end{aligned} \quad (1)$$

We also predict the masses of some narrow states with pronounced decays into nucleon-antihyperon and hyperon-antihyperon.

Consider then a system with two quarks and two antiquarks bound together. We shall leave the confining potential unspecified, apart from assuming as usual that it binds only colour singlets and is independent of both spin and flavour. Such a system will, in general, dissociate easily into $(q\bar{q})$ ($q\bar{q}$) pairs, i.e. into mesons, and will not represent the baryonium states we are interested in. The only way we know to prevent it from doing so is to select the configurations where the quarks and antiquarks are separated by a high orbital angular momentum L . We shall restrict ourselves, therefore, only to such configurations. Further, for simplicity, we shall limit our considerations to the case when the two quarks (antiquarks) in the diquark (antidiquark) are in relative s wave state, which we believe is most favourable for producing a high-lying baryonium Regge trajectory.

Each quark is a triplet in colour $SU(3)$ and a doublet in spin $SU(2)$, which we denote by $(3,2)$. The diquark system can thus be $3 \times 3 = \bar{3} + 6$ in colour, and $2 \times 2 = 1 + 3$ in spin, namely,

$$(\bar{3}, 1), (\bar{3}, 3), (6, 1), (6, 3) \quad (2)$$

Since our s wave diquark has an orbital wave function symmetric under quark permutation, it must be antisymmetric in flavour-colour-spin. The diquark states in (2) must, therefore, be assigned the following flavour $SU(3)$ representations, respectively,

$$\bar{3}, \underline{6}, \underline{6}, \bar{3}. \quad (3)$$

The antiquark is in the conjugate representations.

The confining potential does not distinguish between these states and therefore leaves them degenerate. They will, however, be split by the colour magnetic forces coming from gluon exchanges. Assume, as usual, that such forces are short-ranged so that by asymptotic freedom perturbation theory applies. We have then from the one-gluon exchange term a colour magnetic spin-spin interaction between the two quarks of the following form ³⁾:

$$V_{12} = -C \sum_{a=1}^8 \lambda_1^a \lambda_2^a \underline{\sigma}_1 \cdot \underline{\sigma}_2 \quad (4)$$

C is proportional to the gluon coupling α_s and depends on the overlap of the quark wave function with the potential, while λ^a and $\underline{\sigma}$ are the usual Gell-Mann and Pauli matrices. There are in general also spin-orbit and tensor forces, but these vanish for the s wave state ⁴⁾. The interaction (4) is diagonal in the diquark states (2) and gives for the mass shifts, respectively:

$$-8C, \quad \frac{8}{3}C, \quad 4C, \quad -\frac{4}{3}C. \quad (5)$$

The same values hold also for the antidiquark.

To estimate the size of these splittings, we calculate with the same interaction (4) the mass difference between the total spin $=\frac{3}{2}$ and $\frac{1}{2}$ states in an s wave qqq system, which gives a value $16C$. Identifying this with the experimental value of $m_\Delta - m_N \sim 300$ MeV, we obtain:

$$C \cong 20 \text{ MeV}. \quad (6)$$

One sees therefore that the splitting in (5) can be quite large.

Next we combine the diquark and antidiquark to form a hadron; i.e. colour singlet. We obtain the following possible combinations:

$$\begin{aligned} (\bar{3}, 1), (3, 1); (1, 1) &= [0, \underline{9}] ; -16C & (a) \\ (\bar{3}, 1), (3, 3); (1, 3) &= [1, \underline{18}] ; -\frac{16}{3}C & (b) \\ (\bar{3}, 3), (3, 1); (1, 3) &= [1, \underline{18}] ; -\frac{16}{3}C & (c) \\ (\bar{3}, 3), (3, 3); (1, 1) &= [0, \underline{36}] ; \frac{16}{3}C & (d) \\ (\bar{3}, 3), (3, 3); (1, 3) &= [1, \underline{36}] ; \frac{16}{3}C & (e) \\ (\bar{3}, 3), (3, 3); (1, 5) &= [2, \underline{36}] ; \frac{16}{3}C & (f) \\ (\bar{6}, 1), (\bar{6}, 1); (1, 1) &= [0, \underline{36}'] ; 8C & (g) \\ (\bar{6}, 1), (\bar{6}, 3); (1, 3) &= [1, \underline{18}'] ; \frac{8}{3}C & (h) \\ (\bar{6}, 3), (\bar{6}, 1); (1, 3) &= [1, \underline{18}'] ; \frac{8}{3}C & (i) \\ (\bar{6}, 3), (\bar{6}, 3); (1, 1) &= [0, \underline{9}'] ; -\frac{8}{3}C & (j) \\ (\bar{6}, 3), (\bar{6}, 3); (1, 3) &= [1, \underline{9}] ; -\frac{8}{3}C & (k) \\ (\bar{6}, 3), (\bar{6}, 3); (1, 5) &= [2, \underline{9}] ; -\frac{8}{3}C & (l) \end{aligned} \quad (7)$$

where the third bracket is (total colour, $2 \times$ total spin + 1), the square bracket is [total spin, dimension of flavour SU(3) representation], and the last entry is the colour magnetic mass contribution due to (4) and (5).

The quarks and antiquarks also interact via gluon exchange, giving a colour magnetic force between the diquark and antidiquark, whose spin-spin term is of the form ³⁾:

$$V' = -C' \sum_{a=1}^8 (\lambda_1^a \underline{\sigma}_1 + \lambda_2^a \underline{\sigma}_2) \cdot (\lambda_1^{a*} \underline{\sigma}_1^* + \lambda_2^{a*} \underline{\sigma}_2^*) \quad (8)$$

where $C' \sim C$ when $L=0$, but depends on the angular momentum L in general. For $L \neq 0$, there will also be spin-orbit and tensor terms ⁴⁾, but we shall not deal quantitatively with them at present. These interactions are not diagonal in the representation (7). For the spin-spin term (8), Jaffe ³⁾ has considered the case $C=C'$ in some detail and has given a general method for diagonalizing such operators.

Now it has been noted already ⁵⁾ that the mass splittings between $q\bar{q}$ mesons due to colour magnetic forces as seen in Fig. 1 decrease rapidly as L increases along a Regge trajectory. This can readily be understood if the interaction is indeed short-ranged, as claimed. Along a Regge trajectory, L increases as s , so that the wave function is at an average distance $R \propto L/\sqrt{s} \propto \sqrt{s}$ from the origin. Its overlap with a short-ranged potential must therefore quickly vanish. The strengths of colour magnetic couplings at each value of L can in principle be determined from the resonance masses. Present information is insufficient, however, to allow an unambiguous separation of the various terms.

The interaction between the diquark and antidiquark, being non-diagonal, will mix the states in (7). For $L=0$, when $C'=C$, one sees that the mixing is big ³⁾. When L increases, however, the mixing weakens, being of the order of the mixing interaction strength divided by the mass splittings in (7), namely $\Delta M(L)/\Delta M(L=0)$ in Fig. 1. The states (7) then become approximate eigenstates of the Hamiltonian.

Consider now the possible decay modes of the states (7), starting first with $B\bar{B}$ decays. We distinguish the states (7) (a) - (f) when the diquark is in the colour representation $\bar{3}$ (T baryonium) from those others (7) (g) - (l) when the diquark is in the representation 6 (M baryonium). The diquark in T baryonium states, being in a $\bar{3}$ representation, can combine with another quark in a 3 representation to form a colour singlet baryon. T baryoniums have therefore normal allowed decays into $B\bar{B}$ states, as

indicated in Fig. 2(a), giving widths typically ~ 100 MeV. M baryoniums, however, cannot decay in this way ⁶⁾, since the diquark in a 6 representation when combined with another quark in 3 does not give a singlet: $6 \times 3 = 8 + 10$. To decay directly into $B\bar{B}$ pair, the diquarks in M baryonium will have to shake off their unwanted colour either by exchanging more gluons in a higher order process such as Fig. 2(b), which is presumably small; or by annihilating a quark-antiquark pair as shown in Fig. 2(c), which is also suppressed for the same reason that meson modes are suppressed as we shall later explain. Also, a M baryonium can decay into $B\bar{B}$ by mixing with a T baryonium, but this becomes increasingly difficult for higher L . Indeed, from Fig. 1 one estimates that the mixing angle $\tan \theta \cong \Delta M(L)/\Delta M(L=0)$ is only of order 0.2 and 0.1 for $L=1,2$, respectively. Assuming widths for T baryoniums of ~ 100 MeV, these mixing angles correspond to $B\bar{B}$ widths for M baryonium of only 1-10 MeV. One sees therefore that M baryoniums are not genuine $B\bar{B}$ states at all and will decay prominently into $B\bar{B}$ pairs only when other channels are equally suppressed.

Consider next baryonium decays into mesons. These can proceed either by quark annihilation or quark interchange, as illustrated in Figs. 2(d) and 2(e), respectively. Now, similar processes have been studied extensively in dual unitarization for the simpler case of $q\bar{q}$ mesons, where they are shown to be responsible for the breaking of exchange degeneracy and violations of the OZI rule ⁷⁾. Generally, one finds that these effects will decrease as $(M^2)^{-\delta}$ where δ depends on the particular reaction but is usually of order 1, which is not enough to explain the small mesonic widths of baryonium states. However, these arguments apply only to resonances below the so-called peripheral curve $J \sim M$. For the leading trajectory, the decrease may be much faster if the processes are as short-ranged as the colour magnetic force. Unfortunately, such calculations as exist ⁸⁾ for the leading trajectory in dual unitarization are credible at most up to $\ell \sim 1$. Nonetheless, one can proceed phenomenologically as for the colour magnetic force to estimate these effects from their manifestations in the $q\bar{q}$ meson spectrum. The quark annihilation process is responsible for the splitting of the $I=0$ and 1 meson trajectories. From the mass differences of such resonances, one may therefore estimate its effects as a function of L . In Fig. 1, one sees that these differences behave very similarly to those due to the colour magnetic interaction and have practically the same values. On the other hand, the quark interchange mechanism breaks the exchange degeneracy between, e.g. ρ and A_2 trajectories, and can be estimated there. Experimentally, this splitting is very small, ΔM being only ~ 0.1 GeV at $L=0$, and not detectable for higher L . (This particular case, however, may not be typical.)

Taken together then, one deduces that the decay widths of both T and M baryonium states into mesons are of the same order (namely 1-10 MeV) as the $B\bar{B}$ widths of M baryoniums due to colour magnetic mixing. T baryoniums will therefore decay predominantly into $B\bar{B}$, while M baryoniums will decay into $B\bar{B}$ or mesons with similar probability, but with small widths in both. The situation is further accentuated when L increases.

It is interesting then to ask how high L M baryoniums will decay. The only possibility would seem to be a cascade down to a lower M baryonium state accompanied by the emission of a $q\bar{q}$ meson, as indicated in Fig. 2(f). Now, cascades of this type are known to be strictly governed by angular momentum conservation⁹⁾. The total widths of such decays are expected to be small and to decrease slowly with increasing J . The decay products are most likely a pion plus an M baryonium state lying on the highest Regge trajectory allowed by quantum number conservation. The cascade will continue until the M baryonium has lost sufficient angular momentum to overcome its reluctance (at about $L \sim 2$) to break up into $B\bar{B}$ or into mesons. There is also the possibility that mesons couple weaker to a state where two quarks are in a colour sextet than when they are in a colour triplet as in ordinary hadrons, so that the decay illustrated in Fig. 2(f) will be suppressed more than expected from angular momentum considerations.

To obtain a clearer picture of the spectrum and decay patterns of baryonium states, we have calculated the leading T and M trajectories with no strange quarks under the following assumptions:

- i) the $(q\bar{q})$ ρ trajectory is linear and fitted to known resonances,
- ii) the $(qq\bar{q}\bar{q})$ trajectory is split from the ρ trajectory by a constant $\Delta M \sim 0.93$ GeV, representing the mass of the extra quark pair and any difference in binding energy. The value of ΔM was chosen to give an intercept of the leading $I=2$ T baryonium trajectory comparable to the crude estimates $\alpha_B \sim -0.5$ obtained from phenomenology and from dual unitarization,
- iii) the splitting in M between members of the multiplet is given by (7) assuming $C = 20$ MeV, and taking mixing into account with $C' = C$ for $L=0$. For $L>0$, colour magnetic interaction between diquark and anti-diquark is neglected.

The result is shown in Fig. 3, where the distinction between T and M baryonium states is unambiguous for $L>0$, but is arbitrary for $L=0$ because of the large mixing. Notice that the apparent different slopes obtained for the baryonium trajectories as compared to ρ is a consequence

of assumption (ii) for constant mass splitting. This is known to work very well for $(q\bar{q})$ mesons in explaining the difference in apparent slopes between, e.g. the ρ , K^* and ϕ trajectories and is consistent with equal asymptotic slopes⁵⁾. In principle, the asymptotic slope of the diquark anti-diquark system can be different, especially for M baryoniums because of the different colour content¹⁰⁾. We shall leave this question to be decided by future phenomenology.

The trajectories shown in Fig. 3 are meant for illustrative purposes only, and no attempt has yet been made on assigning predicted states to empirical peaks, or on the optimization of parameters. Clearly, a fair amount of phenomenology remains to be done. We note here only two interesting points.

a) The leading exotic T baryonium is apparently $[2; \underline{36}]$ which for $L=0$ can couple to two pseudoscalars only in a relative d state. This decay is therefore not superallowed in the sense of Jaffe³⁾ and might explain the small cross-sections observed for exotic exchange in pseudoscalar meson-baryon, as compared with baryon-baryon reactions¹¹⁾. One needs not then invoke a Zweig rule for $M^2 = t < 0$, the validity of which would indeed be surprising, in view of the large mixing expected. One predicts also that for vector meson-baryon reactions there will be no such suppression.

b) Cascades of M baryonium proceed via pion emission, which means $\Delta I = 0^{\pm 1}$ and a change in G parity. Previous remarks then imply the most prominent decays would involve the leading trajectories with $I=0$ and 1, namely $[2, \underline{9}]$, $[1, \underline{9}]$ and $[1, \underline{18}]$. Take the $I^G_{JP} = 1^+ 4^-$ resonance at $M = 2.95$ GeV on the $[1, \underline{18}']$ trajectory for example. It has favoured decays by pion emission to $I^G_{JP} = 1^- 3^+$ on $[1, \underline{18}']$, to $I^G_{JP} = 0^- 3^+$ on $[1, \underline{9}]$ and to $I^G_{JP} = 0^- 3^-$ on $[2, \underline{9}]$ with masses 2.65, 2.55 and 2.17, GeV which decay into $p\bar{p}$. An alternative interpretation of (1) is $I^G_{JP} = 1^- 4^-$ on $[1, \underline{18}']$ decaying into $I^G_{JP} = 0^+ 2^-$ on $[1, \underline{9}]$ and $I^G_{JP} = 0^+ 4^+$ on $[2, \underline{9}]$. The close agreement in masses might be accidental but the similarity in decay patterns and absence of $p\bar{\Delta}$ decay of the $M = 2.95$ GeV state is well explained in our picture.

Remarks

- i) In addition to the states listed in Fig. 3, there are many more containing strange quarks, which decay prominently into hyperons. In general, for every narrow peak seen in $N\bar{N}$ at mass M_R say, we expect further narrow peaks at $M_R + 115$ MeV in NY or $Y\bar{N}$, and at

$M_R + 230 \text{ MeV}$ in $Y\bar{Y}$. If previous experience for $q\bar{q}$ resonances should be taken as a guide, these predictions could be good to within $\pm 20 \text{ MeV}$.

- ii) Although M baryoniums are weakly coupled to ordinary hadrons in their decay, it should not be assumed that they are as weakly produced or weakly scattered from ordinary hadrons. The fallacy of such an assumption is already apparent in J/ψ physics. $\sigma_T(J/\psi p)$ and J/ψ production cross-sections are both considerably bigger than naïvely expected from its narrow width into hadrons. At least one reason is that production and scattering involve mixing of states at momentum transfers $t < 0$. Given the strong dependence of the mixing on the particle mass $M^2 = t$, it is not surprising that ψ 's and M baryoniums can be prominently scattered and produced. This is a familiar result of dual unitarization ^{7),8)}.
- iii) In our view, the existence of narrow baryonium states is strongly dependent on the possibility of keeping quarks and antiquarks apart by a high orbital angular momentum. The difference then between a $qq\bar{q}\bar{q}$ model for baryonium ^{*)} and that of a $B\bar{B}$ bound state ¹²⁾ is that the former system has to keep together because of colour, however, high the angular momentum. Whereas, in the second case, since the system is held together merely by nuclear forces, high L states will simply break apart into $B\bar{B}$ pairs.
- iv) We can consider also states in $qq\bar{q}\bar{q}$ in which the $(q\bar{q})(q\bar{q})$ systems are separated by high L . There are again two cases, corresponding to $q\bar{q}$ in a colour singlet or octet representation, respectively. The singlet case (T mesonium) dissociates readily into mesons and is probably not recognizable as resonance states. The octet case (M mesonium), however, have difficulties again decaying into mesons and baryons. Their properties are similar to M baryoniums, though with smaller branching ratios into $B\bar{B}$. They decay again mainly by cascade, but only to M mesonium, not to M baryoniums.

*) A more appropriate name for these states would in fact be "diquonium" to distinguish them from genuine $B\bar{B}$ bound states if such exist.

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FIGURE CAPTIONS

Fig. 1: Mass differences measuring effects of colour magnetic force and quark annihilation process along $q\bar{q}$ trajectories. For $L < 1$ in the annihilation process, we show the values of $A_2 - f$ calculated from dual unitarization ⁸⁾ instead of the experimental mass difference $\omega - \rho$, since this latter case is known to be complicated by strong mixing with the $\lambda\bar{\lambda}$ and $qq\bar{q}\bar{q}$ systems [see e.g. Ref. 7)]

Fig. 2: Decay modes of T and M baryoniums.

Fig. 3: Masses of (a) T baryonium, and (b) M baryonium states on the leading Regge trajectory, calculated in the manner described in the text. The states $L=0$ are inserted only for completeness; they have large widths into mesons and bear generally no resemblance to baryonium states.

FIG.1

FIG.2

FIG.3a

FIG.3b