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BARYONIUM - NUCLEAR ATOM OR COLOUR MOLECULE?

CHAN Hong-Mo

CERN, Geneva, Switzerland

1. INTRODUCTION

An atom is a bound state of a nucleus with electrons held together by electromagnetic forces. When some of the electrons are replaced by other particles such as μ^- or $\bar{p}$, one gets an exotic atom which is the subject of this school. If one now further generalizes the concept by replacing the particles by nucleons or quarks, and the forces as well by nuclear or quark confinement forces, one obtains nuclei or hadrons, respectively. Although these latter systems have very different properties from atoms because of the very different dynamics, there are some similarities in both the experimental and the theoretical techniques involved in their study which are well worth a careful comparison. It is the foresight of the organizers in inviting such a comparison that gives me the pleasure of speaking to you now.

The subject of my lectures is "baryonium". As now used in the literature the name "baryonium"*) refers to a group of empirical

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*) The name unfortunately also carries with it some theoretical prejudices which are not entirely appropriate and which we shall do our best in these lectures to ignore.

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resonances with baryon number zero which have the unusual property of preferring to decay into baryon-antibaryon pairs instead of the kinematically more favourable mesonic final states. For a summary of the present experimental situation, I refer you to the excellent reviews by Montanet¹⁾. Here I shall content myself just with quoting some examples as illustrations.

The best-established of these states were found in formation experiments in $p\bar{p}$ collisions. From accurate measurements of the total and elastic cross-sections, structures are seen at several energies which are usually referred to as the S, T and U bumps. Their properties are summarized in Table 1. One sees that if one interprets these structures as resonances they have large branching ratios into the elastic channel $p\bar{p}$. This is very surprising in view of the large number of competing mesonic channels which are more favourable kinematically. The fact that they prefer instead to decay into $p\bar{p}$ presumably means that their couplings to mesons

Table 1

	S	T	U
Mass (MeV)	1936 ± 1	2185 ± 5	2355 ± 5
Width (MeV)	$8 - 4$	130 ± 30	180 ± 20
$\sigma_T(\text{mb})$	10.6 ± 2.4	5.0	3.1
$\sigma_{el}(\text{mb})$	7.0 ± 1.4	2.9	2.8
$\sigma(p\bar{p} \rightarrow n\bar{n})(\text{mb})$	0.3 ± 0.3	1.2	0.4
$\sigma(p\bar{p} \rightarrow \pi^+\pi^-)(\text{mb})$		$\lesssim 0.24$	$\lesssim 0.12$

are for some reason strongly suppressed. The effect is particularly striking for the S bump, which, being close to the $N\bar{N}$ threshold, should have little chance for $p\bar{p}$ decay and yet was found in this channel with a branching ratio of at least 50 per cent. Such limits as are known for the cross-sections into specific meson channels are indeed smaller than one would have expected (see Table 1) which, if not due to selection rules forbidding decays into these channels, give a measure for these suppressed mesonic couplings. On the other hand, the total widths of the T and U -bumps well above the $N\bar{N}$ threshold are big, of the order 100-200 MeV, which, when combined with the large branching ratios into $N\bar{N}$, implies that they have normal hadronic couplings to $N\bar{N}$ channels of strengths similar, for example, to $f \rightarrow \pi\pi$.

In order to establish that these structures are indeed resonances, one would need to perform some amplitude analysis on the data which is at present impossible for $p\bar{p}$ elastic scattering because of the large number of partial waves involved. It can be done however for $p\bar{p} \rightarrow \pi^+\pi^-$ whose quantum numbers are more selective and for which good data are available in spite of the small cross-sections. The result is listed in Table 2. The masses of the first two states are in rough agreement with the T and U bumps seen in σ_T and σ_{el} , from which we may infer that these structures do indeed contain resonant components in the partial waves $1^+(3^-)$ and $0^+(4^+)$, respectively.

Table 2

$I^G(J^P)$	Mass (GeV)	Width (MeV)
$1^+(3^-)$	2.15	200
$0^+(4^+)$	2.31	210
$1^+(5^-)$	2.48	280

It is clear however that these S, T, U structures are not pure resonances in a single partial wave. If they have pure isospin, whether 0 or 1, then their branching ratios into $p\bar{p}$ and $n\bar{n}$ must be the same, whereas they are either not seen at all (S) in the charge exchange reaction ($p\bar{p} \rightarrow n\bar{n}$), or, if seen, much smaller (T,U) than in elastic scattering (Table 1) implying that they each contain both isospins 0 and 1. Indeed, from a tentative analysis of some data on the reaction $p\bar{p} \rightarrow K^+K^-$, Carter was able to show that in addition to the states in Table 2, there are probable states also in the partial waves $0^-(3^-)$ and $0^-(5^-)$ at about the same masses as the $1^+(3^-)$ and $1^+(5^-)$, respectively. Moreover, there are indications from the reaction $p\bar{p} \rightarrow \pi^0\eta$ that the state $0^+(4^+)$ listed in Table 2 may also have an isospin partner with $I^G(J^P) = 1^-(4^+)$ at a similar mass. We have thus here an apparent isospin degeneracy which is reminiscent of the familiar degeneracy between ρ and ω , f and A_2 , etc., in ordinary meson spectroscopy.

Apart from the S, T, U and V structures just discussed there are suggestions of further states both above and below the $p\bar{p}$ threshold obtained mainly from $\bar{p}d \rightarrow pX$ and $\bar{p}p \rightarrow \gamma X$ experiments. The impression that one gets is a great richness of new states with a particular affinity to the $B\bar{B}$ channel.

That there should exist mesonic states preferring to decay into $B\bar{B}$ is already of considerable interest but this is not all. An examination of the $p\bar{p}$ spectrum from production experiments reveals further surprises. The most popular of these spectra is shown in Fig. 1 which was obtained from backward production in π^-p collisions²⁾. Apart from a small narrow peak at 1930 MeV which presumably corresponds to the S bump already discussed, there are two further spikes with the parameters listed in Table 3. They have appreciable branching ratios into $p\bar{p}$ and are not seen in mesonic channels. In particular, the upper limits for their decays into $\pi^+\pi^-$ and K^+K^- as shown in Table 3 are small, which if

Fig. 1 The empirical $p\bar{p}$ spectrum obtained from a π^-p backward production experiment

Table 3

Mass (MeV)	Width (MeV)	$\frac{M^0 \rightarrow p\bar{p}}{M^0 \rightarrow \text{total}}$	$\frac{M^0 \rightarrow \pi^+\pi^-}{M^0 \rightarrow p\bar{p}}$	$\frac{M^0 \rightarrow K^+K^-}{M^0 \rightarrow p\bar{p}}$
2020 ± 3	24 ± 12	$> 14\%$	$< 15\%$	$< 11\%$
2204 ± 5	16^{+20}_{-16}	$> 16\%$	$< 17\%$	$< 16\%$

not due to their quantum numbers being just such as to forbid these decays, imply that they have suppressed coupling to mesons as compared to $B\bar{B}$. These characteristics are sufficient to qualify them as candidates for baryonium.

However, they have some further peculiar properties which distinguish them from the baryonium states considered above. First, they have small total widths even though they are well above the $N\bar{N}$ threshold. For example, the peak at 2204 has a width of only $\lesssim 20$ MeV as compared with 130 MeV for the T-state of similar mass in Table 1. This means that they must have suppressed couplings also to $B\bar{B}$ channels as compared with normal hadron couplings such as $f \rightarrow \pi\pi$. Second, the production cross-sections of these two states are comparable to that for the S(1936), as can be seen from the experimental numbers quoted in Table 4. Yet, these states have not been seen in formation experiments from $p\bar{p}$ collisions, the upper limits on their formation cross-sections being < 0.4 mb in both σ_{el} and σ_T . When compared with the formation cross-sections for S, T and U quoted in Table 1, one has a suppression factor of at least an order of magnitude. Now, in formation experiment, a resonance is coupled to an $N\bar{N}$ pair where both N and $\bar{N}$ are on mass-shell, whereas in the backward production experiment in Fig. 1, the

Table 4

	$\sigma : \pi^- p \rightarrow \Delta^0 M^0$ $\Delta^0 \rightarrow p\pi^- ; M^0 \rightarrow p\bar{p}$	$\sigma : \pi^- p \rightarrow N^*(1520)M^0$ $N^*(1520) \rightarrow p\pi^- , M^0 \rightarrow p\bar{p}$
S(1936)	9 ± 5 nb	7 ± 5 nb
(2020)	10 ± 4 nb	26 ± 8 nb
(2204)	21 ± 5 nb	

resonance M^0 is coupled to an off-shell " $\bar{N}$ " which is exchanged. The phenomenon noted above may thus be interpreted as saying that the couplings of these states to on-mass-shell nucleons are indeed suppressed, as expected already from their narrow decay widths, but that their couplings to off-mass-shell nucleons are comparable to that of the S. These properties are not only intriguing in themselves but are also quite dramatically different from those of our earlier examples in Table 1. It would thus appear that we are dealing here with yet another new type of objects.

Further narrow states are reported in other production experiments with masses up to 3 GeV with widths comparable to the experimental resolution of order 20 MeV, and decaying into $p\bar{p}$ accompanied in some cases by pions. One of them was even seen to cascade by pion emission into the states (2020) and (2204) just discussed. If taken seriously, they seem to form another family of baryoniums quite as numerous and at least as rich in interest as the first found in formation experiments.

It should be emphasized that few, or perhaps even none, of the baryonium states reported in experiment can yet be regarded as firmly established. Indeed, every one of them is still being hotly disputed among experimenters, not only by opponents but sometimes even by the discoverers themselves. However, assuming even that only some of them survive, they would still represent an exciting new branch of spectroscopy whose unusual characteristics are likely to provide us with new insight into hadron dynamics.

In order to appreciate more fully the challenge presented, let us summarize what we have learned. The two outstanding features of baryoniums are:

- a) They have appreciable branching ratios into $B\bar{B}$ and suppressed couplings to mesonic final states. Yet the $B\bar{B}$ and purely mesonic channels are not, as far as we know, distinguished by any selection rules.

b) There appears to be two families of baryoniums. One family has large couplings to $B\bar{B}$ with widths above $N\bar{N}$ threshold of order 100-200 MeV. The other has widths of order only 10 MeV even high above $N\bar{N}$ threshold indicating suppressed couplings to $N\bar{N}$ pairs as well as to mesons, but yet these resonances can be readily produced.

Apart from these outstanding characteristics there are two further hints which may help in unravelling the mystery:

- c) Those few baryonium states with identified quantum numbers all have fairly high spins ($J \geq 3$).
- d) Some of them are degenerate in isospin 0 and 1.

Obviously, the first question to answer is: "What are these baryonium states?" Presumably they are not ordinary $Q\bar{Q}$ hadrons since these are known to decay into mesons copiously unless forbidden to do so by kinematics or selection rules. They will have to be something else. There are two suggestions:

- A) They are bound or resonant states of a baryon and an anti-baryon held together by ordinary nuclear forces ("nuclear atom") in much the same way as a deuteron is a nuclear bound state of a proton and a neutron.
- B) They are hadrons with constituents $QQ\bar{Q}\bar{Q}$ held together by confinement forces which are supposedly responsible for the formation of all hadronic particles from their quark constituents, e.g., $Q\bar{Q}$ (mesons) and QQQ (baryons).

Here, we have to emphasize one point which is often not made clear in the literature. The assumptions (A) and (B) do not by themselves imply the existence of baryoniums. For example, the $B\bar{B}$ nucleus in (A) may annihilate readily into mesons thus giving it a large mesonic width in contrast to what is experimentally observed (see (a)). Indeed, a low energy $N\bar{N}$ system is known empirically to have large annihilation cross-sections. Similarly,

the $(QQ\bar{Q}\bar{Q})$ system in (B) may also simply recombine into $(Q\bar{Q})(Q\bar{Q})$ dissociating thus into two $(Q\bar{Q})$ mesons. Since no quark pairs need be created, in contrast to decays such as $f \rightarrow \pi\pi$, i.e., $(Q\bar{Q}) \rightarrow (Q\bar{Q}) + (Q\bar{Q})$, one may even expect mesonic widths much in excess of the normal hadronic widths of order 100 MeV.

Therefore, in order to make either (A) or (B) into baryoniums, they have to be supplemented by further conditions. Again, two solutions have been suggested:

- $\alpha)$ There is a new selection rule hitherto unrecognized which forbids those resonances decaying into mesons.
- $\beta)$ Baryoniums have "high" spins; mesonic decays are suppressed kinematically by the angular momentum barrier effect.

Now the solution (α) cannot apply to the $B\bar{B}$ nuclear model (A) simply because nuclear forces at long and medium ranges are quite well known and there is no room for the assumption of a new selection rule. It can, however, be imposed on the $(QQ\bar{Q}\bar{Q})$ hadron model (B) if one so wishes, since "confinement forces" are in any case so badly understood that assumptions made here cannot easily be disproved. In essence, what one would say here is that the "di-quark" QQ in the $(QQ\bar{Q}\bar{Q})$ hadron is not simply just two quarks, but is a two-quark system distinguished by a new characteristic, e.g., a "junction". This characteristic is conserved, thus behaving like a new quantum number and giving rise to a new selection rule. With such a device one can clearly suppress the mesonic decays of baryoniums -- in other words, the effect is not explained but simply assumed.

Now if there exist baryonium states with low spins very stable against mesonic decays which cannot be explained by the angular momentum barrier alone, we may be forced to introduce a new selection rule as in (α) . Notice that what actually matters is not the total spin J but the orbital angular momentum L between B and $\bar{B}$ in (A)

and between (QQ) and $(\bar{Q}\bar{Q})$ in (B). For example, for $L = 2$, one can still build a $J = 0$ state from a spin 1 diquark and a spin 1 anti-diquark, which can already be quite stable due to the angular momentum barrier. In any case, as noted in (c) above, all those baryoniums with identified quantum numbers have $J \geq 3$. It seems therefore uneconomical at this stage to introduce a new selection rule which has also little theoretical justification. For this reason, I shall not discuss the solution (α) any further and concentrate on the alternative (β).

2. NUCLEAR FORCES VERSUS QUARK CONFINEMENT FORCES

In order to distinguish $B\bar{B}$ nuclei from $QQ\bar{Q}\bar{Q}$ hadrons as candidates for baryoniums, let us first examine the forces responsible for holding them together in each case.

A) Nuclear forces between nucleons are quite well understood for distances ≥ 0.8 fermi³⁾. They can be described in terms of a potential due to the exchange of particles, e.g., pions, as illustrated in Fig. 2. The longest range part due to single pion exchange is familiar. The 2π exchange term, effective at medium ranges is sometimes approximated by the exchange of a fictitious σ resonance, but can be more accurately calculated using dispersion relations which has been done, for example, in great detail by the "Paris Group". The potential due to 3π exchange is much harder to calculate and is usually just approximated by ω exchange. Taking

Fig. 2 The nucleon-nucleon potential due to particle-exchange

these exchanges into account, one believes that a fair description is obtained for the interaction between nucleons for $r \geq 0.8$ fermi, but it is unlikely that the procedure can be extended further to smaller distances.

The potential between a nucleon and an antinucleon can be obtained from that between two nucleons by the "G-parity" rule, as illustrated in Fig. 3. Notice that since ω has $G = -$, the potential due to ω exchange though repulsive in NN becomes attractive in $N\bar{N}$. Together with the 2π exchange term which remains attractive here, one has then a very strong attraction at intermediate ranges. As example, we show in Fig. 4 the central potential calculated in this way by the Paris group.

The interaction in $N\bar{N}$, however, is complicated by the presence of annihilation channels as illustrated in Fig. 5. The range of annihilation forces is expected naïvely from the mass of the two nucleons exchanged to be $\sim 1/2m_N \sim 0.1$ fermi, i.e., well below the region of validity of the potential in Fig. 3, but they are likely to be very strong. In any case, because of relativistic effects at short distances, one cannot hope to describe them by a potential. For this reason, they remain largely unknown, and, as we shall see, are the main source of uncertainty in the nuclear model for baryonium.

We note in particular the following points relevant for baryonium spectroscopy.

Fig. 3 The nucleon-antinucleon potential due to particle exchange.

Fig. 4 The central component of the "Paris" NN potential.

Fig. 5 Annihilation effects in nucleon-antinucleon interactions.

A(i) The strong attraction at intermediate ranges can give rise to many bound states and resonances, but

A(ii) Some (or all?) may be very unstable because of strong annihilations at short distances.

A(iii) At large distances, the potential vanishes exponentially,
 $V \sim \exp [-m_\pi r]$.

A(iv) Since the forces arise from exchange of particles carrying isospin, they are strongly isospin dependent.

A(v) They are also strongly spin dependent due especially to the tensor force.

B) The idea of confinement forces comes of course from the seemingly absurd assumption that quarks exist but are not seen. The only reason why one should entertain such a conjecture at all seriously is that there is now a quite attractive theory for strong interaction which makes confinement somewhat plausible. This theory, called quantum chromodynamics (QCD) is modelled on QED and based on a new quantity called "colour" which takes the place of the electric charge.

The concept of colour was first proposed to resolve an inconsistency in the quark model of hadron spectroscopy. Take, for example, Δ^{++} with $J_z = 3/2$ which in the quark model is a state consisting of uuu all in an s-wave and all with spin up. The total wave function is then symmetric under the interchange of any two quarks, whereas quarks being spin $\frac{1}{2}$ fermions, ought to have antisymmetric wave functions. One can bypass this difficulty by assuming that quarks have an additional degree of freedom, namely colour, which can take three values. One further assumes that quark interactions are invariant under the colour group $SU(3)$ in which quarks belong to the fundamental representation 3, and hadrons to the identity 1. The wave function of the $\Delta^{++}(J_z = 3/2)$ considered above has now to be multiplied by a colour wave function: $3 \times 3 \times 3 = 1$ which is antisymmetric, making thus the total wave function also antisymmetric.

At first this assumption looked unattractive and artificial, but gradually it was realized that one can now construct a theory of quark interactions in close analogy to QED. In electrodynamics, electrons interact by exchanging photons; analogously, quark interactions in chromodynamics are mediated by vector particles called

gluons which belong to the adjoint representation 8 and carry colour charges. Further, QED is invariant under local phase transformations $\exp i \Lambda(x)$; similarly, QCD has local gauge invariance under transformations belonging to the group $SU(3)$.

The resultant theory has the following attractive features:

- i) It is renormalizable.
- ii) Due to renormalization effects, the effective coupling α_s decreases with q^2 (asymptotic freedom) so that at high q^2 (i.e., short distances) quarks are effectively free inside hadrons, as suggested by the parton model of deep inelastic scattering.
- iii) Conversely, for low q^2 (i.e., large distances) α_s becomes large (infra-red slavery) so that there is a possibility that quarks may indeed be confined.

Unfortunately, the theory is complicated and one has not yet learned to calculate with it many observable quantities so that at present it has few direct experimental tests. This is particularly true for the quark confinement forces which interest us. Nevertheless, a number of their properties have been conjectured based partly on theoretical arguments and partly on phenomenological observations. We list some of their properties below which are relevant for baryonium spectroscopy.

B(i) Confinement forces are by assumption attractive only in colour singlet states.

B(ii) At large distances, the confining potential increases indefinitely (infra-red slavery), e.g., $V(r) \sim a|r|$, which means that it can lead to an infinite number of bound states and resonances with unlimited angular momentum. This has some support from the empirical hadron spectrum which was indeed at least partly responsible for such a conjecture.

B(iii) Since gluons carry no flavour, confinement forces due to their exchange are independent of isospin so that

for example, the meson spectrum is expected to be isospin degenerate unless modified by other effects such as quark-pair annihilations which are isospin dependent. This also has some support from experiment for resonances with high angular momentum for which the short-ranged annihilation forces are presumably negligible.

B(iv) Confinement forces are probably spin-independent at large distances. Phenomenologically this may be demonstrated as follows. In the quark model, π and ρ are both s-wave $q\bar{q}$ states ($q = u, d$), differing only by the total quark spin S : for π , $S = 0$ (singlet) for ρ , $S = 1$ (triplet). The mass difference $\Delta m(0) = m_\rho - m_\pi$ measures therefore the strength of the spin-dependent interactions between q and $\bar{q}$ in the s-wave, i.e., $\ell = 0$. Similarly, B and A_2 are p-wave $q\bar{q}$ states differing only by the total quark spin S . $\Delta m(1) = m_{A_2} - m_B$ measures the strength of the spin-dependent interaction between q and $\bar{q}$ for $\ell = 1$. Hence, plotting the mass difference Δm between the triplet and singlet states along Regge trajectories gives the variation of spin-dependent forces as functions of ℓ , or, equivalently, of the distance between q and $\bar{q}$. Figure 6 shows such a plot for the $q\bar{q}$ ($q = u, d$) and $q\bar{s}$ mesons. In both cases Δm is seen to decrease rapidly with ℓ .

These properties are very different from those of nuclear forces listed before and the differences will be reflected in the bound state and resonance spectrum. For instance:

i) Nuclear forces are short ranged and cannot bind resonances with high orbital angular momentum, whereas confinement forces are supposedly under no such restrictions. Hence if there exist high mass baryoniums with very high spin, then they are most likely $QQ\bar{Q}\bar{Q}$ hadrons rather than NN nuclei.

ii) Nuclear forces are strongly isospin dependent so that any isospin degeneracy in the spectrum of their bound and resonant states is purely accidental, whereas for confinement forces such a degeneracy is, in general, to be expected. The empirical observation (d) in §1 is therefore a point in favour of the $QQ\bar{Q}\bar{Q}$ hadron model.

Fig. 6 The variation of spin dependent $Q\bar{Q}$ interactions as a function of the orbital angular momentum ℓ .

In order to distinguish further the two models, we shall have to calculate the spectrum of baryoniums in each case and to compare in detail the result with experiment.

3. BARYONIUMS AS NUCLEAR ATOMS

As the $N\bar{N}$ potential is well understood only for $r \gtrsim 0.8$ fermi some assumption has to be made for short distances before one can proceed to a calculation of the resonance spectrum. On this point there are some differences between various authors in this field, but hopefully such differences are not of vital importance for the states of interest with appreciable orbital angular momentum ℓ . For example, the Paris group assumed simply³⁾:

$$U = \begin{cases} U^P(r) & r \geq r_c \\ U^P(r_c) & r < r_c \end{cases}$$

where $U^P(r)$ is the "Paris potential" the central part of which was shown in Fig. 4, and r_c is treated as a free parameter expected to be ~ 0.8 fermi. With such an assumption, the calculation of the spectrum is then explicit.

The result for $r_c = 0.7$ fermi is quoted in Table 5 as an example. Because of the strong attraction at intermediate ranges there are many states both below and above the NN threshold. One notes however the following:

- i) The spectrum is quite sensitive to r_c although the ordering of the states is not.
- ii) The masses and even the existence of the tightly bound states are unreliable because of the strong annihilation effects at short distances.
- iii) States far above the NN threshold are probably too unstable to be regarded as resonances.

Experimentally, there are indeed a large number of candidates for baryoniums reported near the NN threshold. Unfortunately, none of them has yet had its quantum numbers identified so that no detailed comparison is at present possible. It is obvious however that all states obtained in this way are strongly coupled to NN and cannot be associated with the narrow states high above the NN threshold found in production experiments as discussed in (b)§1. There are some theoretical suggestions that such states may still be nuclear bound states not of NN but of nucleon resonances such as $\Delta\bar{\Delta}$. However, since nucleon resonances themselves decay thus: $\Delta \rightarrow N\pi$, with widths ~ 100 MeV, narrow states obtained in this way cannot be a general occurrence.

Table 5

$2I+1, 2S+1L_J$	Binding energy (MeV)	Mass (MeV)	$2I+1, 2S+1L_J$	Binding energy (MeV)	Mass (MeV)
$^3S_0$	-262	1616	$^3P_0$ (n = 0) $^3P_1$ (n = 1)	< -800 -102	< 1078 1776
$^3P_1$	-55	1823	$^3S_1 - ^3D_1$ (n = 0)	< -800	< 1078
$^3P_1$	-116	1762	$^3S_1 - ^3D_1$ (n = 1)	-39	1839
$^3S_1 - ^3D_1$	-202	1676	$^3P_2 - ^3F_2$ (n = 0)	-788	1090
$^3D_2$	165	2043	$^3D_3 - ^3G_3$	-457	1421
$^3D_2$	85	1963	$^3F_4 - ^3H_4$	-89	1789
$^3F_3$	> 200	> 2078	$^3G_5 - ^3I_5$	307	2185

Cf., Experimental states: 1394, 1648, 1680, 1794, 1897, 1936, 1940, 2185 ...

4. BARYONIUMS AS COLOUR MOLECULES^{4),5)}

We shall denote the $QQ\bar{Q}\bar{Q}$ hadron states of interest by

$$(QQ)_{2S+1}^X \xrightarrow{L} (\bar{Q}\bar{Q})_{2\bar{S}+1}^{\bar{X}} \quad (1)$$

where (QQ) denotes a QQ system in an s-wave state (diquark) and X = its total colour representation and S = its total spin. Clearly X can be $3 \times 3 = \bar{3}$ or 6 and $S = 0$ or 1. Also since the hadron (1) is a colour singlet, $\bar{X}$ must be the conjugate representation of X .

In general, states with different diquark colours and spins can mix since there is a priori no reason why the total Hamiltonian should be diagonal in the diquark colour or spin. One believes however that such mixing becomes negligible when L is large. The arguments go as follows: Since quarks are fermions, the diquark wave function must be totally antisymmetric. Hence for $q = u, d$ one must have the following isospin assignments for our diquark states:

$$\begin{array}{cccc} (qq)_1^{\bar{3}} & (qq)_3^{\bar{3}} & (qq)_1^6 & (qq)_3^6 \\ I = 0 & I = 1 & I = \frac{1}{2} & I = 0 \end{array} \quad (2)$$

Now the gluons which mediate interactions between the diquark and antidiquark carry no flavour and can therefore mix only those states with the same isospin, i.e.,

$$(qq)_1^{\bar{3}} \longleftrightarrow (qq)_3^6 ; (qq)_3^{\bar{3}} \longleftrightarrow (qq)_1^6 \quad (3)$$

In other words, to change the colour of a diquark one must also change its spin. However, by (B)(iv) in §2, spin dependent forces decrease rapidly with the separation L so that for large L the mixing between states with different X and S becomes negligible. We may then regard states with different X and S as essentially distinct.

At this stage it is convenient to draw an analogy between our hadron states (1) with ionic molecules such as $\text{Na}^+ - \text{Cl}^-$, where two ions of opposite charges are neutralized by an ionic bond linking them. Here we have two quark clusters with opposite colour charges ("chromions") neutralized by a tube of colour flux lines ("chromionic bond") and linked together to form a ("chromionic") colour molecule. The analogy is clearly only formal since the dynamics is very different, but it teaches us the wisdom of the chemists in treating ions and bonds as separate entities thereby reducing labour and simplifying concepts when considering the properties of molecules. We shall do the same for our colour equivalents.

Consider then first the diquark chromions. Since they cannot exist free, they cannot be assigned a mass, but one can still talk about mass splittings between, for example, their various spin states. Since spin-dependent effects are supposed to be short-ranged, B(iv), one might argue by asymptotic freedom that perturbation theory applies so that to leading order in α_s the "hyperfine splitting" between diquarks is given just by the one-gluon exchange diagram of Fig. 7. This assumption has some support in the spectroscopy of $Q\bar{Q}$ mesons and QQQ baryons. Now the diagram

Fig. 7 The one-gluon exchange term in quark-quark interactions

is very similar to one-photon-exchange between electrons in QED except that the electric charge is now replaced by the colour charge λ_i which is a Gell-Mann 3×3 matrix of SU(3). By means of a Breit reduction as in QED, one obtains from Fig. 7 a spin-spin, a spin-orbit, and a tensor term, of which only the spin-spin term is relevant for our diquark in s-wave. This can be written as

$$H_1 = -C \sum_{a=1}^8 \sum_{r=1}^3 (\lambda_i^a \sigma_i^r)(\lambda_j^a \sigma_j^r) \quad (4)$$

where σ_i are Pauli matrices for the quark spins. The coefficient C is essentially $\alpha_s |\psi(0)|^2$ where ψ is the relative diquark wave function. Since the details of confinement forces are unknown, C cannot be calculated in a model-independent manner, but it can be estimated phenomenologically from QQQ baryon spectroscopy, for example: $\Delta-N = 16 C \sim 300 \text{ MeV}$ gives $C \sim 20 \text{ MeV}$. The hyperfine splitting between diquark ions can then be obtained by diagonalizing H_1 , which is easily done and the result is listed in Table 6.

Table 6

Diquark	$\langle H_1 \rangle$
$(Q_1 Q_2)_1^{\bar{3}}$	$-8 C$
$(Q_1 Q_2)_3^{\bar{3}}$	$\frac{8}{3} C$
$(Q_1 Q_2)_1^6$	$4 C$
$(Q_1 Q_2)_3^6$	$-\frac{4}{3} C$

Next, the chromionic bond also carries a mass by virtue of the colour field energy confined inside the flux tube. Unfortunately this field energy cannot be calculated without an explicit model for confinement, which at present means essentially the MIT bag model. The result is best stated in terms of the angular momentum L carried by the bond, which is asymptotically linear in the bond mass squared, thus

$$L = L_0 + \alpha' M_B^2 \quad (5)$$

where both L and α' depend on the colour X . The Regge slope α'_X is found to be inversely proportional to the square root of the quadratic Casimir operator $\mathcal{C}_X$, thus

$$\alpha'_X \sim 1 / \sqrt{\mathcal{C}_X} \quad (6)$$

which, when combined with the information $\alpha'_3 \sim 0.9 \text{ GeV}^{-2}$ obtained phenomenologically from, for example, the ρ trajectory, gives α'_X for other colours X also, e.g., $\alpha'_8 \sim 0.6 \text{ GeV}^{-2}$. The intercept L_0 however is largely unknown and has to be determined phenomenologically from the baryonium spectrum.

Combining (5) with Table 6, one then has the following mass formula for the colour molecules (1):

$$M = \langle H_1 \rangle_{qq} + \langle H_1 \rangle_{\bar{q}\bar{q}} + \sqrt{(L - L_0^X) / \alpha'_X} \quad (7)$$

depending on one parameter L_0^X for each family with colour X . One sees then that in spite of our ignorance of confinement forces compared with nuclear forces, one has managed by various means to construct a spectrum for $QQ\bar{Q}\bar{Q}$ hadrons quite as detailed as for the NN nuclei of Table 5.

Before we examine the predicted spectrum and compare it with experiment, let us first consider the decay properties of our colour molecules. By construction, all the colour molecules (1) are inhibited in decaying into mesonic final states by virtue of the angular momentum barrier. Their decays into $B\bar{B}$ pairs, however, depends on the colour X of the diquark. For $X = 3$, the molecule can decay by rupturing the colour bond linking the two chromions. Since colour is confined, a $Q\bar{Q}$ pair has to be created so as to neutralize the colour charge exposed. The result is a $B\bar{B}$ pair, thus:

$$(QQ)^3 \not\rightarrow (\bar{Q}\bar{Q})^3 \longrightarrow ((QQ)^3 Q^3)' + ((\bar{Q}\bar{Q})^3 \bar{Q}^3)' = B + \bar{B} \quad (8)$$

For $X = 6$, however, the creation of one $Q\bar{Q}$ pair is insufficient to neutralize the colour 6 exposed on rupture of the bond. Indeed, $6 \times 3 = 8 + 10$ does not contain a colour singlet. At least two $Q\bar{Q}$ pairs have to be created in order to neutralize the colour, thus:

$$(QQ)^6 \not\rightarrow (\bar{Q}\bar{Q})^6 \longrightarrow ((QQ)^6 (\bar{Q}\bar{Q})^6)' + ((\bar{Q}\bar{Q})^6 (QQ)^6)' \quad (9)$$

and the result is not $B\bar{B}$ but two baryoniums. Therefore, for decaying into $B\bar{B}$, a molecule with $X = 6$ has first to mix with $X = 3$ and then decay. Since the mixing effect is believed to be strongly suppressed for large L , the coupling to $B\bar{B}$ of $X = 6$ molecules has to be suppressed accordingly.

We therefore distinguish two families as candidates for baryoniums. By virtue of the large coupling to $B\bar{B}$ colour molecules with $X = 3$ acquire large widths above the $B\bar{B}$ threshold. Their decay into $B\bar{B}$ is similar in mechanism to the decay of, e.g., $f \rightarrow \pi\pi$ and is expected to have a similar width of order 100 MeV. For the same reason they are also readily formed in the direct channel in $N\bar{N}$

collisions. On the other hand, molecules with $X = 6$ couple reluctantly to $B\bar{B}$ as well as to mesons and therefore remain narrow even above the $B\bar{B}$ threshold. In addition, they cannot easily be formed in $N\bar{N}$ collisions but have to be produced along with other particles. This is strikingly reminiscent of the experimental observations in § 1.

In Figs. 8 and 9 we quote the spectrum of molecules with $X = 3$ and 6 as calculated in Ref. 4) by normalizing L_0 for each X to an empirical baryonium state, and compare them with the experimental structures observed. Further, in Table 7, the predicted quantum numbers of the dominant $X = 3$ states are compared with whatever is known experimentally. There is a rather close correspondence between the experimental and predicted spectra, especially for the $X = 3$ sector.

Fig. 8 The predicted spectrum of baryoniums in the colour molecular model with $X = 3$ compared with experimental structures in formation experiments.

Fig. 9 The predicted spectrum of baryoniums in the colour molecular model with $X = 6$ compared with experimental structures observed in production experiments.

7. NUCLEAR ATOMS OR COLOUR MOLECULES?

It should be emphasized once more that the experimental evidence for baryonium states is still far from established especially for those belonging to the narrow family found in production experiments. It may well happen that under scrutiny, most, if not all, will disappear. Under such circumstances it would be unwise at present to conclude definitively in favour of either alternative. Superficially, however, at least to my perhaps biased mind, the existing experimental picture seems to favour slightly the colour molecules.

However, the two pictures need not be mutually exclusive and $N\bar{N}$ nuclear atoms could co-exist with $QQQ\bar{Q}$ colour molecules. We have

Table 7

Energy Region		S		T		U		V	
Traj.	L	$(qq)_1^3 - (\bar{q}\bar{q})_1^3$	$(qq)_3^3 - (\bar{q}\bar{q})_3^3$	$(qq)_1^3 - (\bar{q}\bar{q})_1^3$	$(qq)_3^3 - (\bar{q}\bar{q})_3^3$	$(qq)_1^3 - (\bar{q}\bar{q})_1^3$	$(qq)_3^3 - (\bar{q}\bar{q})_3^3$	$(qq)_1^3 - (\bar{q}\bar{q})_1^3$	$(qq)_3^3 - (\bar{q}\bar{q})_3^3$
		2	0	3	1	4	2	5	3
σ_{el}	$I_{(J^P)}^G$	$0^+(2^+)$	$1^-(2^+)$	$0^-(3^-)$	$1^+(3^-)$	$0^+(4^+)$	$0^+(4^+)$	$0^-(5^-)$	$1^+(5^-)$
	Mass	1.92	1.94	2.13	2.11	2.33	2.34	2.51	2.56
$\sigma_{ch.ex.}$	$I_{(J^P)}^G$	$J^{PC} = 2^{++}$	$J^{PC} = 2^{++}$ or 4^{++}	$I = 0, 1$		$I = 0, 1$		$I = 0, 1$	
	Mass	1.936	1.9 - 2.0	2.155		2.345		2.5	
$\bar{p}p \rightarrow \pi^+ \pi^-$	$I_{(J^P)}^G$	—	$[0^+(2^+)]$	$1^+(3^-)$		$0^+(4^+)$		$1^+(5^-)$	
	Mass	—	$[2.0]$	2.150		2.310		2.5	
$\bar{p}p \rightarrow \pi^0 \pi^0$	$I_{(J^P)}^G$	—	—	$0^+(4^+)$		$0^+(4^+)$		$1^+(5^-)$	
	Mass	—	—	2.35		2.35		$1^+(5^-)$	
$\bar{p}p \rightarrow K^+ K^-$	$I_{(J^P)}^G$	—	—	$0^-(3^-)$		$1^+(4^+)$		$0^-(5^-)$	
	Mass	—	—	2.15		2.34		2.35	
$\bar{p}p \rightarrow \pi^0 \eta$	$I_{(J^P)}^G$	—	—	$1^-(4^+)$		$1^-(4^+)$		$1^+(5^-)$	
	Mass	—	—	2.32		2.32		2.5	
$\bar{p}p \rightarrow \rho^0 \pi^0$	$I_{(J^P)}^G$	—	$1^-(2^+)$	$1^+(3^-)$		$0^+(4^+)$		$0^-(5^-)$	
	Mass	—	1.940	2.15		2.34		2.5	
SUM TOTAL	$I_{(J^P)}^G$	$0^+(2^+)$	$1^-(2^+)$	$0^-(3^-)$		$0^+(4^+)$		$0^-(5^-)$	
	Mass	1.936	~ 2.0	2.15		2.34		2.5	

deliberately taken a controversial approach here only for the sake of economy. Also, the view has sometimes been expressed that the two pictures may be just different descriptions of the same thing. I do not understand this viewpoint. It is of course true that since nucleons are made of quarks, an NN nucleus can also be described by a multiquark wave function. Conversely, so long as quarks are confined, a multiquark state can also be expressed as a composite of the hadrons from which it is formed and into which it decays. In other words, there has to be an equivalence transformation relating the two state vector spaces in terms of quarks on the one hand and of hadrons on the other. However, the baryonium states we are considering are not just any states in the state vector space. They have well-defined structures which are very different for the NN nuclei and the $QQ\bar{Q}\bar{Q}$ hadrons, so that their overlap in either picture must be very small. It just so happens that the NN nucleus is easy to describe in the hadron representation and the $QQ\bar{Q}\bar{Q}$ hadrons in the other. But in either representation they can be regarded as different states for all practical purposes. After all, a helium atom contains the same number of nucleons and electrons as a deuterium molecule, and both can be described in terms of a four-nucleon four-electron wave function, but no one has yet claimed that they are just different descriptions of the same thing.

In any case, both the models are sufficiently explicit that when the data clarify, one should be able to determine which, if any, of the pictures is correct.

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Further details can be found in the literature referred to in these papers.