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NARROW BARYONIUM STATES IN QCDH.G. Dosch^{*)} and M.G. Schmidt

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ABSTRACT

The unconventional states of $qq\bar{q}\bar{q}$, $q\bar{q}$ and pure gluon type, appearing in $B\bar{B}$ dual diagrams with non-planar baryons are discussed, treating their interpolating fields in a confinement approach to QCD due to Migdal. The mass spectrum is given to lowest order and the quantum numbers are discussed. It is argued that some of the resonances should have narrow width and hence could be identified with the narrow resonances in $p\bar{p}$ scattering.

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The observation of narrow ($\Gamma \sim 3-20$ MeV) resonances in $B\bar{B}$ scattering¹⁾ has renewed the interest in unconventional meson states. It is an attractive idea that the binding of such states is closely related to that of the baryons²⁾. Baryons are expected to be non-planar objects in space time in the (non Abelian) quark-gluon gauge theory of strong interactions (QCD). The three quarks of a baryon are argued to be confined by three strings -- realized by gauge gluons -- which are coupled together in a central region antisymmetrically in the colour indices. The $1/N_c$ topological expansion³⁾ ($N_c = 3$ in QCD) breaks down in that central space-time tube as elaborated by Rossi and Veneziano in Ref. 4).

Dual diagrams first introduced as a convenient means to handle group theory, obtain a dynamical interpretation in the string picture. In $B\bar{B}$ scattering with triple string baryons there are then two types⁴⁾ of duality diagrams instead of one conventional diagram (the latter one leads to inconsistencies, as is well known) (Fig. 1). In graphs 1a and 1b there appear unconventional meson states of type $qq\bar{q}\bar{q}$ and $q\bar{q}$; graph 1c contains a pure glue state.

In this note we present estimates for the mass spectra and quantum numbers of these three new types of states as well as some considerations about their decay properties. This is done by constructing gauge invariant local field operators interpolating these states and treating these operators in an infrared regularization procedure recently proposed by Migdal⁵⁾.

The essential idea of this approach is to extract in a unique way a set of particle poles from vacuum expectation values of time-ordered products of gauge invariant local field operators calculated in QCD perturbation theory (in g^2). Confinement, not present in finite order perturbation theory, is enforced by a Padé approximation in the Mandelstam variables, generalized to tensor functions around a point Λ in the deep Euclidean region where asymptotic freedom justifies perturbation theory. The Padé approximants of order (M,N) turn out to have a unique limit for $M,N,\Lambda \rightarrow \infty$ if $(MN/\Lambda^2) = R^2$ is kept fixed in a fixed order of perturbation theory. R (a kind of "bag radius") is supposed to be pushed to infinity in higher orders in g^2 if the parameters (R,g^2) are fixed by two experimental mass ratios.

In zeroth order the QCD structure does not enter explicitly, but only as a justification of the method starting from asymptotic freedom. The mass spectrum appearing in the two-point functions of local gauge invariant operators of dimension d is in this approach given by the zeros of the Bessel function

$$J_{d-2}(2RM) = 0 \quad (1)$$

The direct influence of the gluon interaction will be discussed later. The parameter R fitted by Migdal⁵⁾ to the meson spectra (in particular the vector-meson trajectories), also gives satisfactory results for the baryon spectra⁶⁾ and will be used here to estimate the spectra of the unconventional states.

Examples of local field operators corresponding to the states M_4^2 and M_2^2 of Figs. 1a and 1b are:

$$\phi_1(x) = N(\bar{\psi}_{a\kappa}(x) \bar{\psi}_{b\lambda}(x) \overleftrightarrow{\mathcal{D}}_{\mu_1}^c d_1 \dots \overleftrightarrow{\mathcal{D}}_{\mu_n}^{d_n} c') \quad (2a)$$

$$\psi_{\kappa'}^{a'}(x) \psi_{\lambda'}^{b'}(x) \varepsilon^{abc} \varepsilon_{a'b'c'} \Gamma_{\kappa\lambda; \kappa'\lambda'}$$

$$\phi_2(x) = N(\bar{\psi}_{a,\kappa}(x) \overleftrightarrow{\mathcal{D}}_{\mu_1 b}^{d_1} \dots \overleftrightarrow{\mathcal{D}}_{\mu_n d_n}^{b'} \mathcal{D}_{\mu_{n+1} c}^{d_{n+1}} \dots \mathcal{D}_{\mu_n d_n}^{c'} \psi_{\kappa'}^{a'}(x)) \varepsilon^{abc} \varepsilon_{a'b'c'} \Gamma_{\kappa\kappa'} \quad (2b)$$

Pure gluon operators have been discussed in Ref. 5). Here κ, λ denote spinor and flavour indices

$$\mathcal{D}_{\mu c}^{c'} = \frac{\partial}{\partial x^\mu} \delta_c^{c'} + ig_2 \vec{A}_\mu \vec{\lambda}_c^{c'} \quad (3)$$

N indicates normal product and normalization factors. ψ is the quark and $\vec{A}_\mu$ the gauge gluon field. In ϕ_1 and ϕ_2 there is also the possibility to attach covariant derivatives to the spinors before contracting with G (corresponding to string pieces directly connected to the quarks).

A great manifold of the states of type ϕ_1 and ϕ_2 is coupled to zeroth order in g (i.e., by overlap) to the conventional meson states^{*}). The operators of lowest dimensionality of type ϕ_1 not coupled to zeroth order to conventional states^{**}) are

$$\phi_1^0(x) = N(\bar{\psi}_{a\kappa}(x) \bar{\psi}_{b\lambda}(x) F_{\mu\nu}(x)_c^{c'} \psi_{\kappa'}^{a'}(x) \psi_{\lambda'}^{b'}(x)) \varepsilon^{abc} \varepsilon_{a'b'c'} \Gamma_{\kappa\lambda; \kappa'\lambda'} \quad (4a)$$

and of type ϕ_2

$$\phi_2^0(x) = N(\bar{\psi}_{a\kappa}(x) F_{\mu\nu}(x)_b^{b'} F_{\rho\sigma}(x)_c^{c'} \psi_{\kappa'}^{a'}(x)) \varepsilon^{abc} \varepsilon_{a'b'c'} \Gamma_{\kappa; \kappa'} \quad (4b)$$

The gluon field tensors $F_{\mu\nu}$ arise from the antisymmetrization of covariant derivatives (3). For the spinor part of $\Gamma_{\kappa\lambda; \kappa'\lambda'}$ we choose the form (without loss in generality)

$$\Gamma_{\kappa\lambda; \kappa'\lambda'} = \Gamma_{\kappa\lambda}^{(1)} C^* C^T \Gamma_{\kappa'\lambda'}^{(2)} \quad (5)$$

with the 16 Dirac matrices Γ . The most prominent s wave couplings correspond to $\Gamma^{(i)} = \gamma_5, \gamma_\mu$ (see the Table).

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*) These states are not supposed to be narrow. Those of type ϕ_1 - starting from

$$N \bar{\psi}_{a\kappa}(x) \bar{\psi}_{b\lambda}(x) \psi_{\kappa'}^{a'}(x) \psi_{\lambda'}^{b'}(x) \varepsilon^{abc} \varepsilon_{a'b'c'}$$

could be used in order to explain the broad bumps in the T, u region, similar to Ref. 2). Their lowest dimension is $d = 6$ (with maximal spin 2), and hence the T and u region would be reached by the third and fourth recurrence, respectively.

**) They are, however, coupled to daughters of highly excited planar meson states with $F_{\mu\nu}$ coupling. We do not consider these as conventional.

The simplest pure gluon operator with $\varepsilon\text{-}\varepsilon$ structure is given by

$$\phi_3^0 = F_{\mu\nu}(x)_a^{a'} F_{\mu'\nu'}(x)_b^{b'} F_{\mu''\nu''}(x)_c^{c'} \varepsilon^{a'b'c'} \varepsilon^{abc} \quad (4c)$$

With massless quarks, all the states of the same type (4), according to Eq. (1) have to 0th order the same mass, determined by the dimensions of $\phi_1^{(0)}$, $\phi_2^{(0)}$ and $\phi_3^{(0)}$, $d_1 = 8$, $d_2 = 7$ and $d_3 = 6$. The first non-trivial zero of the "Bessel function" yields the masses

$$\begin{aligned} M_1^0 &= 1990 \text{ MeV} \\ M_2^0 &= 1760 \text{ MeV} \\ M_3^0 &= 1520 \text{ MeV} \end{aligned} \quad (6)$$

Here we have used $R = 2.500 \text{ GeV}^{-1}$, fitted in Ref. 5) in order to give the correct ρ mass [for comparison: the lowest nucleon mass (N, Δ) with this value of R is 1100 MeV, the $N\text{-}\Delta$ splitting coming out correctly with the gluon corrections]. The zeroth order trajectories are given in Fig. 2.

The influence of gluon interaction up to order g^2 - for normal mesons⁵⁾ and baryons⁶⁾ simply to be taken into account by adding in Eq.(1) anomalous dimensions of the operators - is more complicated in the case considered here. There are transitions between operators of types and (1) and (2) to order g and also decays into conventional mesons to that order not suppressed in an $1/N_c$ expansion. We hence give only some plausible arguments on the effects of gluon exchange.

- 1) The combined self-interaction of the fermions and gluons will decrease d and hence the mass. The anomalous dimension (in Feynman gauge) of the gluon is $\gamma_g = -13/3 \cdot g^2/8\pi^2$; of the fermion $\gamma_F = 2/3 \cdot g^2/8\pi^2$; $g^2/8\pi^2 = 0.31$ fitted in Ref. 5) to the ρ trajectory.
- 2) Gluon exchange inside the $qq(\bar{q}\bar{q})$ system decreases the dimension most for $\Gamma_4^{(i)} = \gamma^5$, less for γ^μ . This leads to a splitting between a pseudoscalar and vector diquark system of about 100 MeV.

Mass corrections to the quarks make s wave (γ_5, γ_1) diquarks lighter than p wave $(1, \gamma_5 \gamma_1)$ diquarks. Hence the states listed in Table 1 should be the lightest states of type $qq\bar{q}\bar{q}$ (with an increase in mass from above to below).

Since in states of types ϕ_2 and ϕ_3 , two and three gluons, respectively, are present, the mass of these states will be lowered considerably due to the large negative anomalous dimension of the gluon.

We now present some arguments on the decay width of these states. In order to calculate these widths in accordance with the regularization scheme of Ref. 5) one had to apply the Padé procedure to the N point functions $\langle T\phi_i, \phi_{M1} \dots \phi_{Mn} \rangle$ governing the decay. Here we argue that the three-point functions of some of the fields of type ϕ_1^0 and interpolating fields of the most prominent mesons vanish up to order g^2 . This is due to a cancellation of the contribution of the (divergent) diagram 3a with that of b. In the last column of Table 1 we indicate the three-point functions which vanish up to order g^2 due to this cancellation or trace identities. Furthermore, by the same argument the 5,7, ..., point functions of the operators listed in the Table with scalar or pseudoscalar fields vanish up to order g^2 , thus forbidding to this order, e.g., $4\pi, 6\pi \dots$, decay.

The two point function of $\phi_1^{(0)}$ with conventional meson fields (diagram 3c) vanishes up to order g^2 because of symmetric integration.

The decay (and production) of states corresponding to ϕ_1 in channels with a baryon-antibaryon pair is allowed to order g and hence should compete with mesonic decay modes in spite of the smaller phase space.

By ordering according to g we have assumed this effective coupling constant to be still small, since the soft gluons are taken care of already in the zeroth order calculation of this procedure⁵⁾. Note that the numerical value of R used here is adopted to a calculation up to g^2 .

The decay and production amplitude of states corresponding to ϕ_2 and ϕ_3 is of order g^2 and g^3 , respectively. In our approach based on low order perturbation theory we do not see the suppression of purely mesonic channels as compared to baryonic ones as proposed in Ref. 4).

It is tempting to attribute the observed narrow resonances in the 1900 MeV region to $\phi_1^{(0)}$ type states since - following the above arguments - these are coupled stronger to $B\bar{B}$ than ϕ_2 and ϕ_3 type states and since their mass is also in the range predicted [e.g., Eq. (6)]. If one also wants to associate the 1795 NN bound state with the $\phi_1^{(0)}$ type one would be urged to identify this state with one of the (γ_5, γ_5) species and the 1935 MeV states with the (γ^μ, γ^μ) species since they both have decay channels suppressed in a similar way. There are enough candidates to have several states in the 1935 MeV region. An analysis of quantum numbers of the states and of their decay channels of course would make an identification according to the Table much more stringent.

There are also further states of predicted narrow width of ϕ_2 and ϕ_3 type with masses below the $p\bar{p}$ threshold. Altogether there is a vast number of new states which wait for an experimental check. The pure glue state is predicted to have the lowest mass of the baryonium ($\epsilon\epsilon$ type) states. It is supposed to couple weakly only as a resonance in $B\bar{B}$, but this does not imply a weak coupling as a "Regge" trajectory in $B\bar{B}$ scattering according to common ideas about the Zweig rule; and hence the proposal⁴⁾ that this type gives the leading Regge trajectory dual to leading three-meson (jet) production is not contradicted.

In our discussion of the lowest baryonium states no arguments about highly excited states on leading straight trajectories corresponding to stretched strings enter. On the contrary, the quarks and gluons seem to sit together as close as possible. Narrow-width effects seem to arise -- not from a long straight string but from "string nodes" $F_{\mu\nu}$.

Therefore a determination of the angular momenta would be especially sensitive to the ideas presented here: all the resonances could have small angular momentum (see the Table) without losing their narrow width.

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Figure captions

- Fig. 1 : Dual diagrams for $B\bar{B}$ scattering and string states corresponding to the dot-lined cross-section $[M_4^2, M_2^2, M_0^2, \text{respectively in the notation of Ref. 4)]$. The solid lines in the dual diagram represent quarks, the dashed ones represent non-planarity lines.
- Fig. 2 : Trajectories with maximal spin ($J_{\max}$) for unconventional meson states: $qq\bar{q}\bar{q}$ (ϕ_1), $q\bar{q}$ (ϕ_2), pure glue (ϕ_3).
- Fig. 3 : Diagrams for the n point function of the field $\phi_1^{(0)}$ with
(a),(b) two (conventional) meson fields $(M_{(1)}, M_{(2)})$
(c) one meson field
(d) two baryon fields.

Table

(Γ_1, Γ_2)	SU_3 flavour	$J^P; C_n$	$\Gamma_{M1} \Gamma_{M2}$
(γ^5, γ^5)	$1 \oplus 8$	$1^{\pm}; -$	$(\gamma^5, \gamma^5), (1, 1), (\gamma^{\mu}, \gamma^{\mu}), (\gamma^5 \gamma^{\mu}, \gamma^5 \gamma^{\mu})$ $(\gamma^5, \gamma^{\mu}), (1, \gamma^{\mu}), (\gamma^5, \gamma^5 \gamma^{\mu}), (1, \gamma^5 \gamma^{\mu})$
$(\gamma^5, \gamma^{\mu}) + (\gamma^{\mu}, \gamma^5)$	$8 \oplus (10 + \bar{10})$	$0, 1, 2^{\pm}; +$	$(\gamma^5, \gamma^{\mu}), (1, \gamma^{\mu})$ $(\gamma^5, \gamma^5), (1, \gamma^5), (1, 1), (\gamma^{\mu}, \gamma^{\mu}), (\gamma^5 \gamma^{\mu}, \gamma^5 \gamma^{\mu})$
$(\gamma^5, \gamma^{\mu}) - (\gamma^{\mu}, \gamma^5)$	$8 \oplus (10 - \bar{10})$	$0, 1, 2^{\pm}, -$	$(\gamma^5, \gamma^5 \gamma^{\mu}), (1, \gamma^5 \gamma^{\mu})$ $(\gamma^5, \gamma^5), (1, \gamma^5), (1, 1), (\gamma^{\mu}, \gamma^{\mu}), (\gamma^5 \gamma^{\mu}, \gamma^5 \gamma^{\mu})$
$(\gamma^{\mu}, \gamma^{\mu})$	$1 \oplus 8 \oplus 27$	$0, 1, 2, 3^{\pm}, -$	as for $(\Gamma_1, \Gamma_2) = (\gamma^5, \gamma^5)$

1
∞
1

First column : $\Gamma^{(1)}$ and $\Gamma^{(2)}$ of interpolating fields of type $\phi_i^{(0)}$ [see Eqs. (4) and (5)], which are argued to lead to the lowest masses of the $\bar{q}qqq$ system.

Second column : SU_3 flavour classification, spin, parity and charge conjugation of the neutral member. Note the parity doubling due to both parities of $F_{\mu\nu}$.

Third column : Matrices $\Gamma_{M1} \Gamma_{M2}$ for which the three-point function $\langle T \phi_i^{(0)} \phi_{M1} \phi_{M2} \rangle$ vanishes up to g^2 ; $\phi_M(x) = N \bar{\psi}(x) \psi \Gamma_{M1} \psi(x)$. For pseudoscalar, scalar, vector and axial vector fields we have $\Gamma_{M1} = \gamma_5, 1, \gamma^{\mu}$ and $\gamma^5 \gamma^{\mu}$, respectively.

(a)

(b)

(c)

FIG.1

FIG.2

FIG.3