SPOTTING GLUEBALLS IN COLLINEAR GLUON JETS

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We suggest looking for pure gluon hadronic states (glueballs) in $\Upsilon \to 3g \to$ low sphericity final state (collinear gluon jets). In these events one gluon has the maximum energy, $M_\Upsilon/2$, favouring fragmentation into a glueball. The signature for a true $C = G = +$glueball is its prominence at the Υ resonance in $e^+e^- \to \Upsilon \to$ (glueball $\to$ four charged pions) + ..., and its absence in $q\bar{q}$ jets off resonance (we do not expect significant quark fragmentation into glueballs).

Colored gauge gluons form an important part of quantum chromodynamics (QCD) [1] – the current best candidate for a theory of strong interactions. The existence of weakly and electromagnetically neutral glue in the nucleon is inferred from the "missing" longitudinal momentum in analyses of deep inelastic scattering [2]. There is no tangible evidence yet for the vector and nonabelian nature of gluons. The verification of three-gluon jets [3] in the hadronic decays of heavy $1\,^3S_1$ $Q\bar{Q}$ bound states J/ψ, Υ, ..., will confirm gluons and the gluon-mediated hadronic decay mechanism [4] of such states in QCD. However, that will shed no light on the self-couplings of gluons – the hall-mark of a nonabelian gauge theory. Gluon self-couplings intuitively suggest the existence of color singlet and flavor singlet hadronic glue states (glueballs [5]). The lowest mass glueball should be narrow because of the same mechanism that generates the Zweig [6] rule in QCD [‡1]. The detection and identification of glueballs would be evidence in favor of gluons and their self-interactions.

The intuitive anticipation of glueballs is strengthened by considerations of gluon bound states in dual dynamics and lattice QCD [7]. These calculations, as well as empirical evidence from radiative decays of J/ψ to $\gamma\eta$, $\gamma\eta'$, $\gamma\pi\pi$, $\gamma K\bar{K}$ [8] make it unlikely that a $C = +$glueball exists much below a mass of about 1.2 GeV. A probable mass value for the lowest $C = +$glueball is 1.5 GeV [‡2]. Its width is expected to lie in the range 1–20 MeV. Indeed, the higher energy domain of the inclusive photon spectrum from heavy $1\,^3S_1$ $Q\bar{Q}$ bound states such as J/ψ, Υ, etc. (via $Q\bar{Q} \to \gamma gg \to \gamma$ + hadrons) is an appropriate region for a glueball search, as recently emphasized by Koller and one of us (T.W.) [9]. This rate is down by a factor $36\alpha e_Q^2/5\alpha_s$ in QCD (our no-

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[‡1] The argument is given in the paper of Robson [5]. Note that the suppression of $gg \to q\bar{q}$ in the partial wave $L = 0$ arises because the process takes place at short distances for high gg or $q\bar{q}$ c.m. energy independent of the q mass.

[‡2] We choose this value because we expect the lightest glueball to be a scalar. A light $C = P = +$ state should have been seen via its $\pi\pi$ or $K\bar{K}$ mode in $J/\psi \to \gamma\pi^+\pi^-$ or γK^+K^- [8]. (Also, a $C = +$, $P = -$ state near η' would have produced a very large $\Gamma(J/\psi \to \eta'\gamma)/\Gamma(J/\psi \to \eta\gamma)$ ratio via mixing with η'.) The absence of such signals points to a fairly high mass for the lightest glueball. On the other hand a value much larger than 1.5 GeV would be difficult to interpret as due to a scalar daughter of the Pomeron [7].

tation has the usual meaning); the mass of the 2g system must also be small, leading to a strong suppression at the Υ. Nevertheless, the discovery of a glueball recoiling against a nearly monoenergetic γ ray in $J/\psi \to \gamma$ + hadrons is an exciting possibility – e.g. for the new Crystal Ball detector at SPEAR.

Here we suggest another fruitful way of looking for glueballs: study two (collinear) jet events in the hadronic decay of a very heavy unflavored $1\,^3S_1$ $Q\bar{Q}$ bound state produced in e^+e^- annihilation. The latter will have a substantial branching ratio into a glueball plus $q\bar{q}$ states. The glueball $\mathcal{G}$ (if $C = +$) will be identifiable as a narrow G-even multipion resonance. These decay pions will lie within one jet, since we expect the glueball to be hard (with high momentum). The signature of a glueball as opposed to a $q\bar{q}$ state is its presence (along with $q\bar{q}$ hadrons) in the decay products of the $Q\bar{Q}$ bound state with a large branching ratio and its absence at other values of the e^+e^- c.m. energy. Off resonance, the events are supposed to consist mainly of $q\bar{q}$ jets, and we expect that the probability for a quark to yield a glueball is small. It is certainly much smaller for a quark-initiated jet than for a gluon-initiated one. This is why glueballs have not been found in processes involving the former and should be looked for in those containing the latter.

Our argument for studying collinear jet events (those with low sphericity) is the following [‡3]. Consider the hadronic decay of the upsilon Υ(9.5), for instance. The collinear events lie near the triangular boundary of the jet Dalitz plot for $\Upsilon \to 3g \to 3$ jets. They thus correspond to one parent gluon having nearly the maximum energy with two less energetic ones opposite. This is advantageous for two reasons. First, it is easy to verify the parent gluon direction by a sphericity analysis. The jet from fragmentation of this gluon should contain a glueball. We expect this glueball to have a substantial fraction of the momentum of its parent gluon. Its decay products will lie near the jet axis, easing experimental identification by reducing the number of wrong mass combinations among the decay products (one need only take particles in one jet). Second, higher energy for the gluon should facilitate the fragmentation of a gluon into a massive glueball. For a given c.m. energy, the collinear configuration allows one gluon to be most energetic.

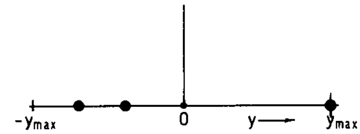

Fig. 1. Rapidity plot for typical collinear three-gluon configuration.

We expand on this last point. We view the lightest gg glueball as a two-constituent state, much as π or K is thought of as a $q\bar{q}$ meson. Consider the rapidity plot (y) of fig. 1 for the collinear case. The energetic single gluon is at one end of the rapidity axis. The two less energetic gluons are usually well separated. We suppose that hadrons are preferentially made when a leading constituent picks up one particle out of a $q\bar{q}$ or gg pair created out of the vacuum [9]. This is a consequence of short range order in rapidity. Hence either the energetic gluon picks up a partner from the glue-polarized vacuum to form a $C = +$ unflavored gg glueball, or it has a transition into a $q\bar{q}$ pair and evolves into $q\bar{q}$ mesons. (We expect the leading production of heavier multi-constituent states such as a three gluon glueball or a $q\bar{q}g$ excited meson to be suppressed.) The former gg option would be favored by simple color charge considerations [10], and would lead to copious glueball production, in fact, where it not that the production of a massive glueball eats up a large rapidity interval in fig. 1. As a result, we expect a kind of threshold effect: only an energetic gluon can fragment into a glueball, but once a gluon is sufficiently energetic, it is expected to fragment preferentially into a glueball. Taking this into account, we only discuss fragmentation into a glueball of the most energetic gluon in a collinear event. For this configuration only, we assume scaling and absence of threshold effects so as to extract an estimate of glueball production.

We now estimate very roughly the glueball production rate at $Q^2 = s = M_\Upsilon^2$, as well as the glueball branching ratio to four charged pions. Taking an energetic gluon and using scaling as described above we can integrate over the Υ peak to write

‡3 Such low sphericity events at Υ are of course contaminated by $\Upsilon \to 1\gamma \to q\bar{q}$. Since we are interested in glueball production from $\Upsilon \to 3g$, this need not worry us. (The $\Upsilon \to 1\gamma \to q\bar{q}$ background may actually be reduced by careful sphericity cuts.)

$$\int_{\Upsilon} \mathrm{d}\sqrt{s}\,\sigma(\mathrm{e\bar{e}} \to \mathcal{G} + \text{anything})$$
$$= \int_{\Upsilon} \mathrm{d}\sqrt{s}\,\sigma(\mathrm{e\bar{e}} \to \Upsilon \to \text{all}) \int_0^1 \frac{\mathrm{d}B_{\Upsilon}^{\mathcal{G}}}{\mathrm{d}z}\,\mathrm{d}z\,, \qquad (1)$$

where $\mathrm{d}B_{\Upsilon}^{\mathcal{G}}/\mathrm{d}z$ is the differential branching ratio and z = glueball momentum/its maximum value = $|\boldsymbol{p}|/|\boldsymbol{p}|_{\max}$. We have $0 \leqslant z \leqslant 1$, important for the normalization. Another possible variable is $z' = E/E_g$; then the normalization of $B_{\Upsilon}^{\mathcal{G}}$ changes with E_g. We might consider this a rough gauge of phase space effects. This turns out not to affect our results seriously for $E_g \approx M_{\Upsilon}/2$ and $M_{\mathcal{G}} \approx 1.5$ GeV, so we use z in the following. The differential branching ratio is

$$\mathrm{d}B_{\Upsilon}^{\mathcal{G}}(z) \equiv \frac{\mathrm{d}\Gamma(\Upsilon \to \mathcal{G} + \text{any})}{\Gamma(\Upsilon \to \text{all})} = \lambda D_{\mathrm{g}}^{\mathcal{G}}(z) B_{\Upsilon}^{3\mathrm{g}}\,\mathrm{d}z\,, \qquad (2)$$

$D_{\mathrm{g}}^{\mathcal{G}}(z)$ being the gluon–glueball fragmentation function, and $B_{\Upsilon}^{3\mathrm{g}}$ and λ stand respectively for $\Gamma_{\Upsilon \to 3\mathrm{g}}/\Gamma_{\Upsilon \to \text{all}}$ and the fraction of low sphericity events for which we believe our estimate makes sense. In the absence of any concrete knowledge of $D_{\mathrm{g}}^{\mathcal{G}}(z)$ we guess that an energetic gluon fragments into a $C = +$ two-gluon glueball analogously to the way a u quark (say) goes into a π^+ containing it. Thus we write

$$\int_0^1 \mathrm{d}z\, D_{\mathrm{g}}^{\mathcal{G}}(z) = 1 - \epsilon\,. \qquad (3)$$

This includes an implicit sum over all types of gg glueballs.

The reason that the total probability for glueball fragmentation of even an energetic gluon is less than unity is that the gluon can turn into a $\mathrm{q\bar{q}}$ pair and then evolve solely into $\mathrm{q\bar{q}}$ mesons. The loss fraction ϵ is a measure of that. We set

$$\epsilon = \int_0^1 \mathrm{d}y\, D_{\mathrm{g}}^{\bar{\mathrm{q}}}(y)\,, \qquad (4)$$

$D_{\mathrm{g}}^{\bar{\mathrm{q}}} = D_{\mathrm{g}}^{\mathrm{q}}$ being the probability for a gluon to evolve to q or $\bar{\mathrm{q}}$. We estimate this by considering the sea of $\mathrm{q\bar{q}}$ pairs in the proton to be generated by gluons. Further, we suppose that the probability for a short-wavelength gluon to evolve to $\mathrm{q\bar{q}}$ is independent of whether it is an isolated hard gluon or a gluon in the nucleon. This is frankly speculative, and may well be far from being true, but this is the best way we can proceed. If $\bar{\mathrm{q}}(x)$, $\mathrm{g}(x)$ are the antiquark and gluon densities in the nucleon, we write a parent–child relation [11] [‡4]

$$\bar{\mathrm{q}}(x) = \int_x^1 \frac{\mathrm{d}y}{y}\,\mathrm{g}(y) D_{\mathrm{g}}^{\bar{\mathrm{q}}}\left(\frac{x}{y}\right) = \int_x^1 \frac{\mathrm{d}y}{y}\,\mathrm{g}\left(\frac{x}{y}\right) D_{\mathrm{g}}^{\bar{\mathrm{q}}}(y)\,. \qquad (5)$$

Now $\mathrm{g}(y)$ is roughly known from QCD applications to deep inelastic ℓN scattering. A popular choice is [12]

$$\mathrm{g}(y) = 3(1-y)^5/y\,. \qquad (6)$$

For any reasonable $D_{\mathrm{g}}^{\bar{\mathrm{q}}}$ [‡5], eqs. (4)–(6) lead to the following determination of ϵ in terms of the deep inelastic function F_2 at $x = 0$:

$$\epsilon = \tfrac{1}{3} \lim_{x \to 0} x\bar{\mathrm{q}}(x) = \tfrac{1}{4} F_2^{\mathrm{eN}}(0)\,, \qquad (7)$$

where an SU_3 symmetric sea has been assumed. With $F_2^{\mathrm{eN}}(0) \sim \frac{1}{2}$ at FNAL energies, one has $\epsilon \sim 1/8$ and eq. (1) may be written as

$$\int_{\Upsilon} \mathrm{d}\sqrt{s}\,\sigma(\mathrm{e\bar{e}} \to \Upsilon \to \mathcal{G} + \ldots)$$
$$\approx \tfrac{7}{8} B_{\Upsilon}^{3\mathrm{g}} \lambda \int_{\Upsilon} \mathrm{d}\sqrt{s}\,\sigma(\mathrm{e\bar{e}} \to \Upsilon \to \text{all})\,. \qquad (8)$$

We take [‡6] $\lambda \approx 0.3$ and the peak integral ≈ 330 nb MeV [13] for Υ a bound state of a charge $-1/3$ quark and its antiquark. From this, the r.h.s. of eq. (8) can be estimated to be ≈ 60 nb MeV.

We favor looking for $\mathcal{G}$ in the four charged pion decay mode, since we expect the fractions into $\pi^+\pi^-$ and

‡4 From the present picture and the parent–child relation the fragmentation function of a sufficiently energetic gluon reads

$$D_{\mathrm{g}}^{\mathrm{h}}(x) = \int_x^1 \frac{\mathrm{d}y}{y} D_{\mathrm{g}}^{\mathcal{G}}(y) D_{\mathcal{G}}^{\mathrm{h}}\left(\frac{x}{y}\right) + \int_x^1 \frac{\mathrm{d}y}{y} D_{\mathrm{g}}^{\mathrm{q}}(y) D_{\mathrm{q}}^{\mathrm{h}}\left(\frac{x}{y}\right) + \int_x^1 \frac{\mathrm{d}y}{y} D_{\mathrm{g}}^{\bar{\mathrm{q}}}(y) D_{\bar{\mathrm{q}}}^{\mathrm{h}}\left(\frac{x}{y}\right).$$

‡5 What is required is that $\lim_{y \to 0} y D_{\mathrm{g}}^{\bar{\mathrm{q}}}(y) = 0$. As an example, on summing over ± 1 g-helicities the $\bar{\mathrm{q}}$ or q into which g materializes satisfies $D_{\mathrm{g}}^{\bar{\mathrm{q}}}(y) \propto y^2 + (1-y)^2$ (footnote 11 of ref. [9]).

‡6 For a numerical estimate we use the fraction of events where one gluon is within 5% of the maximum energy, assuming a constant Dalitz plot density. This is an underestimate [10], since the Dalitz plot density increases slightly toward the boundary in QCD.

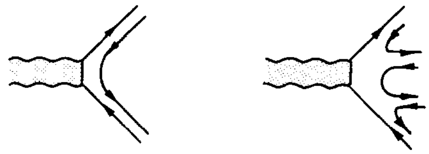

Fig. 2. Decay of a glueball of low mass and spin to $q\bar{q}$ mesons.

$3(\pi^+\pi^-)$ to be small. Establishing the $C = +$ glueball spin from $\mathcal{G} \to 2(\pi^+\pi^-)$ may be difficult; we only note that $\mathcal{G} \to \pi^+\pi^-$ excludes $J^P = 0^-$ and $\mathcal{G} \to \pi A_2$ excludes 0^+. Though we do not discuss $C = -$glueballs here, one might be seen via the allowed $K^+K^-\pi^+\pi^-$ mode. Its presence in this channel and its absence in $2(\pi^+\pi^-)$ excludes $C = +$.

To get an estimate of the $\pi^+\pi^-\pi^+\pi^- \equiv 2(\pi^+\pi^-)$ branching ratio of a glueball of mass ≈ 1.5 GeV, we start by drawing an analogy to $p\bar{p}$ annihilation to mesons. From data [14] on $p\bar{p}$ annihilation near threshold and the plausible notion that the fraction of events with $K\bar{K}$ and η are roughly equal [15], we estimate that $\mathcal{G}$ decays into $N\pi$, N even, about 80% of the time and into $K\bar{K}$ + pions, η + pions 10% of the time each. We expect a low mass and low spin glueball to decay to $q\bar{q}$ states as in fig. 2. The decay $\mathcal{G} \to \pi^+\pi^-$ will be suppressed by form factor effects, as it is for $e^+e^- \to \pi^+\pi^-$, $D^0 \to K^-\pi^+$, or $p\bar{p} \to \pi^+\pi^-$ at threshold. So we concentrate on the four-pion final state. To estimate $B(\mathcal{G} \to 4\pi)$ we take over statistical model results [16] and data [17] on charmed D^0 decay. We set $B(\mathcal{G} \to \pi^+\pi^-) \approx B(D^0 \to K^-\pi^+) \approx 2\%$ (this may be too low actually) and put the branching ratios to 4π and 6π in the same ratio as in Quigg and Rosners' [16] calculation of $D \to K3\pi$ and $K5\pi$ (remembering that $\mathcal{G}$ decays mainly to an even number of pions, by G parity). Then we truncate $\mathcal{G}$ decays at 8π and find $B(\mathcal{G} \to 4\pi) \approx 48\%$.

The main constraint on $B(\mathcal{G} \to 2(\pi^+\pi^-))$ is the existence of a well-known isospin bound on $4\pi^\pm$/all 4π (e.g. ref. [15]). We have

$$\tfrac{1}{3}B(\mathcal{G} \to \text{all } 4\pi) \leqslant B(\mathcal{G} \to 2(\pi^+\pi^-)) \leqslant \tfrac{8}{15}B(\mathcal{G} \to \text{all } 4\pi),$$

$$16\% \leqslant B(\mathcal{G} \to 2(\pi^+\pi^-)) \lesssim 25\% . \tag{9}$$

Our estimate of $B(\mathcal{G} \to \pi^+\pi^-) \approx B(D^0 \to K^-\pi^+) \approx 2\%$ is quite a bit smaller; the estimates mentioned above plus an isospin statistical model gives $B(\mathcal{G} \to 3(\pi^+\pi^-)) \approx 4\%$, also small. Thus,

$$\int_{\Upsilon} d\sqrt{s}\, \sigma(e\bar{e} \to \mathcal{G} + ...)_{\hookrightarrow 2(\pi^+\pi^-)} \approx 10\text{–}15 \text{ nb MeV} . \tag{10}$$

This appears large enough so that even a low statistics Υ experiment has a chance of finding a glueball. Observation of a large branching ratio of Υ to it should clinch the issue. It should appear as a narrow $2(\pi^+\pi^-)$ resonance at Υ but not off resonance. If there are several narrow low mass (≈ 1.5–2 GeV) glueballs, eq. (10) should, of course, be distributed among them.

We do not think of our estimate as unduly optimistic. It would certainly have been optimistic to assume that each g in $\Upsilon \to 3g$ gives rise to a glueball (also implying $\lambda \to 1$). This would make $\mathcal{G}$ production an order of magnitude larger than our estimate – possibly sensible for the next massive $Q\bar{Q}$ state, but not for Υ. As for J/ψ, glueball production should be suppressed by phase space. But we are not aware of any experimental limits.

Nevertheless, there are several places where we could have erred. Phase space effects may be even more important than we judge. Our simple gg constituent picture of a glueball may be wrong. It may be that internal glueball structure is so complex that $D_g^{\mathcal{G}}$ is much suppressed. Or, our estimate of ϵ may be too small (it would be useful to have other ways of obtaining ϵ). Nevertheless, even if the skeptical reader divides our estimate by a factor 5, the resulting $\Upsilon \to \mathcal{G} + ... \to 2(\pi^+\pi^-) + ...$ rate is still about 1% of the peak integral for Υ production. This is comparable to the level at which D mesons were seen [17] in charged decay modes, relative to the total charm production cross section in e^+e^- annihilation. We judge that the prospects of glueball detection are good in collinear gluon jets at Υ. They will be even better at the next heavier $Q\bar{Q}$ bound state.

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