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Massive color-octet bosons and pairs of resonances at hadron colliders

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We analyze collider signatures of massive color-octet bosons whose couplings to quarks are suppressed. Gauge invariance forces the octets to couple at tree level only in pairs to gluons, with a strength set by the QCD gauge coupling. For a spin-1 octet, the cross section for pair production at hadron colliders is larger than that for a quark of equal mass. The octet decays into two jets, leading to a 4-jet signature with two pairs of jets forming resonances of the same mass. For a spin-0 octet the cross section is smaller, and the dominant decay is into $b\bar{b}$, or $t\bar{t}$ if kinematically allowed. We estimate that discovery of spin-1 octets is possible for masses up to 330 GeV at the Tevatron, and 1 TeV at the LHC with 1 fb^{-1}, while the reach is somewhat lower for spin-0 octets.

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1 Introduction

The first discovery of physics beyond the standard model could consist of signals from an effective theory that includes a single new particle. The number of interesting theories of this type is limited because particles are identified by a few quantum numbers (especially spin and gauge charges) which take only a small number of discrete values. Well-known examples include Z' bosons, vectorlike quarks or singlet scalars.

In this letter we study a type of particle present in a variety of theories, whose collider signatures have been less intensely investigated: massive color-octet vector bosons. These arise from an $SU(3)_1 \times SU(3)_2$ gauge group spontaneously broken down to its diagonal subgroup $SU(3)_c$, which is identified with the QCD gauge symmetry. Such a pattern of gauge symmetry breaking, earlier invoked for phenomenological or aesthetic reasons [1], has been used in topcolor models of Higgs compositeness [2,3], and was one of the motivations for studying models with extra dimensions [4]. Color-octet bosons may also be composite particles, such as one of the ρ technimesons [3]. Given that such particles may be associated with many extensions of the standard model, we explore their properties independently of the underlying theory that might justify their existence. In particular, we show that vector octets can give rise to striking multi-jet signatures at hadron colliders and that these could be experimentally observed even within existing data sets once a dedicated analysis is performed.

Distinguished by its coupling to quarks, the spin-1 color octet has been called a coloron (vector coupling [5–7]), an axigluon (axial vector coupling [7,8]), or a topgluon (preferential coupling to top quarks [3, 7], which is also the case for Kaluza-Klein gluons from a warped extra dimension [9]). This illustrates the model dependence of couplings between quarks and color-octet vector bosons, which may even vanish at tree-level, as in the case of level-1 gluons in models with universal extra dimensions. We will present some simple 4-dimensional extensions of the standard model in which such couplings are suppressed, and as a result single production of the color-octet boson is small.

We focus on pair production of spin-1 octets at hadron colliders, which has a large and mostly model-independent rate (see also [10, 11]). The main decay mode is into two quarks, so that the signature is a pair of dijet resonances, for which the backgrounds are highly reducible. Spin-0 octets have similar properties [3,12,13], with the distinction that their couplings to quarks are generically proportional to the quark mass. We study the prospects for observing these particles at the Tevatron and the LHC.

2 Interactions of the vector octet

Consider an effective field theory that includes a vector color-octet field G'_μ of mass M_G, where μ is a

Lorentz index. Any massive spin-1 particle may be identified with the gauge boson of a spontaneously broken gauge symmetry, provided higher-dimensional operators are included. We assume that this effective theory is valid over a range of scales above M_G, so that the higher-dimensional operators are suppressed and may be neglected. The gauge symmetry breaking pattern that gives only a massive spin-1 octet and the massless gluon is $SU(3)_1 \times SU(3)_2 \to SU(3)_c$. Any other gauge group that gives rise to G'_μ should embed this minimal gauge group. The interactions of the massive octet with gluons may be derived by rotating the $SU(3)_1$ and $SU(3)_2$ gauge kinetic terms to the mass eigenstate basis. A pair of G'_μ bosons couples at tree level to one or two gluons:

$$\frac{g_s^2}{2} f^{abc} f^{ade}\, G'^{\mu b} \left[G^{\nu d} \left(G'^c_\nu G^e_\mu + G'^e_\mu G^c_\nu \right) + G'^e_\nu G^{\nu c} G^d_\mu \right] + g_s f^{abc} G'^a_\mu \left[\left(\partial^\mu G'^{\nu b} - \partial^\nu G'^{\mu b} \right) G^c_\nu - G'^b_\nu\, \partial^\mu G^{\nu c} \right] . \tag{1}$$

Here g_s is the QCD gauge coupling, f^{abc} are the $SU(3)_c$ structure constants, and G_μ is the gluon field. The above interactions have an accidental Z_2 symmetry because G'_μ appears only in pairs.

The most general Lorentz-invariant dimension-4 interactions of G'_μ with quarks are of the type $G'^a_\mu \overline{q} \gamma^\mu T^a q'$, where q and q' are quarks carrying the same electroweak charges, and T^a are the generators of the fundamental representation of $SU(3)_c$. The coefficients of these operators form three different 3×3 Hermitian matrices. To avoid a lengthy discussion of flavor-changing processes, we assume that these three matrices have diagonal elements approximately equal (up to a sign) to a parameter $h_q > 0$, and negligible off-diagonal elements. There is an upper limit on h_q set by dijet searches. For a G'_μ mass M_G between 150 and 200 GeV the limit is at most $h_q < g_s/4$ [14], while for some values of M_G above 200 GeV the limit is more stringent, around $h_q < g_s/7$ [15].

Such suppressed couplings of G'_μ to quarks can arise in simple renormalizable models. For example, let us consider an $SU(3)_1 \times SU(3)_2$ gauge theory where the breaking down to $SU(3)_c$ is due to the vacuum expectation value of a complex scalar field Φ (or of a fermion-antifermion pair, induced by some technicolor-like interaction), which transforms as a bifundamental under the product gauge group. After diagonalizing the gauge boson mass matrix, the massless gluon has a gauge coupling $g_s = h_1 h_2/\sqrt{h_1^2 + h_2^2}$, where h_1 and h_2 are the $SU(3)_1 \times SU(3)_2$ gauge couplings. This is to be identified with the QCD coupling at the scale M_G: $g_s \approx 1.1$ for M_G of a few hundred GeV. Imposing perturbativity of both $SU(3)$ interactions at the symmetry breaking scale gives $g_s < h_1, h_2 \lesssim \sqrt{4\pi}$. In the gauge eigenstate basis, the quarks that are triplets under $SU(3)_1$ couple to G'_μ with a strength $h_q = g_s h_1/h_2$, with h_1 and h_2 interchanged for triplets under $SU(3)_2$. A simple choice is that all observed quarks transform as triplets under $SU(3)_1$ [6]. By itself, this would lead to a large coupling of G'_μ to quarks, $h_q \gtrsim g_s^2/\sqrt{4\pi} \approx 0.3$, which would imply that most G'_μ masses between 250 and 750 GeV are ruled out by the CDF search [15]. However, in the presence of new heavy quarks which mix with the observed ones, the couplings of G'_μ may change dramatically. Consider a vectorlike quark Q whose left- and right-handed components transform as a 3 under $SU(3)_2$, and like standard model left-handed quarks (q_L) under $SU(2)_W \times U(1)_Y$. In addition to a mass term for Q, the Lagrangian includes Yukawa couplings of the vectorlike quark to q_L and Φ. The off-diagonal mass term induced by $\langle \Phi \rangle$ requires a rotation of q_L and Q_L by an angle θ, such that the coupling of G'_μ to $\overline{q}_L \gamma^\mu T^a q_L$ becomes

$$h_q = g_s \left(\frac{h_1}{h_2} \cos^2\theta - \frac{h_2}{h_1} \sin^2\theta \right) , \tag{2}$$

and an "off-diagonal" interaction $G'^a_\mu \overline{Q}_L \gamma^\mu T^a q_L$ is induced. For $\tan\theta = h_1/h_2$, we find $h_q = 0$, while the coefficient of the off-diagonal interaction becomes g_s. By including a vectorlike quark for each standard model quark, one may in principle arrange that all tree-level couplings of G'_μ to standard model currents vanish. Such cancellations require fine-tuning, and are unlikely to be realized precisely in nature. Nevertheless, cancellations at the 15% level are sufficient to free G'_μ from the existing limits on dijet resonances. Note that the mixings between standard model quarks and vectorlike ones may be approximately flavor independent, so that the induced flavor-changing neutral currents are not necessarily large.

A more sophisticated model includes an additional $SU(3)$ gauge group, and invariance under a Z_2 symmetry that interchanges two of the groups. There are in this case two heavy spin-1 octets. By virtue of the Z_2 symmetry, the couplings of the lighter octet to standard model quarks vanish exactly. The heavier octet has sizable couplings to quarks in the gauge eigenstate basis, but as in the previous model, in the presence of some vectorlike quarks the cou-

plings to mass eigenstates may be partially canceled. The field content of this model resembles the colored Kaluza-Klein modes of the first two levels in theories with universal extra dimensions.

3 G'_μ production at hadron colliders

The Feynman rules for G'_μ interactions with gluons [see Eq. (1)] is given in Appendix A of Ref. [16]. To leading order in α_s, the partonic processes that lead to G'_μ pair production at hadron colliders have gluons (Fig. 1) or quark-antiquark pairs (Fig. 2) in the initial state. The production cross section may be computed at tree level using CalcHEP [17]. For the gluon-gluon initial state, we find (in agreement with Ref. [10])

$$\sigma(gg \to G'_\mu G'_\mu) = \frac{9\pi\alpha_s^2}{16\hat{s}^3}\left[\beta\hat{s}\left(\frac{8\hat{s}^2}{M_G^2} + 13\hat{s} + 34M_G^2\right) - 8\left(\hat{s}^2 + 3M_G^2\hat{s} - 3M_G^4\right)\ln\left(\frac{1+\beta}{1-\beta}\right)\right] , \quad (3)$$

where $\beta = (1 - 4M_G^2/\hat{s})^{1/2}$ is the boost of G'_μ, and $\hat{s}$ is the center-of-mass energy of the partonic collision. This cross section is independent of $\hat{s}$ for $\hat{s} \gg M_G^2$ as a consequence of spin-1 exchange in the t and u channels. Unitarity is preserved in this process independent of the gauge symmetry breaking sector because the radial modes of the Φ field or whatever else unitarizes longitudinal $G'_\mu G'_\mu$ scattering do not contribute to $gg \to G'_\mu G'_\mu$. Note that Eq. (3) has the same large-$\hat{s}$ behaviour as the cross section for the standard model process $\gamma\gamma \to W^+W^-$ [18].

Assuming negligible couplings of G'_μ to standard model quarks ($h_q \ll 1$), the cross section for $q\overline{q} \to G'_\mu G'_\mu$ depends only on M_G and on the masses of the vectorlike quarks exchanged in the t and u channels. We take these to be of the order of or larger than M_G so that the vectorlike quarks do not affect the G'_μ decays. For vectorlike quark masses equal to M_G, the process with quark-antiquark initial state has a cross section

$$\sigma(q\overline{q} \to G'_\mu G'_\mu) = \frac{\pi\alpha_s^2}{27\hat{s}^2}\left[-\beta\left(83\hat{s} + 72M_G^2\right) + 2\left(20\hat{s} + 49M_G^2\right)\ln\left(\frac{1+\beta}{1-\beta}\right)\right] . \quad (4)$$

This agrees with the cross sections obtained from the squared matrix elements computed in [19]. For vectorlike quark masses (M_Q) larger than M_G or $\sqrt{\hat{s}}$ we find that $\sigma(q\overline{q} \to G'_\mu G'_\mu) \approx \pi\alpha_s^2\hat{s}/(18M_G^4)$ up to corrections of order $1/M_Q^2$. This cross section grows with $\hat{s}$ because the Goldstone boson eaten by G'_μ has a coupling to $\bar{Q}\gamma^\mu T^a Q$ proportional to M_Q, so that Q cannot be much heavier than M_G if we keep $h_q \ll 1$.

4 Signal and background at the Tevatron

Taking the factorization and renormalization scale to be $\sqrt{\hat{s}}/2$, and using the CTEQ6L parton distributions [20], we obtain the leading-order cross section for $G'_\mu G'_\mu$ production at the Tevatron shown in Fig. 3. The shaded (yellow) band denotes the uncertainty in the cross section from varying the factorization and renormalization scale between M_G and $\sqrt{\hat{s}}$. The uncertainty from varying M_Q within the range M_G to $2M_G$ is even smaller. If vectorlike quarks are not included, then the cross section decreases (by a factor of two in the perturbative window allowed by dijet searches, namely $h_q \approx 0.3$ and $M_G \approx 200$ GeV). Next-to-leading order corrections are likely to be large (they are of the order of 50% for $t\bar{t}$ production [21], and G'_μ has larger spin and color representation than the top quark), but computing them is beyond the scope of this letter.

The main G'_μ decays are into $q\bar{q}$, as the G'_μ decays into gluons require higher-dimension operators which we neglect. The signal due to the decay of a G'_μ pair is 4 jets. Assuming equal branching fractions to all quark flavors, and given that decays to

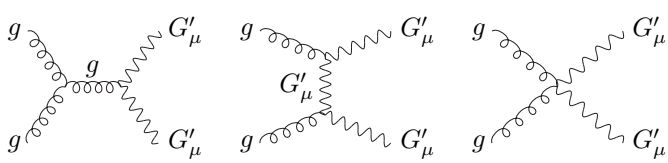

Figure 1: $G'_\mu G'_\mu$ production from gg initial state (u-channel G'_μ exchange is not shown). Curly lines represent gluons, while wavy lines represent massive vector octets.

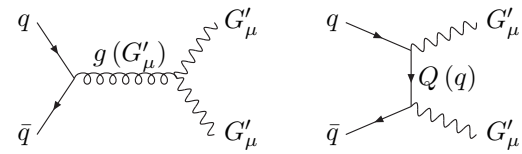

Figure 2: $G'_\mu G'_\mu$ production from $q\bar{q}$ initial state (u-channel diagram is not shown). If couplings of standard model quarks (q) to G'_μ are suppressed due to mixing with vectorlike quarks (Q), then G'_μ or q exchange contributions may be negligible.

$t\bar{t}$ pairs are kinematically forbidden within the mass range accessible at the Tevatron, decays to 4 b-jets, $jjbb$, and $4j$ will occur 4%, 32% and 64% of the time, respectively.

The dominant background is QCD multijet production. Simulating this with both NJETS [22] and MadGraph/MadEvent [23] yielded consistent results of about 59 nb for the cross section with standard cuts for jets (invariant mass $M_{jj} > 10$ GeV, transverse momentum $P_T > 20$ GeV, jet separation $\Delta R > 0.4$, and pseudorapidity $|\eta| < 2.5$). This huge background can be dramatically reduced by requiring that the 4 jets form two dijet resonances with equal invariant mass [13]. Of all possible pairings of the four jets, we choose the one with two pairs closest in invariant mass, thus reducing combinatorial backgrounds. Since the intrinsic width of G'_μ is small (its couplings to quarks are assumed to be smaller than g_s), the width of the visible resonance is set by the detector resolution. Hence, an invariant mass cut of $|M_{jj} - M_G| < 0.2 M_G$ suppresses the background by orders of magnitude (see the curve labelled $jjjj$ in Fig. 3). Final states involving b-jets have even smaller background, but their rates are also suppressed by the b-tagging efficiency, which we assume to be 50% (the b mistag rate is taken to be 1%). Fig. 3 indicates that CDF and DØ could make a 5σ discovery of a G'_μ of mass below approximately 340 GeV (320 GeV) in the $jjbb$ ($4j$) final state with an integrated luminosity of 4 fb^{-1}. The significance increases fast for lower masses (40σ at $M_G = 200$ GeV), allowing flexibility in the choice of cuts. Note that although the QCD background does not include resonant structures, the cuts could induce a false signal with two dijet resonances [13]. However, this false signal can be eliminated because changes in cuts would shift the position of the cut-induced resonances without affecting the location of the resonances due to G'_μ decays. It is likely though that a detailed analysis that includes systematic errors on background would yield a smaller significance, especially at low invariant mass.

5 Spin-0 color octet

We now discuss the case of spin-0 octets. As a consequence of $SU(3)_c$ gauge invariance, at tree level G_H couples to gluons only in pairs (see Eq. (2.4) of [16]). In general, there are no renormalizable interactions of the spinless gluon to standard model fermions (only if the color octet is an $SU(2)_W$ doublet and

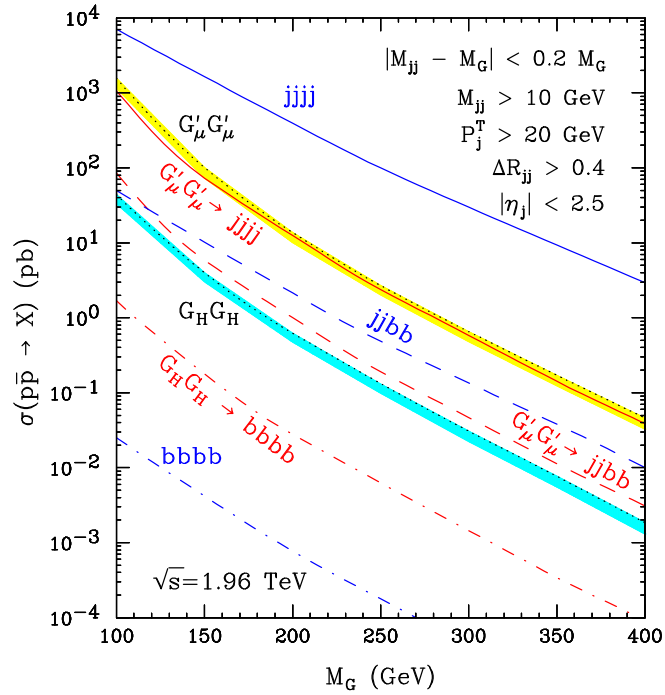

Figure 3: Tevatron cross sections for pair production of spin-1 (G'_μ) and spin-0 (G_H) octets and for backgrounds, as a function of octet mass. Four-jet final states with 0, 2 and 4 b-tags are shown (for both signal and background) as solid, dashed, and dot-dashed lines respectively, for given cuts. Dotted lines represent signal cross sections without cuts, with uncertainties indicated by shaded bands.

carries hypercharge $\pm 1/2$ does it have Yukawa interactions with quarks [24]). However, the following dimension-5 operator can exist

$$\frac{ic_s}{M_G} \left(\bar{q}\gamma^\mu\gamma_5 T^a q\right) \partial_\mu G_H^a \ , \tag{5}$$

where c_s is a dimensionless parameter. For simplicity we ignore flavor off-diagonal operators. The operator (5) can be induced in a renormalizable theory that includes tree-level exchange of either a vector-like quark or a weak-doublet color-octet scalar, or it may be induced at loop level provided that G_H has additional interactions, such as a cubic self-coupling. After integration by parts and use of the field equation for the quark, the operator (5) becomes proportional to the quark mass. As a result, G_H decays into $q\bar{q}$ pairs with widths proportional to the quark mass squared. If the coefficient c_s is flavor independent, G_H decays predominantly into $b\bar{b}$, or $t\bar{t}$ for masses above 350 GeV. However, it is possible that c_s is nonzero only for down-type quarks, so that even above the $t\bar{t}$ threshold the dominant decay is into $b\bar{b}$ (conversely, G_H could couple exclusively to up-type quarks).

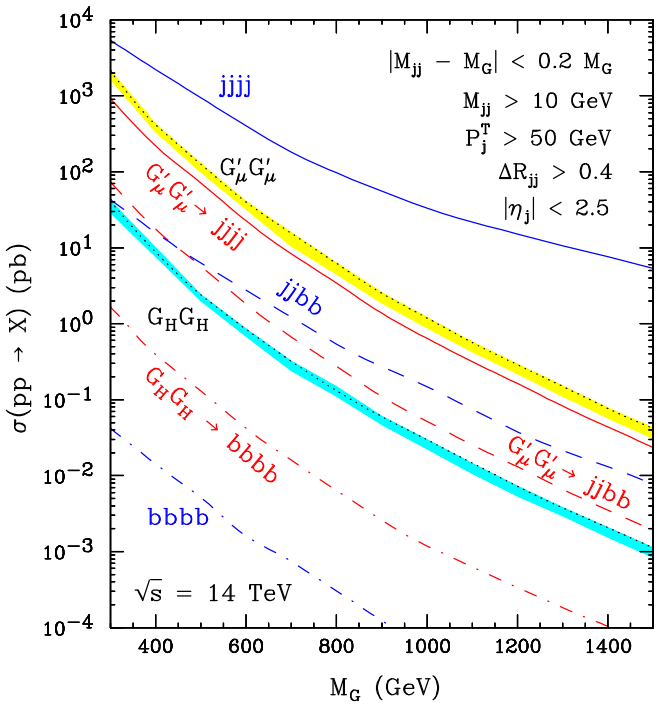

Figure 4: Same as Fig. 3, for the LHC.

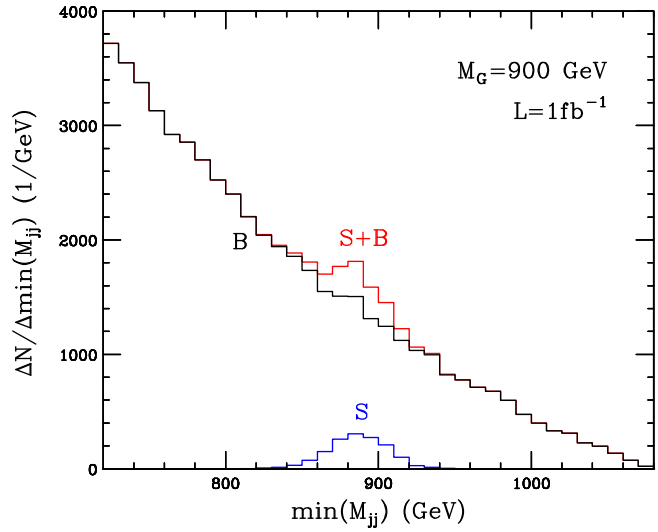

Figure 5: Invariant mass distributions for background (B) and signal (S) at the LHC, for $M_G = 900$ GeV, as a function of the smaller of the two dijet masses. The dijets chosen have the closest invariant masses of all possible pairings of the four jets in the final state. The number of events is normalized to a luminosity of 1 fb^{-1}.

G_H may be produced only in pairs, and the tree-level partonic processes at hadron colliders are analogous to the ones shown in Figs. 1 and 2, with G'_μ replaced by G_H (vectorlike quarks are not included here). The partonic cross sections are given in [13, 24]. In Fig. 3 we show the leading-order $G_H G_H$ production cross section at the Tevatron, and the $4b$ signal after cuts, taking the branching fraction of $G_H \to b\bar{b}$ to be 100%. A 5σ discovery of a G_H with at least 10 signal events can be made for mass below 280 GeV in the $4b$ final state with 4 fb^{-1}. The spin of the octet can be determined using angular distributions, as the events can be fully reconstructed in the center-of-mass frame.

6 Color octets at the LHC

Production cross sections for octets at the LHC are dominated by the gluon initial state [16,24], making them almost entirely model-independent. We plot these in Fig. 4, for equal branching fractions of G'_μ to all quark flavors and a 100% branching fraction of G_H to $b\bar{b}$ (next-to-leading order corrections, which may be sizable, are not included). The production cross section is huge, more than 20 times that for a quark of equal mass, allowing for early discovery. Repeating our Tevatron analysis with a stronger P_T cut of 50 GeV at the LHC, we find a 5σ reach in the $4j$ and $jjbb$ final states for G'_μ of mass below 0.92 and 0.96 TeV (1.2 and 1.3 TeV), respectively, with 1 fb^{-1} (10 fb^{-1}); in the $4b$ channel, requiring 5σ and at least 10 events, the G_H mass reach is 0.75 TeV (1.0 TeV).

Given the large background, one might be concerned that the signal is hard to isolate. We have checked that generally the background distribution falls rapidly with the invariant dijet mass. In some cases a peak may appear near the boundary of the invariant mass distribution, but it is easily eliminated by changing the cuts or binning. In Fig. 5 we show the invariant mass distribution, after cuts, for background and signal in the $4j$ channel for $M_G = 900$ GeV and a luminosity of 1 fb^{-1}. We simulated the response of the hadronic calorimeter using Gaussian smearing with an energy resolution given by [25]:

$$\left(\frac{\Delta E}{E}\right)^2 = \frac{0.5^2}{E\ (\mathrm{GeV})} + 0.03^2\,. \qquad (6)$$

This 6σ excess is clearly visible over the background for the smaller of the two dijet masses, $\min(M_{jj})$. The shift of the peak to a slightly lower mass is due to smearing, and may be corrected, once an excess is observed, by using an alternative variable such as the average of the two dijet masses.

For G'_μ or G_H mass sufficiently above the $t\bar{t}$ threshold there are additional signals: two $t\bar{t}$ resonances [10], or a $t\bar{t}jj$ final state with $t\bar{t}$ and jj separately reconstructed as resonances of same mass. Note also

that if the octet has flavor off-diagonal couplings, then other pairs of resonances would be seen ($t\bar{t}$ + tj, tj + tj, or jj + tj, and similar combinations involving b jets).

To conclude, pair production of heavy color-octet bosons at the Tevatron and LHC is copious and mostly model-independent. It leads to spectacular signatures with jets and b-jets, and possibly top quarks, that reconstruct two narrow resonances of equal mass. Hence the substantial QCD backgrounds are greatly reducible, and dedicated searches can discover such particles for a large range of masses.

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