

Weak-triplet, color-octet scalars and the CDF dijet excess

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We extend the standard model to include a weak-triplet and color-octet scalar. This ‘octo-triplet’ field consists of three particles, two charged and one neutral, whose masses and renormalizable interactions depend only on two new parameters. The charged octo-triplet decay into a W boson and a gluon is suppressed by a loop factor and an accidental cancellation. Thus, the main decays of the charged octo-triplet may occur through higher-dimensional operators, mediated by a heavy vectorlike fermion, into quark pairs. For an octo-triplet mass below the $t\bar{b}$ threshold, the decay into $Wb\bar{b}$ through an off-shell top quark has a width comparable to that into $c\bar{s}$ or $\bar{c}b$. Pair production with one octo-triplet decaying to two jets and the other decaying to a W and two soft b jets may explain the dijet-plus- W excess reported by the CDF Collaboration. The same higher-dimensional operators lead to CP violation in $B_s - \bar{B}_s$ mixing.

1 Introduction

Scalar fields transforming as octets under $SU(3)_c$, the color group of the strong interactions, have been studied in various contexts. The simplest type of color-octet scalar is a singlet under the $SU(2)_W$ group of the weak interactions, and may be referred to as an ‘octo-singlet’. These lead to pairs of dijet resonances at hadron colliders [1, 2], and may explain [3] some deviations from the standard model predictions in the $3b$ search performed by the CDF collaboration [4]. They also enhance the standard model Higgs boson production through gluon fusion [5]. Octo-singlets appear as composite particles due to technicolor [6] and other strong-coupling dynamics [7, 8], or as elementary particles in 6-dimensional extensions of the standard model [9] and in theories with an extended color group [3, 10].

Weak-doublet color-octet scalars (*i.e.*, ‘octo-doublets’) differ dramatically from octo-singlets because the standard model gauge symmetry allows renormalizable couplings of octo-doublets to the standard quarks [11]. Only if these couplings are highly suppressed or aligned with the standard model Yukawa couplings can the octo-doublets be light enough to be produced at the LHC [12]. An octo-doublet field

includes four color-octet states: a charged particle, a neutral one and their antiparticles. The hadron collider signatures of octo-doublets have been explored in [13, 14].

In this paper we study ‘octo-triplets’: real scalar fields that transform in the adjoint representation, $(8, 3, 0)$, of $SU(3)_c \times SU(2)_W \times U(1)_Y$. An octo-triplet field includes three color-octet states: a particle of charge $+1$, its antiparticle, and a neutral real particle. Akin to octo-singlets, octo-triplets are pair produced at hadron colliders through their couplings to gluons, and cannot decay into standard model fermions at renormalizable level because the Yukawa couplings are not $SU(2)_W \times U(1)_Y$ invariant. Unlike octo-singlets, octo-triplets cannot decay into gluons unless there are additional fields that generate certain dimension-7 operators.

One-loop decays of octo-triplets into a gluon and an electroweak boson are allowed, leading to interesting collider signatures involving two gluons and two electroweak bosons. We will show, however, that the rate of these decays is accidentally suppressed by two orders of magnitude compared to usual 1-loop estimates. Thus, new heavy particles could induce the dominant octo-triplet decay modes.

In the presence of some vectorlike quark of mass in the TeV range, the charged octo-triplet may decay into a pair of standard model quarks, or into a W boson and a pair of quarks if it is lighter than the top quark. This leads to a variety of collider signatures, including a dijet resonance, a W boson and two softer jets. If the octo-triplet mass is around 150 GeV this signature may explain the 3.2σ excess observed by the CDF Collaboration in the dijet resonance plus W final state [15]. Some alternative explanations can be found in [3],[16]-[18]. At the LHC, octo-triplets with much larger masses (~ 1 TeV) may be probed in several final states.

Octo-triplets may be elementary particles (*e.g.*, part of the 75 representation of $SU(5)$ grand unification), or may arise as composite ones, for example as fermion-antifermion bound states [8]. We treat the octo-triplets as point-like particles, which is a good approximation only when the compositeness scale is substantially higher than the octo-triplet mass.

In Section 2 we analyze the extension of the standard model by one real octo-triplet field. Section 3 introduces a heavy vectorlike quark which mediates octo-triplet decays into standard model quarks. Our conclusions are summarized in Section 4. In the appendices we present the Feynman rules for octo-triplets, and then we compute the rates for the 3-body weak decay of the charged octo-triplet and for the 2-body 1-loop decays of color-octet scalars.

2 Octo-triplet scalar

We consider the standard model plus an octo-triplet, $\Theta^{a\alpha}$, which is a real field of spin 0 transforming as $(8, 3, 0)$ under the standard $SU(3)_c \times SU(2)_W \times U(1)_Y$ gauge group. We use indices from the beginning of the Roman and Greek alphabets to label the $SU(3)_c$ and $SU(2)_W$ generators, respectively: $a, b, c = 1, \dots, 8$ and $\alpha, \beta, \gamma = 0, 1, 2$.

2.1 Interactions and masses

All interactions of the octo-triplet with standard model gauge bosons are contained in the kinetic term

$$\frac{1}{2} (D_\mu \Theta^{a\alpha}) (D^\mu \Theta^{a\alpha}) \quad , \quad (2.1)$$

where μ is a Lorentz index and the covariant derivative is given by

$$D_\mu \Theta^{a\alpha} = \partial_\mu \Theta^{a\alpha} + g_s f^{abc} G_\mu^b \Theta^{c\alpha} + g \epsilon^{\alpha\beta\gamma} W_\mu^\beta \Theta^{a\gamma} . \quad (2.2)$$

Here f^{abc} and $\epsilon^{\alpha\beta\gamma}$ are the totally antisymmetric tensors of the $SU(3)_c$ and $SU(2)_W$ groups, respectively, g_s and g are the $SU(3)_c \times SU(2)_W$ gauge couplings, G_μ^a is the gluon field, and W_μ^α is the weak gauge field. The octo-triplet field includes three particles: an electrically-neutral color-octet real scalar Θ^{a0} , a color-octet scalar of electric charge +1, Θ^{a+} , and its antiparticle Θ^{a-} :

$$\Theta^{a\pm} = \frac{1}{\sqrt{2}} (\Theta^{a1} \mp i\Theta^{a2}) \quad . \quad (2.3)$$

When referring informally to the octo-triplet particles we use the Θ^\pm and Θ^0 symbols without displaying the color index a .

The kinetic term (2.1) includes interactions of the W boson with two octotriplet particles,

$$-igW_\mu^- [(\partial_\mu \Theta^{a+})\Theta^{a0} - \Theta^{a+}\partial_\mu \Theta^{a0}] + \text{H.c.} \quad , \quad (2.4)$$

and also with an additional gluon:

$$2igg_s f^{abc} G^{\mu a} (W_\mu^+ \Theta^{b-} - W_\mu^- \Theta^{b+}) \Theta^{c0} \quad . \quad (2.5)$$

Similar interactions involve a Z boson and two octo-triplet particles of the same charge, with or without an additional gluon. The interactions of like-sign octo-triplets with one or two gluons (photons) are completely specified by QCD (QED) gauge invariance. The Feynman rules for octo-triplets are given in Appendix A.

The mass of the octo-triplet arises from two terms in the Lagrangian:

$$-\frac{1}{2} (M_0^2 - \lambda_H H^\dagger H) \Theta^{a\alpha} \Theta^{a\alpha} \quad , \quad (2.6)$$

where λ_H is a real dimensionless parameter. The VEV of the standard model Higgs doublet H has a value $v_H \simeq 174$ GeV, so that the mass of the octo-triplet field is

$$M_\Theta = \sqrt{M_0^2 - \lambda_H v_H^2} \quad . \quad (2.7)$$

We require $M_\Theta > 0$ (*i.e.*, $\Theta^{a\alpha}$ does not acquire a VEV) in order to preserve $SU(3)_c$ gauge invariance.

The commutation relations of the Pauli matrices σ^α imply that other operators contributing to

the octo-triplet mass, such as $(H^\dagger \sigma^\alpha \sigma^\beta H) \Theta^{a\alpha} \Theta^{a\beta}$, are either identical to the last one in Eq. (2.6) or vanish. Thus, at tree level Θ^\pm and Θ^0 are degenerate states, having masses equal to M_Θ . At one loop, the electroweak interactions break this degeneracy. The mass splitting between the charged and neutral octo-triplets is [19, 20]

$$\delta M \equiv M_{\Theta^+} - M_{\Theta^0} \simeq \frac{1 - \cos \theta_W}{2 \sin^2 \theta_W} \alpha M_W \quad (2.8)$$

up to corrections of order $(M_W/M_\Theta)^2$. We will see shortly that the octo-triplets have lifetimes much longer than the QCD scale, so that they hadronize. The lightest physical states are ‘‘octo-hadrons’’ given by a Θ^0 or Θ^\pm bound to gluons or quark-antiquark pairs. The mass difference δM between the lightest charged and neutral octo-hadrons is of the same sign and order of magnitude as $M_{\Theta^+} - M_{\Theta^0}$, so that $\delta M \sim 0.2$ GeV.

$SU(2)_W \times U(1)_Y$ gauge-invariance forbids any renormalizable interaction of the octo-triplet with standard model fermions. The most general renormalizable Lagrangian (\mathcal{L}_Θ) for the octo-triplet scalars is given by the kinetic term (2.1), the potential terms quadratic in Θ given in Eq. (2.6), as well as a quartic term and a cubic term:

$$- \lambda_\Theta (\Theta^{a\alpha} \Theta^{a\alpha})^2 + \mu_\Theta f^{abc} \epsilon^{\alpha\beta\gamma} \Theta^{a\alpha} \Theta^{b\beta} \Theta^{c\gamma} \quad , \quad (2.9)$$

where $\lambda_\Theta > 0$ is a dimensionless parameter. The mass parameter μ_Θ may be positive or negative, but its size should not be larger than $O(M_\Theta \lambda_\Theta^{-1/2})$ in order to prevent a Θ VEV. The above cubic term gives the following interaction among the charged and neutral octo-triplet particles:

$$2i\mu_\Theta f^{abc} \Theta^{a+} \Theta^{b-} \Theta^{c0} \quad . \quad (2.10)$$

2.2 Collider signals of octo-triplets

In the $\mu_\Theta \rightarrow 0$ limit, the Lagrangian \mathcal{L}_Θ has an accidental \mathcal{Z}_2 symmetry that makes the lightest octo-triplet (*i.e.*, Θ^0) stable. The charged octo-triplet decays at tree level into Θ^0 and an off-shell W boson. Computing the 3-body width to leading order in δM (Appendix B) we find

$$\Gamma(\Theta^\pm \rightarrow \Theta^0 e^\pm \nu) \simeq \frac{\alpha^2}{15\pi \sin^4 \theta_W} \frac{(\delta M)^5}{M_W^4} \quad . \quad (2.11)$$

Given the small mass splitting $\delta M \sim 0.2$ GeV [see the comment after Eq. (2.8)], the only other relevant decay mode is $\Theta^\pm \rightarrow \Theta^0 \mu^\pm \nu$ with a decay width further phase-space suppressed compared to Eq. (2.11). Hence, the total tree-level width of Θ^\pm is $\Gamma_{\text{tree}}(\Theta^\pm) \simeq 1.8 \times 10^{-16}$ GeV. This 3-body decay width corresponds to a decay length of 1.1 cm.

For $\mu_\Theta \neq 0$, the charged octo-triplet decays into gauge bosons at one loop, with Wg being the only 2-body final state allowed by charge conservation. The diagrams responsible for this decay are shown in Figure 1. The computation of the decay width described in Appendix C gives

$$\Gamma(\Theta^\pm \rightarrow W^\pm g) \simeq \frac{\alpha \alpha_s \mu_\Theta^2}{\pi^3 \sin^2 \theta_W M_\Theta} f(M_W/M_\Theta) \quad , \quad (2.12)$$

where the function $f(R)$ is defined in Eq. (C.8). For M_Θ varying between 150 GeV and 1 TeV, $f(R)$ grows from $(4.0 - 10.3) \times 10^{-3}$, corresponding to $\Gamma(\Theta^\pm \rightarrow W^\pm g)$ in the $(4.3 - 11.2) \times 10^{-7} \mu_\Theta^2/M_\Theta$ range. While one might naively expect $f(R)$ to be of order one, this function is accidentally suppressed: $f(0) \propto (\pi^2/9 - 1)^2$ as shown in Eq. (C.10), while for larger values of R , the function decreases further due to phase space suppression.

The neutral octo-triplet also decays at one loop, into a gluon and Z boson or photon, with partial widths

$$\begin{aligned} \Gamma(\Theta^0 \rightarrow Zg) &\simeq \frac{\alpha \alpha_s \mu_\Theta^2}{\pi^3 \tan^2 \theta_W M_\Theta} f(M_Z/M_\Theta) \quad , \\ \Gamma(\Theta^0 \rightarrow \gamma g) &\simeq \frac{\alpha \alpha_s \mu_\Theta^2}{\pi^3 M_\Theta} f(0) \quad . \end{aligned} \quad (2.13)$$

When M_Θ varies between 150 GeV and 1 TeV, the branching fraction for $\Theta^0 \rightarrow \gamma g$ decreases from 53% to 24%. The decay $\Theta^0 \rightarrow gg$ does not occur at one

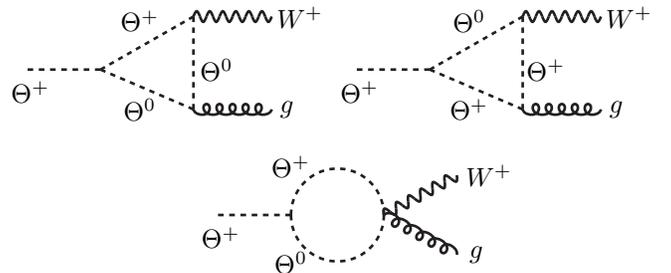


Figure 1: Charged octo-triplet decay to a W boson and a gluon.

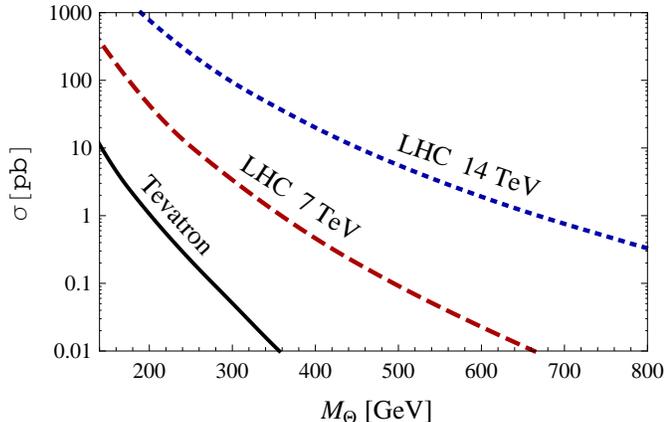


Figure 2: Leading order cross section for charged octo-triplet pair production at the Tevatron (black solid line) and LHC at $\sqrt{s} = 7$ TeV (red dashed line) and at $\sqrt{s} = 14$ TeV (blue dotted line).

loop due to $SU(2)_W$ invariance (this decay requires a dimension-7 operator involving two Higgs fields).

At hadron colliders, octo-triplets are copiously pair produced due to their QCD couplings to gluons. The rate for $\Theta^0\Theta^0$ production is the same as for an octo-singlet of same mass [2, 3], while $\Theta^+\Theta^-$ production is twice as large. In Figure 2 we show the leading order $\Theta^+\Theta^-$ production cross section at the Tevatron and LHC, computed with MadGraph [21] using the CTEQ 6 parton distribution functions [22]. The QCD corrections are not included in this plot; we expect their inclusion to shift these curves upwards by $O(50\%)$.

The $\Theta^0\Theta^0$ pair leads to $(Zj)(Zj)$, $(\gamma j)(Zj)$ and $(\gamma j)(\gamma j)$ final states, where j is a gluonic jet and the parantheses indicate that the two objects form a resonance of mass M_Θ . The $\Theta^+\Theta^-$ pair leads to $(W^+j)(W^-j)$ final states, unless $\mu_\Theta^2/M_\Theta \lesssim O(10^{-9})$ GeV which leads to a large branching fraction for the $\Theta^\pm \rightarrow \Theta^0 e^\pm \nu$ decay. This latter case gives the same final state as in $\Theta^0\Theta^0$ production because the electron and neutrino are very soft and most likely do not pass the cuts (even when Θ^\pm is boosted the electron is not isolated). If $\mu_\Theta^2/M_\Theta \simeq 4 \times 10^{-10}$ GeV, then the 2- and 3-body decays of Θ^\pm have comparable widths, so that the $\Theta^+\Theta^-$ pair leads to $(Wj)(Zj)$ and $(Wj)(\gamma j)$ final states, with the (Zj) and (γj) vertices originating from displaced vertices.

Let us briefly discuss the $(W^+j)(W^-j)$ signal at

the LHC, where the production cross section, Figure 2, can be very large. This final state is most easily identified when one W decays leptonically while the other decays hadronically with an overall $W+4j$ signature. The W decay products reconstruct W resonances and can thereby be isolated from the other jets in the event. These remaining jets can then be grouped alongside the known W decay products and used to reconstruct pairs of octo-triplet resonances. Backgrounds to this signature include W +jets production and $t\bar{t}$ pair-production, the latter of which can be substantially reduced by anti- b tagging.

Besides nonresonant pair production, octo-triplet scalars may induce resonant signatures at hadron colliders because, like other long-lived colored particles [23, 24], they form bound states. If the octo-triplet width is much less than the binding energy E_B due to gluon exchange, then bound states form before either particle decays. This is the case for all values of M_Θ and μ_Θ , as the dominant 2-body octo-triplet width given in Eq. (2.12) easily satisfies

$$\Gamma(\Theta^\pm \rightarrow W^\pm g) \ll E_B = \frac{9}{4}\alpha_s^2 M_\Theta, \quad (2.14)$$

For simplicity, we consider only the formation of color-singlet bound states \mathcal{B} ; this is the dominant channel and our qualitative conclusions apply to different color-representations of bound states.

The bound states annihilate into gauge-boson pairs before either constituent decays. Bound states of neutral octo-triplets can only annihilate through the processes $\Theta^0\Theta^0 \rightarrow \mathcal{B} \rightarrow gg, W^+W^-$, while the annihilation of bound states of charged octo-triplets yields a rich variety of vector boson pairs: $\Theta^+\Theta^- \rightarrow \mathcal{B} \rightarrow gg, W^+W^-, ZZ, \gamma\gamma, \gamma Z$.

Given that the bound-state effects are mainly due to gluon exchange, the production of the $\Theta^+\Theta^-$ bound state is approximately equal to that of octo-singlets computed in [8]: for $M_\Theta \approx 150$ GeV the Tevatron cross section is $\sigma(p\bar{p} \rightarrow \mathcal{B}) \simeq O(100)$ fb. Since the width $\Gamma(\Theta^\pm \rightarrow W^\pm g)$ is orders of magnitude smaller than the main channel for bound states $\Gamma(\mathcal{B} \rightarrow gg) \simeq 0.04$ GeV [8], the annihilation dominates and yields dijet resonances with invariant mass $M_B = 2M_\Theta - E_B$. This cross section is too small to be observed at the Tevatron.

At the LHC, the bound state production cross-section can be considerably larger. For $M_B \sim 1$ TeV,

the cross section is $\sigma(pp \rightarrow \mathcal{B}) \sim \mathcal{O}(1\text{ pb})$ [24] at $\sqrt{s} = 14\text{ TeV}$, which might allow the annihilation signal to compete with the QCD background and give an observable resonance. The electroweak diboson channels are suppressed relative to gg , but give cleaner signals, which contribute to standard Higgs searches.

While the above discussion has been limited to the dominant color-singlet bound state, octo-triplets can also form bound states in higher color representations with exotic annihilation signatures. For instance, color-octet bound states necessarily annihilate into either γg or Zg regardless of whether the bound state comprises charged or neutral scalars.

3 Octo-triplet decays via higher-dimensional operators

Since the octo-triplet widths in Eqs. (2.11)-(2.12) are tiny, higher-dimensional operators induced at the TeV scale could lead to other decays with substantial branching fractions.

Dimension-5 operators allow the coupling of an octo-triplet to a pair of standard model quarks involving a derivative,

$$\frac{c_{ij}}{m_\psi} \Theta^{a\alpha} \bar{Q}_L^i T^a \frac{\sigma^\alpha}{2} \gamma^\mu D_\mu Q_L^j + \text{H.c.} \quad , \quad (3.1)$$

or in the presence of the Higgs doublet,

$$\Theta^{a\alpha} \bar{Q}_L^i \gamma^\mu T^a \frac{\sigma^\alpha}{2} \left(\frac{c'_{ij}}{m_\psi} H u_R^j + \frac{c''_{ij}}{m_\psi} \tilde{H} d_R^j \right) \quad . \quad (3.2)$$

Q_L^i , u_R^j and d_R^j are the quark fields in the gauge eigenstate basis; $i, j = 1, 2, 3$ label the fermion generation; σ^α is a Pauli matrix; c_{ij} , c'_{ij} and c''_{ij} are dimensionless coefficients, and m_ψ is the mass of some heavy field that has been integrated out.

3.1 Octo-triplet plus a vectorlike quark

The dimension-5 operators (3.1) and (3.2) can be induced, for example, by a heavy vectorlike quark Ψ that transforms as $(3, 2, 1/6)$ under $SU(3)_c \times SU(2)_W \times U(1)_Y$, *i.e.* the same way as SM quark doublets Q_L^i . Renormalizable interactions of Ψ with the octo-triplet,

$$\mathcal{L}_{\Theta\Psi} = \Theta^{a\alpha} \bar{\Psi}_R T^a \frac{\sigma^\alpha}{2} (\eta_i Q_L^i + \eta_\psi \Psi_L) + \text{H.c.} \quad , \quad (3.3)$$

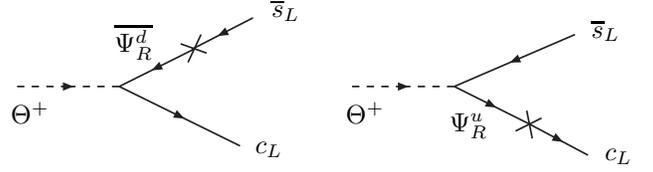


Figure 3: Charged octo-triplet decay to quarks in the presence of a vectorlike quark doublet $\Psi = (\Psi^u, \Psi^d)$. The mass mixing of Ψ with the standard model quarks is depicted by \times . Similar diagrams lead to the decay of the neutral Θ^0 scalar into quark pairs.

and with the Higgs doublet,

$$\mathcal{L}_{H\Psi} = \bar{\Psi}_L \left(\lambda_i^u H u_R^i + \lambda_i^d \tilde{H} d_R^i \right) \quad , \quad (3.4)$$

are allowed. Here η_i , η_ψ , λ_i^u and λ_i^d are dimensionless couplings and $\tilde{H} = i\sigma^2 H^\dagger$. Gauge-invariant fermion mass terms are also allowed:

$$- m_\psi \bar{\Psi}_L \Psi_R - \mu_i \bar{Q}_L^i \Psi_R + \text{H.c.} \quad (3.5)$$

For $m_\psi \gg M_\Theta$, the Ψ fermion can be integrated out, giving rise to the operators (3.1) through the $\mathcal{L}_{\Theta\Psi}$ interactions, and to the operators (3.2) through a combination of $\mathcal{L}_{\Theta\Psi}$ and $\mathcal{L}_{H\Psi}$ interactions.

Let us assume for simplicity that λ_i^u and λ_i^d are negligible, and that the mass mixing parameters satisfy $\mu_i \ll m_\psi$. In this case, only the operators (3.1) are generated, as shown in Figure 3, and their coefficients can be computed in the mass insertion approximation:

$$c_{ij} = - \frac{i \mu_j \eta_i^*}{m_\psi} \quad . \quad (3.6)$$

These are the coefficients at the scale m_ψ ; running down from m_ψ to M_Θ may change c_{ij} at M_Θ by an $\mathcal{O}(1)$ factor, which we can absorb into the definition of η_i .

The width for the decay of the charged octo-triplet into a quark pair is

$$\Gamma(\Theta^+ \rightarrow c \bar{s}) \simeq \frac{m_c^2 + m_s^2}{64 \pi m_\psi^4} |\eta_2 \mu_2|^2 M_\Theta \quad , \quad (3.7)$$

where we have omitted $\mathcal{O}(m_q^4)$ terms and have not included QCD corrections. Taking the charm quark

mass $m_c = 1.3 \text{ GeV}$, we find

$$\Gamma(\Theta^+ \rightarrow c\bar{s}) \simeq 1.3 \times 10^{-6} \text{ GeV} \frac{|\eta_2 \mu_2|^2}{m_\psi^2} \times \frac{M_\Theta}{150 \text{ GeV}} \left(\frac{1 \text{ TeV}}{m_\psi} \right)^2. \quad (3.8)$$

Compared with the decay into Wg computed in Eq. (2.12), the above Θ^+ decay into a pair of jets can easily dominate. For example, for $M_\Theta = 150 \text{ GeV}$, $\mu_\Theta = 1 \text{ GeV}$, $\mu_2/m_\psi = 0.1$, $\eta_2 = 1$, and $m_\psi = 1.1 \text{ TeV}$, we find $\Gamma(\Theta^+ \rightarrow c\bar{s}) \simeq 3.7 \Gamma(\Theta^+ \rightarrow Wg)$.

The width for the decay into $c\bar{b}$ is sensitive to different η and μ parameters:

$$\frac{\Gamma(\Theta^+ \rightarrow c\bar{b})}{\Gamma(\Theta^+ \rightarrow c\bar{s})} \simeq \frac{m_b^2}{m_c^2} \left| \frac{\eta_3^*}{\eta_2} - \frac{\mu_3^*}{\mu_2} \right|^2, \quad (3.9)$$

where $m_b \approx 4.2 \text{ GeV}$ is the b quark mass.

If $M_\Theta > m_t + m_b$, the decay involving a top quark opens up:

$$\Gamma(\Theta^+ \rightarrow t\bar{b}) \simeq 2.2 \times 10^{-2} \text{ GeV} \frac{|\eta_3 \mu_3|^2}{m_\psi^2} \times \frac{M_\Theta}{150 \text{ GeV}} \left(\frac{1 \text{ TeV}}{m_\psi} \right)^2, \quad (3.10)$$

where we have set $m_t = 173 \text{ GeV}$. Due to an m_t^2/m_c^2 enhancement compared to Eq. (3.7), this decay dominates unless $|\eta_3 \mu_3| < 10^{-2} |\eta_2 \mu_2|$.

Similar expressions give the 2-body widths for the neutral octo-triplet decaying to $c\bar{c}$, $b\bar{b}$, or top pairs if $M_\Theta > 2m_t$.

The m_t^2/m_c^2 enhancement in Eq. (3.10) is so large that even for $M_\Theta < m_t + m_b$ the 3-body decay through an off-shell top quark, $\Theta^+ \rightarrow W^+ b\bar{b}$, needs to be taken into account. Its width is

$$\Gamma(\Theta^+ \rightarrow W^+ b\bar{b}) = \frac{\alpha M_\Theta^3 |\eta_3 \mu_3|^2}{(16\pi)^2 \sin^2 \theta_W m_\psi^4} \mathcal{F}(M_\Theta). \quad (3.11)$$

The dimensionless function \mathcal{F} is given by integrating the matrix element over phase space:

$$\mathcal{F}(M_\Theta) = \frac{1}{M_\Theta} \int_0^{E_0} d\bar{E}_b \int_{E_0 - \bar{E}_b}^{E_b^{\max}} dE_b (M_\Theta - 2\bar{E}_b)^2 \times \frac{E_b + (E_0 - \bar{E}_b) \left[\frac{2M_\Theta}{M_W^2} (E_0 - E_b) - 1 \right]}{(m_t^2 - M_\Theta^2 - m_b^2 + 2M_\Theta \bar{E}_b)^2 + m_t^2 \Gamma_t^2}, \quad (3.12)$$

where E_0 is the maximum energy of the \bar{b} or b jet,

$$E_0 = \frac{M_\Theta^2 - M_W^2}{2M_\Theta}, \quad (3.13)$$

and E_b^{\max} is the largest b energy for a fixed \bar{b} energy \bar{E}_b ,

$$E_b^{\max} = \frac{E_0 - \bar{E}_b}{1 - 2\bar{E}_b/M_\Theta}. \quad (3.14)$$

In Eq. (3.12) we neglected m_b everywhere (which is a good approximation for $m_b^2 \ll E_0^2$) with the exception of the denominator where the m_b^2 term becomes important for M_Θ near the 2-body threshold, $m_t + m_b$. To cover that case we also included the top quark width, $\Gamma_t \approx 1.3 \text{ GeV}$, in the propagator. For $M_\Theta = 150 \text{ GeV}$,

$$\Gamma(\Theta^+ \rightarrow W^+ b\bar{b}) \simeq 1.8 \times 10^{-7} \text{ GeV} \frac{|\eta_3 \mu_3|^2}{m_\psi^2} \times \frac{\mathcal{F}(M_\Theta)}{\mathcal{F}(150 \text{ GeV})} \left(\frac{M_\Theta}{150 \text{ GeV}} \right)^3 \left(\frac{1 \text{ TeV}}{m_\psi} \right)^2. \quad (3.15)$$

The ratio $\mathcal{F}(M_\Theta)/\mathcal{F}(150 \text{ GeV})$ is given by 1.6 for $M_\Theta = 155 \text{ GeV}$, and by 2.5 for $M_\Theta = 160 \text{ GeV}$.

It is remarkable that the above 3-body decay through a virtual top quark has a width close to that for the 2-body decay into $c\bar{s}$, given in Eq. (3.8). Ignoring the $c\bar{b}$ width, and assuming for illustration that $\eta_3 \mu_3 = \eta_2 \mu_2$ we find that the branching fraction into $Wb\bar{b}$ is 12, 19, 29% for $M_\Theta = 150, 155, 160 \text{ GeV}$, respectively.

3.2 $B_s - \bar{B}_s$ mixing

Since Ψ has flavor-dependent couplings, its interactions can contribute to flavor-changing neutral processes. The largest couplings are to the 3rd and perhaps 2nd generation quarks, so that we expect that the most prominent effect is in $B_s - \bar{B}_s$ meson mixing. This proceeds through the tree-level diagram in Figure 4. Integrating out Θ and Ψ generates the effective four-Fermi operator

$$\mathcal{L}_{B_s - \bar{B}_s} = \left(\frac{\eta_2 \mu_3 m_b}{2M_\Theta m_\psi^2} \right)^2 (\bar{b}_R T^a s_L)^2 + \text{H.c.} \quad (3.16)$$

Here we have used the fermion field equations and ignored terms suppressed by factors of m_s/m_b .

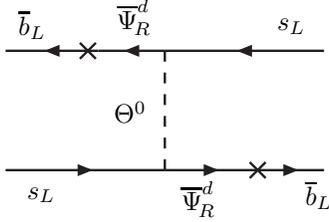


Figure 4: Leading contribution to $B_s - \bar{B}_s$ mixing through Ψ and Θ interactions. Other diagrams differ only by the placement of $\bar{\Psi}_R Q_L$ mass insertions, and are suppressed by additional powers of m_s .

The matrix element of the Hamiltonian due to Θ^0 exchange is

$$\langle \bar{B}_s | \mathcal{H}_\Theta | B_s \rangle \simeq \left(\frac{\eta_2 \mu_3}{M_\Theta m_\psi^2} \right)^2 M_{B_s}^4 f_{B_s}^2 \eta_{\text{QCD}} \frac{5B_2 + 3B_3}{288}, \quad (3.17)$$

where M_{B_s} and $f_{B_s} = (231 \pm 15) \text{ MeV}$ [25] are the B_s meson mass and decay constant respectively; $\eta_{\text{QCD}} \simeq 1.7$ is the QCD correction for the operator in Eq. (3.16) due to running from the scale M_Θ down to M_{B_s} [26]; $B_2 \simeq 0.80$ and $B_3 \simeq 0.93$ are lattice “bag” parameters [27] for the singlet-singlet and octet-octet color structures arising from operator (3.16), respectively.

It is convenient to parametrize the contribution to B_s mixing from Θ^0 relative to the standard model one as

$$\frac{\langle \bar{B}_s | \mathcal{H}_{\text{SM}} + \mathcal{H}_\Theta | B_s \rangle}{\langle \bar{B}_s | \mathcal{H}_{\text{SM}} | B_s \rangle} \equiv C_{B_s} e^{-i\phi_s}, \quad (3.18)$$

where C_{B_s} is a positive parameter and $-\pi < \phi_s < \pi$ is a phase. The standard model contribution can be extracted from the estimate given in [28, 29]:

$$\langle \bar{B}_s | \mathcal{H}_{\text{SM}} | B_s \rangle \approx (8.0 \times 10^{-6} \text{ GeV})^2 (1 \pm 0.15). \quad (3.19)$$

The 15% theoretical uncertainty shown above loosens the constraint on C_{B_s} set by the measured B_s mass difference: $C_{B_s} \approx 0.98 \pm 0.15$. Comparing Eqs. (3.17) and (3.18) we find

$$m_\psi = 1.1 \text{ TeV} \times |\eta_2| \frac{|\mu_3|}{m_\psi} \left(\frac{150 \text{ GeV}}{M_\Theta} \right) \times (C_{B_s}^2 + 1 - 2C_{B_s} \cos \phi_s)^{-1/4}, \quad (3.20)$$

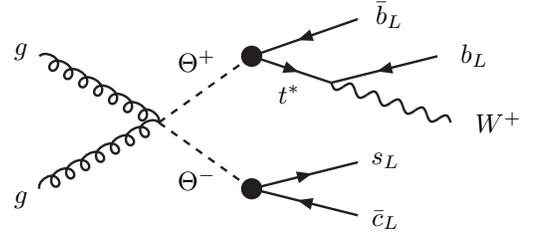


Figure 5: Representative diagram for pair production of charged octo-triplets through gluon fusion with a $(jj)(Wb\bar{b})$ final state. The \bullet symbol denotes dimension-5 operators induced by the Ψ fermion. Similar diagrams lead to $4j$, $(W^+b\bar{b})(W^-b\bar{b})$, or $(jb)(Wb\bar{b})$ final states.

and a less illuminating expression of ϕ_s in terms of the phase of $\eta_2 \mu_3$. For $M_\Theta = 150 \text{ GeV}$, $\mu_3/m_\psi = 0.2$, $\eta_2 = 1$, $C_{B_s} = 0.9$, and a small CP-violating phase $\phi_s = -5^\circ$, we get $m_\psi = 568 \text{ GeV}$. However, if the phase is large, as suggested by the $D0$ like-sign dimuon asymmetry [30], then m_ψ is below the electroweak scale; for example $\phi_s = -45^\circ$ gives $m_\psi = 260 \text{ GeV}$. Such a light vector-like quark is not ruled out. Note that the main decay mode is likely to be $\Psi^d \rightarrow \Theta^+ \bar{c} \rightarrow (c\bar{s})\bar{c}$, so that Ψ pair production leads to a 6-jet final state. The CDF search [31] in a similar channel gives a lower limit on the $3j$ resonance mass below 200 GeV .

This model also contributes to $b \rightarrow s\gamma$ decays at the 1-loop level. Since these diagrams involve two mass insertions and suffer additional loop suppression, we expect their contributions to be small.

3.3 Dijet resonance plus a W boson at the Tevatron

Pair production of octo-triplets, through their QCD couplings to gluons, gives rather large cross sections at the Tevatron, as shown in Figure 2. In the presence of the vectorlike quark Ψ and assuming that the trilinear coupling μ_Θ is small enough (see section 3.1), the main decay modes of Θ^+ are into a pair of jets ($c\bar{b}$ or $c\bar{s}$) and into $Wb\bar{b}$ ($Wb\bar{s}$ is also possible, but at least one b quark is always present due to the decay through the off-shell top quark). One of the final states (see Figure 5) arising from $\Theta^+\Theta^-$ production is then $(jj)(Wb\bar{b})$, where j is any jet and the parantheses indicate a resonance at M_Θ . The branching fractions depend on the $|\eta_2 \mu_2|$ and $|\eta_3 \mu_3|$

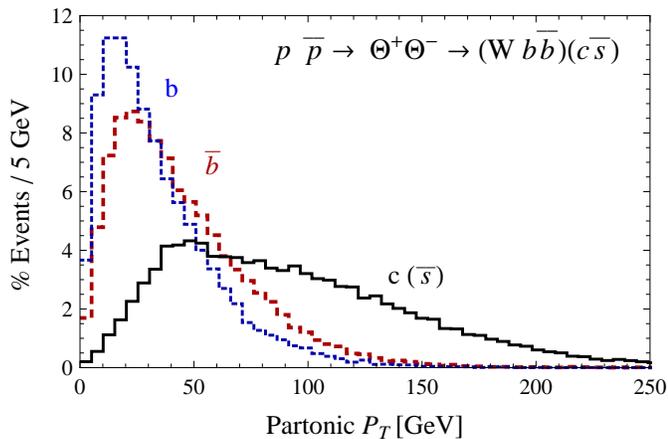


Figure 6: Partonic p_T distributions for the quarks arising from the $\Theta^+\Theta^- \rightarrow (c\bar{s})(W^-b\bar{b})$ process (see Figure 5) with $M_\Theta = 155$ GeV. The c or \bar{s} distribution (black solid line) peaks at higher p_T and has a longer tail than the b and \bar{b} distributions (blue dotted and red dashed lines).

combinations of parameters, and are also quite sensitive to M_Θ , as discussed at the end of section 3.1. We expect that next-to-leading order QCD corrections to this process, which affect both production and decays, increase the rate by a factor $K \sim 1.5$.

The two b jets arising from the decay through an off-shell top quark typically have energies below $(M_\Theta - M_W)/2 - m_b$, so are softer than those arising from the 2-body decay, which typically have energies around $M_\Theta/2$. In Figure 6 we show the transverse momentum of each quark in the $(c\bar{s})(Wb\bar{b})$ final state for $M_\Theta = 155$ GeV, computed with MadGraph [21]. Given that the quarks from the 2-body decay have the highest p_T , the invariant-mass distribution of the two leading jets from the $p\bar{p} \rightarrow \Theta^+\Theta^- \rightarrow (jj)(Wb\bar{b})$ process will exhibit a peak near M_Θ .

In order to compare with the CDF dijet excess [15], we generate events using MadGraph for the partonic level process, PYTHIA [32] for the hadronization and parton shower, and PGS [33] for detector-level effects. We impose the same cuts as CDF [15]: lepton $p_T^\ell > 20$ GeV and $|\eta^\ell| < 1$, $\cancel{E}_T > 25$ GeV, transverse W mass $M_T(W) > 30$ GeV, jet $p_T^j > 30$ GeV and $|\eta^j| < 2.4$, separation between jets $|\Delta\eta_{jj}| < 2.5$, azimuthal separation between the missing E_T and the leading jet $|\Delta\phi| > 0.4$, and $p_T^{jj} \geq 40$ GeV for the leading dijet system. The

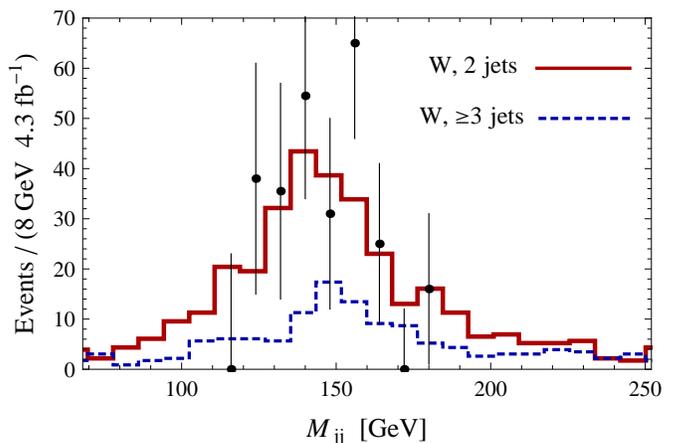


Figure 7: Invariant mass distributions for the leading two jets arising from the $\Theta^+\Theta^- \rightarrow (jj)(Wb\bar{b}) \rightarrow \ell\nu + 4j$ process, where $\ell = e, \mu$, at the Tevatron, for $\mathcal{B}(\Theta^\pm \rightarrow W^\pm b\bar{b}) = 40\%$, $M_\Theta = 155$ GeV and $K = 1.5$. The red solid line represents events where exactly 2 jets pass the cuts, while the blue dashed line includes all events with ≥ 3 jets. The data points with 1σ error bars are taken from the CDF [15] excess region after background subtraction.

resulting dijet invariant mass distribution for events with exactly 2 jets is shown in Figure 7 (solid red line) for $M_\Theta = 155$ GeV and a branching fraction $\mathcal{B}_3 \equiv \mathcal{B}(\Theta^+ \rightarrow W^+ b\bar{b}) = 40\%$. The rate for this process (before cuts and without including the $W \rightarrow \ell\nu$ branching fraction) is

$$2\mathcal{B}_3(1 - \mathcal{B}_3) K \sigma(p\bar{p} \rightarrow \Theta^+\Theta^-) \simeq 3.0 \text{ pb} , \quad (3.21)$$

where $K \approx 1.5$ is our estimate for the QCD corrections. The mass peak fits the CDF dijet excess [15] rather well. The high-mass tail of this distribution is mainly due to events in which the two hardest jets come from different octo-triplets. For $M_\Theta = 150$ GeV the cross section is larger by a factor of 1.7, so that a similar fit is obtained for a smaller $\mathcal{B}_3 \approx 30\%$.

There are a few experimental tests of this interpretation of the CDF excess. Even though the two b jets are relatively soft, the fraction of events that have a 3rd jet that passes all the cuts is large enough to allow the b tagging of the 3rd hardest jet. Furthermore, the additional two b jets allow the reconstruction of the full event. One complication here is that there is a large background from semileptonic $t\bar{t}$ events. Nevertheless, the signal has the property that the reconstructed W boson together with the

two b jets form an invariant mass peak at M_Θ , so that it can be separated from the background.

Another property of this signal is that the invariant mass distribution of the two leading jets also has a peak when three or more jets are allowed. The dashed blue line in Figure 7 includes all $W + n$ jet events with $n \geq 3$. The shape of that distribution is likely to be sensitive to next-to-leading order effects due to additional jets from initial state radiation.

Given that the W in the $(jj)(Wb\bar{b})$ signal originates from a decay through an off-shell top quark, there is no similar signal involving a Z boson.

3.4 Resonant production of $\Theta^+\Theta^-$

It is intriguing that a mass near 150 GeV also appears in a different deviation from the standard model, namely the $3b$ CDF search [4]. That deviation could be due to the $\Theta^0\Theta^0 \rightarrow (b\bar{b})(b\bar{b})$ process [3]. The transverse energy distributions of the jets in that case appear to favor a pair production mechanism through an s -channel resonance rather than through QCD. The simple renormalizable coloron model presented in [3] can be easily adapted to include the octo-triplet discussed here. It is sufficient to charge the scalar field Σ (responsible for breaking the $SU(3) \times SU(3)$ extension of the QCD gauge group [34]) under $SU(2) \times U(1)_Y$, as proposed in [35]. The color-octet scalars present in the spectrum can be identified with our Θ^\pm and Θ^0 (although a small mass splitting can be induced by the Higgs VEV), and they couple to the coloron field G'_μ as follows

$$g_s \frac{1 - \tan^2 \theta}{2 \tan \theta} f^{abc} G'_\mu{}^{ca} \left[\left(\Theta^{b+} \partial^\mu \Theta^{c-} + \text{H.c.} \right) + \Theta^{b0} \partial_\mu \Theta^{c0} \right] . \quad (3.22)$$

Here $\tan \theta$ is a parameter in the 0.1–0.3 range. The coloron couples to quarks proportional to $g_s \tan \theta$, while it couples only in pairs to gluons at renormalizable level [2]. Thus, single G'_μ production proceeds entirely through quark-antiquark collisions.

The resonant $p\bar{p} \rightarrow G'_\mu \rightarrow \Theta^+\Theta^-$ production may be an order of magnitude larger than QCD pair production [3]. The good agreement with the CDF data shown in Figure 7 is then preserved if the branching fraction for $\mathcal{B}(\Theta^+ \rightarrow W^+b\bar{b})$ is decreased accordingly.

3.5 LHC Signals

QCD $\Theta^+\Theta^-$ production, which proceeds through gluon-gluon collisions, is roughly 30 times larger at the 7 TeV LHC than at the Tevatron (see Figure 2), so that the $(jj)(Wb\bar{b})$ signal discussed in section 3.3 will soon be within the reach of the CMS and ATLAS Collaborations. Using typical parameters that explain the CDF dijet resonance, $\mathcal{B}_3 = 40\%$ and $M_\Theta = 155$ GeV, we find that the process in Figure 5 has a leading-order rate (before cuts) of

$$\sigma(pp \rightarrow \Theta^+\Theta^- \rightarrow (jj)(\ell\nu b\bar{b})) \simeq 20 \text{ pb} , \quad (3.23)$$

where $\ell = e, \mu$. Furthermore, the $(W^+b\bar{b})(W^-b\bar{b})$ process also has a large rate, suppressed only by a factor of $2(1/B_3 - 1) \approx 3$, should also be soon within the reach of the LHC.

However, if the resonant production discussed in section 3.4 is the main processes responsible for the CDF excess, then the corresponding rate at the LHC is reduced by approximately an order of magnitude due to the smaller parton distributions for quark-antiquark collisions.

While our analysis has emphasized the region around $M_\Theta = 150$ GeV, future searches could discover much heavier octo-triplets which decay into final states involving top quarks. For non-negligible values of η_3 and μ_3 , the processes $\Theta^+\Theta^- \rightarrow (t\bar{b})(t\bar{b})$ and $\Theta^0\Theta^0 \rightarrow 4t$ are important tests of the octo-triplet decaying through higher-dimensional operators (these final states have been studied in [14]).

4 Conclusions

We have shown that the renormalizable extension of the standard model with one octo-triplet (*i.e.*, a scalar in the adjoint representation of the standard model gauge group) involves two new parameters: the octo-triplet mass M_Θ and cubic self-coupling μ_Θ . For $\mu_\Theta^2/M_\Theta \gg 10^{-9}$ GeV the charged octo-triplet almost always decays into Wg in the absence of other new particles. The rate for this 1-loop process is accidentally suppressed (see Appendix C), but the decay is prompt as long as $\mu_\Theta^2/M_\Theta \gtrsim 10^{-7}$ GeV. The neutral octo-triplet decays to Zg or γg , with widths comparable to that for $\Theta^\pm \rightarrow W^\pm g$. For $\mu_\Theta \rightarrow 0$, the main decay is a tree-level 3-body process, $\Theta^\pm \rightarrow \Theta^0 e^\pm \nu$, with a displaced vertex, while Θ^0 is stable.

At the Tevatron and the LHC, octo-triplets are produced in pairs with relatively large cross sections (see Figure 2). The main signatures are

$$\begin{aligned}\Theta^+\Theta^- &\rightarrow (W^+g)(W^-g) \ , \\ \Theta^0\Theta^0 &\rightarrow (Zg)(Zg) \ , (Zg)(\gamma g) \ , (\gamma g)(\gamma g) \ .\end{aligned}\quad (4.1)$$

The rates for these processes suggest that Tevatron experiments can be sensitive to M_Θ up to a few hundred GeV, and LHC experiments to M_Θ above 1 TeV; however, more precise sensitivity estimates require detailed studies of the backgrounds.

Since octo-triplets have very small widths, decays through higher-dimensional operators may compete with the 1-loop processes. Operators of the type $\Theta\bar{Q}\not{\partial}Q$ may be induced by a heavy vectorlike quark, and lead to the $\Theta^+ \rightarrow t\bar{b}$ decay for $M_\Theta \gtrsim 175$ GeV. For a lighter octo-triplet, there is competition between the 2-body decays $\Theta^+ \rightarrow c\bar{s}$ or $c\bar{b}$ and the 3-body decay $\Theta^+ \rightarrow t^*\bar{b} \rightarrow Wb\bar{b}$ through an off-shell top quark. The neutral octo-triplet decays mainly to $b\bar{b}$, $c\bar{c}$, $W^+b\bar{c}$ and $W^-\bar{b}c$ for $M_\Theta \lesssim m_t + m_c$, to $t\bar{c}$ and $\bar{t}c$ for larger masses below $2m_t$, and to $t\bar{t}$ for masses above $2m_t$. For a range of parameters, the branching fractions for these decays are larger than the ones into a gluon plus an electroweak boson mentioned above. The collider signatures for $M_\Theta \lesssim 175$ GeV then include

$$\begin{aligned}\Theta^+\Theta^- &\rightarrow (jc)(Wb\bar{b}) \ , (bc)(Wb\bar{b}) \ , (W^+b\bar{b})(W^-b\bar{b}) \ , \\ &\quad (jc)(j\bar{c}) \ , (jb)(j\bar{b}) \ , \\ \Theta^0\Theta^0 &\rightarrow (b\bar{b})(b\bar{b}) \ , (b\bar{b})(c\bar{c}) \ , (b\bar{b})(Wbc) \ , \dots \ ,\end{aligned}\quad (4.2)$$

For a heavier octo-triplet the signatures are mainly $(t\bar{b})(\bar{t}b)$ and $(t\bar{c})(\bar{t}c)$, while above 350 GeV the $4t$ final state also opens up.

Signatures of pair production followed by one octo-triplet decaying through higher-dimensional operators and the other decaying into a gluon and an electroweak gauge boson at one loop are also possible. These include $\Theta^+\Theta^- \rightarrow (Wg)(jc)$, $\Theta^0\Theta^0 \rightarrow (Zg)(c\bar{c})$ or $(\gamma g)(c\bar{c})$ and similar processes involving b quarks (or t quarks if kinematically allowed).

Four of the final states mentioned above, namely $(jc)(Wb\bar{b})$, $(bc)(Wb\bar{b})$, $(b\bar{b})(Wbc)$ and $(Wg)(jc)$, may be relevant to the CDF excess [15] in the dijet resonance plus W search. In the case where Θ^+ decays mostly into $c\bar{s}$ and $W^+b\bar{b}$, so that the processes is $p\bar{p} \rightarrow \Theta^+\Theta^- \rightarrow (cj)(Wb\bar{b})$, we have shown that the

b jets are substantially softer than the jets originating from the $\Theta^+ \rightarrow c\bar{s}$ decay. Events where these b jets do not pass the CDF cuts could explain the dijet resonance plus W signal if Θ has a mass in the 150–160 GeV range. It is intriguing that almost the same mass appears in a different deviation from the standard model, namely the $3b$ CDF search [4], which could be attributed to the $\Theta^0\Theta^0 \rightarrow (b\bar{b})(b\bar{b})$ process [3]. The interpretation of the dijet plus W signal in terms of an octo-triplet and vectorlike quark can be tested through the presence of a peak in the invariant mass distribution of the leading two jets from the $W + 3j$ sample, and also through the reconstruction of the $Wb\bar{b}$ peak at the same mass as the dijet peak.

If the couplings of the vectorlike quark are complex, then tree-level Θ^0 exchange induces CP violation in $B_s - \bar{B}_s$ mixing. For a vectorlike quark mass of a few hundred GeV (which is allowed because its main decay is into three jets), this effect can be large enough to produce a significant part of the like-sign dimuon asymmetry observed by the D0 Collaboration [30].

We note that similar final states can also arise from a fermiophobic octo-doublet field. However, this has nontrivial couplings to the standard model Higgs doublet, so the mass splitting between charged and neutral components may be large. By contrast, the tiny mass splitting between charged and neutral octo-triplets suppresses the tree-level 3-body decays. Furthermore, the neutral octo-doublet decays into gluons, while $SU(2)_W$ symmetry forbids pure gluonic decays of the neutral octo-triplet.

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Appendix A: Feynman rules

The Feynman rules for octo-triplets, derived from Eqs. (2.1)-(2.5), are given by:

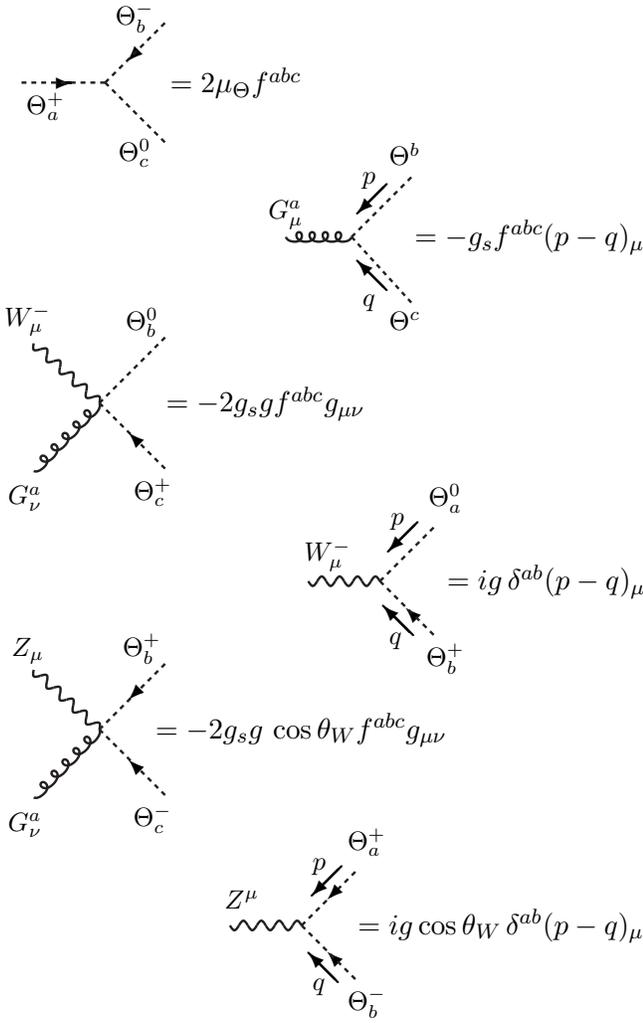


Figure 8: 3-body decay of the charged octo-triplet scalar through an off-shell W boson.

The decay width in the Θ^+ rest frame is then given by

$$\Gamma(\Theta^+ \rightarrow \Theta^0 e^+ \nu) = \frac{g^4 M_{\Theta^+}}{16\pi^3} \int_0^\varepsilon dE_\nu \int_{\varepsilon - E_\nu}^{E_e^{\max}} dE_e \times \frac{M_{\Theta^+} (\varepsilon - E_\nu - E_e) + 2E_\nu E_e}{[M_W^2 - 2M_{\Theta^+} (\varepsilon - E_\nu - E_e)]^2}, \quad (\text{B.3})$$

where ε is the maximum lepton energy,

$$\varepsilon = \frac{M_{\Theta^+}^2 - M_{\Theta^0}^2}{2M_{\Theta^+}}, \quad (\text{B.4})$$

and E_e^{\max} is the largest positron energy for a fixed neutrino energy E_ν ,

$$E_e^{\max} = \frac{\varepsilon - E_\nu}{1 - 2E_\nu/M_{\Theta^+}}. \quad (\text{B.5})$$

The integrals in Eq. (B.3) can be performed analytically, with the result

$$\Gamma(\Theta^+ \rightarrow \Theta^0 e^+ \nu) = \frac{\alpha^2 M_W^4}{16\pi \sin^4 \theta_W M_{\Theta^+}^3} \mathcal{G}\left(\frac{\varepsilon}{M_{\Theta^+}}, \frac{2M_{\Theta^+}^2}{M_W^2}\right) \quad (\text{B.6})$$

where we introduced a function

$$\begin{aligned} \mathcal{G}(x, r) = & -\xi (1+r-rx) \ln(1-x+rx^2+\xi x) \\ & - \frac{1}{2} \left[\left(\frac{1}{2} - x \right) r^2 + \xi^2 - \xi (1+r-rx) \right] \ln(1-2x) \\ & + \frac{r^3 x^3}{3} + \frac{3}{2} r^2 x (1-x) + rx, \end{aligned} \quad (\text{B.7})$$

with

$$\xi \equiv \xi(x, r) = [2r + (1-rx)^2]^{1/2}. \quad (\text{B.8})$$

Interestingly, the expansion of $r^{-4} \mathcal{G}(x, r)$ for $|x| < 1$ starts at x^5 , and the leading r -dependent term arises even later, at x^7 :

$$\mathcal{G}(x, r) = \frac{r^4 x^5}{15} \left[1 + x + \frac{2}{7} (4-r) x^2 \right] + O(x^8). \quad (\text{B.9})$$

The Feynman rules involving photons are identical to those involving a Z boson shown here but with the replacement $g \cos \theta_W \rightarrow e$.

Appendix B: Tree-level octo-triplet decay

In this appendix we compute width for the 3-body decay of the charged octo-triplet through an off-shell W , as shown in Figure 8. We define p_e, p_ν , and p_0 to be the outgoing momenta for e, ν and Θ^0 respectively. Using the Feynman rule from Appendix A, the amplitude for this process is

$$\mathcal{M} \simeq \frac{\sqrt{2} g^2}{M_W^2 - 2p_e \cdot p_\nu} \bar{u}(p_e) \not{p}_0 P_L v(p_\nu), \quad (\text{B.1})$$

where we have used the e, ν equations of motion in the massless lepton limit. Squaring the amplitude and summing over helicities we find

$$|\overline{\mathcal{M}}|^2 = 4g^4 \frac{2(p_e \cdot p_0)(p_\nu \cdot p_0) - p_e \cdot p_\nu M_{\Theta^0}^2}{(M_W^2 - 2p_e \cdot p_\nu)^2}. \quad (\text{B.2})$$

Translating this expansion into a power series in $\delta M \equiv M_{\Theta^+} - M_{\Theta^0}$ we find that the exact tree-level width of Eq. (B.6) is given, up to corrections of order $(\delta M)^8$, by

$$\Gamma(\Theta^+ \rightarrow \Theta^0 e^+ \nu) \simeq \frac{\alpha^2 (\delta M)^5}{15\pi \sin^4 \theta_W M_W^4} \left[1 - \frac{3(\delta M)}{2M_{\Theta^+}} + \frac{4}{7} \left(\frac{9}{8} - \frac{M_{\Theta^+}^2}{M_W^2} \right) \left(\frac{\delta M}{M_{\Theta^+}} \right)^2 \right]. \quad (\text{B.10})$$

Appendix C: One-loop decay of a scalar octet into gauge bosons

In this Appendix we compute the width for a color-octet scalar decaying to a gluon and a (massive or massless) vector boson V_μ , which proceeds through scalar 1-loop diagrams like those of Figure 1. In particular, this computation applies to the process $\Theta^\pm \rightarrow W^\pm g$, $\Theta^0 \rightarrow \gamma g$ or Zg .

We label the V_μ and gluon 4-momenta (polarizations) by p_1 (ϵ_1) and p_2 (ϵ_2), respectively. Since the gluon is always transversely polarized ($\epsilon_2 \cdot p_2 = 0$), angular momentum conservation demands that the other vector also be transverse, so $\epsilon_1 \cdot p_1 = 0$. Given that the contraction $\epsilon_{\mu\nu\rho\sigma} p_1^\mu \epsilon_1^\nu p_2^\rho \epsilon_2^\sigma$, cannot arise from scalar triangle diagrams, the amplitude contains only two terms: $\epsilon_1 \cdot \epsilon_2$ and $(\epsilon_1 \cdot p_2)(\epsilon_2 \cdot p_1)$. Furthermore, by the Ward-Takahashi identity the amplitude vanishes upon replacing ϵ_2 with p_2 , so that the most general amplitude due to scalar loops is given by

$$\mathcal{M} = \frac{\mu_\Theta g_s \tilde{g}}{\pi^2} \mathcal{C} \left(\frac{1}{2} \epsilon_1 \cdot \epsilon_2 - \frac{(\epsilon_1 \cdot p_2)(\epsilon_2 \cdot p_1)}{M_\Theta^2 - M_V^2} \right). \quad (\text{C.1})$$

where μ_Θ is the scalar trilinear coupling, \tilde{g} is the scalar-vector gauge coupling, M_V and M_Θ are the V_μ and scalar masses. The dimensionless coefficient \mathcal{C} is the only quantity that needs to be computed from loop integrals.

We now compute the coefficient \mathcal{C} for the process $\Theta^+ \rightarrow W^+ g$ by evaluating the 1-loop diagrams

$$\mathcal{M} = \int \frac{d^4 k}{(2\pi)^4} \frac{48i \mu_\Theta g_s g}{(k^2 - M_\varphi^2)[(p_1 + p_2 + k)^2 - M_\varphi^2]} \times \left(\frac{1}{4} \epsilon_1 \cdot \epsilon_2 - \frac{\epsilon_1 \cdot (p_2 + k)(\epsilon_2 \cdot k)}{(p_2 + k)^2 - M_\varphi^2} \right), \quad (\text{C.2})$$

where k is the loop 4-momentum and M_φ is the mass of the scalars running in the loop. The logarithmic divergences from the three diagrams cancel, and \mathcal{C} in Eq. (C.1) can be written as a Feynman parameter integral

$$\mathcal{C} = \int_0^1 dx \int_0^{1-x} dy \frac{-3(1-R^2)xy}{1-xy-R^2x(1-x-y)}. \quad (\text{C.3})$$

with $R \equiv M_W/M_\varphi$. After integration over y we obtain

$$\mathcal{C} = \frac{-3}{2(1-R^2)} \left[\frac{\pi^2}{9} - 1 + R^2 \left(\frac{\pi}{\sqrt{3}} - 1 \right) + 2J(R) \right], \quad (\text{C.4})$$

where we have defined

$$J(R) = \int_0^1 dx \left(\frac{1}{x} - R^2 x \right) \ln [1 - R^2 x(1-x)]. \quad (\text{C.5})$$

For $R \ll 1$ the function J has the form

$$J(R) = -\frac{R^2}{2} \left(1 - \frac{R^2}{12} - \frac{R^4}{180} \right) + \mathcal{O}(R^8). \quad (\text{C.6})$$

After squaring the amplitude (C.1) and summing over final state polarizations, we find the following decay width:

$$\Gamma(\Theta^+ \rightarrow W^+ g) = \frac{\alpha_s \alpha \mu_\Theta^2}{\pi^3 \sin^2 \theta_W M_\Theta} f(R), \quad (\text{C.7})$$

where we have assumed $M_\varphi = M_\Theta$, as is the case for the octo-triplet (see Figure 1), and

$$f(R) = \frac{1}{2} \mathcal{C}^2 (1 - R^2), \quad (\text{C.8})$$

with \mathcal{C} depending on R as shown in Eq. (C.4). This function, which appears in all 1-loop decays discussed in this paper, is accidentally suppressed by cancellations between terms involving various powers of π . To see this, consider the expansion around $R \rightarrow 0$:

$$f(R) = f(0) + f_1 R^2 + f_2 R^4 + \mathcal{O}(R^6), \quad (\text{C.9})$$

Each of the above coefficients happens to be much

smaller than order one:

$$f(0) = \frac{9}{8} \left(\frac{\pi^2}{9} - 1 \right)^2 \simeq 1.05 \times 10^{-2} \quad , \quad (\text{C.10})$$

$$f_1 = f(0) + \frac{9}{4} \left(\frac{\pi^2}{9} - 1 \right) \left(\frac{\pi}{\sqrt{3}} - 2 \right) \\ \simeq -3.00 \times 10^{-2} \quad ,$$

$$f_2 = \frac{9}{8} \left(\frac{\pi^2}{9} + \frac{\pi}{\sqrt{3}} - 3 \right)^2 + \frac{3}{16} \left(\frac{\pi^2}{9} - 1 \right) \\ \simeq 2.71 \times 10^{-2} \quad .$$

Thus, the 2-body decays of the octo-triplet are suppressed by two orders of magnitude compared to estimates based on dimensional analysis.

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