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Collider Searches for Fermiophobic Gauge Bosons

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Abstract

We explore the phenomenology of an extra U(1) gauge boson which primarily couples to Standard Model gauge bosons. We classify all possible parity-odd couplings up to dimension six operators. We then study the prospects for the detection of such a boson at the LHC and show that the electroweak decay channels lead to very clean signals, allowing us to probe couplings well into the TeV scale.

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I. INTRODUCTION

One of the most natural possibilities for physics beyond the Standard Model is the existence of new gauge groups. In particular, new $U(1)$ gauge groups which are Higgsed at the TeV scale can lead to new massive gauge bosons (see e.g.[1–4]). Such massive gauge bosons are a generic feature of many extensions of the Standard Model like grand unified theories [5]. String theoretic constructions can also lead to a plethora of new gauge groups [6–15].

The new gauge bosons can couple to the Standard Model in many ways. Usually they are assumed to have direct couplings to the Standard Model fermions, and they can then be directly produced as resonances in colliders. There has been great interest in collider searches for such Z' gauge bosons, and strong constraints have been placed on such resonances [16–18].

A more interesting possibility is if the new gauge boson has no direct couplings to the Standard Model fermions (we will refer to such a gauge boson as being *fermiophobic*). The new gauge boson (hereafter referred to as X) may then have loop-induced couplings to the Standard Model if there are fermions charged under both the new gauge group and the Standard Model. If the fermions are very heavy, then it may be kinematically impossible to produce them on-shell; they would instead be integrated out to yield effective higher dimensional operators coupling X to Standard Model gauge bosons. We will focus here on this possibility.

There are several scenarios for a fermiophobic X . One commonly studied possibility is that of kinetic mixing [19–26], in which there is a dimension 4 operator which mixes the kinetic terms of X and the hypercharge gauge boson. This kinetic mixing induces suppressed couplings between X and the Standard Model fermions, and the X then appears as a Z' with a small coupling. There are, however, many models where such a kinetic mixing term is absent; for example if the heavy fermions are coupled to a non-Abelian Standard Model group, then the kinetic mixing diagrams are forbidden. Effective operators must then couple X to at least two Standard Model gauge bosons [27]. We would then need to search for X through its couplings to two gauge bosons.

If the X couples only to electroweak gauge bosons, X can be produced at hadron colliders through vector boson fusion, followed by the decay $X \rightarrow ZZ \rightarrow 4l$. This possibility was considered in [14], where the authors considered a fermiophobic gauge boson coupled

to electroweak gauge bosons through dimension six operators. This was further extended in [28], where it was pointed out that X can couple to electroweak gauge bosons through dimension four operators as well, enhancing the production cross-section.

Here we consider the more general case where X couples both to gluons as well as to electroweak gauge bosons (as would happen if the heavy fermions couple to $SU(3)_{qcd}$ as well as $SU(2)_L$). We examine the prospects for an LHC search for a massive spin-1 boson coupled to gluons and electroweak gauge bosons through all possible parity-odd couplings up to dimension 6. We find that the on-shell production of X arises through a unique dimension 6 operator coupling X to gluons. However, there is greater freedom in writing operators coupling X to electroweak gauge bosons. As a result, decay can arise through a variety of dimension 4 and 6 operators, the coefficients of which determine the branching fraction to the final states ZZ , $Z\gamma$ and W^+W^- . Interestingly, X cannot decay to $\gamma\gamma$. This follows from the Landau-Yang theorem [29], which asserts that a massive spin-1 boson cannot decay to two massless vector bosons.

The organization of the paper is as follows. In section II we present the effective operator description of the coupling of the hidden sector gauge boson to Standard Model gauge bosons. In section III we describe our analysis of LHC detection prospects for this signal, assuming $\sqrt{s} = 7$ TeV. We conclude with a discussion of our results in section IV.

II. EFFECTIVE THEORY OF THE FERMIPHOBIC GAUGE BOSON

We consider a theory with a new gauge group $U(1)_X$ spontaneously broken by the expectation value of a charged scalar field Φ , which is eaten by the Higgs mechanism giving the gauge boson X a mass. We will consider the case where the gauge boson X has only negligible couplings to Standard Model fermions, but couples non-trivially to Standard Model gauge bosons. We will further specialize to the case where X is a pseudovector; the vector case will be considered elsewhere.

$SU(3)$ gauge invariance constrains the coupling of X to gluons to be a combination of

three effective operators:

$$\begin{aligned}
\mathcal{O}_{Xgg}^1 &= \frac{1}{\Lambda^2} \epsilon^{\mu\rho\alpha\beta} X_\mu D^\nu G_{\alpha\nu}^a G_{\beta\rho}^a \\
\mathcal{O}_{Xgg}^2 &= \frac{1}{\Lambda^2} \epsilon^{\mu\rho\alpha\beta} \partial^\nu X_\mu G_{\alpha\nu}^a G_{\beta\rho}^a \\
\mathcal{O}_{Xgg}^3 &= \frac{1}{\Lambda^2} \epsilon^{\alpha\beta\nu\rho} \partial_\mu X^\mu G_{\alpha\beta}^a G_{\nu\rho}^a,
\end{aligned} \tag{1}$$

where D_μ is a covariant derivative and $G_{\alpha\beta}^a$ is a gluon field-strength. The operator \mathcal{O}_{Xgg}^3 cannot contribute to any process where the X is on-shell, since the momentum of X is orthogonal to its physical polarizations. Thus we can ignore this term if the narrow-width approximation is valid (and we will find that it is). \mathcal{O}_{Xgg}^2 also cannot contribute to any process where the X is on-shell. One can see this by assuming without loss of generality that X is in the rest frame ($p_X = (M_X, 0, 0, 0)$) with polarization $\epsilon_X = (0, 1, 0, 0)$. The only non-vanishing terms are thus $\epsilon^{1\rho\alpha\beta} \partial^0 X_1 G_{\alpha 0}^a G_{\beta\rho}^a$, and it is easy to verify that this expression will vanish due to the antisymmetric property of the epsilon tensor.

The only operator which contributes to on-shell production of X is \mathcal{O}_{Xgg}^1 . The corresponding vertex for this operator is

$$\Gamma_{\mu\nu\rho}^{Xgg}(k_X, k_1, k_2) = \frac{1}{\Lambda^2} [\epsilon_{\mu\nu\rho\sigma} (-k_1^2 k_2^\sigma + k_2^2 k_1^\sigma) + \epsilon_{\mu\rho\sigma\tau} k_{1\nu} k_2^\sigma k_1^\tau - \epsilon_{\mu\nu\sigma\tau} k_{2\rho} k_2^\sigma k_1^\tau]. \tag{2}$$

Note that in this case the vertex is only non-vanishing if at least one gluon is off-shell. This is a consequence of the Landau-Yang Theorem.

Since electroweak symmetry is broken, it is not necessary for operators to exactly satisfy the $SU(2)_L$ Ward Identity. As a result, we may write operators in the effective Lagrangian in terms of the Z and W gauge fields as well as the field strengths. The most general XZZ coupling can be derived from 4 effective operators (see also [30]):

$$\begin{aligned}
\mathcal{O}_{XZZ}^1 &= \epsilon^{\mu\nu\rho\sigma} X_\mu Z_\nu Z_{\rho\sigma} = \epsilon^{\mu\nu\rho\sigma} \frac{X_\mu H^\dagger D_\nu H Z_{\rho\sigma}}{|H|^2} \\
\mathcal{O}_{XZZ}^2 &= \frac{1}{\Lambda^2} \epsilon^{\mu\rho\alpha\beta} X_\mu \partial^\nu Z_{\alpha\nu} Z_{\beta\rho} \\
\mathcal{O}_{XZZ}^3 &= \frac{1}{\Lambda^2} \epsilon^{\mu\rho\alpha\beta} \partial^\nu X_\mu Z_{\alpha\nu} Z_{\beta\rho} \\
\mathcal{O}_{XZZ}^4 &= \frac{1}{\Lambda^2} \epsilon^{\alpha\beta\rho\sigma} \partial_\mu X^\mu Z_{\alpha\beta} Z_{\rho\sigma},
\end{aligned} \tag{3}$$

where $Z_{\alpha\beta}$ is the Z -boson field strength.

Using the same arguments as for the gluon coupling, it is clear that \mathcal{O}_{XZZ}^3 and \mathcal{O}_{XZZ}^4 cannot contribute to any process involving an on-shell X . The vertices for the other two

effective operators are

$$\begin{aligned}\Gamma_{\mu\nu\rho}^{XZZ,1}(k_X, k_1, k_2) &= \epsilon_{\mu\nu\rho\sigma}(k_2^\sigma - k_1^\sigma) \\ \Gamma_{\mu\nu\rho}^{XZZ,2}(k_X, k_1, k_2) &= \frac{1}{\Lambda^2} [\epsilon_{\mu\nu\rho\sigma}(-k_1^2 k_2^\sigma + k_2^2 k_1^\sigma) + \epsilon_{\mu\rho\sigma\tau} k_{1\nu} k_2^\sigma k_1^\tau - \epsilon_{\mu\nu\sigma\tau} k_{2\rho} k_2^\sigma k_1^\tau]\end{aligned}\quad (4)$$

Note that if the Z s are on-shell, as we require, the dimension 4 operator yields the same vertex as the dimension 6 operator:

$$\Gamma_{\mu\nu\rho}^{XZZ,2} \approx -\frac{M_Z^2}{\Lambda^2} \Gamma_{\mu\nu\rho}^{XZZ,1}.\quad (5)$$

Thus we need only consider the dimension 6 operator in the remainder of this paper. In the case where interactions are mediated by a dimension 4 operator, the coupling of X to electroweak states can be easily obtained using the expression above.

The $XZ\gamma$ vertex does not have a symmetry between the two field strengths. For the photon only the field strength can appear, while the field Z_μ can appear by itself. The most general such interaction is a combination of the operators

$$\begin{aligned}\mathcal{O}_{XZ\gamma}^1 &= \epsilon^{\mu\nu\rho\sigma} X_\mu Z_\nu F_{\rho\sigma} \\ \mathcal{O}_{XZ\gamma}^2 &= \frac{1}{\Lambda^2} \epsilon^{\mu\rho\alpha\beta} \partial^\nu X_\mu (Z_{\alpha\nu} F_{\beta\rho} + F_{\alpha\nu} Z_{\beta\rho}) \\ \mathcal{O}_{XZ\gamma}^3 &= \frac{1}{\Lambda^2} \epsilon^{\mu\rho\alpha\beta} \partial^\nu X_\mu (Z_{\alpha\nu} F_{\beta\rho} - F_{\alpha\nu} Z_{\beta\rho}) \\ \mathcal{O}_{XZ\gamma}^4 &= \frac{1}{\Lambda^2} \epsilon^{\mu\rho\alpha\beta} X_\mu \partial^\nu Z_{\alpha\nu} F_{\beta\rho} \\ \mathcal{O}_{XZ\gamma}^5 &= \frac{1}{\Lambda^2} \epsilon^{\mu\rho\alpha\beta} X_\mu \partial^\nu F_{\alpha\nu} Z_{\beta\rho} \\ \mathcal{O}_{XZ\gamma}^6 &= \frac{1}{\Lambda^2} \epsilon^{\alpha\beta\nu\rho} X^\mu \partial_\mu Z_{\alpha\beta} F_{\nu\rho} \\ \mathcal{O}_{XZ\gamma}^7 &= \frac{1}{\Lambda^2} \epsilon^{\alpha\beta\nu\rho} \partial_\mu X^\mu Z_{\alpha\beta} F_{\nu\rho}\end{aligned}\quad (6)$$

where $F_{\alpha\beta}$ is an electromagnetic field strength. The operators $\mathcal{O}_{XZ\gamma}^2$, $\mathcal{O}_{XZ\gamma}^5$ and $\mathcal{O}_{XZ\gamma}^7$ do not contribute to any process in which X and the photons are on-shell.

We can further assume that the only operators we generate are at most dimension 6 when written in manifestly $SU(2)$ -covariant notation. In this case, the only electroweak operators we can write are

$$\begin{aligned}\mathcal{O}^1 &= \frac{C_1}{\Lambda^2} \epsilon^{\mu\rho\alpha\beta} X_\mu \text{Tr}[\partial^\nu C_{\alpha\nu} C_{\beta\rho}] \\ \mathcal{O}^2 &= \frac{C_2}{2\Lambda^2} \epsilon^{\mu\rho\alpha\beta} X_\mu \partial^\nu B_{\alpha\nu} B_{\beta\rho}\end{aligned}\quad (7)$$

where C is the $SU(2)$ gauge field-strength, and B is the hypercharge field-strength.

These operators then completely determine the vertices for XZZ , $XZ\gamma$, XWW and $X\gamma\gamma$ (for on-shell X). Defining

$$\Gamma_{\mu\nu\rho}(k_X, k_1, k_2) = (k_{2\rho}\epsilon_{\mu\nu\sigma\tau}k_1^\sigma k_2^\tau - k_{1\nu}\epsilon_{\mu\rho\sigma\tau}k_1^\sigma k_2^\tau + \epsilon_{\mu\nu\rho\sigma}k_1^\sigma k_2 \cdot k_2 - \epsilon_{\mu\nu\rho\sigma}k_2^\sigma k_1 \cdot k_1),$$

we have

$$\Gamma_{\mu\nu\rho}^{XZZ}(k_X, k_1, k_2) = \frac{1}{\Lambda^2}(C_1 \cos^2 \theta_W + C_2 \sin^2 \theta_W)\Gamma_{\mu\nu\rho}(k_X, k_1, k_2) \quad (8)$$

$$\Gamma_{\mu\nu\rho}^{XZ\gamma}(k_X, k_1, k_2) = \frac{1}{\Lambda^2}(C_1 - C_2) \sin \theta_W \cos \theta_W \Gamma_{\mu\nu\rho}(k_X, k_1, k_2) \quad (9)$$

$$\Gamma_{\mu\nu\rho}^{XW^+W^-}(k_X, k_1, k_2) = \frac{C_1}{\Lambda^2}\Gamma_{\mu\nu\rho}(k_X, k_1, k_2) \quad (10)$$

$$\Gamma_{\mu\nu\rho}^{X\gamma\gamma}(k_X, k_1, k_2) = \frac{1}{\Lambda^2}(C_1 \sin^2 \theta_W + C_2 \cos^2 \theta_W)\Gamma_{\mu\nu\rho}(k_X, k_1, k_2). \quad (11)$$

If all particles are on-shell, these vertices simplify considerably;

$$\Gamma_{\mu\nu\rho}^{XZZ}(k_X, k_1, k_2) = \frac{M_Z^2}{\Lambda^2}(C_1 \cos^2 \theta_W + C_2 \sin^2 \theta_W)\epsilon_{\mu\nu\rho\sigma}(k_1^\sigma - k_2^\sigma) \quad (12)$$

$$\Gamma_{\mu\nu\rho}^{XZ\gamma}(k_X, k_1, k_2) = \frac{M_Z^2}{\Lambda^2}(C_2 - C_1) \sin \theta_W \cos \theta_W \epsilon_{\mu\nu\rho\sigma}k_2^\sigma \quad (13)$$

$$\Gamma_{\mu\nu\rho}^{XWW}(k_X, k_1, k_2) = C_1 \frac{M_W^2}{\Lambda^2} \epsilon_{\mu\nu\rho\sigma}(k_1^\sigma - k_2^\sigma) \quad (14)$$

$$\Gamma_{\mu\nu\rho}^{X\gamma\gamma}(k_X, k_1, k_2) = 0. \quad (15)$$

III. X PRODUCTION AND DECAY

We will be considering processes in which the X boson is produced on-shell in hadron collisions. As we have seen, the Landau-Yang theorem prohibits the decay of a massive spin-1 particle to two massless vector particles and also prohibits resonance production of a massive spin-1 particle from two massless vectors. QCD processes therefore always produce the X boson in association with a jet. Note that this is only true for on-shell production of X ; if X is not on-shell, it can be produced without extra jets. For the moment we neglect this possibility; it would be interesting to see if off-shell production of X can lead to nontrivial results.

The parton-level process $gg \rightarrow Xg$ also vanishes. The only relevant parton-level production channels are therefore $qg \rightarrow qX$, $\bar{q}g \rightarrow \bar{q}X$ and $q\bar{q} \rightarrow gX$.

The branching fractions for X decay can also be calculated. The branching fraction for $X \rightarrow gg$ and $X \rightarrow ggg$ turn out to be zero. As a result, the only hadronic decay of X

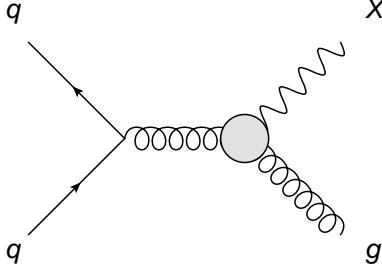


FIG. 1. X production through qg , $\bar{q}g$, and $q\bar{q}$.

to fewer than four jets is through the process $X \rightarrow gq\bar{q}$. Depending on the relative values of the coefficients for the gluon and electroweak operators, this can be an important decay channel.

In this paper we are interested in the electroweak decay channels only. For the purposes of illustrating relative branching fractions to these channels, we will assume that the operator coefficients are chosen such that the partial width for $X \rightarrow gq\bar{q}$ is negligible. (Our final result will be independent of this assumption.) In this case the primary decay modes are ZZ , $Z\gamma$ and WW . We find

$$\begin{aligned}\Gamma(X \rightarrow WW) &= (42 \text{ MeV}) \left(\frac{\text{TeV}}{\Lambda}\right)^4 \left(\frac{M_X}{\text{TeV}}\right)^3 \left(1 - \frac{4M_W^2}{M_X^2}\right)^{5/2} C_1^2 \\ \Gamma(X \rightarrow ZZ) &= (16 \text{ MeV}) \left(\frac{\text{TeV}}{\Lambda}\right)^4 \left(\frac{M_X}{\text{TeV}}\right)^3 \left(1 - \frac{4M_Z^2}{M_X^2}\right)^{5/2} (C_1 + C_2 \tan^2 \theta_W)^2 \quad (16) \\ \Gamma(X \rightarrow \gamma Z) &= (4.9 \text{ MeV}) \left(\frac{\text{TeV}}{\Lambda}\right)^4 \left(\frac{M_X}{\text{TeV}}\right)^3 \left(1 - \frac{M_Z^2}{M_X^2}\right)^3 \left(1 + \frac{M_Z^2}{M_X^2}\right) (C_2 - C_1)^2.\end{aligned}$$

Note that for $M_X, \Lambda \sim \text{TeV}$, the decay width of X is indeed much smaller than its mass, justifying our use of the narrow-width approximation.

In Fig. 2 we plot the branching fractions $BR(X \rightarrow ZZ, W^+W^-, Z\gamma)$ as a function of C_2/C_1 for $M_X = 250 \text{ GeV}$ and $M_X = 1000 \text{ GeV}$.

IV. COLLIDER ANALYSIS

In this analysis we will study potential signals at the 7 TeV LHC. We will focus on the case of X production through QCD couplings via the operator

$$\mathcal{O}_{Xgg} = \mathcal{O}_{Xgg}^1 = \frac{1}{\Lambda^2} \epsilon^{\mu\rho\alpha\beta} X_\mu D^\nu G_{\alpha\nu}^a G_{\beta\rho}^a \quad (17)$$

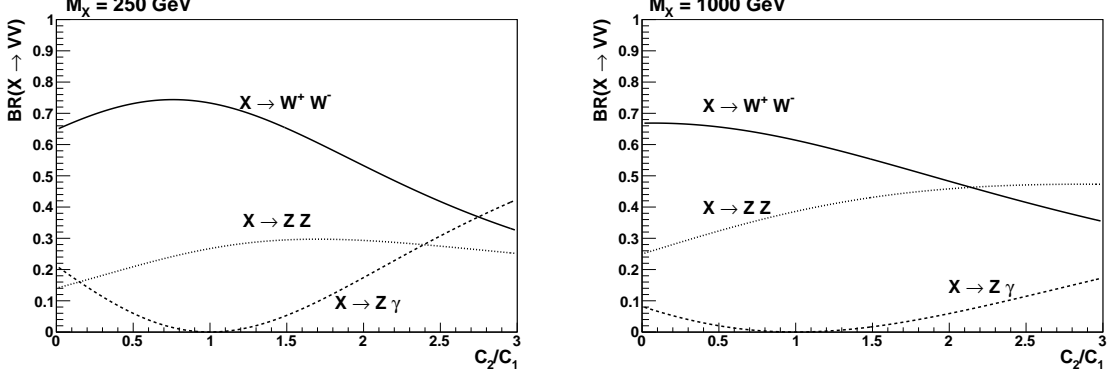


FIG. 2. Branching ratios for X decaying to Standard Model electroweak gauge bosons for $M_X = 250$ GeV (left) and for $M_X = 1000$ GeV (right). We have assumed that the branching fraction to $gq\bar{q}$ is negligible.

followed by $X \rightarrow ZZ$ and $X \rightarrow Z\gamma$ decays, which are the cleanest. We will further specialize to the case where the Z decays to leptons. We have simulated the signal and Standard Model background in MadGraph 5 [31], showered the partons using Pythia 6.4.22 [32], and performed a detector simulation in PGS4 [33]. We consider each final state separately.

A. Cuts

(a) *ZZ decays*: For $X \rightarrow ZZ$ decays the signal is 4 leptons plus a jet. The primary background is $ZZ + \text{jet}$ production. We impose the following cuts:

- One jet with $p_T \geq 50$ GeV, $|\eta| < 2.5$
- 4 leptons with $p_T \geq 20$ GeV, $|\eta| < 2.5$, and pairwise invariant masses in the range 80-100 GeV

(b) *Z γ decays*: For $X \rightarrow Z\gamma$ decays the signal is 2 leptons, a photon and a jet. The primary background is $Z\gamma + \text{jet}$ production. We impose the following cuts:

- One jet with $p_T \geq 50$ GeV, $|\eta| < 2.5$
- 2 leptons with $p_T \geq 20$ GeV, $|\eta| < 2.5$, and invariant mass in the range 80-100 GeV
- 1 photon with $p_T \geq 10$ GeV

To look for the X resonance, we can study the total invariant mass of the 4 leptons (or 2 leptons and photon). The invariant mass distributions for the signal vs. background (assuming the only electroweak coupling is through operator \mathcal{O}^1) are shown in Fig. 3 for $M_X = 250$ GeV and for various values of Λ . The Standard Model background events give a smooth distribution over the relevant invariant mass combinations (see also [34, 35]). The cross sections for the signal are well above background for Λ as high as 2 TeV.

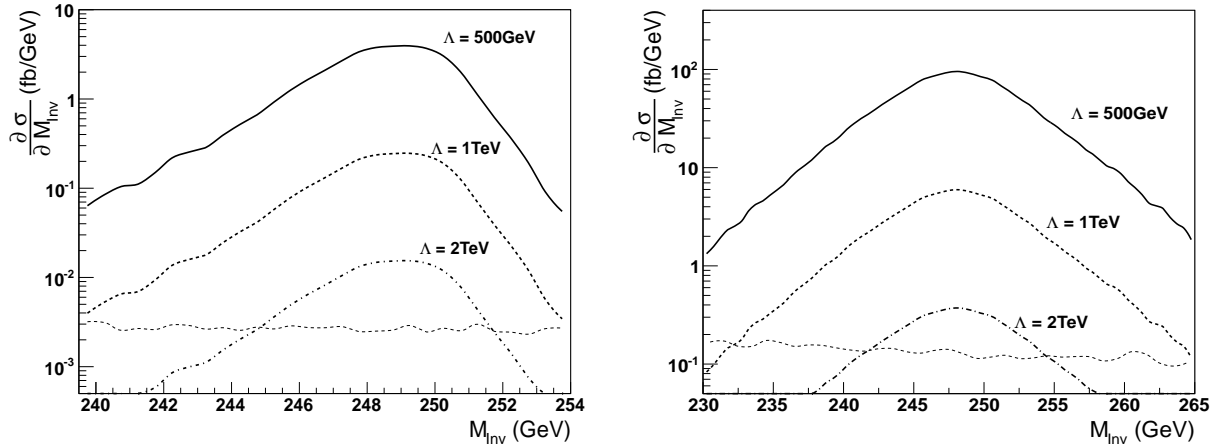


FIG. 3. Invariant mass spectrum for signal ($M_X = 250$ GeV) and background for the LHC at $\sqrt{s} = 7$ TeV, assuming the only electroweak coupling is through operator \mathcal{O}^1 . *Left:* Signal and background for $X \rightarrow ZZ$, for different values of Λ . *Right:* Signal and background for $X \rightarrow Z\gamma$. Both signal and background cross-sections are generally lower for the ZZ process due to the extra factor of the dilepton branching ratio.

Since signal events will exhibit a narrow invariant mass peak, our analysis will compare the number of observed events to the number of expected background events with an invariant mass within $\pm 10\%$ of a given central value $m_{central}$. For both ZZ and $Z\gamma$ channels, this invariant mass cut drastically lowers background cross-sections. In Table I we present the signal and Standard Model background cross-sections for events satisfying the cuts with 4 lepton (or 2 lepton plus photon) invariant mass within 10% of the given $m_{central}$.

B. Detection Prospects

We find the number B of background events with $4l$ ($2l + \gamma$) invariant mass within $\pm 10\%$ of any given $m_{central}$ and compare this to the number S of signal events within the same

TABLE I. Table of signal and Standard Model background production cross-sections (in fb) for the LHC at $\sqrt{s} = 7$ TeV, for 4 leptons and for 2 leptons and a photon, as labeled. The signal cross-sections are normalized by taking $\Lambda/BR(X \rightarrow VV)^{\frac{1}{4}} = \text{TeV}$ for each final state. We assume the cuts described in the text and the assumption that the invariant mass is within 10% of $m_{central}$.

$m_{central}$ (GeV)	$\sigma_{BG}(\text{fb})$ $pp \rightarrow j l^+ l^- l^+ l^-$	$\sigma_{BG}(\text{fb})$ $pp \rightarrow j \gamma l^+ l^-$	$\sigma_{sig}/BR(ZZ)(\text{fb})$ $pp \rightarrow j l^+ l^- l^+ l^-$	$\sigma_{sig}/BR(Z\gamma)(\text{fb})$ $pp \rightarrow j \gamma l^+ l^-$
250	0.26	6.4	18	690
500	0.050	0.76	5.5	140
750	0.010	0.17	1.3	36
1000	0.0021	0.034	0.26	9.6
1250	0.0004	0.014	0.054	2.5
1500	0.0001	0.0051	0.012	0.66
1750	<0.0001	<0.0010	0.0032	0.18
2000	<0.0001	<0.0010	0.0008	0.049

invariant mass window, assuming $M_X = m_{central}$. The significance is defined as $s = \frac{S}{\sqrt{B}}$. For each point in parameter space, we can find the luminosity required to achieve discovery. When the number of expected background events at a certain luminosity is less than one, we define discovery as $S \geq 5$; otherwise, we define discovery as $s \geq 5$. For all of the parameter space considered one finds $S/B \geq 0.2$ at discovery.

In a realistic experimental analysis the actual signal significance would be reduced by a trials factor associated with the freedom in choosing $m_{central}$, the center of the invariant mass analysis window.

Since we have seen that the narrow-width approximation is valid for X , the detection prospects of the LHC depend on the electroweak coupling operator coefficients (C_1 and C_2) only through the branching fraction for X to decay to each channel. Since the minimum cross-section for discovery scales as the production cross-section times the branching ratio,

$$\sigma_{pp \rightarrow X + jet \rightarrow VV jet} = \sigma_{prod} \times B(X \rightarrow VV) \quad (18)$$

$$\propto \Lambda^{-4} \times B(X \rightarrow VV), \quad (19)$$

we define the mass reach in terms of the quantity $\Lambda/[B(X \rightarrow VV)]^{\frac{1}{4}}$. This mass reach is then independent of the relative strengths of couplings to the strong and electroweak sectors,

as encoded by the factors C_1 and C_2 . We plot the mass reach accessible at the LHC for various luminosities at collider energy 7 TeV in both the ZZ and $Z\gamma$ channel in Fig. 4.

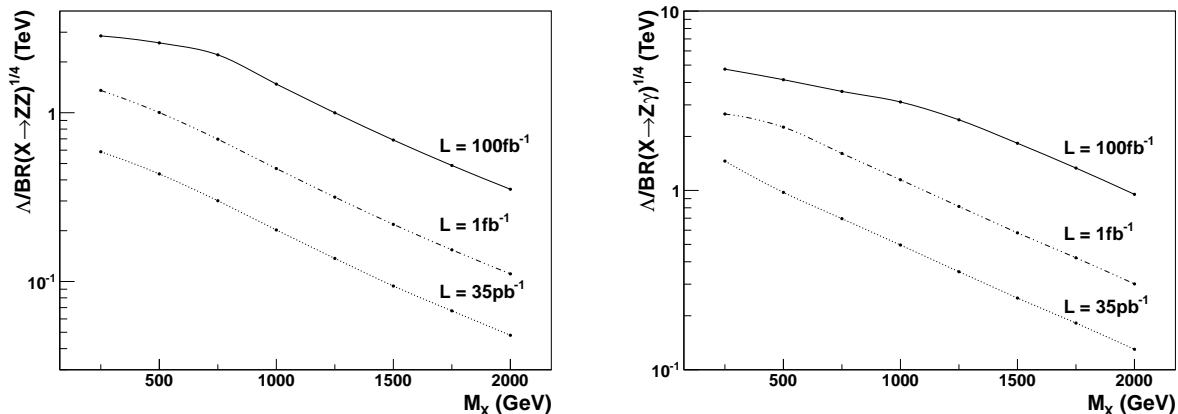


FIG. 4. Discovery reach of the LHC at $\sqrt{s} = 7$ TeV for the $X \rightarrow ZZ$ channel (left) and the $X \rightarrow Z\gamma$ channel (right). Λ is the mass scale of the dimension 6 operator coupling X to gluons.

Note that, for large M_X and $\mathcal{L} = 35 \text{ pb}^{-1}$, the discovery reach can drop as low as $\Lambda \sim 40 \text{ GeV}$. For typical models, this would imply that particles which have been integrated out to generate the higher-dimensional effective operator are in fact lighter than the energy of the hard process, rendering the effective operator analysis inconsistent. Moreover, one might expect additional operators with dimension greater than 6 to provided contributions which are suppressed by additional powers of M_Z^2/Λ^2 ; if this factor is large, then one cannot ignore the effect of these contributions. While these issues would hold with typical models, one can also have models in which the degrees of freedom which have been integrated out are indeed heavy, but have a large multiplicity which serves to reduce Λ (see also [14]). Moreover, the contribution of operators with dimension > 6 depends on the details of the heavy degrees of freedom. Thus, we have plotted the discovery reach even if $\Lambda < M_{Z,X}$ for completeness, but in those cases the effective operator analysis may not be valid.

V. CONCLUSIONS

We see that the best detection prospects arise from the $X \rightarrow Z\gamma$ channel, because the $X \rightarrow ZZ$ channel suffers from the small branching fraction for two Z s to decay to 4 leptons. In both channels Standard Model backgrounds become significant for relatively light M_X

and large luminosities. As expected, the sensitivity of the LHC to resonant X production is greatly enhanced when production through QCD processes is possible. Comparing to the results in [14], we see that if $M_X = 1000$ GeV, then gluon couplings allow the LHC to probe operators suppressed by a mass scale Λ which is 10 times larger than the scale which could be probed if only electroweak couplings were allowed. Note again that our analysis has focused on a 7 TeV center of mass energy; detection prospects would be improved significantly if the center of mass energy were upgraded to $\sqrt{s} = 14$ TeV. For example, for $M_X = 1000$ GeV, a 100 fb^{-1} LHC run at $\sqrt{s} = 14$ TeV will have roughly twice the mass reach of a run at 7 TeV (in either the ZZ or $Z\gamma$ channels).

It is worth noting that a resonance which decays to ZZ and W^+W^- is a characteristic signature of a relatively heavy Higgs boson. A new pseudovector coupling to Standard Model gauge bosons can thus “counterfeit” standard Higgs signals [36]. It may be especially difficult to distinguish these possibilities, since a heavy Higgs decays to light fermions with a very small branching fraction. The features of a pseudovector which can be used to distinguish it from a Higgs include the absence of the $\gamma\gamma$ channel and the fact that production through QCD processes requires the presence of an additional jet. It would be interesting to study in detail the prospects for distinguishing the spin of any resonance which couples to Standard Model gauge bosons.

It is also worth noting that we have focused only on effective coupling operators which are non-trivial when the X and the Standard Model gauge bosons are both on-shell. If these operators vanish (or have very small couplings), then the production and decay of X may be dominated by operators which only yield non-trivial vertices when X is off-shell. This would imply that X production is not associated with a resonance peak. A detailed collider study of such operators would be very interesting.

Acknowledgments

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