Anomaly driven signatures of new invisible physics at the Large Hadron Collider

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Abstract

Many extensions of the Standard Model (SM) predict new neutral vector bosons at energies accessible by the Large Hadron Collider (LHC). We study an extension of the SM with new chiral fermions subject to non-trivial anomaly cancellations. If the new fermions have SM charges, but are too heavy to be created at LHC, and the SM fermions are not charged under the extra gauge field, one would expect that this new sector remains completely invisible at LHC. We show, however, that a non-trivial anomaly cancellation between the new heavy fermions may give rise to observable effects in the gauge boson sector that can be seen at the LHC and distinguished from backgrounds.
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1 Introduction: Mixed Anomalies in Gauge Theory

It is well known that theories in which fermions have chiral couplings with gauge fields suffer from anomalies – a phenomenon of breaking of gauge symmetries of the classical theory at one-loop level. Anomalies make a theory inconsistent (in particular, its unitarity is lost). The only way to restore consistency of such a theory is to arrange the exact cancellation of anomalies between various chiral sectors of the theory. This happens, for example, in the Standard Model (SM), where the cancellation occurs between quarks and leptons within each generation [1, 2, 3].

Another well-studied example is the Green-Schwarz anomaly cancellation mechanism [4] in string theory. In this case the cancellation happens between the anomalous contribution of chiral matter of the closed string sector with that of the open string.1

Particles involved in anomaly cancellation may have very different masses – for example, the mass of the top quark in the SM is much higher than the masses of all other fermions. On the other hand, gauge invariance should pertain in the theory at all energies, including those which are smaller than the mass of one or several particles involved in anomaly cancellation. The usual logic of renormalizable theories tells us that the interactions, mediated by heavy fermions running in loops, are generally suppressed by the masses of these fermions [6]. The case of anomaly cancellation presents a notable counterexample to this famous “decoupling theorem” – the contribution of a priori arbitrary heavy particles should remain unsuppressed at arbitrarily low energies. As was pointed out by D’Hoker and Farhi [7, 8], this is

1Formally, the Green-Schwarz anomaly cancellation occurs due to the anomalous Bianchi identity for the field strength of a 2-form closed string. However, this modification of Bianchi identity arises from the 1-loop contribution of chiral fermions in the open string sector. A toy model, describing microscopically Green-Schwarz mechanism was studied e.g. in [5].
possible because anomalous (i.e. gauge-variant) terms in the effective action have topological nature and are therefore scale independent. As a result, they are not suppressed even at energies much smaller than the masses of the particles producing these terms via loop effects. This gives hope to see some signatures at low energies generated by new high-energy physics.

One possibility is to realize non-trivial anomaly cancellation in the electroweak (EW) sector of the SM. Here the electromagnetic $U(1)$ subgroup is not anomalous by definition. However, the mixed triangular hypercharge $U_Y(1) \times SU(2)^2$ anomalies and gravitational anomalies are non-zero for a generic choice of hypercharges. If one takes the most general choice of hypercharges, consistent with the structure of the Yukawa terms, one sees that it is parametrized by two independent quantum numbers $Q_e$ (shift of hypercharge of left-handed lepton doublet from its SM value) and $Q_q$ (corresponding shift of quark doublet hypercharge). All the anomalies are then proportional to one particular linear combination: $\epsilon = Q_e + 3Q_q$. Interestingly enough, $\epsilon$ is equal to the sum of electric charges of the electron and proton. The experimental upper bound on the parameter $\epsilon$, coming from checks of electro-neutrality of matter is rather small: $\epsilon < 10^{-21}e$ [38, 39]. If it is non-zero, the anomaly of the SM has to be cancelled by additional anomalous contributions from some physics beyond the SM, possibly giving rise to some non-trivial effects in the low energy effective theory.

In the scenario described above the anomaly-induced effects are proportional to a very small parameter, which makes experimental detection very difficult. In this paper we consider another situation, where anomalous charges and therefore, anomaly-induced effects, are of order one. To reconcile this with existing experimental bounds, such an anomaly cancellation should take place between the SM and “hidden” sector, with the corresponding new particles appearing at relatively high energies. Namely, many extensions of the SM add extra gauge fields to the SM gauge group (see e.g. [42] and refs. therein). For example, additional $U(1)$s naturally appear in models in which $SU(2)$ and $SU(3)$ gauge factors of the SM arise as parts of unitary $U(2)$ and $U(3)$ groups (as e.g. in D-brane constructions of the SM [43, 44, 45]). In this paper, we consider extensions of the SM with an additional $U_X(1)$ factor, so that the gauge group becomes $SU(3)_c \times SU(2)_W \times U_Y(1) \times U_X(1)$. As the SM fermions are chiral with respect to the EW group $SU(2)_W \times U_Y(1)$, even choosing the charges for the $U_X(1)$ group so that the triangular $U_X(1)^3$ anomaly vanishes, mixed anomalies may still arise: $U_X(1)U_Y(1)^2$, $U_X(1)^2U_Y(1)$, $U_X(1)SU(2)^2$. In this work we are interested in the situation when only (some of these) mixed anomalies with the electroweak group $SU(2) \times U_Y(1)$ are non-zero. A number of works have already discussed such theories and their signatures (see e.g. [11, 12, 43, 46, 47, 48, 49]).

The question of experimental signatures of such theories at the LHC should be addressed differently, depending on whether or not the SM fermions are charged with respect to the $U_X(1)$ group:

- If SM fermions are charged with respect to the $U_X(1)$ group, and the mass of the new $X$ boson is around the TeV scale, we should be able to see the corresponding
resonance in the forthcoming runs of LHC in e.g., $q\bar{q} \rightarrow X \rightarrow f\bar{f}$. The analysis of this is rather standard \(Z'\) phenomenology, although in this case an important question is to distinguish between theories with non-trivial cancellation of mixed anomalies, and those that are anomaly free.

- On the other hand, one is presented with a completely different challenge if the SM fermions are not charged with respect to the \(U_X(1)\) group. This makes impossible the usual direct production of the \(X\) boson via coupling to fermions. Therefore, the question of whether an anomalous gauge boson with mass \(M_X \sim 1\) TeV can be detected at LHC becomes especially interesting.

A theory in which the cancellation of the mixed \(U_X(1)SU(2)^2\) anomaly occurs between some heavy fermions and Green-Schwarz (i.e. \textit{tree-level gauge-variant}) terms was considered in [49]. The leading non-gauge invariant contributions from the triangular diagrams of heavy fermions, unsuppressed by the fermion masses, cancels the Green-Schwarz terms. The triangular diagrams also produce subleading (gauge-invariant) terms, suppressed by the mass of the fermions running in the loop. This leads to an appearance of dimension-6 operators in the effective action, having the general form \(F_{\mu\nu}^X/A_X^2\), where \(F_{\mu\nu}\) is the field strength of \(X\), \(Z\) or \(W^\pm\) bosons. Such terms contribute to the \(XZZ\) and \(XWW\) vertices. As the fermions in the loops are heavy, such vertices are in general strongly suppressed by their mass. However, motivated by various string constructions, [49] assumed two things: (a) these additional massive fermions are above the LHC reach but not too heavy (e.g., have masses in tens of TeV); (b) there are many such fermions (for instance Hagedorn tower of states) and therefore the mass suppression can be compensated by the large multiplicity of these fermions.

In this paper we consider another possible setup, in which the anomaly cancellation occurs \textit{only within a high-energy sector} (at scales not accessible by current experiments), but at low energies there remain contributions unsuppressed by masses of heavy particles. A similar setup, with completely different phenomenology, has been previously considered in [11, 12].

The paper is organized as follows. We first consider in section 2 the general theory of D’Hoker-Farhi terms arising from the existence of heavy states that contribute to anomalies. We illustrate the theory issues in the following section 3 with a toy model. In section 4 we give a complete set of charges for a realistic \(SU(2) \times U(1)_Y \times U(1)_X\) theory. In section 5 we bring all these elements together to demonstrate expected LHC phenomenology of this theory.

## 2 D’Hoker-Farhi Terms from High Energies

In this Section we consider an extension of the SM with an additional \(U_X(1)\) field. The SM fields are neutral with respect to the \(U_X(1)\) group, however, the heavy fields are charged with respect to the electroweak (EW) \(U_Y(1) \times SU(2)\) group. This leads
to a non-trivial mixed anomaly cancellation in the heavy sector and in this respect our setup is similar to the work [49]. However, unlike the work [49], we show that there exists a setup in which non-trivial anomaly cancellation induces a dimension 4 operator at low energies. The theories of this type were previously considered in [11, 12].

At energies accessible at LHC and below the masses of the new heavy fermions, the theory in question is simply the SM plus a massive vector boson $X$:

$$\mathcal{L} = \mathcal{L}_{SM} - \frac{1}{4g_X^2} |F_X|^2 + \frac{M_X^2}{2} |D\theta_X|^2 + \mathcal{L}_{int}$$  \hspace{1cm} (1)

where $\theta_X$ a pseudo-scalar field, charged under $U_X(1)$ so that $D\theta_X = d\theta_X + X$ remains gauge invariant (Stückelberg field). One can think about $\theta_X$ as being a phase of a heavy Higgs field, which gets “eaten” by the longitudinal component of the $X$ boson. Alternatively, $\theta_X$ can be a component of an antisymmetric $n$-form, living in the bulk and wrapped around an $n$-cycle. The interaction term $\mathcal{L}_{int}$ contains the vertices between the $X$ boson and the $Z, \gamma, W^\pm$:

$$\epsilon^{\mu\nu\lambda\rho} Z_\mu X_\nu \partial_\lambda Z_\rho, \quad \epsilon^{\mu\nu\lambda\rho} X_\mu \partial_\lambda \gamma_\rho, \quad \epsilon^{\mu\nu\lambda\rho} W^+_{\mu} X_\nu \partial_\lambda W^-_\rho.$$  \hspace{1cm} (2)

We wish to generalize these terms into an $SU(2) \times U_Y(1)$ covariant form. One possible way would be to have them arise from

$$\epsilon^{\mu\nu\lambda\rho} X_\mu Y_\nu \partial_\lambda Y_\rho \quad \text{and} \quad \epsilon^{\mu\nu\lambda\rho} \omega_{\nu\lambda\rho}(A^a)$$  \hspace{1cm} (3)

where $\omega_{\nu\lambda\rho}(A^a)$ is the Chern-Simons term, built of the $SU(2)$ fields $A^a_\mu$:

$$\omega_{\nu\lambda\rho}(A^a) = A^a_\nu \partial_\lambda A^a_\rho + \frac{2}{3} \epsilon_{abc} A^a_\nu A^b_\lambda A^c_\rho$$  \hspace{1cm} (4)

However, apart from the desired terms of eq. (2) they contain also terms like $\epsilon^{\mu\nu\lambda\rho} X_\mu \gamma_\nu \partial_\lambda \gamma_\rho$ which is not gauge invariant with respect to the electro-magnetic $U(1)$ group, and thus unacceptable.

To write the expressions of (2) in a gauge-covariant form, we should recall that it is the SM Higgs field $H$ which selects massive directions through its covariant derivative $D_\mu H$. Therefore, we can write the interaction term in the following, explicitly $SU(2) \times U_Y(1) \times U_X(1)$ invariant, form:

$$\mathcal{L}_{int} = c_1 \frac{H^\dagger DH}{|H|^2} D\theta_X F_Y + c_2 \frac{HF_W D\theta_X}{|H|^2} D\theta_X$$  \hspace{1cm} (5)

The coefficients $c_1, c_2$ are dimensionless and can have arbitrary values, determined entirely by the properties of the high-energy theory. In eq. (5) we use the differential form notation (and further we omit the wedge product symbol $\wedge$) to keep expressions
1 is such that triangular anomalies cancel separately for the mixed anomalies occurs only between chiral with respect to the $U$ of the charges in Table 1: A simple choice of charges for all fermions, leading to the low-energy effective action (8). The charges are chosen in such a way that all gauge anomalies cancel. The cancellation of $U(1)^3_A$ and $U(1)^3_B$ anomalies happens for any value of $e_i, q_i$. Cancellation of mixed anomalies requires $q_2 = \frac{q_1(e_1^2 - e_2^2)}{2(e_3^2 - e_4^2)}$.

more compact. We will often call the terms in eq. (5) as the D'Hoker-Farhi terms [7, 8].

What can be the origin of the interaction terms (5)? The simplest possibility would be to add to the SM several heavy fermions, charged with respect to $SU(2) \times U_Y(1) \times U_X(1)$. Then, at energies below their masses the terms (5) will be generated. Below, we illustrate this idea in a toy-model setup.

Consider a theory with a set of chiral fermions $\psi_{1,2}$ and $\chi_{1,2}$, charged with respect to the gauge groups $U(1)_A \times U(1)_B$. As the fermions are chiral, they can obtain masses only through Yukawa interactions with both $\Phi_1$ and $\Phi_2$ scalar fields. $\Phi_1$ is charged with respect to $U(1)_B$, and $\Phi_2$ is charged with respect to $U(1)_A$:

$$\mathcal{L}_{Yukawa} = i \sum_{i=1,2} \bar{\psi}_i \not{\partial} \psi_i + (f_1 v_1) \bar{\psi}_1 e^{i\gamma^5 \theta_B} \psi_1 + (f_2 v_2) \bar{\psi}_2 e^{-i\gamma^5 \theta_B} \psi_2$$

$$+ i \sum_{i=1,2} \bar{\chi}_i \not{\partial} \chi_i + (\lambda_1 v_2) \bar{\chi}_1 e^{i\gamma^5 \theta_A} \chi_1 + (\lambda_2 v_2) \bar{\chi}_2 e^{-i\gamma^5 \theta_A} \chi_2 + h.c. \quad (6)$$

Here we have taken $\Phi_1$ in the form $\Phi_1 = v_1 e^{i\theta_B}$, where $v_1$ is its vacuum expectation value (VEV) and $\theta_B$ is charged with respect to the $U(1)_B$ group with charge $2q_1$, and $\Phi_2 = v_2 e^{i\theta_A}$, where $\theta_A$ is charged with respect to $U(1)_A$ group with charge $e_3 - e_4$.

The structure of the Yukawa terms restricts the possible charge assignments, so that the fermions $\psi_{1,2}$ should be vector-like with respect to the group $U(1)_B$ and chiral with respect to the $U(1)_B$ (and vice versa for the fermions $\chi_{1,2}$). The choice of the charges in Table 1 is such that triangular anomalies $[U(1)_A]^3$ and $[U(1)_B]^3$ cancel separately for the $\psi$ and $\chi$ sector for any choice of $e_i, q_i$. The cancellation of mixed anomalies occurs only between $\psi$ and $\chi$ sectors. It is instructive to analyze it at energies below the masses of all fermions. The terms in the low-energy effective action, not suppressed by the scale of fermion masses are given by

$$S_{cs} = \int \frac{(e_1^2 - e_2^2) q_4}{16\pi^2} \theta_B F_A \wedge F_A + \frac{(e_3^2 - e_4^2)(2q_2)}{16\pi^2} \theta_A F_A \wedge F_B + \alpha A \wedge B \wedge F_A \quad (7)$$

The diagrammatic expressions for the first two terms are shown in Fig. 1, while the Chern-Simons (CS) term is produced by the diagrams of the type presented in Fig. 2.
The contribution to the CS term $A \wedge B \wedge F_A$ comes from both sets of fermions. Only fermions $\psi$ contribute to the $\theta_B$ terms and only fermions $\chi$ couple to $\theta_A$ and thus contribute to $\theta_A F_A \wedge F_B$. Notice that while coefficients in front of the $\theta_A$ and $\theta_B$ terms are uniquely determined by charges, the coefficient $\alpha$ in front of the CS term is regularization dependent. As the theory is anomaly free, there exists a choice of $\alpha$ such that the expression (7) becomes gauge-invariant with respect to both gauge groups. Notice, however, that in the present case $\alpha$ cannot be zero, as $\theta_A F_A \wedge F_B$ and $\theta_B F_A \wedge F_A$ have gauge variations with respect to different groups. For the choice of charges presented in Table 1, the choice of $\alpha$ is restricted such that expression (7) can be written in an explicitly gauge-invariant form:

$$S_{cs} = \int \kappa D\theta_A \wedge D\theta_B \wedge F_A$$

where the relation between the coefficient $\kappa$ in front of the CS term and the fermion charges is given by

$$\alpha \equiv \kappa = \frac{q_1(e_1^2 - e_2^2)}{16\pi^2}$$

For the anomaly cancellation, it is also necessary to impose the condition

$$q_2 = \frac{q_1(e_1^2 - e_2^2)}{2(e_3^2 - e_4^2)}$$

as indicated in Table 1.

The term (8) was obtained by integrating out heavy fermions (Table 1). The resulting expression is not suppressed by their mass and contains only a dimensionless coupling $\kappa$. Unlike the case of [7, 8], the anomaly was cancelled entirely among the fermions which we had integrated out. The expression (8) represents therefore an apparent counterexample of the “decoupling theorem” [6]. Note that the CS term (8) contains only massive vector fields. This effective action can only be valid at energies above the masses of all vector fields and below the masses of all heavy fermions, contributing to it. However, masses of both types arise from the same Higgs fields. Therefore a hierarchy of mass scales can only be achieved by making gauge couplings smaller than Yukawa couplings. On the other hand, the CS coefficient $\kappa$ is proportional to the (cube of the) gauge couplings. Therefore we can schematically write a dimensionless coefficient $\kappa \sim (M_V/M_f)^3$, where $M_V$ is the mass of the vector fields and $M_f$ is the mass of the fermions (with their Yukawa couplings $\sim 1$). In the limit when $M_f$ is sent to infinity, while keeping $M_V$ finite, the decoupling theorem holds, as the CS terms get suppressed by the small gauge coupling constant. However, a window of energies $M_V \lesssim E \lesssim M_f$, at which the term (8) is applicable, always remains and this opens interesting phenomenological possibilities, which are absent in the situation when the corresponding terms in the effective action are suppressed as $E/M_f$ (as in [6]) and not as $M_V/M_f$.  

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Finally, it is also possible that the fermion masses are not generated via the Higgs mechanism, (e.g. coming from extra dimensions) and are not directly related to the masses of the gauge fields. In this case, the decoupling theorem may not hold and new terms can appear in a wide range of energies (see e.g. [9, 10] for discussion).

3 A Standard Model Toy Example

Let us now generalize this construction to the case of interest, when one of the scalar fields generates mass for the chiral fermions and is the SM Higgs field, while at the same time the masses of all new fermions are higher than about 10 TeV.

Note that previously, in the theory described by (6) the mass terms for fermions were diagonal in the basis $\psi, \chi$ and schematically had the form $m_1 \bar{\psi} \psi + m_2 \bar{\chi} \chi$. To make both masses for $\psi$ and $\chi$ heavy (i.e., determined by the non-SM scalar field), while still preserving a coupling of the fermions with the SM Higgs, we consider a non-diagonal mass term which (schematically) has the following form:

$$L_{\text{mass}} = m \bar{\psi} \psi + M (\bar{\psi} \chi + \bar{\chi} \psi)$$

Computing the eigenvalues of the mass matrix, we find that the two mass eigenstates have masses $M \pm \frac{m}{2}$ (in the limit $m \ll M$).

Now, we consider the case when the mass terms, similar to those of Eq. (11) are generated through the Higgs mechanism. We introduce two complex scalar Higgs
Table 2: Charge assignment for the $U_Y (1) \times U_X (1)$ with 4 Dirac fermions. Charges of the scalar fields $H$ and $\Phi$ are equal to (1,0) and (0,1), respectively.

| $Q_X$ | $Q_Y$ |
|-------|-------|
| $x$   | $y$   |
| $x$   | $y + 1$ |
| $x - 1$ | $y_1$ |
| $x + 1$ | $y_2$ |
| $x + 1$ | $y_1$ |
| $x - 1$ | $y_2$ |

Table 2: Charge assignment for the $U_Y (1) \times U_X (1)$ with 4 Dirac fermions. Charges of the scalar fields $H$ and $\Phi$ are equal to (1,0) and (0,1), respectively.

fields: $H = H_1 + iH_2$ and $\Phi = \Phi_1 + i\Phi_2$. $H$ is charged with respect to the $U_Y (1)$ only (with charge 1), while $\Phi$ is $U_Y (1)$ neutral, but has charge 1 with respect to the $U_X (1)$. We further assume that both Higgs fields develop non-trivial VEVs:

$$\langle H \rangle = v \quad ; \quad \langle \Phi \rangle = V \quad ; \quad v \ll V \quad (12)$$

Then, we may write

$$H = ve^{i\theta_H} ; \quad \Phi = Ve^{i\theta_X} \quad (13)$$

neglecting physical Higgs field excitations ($H(x) = (v + h(x))e^{i\theta_H}$, etc.).

Let us suppose that the full gauge group of our theory is just $U_Y (1) \times U_X (1)$. Consider 4 Dirac fermions ($\psi_1, \psi_2, \chi_1, \chi_2$) with the following Yukawa terms, leading to the Lagrangian in the form, similar to (11):

$$L_{\text{Yukawa}} = m_1 \bar{\psi}_1 e^{i\gamma^5 \theta_H} \psi_1 + M_1 (\bar{\psi}_1 e^{i\gamma^5 \theta_X} \chi_1 + c.c.) + M_2 (\bar{\psi}_2 e^{-i\gamma^5 \theta_X} \chi_2 + c.c.) \quad (14)$$

Here we introduced masses $m_1 = f_1 v$ and $M_{1,2} = F_{1,2} V$, with $f_1$ and $F_{1,2}$ the corresponding Yukawa couplings.

The choice of fermion charges is dictated by the Yukawa terms (14). The $\psi$ fermions are vector-like with respect to $U_X (1)$ group, although chiral with respect to the $U_Y (1)$. The fermions $\psi_1, \chi_1$ (and $\psi_2, \chi_2$) have charges with respect to $U_Y (1)$ group, such that

$$Q_Y (\psi_1 L) = Q_Y (\chi_1 R) \quad \text{and} \quad Q_Y (\psi_1 R) = Q_Y (\chi_1 L) \quad (15)$$

and similarly for the pair $\psi_2, \chi_2$. Unlike $\psi_1$, the fermions $\psi_2$ do not have Yukawa term $m_2 \bar{\psi}_2 e^{i\gamma^5 \theta_H} \psi_2$, as this would make the choice of charges too restrictive and does not allow us to generate terms similar to (8). The resulting charge assignment is shown in Table 2.

It is clear that the triangular anomalies $XXX$ and $YYY$ cancel as there is equal number of left and right moving fermions with the same charges. Let us consider the mixed anomaly $XYY$. The condition for anomaly cancellation is given by

$$A_{XYY} = \sum Q_X^L (Q_Y^L)^2 - Q_X^R (Q_Y^R)^2 = y_1^2 + y_2^2 - 1 - 2y - 2y^2 = 0 \quad (16)$$

The other mixed anomaly $XXY$ is proportional to

$$A_{XXY} = 1 - y_1 + y_2 + 2x(-2y + y_1 + y_2 - 1) = 0 \quad (17)$$
Table 3: An example of charge assignments for the $U_Y(1) \times U_X(1)$ of the 4 Dirac fermions. The anomaly coefficient $\kappa$ (Eq. (21)) is nonzero and equal to 6.

|   | $\psi_1$ | $\psi_2$ | $\chi_1$ | $\chi_2$ |
|---|----------|----------|----------|----------|
| $Q_X$ | 1        | 1        | 1        | 1        |
| $Q_Y$ | 1        | 2        | $-1$     | 2        |

and should also cancel.

In analogy with the toy-model, described above, Table 3 presents an anomaly free assignment for which the mixed anomalies cancel only between the $\psi$ and $\chi$ sectors and lead to the following term in the effective action (similar to (8)):

$$\mathcal{L}_A = \kappa D\theta_H \wedge D\theta_X \wedge F_Y$$

Here the parameter $\kappa$ is defined by the $XYY$ anomaly in the $\psi$ or $\chi$ sector, in analogy with Eq. (9):

$$\kappa = -\frac{x (-y_1^2 + y_2^2 + 2y + 1)}{32\pi^2}$$

To have $\kappa \neq 0$ we had to make two mass eigenstates in the sector $\psi_2, \chi_2$ degenerate and equal to $M_2$. The charges $x, y$ become then arbitrary, while $y_{1,2}$ should satisfy the constraints (16) and (17). It is easy to see that indeed this can be done together with the inequality $\kappa \neq 0$. The solution gives:

$$y_1 = \frac{4yx^2 - 4yx - 4x - y}{4x^2 + 1}; \quad y_2 = \frac{4yx^2 + 4x^2 + 4yx - y - 1}{4x^2 + 1}$$

The choice (20) leads to the following value of $\kappa$:

$$\kappa = -\frac{2x (4x^2 - 1) ((8y + 4)x^2 + 8y(y + 1)x - 2y - 1)}{(4x^2 + 1)^2}$$

One can easily see that $\kappa$ is non-zero for generic choices of $x$ and $y$. One such a choice is shown in Table 3 (recall that all $U_X(1)$ charges are normalized so that $\theta_X$ has $Q_X(\theta_X) = 1$ and all $U_Y(1)$ charges are normalized so that $Q_Y(\theta_H) = 1$).

To make the anomalous structure of the Lagrangian (14) more transparent, we can perform a chiral change of variables, that makes the fermions vector-like. Namely, let us start with the term $m_1 \bar{\psi}_1 e^{i\theta_H \gamma^5} \psi_1$. We want to perform a change of variables to a new field $\tilde{\psi}$, which will turn this term into $m_1 \bar{\tilde{\psi}} \tilde{\psi}_1$. This is given by

$$\left(\begin{array}{c} \psi_{1L} \\ \psi_{1R} \end{array}\right) \rightarrow \left(\begin{array}{c} e^{-\frac{i}{2}\theta_H \gamma^5} \psi_{1L} \\ e^{\frac{i}{2}\theta_H \gamma^5} \psi_{1R} \end{array}\right) \quad \text{or} \quad \psi \rightarrow e^{-\frac{i}{2}\gamma^5 \theta_H} \tilde{\psi}$$
Table 4: Charge assignment for the $SU(2) \times U_Y(1) \times U_X(1)$ gauge group. Fermions, which are doublets with respect to the $SU(2)$ are marked with the superscript $a$. Charges of the SM Higgs field $H$ and of the heavy Higgs $\Phi$ are equal to (1,0) and (0,1) with respect to $U_Y(1) \times U_X(1)$.

| $Q_X$ | $x$ | $x$ | $x$ | $x$ | $x - 1$ | $x + 1$ | $x + 1$ | $x - 1$ |
|-------|-----|-----|-----|-----|---------|---------|---------|---------|
| $Q_Y$ | $y$ | $y + 1$ | $y_1$ | $y_2$ | $y + 1$ | $y$ | $y_2$ | $y_1$ |

so that the Yukawa term becomes $m_1 \bar{\psi}_1 \tilde{\psi}_1$. The field $\tilde{\psi}_1$ has vector-like charge $x$ with respect to $U_X(1)$ and vector-like charge $y + \frac{1}{2}$ with respect to $U_Y(1)$. As the change of variables is chiral, it introduces a Jacobian $J_{\psi_1}$ [50]. The transformation (22) turns the term $M_1(\bar{\psi}_1 e^{i\gamma^5\theta_x} \chi_1 + c.c.)$ into $M_1(\bar{\psi}_1 e^{i\gamma^5(\theta_x - 2\theta_5)} \chi_1 + c.c.)$. By performing a change of variables from $\chi_1$ to $\tilde{\chi}_1$,

$$\chi_1 \to e^{-\frac{i}{2} \gamma^5(\theta_x - \theta_5)} \tilde{\chi}_1,$$

we make the sector $\tilde{\psi}_1, \tilde{\chi}_1$ fully vector-like, and generate two anomalous Jacobians $J_{\psi_1}$ and $J_{\chi_1}$. Similarly, for the last term in eq. (14), we perform the change of variables $\chi_2 \to e^{i\theta_x/2} \chi_2$ and $\bar{\psi}_2 \to e^{i\theta_x/2} \bar{\psi}_2$, generating two more Jacobians. By computing the Jacobians, one can easily show that performing the above change of variables for all 4 fermions, we arrive to a vector-like Lagrangian with the additional term (18).

### 4 Charges in a Realistic $SU(2) \times U_Y(1) \times U_X(1)$ Model

The above example shows us how to construct a realistic model of high-energy theory, whose low-energy effective action produces the terms (5). We consider the following fermionic content (iso-index $a = 1, 2$ marks $SU(2)$ doublets): two left $SU(2)$ doublets $\psi_{1L}^a$ and $\chi_{2L}^a$, two right $SU(2)$ doublets $\psi_{2R}^a$ and $\chi_{1R}^a$, as well as two left $SU(2)$ singlets $\psi_{2L}$ and $\chi_{1L}$, and two right $SU(2)$ singlets $\psi_{1R}$ and $\chi_{2R}$. The corresponding charge assignments are shown in Table 4.

The Yukawa interaction terms have the form:

$$\mathcal{L}_{\text{Yukawa}} = f_1(\bar{\psi}_{1L}^a H_a)\psi_{1R} + F_1 \left( \bar{\psi}_{1L}^a (\Phi_1 - i\gamma^5 \Phi_2) \chi_{1R}^a + c.c. \right)$$

$$+ F_2 \left( \bar{\psi}_{2R}^a (\Phi_1 + i\gamma^5 \Phi_2) \chi_{2L}^a + c.c. \right)$$

$$+ \tilde{F}_1 \left( \bar{\psi}_{1R} (\Phi_1 - i\gamma^5 \Phi_2) \chi_{1L} + c.c. \right) + \tilde{F}_2 \left( \bar{\psi}_{2L} (\Phi_1 + i\gamma^5 \Phi_2) \chi_{2R} + c.c. \right)$$

(24)

where $H$ is the SM Higgs boson and $\Phi_{1,2}$ are $SU(2) \times U(1)_Y$ singlets. Here again $\langle H \rangle = v \ll \langle \Phi \rangle$, and all states have heavy masses $\sim F\langle \Phi \rangle$ (plus possible corrections of order $O(fv)$).
Table 5: Explicit charge assignment for the $SU(2) \times U_Y(1) \times U_X(1)$ gauge group.

|       | $\psi_1$ | $\psi_2$ | $\chi_1$ | $\chi_2$ |
|-------|----------|----------|----------|----------|
| $Q_X$ | $-\frac{1}{6}$ | $-\frac{1}{6}$ | $-\frac{1}{6}$ | $-\frac{1}{6}$ |
| $Q_Y$ | $\frac{1}{2}$ | $\frac{1}{2}$ | $\frac{1}{2}$ | $\frac{1}{2}$ |

Figure 3: $\Gamma_{XZZ}$ and $\Gamma_{XZ\gamma}$ interaction vertices, generated by (25)

The anomaly analysis is similar to the one performed in the previous section. The only difference being of course two isospin degrees of freedom in the $SU(2)$ doublets. The resulting choice of charges is shown in Table 5 (we do not write the general expression as it is too cumbersome and provides only an example when $x = -Q_H/6$, $y = Q_\phi/2$). One may check that for this choice of charges the resulting coefficients $c_{1,2}$ in the interaction terms (5) are non-zero, which leads to interesting phenomenology to be discussed in the next section.

5 Phenomenology

The analysis of the previous sections puts us in position to now discuss the phenomenology of the $X$ boson. To do this, we first detail the relevant interactions it has with the SM gauge bosons.

The first term in (5) generates two interaction vertices: $XZZ$ and $XZ\gamma$ (Fig.3). In the EW broken phase one can think of the first term in expression (5) as being simply

$$\mathcal{L}_{XZY} = c_1(d\theta_Z + Z)F_YD\theta_X + \mathcal{O}\left(\frac{\partial h}{v}\right)$$ (25)

where we parametrized the Higgs doublet as

$$H = e^{i(\tau^+ \theta_+(x)+\tau^- \theta_-(x)+\frac{1}{2}+\tau^3)\theta_Z} \begin{pmatrix} 0 \\ v + h(x) \end{pmatrix}$$ (26)
Here the phases $\theta_{\pm}$, $\theta_{Z}$ will be “eaten” by $W^\pm$ and $Z$ bosons correspondingly, $v$ is the Higgs VEV and the real scalar field $h$ is the physical Higgs field.

The vertices $\Gamma_{XZZ}$ and $\Gamma_{XZ\gamma}$ are given correspondingly by

$$
\Gamma^{\mu\nu\rho}_{XZZ}(k_1, k_2 | k_3) = \frac{1}{2} c_1 \sin \theta_w \epsilon^{\mu\nu\lambda\rho}(k_{2\lambda} - k_{1\lambda}) 
$$

$$
\Gamma^{\mu\nu\rho}_{XZ\gamma}(k_1, k_2 | k_3) = c_1 \cos \theta_w \epsilon^{\mu\nu\lambda\rho} k_{2\rho}
$$

(27)

Similarly to above one can analyze the second term in (5). It leads to the interaction $XW^+W^-:

$$
\Gamma^{\mu\nu\rho}_{XW^+W^-}(k_1, k_2 | k_3) = c_2 \epsilon^{\mu\nu\lambda\rho}(k_{2\lambda} - k_{1\lambda})
$$

(28)

The most important relevant fact to phenomenology is that the $X$ boson is produced by and decays into SM gauge bosons. We shall discuss in turn the production mechanisms and the decay final states of the $X$ boson and then estimate the discovery capability at colliders.

### 5.1 Production of $X$ boson

Producing the $X$ boson must proceed via its coupling to pairs of SM gauge bosons. One such mechanism is through *vector-boson fusion*, where two SM gauge bosons are radiated off initial state quark lines and fused into an $X$ boson:

$$
pp \rightarrow qq'VV' \rightarrow qq'X \text{ or } VV' \rightarrow X \text{ for short,}
$$

(29)

where $VV'$ can be $W^\pm W^-$, $ZZ$ or $Z\gamma$. This production mechanism was studied in ref. [49]. One of the advantages is that if the decays of $X$ are not much different than the SM, the high-rapidity quarks that accompany the event can be used as “tagging jets” to help separate signal from the background. This production mechanism is very similar to what has been exploited in the Higgs boson literature.

A second class of production channels is through *associated production*:

$$
pp \rightarrow qq' \rightarrow V^* \rightarrow XV'
$$

(30)

where an off-shell vector boson $V^*$ and the final state $V'$ can be any of the SM electroweak gauge bosons: $XZ$, $XW^\pm$ or $X\gamma$. It turns out that this production class has a larger cross-section than the vector boson fusion class. This is opposite to what one finds in SM Higgs phenomenology, where $VV' \rightarrow H$ cross-section is by $O(10^2)$ greater than $HV'$ associated production. The reason for this is that both vector bosons can be longitudinal when scattering into $H$, thereby increasing the $VV' \rightarrow H$ cross-section over $HV'$. This is not the case for the $X$ boson production, in which only one longitudinal boson can be present at the vertex. This leads to a suppression by $\sim (\sqrt{s}/M_V)^2$ of the process (29) as opposed to the similar process for the Higgs boson. For LHC energies ($\sqrt{s} \sim 10$ TeV) this suppression is of the order $10^{-4}$. 13
Without special longitudinal enhancements, the two body final state $XV'$ dominates over the three-body final state $qq'X$, which makes the associated production (30) about 2 orders of magnitude stronger than the corresponding vector-boson fusion. As we shall see below, the decays of the $X$ boson are sufficiently exotic in nature that background issues do not change the ordering of the importance of these two classes of diagrams. Thus, we focus our attention on the associated production $XV'$ to estimate collider sensitivities.

In figs. 4 and 5 we plot the production cross-sections of $XV$ for various $V = W^\pm, Z, \gamma$ at $\sqrt{s} = 14$ TeV $pp$ LHC, $\sqrt{s} = 2$ TeV $p\bar{p}$ Tevatron and $\sqrt{s} = 200$ GeV $e^+e^-$ LEP.

### 5.2 Decays of $X$ boson

The $X$ boson decays primarily via its couplings to SM gauge boson pairs. The important decay channels are computed from the interaction vertices computed above.
The corresponding decay widths are:

\[
\Gamma_{X \rightarrow ZZ} = \frac{c_1^2 \sin^2 \theta_w M_X^3}{192\pi M_Z^2} \left(1 - \frac{4M_Z^2}{M_X^2}\right)^{5/2} \approx c_1^2 \left(\frac{45 \text{ GeV}}{\text{TeV}}\right)^3 + \ldots, \\
\Gamma_{X \rightarrow W^+W^-} = \frac{c_2 M_X^2}{48\pi M_W^2} \left(1 - \frac{4M_W^2}{M_X^2}\right)^{5/2} \approx c_2^2 \left(\frac{1.03 \text{ TeV}}{\text{TeV}}\right)^3 + \ldots \quad (31) \\
\Gamma_{X \rightarrow Z\gamma} = \frac{c_1^2 \cos^2 \theta_w M_X^3}{96\pi M_Z^2} \left(1 - \frac{M_Z^2}{M_X^2}\right)^3 \left(1 + \frac{M_Z^2}{M_X^2}\right) \approx c_1^2 \left(\frac{307 \text{ GeV}}{\text{TeV}}\right)^3 + \ldots.
\]

where \ldots denote corrections of the order \((M_V/M_X)^2\). The interaction term of eq. (25) also allows interaction of the \(X\) boson with \(ZH\) and \(\gamma H\), which are generically small.

At leading order in \(M_Z/M_X\) the decay width into \(Z\gamma\) exceeds that of \(ZZ\) by

\[
\frac{\Gamma_{X \rightarrow Z\gamma}}{\Gamma_{X \rightarrow ZZ}} = \frac{2\cos^2 \theta_w}{\sin^2 \theta_w} \approx 6.7 \quad (32)
\]

The branching ratio into \(WW\) is the largest over much of parameter space where \(c_2 \gtrsim c_1\), and exceeds that of \(ZZ\) by

\[
\frac{\Gamma_{X \rightarrow W^+W^-}}{\Gamma_{X \rightarrow ZZ}} = \frac{4}{\sin^2 \theta_w} \frac{c_2^2}{c_1^2} \approx 17.4 \frac{c_2^2}{c_1^2}. \quad (33)
\]
Figure 6: Branching fractions of $X$ boson decays into $W^+W^-$ (blue), $ZZ$ (yellow-green) and $Z\gamma$ (purple) as a function of $c_2/c_1$ assuming $M_X \gg M_Z$.

This ratio depends on the a priori unknown ratio of couplings $c_2/c_1$. In Fig. 6 we plot the branching fractions of $X$ into the $WW$ (blue), $Z\gamma$ (purple) and $ZZ$ (yellow-green) as a function of $c_2/c_1$.

Let us compare decay widths (31) with analogous expressions from [49]. Schematically, decay widths can be obtained in our case as

$$\Gamma_{X \rightarrow VV} \sim c_{1,2}^2 \frac{M_X^3}{M_V^2}$$

where we denote by $V$ both $Z$ and $W^\pm$ vector bosons and $M_V = \{M_Z, M_W\}$. In case of setup of Ref. [49] the interaction is the dimension 6 operators, suppressed by the cutoff scale $\Lambda_X^2$. Therefore, the decay width is suppressed by $\Lambda_X^4$ and the whole expression is given by

$$\Gamma_{X \rightarrow VV} \sim \frac{M_X^4}{\Lambda_X^4} \frac{M_X^3}{M_V^2} \frac{M_X^4}{M_X^4} = \frac{M_X^3 M_V^2}{\Lambda^4}$$

The presence of the factor $\frac{M_X^4}{M_X^4}$, appearing in the first equation of (35), can be explained as follows. The vector boson current is conserved in the interaction, generated by the higher-dimensional operators of Ref. [49]. Therefore the corresponding probability for emitting on-shell $Z$ or $W$ boson is suppressed by the $(\frac{M_V}{E})^4$ where the energy $E \sim M_X$. In case of the interaction (5) the vector current is not conserved in the vertex and therefore such a suppression does not appear.
5.3 Collider Searches

Combining the various production modes and branching fractions yields many permutations of final states to consider at high energy colliders. All permutations, after taking into account $X$ decays, give rise to three vector boson final states such as $ZZZ$, $W^+W^−γ$, etc. The collider phenomenology associated with these kinds of final states is interesting, and we focus on a few aspects of it below.

Our primary interest will be to study how sensitive the LHC is to finding this kind of $X$ boson. The limits that one can obtain from LEP 2 and Tevatron are well below the sensitivity of the LHC, and so we forego a more thorough analysis of their constraining power. Briefly, in the limit of no background, the Tevatron cannot do better than the mass scale at which at least a few events are produced. This implies from fig. 4b that $M_X \gtrsim 750$ GeV (for $c_i = 1$) is inaccessible territory to the Fermilab with up to 10 fb$^{-1}$ of integrated luminosity. The LHC can do significantly better than this, as we shall see below.

Moving to the LHC, the energy is of course an important increase as is the planned luminosity. After discovery is made a comprehensive study programme to measure all the final states, and determine production cross-sections and branching ratios would be a major endeavor by the experimental community. However, the first step is discovery. In this section we demonstrate one of the cleanest and most unique discovery modes to this theory. As has been emphasized earlier and in ref. [43], the $X \rightarrow γZ$ decay mode is especially important for this kind of theory. Thus, we study that decay mode. Consulting the production cross-sections results for LHC, we find that producing the $X$ in association with $W^\pm$ gives the highest rate. Thus, we focus our attentions on discovering the $X$ boson through $XW^\pm$ production followed by $X \rightarrow γZ$ decay.

The $γZW^\pm$ signature is an interesting one since it involves all three electroweak gauge bosons. If the $Z$ decays into leptons, it is especially easy to find the $X$ boson mass through the invariant mass reconstruction of $γl^+l^−$. The additional $W$ is also helpful as it can be used to further cut out background by requiring an additional lepton if the $W$ decays leptonically, or by requiring that two jets reconstruct a $W$ mass.

In our analysis, we are very conservative and only consider the leptonic decays of the $Z$ and the $W$. Thus, after assuming $X \rightarrow γZ$ decay, 1.4 percent of $γZW^\pm$ turn into $γl^+l^−l'^\pm$ plus missing $E_T$ events. These events have very little background when cut around their kinematic expectations. For example, if we assume $M_X = 1$ TeV we find negligible background while retaining 0.82 fraction of all signal events when we making kinematic cuts $η(γ,l) < 2.5$, $m_{l^+l^−} = m_Z \pm 5$ GeV, $p_T(γ) > 50$ GeV, $p_T(l^+,l^−,l') > 10$ GeV, missing $E_T$ greater than 10 GeV and $m_{γl^+l^−} > 500$ GeV. Thus, for 10 fb$^{-1}$ of integrated luminosity at the LHC, when $c_i = 1$ ($c_i = 0.1$) we get at least five events of this type, $γl^+l^−l'$ plus missing $E_T$, if $M_X > 4$ TeV ($M_X > 2$ TeV). This would be a clear discovery of physics beyond the SM and would point to a new
resonance, the $X$ boson.

One subtlety for this signal is the required separation of the leptons from the $Z$ decay in order to distinguish two leptons and be able to reconstruct the invariant mass well. The challenge arises because the $Z$ is highly boosted if its parent particle has mass much greater than $m_Z$, and thus the subsequent leptons from $Z$ decays are highly boosted and collimated in the detector. One does not expect this to be a problem for $Z \to \mu^+\mu^-$ decays, as muon separation is efficient. Separation of electron and positron in the electromagnetic calorimeter in highly boosted $Z \to e^+e^-$ final states is expected to be more challenging. We do not attempt to give precise numbers of separability for $e^+e^-$. Instead, we only make two relevant comments. First, one is safe restricting to muons. Second, once separability of $e^+e^-$ is better understood, it can be compared with the kinematic distributions of this example to estimate the number of events that are cut out due to the inability to resolve $e^+e^-$. In Fig. 7 we show the $\Delta R$ separation of $e^+e^-$ for a parent $M_X = 500$ GeV, 1 TeV and 2 TeV. For example, if it turns out that $\Delta R > 0.2 (0.1)$ is required, then one can expect about $2/3 (1/4)$ of the $e^+e^-$ events are cut out by this separation criterion.

After discovery, in addition to doing a comprehensive search over all possible final states, each individual final state will be studied carefully to see what evidence exists for the spin of the $X$ boson. The topology of $\gamma W^\pm$ exists within the SM for $HW^\pm$ production followed by $H \to \gamma Z$ decays. However, the rate at which this happens is very suppressed even for the most optimal mass range of the Higgs boson [51]. A heavy resonance that decays into $\gamma Z$ would certainly not be a SM Higgs boson, but nevertheless a scalar origin would be considered if a signal were found. Careful studying of angular correlations among the final state particles can help determine this question directly. For example, distinguishing between the scalar and vector spin possibilities of the $X$ boson is possible by carefully analyzing the photon’s
cos $\theta$, distribution with respect to the $X$ boost direction in $X \rightarrow \gamma Z$ decays in the rest frame of the $X$. If $X$ is a scalar its distribution is flat in $\cos \theta$, whereas if it is a vector it has a non-trivial dependence on $\cos \theta$. With enough events (several hundred) this distribution can be filled in, and the spin of the $X$ resonance can be discerned among the possibilities.

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