



Universal Extra-Dimensional models with boundary terms: Probing at the LHC

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Abstract

In universal extra-dimensional models a conserved Z_2 parity stabilizes the lightest Kaluza-Klein particle, a dark-matter candidate. Boundary-localized kinetic terms, in general, do not preserve this symmetry. We examine, in the presence of such terms, the single production of Kaluza-Klein excitations of the neutral electroweak gauge bosons and their decay to zero-mode fermion-antifermion pairs. We explore how experiments at the Large Hadron Collider constrain the boundary-localized kinetic terms for different compactification radii.

Keywords:

Extra dimension, Boundary-localized terms

1. Introduction

The Large Hadron Collider (LHC) at CERN is a valuable window for exploring particle physics models. Here we look at a class of models where *all* particles have access to an extra dimension [1]. The models which we examine can be termed ‘non-minimal universal extra-dimensions’ (nmUED) for reasons discussed below. We show that LHC running at 8 TeV may exclude significant segments of the parameter space of this class of models.

We consider one extra spacelike dimension, y , which is flat and compactified. This coordinate runs from 0 to $2\pi R$, where R is the radius of compactification. All particles – scalars, spin-1/2 fermions, and gauge bosons – are five-dimensional fields or equivalently towers of four-dimensional Kaluza-Klein (KK) states. The n -th level KK states for all particles have the same mass n/R . Further, to retain fermion chirality a Z_2 symmetry ($y \leftrightarrow -y$) is imposed. Thus the extra dimension is compactified on an orbifold S^1/Z_2 with $y \equiv y + \pi R$. This symmetry results in a conserved KK-parity given by $(-1)^n$ where n is the KK-level. The standard model (SM) particles have $n = 0$ and are of even parity while the KK-states of the first level are odd. KK-parity en-

sures that the lightest $n = 1$ particle is absolutely stable and hence a dark matter candidate, the Lightest Kaluza-Klein Particle (LKP). This constitutes the Universal Extra Dimension (UED) Model.

S^1/Z_2 has two fixed points at $y = 0$ and $y = \pi R$. One can admit additional interaction terms between the KK-states at these points. These are commonly introduced as counterterms for 5-dimensional loop-induced effects. In the minimal Universal Extra-Dimensional Models (mUED) [2] these terms are so chosen that 5-dimensional loop contributions at the cutoff scale Λ are exactly compensated and the corrections, e.g., logarithmic contributions to masses of KK particles, can be taken to be zero at Λ . These contributions remove the mass degeneracy among states at the same KK-level n .

In this work [3] we examine non-minimal UED (nmUED). Here the boundary terms are not restricted to the special choice in mUED. The two departures are: the two boundary terms are allowed to be unequal – this breaks a remnant Z_2 symmetry which exchanges $y \longleftrightarrow (y - \pi R)$ – and also the strengths of the terms are not the special values chosen in mUED. The breaking of the Z_2 symmetry has far-reaching consequences. For example, the $n = 1$ KK-modes of the neutral gauge

bosons, B^1 and W_3^1 , may be produced singly at the LHC and can decay to zero-mode fermion-antifermion pairs. We explore the prospects of detecting a signal of KK-particles through this route at the LHC.

2. Boundary-localized terms, nmUED

When boundary-localized terms (BLT) come into play it is convenient to express four-component 5-dimensional fermion fields using two component chiral spinors¹ as [4]:

$$\Psi_L(x, y) = \begin{pmatrix} \phi_L(x, y) \\ \chi_L(x, y) \end{pmatrix} = \sum_{n=0}^{\infty} \begin{pmatrix} \phi_n(x) f_L^n(y) \\ \chi_n(x) g_L^n(y) \end{pmatrix}, \quad (1)$$

$$\Psi_R(x, y) = \begin{pmatrix} \phi_R(x, y) \\ \chi_R(x, y) \end{pmatrix} = \sum_{n=0}^{\infty} \begin{pmatrix} \phi_n(x) f_R^n(y) \\ \chi_n(x) g_R^n(y) \end{pmatrix}. \quad (2)$$

Above, the $SU(2)_L$ behaviour of the fields are suppressed. In mUED $f_i^n(y)$, $g_i^n(y)$, ($i = L, R$) are either a sine or a cosine function of y . For nmUED this will no longer be the case. Further, the mass of the KK-excitations will deviate from the simple n/R formula.

In UED the scalar and vector boson fields are obviously also 5-dimensional. Their KK expansions can be similarly written down.

In nmUED one may additionally consider kinetic and mass terms localized at the fixed points of the orbifold [5, 6, 7, 8, 9]. We restrict ourselves to boundary-localized kinetic terms (BLKT) only. Specifically, we examine the interaction of quarks and leptons with electroweak gauge bosons in a 5-dimensional theory with BLKT at $y = 0$ and $y = \pi R$.

We consider fermion fields $\Psi_{L,R}$ whose zero-modes are the chiral projections of the SM fermions. The five-dimensional free fermion action with BLKT is written as [4]:

$$S = \int d^4x dy \left[\bar{\Psi}_L i \Gamma^M \partial_M \Psi_L + r_f^a \delta(y) \phi_L^\dagger i \bar{\sigma}^\mu \partial_\mu \phi_L + r_f^b \delta(y - \pi R) \phi_L^\dagger i \bar{\sigma}^\mu \partial_\mu \phi_L + (\Psi_L, \phi_L, \bar{\sigma} \leftrightarrow \Psi_R, \chi_R, \sigma) \right] \quad (3)$$

with $\sigma^\mu \equiv (I, \vec{\sigma})$ and $\bar{\sigma}^\mu \equiv (I, -\vec{\sigma})$; $\vec{\sigma}$ being the Pauli matrices. r_f^a, r_f^b are the strengths of the boundary terms chosen to be the same for Ψ_L and Ψ_R . In the following only the Ψ_L, ϕ_L part of the lagrangian is considered. The results for Ψ_R, χ_R are similar.

Variation of the above action leads to

$$\left[1 + r_f^a \delta(y) + r_f^b \delta(y - \pi R) \right] m_n f_L^n - \partial_y g_L^n = 0, \\ m_n g_L^n + \partial_y f_L^n = 0, \quad (n = 0, 1, 2, \dots). \quad (4)$$

Eliminating g_L^n one obtains:

$$\partial_y^2 f_L^n + \left[1 + r_f^a \delta(y) + r_f^b \delta(y - \pi R) \right] m_n^2 f_L^n = 0. \quad (5)$$

The boundary conditions at $y = 0$ we impose are [6]:

$$f^n(y)|_{0^-} = f^n(y)|_{0^+}, \quad \left. \frac{df^n}{dy} \right|_{0^+} - \left. \frac{df^n}{dy} \right|_{0^-} = -r_f^a m_n^2 f^n(y)|_{0^+}, \quad (6)$$

and similar conditions at $y = \pi R$ with $r_f^a \rightarrow r_f^b$. Here we consider the alternative² $r_f^a \equiv r_f \neq 0, r_f^b = 0$. The solution for $0 \leq y < \pi R$ is:

$$f^n(y) = N_n \left[\cos(m_n y) - \frac{r_f^a m_n}{2} \sin(m_n y) \right], \quad (7)$$

where m_n for $n = 0, 1, \dots$ satisfy the transcendental equation [6]:

$$\tan(m_n \pi R) = -\frac{r_f^a m_n}{2}. \quad (8)$$

The solutions satisfy the *orthonormality* relations:

$$\int dy \left[1 + r_f^a \delta(y) \right] f^n(y) f^m(y) = \delta^{nm}. \quad (9)$$

The constant N_n is determined from orthonormality:

$$N_n = \sqrt{\frac{2}{\pi R}} \left[\frac{1}{\sqrt{1 + \frac{r_f^2 m_n^2}{4} + \frac{r_f}{2\pi R}}} \right]. \quad (10)$$

The reader is urged to note that (a) solutions in eq. (7) are combinations of sine and cosine functions rather than any one of them alone and that (b) the KK masses are solutions of eq. (8) not simply n/R . These features provide the novelty of nmUED.

In our work we only require the zero-modes and the $n = 1$ excitations of the fermion fields. We also deal with the five-dimensional SM Higgs scalar and the electroweak gauge bosons. For these boson fields one can undertake a similar KK decomposition. For the Higgs field we take $r_h^a \equiv r_h \neq 0, r_h^b = 0$ while for the gauge fields $r_G^a \equiv r_G \neq 0, r_G^b = 0$. The upshot is that the y -dependent part of the wave functions and the masses satisfy eqns. (7) and (8).

¹The Dirac gamma matrices are in the chiral representation with $\gamma_5 = \text{diag}(-I, I)$.

²For another alternative see [3].

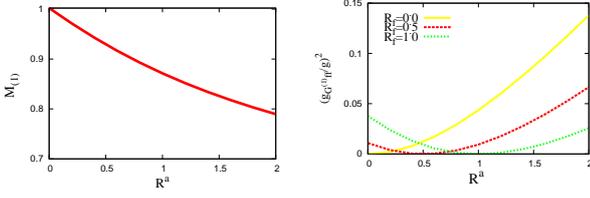


Figure 1: Left panel: Dependence of the mass of the $n = 1$ KK-mode $M_{(1)} \equiv m^{(1)}R$ on the BLKT strength $R^a \equiv r^a/R$ for any particle. Note that larger R^a yields a smaller mass. Right panel: The square of the KK-parity violating coupling (between $G^1 \equiv B^1$ or W_3^1 and a pair of fermions) with $R^a \equiv R_G$ for different choices of R_f .

The roots of equation (8), which may be obtained numerically, are the extra-dimensional contributions to the masses of the KK modes³. In the left panel of Fig. 1, we plot the dimensionless quantity $M_{(1)} \equiv m^{(1)}R$. It applies to all fields, namely, scalars, fermions, and gauge bosons. We show the variation of $M_{(1)}$ with⁴ $R^a \equiv r^a/R$. When $R^a = 0$, i.e., no BLKT at all, one gets $m^{(1)} = R^{-1}$, as expected. One finds that $m^{(1)}$ monotonically decreases as R^a increases asymptotically approaching $0.5R^{-1}$.

It is to be noted that the mass $M_{(1)}$ is determined entirely by the BLKT parameter, R^a , and the compactification radius R . The gauge coupling is not involved. Therefore this discussion applies for both W_3^1 and B^1 so long as the appropriate BLKT parameters are used.

We have checked that for the range of BLKT which we entertain the mixing between states of different KK level, n , is very small and may be ignored. Further, only if the BLKT parameters for the B and W gauge bosons are equal or very nearly equal the mixing between B^1 and W_3^1 is substantial, it being equal to the zero-mode weak mixing angle in the case of equality. If $(r_B - r_W)/R$ is as small as 0.1 this mixing is already negligible. This follows from the mass matrix as we outline below.

The mass matrix for the $n = 1$ neutral electroweak gauge bosons receives contributions from two sources: one originates from the spontaneous breaking of the electroweak symmetry and the other due to the extra-dimension discussed above. The contributions from symmetry breaking are $\mathcal{O}(v^2)$. The order of the extra-dimensional contributions, $(m_G^{(1)})^2$, is set by $(1/R)^2$ and is always much larger by far. Thus effectively these

terms determine the mass eigenvalues and the mixing is negligible for $(R_W - R_B) \sim 0.1$ or larger⁵. So, we take B^1 and W_3^1 to be the neutral electroweak gauge eigenstates for the $n = 1$ KK-level.

The five-dimensional gauge couplings g_5, g'_5 and the vacuum expectation value (vev) v_5 are related to the standard couplings g, g' respectively and the vev v defined in four dimensions by:

$$g_5 = g \sqrt{\pi R S_W}, \quad g'_5 = g' \sqrt{\pi R S_B}, \quad v_5 = v / \sqrt{\pi R S_H}, \quad (11)$$

where

$$S_W = \left(1 + \frac{R_W}{2\pi}\right), \quad S_B = \left(1 + \frac{R_B}{2\pi}\right), \quad S_H = \left(1 + \frac{R_h}{\pi}\right). \quad (12)$$

3. Coupling of B^1 and W_3^1 with zero-mode fermions

We have now all the ingredients needed to calculate the coupling of the states W_3^1 and B^1 to two zero-mode fermions $f^{(0)}$. Here $f^{(0)}$ could be SM quarks or leptons. We find that the results are not very sensitive to the exact value of R_h . Our chosen value ensures that the H^1 is always heavier than the B^1 and W_3^1 .

Below we discuss the case for a generic gauge boson G^1 , which could be either of B^1 or W_3^1 . The y -dependent wave-functions of our interest here are found to be

$$f_L^0 = g_R^0 = \frac{1}{\sqrt{\pi R(1 + R_f/2\pi)}}, \quad (13)$$

for fermions and for the $n = 1$ gauge boson

$$a^1 = \sqrt{\frac{1}{\pi R}} \sqrt{\frac{2}{1 + \left(\frac{R^a M_{(1)}}{2}\right)^2 + \frac{R^a}{2\pi}}} \left[\cos\left(\frac{M_{(1)}y}{R}\right) - \frac{R^a M_{(1)}}{2} \sin\left(\frac{M_{(1)}y}{R}\right) \right], \quad (14)$$

where $R^a \equiv r^a/R$. Using the above we find

$$g_{G^1 f^0 f^0} = \frac{\sqrt{2} g(G) \sqrt{S_G}}{\left(1 + \frac{R_f}{2\pi}\right) \sqrt{1 + \left(\frac{R^a M_{(1)}}{2}\right)^2 + \frac{R^a}{2\pi}}} \left(\frac{R_f - R^a}{2\pi}\right). \quad (15)$$

where we have used as earlier $M_{(1)} \equiv m_G^{(1)}R$.

In the right panel of Fig. 1 we plot the square of the above coupling strength as a function of R^a for several choices of R_f . A noteworthy feature is that the coupling vanishes when $R^a = R_f$.

³The KK excitations also receive a contribution to their masses from spontaneous breaking of the electroweak symmetry.

⁴Here $R_f = r_f/R$, $R_G = r_G/R$.

⁵If $R_W = R_B$ then the dominant diagonal terms are equal and do not contribute to the mixing and simply shift the masses of the eigenstates. In this case the mixing between W_3^1 and B^1 is just as in the Standard Model with $\tan \theta = g'/g$.

4. B^1 or W_3^1 production at the LHC and decay

We now turn to discuss a signal of nmUED at the LHC. We are interested in the resonant production of the $n = 1$ KK-modes of neutral EW gauge bosons, by the process $pp (q\bar{q}) \rightarrow G^1$ followed by $G^1 \rightarrow l^+l^-$ where G^1 is either of B^1 and W_3^1 and l^\pm could be either e^\pm or μ^\pm .

From here onwards for the SM particles we will not explicitly write the KK-number ($n = 0$) as a superscript. The final state leads to two leptons (e or μ), with invariant mass peaked at m_{G^1} . Note that the production as well as the decay of these $n = 1$ KK-excitations are driven by KK-parity violating couplings. If such a signature is found at the LHC, then it would be a good channel for the determination of such KK-parity violating couplings.

The production cross section in proton proton collisions can be written in a compact form :

$$\sigma(pp \rightarrow G^1 + X) = \frac{4\pi^2}{3m_{G^1}^3} \sum_i \Gamma(G^1 \rightarrow q_i\bar{q}_i) \int_\tau^1 \frac{dx}{x} \left[f_{q_i}^p(x, m_{G^1}^2) f_{\bar{q}_i}^p(\tau/x, m_{G^1}^2) + q_i \leftrightarrow \bar{q}_i \right] \quad (16)$$

Here, q_i and \bar{q}_i represent a generic quark and the corresponding antiquark of the i -th flavour respectively. $\Gamma(G^1 \rightarrow q_i\bar{q}_i)$ is the decay width of G^1 into a quark and antiquark pair of the i th flavour. $\tau \equiv m_{G^1}^2/S_{pp}$, where $\sqrt{S_{pp}}$ is the proton proton c.m. energy. f stands for the quark or antiquark distribution functions within a proton.

For B^1 production,

$$\Gamma = (g_{G^1 q\bar{q}}^2/32\pi) \left[(Y_L^q)^2 + (Y_R^q)^2 \right] m_{B^1} \quad (17)$$

(with Y_L^q and Y_R^q the weak-hypercharges for the left- and right-chiral quarks), while for producing W_3^1 one must use

$$\Gamma = (g_{G^1 q\bar{q}}^2/128\pi) m_{W_3^1}. \quad (18)$$

$g_{G^1 q\bar{q}}^{(\prime)}$ is the KK-parity violating coupling among SM quarks and W_3^1 (B^1) as given in eq. (15). In the above equations m_{G^1} stands for the mass eigenvalue of the gauge boson $n = 1$ excitation.

The numerical results for the cross sections are obtained using a parton-level Monte Carlo with parton distribution functions parametrized as in CTEQ6L [10]. We take the pp c.m. energy to be 8 TeV. The renormalisation (for α_s) and factorisation scales (in the parton distributions) are set at m_{G^1} .

To make a realistic estimate of the signal cross section, a simple calorimeter simulation has been implemented with:

- $p_T^\ell > 20$ GeV.
- The calorimeter rapidity coverage (for leptons) is $|\eta| < 3.0$.
- A cone algorithm with $\Delta R = \sqrt{\Delta\eta^2 + \Delta\phi^2} = 0.5$ has been used for lepton isolation.

Once produced, B^1 (W_3^1) will decay *via* KK-parity violating couplings to SM quarks and leptons and, if kinematically allowed, to $f^1\bar{f}$ (or to $f\bar{f}^1$) through KK-conserving couplings. The KK-parity violating leptonic decays provide a cleaner environment at the LHC. For simplicity we assume an universal coefficient r_f for the BLKTs involving all SM fermions. If kinematically allowed (broadly when $R_f > R_G$), the KK-conserving decay rates can be substantial and in such situations the branching ratios (BRs) for KK-violating decay rates are small. However, when the KK-conserving decays are kinematically disallowed, decays to a SM fermion anti-fermion pair are the only possible modes and hence the branching ratios become independent of the input BLKT strengths. Consequently, the decay rate of B^1 to any fermion species is proportional to the sum of the squares of the respective weak hypercharges $\left[(Y_L^f)^2 + (Y_R^f)^2 \right]$ of the left- and right-chiral components. W_3^1 , on the other hand, decays democratically with branching ratio of $\frac{1}{21}$ to each type of left-handed zero-mode fermions. This immediately implies that the decay branching ratio of B^1 (W_3^1) to e^+e^- or $\mu^+\mu^-$ is approximately $\frac{30}{103} \left(\frac{2}{21} \right)$.

Final states with dileptons arise in the SM mainly from resonant Z -production or Drell-Yan (DY). The first of these can be vetoed as in this case the dilepton invariant mass peaks around m_Z . We find that for 10 GeV bins around 700 (800) GeV the DY cross section is 2.29 (1.56) fb . This background though non-negligible is such that for the W_3^1 and B^1 masses which we consider $S/\sqrt{B} \geq 5$ can be reached for 20 fb^{-1} integrated luminosity.

We now present the results. In Fig. 2 we have plotted in the $R_B - R_W$ plane the iso-event curves⁶ (40 events with 20 fb^{-1} luminosity for LHC running at 8 TeV).

The essence of Fig. 2 can be understood from the earlier discussions. The KK-parity violating couplings

⁶The requirement here is that electron *plus* muon events together resulting from W_3^1 and B^1 production and decay add up to 40.

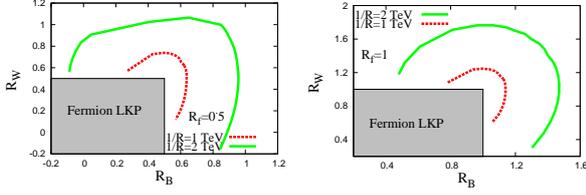


Figure 2: Iso-event curves (40 signal events with 20 fb^{-1} data at the LHC running at 8 TeV) for combined W_3^1 and B^1 signals in the $R_W - R_B$ plane. The region enclosed by the solid green (dashed red) curve and the fermion LKP box yields less than 40 events for $R^{-1} = 2 \text{ TeV}$ (1 TeV). The left (right) panel corresponds to $R_f = 0.5$ (1.0). Note difference in scales in the two panels.

vanish when $R_f = R_G$ where $G \equiv W$ or B as noted from eq. (15). This is seen in the right panel of Fig. 1 for different choices of R_f . So, the production mode we consider becomes unavailable. Further, in this situation, the $n = 1$ gauge boson and fermion states are mass degenerate and KK-number conserving decay modes are also not allowed. In the neighbourhood of this point there is an important asymmetry, between whether $R_{B,W}$ is more than R_f or less. In the former case, the gauge boson $n = 1$ state is lighter than the corresponding fermion state and KK-number conserving decays are not kinematically possible. If both R_B and R_W are smaller than R_f then the fermion state is the LKP.

For any iso-event contour the largest value of R_B will be for $R_W = R_f$ and *vice-versa*. This is because when $R_W = R_f$ there is no contribution to the signal from the W_3^1 as noted above. As R_W moves from this value the contribution from W_3^1 reduces the needed R_B . This behaviour is not symmetric because of the difference in the branching ratios of W_3^1 on the two sides of R_f .

In the figure we also indicate the region where the $n = 1$ fermion is the LKP. It is readily seen that for the point $R_W = R_B = R_f$ the KK-number violating coupling is zero and also the $n = 1$ states are degenerate. So, no decays, neither KK-number violating nor conserving, are permitted.

Data have been collected at LHC at 8 TeV proton-proton c.m. energy. Observation or otherwise of the proposed dilepton signal would help in exploring or excluding the parameter space of non-minimal UED models. The projected exclusion limits can be directly read off from the plots in Fig. 2. Any point above a particular iso-event curve can be excluded from the non-observation of such events. The r_f dependence of the cross section comes only through the coupling. A careful examination of the coupling in eqs. (15) shows that

it tends to a constant as $R_f \rightarrow \infty$.

5. Conclusions

To summarize, we have considered the effects of boundary localized kinetic terms in models where all SM fields can propagate in a spacetime with four spatial and one timelike dimensions. The extra spatial dimension y is compact and can be considered as a circle of radius R with a $y \leftrightarrow -y$ symmetry. This results in two fixed points at $y = 0$ and $y = \pi R$. At these points one can include terms consistent with 4-dimensional Lorentz symmetry. These are either kinetic or mass terms. We examine the former.

In the minimal Universal Extra Dimensional model, radiative corrections play an important role in removing the near-degeneracy of the masses of the KK-modes of all SM particles of the same KK-level, n . UED, being an effective theory, is valid only up to a cut-off, Λ . In mUED the boundary terms are fixed in a manner such that at Λ the contribution due to radiative corrections is exactly compensated. In fact, instead of calculating the radiative correction in a 5d set up one may also parametrize these effects by incorporating a set of BLKTs. These are somewhat similar in spirit to the high-scale universal mass parameters m_0 and $m_{\frac{1}{2}}$ often introduced in SUSY which serve as boundary conditions for the estimation of the low-energy masses.

There are two possibilities of choosing the BLKTs with rather different physics consequences. In the first, the BLKTs are of equal strength at the two boundary points ($y = 0, \pi R$). Here, a Z_2 symmetry $y \leftrightarrow (y - \pi R)$ remains. This results in a theory where the spectrum of KK-particles and the couplings are drastically different from mUED. The lightest among the $n = 1$ KK particles can be a dark matter candidate⁷. The other alternative is to permit the BLKTs at $y = 0$ and $y = \pi R$ to be of unequal strengths. This leads to a breakdown of KK-parity and will allow, for example, the decay we have examined, $B^1(W_3^1) \rightarrow e^+ e^-, \mu^+ \mu^-$, and production of the $B^1(W_3^1)$ singly.

In this presentation, we have examined the possible BLKTs for an interacting theory of fermions and the neutral electroweak gauge bosons. We have considered the situation where the BLKTs are present *only* at the $y = 0$ fixed point and they are vanishing at $y = \pi R$. The boundary terms modify the field equations for all particles in the y -direction. Consistency conditions of

⁷An analysis of Dark Matter in the context of these nmUED models is carried out in [11]

the solutions of these equations lead to masses of the higher KK-modes of fermions and the photon.

For purpose of illustration, we have calculated the coupling of W_3^1 and B^1 , the $n = 1$ KK-excitations of the neutral electroweak gauge bosons, to a pair of zero-mode fermions (i.e., SM fermions) as a function of r_f, r_G , and R^{-1} . The production and decay of W_3^1 and B^1 at the LHC, *via* these KK-parity violating couplings, have been considered. We have explored the viability of the *dilepton* signature at the LHC running at 8 TeV pp center of mass energy. It is revealed that non-observation of such a high mass *dilepton* signal with 20 fb⁻¹ integrated luminosity in the 8 TeV run of LHC will disfavour a large part of the parameter space (spanned by r_f, r_G , and R^{-1}).

The same signal can also arise if there are extra Z-like bosons as in several popular extensions of the SM, e.g., the Left-Right symmetric models or models with an extra $U(1)$ symmetry. We have not compared the predictions of the model under consideration with those in these other scenarios.

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