

LHC signals for singlet neutrinos from a natural warped seesaw mechanism. II

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A natural seesaw mechanism for obtaining the observed size of SM neutrino masses can arise in a warped extra-dimensional/composite Higgs framework. In a previous paper, we initiated the study of signals at the LHC for the associated \sim TeV mass SM singlet neutrinos, within a canonical model of $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ (LR) symmetry in the composite sector, as motivated by consistency with the EW precision tests. Here, we investigate LHC signals in a different region of parameter space for the same model, where production of singlet neutrinos can occur from particles beyond those in the usual LR models. Specifically, we assume that the composite $(B-L)$ gauge boson is lighter than all the others in the EW sector. We show that the composite $(B-L)$ gauge boson can acquire a significant coupling to light quarks simply via mixing with elementary hypercharge gauge boson. Thus, the singlet neutrino can be pair-produced via decays of the $(B-L)$ gauge boson, without a charged current counterpart. Furthermore, there is no decay for the $(B-L)$ gauge boson directly into dibosons, unlike for the usual case of W_R^\pm and Z' . Independently of the above extension of the EW sector, we analyze production of singlet neutrinos in decays of composite partners of $SU(2)_L$ doublet leptons, which are absent in the usual LR models. In turn, these doublet leptons can be produced in composite W_L decays. We show that the $4-5\sigma$ signal can be achieved for both cases described above for the following spectrum with 3000 fb^{-1} luminosity: 2–2.5 TeV composite gauge bosons, 1 TeV composite doublet lepton (for the second case) and 500–750 GeV singlet neutrino.

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I. INTRODUCTION

Numerous versions of the seesaw mechanism for explaining the extreme smallness of standard model (SM) neutrino masses have been proposed over the last few decades. In a previous paper [1], we discussed how a natural avatar arises with the SM fields propagating in a warped extra dimension. By the AdS/CFT correspondence, this framework is dual to the SM Higgs boson being composite of some new strong dynamics, while rest of the SM fields are partially composite, i.e., admixtures of composites and elementary particles external to the strong dynamics. In fact, this realization of the seesaw paradigm combines two specific models therein. Namely, it is

basically an inverse seesaw [2]¹; i.e., a SM neutrino mass is generated by exchange of weak scale pseudo-Dirac singlet neutrinos, but with the tiny Majorana mass term for the relevant component of the singlet being the result of a type I high-scale seesaw [4]. Finally, the effective seesaw scale being several orders of magnitude below Planck scale can also be accommodated naturally in this scenario: for more explanation and more references, see Ref. [1].

All in all, a study of the signals at the Large Hadron Collider (LHC) from production of these TeV-mass singlets, which play a crucial role in generating SM neutrino masses, is then highly motivated. We would like to emphasize here that earlier analyses [5] of the same framework did not use the mass basis for these singlet modes, which resulted in a suggestion that it is a (purely) high-scale seesaw instead. With this incorrect impression, even though a KK tower of singlet particles starting at \sim TeV is still present, the focus would be instead on a super-heavy singlet mode, i.e., seemingly beyond direct reach of

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¹Linear seesaw [3] has some features which are similar to inverse seesaw.

current and future experiments. In this sense, as a consequence of the realization in Ref. [1] that it is physically an inverse seesaw, the status of the TeV-mass singlets changed from being mere “vestiges” of SM neutrino mass generation to central players therein.

In a very recent paper [6], we took the first step in this direction. There, we focused on a specific five-dimensional model with a $SU(2)_L \times SU(2)_R \times U(1)_X$ bulk electroweak (EW) gauge symmetry, where $SU(2)_R \times U(1)_X$ is broken down to $U(1)_Y$ on the Planck brane [7]. This choice is dual to the composite sector respecting a $SU(2)_L \times SU(2)_R \times U(1)_X$ global symmetry, whereas only its $SU(2)_L \times U(1)_Y$ subgroup is gauged by the elementary sector. The purpose of this extension of the EW symmetry to a left-right symmetric (LR) structure was to ameliorate constraints from EW precision tests. We analyzed the production of the composite singlet neutrinos via decays of on-shell composite gauge bosons, namely, the W_R^\pm and Z' . We assumed there that these neutrinos, along with the $SU(2)_L$ singlet composite charged leptons, are doublets of $SU(2)_R$, denoted by $(N^{(1)}, \tilde{\ell}^{(1)})$.²

In addition, in the above-mentioned work, we made two assumptions mostly for simplicity. The first one was that (i) all the EW spin-1 composites are approximately degenerate, with the composite sector taken in isolation, i.e., neglecting the small mixing with the elementary sector, and considering EW symmetry breaking (EWSB) effects also as a perturbation.

As we will explain below, the role of the above choice was crucial in inducing a significant coupling of light quarks to W_R^\pm and Z' , as required for their production at the LHC. The net result is that for 2 TeV composite gauge boson mass, 750 GeV singlet neutrinos and with 300 fb^{-1} luminosity, we can have discovery of singlet neutrinos via decay of W_R^\pm in this case [6].

The second choice we implicitly made was that

- (i) $SU(2)_L$ doublet composite leptons, which are mandatory in this framework, are heavier than one-half of the composite gauge boson; i.e., only singlets are in the game.

In summary, we see that the above signals are roughly similar to usual LR models,³ even though quantitative details are different, that too significantly. For example, note that W_R^\pm decay produces composite singlet charged lepton (in association with singlet neutrino), which, in turn, decays into SM charged lepton and Higgs (including longitudinal W/Z). That is, we get an extra identifiable final state particle as compared to usual LR model, where W_R^\pm directly decays into SM charged lepton (and singlet neutrino).

²We use this notation for these composites since they correspond to KK modes of the five-dimensional model: the qualifier “tilde” on ℓ will be explained later.

³For a review, see [8].

In the present paper, which is to be considered as the second installment of this series, we study LHC signals of singlet neutrinos in the same setup as above, but now moving on to different region of its parameter space. As we will see, even though this step looks simply like a “quantitative” change, we show that it will lead to variations in the qualitative features of the signals, in particular, involving particles for which there is no counterpart in the usual LR models. In this sense, the search channels will be even more different from the usual LR models than those discussed in the earlier paper.

In the first part of this paper,

- (i) we will relax the earlier assumption of near-degeneracy of the composite spin-1 states. In particular, we keep some at 2 TeV so that we get sufficient production rate at the LHC, but raise the others above this value.

Indeed, the spirit here is simply to explore other options as compared to the previous paper. However, at first sight, this seems like a drastic step to take, namely, this could reduce the coupling of light quarks to W_R^\pm , thus seems to render negligible the W_R^\pm signal (similarly for the Z'). Hence, naively we might then have to look for alternate avenues for production of singlet neutrino in this case, recalling that the singlet neutrino couples only to W_R^\pm and Z' among the spin-1 composites.⁴ Indeed, this is what actually happens for W_R^\pm .

- (i) Remarkably, we discover that in the process of making composite $(B-L)$ boson lighter⁵ than composite W_R^3 (or vice versa), a “new” neutral current channel can emerge. Namely, the lighter of these two composite gauge boson potentially still has significant production rate, with out requiring EWSB.

This “twist” arises in a subtle manner as follows. For this purpose, it is worthwhile recapping what are the various couplings of composite spin-1 particles to fermions. First of all, the matter particles in a given sector, either elementary or composite, only couple to the corresponding gauge bosons. However, there is mass mixing between these two sectors, both for fermions and gauge bosons, resulting in modifications of these couplings as well. The light SM quarks are mostly elementary so that the coupling of these quarks to the heavy gauge bosons induced by the elementary-composite fermionic mixing is negligible. However, a sizable (even if mildly suppressed as compared to SM) such coupling can result from mixing of elementary and

⁴On the other hand, the composite charged lepton $\tilde{\ell}^{(1)}$ couples also to composite hypercharge so that only external-composite mixing suffices for its production via light quark initial state; i.e., the EWSB induced mixing, which is suppressed for this case of nondegeneracy is not needed here. However, $\tilde{\ell}^{(1)}$ is not directly related to the mechanism of generation of SM neutrino mass.

⁵A factor of ~ 1.5 is enough here.

composite gauge bosons⁶: straightforwardly, we have elementary W_L (charged and neutral) mixing with their composite counterparts. Obviously, this opens the door for production of these composite particles at the LHC. Needless to say, the composite singlet neutrino does not couple to these composite W_L 's; on the other hand, it does have a coupling to composite W_R^\pm . However, there is no mixing effect analogous to W_L in the W_R^\pm sector, where there is only the composite side, thus making W_R^\pm (and, in turn, the singlet neutrino) inaccessible to the LHC at this level. Finally, there remain the neutral gauge bosons W_R^3 and $(B - L)$, where the situation is rather subtle as follows. The elementary hypercharge (denoted by B henceforth) gauge boson mixes with a specific combination of the composite W_R^3 and $(B - L)$. The point is that, in the approximation of composite W_R^3 and $(B - L)$ being degenerate, this “superposition” is a mass eigenstate, dubbed “composite” hypercharge. Thus, the composite B does couple to light quarks via this mixing, just like for the case of W_L 's discussed above. Whereas, the orthogonal combination of composite W_R^3 and $(B - L)$, usually denoted as Z' , does not mix with elementary hypercharge, thus being decoupled from light quarks at this order. Recall that the composite SM singlet neutrino couples to both the composite W_R^3 and $(B - L)$ gauge bosons, in such a manner that in the degenerate case, there is then only a coupling of singlet neutrino to Z' , but not to composite hypercharge. The upshot here then seems to be that, including both neutral and charged channels, there is no coupling of singlet neutrino to light quarks for its production at the LHC via spin-1 intermediaries at this stage.

Turning on the Higgs VEV mixes the various composites amongst themselves, in both charged and neutral sectors. Combined with elementary-composite mixing, this effect of EWSB then does induce a coupling of light quarks to Z' : explicitly, this proceeds via Z' mixing with composite B and W_L^3 , followed by elementary-composite B , W_L^3 mixing. A similar argument applies to W_R^\pm , now involving mixing with various W_L^\pm 's. Hence, singlet neutrino production at the LHC can then take place in both neutral and charged channels. However, naively this coupling still seems to be suppressed, since the Higgs VEV is somewhat smaller than the compositeness scale. Remarkably, the same composite degeneracy comes to the rescue here, since it can result in a large mixing angle between Z' and composite B and W_L^3 , even with a smaller mass mixing term. Once again, a parallel consideration holds for $W_R^\pm - W_L^\pm$ mixing. Thus, we finally get a non-negligible coupling of Z' and W_R^\pm to light quarks inside the proton. Furthermore, the above

argument suggests that a significant nondegeneracy of spin-1 composites might reduce these signals.

However, when the spin-1 composites are not degenerate, we have to go back to their “original” identities, chosen as per the symmetries of the strong dynamics in order to be in the physical, mass basis (this point was also noted in the third reference in [9]). Namely, we have a degenerate triplet of composite W_L 's, another one for W_R and finally a neutral composite $(B - L)$, with three different masses in general. Of course, compared to the above one, there is *no* change in the basis for composite W_L 's and W_R^\pm here, but there is a crucial difference for composite W_R^3 , $(B - L)$ as follows. As before, the external B mixes with both composite W_R^3 and $(B - L)$, with $N^{(1)}$ coupling to both of them. However, the composite W_R^3 and $(B - L)$ now have different masses so that this combination of composites is not even close to a mass eigenstate. This situation is to be contrasted with the degenerate case, where this admixture is a mass eigenstate, i.e., the composite B , and in fact the singlet neutrino decouples from it, due to a “cancellation” between couplings to the constituent W_R^3 and $(B - L)$. We can then contemplate two cases, i.e., composite $(B - L)$ is lighter or heavier than composite W_R^3 , say, ~ 2 vs ~ 3 TeV.

So, as anticipated earlier, we can have non-negligible production of composite $(B - L)$ (assuming that is lighter) simply via its mixing with external hypercharge gauge boson which, in turn, couples to light quarks inside proton. In this way, there is no need to involve EWSB for generating this coupling, cf. for composite Z' in the degenerate case. Clearly, production of the heavier composite W_R^3 (or composite $(B - L)$ in the other case) via a similar mechanism can then be neglected in comparison: coupling of this heavier state to light quarks is similar to that of the lighter one so that suppression in the rate is simply due to the masses. Furthermore, composite $(B - L)$ (or W_R^3 in the other case) will decay into pair of $N^{(1)}$ with a sizable branching ratio.

Even though the final state looks similar to production of singlet neutrinos from decay of Z' in usual LR models, the composite $(B - L)$ (or W_R^3) production involves different couplings/branching ratios etc. and thus the two should be distinguishable. In fact, in this regard

- (i) the composite $(B - L)$ production is extremely interesting, since it does not decay into Higgs/ W/Z_{long} : only channels are SM fermions,⁷ the singlet neutrino and composite RH charged lepton [assuming other composite fermions are heavier than one-half the mass of composite $(B - L)$].

So, the singlet production could serve not only as a test of the seesaw mechanism, but also as a first signal for this EW

⁶The amount of elementariness of the heavy gauge bosons, even if small, is much larger than the degree of compositeness of light quarks.

⁷In this case, it is dominated by SM top simply because other SM fermions are mostly elementary and hence couple weakly to composite $(B - L)$.

composite boson. Recall that usually this role is played by the dibosons, i.e., Higgs/ W/Z_{long} instead (including in the degenerate case studied earlier).

In addition, suppose we also make composite W_L 's heavier than composite W_R 's or vice versa. In this case, production of W_R^\pm from light quarks might become negligible, since this coupling requires EWSB mixing, whose effect is damped by the nondegeneracy. Note that elementary W_L^\pm mixes only with its composite counterpart, cf. neutral case, where elementary B mixes with both composite W_R^3 and $(B-L)$.⁸ In other words, there is no analog of above neutral channel effect in the charged case. So,

- (i) another striking feature of this case is singlet neutrino production via neutral gauge boson—either composite $(B-L)$ or W_R^3 —with out being accompanied by similar contribution from the *charged*, W_R^\pm , channel, cf. usual LR models, where W_R^\pm signal is typically larger than Z' due to the former's smaller mass.

For simplicity and for clearly illustrating the above two distinctive signatures, we will focus our analysis on the case of composite $(B-L)$ gauge boson being light (say, ~ 2 TeV), with *both* the W_R and W_L (charged and neutral) being heavier ($\gtrsim 4$ TeV) so that latter's production at the LHC is negligible. Hence, we will only observe a neutral spin-1 heavy particle in the EW sector, which also does not decay into dibosons, but with sizable singlet neutrino production from it.

Independent of above nondegeneracy of spin-1 composites,

- (i) we consider the possibility of composite singlet neutrino signal from production and decay of composite lepton $SU(2)_L$ doublet, denoted by $L_L^{(1)}$. Note that $L_L^{(1)}$ couples both to composite W_L , via a gauge coupling, and to $N^{(1)}$, $\tilde{\nu}^{(1)}$ and the Higgs doublet, via the coupling that is a central ingredient for seesaw mechanism. For example, we could choose a spectrum where $L_L^{(1)}$ is at

⁸We lose $Zb\bar{b}$ custodial symmetry in this case, since that requires invariance under $L \leftrightarrow R$ exchange; i.e., $(t, b)_L$ is bi-doublet of $SU(2)_L \times SU(2)_R$ and degenerate composite W_L and W_R [10]. However, we can still ensure that shift in $Zb\bar{b}$ is small as follows. In the case of composite $(B-L)$ being light, say, ~ 2 TeV, it suffices to assume that *both* $W_{L,R}^3$ are heavy, say, ~ 5 TeV, irrespective of representations of b_L under $SU(2)_R$, since these are the only states which couple to Higgs VEV and thus can cause $Zb\bar{b}$ shift. Whereas in the case of ~ 2 TeV composite W_R^3 (with composite $(B-L)$ at $\gtrsim 3$ TeV), we might have to revert to canonical representations instead; i.e., $(t, b)_L$ is the singlet of $SU(2)_R$ so that only composite W_L^3 couples to both Higgs and b_L causing shift in $Zb\bar{b}$, but which can be heavy enough (again, ~ 5 TeV) in order to make this shift small. In other words, light (but still $\gtrsim 2$ TeV) composite W_R^3 or $(B-L)$ can be consistent with $Zb\bar{b}$ shift, since these particles can decouple either from Higgs or b_L .

~ 1 TeV so that it can be pair-produced at the LHC in decays of on-shell ~ 2.5 TeV composite W_L (either charged or neutral), which couple to light quarks using elementary-composite W_L mixing. And, we take $N^{(1)}$, $\tilde{\nu}^{(1)}$ mass to be a bit smaller than $L_L^{(1)}$, say, 500 GeV so that the dominant decays of $L_L^{(1)}$ will be to $N^{(1)}$, $\tilde{\nu}^{(1)}$ and Higgs/ W/Z_{long} . It may be worthwhile to emphasize that

- (i) such composite lepton doublets (i.e., possibly heavier than the RH neutrino which induces the seesaw)—and thus the associated signals mentioned above—are absent in usual/four-dimensional LR models,⁹ whereas their presence is required in the composite (or five-dimensional) seesaw being studied here.

Also, we would like to reiterate that we have not at all changed the model compared to the one used in our previous paper. The composite lepton doublet was still present even earlier, but was simply assumed to be heavy.

Finally, we emphasize the following model-independence of the above signal. Suppose the production of singlet neutrino via direct decays of spin-1 composites, e.g. $W_R^{3,\pm}$ or $(B-L)$, is suppressed, for example, due to latter being heavy, say, $\gtrsim 3$ TeV. An even more extreme case is the “absence” of the singlet neutrino gauge couplings altogether, either because the RH neutrino is singlet even under extended EW symmetry of the strong dynamics or we do not even have such an extension in the first place. Even in these cases, we will still have available the above-mentioned avenue of singlet neutrino production via decay of doublet composite lepton, in turn, originating from decays of composite W_L . The point is that this entire reaction, production of W_L^\pm followed by decay into composite lepton doublet and decay of composite lepton doublet into singlet neutrino, proceeds via couplings (gauge and Yukawa, respectively) which are always present and sizable. This feature is to be contrasted with the singlet neutrino signal via $W_R^{3,\pm}$ or $(B-L)$. In the latter case, the size of couplings in production of spin-1 states depends on the amount of degeneracy when EWSB effects are important, whereas their decay into singlet neutrino is dictated by the choice of the representation of the singlet neutrino under the extended gauge symmetry. Of course if composite W_L is also heavy, then we can resort to SM W_L exchange for production of composite lepton doublet, followed by its decay into singlet neutrino as above, although in this case, the cross section will be smaller since it will be nonresonant.¹⁰

⁹Similarly, while the four-dimensional type-III seesaw features EW-charged extra fermions, these are degenerate with the RH neutrino giving the seesaw.

¹⁰An even more model-independent production of singlet neutrino, involving neither heavy spin-1 nor other spin-1/2 states, is via its coupling to SM W and LH electron. This coupling arises from mixing of SM doublet and singlet neutrinos induced by the Higgs VEV, i.e. using the same Yukawa coupling which gives the SM neutrino Dirac mass term in the seesaw [11].

Here is the outline of the rest of this paper. We begin in Sec. II with a brief review of the basic seesaw model in the warped extra dimensional framework, emphasizing the region of parameter space which is new compared to the one studied in our earlier paper. In Sec. III, we present details of the “simplified,” two-site approach [9] to studying the five-dimensional model. We will employ this two-site model in our actual analysis of LHC signals. We then discuss our main results, starting with production mechanism and decay widths of various heavy particles in Sec. IV, followed by analyses of SM backgrounds and thus the discovery potential for the new particles in Sec. V. We conclude and present some directions for future work in Sec. VI. Appendixes contain the more technical details of the mixing between elementary hypercharge and composite W_R^3 , $(B - L)$ gauge bosons.

II. REVIEW OF THE FIVE-DIMENSIONAL MODEL

We study a model implementing the seesaw mechanism for SM neutrino mass in the context of SM fields propagating in a warped extra dimension. As mentioned in the introduction, the present paper is part of a series on this topic. We will continue using the same basic model with left-right symmetry structure in the EW sector as our previous paper on LHC signals [6]. So in this section, we just give a brief review, simply referring to [6] for further details and more references. The main purpose is to discuss new parts of parameter space or features which were not elaborated upon in the earlier paper. Also, the reader is just referred to [1,5] for the basic warped seesaw model without extended EW gauge symmetry.

The bulk EW gauge symmetry is $SU(2)_L \times SU(2)_R \times U(1)_X$. The $SU(2)_L \times SU(2)_R$ subgroup is broken by the Higgs VEV on the IR brane down to $SU(2)_V$ custodial symmetry, while $U(1)_X$ is unbroken here. On the UV brane, $SU(2)_R \times U(1)_X$ broken down by boundary conditions to $U(1)_Y$, where $Y = T_{3R} + X$. The bulk fermions are taken to be similar to before. In particular, the SM $SU(2)_L$ doublet lepton is a singlet of $SU(2)_R$, whereas the SM right-handed [$SU(2)_L$ singlet] charged lepton is embedded in a doublet of $SU(2)_R$.

In this framework, the new states beyond SM are the Kaluza-Klein (KK) excitations of the SM particles, as well as singlet neutrino modes, all with mass at the \sim TeV scale. In this and the previous paper, we study the LHC signals from production and decay of these singlet neutrinos arising from the decay of other heavier KK particles, in particular, belonging to the extended EW sector alluded to above. In the current paper, we actually perform two signal analyses. In both of cases, as for the most general possibility, we assume presence of non-negligible brane-localized kinetic terms (BKTs) for bulk gauge fields, which change their masses and

couplings [12].¹¹ In particular, this will result in essentially three independent KK masses for the $U(1)_X$, $SU(2)_L$ and $SU(2)_R$ gauge bosons. In this work, all we need is these masses differing by $O(1)$ factors for which $O(1)$ BKTs suffices (i.e., no larger hierarchy is called for here). This is to be contrasted with our earlier paper where BKTs for gauge fields were implicitly neglected (or assumed same for all) so that these KK fields would instead be approximately degenerate up to boundary conditions on UV brane and EWSB effects. Note that in the earlier paper, BKTs were invoked for the singlet neutrino¹² in order to make it lighter than the EW KK modes such that former can be produced in decays of the latter: we continue to do so in this paper also.

We will work out details in two-site model in the next section, but here we would like just to summarize the impact of the above nondegeneracy. Firstly, the KK W_R^\pm -KK W_L^\pm mixing due to Higgs VEV is now suppressed as $\propto v^2/(\text{KK mass splitting})^2$, with KK mass splitting of order KK mass itself. A more subtle effect happens in the extended neutral gauge sector which couples to singlet neutrino. Namely, without degeneracy in their masses, reorganizing KK modes of W_R^3 and X into KK hypercharge and Z' is no longer valid. Instead, KK modes of W_R^3 and KK X are approximately mass eigenstates separately. Although they mix on UV brane, where $U(1)_R \times U(1)_X$ is broken down to $U(1)_Y$, this UV-brane localized mixing effect can be treated as a perturbation given that the KK mode profiles are peaked near the IR brane. Note that the light quarks are effectively localized on the UV brane. In this way, both approximated mass eigenstates, mostly made of KK X and KK W_R^3 , couple to light quarks (\propto hypercharge coupling), thus can be produced at the LHC, and decay into pair of singlet neutrinos along with other channels. This scenario is the focus of the first part of this paper, where we assume the canonical choice $X = \frac{1}{2}(B - L)$. Remarkably, KK X does not decay to dibosons, cf., typical EW KK gauge bosons.

In the second study of LHC signals in this paper, the main new feature compared to before is to assume KK excitations of SM $SU(2)_L$ doublet are light due to BKTs. In particular, the mass of KK doublet lepton is taken to be $\lesssim \frac{1}{2}$ of KK W_L so that these can be pair-produced in KK W_L decays. These $SU(2)_L$ KK doublet leptons can then decay into the SM singlet neutrinos, assuming it is kinematically allowed. Note that these states always exist in the five-dimensional models, but are simply assumed to be heavy in previous paper.

Having pointed out new features of the five-dimensional models we consider in this paper, we now switch to two-site model for an actual analysis.

¹¹Another possibility for such modifications of properties of gauge KK is to assume that the various gauge fields propagate in different bulk regions [13].

¹²For a discussion of fermion BKTs in general, see Ref. [14].

III. TWO-SITE APPROACH TO THE FIVE-DIMENSIONAL MODEL

The two-site (or sector) model is an economical, effective description of the above five-dimensional model, motivated by deconstruction and AdS/CFT correspondence. It is roughly equivalent to keeping only first KK and zero modes of the five-dimensional model. We start with a brief review of the two-site model used in our previous paper [6]. The composite sector has $SU(2)_L \times SU(2)_R \times U(1)_X$ global symmetry so that there are massive spin-1 composites in the adjoint representation of this group. Only the SM subgroup of this symmetry, $SU(2)_L \times U(1)_Y$, is gauged. Namely, in elementary sector, $W_L^{3,\pm}$ and hypercharge B gauge bosons mix with the appropriate spin-1 composites. Before EWSB effects, which we discuss in the next subsection, the SM EW gauge fields are the resulting massless eigenstates, which have a small admixture of the composites. Similarly, the heavy spin-1 fields are mostly the composites, but also have a bit elementary components.

Moving on to the fermionic sector, we have massive Dirac composites in suitable representations of the global symmetry, having couplings to the relevant composite gauge bosons. Elementary *chiral* fermions mix with these composites, producing the SM fermions (before EWSB) as the massless eigenstates after diagonalization, plus the heavy fermions, which are mostly composite, but have a small admixture of the elementary fermions.

Next, we will elaborate further on the sub-sector whose LHC signals will be analyzed, again highlighting the difference from before. As mentioned above, we work out two different cases in this regard so that it is convenient to spell out the Lagrangians separately.

(a) Composite $(B-L)$ gauge boson

Here, we assume the canonical LR structure in the lepton sector as in [6], i.e., identify X with $\frac{1}{2}(B-L)$, taking the two sets of composite leptons to be $(2, 1)_{-1/2}$ and $(1, 2)_{-1/2}$ under $SU(2)_L \times SU(2)_R \times U(1)_X$, respectively. The modification from [6] is allowing for nondegeneracy of EW spin-1 composites, $SU(2)_L$, $SU(2)_R$, and $U(1)_X$. Such nondegeneracy is “inspired” by the effect of BKT’s in the five-dimensional model mentioned above.

One consequence of this departure from our earlier work takes place after EWSB which we will briefly consider in the next section. Here, we focus on an

interesting effect which takes place even before EWSB in the spin-1 sector coupled to the SM singlet neutrino, i.e., W_R^3 and $(B-L)$ gauge bosons. The details of this process are given in the Appendix, from which we extract the bottom lines as follows.

The elementary hypercharge mixes with a combination of composite $(B-L)$ and W_R^3 as before,¹³ but the crucial point is that these two composites are not degenerate anymore. The two heavy mass eigenstates resulting from this procedure thus contain mostly composite $(B-L)$ and W_R^3 , respectively, with small admixture of elementary hypercharge, followed by even smaller bit of the other composite. Such small admixture of elementary hypercharge, however, turns out to be large enough to secure significant production of heavy gauge bosons at the LHC. These are then dubbed simply (with slight abuse of notation) as “heavy $(B-L)$ ” and “heavy W_R^3 .” Note that light quarks essentially couple only to the elementary hypercharge. As a result, both the heavy spin-1 fields acquire a coupling to light quarks, roughly at similar level as in the W_L sector, i.e., $\sim g^2/g_*$, where g and g_* are the appropriate SM and composite gauge couplings, respectively. And, both heavy $(B-L)$ and W_R^3 can decay into singlet neutrinos.

The same analysis in the degenerate limit shows that one heavy mass eigenstate decouples from light quarks, but not from the singlet neutrino, and vice versa for the other. This is in agreement with [6]. Former would be combination of composite W_R^3 and $(B-L)$ which does not mix with elementary hypercharge (called Z'), whereas latter is simply the “heavy hypercharge boson” (corresponding to KK hypercharge of the five-dimensional model).

For simplicity and because the signal is then more dramatic, we assume composite W_R^3 (and W_R^\pm) is heavier than composite $(B-L)$, thus keeping only the heavy $(B-L)$. Similarly, composite W_L ’s are also taken to be heavier than composite $(B-L)$ and therefore neglected. Finally, the composite $SU(2)_L$ doublet is made heavier than one-half the mass of heavy $(B-L)$, whereas composite SM singlet neutrino lighter so that it can be pair-produced in decays of heavy $(B-L)$.

Hence, before EWSB, the Lagrangian in the mass eigenstate basis for the fermion-gauge sector relevant for our first analysis is given by

$$\begin{aligned} \mathcal{L}_{\text{gauge-fermion}} = & -\frac{1}{4} [F_{\mu\nu}^{(0)2} + \rho_{B-L}^{\mu\nu 2}] + \frac{1}{2} m_*^2 \rho_{B-L\mu}^2 + \overline{\psi}^{(0)} i \not{D} \psi^{(0)} + \overline{\tilde{L}}_R^{(1)} (i \not{D} - m_R) \tilde{L}_R^{(1)} - Q_Y \frac{g_Y^2}{g_*} \overline{\psi}_{\text{light}}^{(0)} \rho_{B-L}^\mu \gamma_\mu \psi_{\text{light}}^{(0)} \\ & + \frac{1}{6} \left(-\frac{g_Y^2}{g_*} \cos^2 \phi_{Q_L^3} + g_* \sin^2 \phi_{Q_L^3} \right) \overline{Q}_L^{(0)3} \rho_{B-L}^\mu \gamma_\mu Q_L^{(0)3} + \frac{1}{6} g_* \overline{t}_R^{(0)} \rho_{B-L}^\mu \gamma_\mu t_R^{(0)} - \frac{1}{2} g_* \overline{\tilde{L}}_R^{(1)} \rho_{B-L}^\mu \gamma_\mu \tilde{L}_R^{(1)}. \quad (1) \end{aligned}$$

¹³This matches mixing between these KK’s on UV brane in the five-dimensional model, which is not an EWSB effect.

An explanation of the notation is in order here. The massless (before EWSB) SM fields, including gauge bosons and fermions, are denoted by superscript “(0)”, since they correspond to zero modes of the five-dimensional model. In particular, $\psi_{\text{light}}^{(0)}$ stands for $(\psi^{(0)} - \{Q_L^{(0)3}, t_R^{(0)}\})$, which are assumed to be mostly elementary. We have assumed $t_R^{(0)}$ is fully composite, whereas $\phi_{Q_L^3}$ is the elementary-composite mixing angle for $Q_L^{(0)3}$. The only heavy fermion is $\tilde{L}_R^{(1)}$, the vectorlike/heavy $SU(2)_R$ doublet fermion which contains the SM singlet neutrino relevant for the SM neutrino mass seesaw (along with a charged leptonlike partner). The superscript “(1)” again is a reminder that this would be the KK mode of the corresponding five-dimensional field, while the “tilde” notation will be explained a bit later. Moving onto spin-1 sector, ρ_{B-L} stands for the heavy $(B-L)$ gauge boson (roughly the KK $(B-L)$ of the five-dimensional model), with g_\star being the associated composite gauge coupling. The elementary hypercharge-composite $(B-L)$ mixing angle is $\approx g_Y/g_\star$ (see Appendix for details), where g_Y denotes SM hypercharge coupling. Since $g_Y/g_\star \ll 1$, we simply set cosine of this mixing angle to be 1 for a good approximation. Finally, D_μ stands for the covariant derivative with respect to the SM gauge group; i.e., $D_\mu = \partial_\mu - igA_\mu^{(0)}$, where $A_\mu^{(0)}$ is in matrix form with appropriate generators.

(b) Composite $SU(2)_L$ doublet lepton

In the second study, we focus on the composite, heavy partner of the $SU(2)_L$ doublet lepton, which is coupled to composite W_L gauge boson. We assume mass of former is less than one-half of the latter, but larger than that of the composite singlet neutrino. Also, this heavy $SU(2)_L$ doublet lepton couples to the Higgs field and singlet neutrino. This coupling is related to that entering neutrino mass seesaw, i.e., with SM replacing the composite doublet lepton. In this manner, the singlet neutrino production can occur via the decay of composite doublet lepton using Yukawa coupling, with the latter produced via decay of heavy W_L in addition to the exchange of SM W_L . Note that the couplings in both production (related to SM gauge) and decay (related to Yukawa) of composite doublet lepton are independent of representation of SM singlet neutrino under $SU(2)_R \times U(1)_X$. So, for simplicity and in order to reiterate the above model independence of this signal, we just drop the extended $(SU(2)_R \times U(1)_X)$ structure, while consider only a $SU(2)_L \times U(1)_Y$ composite sector model with composite singlet neutrino for this study of the composite $SU(2)_L$ doublet. We also assume composite hypercharge is heavy. Thus, the Lagrangian in the fermion-gauge sector for this signal is

$$\begin{aligned} \mathcal{L}_{\text{gauge-fermion}} = & -\frac{1}{4}F_{\mu\nu}^{(0)2} + \frac{1}{2}(D_\mu\rho_{W_L\nu} - D_\nu\rho_{W_L\mu})^2 + m_\star^2\rho_{W_L\mu}^2 + \overline{\psi^{(0)}}i\not{D}\psi^{(0)} + \overline{L_L^{(1)}}(i\not{D} - m_L)L_L^{(1)} \\ & + \overline{N^{(1)}}(i\not{D} - m_N)N^{(1)} - \frac{g_W^2}{g_\star}\overline{\psi_{\text{light}}^{(0)}}\rho_{W_L}^\mu\gamma_\mu\psi_{\text{light}}^{(0)} + \left(-\frac{g_W^2}{g_\star}\cos^2\phi_{Q_L^3} + g_\star\sin^2\phi_{Q_L^3}\right)\overline{Q_L^{(0)3}}\rho_{W_L}^\mu\gamma_\mu Q_L^{(0)3} \\ & + g_\star\overline{L_L^{(1)}}\rho_{W_L}^\mu\gamma_\mu L_L^{(1)}, \end{aligned} \quad (2)$$

where ρ_{W_L} denotes the heavy W_L gauge boson with coupling g_\star , which corresponds to KK W_L of the five-dimensional model. Note that in the above Lagrangian, ρ_{W_L} is the matrix of gauge bosons with appropriate generators. One should take the trace of the pure gauge part (the first row) of the above Lagrangian to get the final answer. Terms with more than two heavy spin-1 fields are dropped, simply because they are much smaller than the dominate interactions we consider. $L_L^{(1)}$ is the vectorlike/heavy $SU(2)_L$ doublet lepton and the singlet neutrino is denoted by $N^{(1)}$ which is also vectorlike, but has no charged lepton partner (cf. composite $(B-L)$ model or previous paper). Also, the elementary-composite W_L mixing angle is given approximately by $\frac{g_W}{g_\star}$, where g_W is SM W coupling. Similarly to the case in

composite $(B-L)$ model, we simply treat cosine of this angle to be 1, provided this mixing angle is small. Note that $N^{(1)}$ in this model does not have any gauge interactions (before EWSB). Finally, although it is not shown explicitly, the isospin charges are to be understood in the gauge couplings.

We would like to emphasize that the above channel exists even in the original model, i.e., with the SM singlet neutrino being the doublet of $SU(2)_R$. Qualitatively the signal will be similar to what we discuss here, although details at the $O(1)$ factors will be different. In fact, we choose a few TeV mass scale for heavy W_L in order to get enough signal for this model. With such a mass, EW precision tests actually require the extended EW structure. In particular, even though T parameter constraint can be satisfied with a suitably heavy

hypercharge, suppressing the shift in $Z \rightarrow b\bar{b}$ coupling with a custodial symmetry mandates composite $SU(2)_R$ gauge bosons which are degenerate with $SU(2)_L$. This will result in modifications of the signal. In this sense, our study could be taken as a toy or simplified version of the fully realistic case, but is clearly sufficient for the purpose of illustrating the basics of this signal.

A. Higgs sector

The SM Higgs doublet is taken to be purely in the composite sector, with *no* direct coupling to the elementary gauge bosons or fermions. In the cases that we study here, it has the following effects. The Higgs VEV gives mass to the SM EW gauge fields and SM fermions. The gauge and Yukawa couplings of the Higgs field also result in decay channels for heavy spin-1 and fermion into physical Higgs, as well as into longitudinal W/Z . In the unitary gauge, this actually arises via mass mixing (induced by the Higgs VEV) between various fields. The relevant cases here are heavy-heavy mixing for fermions and heavy-massless mixing for fermions and spin-1 states. In principle, the Higgs VEV induces heavy-heavy spin-1 mixing. Indeed, in [6], this effect was enhanced by degeneracy of spin-1 composites and played a crucial role in the LHC signals, in particular, resulting in significant coupling of composite W_R^\pm to light quarks. However, this particular Higgs induced mixing is rendered negligible for the purposes of LHC signals in this paper by our assumption of significant nondegeneracy among composite spin-1 states.

For a detailed explanation of the above features, it is better to treat the two LHC signal models separately.

(a) Composite ($B - L$) gauge boson

The bosonic part of Higgs Lagrangian is

$$\mathcal{L}_{\text{gauge-Higgs}} = |D_\mu \Phi|^2 - V(\Phi), \quad (3)$$

where $\Phi \equiv (\tilde{\Phi}, \Phi)$ is Higgs bi-doublet with $\tilde{\Phi} \equiv i\sigma_2 \Phi^*$ and Φ being Higgs doublet. We will use H for the physical Higgs boson. We can show (see Appendix) that ρ_{B-L} coupling to Higgs is doubly suppressed in elementary-composite mixing angles, so that it is neglected here. Thus, the only effect of Higgs here is to give masses to SM EW gauge fields.

Next, we consider the Yukawa couplings. While this is similar to [6], for the sake of completeness, we repeat it here. As already mentioned above, in this case study, we assume that the composite $SU(2)_L$ doublet $L_L^{(1)}$ is too heavy to be relevant for LHC signal, keeping only singlet neutrino (and its charged lepton-like partner). Thus, the relevant part of Lagrangian is

$$-\mathcal{L}_{\text{Yukawa}} = y_{00} \overline{L_L^{(0)}} \Phi \ell_R^{(0)} + y_{01} \overline{L_L^{(0)}} \Phi \tilde{L}_R^{(1)}, \quad (4)$$

where $y_{00}(y_{01})$ denotes Yukawa coupling between two massless (one massless and one heavy) modes. Here we drop flavor indices for simplicity, but there should be one copy of the above form for each generation of leptons. Also, $L_L^{(0)}$ and $\tilde{L}_R^{(1)}$ are given in doublet of $SU(2)_L$ and $SU(2)_R$ respectively. Their components are

$$\begin{aligned} L_L^{(0)} &= (\nu_L^{(0)}, \ell_L^{(0)}) \\ \tilde{L}_R^{(1)} &= (N^{(1)}, \tilde{\ell}_R^{(1)}), \end{aligned} \quad (5)$$

where first line is chiral $SU(2)_L$ doublet lepton, while the second line is heavy and vectorlike. Note that we are invoking ‘‘split’’ multiplets in the $SU(2)_R$ doublet lepton sector. That is, we introduce separate multiplets for SM right-handed charged lepton and for the singlet neutrino involved in seesaw, so that $\tilde{\ell}_R^{(1)}$ in second line above is not the heavy version (or KK excitation) of $\ell_R^{(0)}$ (see [6] for more details).

Clearly, the above couplings lead directly to decay of heavy fermions into SM fermions and the physical Higgs boson. In addition, after Higgs VEV, we get heavy-massless fermion mass mixing as follows. We begin with the neutrino sector, where these mass terms (along with vectorlike masses) can be written as

$$-\mathcal{L}_{\text{mass}} = \frac{y_{01} v}{\sqrt{2}} \overline{\nu_L^{(0)}} N_R^{(1)} + m_N \overline{N_L^{(1)}} N_R^{(1)}. \quad (6)$$

The resulting mass eigenstates then (approximately) are

$$\begin{aligned} N_R &\approx N_R^{(1)} \\ N_L &\approx N_L^{(1)} + V_{\ell N} \nu_L^{(0)} \\ \nu_L &\approx \nu_L^{(0)} - V_{\ell N} N_L^{(1)}, \end{aligned} \quad (7)$$

where N is the heavy, mostly SM singlet, mass eigenstate, while ν_L is the SM/massless one. The mixing angle is given by $V_{\ell N} \approx y_{01} v / (\sqrt{2} m_N)$. Here, we are treating the Higgs VEV effect as a perturbation, i.e., assuming the mass mixing terms above to be smaller than the vectorlike mass.

For the charged lepton sector, the relevant mass terms, including Higgs VEV-induced mixing, but dropping the negligible SM charged lepton Yukawa, i.e. the first term in Eq. (4), are given by

$$\frac{y_{01} v}{\sqrt{2}} \overline{\ell_L^{(0)}} \tilde{\ell}_R^{(1)} + m_N \overline{\tilde{\ell}_L^{(1)}} \tilde{\ell}_R^{(1)}. \quad (8)$$

Thus, the mass eigenstates are

$$\begin{aligned}
\tilde{\ell}_L &\approx \tilde{\ell}_L^{(1)} + V_{\ell N} \ell_L^{(0)} \\
\tilde{\ell}_R &\approx \tilde{\ell}_R^{(1)} \\
\ell_L &\approx \ell_L^{(0)} - V_{\ell N} \tilde{\ell}_L^{(1)}.
\end{aligned} \tag{9}$$

Again, ℓ_L here is the SM field, whereas $\tilde{\ell}$ is the heavy, mostly $SU(2)_L$ singlet, mass eigenstate.

Recall that the fields with superscripts “(0)” and “(1)” on the RHS of the mass eigenstate equations above are the mass eigenstates before EWSB. Equivalently, they are the weak/gauge eigenstates, with gauge couplings given in Eq. (1). So, we need to reexpress the gauge couplings in Eq. (1) in terms of the final mass eigenstate via the inverse of the above mass eigenstate equations. In particular, the EWSB-induced $SU(2)_L$ doublet-singlet mass mixing will then result in the N (similarly for $\tilde{\ell}$) coupling to (and thus decaying into) the SM lepton $SU(2)_L$ doublet and W/Z . Of course, this unitary gauge effect (i.e., via gauge coupling) is equivalent to thinking of longitudinal W/Z as unphysical Higgs so that this decay proceeds simply via the Yukawa coupling instead.

We will return to all these couplings in the next section.

(b) Composite $SU(2)_L$ doublet lepton

As a reminder, for simplicity, we neglect here the $SU(2)_R \times U(1)_X$ structure so that the bosonic part of Higgs Lagrangian is

$$\mathcal{L}_{\text{gauge-Higgs}} = |D_\mu \Phi + ig_* \rho_{W_L} \Phi|^2 - V(\Phi). \tag{10}$$

We thus get decay of ρ_{W_L} into physical Higgs boson and W/Z simply from the cross term in the above covariant derivative. In addition, the Higgs VEV leads to mass mixing between $W_L^{(0)}$ and ρ_{W_L} , resulting in decay of the heavier mass eigenstate into a pair of W/Z via the SM trilinear coupling of the $W_L^{(0)}$ component of the heavy state/combination¹⁴: this is the same phenomenon as in [6] and earlier studies, so we will just refer to it for details. However, the mass of the final heavy spin-1 state (and its couplings to fermions) are negligibly shifted compared to those of ρ_{W_L} . Thus for simplicity of notation, we will continue to denote this particle by “ ρ_{W_L} .”

Moving onto Yukawa couplings, the heavy $SU(2)_L$ doublet lepton in weak/gauge basis is denoted in component form as

$$L_L^{(1)} = (\nu^{(1)}, \ell^{(1)}), \tag{11}$$

which is vectorlike. Once again, the heavy/vectorlike SM singlet neutrino is labeled as $N^{(1)}$. The relevant Lagrangian becomes

$$\begin{aligned}
-\mathcal{L}_{\text{Yukawa}} &= y_{00} \overline{L_L^{(0)}} \Phi \ell_R^{(0)} + y_{01} \overline{L_L^{(0)}} \tilde{\Phi} N_R^{(1)} \\
&\quad + y_{10} \overline{L_L^{(1)}} \Phi \ell_R^{(0)} + y_{11} \overline{L_L^{(1)}} \tilde{\Phi} N_R^{(1)},
\end{aligned} \tag{12}$$

where $y_{10}(y_{11})$ denotes Yukawa coupling between one massless and one heavy (two heavy) modes. y_{11} is naturally greater than y_{10} or y_{01} due to the later being suppressed by one power of elementary-composite mixing. Including mass mixing terms (along with the dominant vectorlike masses) in the neutrino sector, one can obtain

$$\begin{aligned}
\frac{y_{01} v}{\sqrt{2}} \overline{\nu_L^{(0)}} N_R^{(1)} + \frac{y_{11} v}{\sqrt{2}} \overline{\nu_L^{(1)}} N_R^{(1)} + m_N \overline{N_L^{(1)}} N_R^{(1)} \\
+ m_L \overline{\nu_L^{(1)}} \nu_R^{(1)}.
\end{aligned} \tag{13}$$

Assuming $y_{11} \gg y_{01}$ and $m_L > m_N > v$, the resulting mass eigenstates can be approximated as

$$\begin{aligned}
\nu_L &\approx \nu_L^{(0)} - V_{\ell N} N_L^{(1)} \\
N_L &\approx N_L^{(1)} - V_{LN} \frac{m_N}{m_L} \nu_L^{(1)} + V_{\ell N} \nu_L^{(0)} \\
N_R &\approx N_R^{(1)} - \nu_R^{(1)} V_{LN} \\
\nu_L^h &\approx \nu_L^{(1)} + V_{LN} \frac{m_N}{m_L} N_L^{(1)} \\
\nu_R^h &\approx \nu_R^{(1)} + V_{LN} N_R^{(1)},
\end{aligned} \tag{14}$$

where $V_{\ell N}$ is the same as before and $V_{LN} \approx y_{11} v m_L / (\sqrt{2}(m_L^2 - m_N^2))$. ν^h denotes heavy, approximately $SU(2)_L$ doublet, mass eigenstate (rest of the notation is as before).

Similarly, in the charged lepton sector, we have the relevant mass terms

$$\frac{y_{10} v}{\sqrt{2}} \overline{\ell_L^{(1)}} \ell_R^{(0)} + m_L \overline{\ell_L^{(1)}} \ell_R^{(1)}, \tag{15}$$

giving

$$\begin{aligned}
\ell_R^h &\approx \ell_R^{(1)} + V_{\ell L} \ell_R^{(0)} \\
\ell_R &\approx \ell_R^{(0)} - V_{\ell L} \ell_R^{(1)} \\
\ell_L &\approx \ell_L^{(0)} \\
\ell_L^h &\approx \ell_L^{(1)},
\end{aligned} \tag{16}$$

where ℓ^h denotes the heavy, approximately $SU(2)_L$ doublet, charged lepton mass eigenstate and $V_{\ell L} \approx y_{10} v / (\sqrt{2} m_L)$.

¹⁴Equivalently, as done above for fermions, we can obtain the same effect by thinking of longitudinal W/Z as unphysical Higgs.

Similarly to the above case, the decay of the heavy $SU(2)_L$ doublet states (i.e., ν^h and ℓ^h) into singlet neutrino state (N) and physical Higgs boson follows directly from Eq. (13), whereas decay into longitudinal W/Z (plus singlet neutrino) is the result of Higgs VEV induced mass mixing in the unitary gauge.

IV. OVERVIEW OF LHC SIGNALS

In this section, we describe our signal channels for each of the two models presented above. They both involve production of the singlet neutrino via decays of heavier states. We first explicitly show the interactions which are directly relevant for our signals. Then, we introduce parameter choices for each process which are consistent with current bounds. Analytic expressions for the decay widths for the new, heavy particles are listed. Since we work on two different models (depending on the immediate parent particle for the singlet neutrino), we will discuss each model in a separate subsection. Composite $(B-L)$ model is analyzed in Sec. IV A, while Sec. IV B is devoted to the composite $SU(2)_L$ doublet lepton model.

A. Composite $(B-L)$

In this section, we consider the scenario where the composite sector only has $(B-L)$ gauge field; i.e., we assume that all $SU(2)_L \times SU(2)_R$ gauge bosons are heavy and, hence, irrelevant for our signal study. For producing this new particle at the LHC, we need a significant coupling to light quarks inside the proton. However, as discussed in detail in Sec. III, the light quarks being elementary do not directly couple to the composite gauge bosons including composite $(B-L)$ gauge boson. Nonetheless, the composite $(B-L)$ gauge boson mixes with the elementary hypercharge gauge boson. The heavy mass eigenstate (denoted by ρ_{B-L}) can then couple to light quarks via this small admixture of elementary hypercharge gauge boson. In short, our signal channel is then the produced ρ_{B-L} decaying to a pair of N 's. Subsequently, each N decaying to ℓ and W or ν and Z/H .

1. Relevant couplings

Here, we summarize the relevant couplings for our signal. For notational details, the reader is referred to Sec. III. There are three types of couplings: (1) couplings between ρ_{B-L} and SM fermions (2) couplings of ρ_{B-L} to a pair of N or a pair of $\tilde{\ell}$, and (3) couplings among $N(\tilde{\ell})$, $H/W/Z$ and SM $\ell(\nu)$ via Yukawa coupling.

(1) The first type of coupling can be obtained from Eq. (1):

$$\begin{aligned} \delta L_{(1)} = & -Q_Y \frac{g_Y^2}{g_\star} \rho_{B-L}^\mu \bar{\Psi}_{\text{light}} \gamma_\mu \Psi_{\text{light}} \\ & + \frac{1}{6} g_\star \sin^2 \phi_{Q_L^3} \rho_{B-L}^\mu \bar{Q}_L^3 \gamma_\mu Q_L^3 \\ & + \frac{1}{6} g_\star \rho_{B-L}^\mu \bar{t}_R \gamma_\mu t_R, \end{aligned} \quad (17)$$

where g_Y is SM hypercharge coupling and g_\star is composite $(B-L)$ coupling. Q_Y is SM hypercharge of corresponding light fermion Ψ_{light} and factor of $\frac{1}{6}$ in the last two terms arises from the $(B-L)$ charge of Q_L^3 and t_R . In the second coupling, we have assumed that $g_\star \sin^2 \phi_{Q_L^3}$, is large enough that the component arising from spin-1 mixing, i.e., $\propto \frac{g_Y^2}{g_\star} \cos^2 \phi_{Q_L^3}$, can be neglected here. All fermions in all Lagrangians shown in this section are mass eigenstates after EWSB. Since EWSB has negligible effects for all $\psi^{(0)}$ s, we simply drop superscript “(0)” to denote their mass eigenstates.

These couplings are responsible for the production of ρ_{B-L} via light quarks inside proton and also constitute decays channels for the ρ_{B-L} .

(2) The second type of coupling can be understood from Eq. (1):

$$\delta \mathcal{L}_{(2)} = -\frac{1}{2} g_\star \rho_{B-L}^\mu \bar{N} \gamma_\mu N - \frac{1}{2} g_\star \rho_{B-L}^\mu \bar{\tilde{\ell}} \gamma_\mu \tilde{\ell}, \quad (18)$$

where $-\frac{1}{2}$ is the $(B-L)$ charge of N and $\tilde{\ell}$. These couplings lead to the dominant decays of ρ_{B-L} to a pair of N and a pair of $\tilde{\ell}$.

(3) The third type of couplings are similarly obtained from Eqs. (1) and (4) and mixing induced by EWSB Eqs. (7) and (9):

$$\begin{aligned} \delta \mathcal{L}_{(3)} = & \frac{g_W}{\sqrt{2}} V_{\ell N} W_\mu^+ \bar{N}_L \gamma^\mu \ell_L + \{N \leftrightarrow \nu; \ell \leftrightarrow \tilde{\ell}\} \\ & + \frac{g_Z}{2} V_{\ell N} Z_\mu \bar{N}_L \gamma^\mu \nu_L - \frac{y_{01}}{\sqrt{2}} H \bar{N}_R \nu_L \\ & + \{N \leftrightarrow \tilde{\ell}; \nu \leftrightarrow \ell\} + \text{H.c.}, \end{aligned} \quad (19)$$

where $g_{W/Z}$ is the SM W/Z gauge coupling and y_{01} is the Yukawa coupling defined in Eq. (4). These couplings lead to the decays of N and $\tilde{\ell}$ to $H/W/Z$ and ℓ/ν : note that this part is the same as in our previous paper on LHC signals.

2. Parameter choice

For the composite $(B-L)$ model, the only relevant composite gauge boson is ρ_{B-L} , thus the only relevant parameters in spin-1 composite sector being g_\star and m_\star . In general, there is a lower bound on these g_\star 's from the requirement that the Landau poles for the SM gauge couplings are above the GUT/Planck scale, which turns out to be ≈ 3 [15] for hypercharge in $SU(5)$ normalization (this was studied in the context of gauge coupling unification, hence this choice of normalization). Converting this to SM normalization for hypercharge used here, we get $g_{\star Y} \geq 3\sqrt{\frac{3}{5}}$. In our model, the composite sector has $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ symmetry, with composite $U(1)_Y$ obtained from a linear combination of $U(1)_R$ (the

$U(1)$ part of $SU(2)_R$ and $U(1)_{B-L}$. This structure leads to the relation among couplings as $g_{\star Y} = \frac{g_{\star R} g_{\star}}{\sqrt{g_{\star R}^2 + g_{\star}^2}}$, where $g_{\star R}$ is the composite gauge coupling for $SU(2)_R$. Since $SU(2)_L \times SU(2)_R$ gauge fields in the model are irrelevant for signal study, we can choose $g_{\star R} \gg g_{\star}$. In this case, the lower bound on $g_{\star Y}$ becomes the lower bound on g_{\star} . Therefore, our model requires $g_{\star} \geq 3\sqrt{\frac{3}{5}}$. For this collider analysis, we choose $g_{\star} = 2.5$ as a benchmark point.

The masses for spin-1 composites are in general constrained by EW precision tests. Since we choose the relevant states, i.e., $SU(2)_L \times SU(2)_R$ composite gauge bosons, to be heavy ≥ 4 TeV, the EW precision tests are not major constraints for the mass m_{\star} of ρ_{B-L} in our model. Moreover, the amount of compositeness of Q_L^3 , i.e., $\sin\phi_{Q_L^3}$, is also not much constrained by EW precision tests for the same reason. We choose $\sin^2\phi_{Q_L^3} = 0.21$ in this study just for consistency with our previous paper on LHC signals [6].

In addition, we have to consider bounds from direct searches at the LHC for these spin-1 states in the various decay channels with two SM particles. However, for ρ_{B-L} , di-lepton and di-jet branching ratios are negligible due to the associated couplings being suppressed by $\frac{g_{\star}^2}{g_{\star}^2}$ compared to the couplings to N and $\tilde{\ell}$. The coupling for top quark is not only unsuppressed but also it gets enhanced by a color multiplicity factor. However, $(B-L)$ charge for top being $\frac{1}{6}$ is smaller than that of N and $\tilde{\ell}$ by a factor of 3. Moreover, heavy leptons are vectorlike, providing a factor of two, and also come in three generations, providing multiple decaying channels. Total branching ratio to the ditop is thus much smaller than the dominant channel N (as discussed in next section): we have checked that the bound from the ditop searches is indeed very weak. The diboson search provides usually one of the most stringent bounds for warped/composite gauge bosons in the EW sector. However, in our model, the light composite gauge boson is ρ_{B-L} , which does not (directly) couple to Higgs. Therefore, we are safe from diboson constraints also. In conclusion, there is no direct bound on m_{\star} in our model, and we simply choose $m_{\star} = 2$ TeV for the rest of study.

Next, we discuss the parameter choices in the lepton sector. The elementary-composite mixing, i.e., $|V_{\ell N}|^2$, is constrained by various experiments. As we discussed previously in [6], we simply choose $|V_{\ell N}|^2 = 0.001$ for all three generations to be consistent with experimental bounds. Finally, the constraint on the mass of N and $\tilde{\ell}$ is correlated with the choice of $|V_{\ell N}|^2$. With the choice we make $|V_{\ell N}|^2 = 0.001$, however, the lower bound on m_N is $O(100)$ GeV [11]. Nevertheless, we require ρ_{B-L} to decay into a pair of N or $\tilde{\ell}$, implying $m_N < \frac{1}{2}m_{\star}$. Besides, since the $(B-L)$ gauge boson and N all come from the same composite sector, there should not be a big hierarchy

between m_N and m_{\star} . Taking all these considerations into account, we choose $m_N = 750$ GeV in this study.

3. Decay widths

In this section we show analytic expressions for decay widths for relevant particles. All the decay widths presented here assume $m_{\star} > 2m_N \gg$ mass of SM particles, thus masses of SM particles being reasonably neglected.

Composite ρ_{B-L} .—Decay width for each ρ_{B-L} decay channel is computed using couplings in Eqs. (17) and (18) and is shown below,

$$\begin{aligned}\Gamma(\rho_{B-L} \rightarrow N_i \bar{N}_i / \tilde{\ell}_i \bar{\tilde{\ell}}_i) &= g_{\star}^2 \left(1 + 2 \frac{m_N^2}{m_{\star}^2}\right) \sqrt{1 - 4 \frac{m_N^2}{m_{\star}^2}} \frac{m_{\star}}{48\pi} \\ \Gamma(\rho_{B-L} \rightarrow t_R \bar{t}_R) &= g_{\star}^2 \frac{m_{\star}}{288\pi} \\ \Gamma(\rho_{B-L} \rightarrow t_L \bar{t}_L / b_L \bar{b}_L) &= g_{\star}^2 \sin^4 \phi_{Q_L^3} \frac{m_{\star}}{288\pi} \\ \Gamma(\rho_{B-L} \rightarrow \psi \bar{\psi}) &= N_c Q_Y^4 \frac{g_Y^4}{g_{\star}^2} \frac{m_{\star}}{24\pi},\end{aligned}\quad (20)$$

where ψ denotes *light* (other than top/bottom) SM fermions, and N_c is the color factor of corresponding fermion ψ : 3 for quarks and 1 for leptons. Also, the subscript i for N and $\tilde{\ell}$ is generation index ($i = e, \mu, \tau$): note that these particles are vectorlike.

Composite $SU(2)_R$ doublet $(N, \tilde{\ell})$.—Their decays proceeded via the couplings in Eq. (19), giving

$$\begin{aligned}\Gamma(N \rightarrow W\ell) &= y_{01}^2 \frac{m_N}{32\pi} \\ \Gamma(N \rightarrow H/Z\nu) &= y_{01}^2 \frac{m_N}{64\pi} \\ \Gamma(\tilde{\ell} \rightarrow W\nu) &= y_{01}^2 \frac{m_N}{32\pi} \\ \Gamma(\tilde{\ell} \rightarrow H/Z\ell) &= y_{01}^2 \frac{m_N}{64\pi}.\end{aligned}\quad (21)$$

In principle, there will be three body decays via virtual ρ_{W_L} or ρ_{W_R} . However, these are suppressed compared to two-body decays given above due to ρ_{W_L} and ρ_{W_R} being heavy.

B. Composite $SU(2)_L$ doublet

In this section, we consider the case where the composite sector has $SU(2)_L \times U(1)_Y$ global symmetry: we assume that the composite hypercharge gauge boson is heavy, thus we keep only the composite $SU(2)_L$ gauge boson. Also, in the composite fermion sector, we add to the first model the composite partner of SM lepton doublet, denoted as (ℓ^h, ν^h) , with the singlet neutrino N now not having a charged partner (unlike in the first model): for details, see Sec. III. Our signal channel involves heavy spin-1 state

(denoted by ρ_{W_L}) being produced via light quarks through mixing with elementary W boson. This is followed by ρ_{W_L} decaying to (ℓ^h, ν^h) pairs. Each $\ell^h(\nu^h)$ could further decay to N and $W(H/Z)$ through Yukawa coupling, with each N decaying to ℓ and W or ν and Z/H .

1. Relevant couplings

There are four types of couplings that we need to consider: (1) couplings between composite ρ_{W_L} and SM fermions; (2) couplings between composite ρ_{W_L} and composite $SU(2)_L$ lepton doublet (ℓ^h, ν^h) ; (3) Yukawa couplings of (ℓ^h, ν^h) to singlet N and Higgs, and (4) Yukawa couplings of N to SM $W\ell$ or $H\nu$.

- (1) The first type of coupling can be obtained by using Eq. (2):

$$\delta\mathcal{L}_{(1)} = -\frac{g_W^2}{\sqrt{2}g_\star} \rho_{W_L\mu}^+ \bar{\psi}_L \gamma^\mu \psi'_L + \frac{g_\star}{\sqrt{2}} \sin^2\phi_{Q_L^3} \rho_{W_L\mu}^+ \bar{t}_L \gamma^\mu b_L + \text{H.c.}, \quad (22)$$

where g_\star is composite ρ_{W_L} coupling. These couplings are responsible for the production of ρ_{W_L} via light quarks inside proton and some of the ρ_{W_L} decay channels.

- (2) The second type of coupling can be understood from Eq. (2):

$$\delta\mathcal{L}_{(2)} = \frac{g_\star}{\sqrt{2}} \rho_{W_L\mu}^+ \bar{\ell}^h \gamma^\mu \nu^h + \text{H.c.} \quad (23)$$

This coupling leads to the decays of composite ρ_{W_L} to ℓ^h and ν^h .

- (3) The third type of coupling are obtained from Yukawa coupling in Eq. (2) and EWSB effect in Eq. (14):

$$\begin{aligned} \delta\mathcal{L}_{(3)} = & -\frac{g_W}{\sqrt{2}} V_{LN} \frac{m_N}{m_L} W_\mu^+ \bar{N}_L \gamma^\mu \ell_L^h \\ & -\frac{g_Z}{2} V_{LN} \frac{m_N}{m_L} Z_\mu \bar{N}_L \gamma^\mu \nu_L^h \\ & -\frac{g_W}{\sqrt{2}} V_{LN} W_\mu^+ \bar{N}_R \gamma^\mu \ell_R^h - \frac{g_Z}{2} V_{LN} Z_\mu \bar{N}_R \gamma^\mu \nu_R^h \\ & -\frac{y_{11}}{\sqrt{2}} H \bar{N}_R \nu_L^h + \text{H.c.}, \quad (24) \end{aligned}$$

where y_{11} is the Yukawa coupling defined in Eq. (12). These couplings lead to the decays of ℓ^h to W and N and ν^h to H/Z and N .

- (4) The fourth type of coupling are obtained from Yukawa coupling in Eq. (2) and EWSB in Eq. (14), which is essentially the same as Eq. (19):

$$\begin{aligned} \delta\mathcal{L}_{(4)} = & \frac{g_W}{\sqrt{2}} V_{\ell N} W_\mu^+ \bar{N}_L \gamma^\mu \ell_L + \frac{g_Z}{2} V_{\ell N} Z_\mu \bar{N}_L \gamma^\mu \nu_L \\ & -\frac{y_{01}}{\sqrt{2}} H \bar{N}_R \nu_L + \text{H.c.} \quad (25) \end{aligned}$$

These couplings lead to the decays of N to W and ℓ or H and ν .

2. Parameter choice

As mentioned in Sec. III, in this model, we choose the global symmetry of composite sector to be $SU(2)_L \times U(1)_Y$ for the purpose of this collider study only. However, the realistic model should have full $SU(2)_L \times SU(2)_R \times U(1)_X$ global symmetry in composite sector, in order to satisfy all EW precision bounds with a few TeV compositeness scale. In other words, the model being studied here is to be viewed as its simplified version. This full model is almost the same as the model in our previous paper on LHC signals, thus a similar estimation of the bounds and parameter choices also apply here. Therefore, we choose the same value of $g_\star = 3$ and $\sin^2\phi_{Q_L^3} = 0.21$. In the present study, m_\star is chosen to be 2.5 TeV, slightly bigger value than the choice previously ($m_\star = 2$ TeV). As mentioned in [6] already, for these parameter choices, even with m_\star being 2.5 TeV, some additional model building might still be needed to be fully safe from EW precision bounds.

As mentioned earlier, diboson searches at the LHC usually set a stringent bound on composite Higgs models. However, in our model, the branching ratio of ρ_{W_L} to diboson is suppressed compared to that of standard KK W . This is because the dominant decay channels of ρ_{W_L} are $\rho_{W_L} \rightarrow \ell^h \nu^h$. As will be shown clearly in Sec. IV B 3, the branching ratio to the diboson is $\sim \frac{1}{13}$. With our choice of $g_\star = 3$, we then checked that the current bound is $m_\star > 1.5$ TeV [16]. Obviously, our choice of m_\star is then safe from this bound.

Moreover, there are two lepton mixings in this model. V_{LN} is the mixing between ν^h and N and $V_{\ell N}$ is the mixing between N and SM ν . $V_{\ell N}$ is the same mixing angle in our previous study of LHC signals (and in composite $(B-L)$ model here). So, we choose $|V_{\ell N}|^2 = 0.001$ for all three generations in order to be safe with the relevant bounds and also for consistency with our other studies. However, there is no direct bound on V_{LN} , i.e., the composite-composite mixing angle, which is the new feature in this model. Also, this mixing will be larger than the mixing between one composite and one elementary ($V_{\ell N}$): this point was also mentioned in Sec. III. In this study, we choose $|V_{LN}|^2 = 0.01$.

Our signal assumes the following cascade decay: composite ρ_{W_L} goes into ℓ^h and ν^h pair, which further decay to N and $W/H/Z$. This requires the spectrum to satisfy $\frac{1}{2}m_\star > m_L > m_N$. Similarly to our previous cases

(in [6] or composite $(B-L)$ model here), there is very weak bound ($O(100)$ GeV) on m_L and m_N provided we make the choices of V_{LN} and $V_{\ell N}$ discussed above. At the same time, too much gap between these states is unnatural. So, we pick $m_L = 1000$ GeV and $m_N = 500$ GeV.

3. Decay widths

All the decay widths presented below are with the assumption $\frac{1}{2}m_* > m_L > m_N \gg$ mass of SM particles, thus masses of SM particles are reasonably neglected.

Composite ρ_{W_L} .—As mentioned above, $SU(2)_L$ doublet composite leptons are produced via decays of composite ρ_{W_L} using the coupling in Eq. (23). In addition, ρ_{W_L} can decay into pair of SM fermions via the couplings in Eq. (22). Decay widths for composite ρ_{W_L} are then given as

$$\begin{aligned}\Gamma(\rho_{W_L} \rightarrow \ell_i^h \nu_i^h) &= g_*^2 \left(1 + 2 \frac{m_L^2}{m_*^2}\right) \sqrt{1 - 4 \frac{m_L^2}{m_*^2}} \frac{m_*}{24\pi} \\ \Gamma(\rho_{W_L} \rightarrow WZ/WH) &= g_*^2 \frac{m_*}{192\pi} \\ \Gamma(\rho_{W_L} \rightarrow t_L b_L) &= g_*^2 \sin^4 \phi_{Q_L} \frac{m_*}{16\pi} \\ \Gamma(\rho_{W_L} \rightarrow \psi\psi') &= N_c \frac{g_W^4}{g_*^2} \frac{m_*}{48\pi}.\end{aligned}\quad (26)$$

Composite $SU(2)_L$ doublet (ℓ^h, ν^h) .—Similarly, using Eq. (24), decay widths for composite (ℓ^h, ν^h) are given as

$$\begin{aligned}\Gamma(\ell^h \rightarrow WN) &= y_{11}^2 \left(1 - \left(\frac{m_N}{m_L}\right)^4\right) \frac{m_L}{32\pi} \\ \Gamma(\nu^h \rightarrow H/ZN) &= y_{11}^2 \left(1 - \left(\frac{m_N}{m_L}\right)^4\right) \frac{m_L}{64\pi}.\end{aligned}\quad (27)$$

Decay widths for N are the same as Eq. (21).

V. DISCOVERY POTENTIAL

In this section, we present our results for the phenomenological studies of the LHC signals for the model described in Sec. III. In Sec. VA, we study the pair production of the singlet neutrino (N) via the on-shell decay of composite gauge boson ρ_{B-L} . Once produced, each N can decay into $\ell + W$ or $\nu + H/Z$, followed by SM W decaying into $\ell\nu$ or jj . Thus, the cascade decays of a pair of N 's can result in several possible final states. As our benchmark, in this paper, we consider the following decay channel.¹⁵

$$\begin{aligned}pp \rightarrow \rho_{B-L} \rightarrow NN, \quad N \rightarrow \ell W, \quad (W \rightarrow jj), \\ (W \rightarrow \ell\nu).\end{aligned}\quad (28)$$

Namely, both N 's decay first into a lepton and W boson. One produces W decays leptonically, and the other decays hadronically. As a result, the final state consists of three leptons, two from direct decay of N and one from leptonic decay of W , two jets from hadronically decaying W and a neutrino in the form of missing energy or MET (thus labeled $3\ell 2j + \text{MET}$). Since the production of N is achieved by the s -channel ρ_{B-L} composite gauge boson, we call this channel the “ ρ_{B-L} -channel.”

Section VB is devoted to describing the study for the pair production of N via decays of the composite $SU(2)_L$ lepton doublet, (ν^h, ℓ^h) , which is produced by the on-shell decay of charged composite $SU(2)_L$ gauge boson ρ_{W_L} . Similarly to the above study, several final states are possible, again, depending on how N and W decay. In the current study, we consider the following cascade decay channel:

$$\begin{aligned}pp \rightarrow \rho_{W_L} \rightarrow \ell^h \nu^h, \quad \ell^h \rightarrow NW, \\ \nu^h \rightarrow NH/Z, \quad N \rightarrow \ell W, \\ H/Z \rightarrow b\bar{b}, \quad \text{Two } W \rightarrow jj, \quad \text{One } W \rightarrow \ell\nu.\end{aligned}\quad (29)$$

In detail, the cascade decay of composite gauge boson ρ_{W_L} produces composite lepton doublet, (ν^h, ℓ^h) , which in turn decay into two singlet neutrino N , one H/Z and one W . Subsequent decay of two N , then, produce two more W 's and two SM leptons. H/Z decays into $b\bar{b}$ and two of the three W 's decay hadronically, rendering four jets, and one W decays leptonically, producing one lepton and missing energy ($4j 2b 3\ell + \text{MET}$). Notice that we combine contributions from processes with H and Z intermediate states. This is because resolutions of LHC detectors may not be able to distinguish the two cases, and at the same time, we achieve a slight increase in the signal rate. Since the singlet neutrino N is produced by the decay of $SU(2)_L$ composite doublet lepton, we call this channel the “composite $SU(2)_L$ doublet-channel.” The tree-level Feynman diagrams for both signal processes are shown in Fig. 1. The topology of our signal processes are characterized by several resonance peaks in various invariant mass distributions: invariant masses of S -channel gauge bosons, which we take to be $M_{\rho_{B-L}} = 2$ TeV and $M_{\rho_{W_L}} = 2.5$ TeV, invariant mass of singlet neutrino N , which we take $M_N = 750$ (500) GeV for ρ_{B-L} (composite $SU(2)_L$ doublet)-channel, and invariant mass of composite $SU(2)_L$ doublet leptons, which we take $M_{\ell^h} = M_{\nu^h} = 1$ TeV. All of these invariant mass distributions, when successfully formed, will draw sharp distinctions between signal and SM backgrounds, allowing to achieve large significance. Notice, however, that both of our signal processes accompany *single* neutrino in the form of missing energy. In order to reconstruct all (or most) of resonance peaks, therefore, the reconstruction of the longitudinal component of the neutrino momentum is a must. The presence of multiple leptons/jets in the signal processes then require that lepton/jet identification, i.e., the

¹⁵To avoid notational clutter, we drop the particle's charge.

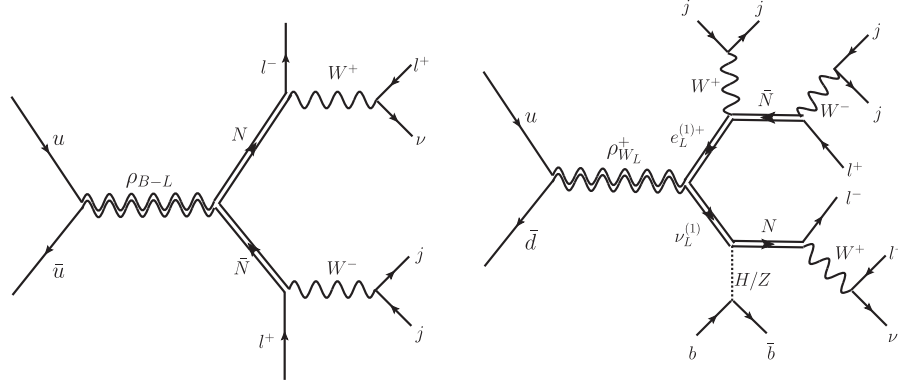


FIG. 1. The left panel shows Feynman diagram for the signal process of ρ_{B-L} -channel. The right panel shows Feynman diagram for the signal process of composite $SU(2)_L$ doublet-channel. Double (single) lines denote composite (SM) particles.

correct pairing of leptons and jets, needs to be figured out along with neutrino momentum reconstruction. We show below that appropriate understanding of kinematics of the signal processes together with the help of symmetries, e.g. $SU(2)_L$, make this seemingly difficult task possible. We wish to emphasize the importance of reconstructing resonance peaks beyond simply achieving large $S/\sqrt{S+B}$. Accomplishing large enough significance is of course crucial in order to be able to reveal new physics out of overwhelming SM backgrounds. However, once discovered, immediate following task will be to understand the underlying physics that gives rise to such events. Successful reconstruction of all resonance peaks will answer the most important part of the questions: new particle content and spectrum.

In this study, we constrain ourselves to $\Delta R_{jj} \geq 0.4$ focusing only on regular (as opposed to boosted/fat) jets from W decay. Inclusion of boosted W -jet contribution will yield increase in signal rate, resulting in larger significance $S/\sqrt{S+B}$. We leave study of boosted jet final states as well as study of several other decay channels mentioned above for future investigation.

Event simulations are performed by employing a sequence of simulation tools. We first created our two-site simplified model files using FEYNRULES [17] based on heavy vector triplets models [18]. Then we used them as input models in a Monte Carlo event generator MG5AMC@NLO [19] to generate parton-level events. In this procedure, parton distribution functions parameterized by NN23LO1 [20] is used. All the simulations are done at the leading order with a $\sqrt{s} = 14$ TeV pp collider. The generated parton-level events are then streamlined to PYTHIA 6.4 [21] to take care of showering and hadronization/fragmentation. Since all our channels contain only regular jets, we directly pass on the output from PYTHIA 6.4 to DELPHES 3 [22]. DELPHES 3, interfaced with FASTJET [23,24], provides a way to incorporate the detector effects and jet formation. This includes b -(miss)tagging and lepton/photon identifications at the detector, which usually

depend on p_T and rapidity of the object under consideration. The jets are constructed with the anti- k_r algorithm [24] with a radius parameter $R = 0.4$.

A. ρ_{B-L} channel

In this section, we present the results for the ρ_{B-L} channel: the relevant couplings for this channel are given in Sec. IV A. In the ρ_{B-L} -channel, as explained above, singlet N is pair-produced via the decay of neutral composite ρ_{B-L} gauge boson. This is possible because ρ_{B-L} is the gauge boson of $U(1)_{B-L}$ and N carries lepton number. The production of ρ_{B-L} becomes possible thanks to the mixing between ρ_{B-L} and elementary hypercharge gauge boson B , which then couples to quarks inside the proton. N decays dominantly to ℓW using Yuakwa coupling. As outlined above, we consider the case where one W decays to jj and the other to $\ell\nu$. Consequently, the final state consists of $2j3\ell + \text{MET}$.

There are multiple invariant mass variables that will turn out to be very efficient and crucial in background suppression. These include M_{jj} , $M_{\ell\nu}$, $M_{\ell\ell\nu}$, $M_{\ell jj}$, $M_{\ell\ell\ell}$, and M_{All} , where M_{All} is the invariant mass constructed from all visible objects, not including MET, and ℓ_ν is the lepton that groups with neutrino to reconstruct leptonically decaying W . Due to the fact that there are three leptons in the process, particle groupings, e.g. $\{\ell jj\}$ v.s. $\{\ell\ell_\nu\nu\}$ which reconstruct N separately, is not obvious *a priori*. We will describe our grouping prescription below and show that we can reconstruct all particle bumps, resolving ambiguity in particle pairings. When properly reconstructed, signal distribution of these variables will be peaked at $M_{jj} = M_{\ell\nu} = M_W$, $M_{\ell\ell\nu} = M_{\ell jj} = M_N$, and $M_{\text{All}} = M_{\rho_{B-L}}$. The distribution of $M_{\ell\ell\ell}$ does not have direct connection to resonance peak, nevertheless, it will provide a very strong cut.

There are several SM backgrounds we need to consider and we describe them one by one now. As it will become relevant soon, we simulate processes to leading order in QCD and QED couplings. Namely, the set of diagrams

considered in the numerical simulations is those with minimum order in QCD and QED couplings.

- (1) $jj\ell\ell W$: The relevant process is $pp \rightarrow jj\ell\ell W$, $W \rightarrow \ell\nu$. This is the dominant background with largest cross section. Since two leptons (not from W decay) in this process arise mostly from near-on-shell Z decay, $M_{\ell\ell}$ distribution is sharply peaked at M_Z . On the other hand, in signal, they are from direct decays of N and hence smoother $M_{\ell\ell}$ distribution. For this reason, the condition $M_{\ell\ell} \neq m_Z$ will achieve significant background suppression. However, we find that practically, $M_{\ell\ell\ell}$ distribution provides slightly sharper distinction between signal and backgrounds. Most of the other invariant mass variables described above will be important in reducing this background.
- (2) $t\bar{t}\ell\ell$: The relevant process is $pp \rightarrow t\bar{t}\ell\ell$, with each top decaying to $b + W$. One of the W 's decays hadronically giving jj , the other decays leptonically to $\ell\nu$. Since we use inclusive event selection criteria, i.e., $2j3\ell + X$, this process, even if it contains an extra two b 's, is indeed one of the backgrounds. Background reduction can be achieved in a similar manner as $jj\ell\ell W$.
- (3) $t\bar{t}W$: The relevant process is $pp \rightarrow t\bar{t}W$, with each top decaying to $b + W$. All three W 's decay leptonically generating $\ell\ell\ell + \text{MET}$. In order to pass the event selection criteria, both b 's must be untagged as b -jets and detected as regular jets. This helps reduce this background.
- (4) $WW\ell\ell$: The relevant process is $pp \rightarrow WW\ell\ell$, with one W decaying to jj , the other to $\ell\nu$. However, naively, one can think that this process is already included in $jj\ell\ell W$: just let one of the W 's decay to jj and notice that the final state is precisely that of $jj\ell\ell W$. Nonetheless, we include this background separately and yet we do not have double counting issue. This is because the leading coupling order for this process is QCD = 0 and QED = 6 in MG5AMC@NLO, and yet that for $jj\ell\ell W$ is QCD = 2 and QED = 4 and hence not included there.¹⁶

Defining N_ℓ and N_j as the number of isolated leptons and regular jets, respectively, we select events using the following selection criteria:

$$\begin{aligned} N_\ell &> 2 \quad \text{with} \quad |\eta_\ell| < 2.5 \\ N_j &> 1 \quad \text{with} \quad |\eta_j| < 3. \end{aligned} \quad (30)$$

In addition, we impose a set of basic cuts like $p_{Tj} > 20$ GeV and $p_{T\ell_1} > 200$ GeV at parton-level event simulation for

¹⁶As a further sanity check, we explicitly checked, using parton-level events, that $jj\ell\ell W$ does not contain any event with two intermediate W 's.

signal and backgrounds, where $P_{T\ell_1}$ means P_T of the hardest lepton. This rather hard cut ($p_{T\ell_1} > 200$ GeV) is designed to improve background statistics for data analysis. We reimpose such cuts, together with $\Delta R > 0.4$ for *all* possible pair of objects chosen, on objects (hardest two jets, and three leptons) in events that pass selection criteria of Eq. (30). We use p_T to evaluate hardness of the reconstructed objects.

Next, we discuss the way we reconstruct all invariant mass of each particles.

- (I) *Reconstruction of $p_{\nu z}$* : For a given choice of lepton, a candidate for ℓ_ν , the z-component of the neutrino's momentum can be obtained by requiring

$$M_W^2 = (p_{\ell_\nu} + p_\nu)^2 \quad (31)$$

where M_W is the mass of the SM W boson. Neutrino four-momentum p_ν^μ is constructed using the missing transverse momentum \cancel{E}_T to get p_x and p_y , and using masslessness of neutrino to compute E_ν from p_x , p_y and $p_{\nu z}$. The equation $(p_{\ell_\nu} + p_\nu)^2 = M_W^2$ is, therefore, a quadratic equation for $p_{\nu z}$. For each choice of ℓ_ν , we vary M_W , starting from the central value $M_W = 80$ GeV, by $M_W \pm \Delta$ in step size of 2 GeV until the quadratic equation finds real solution(s). We choose the maximum step to be $|\Delta| \leq 80$ so that the net mass ($M_W - \Delta$) is still positive semidefinite. We repeat this for all three leptons and keep all possible solutions, if they exist, for each lepton. If no solution exists after all three leptons, we drop the corresponding event. Notice that this procedure may end up giving more than one ℓ_ν , and furthermore, for each ℓ_ν , there may exist more than one solution for $p_{\nu z}$. In step II, we will specify the criteria by which we pick up one unique solution.

- (II) *Reconstruction of M_N* : Once the longitudinal component of the neutrino's momentum (or equivalently full p_ν^μ) is reconstructed by step I, for a given choice of ℓ_ν , we then determine ℓ_j and ℓ_W , the lepton combined with jet pair to reconstruct N (ℓ_j) and the lepton grouped with W (i.e., $\ell_\nu\nu$) to form N (ℓ_W), respectively, by minimizing

$$|M_{jj\ell_j} - M_{\ell_\nu\ell_W\nu}|. \quad (32)$$

This is motivated by the fact that both $\{jj\ell_j\}$ and $\{\ell_W\ell_\nu\nu\}$ comes from the decay of on-shell N . If correctly chosen, $\{jj\ell_j\}$ and $\{\ell_W\ell_\nu\nu\}$ should have invariant mass peaked at the same value (M_N), thus giving a small difference of two invariant masses. Whereas wrong pairing would tend to give much larger difference. We repeat this procedure for all three possible choices of ℓ_ν . Final decision is made

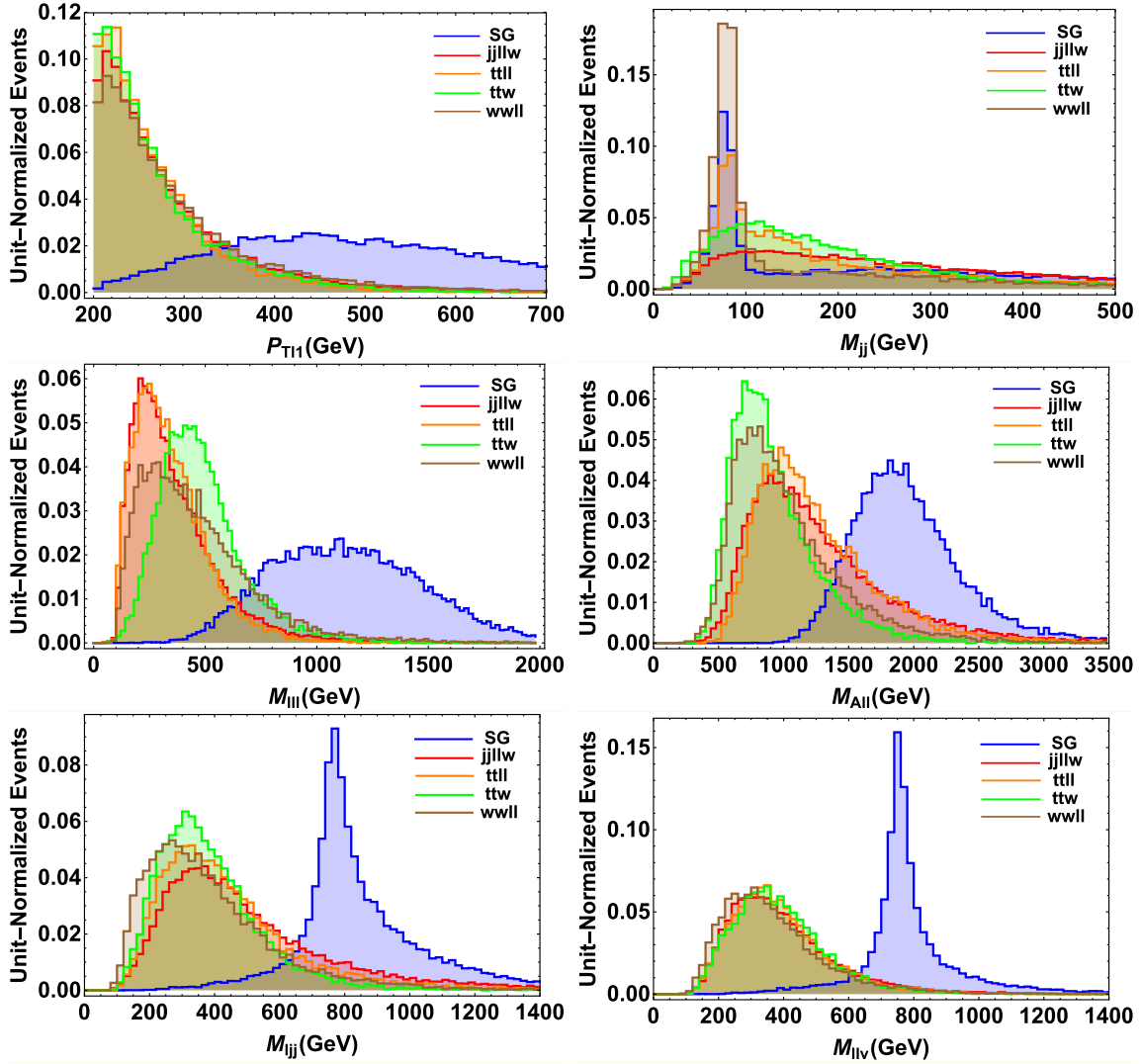


FIG. 2. ρ_{B-L} -Channel: p_{Tl} (top row, left), M_{jj} (top row, right), $M_{\ell\ell\ell}$ (middle row, left), M_{All} (middle row, right), $M_{\ell jj}$ (bottom row, left), $M_{\ell\ell v}$ (bottom row, right) for signal (solid blue) and backgrounds (solid, $jj\ell\ell W$ -red, $t\bar{t}\ell\ell$ -orange, $t\bar{t}W$ -green, $\ell\ell WW$ -brown)

for the combination $\{\ell_\nu, \ell_W, \ell_j\}$ that renders minimum value for Eq. (32).

In Fig. 2, we show distributions of various kinematic variables for signal and background events that pass selection criteria and basic cuts. Invariant mass variables are reconstructed following the prescription described above. The invariant mass distributions of signal events evidently show that our reconstruction prescription for invariant mass distributions is very successful. In particular, the distributions of M_{jj} (top row, right), M_{All} (middle row, right), $M_{\ell jj}$ (bottom row, left), $M_{\ell\ell v}$ (bottom row, right) are peaked at the position expected from the input values. Also shown is $M_{\ell\ell\ell}$ which reveals apparent separation between signal and background, supplying a very efficient cut for background reduction.

We employed a set of cuts to achieve significant $S/\sqrt{S+B}$ for our analysis. We provide the cut flows for signal and the major SM backgrounds in Table I. We

find that the ρ_{B-L} -channel may provide a sensitivity to discover N and ρ_{B-L} by $\sim 4.8\sigma$ with an integrated luminosity of $\mathcal{L} = 3000 \text{ fb}^{-1}$.

B. Composite $SU(2)_L$ doublet channel

As we briefly described at the beginning of the section, in the $SU(2)_L$ doublet channel, the singlet neutrino N is produced via the on-shell decay of composite $SU(2)_L$ lepton doublet, which, in turn, is produced by the on-shell decay of composite gauge boson ρ_{WL} . The relevant couplings for this channel are given in Sec. IV B. As already mentioned in the introduction, we want to emphasize that this production channel for singlet N is largely model independent, making its study important and well motivated. It is model independent in a sense that (i) the production is via couplings that always exist and sizable and (ii) it is largely independent of the representation of the singlet N under global symmetries of the composite sector.

TABLE I. Cut flows for signal and major background events in terms of their cross sections. The cross sections are in fb. The cross sections in the first row are such that for signal and for backgrounds, a set of minimum cuts ($p_{Tj} > 20$ GeV, $p_{T\ell} > 10$ GeV, $\Delta R > 0.4$ etc.) at event generation level are applied to avoid IR-divergence. In the second row, the same basic cuts ($P_{Tj} > 20$ GeV, $P_{T\ell_1} > 200$ GeV, $|\eta_j| < 3$, $|\eta_\ell| < 2.5$, $\Delta R \geq 0.4$ for *all* pairs of objects) are reimposed on both signal and backgrounds events.

Cuts	Signal	$jj\ell\ell W$	$t\bar{t}\ell\ell$	$t\bar{t}W$	$WW\ell\ell$
Generator-level cuts	4.88×10^{-2}	1.39×10^2	7.26	2.50	7.61×10^{-1}
$N_\ell \geq 3$, $N_j \geq 2$ with basic cuts	2.88×10^{-2}	4.26	2.14×10^{-1}	5.32×10^{-2}	3.95×10^{-2}
$M_{\ell\ell\ell} \in [600, \infty]$ GeV	2.76×10^{-2}	4.49×10^{-1}	1.62×10^{-2}	1.16×10^{-2}	9.01×10^{-3}
$M_{\ell jj} \in [700, 1200]$ GeV	2.14×10^{-2}	1.01×10^{-1}	3.21×10^{-3}	1.08×10^{-3}	1.03×10^{-3}
$M_{\ell\ell\nu} \in [700, 1000]$ GeV	1.79×10^{-2}	4.51×10^{-2}	1.39×10^{-3}	3.48×10^{-4}	3.97×10^{-4}
$p_{T\ell_1} \in [300, \infty]$ GeV	1.68×10^{-2}	3.12×10^{-2}	5.99×10^{-4}	2.04×10^{-4}	2.99×10^{-4}
$M_{jj} \in [0, 500]$ GeV	1.46×10^{-2}	1.40×10^{-2}	3.42×10^{-4}	1.39×10^{-4}	1.96×10^{-4}
$M_{\text{All}} \in [0, 2500]$ GeV	1.40×10^{-2}	1.13×10^{-2}	2.78×10^{-4}	1.39×10^{-4}	1.70×10^{-4}
S/B	1.18
$S/\sqrt{S+B}$ ($\mathcal{L} = 300 \text{ fb}^{-1}$)	1.51
$S/\sqrt{S+B}$ ($\mathcal{L} = 3000 \text{ fb}^{-1}$)	4.77

In order to see this, first notice that composite $SU(2)_L$ doublet lepton (ν^h, ℓ^h) is present in any composite Higgs models (or its five-dimensional dual Randall-Sundrum model) whenever leptons exhibit partial-compositeness. The same is true for composite ρ_{W_L} , even when we do not extend our electroweak symmetry to LR symmetric version. Then, the coupling between composite lepton doublet and composite ρ_{W_L} exists by global symmetry of the composite sector (or bulk gauge symmetry in the five-dimensional picture). Moreover, the coupling of composite lepton doublet with singlet N and Higgs must be present simply because this is the coupling that generates the mass of the SM neutrino mass by the seesaw mechanism. Finally, the production of the composite ρ_{W_L} can be achieved simply by the composite-elementary mixing between composite ρ_{W_L} and its elementary partner, which then couples to quarks inside the proton. Therefore, we see that, provided masses of composite particles are light enough and are such that their on-shell decays are kinematically allowed (hence resonance-enhancement), we can ensure enough rate for a signal process at the LHC, which is, again, independent of specifics of the model.

Once produced, singlet N decays (using Yuakwa coupling) to a lepton and W boson. As can be seen from the Feynman diagram shown in Fig. 1, the signal process contains three W 's, two leptons and one either H or Z . We choose to study the case when two of the three W 's decay hadronically producing four jets and one leptonically producing lepton + MET. Since there are three W 's, there will be three ways that this occurs. H/Z decays to $b\bar{b}$. The final state, therefore, consists of $4j2b3\ell + \text{MET}$.

There are several invariant mass variables which are, as we will explain momentarily, all fully reconstructible and crucial in revealing signal process out of large SM backgrounds. Those invariant mass variables include, $M_{\rho_{W_L}}$, M_{ℓ^h} , M_N , M_H , and M_W , where $M_{\rho_{W_L}}$ is the reconstructed

invariant mass of composite ρ_{W_L} , and similarly for others as is evident from their names. Due to $SU(2)_L$ symmetry, ℓ^h and its partner ν^h are degenerate and we use M_{ℓ^h} to denote the mass for $SU(2)_L$ composite doublet lepton. Because of combinatorics in W decays, i.e., two $W \rightarrow jj$ and one $W \rightarrow \ell\bar{\nu}$, the identification of a set of particles coming from decay of N (and similarly for ℓ^h/ν^h) is not fixed. For example, in case when W from direct decay of ℓ^h decays leptonically, both N will decay eventually into ℓjj . On the other hand, if leptonically decaying W is from N , then one N decays into ℓjj , but the other into $\ell\ell + \text{MET}$. For this reason, the sets of particles that reconstruct invariant masses of composite particles will vary event by event and this explains why we denote invariant mass variables by the name of the corresponding particle, e.g. M_N , instead of its daughter SM particles, say, $M_{jj\ell}$ like in our earlier paper [6]. If successfully reconstructed, the signal distributions of these variables will be peaked at $M_{\rho_{W_L}} = 2.5$ TeV, $M_{\ell^h} = 1$ TeV, $M_N = 500$ GeV, $M_H \approx 125$ GeV, and $M_W \approx 80$ GeV, respectively.

There are several SM backgrounds we need to consider, and we now describe them one by one in order of significance.

- (1) $t\bar{t}jj\ell\ell$: The relevant process is $pp \rightarrow t\bar{t}jj\ell\ell$, followed by subsequent decay of $t \rightarrow bW$, and similarly for \bar{t} . One of the W 's from the top decays hadronically, and the other decays leptonically. In principle, one may consider even more inclusive background, $pp \rightarrow 4j2b\ell\ell W$, $W \rightarrow \ell\nu$, where we need W to obtain the third lepton from its decay without violating lepton number conservation. However, we found that generating such an inclusive process in a Monte Carlo event generator like MG5AMC@NLO is rather impractical mainly due to proliferation of Feynman diagrams. As the next best plan, we decided to consider the above

described process, which we think is inclusive enough to capture relevant SM backgrounds for our study. The kind of processes that we are missing by not considering the most inclusive one will be those with soft pure QCD jets (including b -jets). In order to strengthen our argument, therefore, we impose hard p_T cut on the hardest jet, i.e., $p_{T_{j_1}} > 100$ GeV, where j_1 is the jet with the largest p_T . This is the background with largest cross section. As already mentioned, we simulate processes to leading order in QCD and QED couplings. For later purpose, we mention that the coupling order relevant for this background is QCD = 4 and QED = 6. Background reduction will be achieved by means of a combination of various invariant mass cuts. Since two leptons (those not from W decay) in this background will mostly come from near-on-shell decay of Z boson, the distribution of the dilepton invariant mass, $M_{\ell\ell}$, will be sharply peaked at the mass of the Z boson, M_Z . However, since two leptons in the signal process arise from direct decay of N , they do not reconstruct M_Z . Consequently, the condition $M_{\ell\ell} \neq M_Z$ provides a very efficient cut. Other useful cuts are $M_{\rho_{W_L}}$, $M_{\ell\ell\ell}$ and M_N cuts. $P_{T\ell_1}$ will also be very useful.

- (2) $t\bar{t}\bar{t}$: The relevant process is $pp \rightarrow t\bar{t}\bar{t}$, where each top decays to $t \rightarrow bW$. Out of four W 's produced in top decays, one decays hadronically and the remaining three decays leptonically, resulting in three leptons and MET. In order to pass the selection criteria, however, two of the four b 's must be tagged as b -jets and the other two must be un-tagged as regular two jets, leading to a large reduction of the background. In addition, $M_{\rho_{W_L}}$ and M_N cuts will be particularly efficient for reduction of this background.
- (3) $t\bar{t}\ell\ell W$: The relevant process is $pp \rightarrow t\bar{t}\ell\ell W$, with subsequent decay of tops to $b + W$. Two W 's decay hadronically rendering four jets and the third one leptonically. Similarly to $t\bar{t}jj\ell\ell$, the lepton pair comes mostly from decay of on-shell Z (and off-shell photon). Therefore, $M_{\ell\ell} \neq M_Z$ will remove most of this background. Several other invariant mass cuts will be useful. One may readily notice that if one of the hadronically decaying W 's is the one not from the top, this process has precisely the same final state (before top decays) as $t\bar{t}jj\ell\ell$ and one would worry about double counting issue. However, we claim there is no such issue. The resolution is that we are working at leading order in QCD and QED couplings, and the leading coupling order that gives rise to this process is QCD = 2 and QED = 8, while that for $t\bar{t}jj\ell\ell$ background was QCD = 4 and QED = 6, hence not captured there.
- (4) *others*: There are several other processes that can contribute to SM backgrounds. One of them is

$pp \rightarrow t\bar{t}WWW$, with three W 's decaying leptonically and the other two hadronically. Another is $pp \rightarrow H/Z\ell\ell WWW$, with two W 's decay hadronically and the third one leptonically. However, these processes (i) require high QED coupling order (QED = 10) leading to parametric suppression compared to the above three cases and (ii) multiplicity of the final state is equal or even greater than the above major processes (i.e., either comparable or additional phase space suppression). As a result, the cross section of these processes is much smaller, leading to at most $O(1)$ events before selection criteria and hence no effects in the final results. For this reason, we do not consider these backgrounds explicitly.

Defining N_ℓ , N_b , and N_j as the number of isolated leptons, b -tagged jets and non- b -tagged jets, respectively, we select events using the following selection criteria:

$$\begin{aligned} N_\ell &> 2 && \text{with } |\eta_\ell| < 2.5 \\ N_b &> 1 && \text{with } |\eta_b| < 3 \\ N_j &> 3 && \text{with } |\eta_j| < 3. \end{aligned} \quad (33)$$

In addition, we impose a set of basic cuts $p_{T_j}/p_{T_b} > 20$ GeV, $p_{T\ell} > 10$ GeV, $\Delta R > 0.4$, and so on at parton-level event simulation, partly to avoid possible IR-divergence issues for background simulations. We reimpose such cuts on objects (hardest four jets, two b -jets, and three leptons) in events that pass selection criteria of Eq. (33). We use p_T to evaluate hardness of the reconstructed objects. We also explicitly impose $\Delta R > 0.4$ for *all* possible pair of objects chosen out of above selected objects.

Now, we describe how we reconstruct invariant masses of each particle.

- (I) *Reconstruction of p_{ν_z}* : The first step towards the reconstruction of all resonance peaks is the reconstruction of longitudinal component of neutrino momentum, p_{ν_z} . There are three leptons in the process and each of these can potentially pair with neutrino to form W . For each choice of lepton, call it ℓ_ν , we first solve the quadratic equation $(p_{\ell_\nu} + p_\nu)^2 = M_W^2$, where p_ν^μ is formed using the missing transverse momentum \cancel{E}_T to get p_x and p_y and using masslessness to compute E_ν from p_x , p_y and p_{ν_z} . The equation $(p_{\ell_\nu} + p_\nu)^2 = M_W^2$ is, therefore, a quadratic equation for p_{ν_z} . For each choice of ℓ_ν , we vary M_W , starting from the central value $M_W = 80$ GeV, by $M_W \pm \Delta$ in step size of 1 GeV until the quadratic equation finds real solution(s). We choose maximum step to be $|\Delta| \leq 80$ so that the net mass $(M_W - \Delta)$ is still positive semidefinite. We repeat this for all three leptons and choose the lepton (s) with minimum Δ ; i.e., we choose the lepton(s) which reconstruct the p_{ν_z} such that the computed W mass, i.e., $(p_{\ell_\nu} + p_\nu)^2$, is closest to the central value

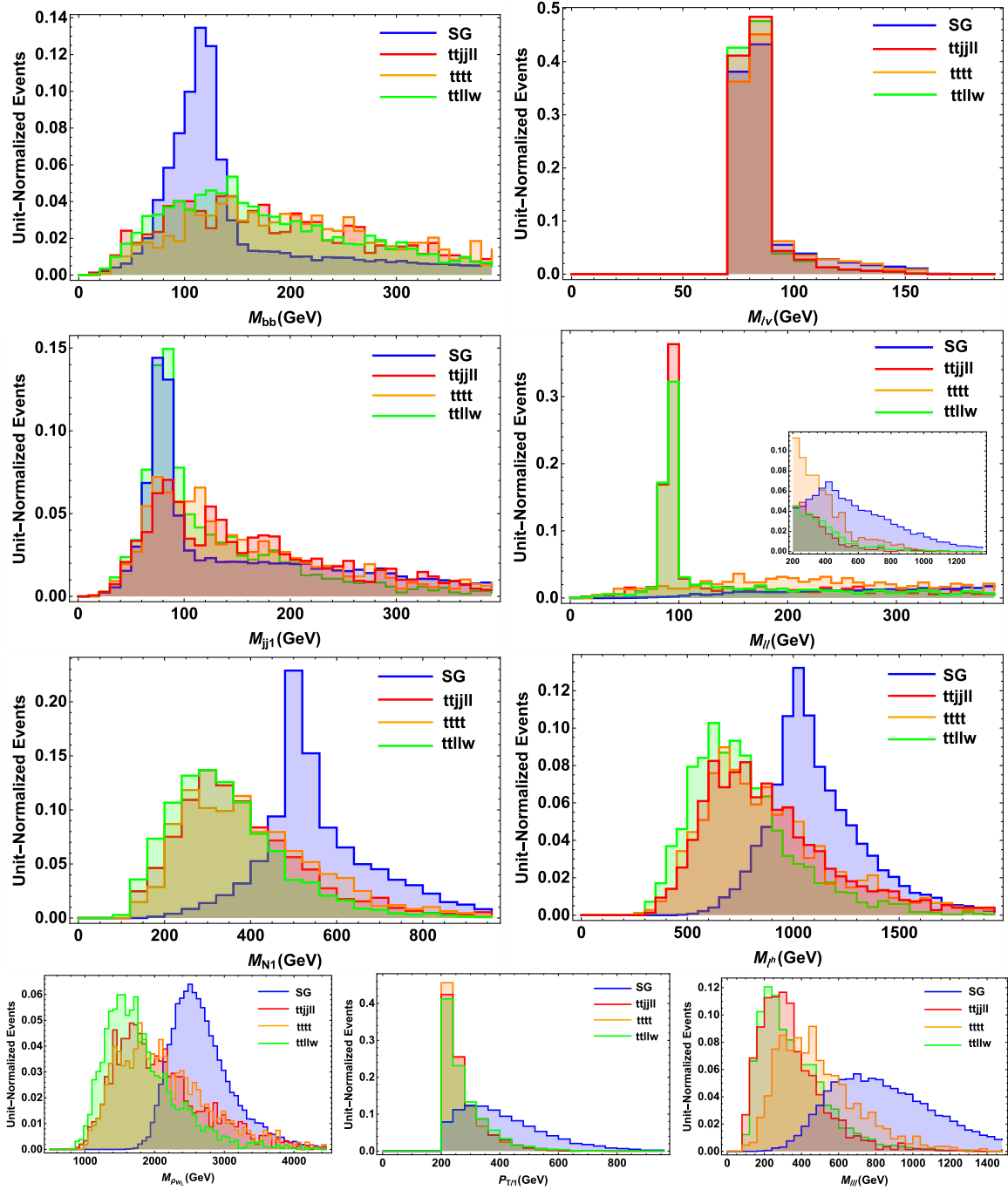


FIG. 3. Composite $SU(2)_L$ doublet-channel: Distributions of variables: M_{bb} (top row, left), $M_{\ell\nu}$ (top row, right), M_{jj_1} (second row, left), $M_{\ell\ell}$ (second row, right), M_{N_1} (third row, left), M_{ρ^h} (third row, right), $M_{\rho_{W_L}}$ (bottom row, left), $P_{T\ell_1}$ (bottom row, middle) and $M_{\ell\ell\ell}$ (bottom row, right) for signal (solid blue) and backgrounds (solid, $t\bar{t}j\ell\ell$ -red, $t\bar{t}t$ -orange, $t\bar{t}\ell\ell W$ -green)

of 80 GeV. If no solution exists after all three leptons, we drop the corresponding event. Notice that this procedure may end up giving more than one ℓ_ν with the same deviation factor Δ , and furthermore, for each ℓ_ν , there may exist more than one solution for $p_{\nu z}$. In step III, we will specify the criteria by which we pick one unique solution.

(II) *Jet pairing*: In signal process, four jets arise from decay of two W 's. So, we determine jet pairing (i.e., figuring out jet pair from each W boson) as follows. We consider all possible jet pairings. Then for each jet pairing jj_1, jj_2 , e.g. $jj_1 \equiv (j_1, j_2)$ and $jj_2 \equiv (j_3, j_4)$, we compute the following quantity (" L^2 -distance" function):

TABLE II. Cut flows for signal and major background events in terms their cross sections. The cross sections are in fb. The numbers in the first row (“generator-level cuts”) are cross sections obtained with basic cuts at the generation level to avoid divergence (for both signal and backgrounds). In the second row, the same basic cuts are reimposed, in addition to $\Delta R > 0.4$ for all selected object pairs, to both signal and background events along with multiplicity requirements for b -jet, non- b -jet, and leptons.

Cuts	Signal	$t\bar{t}jj\ell\ell$	$t\bar{t}t\bar{t}$	$t\bar{t}\ell\ell W$
Generator-level cuts	7.87×10^{-2}	6.86	3.45×10^{-1}	1.65×10^{-2}
$N_\ell > 2, N_j > 3, N_b > 1$ with basic cuts	7.62×10^{-3}	4.59×10^{-1}	1.38×10^{-2}	1.22×10^{-3}
$p_{T\ell_1} \geq 200$ GeV, $p_{Tj_1} \geq 100$ GeV	6.94×10^{-3}	6.61×10^{-2}	1.49×10^{-3}	1.66×10^{-4}
$M_{\ell\ell} \in [110, \infty]$ GeV	6.87×10^{-3}	2.46×10^{-2}	1.33×10^{-3}	7.19×10^{-5}
$M_{\rho_{wL}} \in [2000, \infty]$ GeV	6.73×10^{-3}	1.14×10^{-2}	7.11×10^{-4}	2.00×10^{-5}
$M_{\ell\ell\ell} \in [450, \infty]$ GeV	6.42×10^{-3}	3.33×10^{-3}	4.72×10^{-4}	9.04×10^{-6}
$M_{N_1} \in [400, \infty]$ GeV	5.97×10^{-3}	2.17×10^{-3}	3.13×10^{-4}	5.08×10^{-6}
S/B	2.40
$S/\sqrt{S+B}$ ($\mathcal{L} = 300$ fb $^{-1}$)	1.12
$S/\sqrt{S+B}$ ($\mathcal{L} = 3000$ fb $^{-1}$)	3.56

$$d_{jj} = (M_{jj_1} - M_W)^2 + (M_{jj_2} - M_W)^2. \quad (34)$$

We choose the pairing with minimum d_{jj} .

- (III) *Reconstruction of M_N* : Next, we want to determine sets of particles coming from N decay. Let’s call each set N_1 and N_2 . Depending on how W from N decays (to jj v.s. to $\ell\nu$), the final set can be either ℓjj or $\ell\ell\nu\nu$, where ℓ_ν is the lepton that pairs with neutrino to form a W .¹⁷ Thus, there are two possibilities we need to consider: (i) both N decay to ℓjj and (ii) one to ℓjj and the other to $\ell\ell\nu\nu$. In step I, we build a set that collects lepton(s), ℓ_ν , and corresponding $p_{\nu z}$ ’s. For each such ℓ_ν , we can now simply figure out the remaining two leptons. In case (i), these will be paired with jet-pairs to form N . Using the results of step II, there are only two combinations of $\{\ell jj\}$. For each such constructed set N_1 and N_2 , we compute the distance function:

$$d_N = (M_{N_1} - M_{N_2})^2. \quad (35)$$

In case (ii), on the other hand, one lepton goes with jet pair (two choices, jj_1 or jj_2), and the other lepton groups with ℓ_ν . For latter, we consider all choices of ℓ_ν and corresponding $p_{\nu z}$ that pass the criteria of step I. Similarly to the case (i), for each set N_1 and N_2 , we compute the distance Eq. (35). For the set containing ℓ_ν , we use fully reconstructed p_ν^μ to calculate invariant mass. At the end, we choose the set N_1 and N_2 with minimum d_N .¹⁸

¹⁷Recall that, at this stage, there still can be several ℓ_ν choices, together with the possibility of two $p_{\nu z}$ solutions. Unique solutions will be determined by the determination of the set N_1 and N_2 .

¹⁸At this stage, we have made a unique decision for ℓ_ν and corresponding $p_{\nu z}$, unless numerical coincidence happens by complete accident. Considering the amount of significant figures in the data, this will be very unlikely and indeed, we have not found one case.

- (IV) *Reconstruction of M_{ℓ^h}* : Finally, we determine the set of particles from the decay of ℓ^h and ν^h . Thanks to $SU(2)_L$ symmetry, we know that $M_{\ell^h} = M_{\nu^h}$ and using this we consider all the combinations and choose the one that minimizes

$$d_{\ell^h} = (M_{\ell^h} - M_{\nu^h})^2. \quad (36)$$

To be more specific, in case (i) of step III, both N_1 and N_2 are set of the form $\{\ell jj\}$. One of these will combine with $\{\ell_\nu\nu\}$ to render ℓ^h and the other $\{b\bar{b}\}$ to make up ν^h . In case (ii) of step III, on the other hand, one of N_1/N_2 will group with $\{jj\}$ to give ℓ^h and the other with $\{b\bar{b}\}$ to construct ν^h . We simply consider all these combinations and compute invariant masses for each choice. At the end, we simply choose the one with minimum d_{ℓ^h} .

Figure 3 shows distributions of various variables for signal and backgrounds constructed using the above described procedure for events that pass selection criteria and basic cuts. Remarkably, as seen clearly from signal distributions, our reconstruction prescription for invariant mass distributions is very successful. For example, the distributions of $M_{\ell_\nu\nu}$ (top row, right), M_{jj_1} (second row, left), M_{N_1} (third row, left), M_{ℓ^h} (third row, right) and $M_{\rho_{wL}}$ (bottom row, left) reveal a well-developed resonance peak with the position matched well with the input values. Considering the number of intermediate states in the process (see Fig. 1) and high multiplicity of leptons and jets, such outcome is rather surprising. We also present distributions for M_{bb} , which is peaked at around M_H for signal, while smoother for backgrounds, $M_{\ell\ell}$ for non- ℓ_ν two leptons, which is sharply peaked at M_Z for $t\bar{t}jj\ell\ell$ and $t\bar{t}\ell\ell W$, $M_{\ell\ell\ell}$, which provides yet another very strong cut, and p_T of the hardest lepton. All of these features are as expected. In particular, we see that $M_{\ell\ell}$, M_{N_1} , M_{ℓ^h} , $M_{\rho_{wL}}$, $M_{\ell\ell\ell}$, and $p_{T\ell_1}$ draw sharp distinctions between signal and

backgrounds, providing strong cuts to attain significant backgrounds reduction.

We performed analysis by applying a series of these kinematic cuts. We provide the cut flows for the signal and the major SM backgrounds in Table II. We find that the composite $SU(2)_L$ doublet channel may provide a sensitivity to uncover warped seesaw nature in largely model-independent way by $\sim 3.6\sigma$ with an integrated luminosity of $\mathcal{L} = 3000 \text{ fb}^{-1}$. Notice, however, that in this analysis, we have selected events with W 's decaying into regular jet pair. Given the heaviness of ℓ^h and N from which it is produced, we expect that including boosted W events with fat jet final state will provide a significant increase in the signal rate, and hence larger significance.

VI. CONCLUSIONS AND OUTLOOK

In [1], we argued that (i) warped seesaw is a natural implementation of the essence of the original seesaw paradigm and that (ii) composite TeV-mass singlet neutrinos play a crucial role in SM neutrino mass generation. Then, in a previous paper on LHC signals [6], we studied production of these singlet neutrinos using mechanisms analogous to four-dimensional LR models, but with important and interesting differences. In this paper, we considered production of singlet neutrinos from decays of particles beyond four-dimensional LR models.

In the earlier work [6], we considered models where the composite sector has an extended EW global symmetry, $SU(2)_L \times SU(2)_R \times U(1)_X$. Such a left-right symmetric structure is motivated in the context of a composite Higgs or five-dimensional warped framework by consistency with EW precision tests. In addition, we assumed degeneracy of masses of composite gauge bosons. The couplings of the light quarks to composite W_R^\pm and Z' required for their production, and hence that of the singlet neutrino via their decay, then were established in part by EWSB induced mixing between composite W_L^\pm and W_R^\pm (and similarly for neutral gauge bosons); mass degeneracy was crucial for enhancing the size of this mixing. This effect has to be combined with the elementary-composite gauge boson mixing. In this case, composite W_R^\pm couples to left-handed quarks. This is to be contrasted with the usual four-dimensional LR models, where W_R^\pm directly couples to right-handed quarks. For more detail, see [6].

In this paper, we show that taking the same model (with $X = \frac{1}{2}(B - L)$) and yet exploring different regions of parameter space can result in different production channels for singlet neutrino with remarkable qualitative differences. In the first part of this paper, as a more general consideration, we choose composite W_R^3 and $(B - L)$ to be nondegenerate. At first sight, such nondegeneracy seems to suppress Z' signal. Remarkably, however, there is actually an emergence of another new neutral channel signal. The point is that composite W_R^3 and $(B - L)$ are now separately

mass eigenstates; i.e., their ‘‘hypercharge’’ and Z' combinations appearing in the degenerate case are nowhere close to being mass eigenstates. Both these spin-1 composites, W_R^3 and $(B - L)$, mix with elementary hypercharge. This effect suffices to couple them to light quarks, without need of EWSB in addition. This is to be contrasted with EWSB mixing also being required for coupling composite W_R^\pm/Z' to light quarks in this model. In fact, in this sense, the status of composite W_R^3 and $(B - L)$ in the nondegenerate case is rather similar to composite W_L 's. Moreover, the singlet neutrino couples to both W_R^3 and $(B - L)$. So, provided the lighter of the two is light enough, it can be produced at the LHC with a significant rate, and thus, producing the singlet neutrino through its decay. We show that 4.8σ signal can be achieved with 3000 fb^{-1} luminosity for 2 TeV composite $(B - L)$ gauge boson and 750 GeV singlet neutrino.

It is worth emphasizing a distinction from four-dimensional LR here: there is no analogue in that model of W_R^3 and $(B - L)$ as separate mass eigenstates because only their heavy combination, Z' , corresponding to the broken gauge symmetry, is in play, orthogonal one being SM hypercharge. Even if we focus only on the lightest state in the composite/five-dimensional case above, note that its couplings are different from that of Z' (whether the four-dimensional LR model or the degenerate case in five dimensions), allowing disambiguation between the two. We studied this interesting channel here, taking as an illustration the case of a composite $(B - L)$ being lighter than W_R^3 .

Then, in the second part of this paper, we demonstrated the possibility of obtaining a signal from singlet neutrino which is independent of gauge couplings of singlet neutrino, e.g. W_R and $(B - L)$ are too heavy to be relevant or neutrino is a singlet of $SU(2)_R \times U(1)_X$ also. Namely, decays of composite $SU(2)_L$ doublet leptons can produce the composite singlet neutrino, accompanied by Higgs boson/longitudinal W/Z . Note that the associated coupling originates purely in the composite sector, thus is sizable. It is also robust feature of this model in the sense that the same coupling is also involved in the Dirac mass term part of the neutrino mass seesaw. At the same time, the composite doublet leptons are ‘‘guaranteed’’ to have significant couplings to light quarks and thus substantial production rate at the LHC, via elementary and composite *left-handed* W . On the other hand, for the singlet neutrino signal either from decays of W_R^3/X as studied in the first part of this paper, or from decays of Z'/W_R^\pm studied in the previous paper, we need certain representations of the SM singlet neutrino under the extended EW symmetry. In addition, degeneracy of spin-1 states was invoked in the earlier work for the purpose of obtaining a significant coupling of light quarks to Z'/W_R^\pm for their production. We show that $\sim 4\sigma$ signal can be achieved with 3000 fb^{-1} luminosity for the following spectrum: 2.5 TeV composite

ρ_{W_L} , 1 TeV composite $SU(2)_L$ doublet lepton and 500 GeV singlet neutrino.

Even though we studied singlet neutrino production via particles not featured in the usual LR models, either W_R^3 or $(B-L)$ or composite lepton doublet, the resulting signals from these have appreciable similarities to the standard ones from usual LR W_R^\pm/Z' . Specifically, we get multi-leptons, in association with SM gauge bosons generically. However, there do exist important differences as well. In particular, we do get extra Higgs/ Z in pair-production of singlet neutrinos from decay of $SU(2)_L$ doublet composite leptons as compared to Z' . The different detailed topologies involved also imply that the kinematic distributions of common part of the final state, say, leptons, can also be rather distinct.

In fact, in part of the parameter space, the situation might be even better as follows. For example, in the case of composite $(B-L)$ being lighter than all others, merely comparing the other decay channels of this spin-1 state and the production/decay of other heavy resonances can readily reveal nature of the underlying model. In other words, we are concerned here with more qualitative differences between the signals for the various models, which can afford model characterization even before any detailed analyses are available (for example, related to extra bosons in final state mentioned above). This is because there is no charged channel at all to be seen at the LHC in this example. Furthermore, it is interesting that there are no diboson decays for the composite $(B-L)$. Both these features are in sharp contrast to the four-dimensional LR models, where Z' (neutral channel) obviously has a charged counterpart, i.e., W_R^\pm . In fact, W_R^\pm is typically easier to discover because it is lighter. Moreover, both Z' and W_R^\pm decay to dibosons with similar rate as singlet neutrinos. Once again, such dramatic differences in the signal afford discrimination between the two models essentially right at the time of discovery.

Relatedly, even within the context of the composite/five-dimensional model, the singlet neutrino decay channel can have significant implications for discovery of EW spin-1 composites. Regardless of the seesaw model, top/EW gauge boson/Higgs are the usual discovery channel for such composite EW gauge bosons. However, the diboson is not available for composite $(B-L)$ as already mentioned above. Composite $(B-L)$ does decay to top quarks, but that might be diluted by the new, singlet neutrino channel. Amazingly then, for this “new” particle, the decay into singlet neutrino might instead be the way to go even for discovery. Whereas, in the earlier case of degenerate composites, W_R^\pm or Z' , the diboson or top channels were available for discovery, and the singlet neutrino channel was more for testing the seesaw itself.

Finally, a word about possible future work related to this framework. One direction would be to explore LHC signals in other regions of parameter space, for example, the (more

challenging) case of heavier composite singlet/doublet neutrinos/charged leptons such that spin-1 composites can not decay into pair of them. Also, the idea of a 100 TeV hadron collider is being discussed a lot: it will be interesting to determine the reach of this proposed collider as far as the singlet neutrino is concerned.

It is worth noting that there are beyond-high-energy collider phenomenological aspects to TeV-mass singlets. For example, we plan to study leptogenesis in this setup, which could occur at $O(\text{TeV})$ temperatures, cf., at super-high scales for original seesaw. In addition, we have constraints from lepton flavor-violation arising from virtual effects of \lesssim TeV-mass singlet neutrino/composite charged leptons. We assumed some sort of flavor symmetries here in order to be consistent with those precision tests. It might be worthwhile exploring these in more detail. The bottom line is that our work has opened up new avenues for studies of the natural seesaw idea for neutrino mass within the framework of a warped extra dimension/composite Higgs, whether at the LHC or other non-(high-energy) colliders.

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APPENDIX: ONE ELEMENTARY-TWO COMPOSITE GAUGE BOSON MIXING

As mentioned in the main text, in the two-site basis, we have elementary hypercharge gauge bosons (denoted by B_{elem}) mixing with composite X (labeled $\tilde{\rho}_X$) and composite W_R^3 (called $\tilde{\rho}_{W_R^3}$). Their gauge couplings (to other particles in the corresponding sector only) are g_{elem} , g_\star , and G_\star , respectively. The diagonal mass terms for $\tilde{\rho}_X$ and $\tilde{\rho}_{W_R^3}$ are denoted by m_\star and M_\star . In addition, there are mass mixing terms between elementary and composites, along with the appropriate diagonal mass term for the elementary gauge boson. Note that, even before EWSB, these are then the gauge/weak eigenstates, i.e., not mass eigenstates.

The mass matrix has the form

$$M^2 = \begin{pmatrix} \frac{g_{\text{elem}}^2}{g_\star^2} m_\star^2 + \frac{g_{\text{elem}}^2}{G_\star^2} M_\star^2 & -\frac{g_{\text{elem}}}{g_\star} m_\star^2 & -\frac{g_{\text{elem}}}{G_\star} M_\star^2 \\ -\frac{g_{\text{elem}}}{g_\star} m_\star^2 & m_\star^2 & 0 \\ -\frac{g_{\text{elem}}}{G_\star} M_\star^2 & 0 & M_\star^2 \end{pmatrix}. \quad (\text{A1})$$

We can check that the above matrix has a vanishing determinant, which is as expected based on it being made up of the “usual” two elementary-composite 2×2 blocks. This results in one field with zero eigenvalue, corresponding to the SM hypercharge gauge boson.

This matrix is diagonalized by the rotation¹⁹

$$U^\dagger M^2 U = M_{\text{diag}}^2, \quad (\text{A2})$$

where we have a product of three effectively 2×2 rotations:

$$U = U_{12} U_{13} U_{23} \quad (\text{A3})$$

with

$$U_{12} = \begin{pmatrix} c & -s & 0 \\ s & c & 0 \\ 0 & 0 & 1 \end{pmatrix}$$

$$U_{13} = \begin{pmatrix} C & 0 & -S \\ 0 & 1 & 0 \\ S & 0 & C \end{pmatrix}$$

$$U_{23} = \begin{pmatrix} 1 & 0 & 0 \\ 0 & c_\star & -s_\star \\ 0 & s_\star & c_\star \end{pmatrix}. \quad (\text{A4})$$

Here, the sines of the first two mixing angles are given by

$$s = \frac{g_{\text{elem}}}{\sqrt{g_{\text{elem}}^2 + g_\star^2}}$$

$$S = \frac{g_{\text{elem}} C}{\sqrt{g_{\text{elem}}^2 c^2 + G_\star^2}}, \quad (\text{A5})$$

and c and C are the corresponding cosines:

$$c = \sqrt{1 - s^2} \left(= \frac{g_\star}{\sqrt{g_{\text{elem}}^2 + g_\star^2}} \right)$$

$$C = \sqrt{1 - S^2} \left(= \frac{G_\star}{\sqrt{g_{\text{elem}}^2 c^2 + G_\star^2}} \right). \quad (\text{A6})$$

Finally, c_\star and s_\star are cosine/sine of last/2-3 rotation angle (θ_\star):

$$\tan 2\theta_\star = \frac{-2M_\star^2 \frac{g_{\text{elem}}}{G_\star} s(C + cS \frac{g_{\text{elem}}}{G_\star})}{M_\star^2(1 + c^2 \frac{g_{\text{elem}}^2}{G_\star^2} - s^2 \frac{g_{\text{elem}}^2}{G_\star^2}) - m_\star^2(1 + \frac{g_{\text{elem}}^2}{g_\star^2})}. \quad (\text{A7})$$

¹⁹In general, this would be unitary, for example, for mass matrix of charged gauge bosons so that we keep that notation.

The relation between the mass (denoted by $B^{(0)}$, ρ_X and $\rho_{W_R^3}$ ²⁰) and weak/gauge eigenstate bases is given by

$$\begin{pmatrix} B^{(0)} \\ \rho_X \\ \rho_{W_R^3} \end{pmatrix} = U_{23}^\dagger U_{13}^\dagger U_{12}^\dagger \begin{pmatrix} B_{\text{elem}} \\ \tilde{\rho}_X \\ \tilde{\rho}_{W_R^3} \end{pmatrix}, \quad (\text{A8})$$

so that we have the massless and two heavy eigenstates:

$$B^{(0)} = cCB_{\text{elem}} + Cs\tilde{\rho}_X + S\tilde{\rho}_{W_R^3}$$

$$\rho_X = -(sc_\star + s_\star Sc)B_{\text{elem}} + (cc_\star - s_\star Ss)\tilde{\rho}_X + s_\star C\tilde{\rho}_{W_R^3}$$

$$\rho_{W_R^3} = (s_\star s - c_\star Sc)B_{\text{elem}} - (s_\star c + c_\star Ss)\tilde{\rho}_X + c_\star C\tilde{\rho}_{W_R^3}. \quad (\text{A9})$$

The elementary fermions have a charge Q_Y under B^{elem} only so that they couple to $B^{(0)}$ with strength $Q_Y g_{\text{elem}} cC$. Thus, if g_Y denotes the SM hypercharge gauge coupling, we can identify

$$g_Y = cCg_{\text{elem}} \approx g_{\text{elem}}, \text{ assuming } g_{\text{elem}} \ll g_\star, G_\star. \quad (\text{A10})$$

Whereas, composite fermions couple only to $\tilde{\rho}_X$ and $\tilde{\rho}_{W_R^3}$ with charges denoted by Q_X and $Q_{W_R^3}$, respectively. One can check that their coupling to $B^{(0)}$ is given by $g_Y(Q_{W_R^3} + Q_X)$, where we have $Q_Y = Q_{W_R^3} + Q_X$. In particular, in the first model that we study, we identify $X = \frac{1}{2}(B - L)$, with the singlet neutrino $N^{(1)}$ having $Q_{W_R^3} = \frac{1}{2}$ and $Q_X = -\frac{1}{2}$ (thus $Q_Y = 0$).

Assuming

$$M_\star \gg m_\star$$

$$g_{\text{elem}} \ll g_\star, G_\star \text{ i.e., } s, S \ll 1, \quad (\text{A11})$$

we get

$$s_\star \approx -sS \left(\approx \frac{g_{\text{elem}} g_{\text{elem}}}{g_\star G_\star} \right); \quad (\text{A12})$$

i.e., in this case, the composite $W - X$ mixing angle is given roughly by the product of $B_{\text{elem}} - \tilde{\rho}_X$ and $B_{\text{elem}} - \tilde{\rho}_{W_R^3}$ mixing angles (second order in elementary-composite mixing, thus negligible for purpose of LHC signals). Actually,

$$M_\star^2 - m_\star^2 \sim O(M_\star^2), \text{ for example } M_\star \sim 2m_\star, \quad (\text{A13})$$

suffices to give the above negligibly small size of the $W - X$ mixing angle, i.e.,

²⁰These models can be identified with zero mode and KK models of the five-dimensional model. This choice of notation for the heavy mass eigenstates will be explained below.

$$s_\star \sim sS \left(\sim \frac{g_{\text{elem}}}{g_\star} \frac{g_{\text{elem}}}{G_\star} \right). \quad (\text{A14})$$

Similarly, in the opposite limit, composite $(B-L)$ is heavier than W_R^3 , and we find that the composite $X-W$ mixing angle is doubly suppressed as above.

Thus, generically, the two composite masses are $O(1)$ different, we see that the two heavy mass eigenstates ($\rho_{W_R^3}$ and ρ_X in Eq. (A9)) are to a good approximation the same as the weak basis, i.e., composite W and $(B-L)$, which, *a posteriori*, explains the labeling of the 2 massive eigenstates in Eq. (A9). The leading deviation from this identification stemming from mixing with elementary hypercharge, with composite-composite mixing being even smaller.

Therefore, the couplings of ρ_X (whether it is the lighter or heavier massive eigenstate) to $N^{(1)}$ and light quarks, which

to a very good approximation only couple to B_{elem} , are given by $\approx -\frac{1}{2}g_\star$ and $-Q_Y \frac{g_Y^2}{g_\star}$ respectively. Similarly, couplings of $\rho_{W_R^3}$ to $N^{(1)}$ and light quarks are given by $\approx \frac{1}{2}G_\star$ and $-Q_Y \frac{g_Y^2}{G_\star}$ respectively. Of course, for the study of the LHC signals, it suffices to keep only the lighter of these two states, since that production will dominate.

Whereas in the special case of degeneracy, we can show that the product of the couplings of the light quark and singlet neutrino to each of the above heavy mass eigenstates vanishes. This is as expected, since one of them corresponds to the ‘‘heavy’’ (or KK of the five-dimensional model) hypercharge which decouples from singlet neutrino, whereas the other one is composite (or KK of the five-dimensional model) Z' which decouples from light quarks instead.

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