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The two scales of new physics in Higgs couplings

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Abstract: Higgs coupling deviations from Standard Model predictions contain information about two scales of Nature: that of new physics responsible for the deviation, and the scale where new bosons must appear. The two can coincide, but they do not have to. The scale of new bosons can be calculated by going beyond an effective field theory description of the coupling deviation. We compute model-independent upper bounds on the scale of new bosons for deviations in Higgs to WW and ZZ couplings, finding that any measured deviation at present or future colliders requires the existence of new bosons within experimental reach. This has potentially interesting implications for naturalness.

Keywords: Anomalous Higgs Couplings, Vector-Like Fermions

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Contents

1 Introduction

The LHC broke our collective heart by discovering one of the most interesting particles in history[1] and nothing else. The prevailing response in the particle community is to abandon model-building in favor of model-independent methods of interrogating the data.

[1]What appears to be the only fundamental scalar discovered so far.

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The wealth of information still hidden in LHC data should be extracted using Effective Field Theory (EFT) techniques and the Standard Model (SM) fields. In this work we strive to find a path in between the model-building feast of the past and the EFT austerity of the present. We go beyond a SM EFT description of the data, but we are still able to make a model-independent statement about Higgs couplings. This could not have been done by considering only operators built out of SM fields.

A deviation in Higgs couplings, compared to the SM prediction, implies the existence of new bosons below a calculable energy scale. We compute this scale for hWW and hZZ couplings. In general, new bosons appear at an energy greater or equal than the new particles responsible for the deviation. However, in the case of hWW and hZZ, we find that any deviation observed at HL-LHC can only be generated by new bosons. Additionally, even deviations as small as those that can be probed at the most precise future lepton colliders require either new bosons roughly below 100 TeV, or new fermions with masses $M \simeq 100\,\text{GeV}$ and weak interactions with the SM, that can be discovered at HL-LHC.

We obtain these results by writing theories that contain only new fermions and generate the coupling deviation. We find that these theories are unstable under renormalization group equation (RGE) running and have to be modified above a given scale. Only new bosons can stabilize the running by avoiding a deep AdS minimum in the Higgs potential that gives rise to an unacceptably large rate of SM vacuum decay. We imagine having already discovered the new fermions responsible for the coupling deviation and we then calculate the finite range of validity of these theories. This logic was already outlined in [1, 2], but never applied to hWW and hZZ coupling deviations that are the most sensitive to the presence of new bosons, but also the most laborious to calculate. Compared to [1, 2] we introduce a second novelty. Instead of considering only a fixed set of low energy fermionic theories in small SM representations, we study the dependence of our results on the new fermions' representations and identify those that give the most conservative upper bound on the scale of new bosons, thus obtaining a model-independent result. Other works that discuss, in different contexts, the instability of the Higgs potential due to new fermions include [3–19].

The hWW and hZZ couplings are special because new fermions can modify them only at loop-level, but they exist at tree-level in the SM. As discussed in section 2, this implies the tightest upper bounds on the scale of new bosons if compared to other couplings exhibiting the same relative deviation. Additionally, hWW and hZZ are among the three Higgs couplings that will be most precisely measured at future colliders in terms of relative coupling deviation. This makes observing a deviation in these couplings particularly interesting, also in view of the possible connection between new bosons and naturalness.

The rest of the paper is organized as follows: in section 2 we outline the conceptual steps that lead to a bound on the scale of new bosons, in section 3 we introduce the fermionic theories that can induce a hWW or hZZ coupling deviation, in section 4 we calculate the coupling deviation and its dependence on the parameters of the fermionic theory, recalling the ingredients necessary to renormalize the electroweak sector of the SM at one loop. In section 5 we use these results to restrict the fermionic theories that we need to consider to obtain an upper bound on the scale of new bosons. The upper bound is then computed in section 6, where we comment on the detectability of the new bosons.

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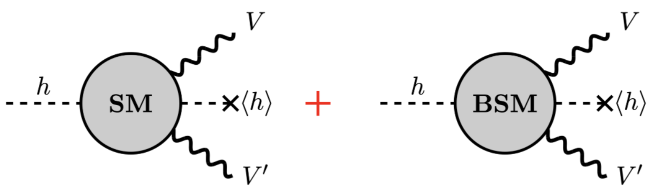
2 Basic idea

New particles that stabilize the Higgs mass via loop diagrams

can affect its production and decay rates through (almost) the same diagrams[2]



As a consequence, theories that are natural in the traditional sense (i.e. those where a new weak-scale symmetry explains the observed value of m_h^2) predict deviations in Higgs couplings compared to the SM expectation. In practice it is extremely hard to turn this suggestive picture into quantitative and model-independent statements relevant for colliders, but in this work we discuss a rare case where this is possible.

Measuring a deviation in Higgs couplings to WW and/or ZZ sets a calculable upper bound on the scale where new bosonic particles must appear in Nature. The bigger the deviation the smaller the upper bound. This is not an explicit statement about naturalness, but it has important implications for it. New bosons can be a definitive sign of unnaturalness if they have spin zero and come without symmetries (or anthropic roles) protecting them. However, they could equally well be the first sign of the long awaited symmetry that explains the value of m_h^2. The arguments in this work do not allow to distinguish between these two possibilities or to eliminate the third option (i.e. new bosons of spin ≥ 1 having nothing to do with m_h^2), but allow to compute the scale of the new bosons. Experiment will then give us the answer about naturalness.

The scale of new bosons can be computed by considering low energy theories that contain only new fermions in addition to the SM and showing that they have a finite range of validity. This work is a "proof by contradiction". We focus on the phenomenology of new fermions to prove that all theories with a hWW or hZZ coupling deviation must contain new bosons below an energy scale that can be calculated.

Note that this is quite different compared to the traditional EFT intuition that associates a new physics scale to a Higgs coupling deviation. Consider for example a $d = 6$ operator that contributes to a Higgs coupling to SM particles as $\delta g_h \simeq c(v^2/\Lambda^2)$. We imagine that we have already probed the scale $\Lambda \simeq v\sqrt{c/\delta g_h}$ and found only new fermions. Then

[2] The cartoon is adapted from [20].

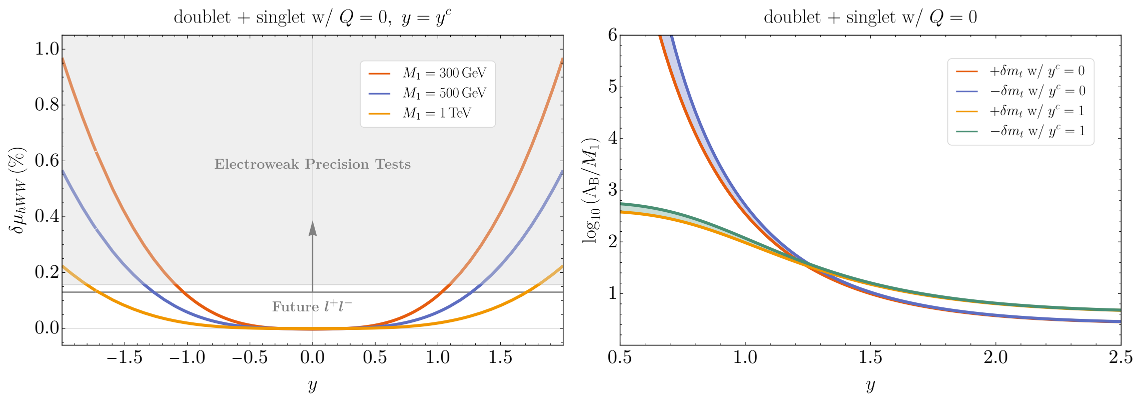


Figure 1. Left panel: relative hWW coupling deviation for a configuration of parameters of the doublet+singlet model (introduced in section 3.2) that maximizes it. The future collider sensitivity is discussed in section 6. Right panel: instability scale Λ_B of the doublet+singlet theory, as a function of one of its two Yukawa couplings (eq. (3.8)). The scale Λ_B is normalized to the lightest new fermion mass M_1. The two pairs of curves correspond to $y^c = 0$ (red and blue) and $y^c = 1$ (green and yellow). Curves with the same y^c differ by the choice of top Yukawa. We vary the top mass by twice its width to show that in these theories the instability scale is not very sensitive to the top Yukawa.

we show that these theories are valid only up to a scale $\Lambda_B > \Lambda$ where new bosons must appear. In practice we find that hWW or hZZ coupling deviations within reach of HL-LHC imply $\Lambda_B \simeq \Lambda$. Even deviations as small as one part in a thousand (that can be probed at future lepton colliders) set tight upper bounds on Λ_B. We give more precise results in section 6, where we also show that in most cases the new bosons are within reach of future hadron colliders or in some cases even the LHC.

Our calculation of Λ_B proceeds as follows: the only renormalizable coupling between new fermions and the SM that can influence SM Higgs couplings is a Yukawa interaction. Higgs couplings to WW and ZZ exist at tree-level in the SM, but the leading contribution from new fermions is at one (weak) loop. Therefore we need a large Yukawa coupling to generate a visible deviation, larger than any coupling in the SM. These large Yukawa couplings y dominate the RGEs of the model and can lead to two forms of instabilities. They can change the sign of the Higgs quartic coupling

$$\frac{d\lambda}{d\log\mu} \sim -\frac{y^4}{16\pi^2}\,, \tag{2.1}$$

and/or hit a Landau pole[3]

$$\frac{dy}{d\log\mu} \sim \frac{y^3}{16\pi^2}\,. \tag{2.2}$$

[3]Here and in the following we use "Landau pole" very loosely. As far as we know, in our theories there is no lattice proof that y hits an actual Landau pole. We instead identify Λ_B with the scale where y becomes non-perturbative. We comment on what we mean by non-perturbative and on the validity of our perturbative calculations in section 6.

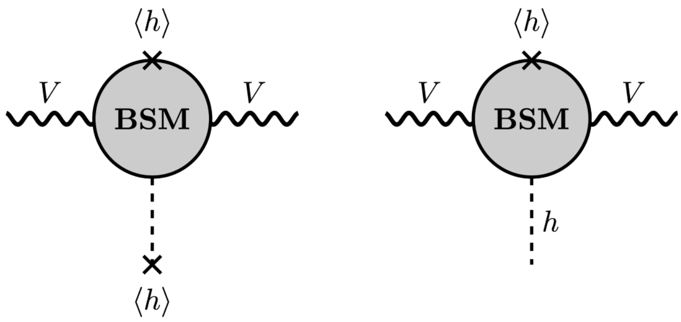


Figure 2. Schematic view of diagrams that generate a hVV coupling deviation (right) and at the same time affect electroweak precision measurements (left).

Therefore in these theories we can associate a cutoff to any given Higgs coupling deviation. The precise definition of the cutoff Λ_B can be found in section 6.1. We often refer to this scale, and sometimes plot it, in previous sections. The reader who prefers to see a definition first can skip ahead to section 6.1 that is self-contained.

To make contact with most studies of Higgs couplings at present and future colliders [21–28], we define an effective coupling from the Higgs partial width[4]

$$g_{hVV}^2 \equiv (g_{hVV}^{\rm SM})^2 \frac{\Gamma(h \to VV^{(*)})}{\Gamma^{\rm SM}(h \to VV^{(*)})}, \tag{2.3}$$

and a relative coupling deviation

$$\delta\mu_{hVV} \equiv \frac{g_{hVV} - g_{hVV}^{\rm SM}}{g_{hVV}^{\rm SM}} = \sqrt{\frac{\Gamma(h \to VV^{(*)})}{\Gamma^{\rm SM}(h \to VV^{(*)})}} - 1\,. \tag{2.4}$$

Since $|\delta\mu_{hVV}| \sim y^n$ with $n > 0$, the bigger $|\delta\mu_{hVV}|$, the faster the running of y and λ, giving a smaller scale at which our theory description must be modified. In figure 1 we summarize this discussion showing the size of the coupling deviation and of Λ_B for a given Yukawa coupling and a representative model discussed in section 3.2.

Figure 1 is just an illustration of the more detailed results in section 6, but it is sufficient to appreciate the main point of the paper. To measure any coupling deviation, even at the most precise e^+e^- colliders that are currently being proposed, we need large Yukawa couplings and relatively light fermions $M_1 \lesssim 500\,$GeV (left panel) corresponding to an instability scale a factor of 10 to a few above M_1 (right panel). Figure 1 contains also a quantitative illustration of the RGE domination of the new Yukawas: changing the top Yukawa by twice its error does not appreciably affect the scale of instability. To obtain all our results we compute the model RGEs using `SARAH` [29–34].

Even if we consider models with large Yukawas and light fermions, it is not easy to generate a large coupling deviation consistent with electroweak precision tests (EWPTs).

[4]When comparing decay widths we choose the following decays $Z^* \to e^+e^-$ and $W^* \to e\bar{\nu}_e$, neglecting the masses of the fermions whenever they give relative corrections smaller than $\delta\mu$.

The loop diagrams responsible for the deviation correct the WW and ZZ two-point functions, as shown schematically in figure 2. The parameters that most conveniently describe these corrections are known at the permille level from LEP [35] and in generic models they already set a bound on hWW and hZZ comparable to the sensitivity of future colliders, as shown in figure 1. Therefore, the impact of EWPTs on the maximally allowed hWW and hZZ deviations might be enough in itself to conclude that any observable deviation is generated by new bosons. For instance, we could have a gauge singlet scalar mixing with the Higgs after electroweak symmetry breaking that affects Higgs couplings at a measurable level, without compromising other electroweak precision observables.

However, even if we always show the EWPTs constraints on our models, we prefer to compute explicitly a conservative upper bound on Λ_B from the more model-independent RGE argument outlined above. The bounds from EWPTs can always be partially evaded with some amount of model-building[5] and the value of Λ_B from RGEs is low enough (for any observable deviation) to be interesting in itself and single out hWW and hZZ couplings as prime direct probes of new bosons and prime indirect probes of naturalness.

Since we are proving our statement by contradiction we want to find the theories with the largest possible Λ_B for any fixed $\delta\mu$. In the next three sections we show that only a handful of new fermions representations need to be considered to achieve this goal.

3 Fermionic low energy theories

We would like to extend the SM including new fermions that satisfy the following requirements: 1) They induce observable modifications in hZZ and hWW couplings, 2) They are consistent with all existing experimental constraints, 3) They are part of a consistent low energy theory (i.e. they do not introduce gauge anomalies). A comprehensive survey of representations that satisfy more general requirements was conducted in [36]. The only difference with respect to our needs is that in [36] deviations in any Higgs coupling to the SM were considered interesting. The new fermions relevant to our purposes are, therefore, a subset of those discussed in [36]. We can exclude the most minimal extensions considered in [36], comprising only one or two new fermions, because they induce unobservably small deviations in hZZ and hWW, as discussed in appendix A. The case of three chiral fermions was phenomenologically viable when [36] was published, but it is currently excluded by measurements of the hgg and $h\gamma\gamma$ couplings. We are left with SM extensions with three or more new fermions.

3.1 Three new fermions

The minimal extension of the SM that is relevant to us contains three new fermions, one vector-like pair and a Majorana fermion. As shown in [36] we can have two distinct possibilities. If we adopt the notation $(a,b)_Y$, with a the dimension of the $\mathrm{SU}(3)_c$ representation, b that of the $\mathrm{SU}(2)_L$ representation and Y the hypercharge, the first possibility is

$$L=(r,2n+1\pm1)_{-1/2}, \qquad N=(r,2n+1)_0, \qquad r=\bar{r}, \tag{3.1}$$

[5] However this comes at a price, it requires adding new Yukawas of the same size as those responsible for the Higgs coupling deviation [1], lowering Λ_B even further.

plus L^c, the vector-like partner of L. Note that we need the color representation to be self-conjugate ($r=\bar{r}$), as explained in [36]. We can then write the interaction Lagrangian

$$\mathcal{L}_3 \supset -yLHN - y^c L^c H^\dagger N - M_L L L^c - \frac{M_N}{2} N^2 + \text{h.c.}. \tag{3.2}$$

Alternatively, we can consider a $SU(2)_L$ representation of even dimension for[6] N

$$L = (r, 2n \pm 1)_{-1/2}, \qquad N = (r, 2n)_0, \qquad r = \bar{r}, \tag{3.3}$$

but in this case we do not have a Majorana mass term and the Lagrangian reads

$$\mathcal{L}_3' \supset -yLHN - y^c L^c H^\dagger N - M_L L L^c + \text{h.c.}. \tag{3.4}$$

These are the minimal extensions of the SM consistent with our requirements. We have listed them for completeness, but in the following we mostly consider theories with two pairs of vector-like fermions. The reason is that we are not looking for a minimal extension of the SM, but for the extension that has the largest possible cutoff Λ_B for a given Higgs coupling deviation. Including another fermion allows us to add another layer of generality to the low energy theory: the hypercharges are not fixed by the three conditions stated at the beginning of this section and we have the extra freedom to dial Y to increase the hZZ coupling deviation. Going beyond four fermions does not introduce any other qualitative difference, but we discuss how our results scale if we include multiple copies of the four fermions presented in the next section.

3.2 The main character: a four fermion extension of the SM

The simplest extension of the SM with four new fermions that can give a potentially large hWW and/or hZZ coupling deviation, without being obviously excluded by current measurements and without introducing gauge anomalies is given by

$$L = (r, n)_Y, \qquad N^c = (\bar{r}, n-1)_{-Y-1/2} \tag{3.5}$$

and their vector-like partners

$$L^c = (\bar{r}, n)_{-Y}, \qquad N = (r, n-1)_{Y+1/2}. \tag{3.6}$$

Their renormalizable Lagrangian (leaving implied kinetic terms and gauge interactions) is[7]

$$\mathcal{L}_4 = -yLHN^c - y^c L^c H^\dagger N - M_L L L^c - M_N N N^c + \text{h.c.}. \tag{3.8}$$

[6] Note that if $2n = 2+4k$, $k \in \mathbb{N}$, we need r to be even to avoid a $SU(2)_L$ anomaly [36].

[7] The case $r=\bar{r}, Y=0$ deserves special attention. If we keep the same particle content as above with N distinct from N^c the theory in eq. (3.8) is consistent and preserves a U(1) "lepton number" with charges $Q_{L,N}=1$, $Q_{L^c,N^c}=-1$. However we can break the U(1) symmetry and include additional Yukawa couplings

$$\mathcal{L}_2 \supset -M_L L L^c - M_N N N^c - yLHN^c - y^c L^c H^\dagger N - y' LHN - y^{c\prime} L^c H^\dagger N^c + \text{h.c.}. \tag{3.7}$$

For N, N^c in $SU(2)_L$ representations of odd dimension we can add also Majorana masses. We do not consider these cases in detail because they do not change our conclusions in section 6.

In the following we often discuss a subset of these models for which it is useful to have a name. We call *doublet+singlet* the model with $r=1, n=2, Y=-1/2$ and *doublet+triplet* the model with $r=1, n=3, Y=0$. In the doublet+singlet model we call ℓ, ℓ^c the two doublets that have the same quantum numbers as SM lepton doublets, and n, n^c the two singlets that have the same quantum numbers as right-handed neutrinos. The particles in the doublet+triplet model have the same quantum numbers as a pair of vector-like winos plus two Higgsinos.

When we think about the limit of heavy new fermions, we have to take $M_{L,N}$ large. Taking the Yukawa couplings large, while leaving $M_{L,N}$ fixed is already excluded by hgg and $h\gamma\gamma$ measurements, as mentioned at the beginning of section 3 when discussing chiral fermions.

Note also that the four parameters y, y^c, M_L, M_N cannot all be taken to be real and positive without loss of generality. There is one physical phase in this model. However, we want to avoid stringent constraints on CP violation that increase $M_{L,N}$ and reduce our cutoff Λ_B. Therefore, we always take the only physical phase to be 0 or π, allowing at most for a relative sign between the two Yukawa couplings.

We are going to study these theories in their perturbative limit. The Yukawas that we consider can be larger than one, but at the mass scale of the new fermions they are perturbative. This gives the most conservative upper bound on the scale of new physics that we want to compute. A simple argument is enough to see this: imagine a composite Higgs theory that in the UV is described by new fermions that condense at low energy. We call f the scale of the σ-model that describes the pions of the confining sector, including the Higgs. A low energy observer measures Higgs coupling deviations at $\mathcal{O}(v^2/f^2)$ from a series of irrelevant operators [37]. The scale where new particles (and new bosons) appear is $m_* = g_* f$, for a strongly coupled theory $m_* \sim 4\pi f$. This should be compared with our weakly coupled theories where the Higgs coupling deviations arise at one-loop at $\mathcal{O}(y^4 v^2/16\pi^2 M^2)$ with M a vector-like mass and y a coupling. We can compare with an effective $f_{\rm eff} \simeq 4\pi M/y^2$ and conclude that these theories can be extrapolated to much higher energies than $4\pi f_{\rm eff}$ and this is why we focus on them to get a conservative upper bound on the scale of new physics. When $y \lesssim \mathcal{O}(1)$ we will indeed find that our perturbative theories have much larger cutoffs than $4\pi f_{\rm eff}$. In some special cases, with large coupling deviations and large y, the running is sufficiently rapid to give cutoffs lower than $4\pi f_{\rm eff}$. These latter cases must be taken with a grain of salt. First of all, our two-loop approximation is not adequate to capture the precise value of the cutoff in these limiting cases. Secondly, one could get a more conservative upper bound in a strongly coupled theory, so in practice any cutoff below $4\pi f_{\rm eff}$ is not the largest one that is possible to achieve.

To conclude this section, note that in most of these models (the only exception being the doublet+singlet model with a neutral singlet) we generate also a deviation in the $h\gamma\gamma$ coupling. This is a much bigger relative effect than the deviation in hWW and hZZ because we are comparing a one-loop new physics effect with a one-loop coupling in the SM. For completeness we compute the scale of new bosons for $h\gamma\gamma$ in section 6.3, but this case was already discussed in [1, 2, 6–8].

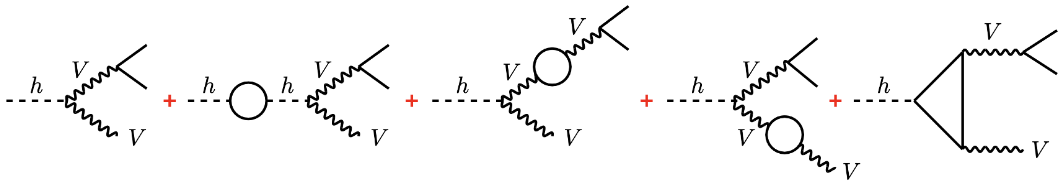


Figure 3. Schematic view of the leading diagrams contributing to the $h \to VV^* \to V\psi_a\psi_a$ decay width.

4 Calculation of the coupling deviations

The coupling deviation that we want to compute receives contributions from the diagrams that are listed schematically in figure 3. We have included in the figure only loop corrections from the new fermions, because SM loop corrections cancel in the ratio of eq. (2.4).

To implement the calculation consistently, in addition to the diagrams in figure 3, we have to consider additional processes, involving the new fermions, that renormalize the electroweak sector of the SM at one loop. We start by considering the diagrams already contained in figure 3. The bubble diagrams on the external legs have two effects: they modify the tree-level relation between the pole mass of SM bosons and the corresponding $\overline{\text{MS}}$ Lagrangian parameters, and they give a field strength renormalization that we need to include when computing S-matrix elements from Feynman diagrams. In the ZZ case they also induce a mixing with the photon that we discuss after the WW calculation.

Following the notation in [38], we call $\hat{m}_Z, \hat{e}, \dots$ the measured parameters and $m_Z, e, \dots$ the $\overline{\text{MS}}$ Lagrangian parameters. We can compute the $\overline{\text{MS}}$ V-boson mass ($V = W, Z$) from the measured mass $\hat{m}_V$ as

$$m_V^2 = \hat{m}_V^2 \left(1 + \frac{\text{Re}[\Pi_{VV}(\hat{m}_V^2)]}{\hat{m}_V^2}\right), \tag{4.1}$$

where

$$\mathcal{M}^{\mu\nu}(V(p) \to V(p)) = g^{\mu\nu}\Pi_{VV}(p^2) + \dots . \tag{4.2}$$

The field strength renormalization is given by

$$Z_V = \frac{1}{1 - \Pi'_{VV}(\hat{m}_V^2)}, \tag{4.3}$$

where the prime denotes derivation by p^2. Note that the sign in front of Π' is opposite for the Higgs boson, if we use the standard convention for scalar propagators [38, 39]. These equations are useful to set the notation, but we do not give further details because this is a standard calculation for EWPTs, that can be found, for instance, in chapter 31 of [38].

The bubble diagram on the internal V-boson leg in figure 3 is slightly more tricky. Summing all 1PI diagrams we obtain the propagator

$$\frac{-ig^{\mu\nu}}{p^2 - m_V^2 - \Pi_{VV}(p^2)} + \mathcal{O}(p^\mu p^\nu). \tag{4.4}$$

We can omit terms proportional to $p^\mu p^\nu$ since the internal gauge boson is contracted with a light SM current. The function $\Pi_{VV}(p^2)$ contains $1/\epsilon$ and μ-dependent terms that might be modified non-trivially by the integral over p^2. To make all the cancellations manifest we rearrange the terms in the propagator as

$$\frac{-ig^{\mu\nu}}{p^2-m_V^2-\Pi_{VV}(p^2)} = \frac{-iZ_V g^{\mu\nu}}{p^2-\hat{m}_V^2-\Xi_{VV}(p^2)}, \tag{4.5}$$

where we have defined the new function $\Xi_{VV}(p^2)$ as

$$\Xi_{VV}(p^2) \equiv \Pi_{VV}(p^2)-\mathrm{Re}[\Pi_{VV}(\hat{m}_V^2)]-(p^2-\hat{m}_V^2)\Pi'_{VV}(\hat{m}_V^2). \tag{4.6}$$

It is easy to show that $\Xi(p^2)$ is UV-finite and μ-independent. The divergence and μ-dependent pieces in Z_V are independent of p^2 and can be factored out of the integral.

The triangle diagram[8] in figure 3 gives a complicated function of external momenta, let us call this function T, then the amplitude for the decay $h\to WW^*$ reads

$$\begin{aligned}&\mathcal{M}(h\to W(p_1)W^*(q_{23})\to W(p_1)\psi_a(p_2)\psi_b(p_3)) = \\ &= i\frac{g}{\sqrt{2}}\varepsilon^*_\nu(p_1)\,[\bar{u}(p_3)\gamma_\mu P_L v(p_2)]\,\frac{m_W(g+T(p_1,q_{23}))}{q_{23}^2-\hat{m}_W^2+\Xi(q_{23}^2)}\sqrt{Z_W^3 Z_h}\,,\end{aligned} \tag{4.7}$$

where we have included also the tree-level contribution from the first diagram in figure 3. To get a physical answer we express g and m_W in terms of measured quantities using the one-loop renormalization conditions [38, 40]

$$\begin{aligned} e^2 &= \hat{e}^2(\hat{m}_Z)\left[1-\frac{\Pi_{\gamma\gamma}(\hat{m}_Z^2)}{\hat{m}_Z^2}\right], \\ m_Z^2 &= \hat{m}_Z^2\left(1-\frac{\mathrm{Re}[\Pi_{ZZ}(\hat{m}_Z^2)]}{\hat{m}_Z^2}\right), \\ s_W^2 &= \hat{s}_W^2\left[1+\frac{\hat{c}_W^2}{\hat{c}_W^2-\hat{s}_W^2}\left(\frac{\mathrm{Re}[\Pi_{ZZ}(\hat{m}_Z^2)]}{\hat{m}_Z^2}-\frac{\Pi_{\gamma\gamma}(\hat{m}_Z^2)}{\hat{m}_Z^2}-\frac{\mathrm{Re}[\Pi_{WW}(0)]}{\hat{m}_W^2}\right)\right]. \end{aligned} \tag{4.8}$$

These conditions incorporate also the relevant effects of the new fermions that are not illustrated by figure 3. The measured parameters that we use are [38, 41]

$$\begin{aligned} \hat{\alpha}(0) &= (137.035999074\pm 0.000000044)^{-1}\to \hat{e}^2(\hat{m}_Z) = 4\pi\hat{\alpha}(m_Z) = 4\pi(127.944\pm 0.014)^{-1}, \\ \hat{m}_Z &= (91.1876\pm 0.0021)\,\mathrm{GeV}, \\ \hat{G}_F &= (1.1663787\pm 0.0000006)\times 10^{-5}\,\mathrm{GeV}^{-2}. \end{aligned} \tag{4.9}$$

Here and in the following we call c_W and s_W the cosine and sine of Weinberg's angle and we neglect the SM running between m_Z and m_h. Note that g and m_W are not the same in the SM and in our vector-like fermion theory because the two-point functions Π_{XX} differ in the two cases. Combining the three previous equations we obtain

$$\mathcal{M}(h\to W(p_1)W^*(q_{23})\to W(p_1)\psi_a(p_2)\psi_b(p_3)) = \hat{\mathcal{M}}\left(1+\frac{\Pi_g}{2}\right), \tag{4.10}$$

[8] For $h\to WW^*$ there is a single diagram, in the case of $h\to ZZ^*$ the figure represents schematically three independent diagrams. Two with two Z-bosons that can be obtained by exchanging the momenta and Lorentz indexes of the two Z's. The last one with an on-shell Z and an off-shell photon.

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where we found convenient to define

$$\Pi_g \equiv \frac{3\left(\Pi'_{WW}(\hat{m}_W^2)\right)}{2} - \frac{\left(\Pi'_{hh}(\hat{m}_h^2)\right)}{2} - \frac{\Pi_{\gamma\gamma}(\hat{m}_Z^2)}{\hat{m}_Z^2} - \frac{\mathrm{Re}\left[\Pi_{ZZ}(\hat{m}_Z^2)\right]}{2\hat{m}_Z^2} - \frac{\Pi_R}{\hat{c}_W^2 - \hat{s}_W^2}\left(\frac{\hat{s}_W^2}{2} + \hat{c}_W^2\right), \tag{4.11}$$

with

$$\Pi_R = -\frac{\Pi_{\gamma\gamma}(\hat{m}_Z^2)}{\hat{m}_Z^2} + \frac{\Pi_{ZZ}(\hat{m}_Z^2)}{\hat{m}_Z^2} - \frac{\Pi_{WW}(0)}{\hat{m}_W^2} \tag{4.12}$$

and $\hat{\mathcal{M}}$ is the matrix element in the second line of eq. (4.7) where all $\overline{\mathrm{MS}}$ Lagrangian parameters are replaced by measured quantities, $Z_W = Z_h = 1$ and we drop terms $\sim T \times \Pi_g$ that are of second order in the loop expansion.

Using eq.s (4.10) and (4.11) we can compute $\delta\mu_{hWW}$ by summing over all SM fermions in the final states (except the top quark). Taking $y^c = y$, $M_L = M_N = M$ and $M \gg yv$ we have, for the doublet+singlet model defined in the previous section ($Y = -1/2$),

$$\begin{aligned}
\delta\mu_{hWW}^{\mathrm{DS}} &= \frac{y^2 v^2}{80\pi^2 M^2}\left[y^2 \frac{(1+c_W^2)}{c_W^2 - s_W^2} - \frac{m_h^2}{v^2} - \frac{g^2}{4}\frac{s_W^2}{c_W^2 - s_W^2}\right] + \frac{1}{2}\sum_{i=1}^{3} c_i^W \frac{m_h^3}{m_W M^2}\frac{P_W^{(i)}(x)}{R_W(x)}, \\
c_1^W &= \frac{13 y^2}{240\pi^2}, \qquad c_2^W = -c_3^W = -\frac{7y^2}{240\pi^2}, \qquad x \equiv \frac{m_W^2}{m_h^2},
\end{aligned} \tag{4.13}$$

where all parameters are the measured ones and we neglected terms without new Yukawas. We left the "hats" implied to improve readability. We have already accounted for one-loop renormalization conditions and one can express the above parameters in terms of the measured quantities in eq. (4.9) using the SM tree-level relations (i.e. $e = g s_W$, $m_W = c_W m_Z$). The functions $P_W^{(i)}(x)$ can be found in appendix C, they correspond to the three higher-dimensional operators: $O_1 = h\partial_\nu W_\mu \partial^\nu W^\mu$, $O_2 = h\partial_\nu W_\mu \partial^\mu W^\nu$, and $O_3 = h(\Box W_\mu)W^\mu$. Note that these operators exist only after EW symmetry breaking and are suppressed by an extra m_W/M compared to their naive scaling dimension. All other operators are either redundant or proportional to $\partial_\mu W^\mu$ which is zero for the on-shell W and when contracted with a massless SM current. Corrections proportional to the SM fermions' masses are smaller than the smallest values of $\delta\mu$ that we consider in this work.

The parametric form of the result can be understood from the two diagrams in figure 4. From the left panel we see that at least two insertions of y are needed to close the loop. Additionally, we have two powers of a SM gauge coupling from the two gauge boson vertices. We might naively conclude that the coupling deviation from the triangle diagrams scales as $\delta g_{hWW} \sim g^2 y^2 v^3/16\pi^2 M^2$. If we restore units to $\hbar$ we see immediately that this is wrong because δg_{hWW} has dimensions of $\hbar^{-1} \times v$. If we take into account the $\hbar$ from the loop, the coupling deviation from these diagrams must scale as $(g^2 y^2 v^3/16\pi^2 M^2) \times g_*^2$ where g_* has the dimensions of a gauge coupling ($\hbar^{-1/2}$). The largest coupling in the theory is y, so we need to evaluate the diagram in the right panel of figure 4 which gives parametrically the leading term in eq. (4.13).[9]

[9]In practice we always work with mass eigenstates and resum all y insertions. We show figure 4 just to illustrate the parametrics.

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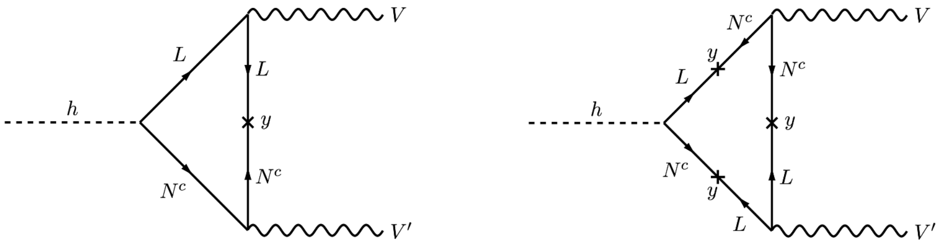

Figure 4. Left panel: naive leading diagram for the hVV width from vector-like leptons. Right panel: leading diagram for the hVV width from vector-like leptons in the limit of large Yukawa couplings ($y \gtrsim g_{\rm SM}$).

For the general $\mathrm{SU}(2)_L$ representations in eq. (3.5) and (3.6) and $Y = -1/2$, for $y^c = (-1)^{n-1} y$ (the choice that makes the mass matrices symmetric) we have

$$\begin{aligned}
\delta\mu_{hWW} &= \frac{n y^2 v^2}{240\pi^2 M^2}\left[y^2 \frac{(2n-1)(c_W^2 + \frac{s_W^2}{2})}{(n-1)(c_W^2 - s_W^2)} - \frac{3m_h^2}{2v^2} - g^2 \frac{(23-10n)s_W^2}{8(c_W^2 - s_W^2)}\right] \\
&\quad + \frac{1}{2}\sum_{i=1}^{3} c_i^W \frac{m_h^3}{m_W M^2}\frac{P_W^{(i)}(x)}{R_W(x)}, \qquad x \equiv \frac{m_W^2}{m_h^2}, \\
c_1^W &= n\frac{12y^2}{640\pi^2} + \left(\frac{1}{2} + \frac{n}{5} - (n-1)^2 - \frac{(n-1)^3}{2}\right)\frac{y^2}{36\pi^2}, \\
c_2^W &= -n\frac{y^2}{480\pi^2} + \left(-\frac{1}{2} - \frac{n}{20} + (n-1)^2 + \frac{(n-1)^3}{2}\right)\frac{y^2}{36\pi^2}, \\
c_3^W &= n\frac{17y^2}{960\pi^2} + n\frac{y^2}{320\pi^2}.
\end{aligned} \tag{4.14}$$

As before we took $M \gg yv$, dropped terms that do not contain powers of y and left the "hats" implied, but all parameters are the measured ones. The growth of $\delta\mu_{hWW}$ as n^3 at large n is expected on general grounds from the Higgs low energy theorems [42–45] that relate the coupling deviation to the W two-point function. The two-point function scales as the Dynkin index of the representation that for SU(2) at large n goes as $T(n) \sim n^3$.

The case of hZZ is similar, but not identical. There are three additional diagrams: 1) the bubble diagram on the internal Z-leg can mix the Z with a photon, 2) the triangle diagram in figure 3 can have a photon in the internal gauge boson line, 3) there are two triangle diagrams with two Z's that differ by the exchange of momenta and Lorentz indexes of the two gauge bosons. Including these diagrams, we have

$$\mathcal{M}(h \to Z^* Z \to Z\bar{\psi}\psi) = \hat{\mathcal{M}}_Z\left(1 + \frac{\Pi_g^Z}{2}\right), \tag{4.15}$$

where

$$\begin{aligned}
\Pi_g^Z &\equiv \frac{3\,(\Pi'_{ZZ}(\hat{m}_Z^2))}{2} - \frac{(\Pi'_{hh}(\hat{m}_h^2))}{2} - \frac{\Pi_{\gamma\gamma}(\hat{m}_Z^2)}{\hat{m}_Z^2} - \frac{\mathrm{Re}\,[\Pi_{ZZ}(\hat{m}_Z^2)]}{2\hat{m}_Z^2} \\
&\quad + \frac{a_2\hat{s}_W^2 + 2a_4\hat{s}_W^4}{2(a_0 + a_2\hat{s}_W^2 + a_4\hat{s}_W^4)}\left[\frac{\hat{c}_W^2}{\hat{c}_W^2 - \hat{s}_W^2}\Pi_R - \frac{\hat{c}_W}{\hat{s}_W}\left(\Pi'_{Z\gamma}(0)\right)\right].
\end{aligned} \tag{4.16}$$

$a_{1,2,3}$ are $\mathcal{O}(1)$ numbers that depend on the SM fermion in the final state. We can obtain them from the SM Z-boson current,

$$(a_0 + a_2 \hat{s}_W^2 + a_4 \hat{s}_W^4) \equiv (T_3 - \hat{s}_W^2 Q)^2 . \tag{4.17}$$

The explicit result for the doublet+singlet model for $Z^* \to e^+e^-$, taking $y^c = y$, $M_L = M_N = M$ and $M \gg v$, is

$$\begin{aligned}
\delta\mu_{hZZ}^{\rm DS} &= \frac{y^4 v^2}{40\pi^2 M^2} \frac{1-2c_W^4+4(-1+2c_W^4)s_W^2+8(1-2c_W^2)s_W^4}{c_{2W}(1-4s_W^2+8s_W^4)} \\
&\quad + \frac{y^2v^2}{480\pi^2 M^2}\left[-6\frac{m_h^2}{v^2}+g^2\frac{s_W^2(4s_W^2-1)}{c_{2W}(1-4s_W^2+8s_W^4)}\right]+\frac{1}{2}\sum_{i=1}^{3} c_i^Z \frac{m_h^2}{M^2}\frac{P_Z^{(i)}(z)}{R_Z(z)} , \\
c_1^Z &= \frac{13y^2}{240\pi^2} , \qquad c_2^Z = -c_3^Z = -\frac{7y^2}{240\pi^2} , \qquad z \equiv \frac{m_Z^2}{m_h^2} .
\end{aligned} \tag{4.18}$$

All parameters are the measured ones and we neglected terms without new Yukawas. We left the "hats" implied to improve readability and neglected terms that do not contain powers of the new Yukawas. The functions $P_Z^{(i)}(x)$ can be found in appendix C. The result for general $\mathrm{SU}(2)_L$ representations and $Y = -1/2$ is

$$\begin{aligned}
\delta\mu_{hZZ} &= \frac{n(2n-1)}{n-1}\frac{y^4v^2}{240\pi^2M^2}\frac{(1-2c_W^4)(1-4s_W^2)+8s_W^4(1-2c_W^2)}{c_{2W}(1-4s_W^2+8s_W^4)} \\
&\quad + \frac{ny^2v^2}{960M^2\pi^2}\left[-6\frac{m_h^2}{v^2}+\frac{g^2}{3}(23-10n)\frac{s_W^2(-1+4s_W^2)}{c_{2W}(1-4s_W^2+8s_W^4)}\right]+\frac{1}{2}\sum_{i=1}^{4} c_i^Z \frac{m_h^2}{M^2}\frac{P_Z^{(i)}(z)}{R_Z(z)} , \\
c_1^Z &= n\frac{y^2}{1440\pi^2 c_W^2}\left(-1+60s_W^2-20s_W^4+20(n-1)c_W^2+20(n-1)^2c_W^4\right) , \\
c_2^Z &= -n\frac{y^2}{1440\pi^2 c_W^2}\left(-19+60s_W^2-20s_W^4+20(n-1)c_W^2+20(n-1)^2c_W^4\right) , \\
c_3^Z &= n\frac{7y^2}{480\pi^2 c_W^2} , \\
c_4^Z &= n(n-2)\frac{y^2}{144\pi^2}\frac{t_W^2\left(2s_W^2-\frac{1}{2}\right)}{2s_W^4-s_W^2+\frac{1}{4}}\left(3-2s_W^2+2(n-1)c_W^2\right) .
\end{aligned} \tag{4.19}$$

The parametric form of the results for hZZ has the same explanation as that of hWW, given above. The expansions for large M are useful for the arguments in section 5, but we often have to compute $\delta\mu$ also when $M \simeq yv$, for example in the doublet+singlet model introduced in the previous section. When the vector-like masses are comparable to yv, we compute numerically the loop function T and Π_g in eq. (4.7) (and the equivalent functions for hZZ), using `Package-X` [46, 47] and then we integrate them numerically over phase space using `Mathematica`.

5 Fermion representations

Before making quantitative predictions on the scale of new bosons, we want to understand what color, $\mathrm{SU}(2)_L$ and hypercharge representations we need to consider to obtain the

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most conservative upper bound on this scale. We can considerably reduce the number of representations that we need to study, compared to eq.s (3.5) and (3.6), by following a few scaling arguments presented in the next sections. We also discuss the impact of adding multiple copies of the four fermions introduced in section 3.

5.1 Color and fermion multiplicity

We can begin by showing that we do not need to consider colored particles. In section 4 we have discussed the parametric form of the expected coupling deviation. It is straightforward to include the scaling with the $\mathrm{SU}(3)_c$ representation dimension r,

$$\delta g_{hVV} \sim r \frac{\alpha_V}{4\pi} \frac{y^4 v^2}{M^2} \lesssim r \frac{\alpha_V}{4\pi} \frac{y^4 v^2}{M_{\mathrm{exp}}^2}, \tag{5.1}$$

where M_{exp} is the experimental bound on the mass of the new states and α_V is an electroweak gauge coupling that depends on the choice of final state. We can use this scaling because LHC bounds require $M_{\mathrm{exp}} \gg yv$ for colored states. We have assumed small $\mathrm{SU}(2)_L$ representations and hypercharge. For simplicity we have also taken $y^c = y$ and $M_L = M_N = M$, more general choices do not affect the conclusions of this section. Eq. (5.1) is valid both for the four fermions extensions in section 3.2 and the three fermions in section 3.1.

As discussed in section 2, above a certain cutoff we do not have perturbative control of the theory or the Higgs potential becomes strongly unstable. We want to find the value of r that gives the largest cutoff, given a fixed $\delta g_{hVV}/g_{hVV}^{\mathrm{SM}}$. Focusing on the representation that gives the largest cutoff allows us to set an upper bound on the scale at which new physics *must* appear for any fixed value of $\delta g_{hVV}/g_{hVV}^{\mathrm{SM}}$.

We can easily conclude that we do not need to consider $r > 8$, since larger representations generate a Landau pole in the QCD coupling g_s a factor of a few above the mass of the new particles [2], independently of the value of $\delta g_{hVV}/g_{hVV}^{\mathrm{SM}}$.

We can do better and show that we need to consider only color singlets ($r = 1$). For small $\mathrm{SU}(3)_c$ representations, $r \leq 8$, the leading effects that make us lose perturbative control of the theory or destabilize the Higgs potential are due to the new couplings y, y^c. If we hit a Landau pole before any Higgs instability, from the RGEs we can conclude that a fixed value for the couplings $y_r^{(c)}$ defined as

$$y_r^{(c)} \equiv y^{(c)} \sqrt{r} \tag{5.2}$$

gives the same scale of the Landau pole for any r. y_r and y_r^c play a similar role as the 't Hooft coupling in gauge theories at large-N [48]. Therefore what is most relevant for us is the scaling of δg_{hVV} at fixed $y_r^{(c)}$ and not at fixed values of the Yukawas $y^{(c)}$. This is given by

$$\delta g_{hVV} \sim r \frac{\alpha_V}{4\pi} \frac{y^4 v^2}{M^2} = \frac{\alpha_V}{4\pi} \frac{y_r^4 v^2}{r M^2} \lesssim \frac{\alpha_V}{4\pi} \frac{y_r^4 v^2}{r M_{\mathrm{exp}}^2}. \tag{5.3}$$

Colored states give a smaller δg_{hVV} at fixed $y_r^{(c)}$ compared to their colorless counterparts. This effect is further enhanced by the dependence of M_{exp} on r. Collider bounds on colored particles masses $M_{\rm exp}$ are stronger due to their larger production cross sections. So it is sufficient to consider vector-like leptons to get the most conservative upper bound on the scale of new bosons.

If a Higgs instability arises before any Landau pole, the couplings that give a fixed scale of the instability for any r are

$$y_r^{(c)\prime} \equiv y^{(c)\prime} r^{1/4}, \tag{5.4}$$

and the Higgs coupling deviation scales as

$$\delta g_{hVV} \sim r \frac{\alpha_V}{4\pi} \frac{y^4 v^2}{M^2} = \frac{\alpha_V}{4\pi} \frac{y_r^{\prime 4} v^2}{M^2} \lesssim \frac{\alpha_V}{4\pi} \frac{y_r^{\prime 4} v^2}{M_{\rm exp}^2}. \tag{5.5}$$

Once again we can conclude that colored states give a smaller δg_{hVV} at fixed $y_r^{(c)}$ compared to their colorless counterparts, since collider bounds on their masses $M_{\rm exp}$ are stronger due to their larger production cross section.

Ultimately in our analysis it is sufficient to consider only vector-like leptons

$$L = (1,n)_{Y-1/2}, \qquad N^c = (1,n-1)_{-Y}, \qquad L^c = (1,n)_{-Y+1/2}, \qquad N = (1,n-1)_Y. \tag{5.6}$$

If a Higgs coupling deviation is measured it will be instructive to repeat our analysis for $r \neq 1$ in order to go beyond our scaling arguments and obtain precise predictions for the scale of new bosons for any r. We leave this study to future work.

Models with new vector-like (or chiral) leptons have already been studied extensively in the literature in relation to neutrino masses, the muon anomalous magnetic moment, flavor models, and a variety of other subjects. A comprehensive review is beyond the scope of this work, but we refer to [1, 2, 4–7, 9–19, 36, 49–315, 345] for a list of relevant studies.

The only possible loophole in our argument is that we focused on the part of the coupling deviation proportional to y^4, but there are also terms proportional to $y^2 g_{\rm SM}^2$. The $y^2 g_{\rm SM}^2$ terms proportional to yy^c grow with the dimension of the SU(2)$_L$ representation as n^3, versus the more modest linear growth of the y^4 terms, so one might worry that states with large r and n invalidate our argument. However, we verified numerically that for $\delta\mu_{hVV} \gtrsim 0.13\%$ (the sensitivity of the most precise future lepton colliders) the y^4 term always dominates if M_1 is compatible with LHC bounds on colored particles. This is shown in figures 5 and 6. The mass of the lightest new states in the figures is fixed at $M_1 = 800\,$GeV to represents a conservative lower bound on masses of new colored states at the LHC. The upper bound on SU(2)$_L$ representations and hypercharges in the figure is determined by low energy Landau poles, as detailed in section 5.2.

Before turning to SU(2)$_L$ and hypercharge quantum numbers, it is worth pointing out that adding multiple copies of the fermions in eq. (5.6) does not allow to raise the maximal UV cutoff. The scaling arguments that lead to this conclusion are exactly the same as those given for the color representation, replacing the dimension of the representation r with the number of copies $\mathcal{N}$. In making this statement we are implicitly assuming that all $\mathcal{N}$ fermions are close in mass. However, if they are not, one generation dominates the

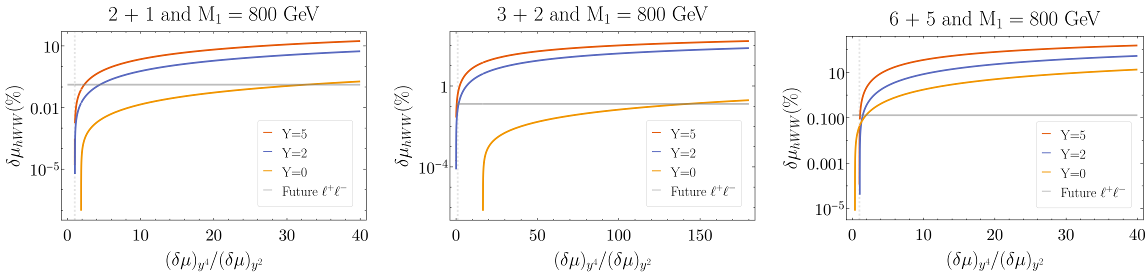

Figure 5. Relative importance of the y^4 and $y^2 g_{\mathrm{SM}}^2$ terms in the hWW coupling deviation. The gray dashed line signals where the two terms are equal. The y^4 terms dominate for coupling deviations large enough to be detected at future lepton colliders (solid gray line) or HL-LHC. The future collider sensitivity is discussed in section 6.

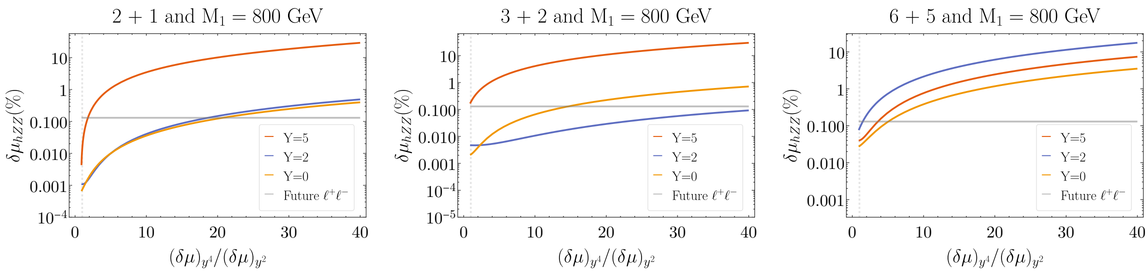

Figure 6. Relative importance of the y^4 and $y^2 g_{\mathrm{SM}}^2$ terms in the hZZ coupling deviation. The gray dashed line signals where the two terms are equal. The y^4 terms dominate for coupling deviations large enough to be detected at future lepton colliders (solid gray line) or HL-LHC. The future collider sensitivity is discussed in section 6.

coupling deviation and we are effectively in the $\mathcal{N} = 1$ case (or a case with even lower cutoff if multiple Yukawas are large).

One might envision adding $\mathcal{N}$ fermions that are close enough in mass to contribute comparably to δg_{hVV}, but still far enough not to show up as a $\mathcal{N}^2$ enhancement of the production cross-section. In this case we cannot strictly say that M_{exp} in the previous equation increases compared to the $\mathcal{N} = 1$ case, as we did for colored fermions, but even if M_{exp} were to remain the same, the above arguments are still valid and show that the UV cutoff with $\mathcal{N} > 1$ copies is not larger than that obtained for $\mathcal{N} = 1$. Note that this is not true for Higgs couplings to massless gauge bosons, due to the different scaling of the coupling deviation with y. In those cases $r > 1$ and/or $\mathcal{N} > 1$ can give a bigger cutoff for a fixed coupling deviation [1].

5.2 SU(2)$_L$ and hypercharge

We can obtain an upper bound on the largest n and Y that we need to consider from Landau poles in gauge couplings. If, for a given choice of fermions representations, we get a Landau pole right above the masses of the new fermions, we do not need to consider larger representations because the theory needs to be extended very close to where the new fermions appear, independently of δg_{hVV}.

In figure 7 we show the location of the Landau pole in g and g' for $y = y^c = 0$. For each plot we take the smallest value of the other parameter that we want to bound, so $n = 2$ for the hypercharge Landau pole and $Y = 0$ for the $\mathrm{SU}(2)_L$ Landau pole. This is conservative, since larger values move the Landau pole to lower energies, as can be seen from the model RGEs. We have verified that taking the new Yukawas to be different from zero does not affect appreciably the location of the Landau pole in the gauge coupling, unless we take them so large that we hit an instability even earlier.

For $n = 8$, i.e. when we add to the theory a vector-like pair of 8's and a vector-like pair of 7's, the $\mathrm{SU}(2)_L$ Landau pole is a factor of 10 above the heaviest new lepton mass, M_{max}, signalling that the theory is barely under control perturbatively. For this reason we consider only $n \leq 7$ in what follows. In section 6 we show explicitly Λ_B only up to $n = 4$, as larger representations give even lower values of Λ_B at fixed $\delta\mu$. We could be more aggressive and consider only smaller representations, but we are trying to be conservative when setting an upper bound on Λ_B.

Our results for the Landau pole are very close to the analytic estimate that one can perform at one loop

$$\mu_{\mathrm{LP}} = \exp\left[\frac{1}{\beta_{\mathrm{SM}} + \Delta\beta}\left(\frac{8\pi^2}{g^2(M_Z)} + \beta_{\mathrm{SM}}\log(M_Z) + \Delta\beta\log(M)\right)\right], \tag{5.7}$$

where β_{SM} is the one-loop $\mathrm{SU}(2)_L$ β-function in the SM and

$$\Delta\beta = \frac{4}{3}\left[T(n) + T(n-1)\right] \tag{5.8}$$

is the contributions from the new leptons, where $T(n)$ is the Dynkin index of a $\mathrm{SU}(2)_L$ representation of dimension n

$$T(n) = \frac{n(n-1)(n+1)}{12}. \tag{5.9}$$

Beyond this discussion, there is no other simple scaling argument that can exclude large $\mathrm{SU}(2)_L$ representations from our analysis. At large n and large M the coupling deviation scales as

$$\delta g_{hVV} \sim \frac{\alpha_V}{4\pi}\left(An^3\frac{g_{\mathrm{SM}}^2 y^2 v^2}{M^2} + Bn\frac{y^4 v^2}{M^2}\right), \tag{5.10}$$

with A and B both $\mathcal{O}(1)$ numbers. The n^3 term in the previous expression is expected on general grounds, since the VV two-point function scales as the Dynkin index $T(n) \sim n^3$ and the coupling deviation for small external momenta can be obtained from it through the Higgs low energy theorem [42–45], that reproduces exactly our results in the appropriate limit. To try to reduce the representations in our analysis we have to compare the scaling in eq. (5.10) to the scaling of the relevant RGEs,

$$16\pi^2\frac{dy}{dt} \sim ny^3, \qquad 16\pi^2\frac{d\lambda}{dt} \sim -ny^4 + n\lambda y^2. \tag{5.11}$$

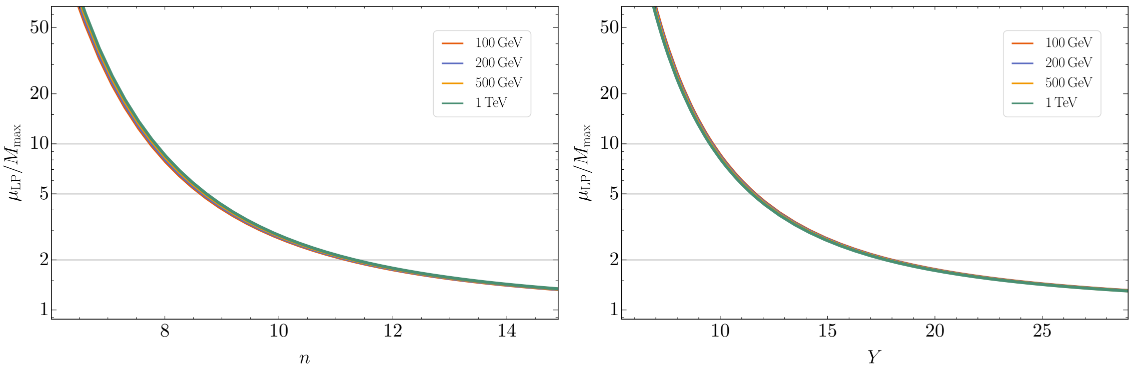

Figure 7. Location of the $SU(2)_L$ Landau pole (left panel) and $U(1)_Y$ Landau pole (right panel), normalized to the heaviest mass of the vector-like fermions $M_{\rm max}$, as a function of the dimension of their representations.

For example, in the case of a Higgs instability we can define

$$y_n \equiv n^{1/4} y\,, \tag{5.12}$$

that gives the same instability scale for any n and write the coupling deviation as

$$\delta g_{hVV} \sim \frac{\alpha_V}{4\pi}\left(A n^{5/2}\frac{g^2 y_n^2 v^2}{M^2} + B\frac{y_n^4 v^2}{M^2}\right), \tag{5.13}$$

so the numerator of the coupling deviation grows with n (if n is large enough that the first term dominates). However, the denominator also grows with n, due to the stronger collider (and indirect) constraints on M for representations with large charges. States in large-n representations have exotic charges and can decay to final states with effectively zero background. Additionally, the doublet+singlet model ($n=2$) is still very unconstrained by the LHC, with LEP setting the most stringent bounds (see section 6.2) and eq. (5.13) does not appropriately describe the coupling deviation in this case (i.e. we cannot take $M \gg v$). So we cannot reach a conclusion on the $SU(2)_L$ representations with the largest Λ_B from these simple arguments. The only case where we can say something definite is when the first problem with the theory is a $SU(2)_L$ Landau pole. In this case the important RGE is $dg/dt \sim n^3 g^2$ and we see by rescaling g that smaller representations give a bigger coupling deviation for a fixed scale of the Landau pole.

In summary our simple scaling arguments are inconclusive and in section 6 we consider both the doublet+singlet model and higher-dimensional $SU(2)_L$ representations. In practice we find that smaller representation in general give a bigger Λ_B, so that we can restrict our analysis to $n=4$ and below.

A similar reasoning holds for the hypercharge of the new fermions, taking now $T(Y) = Y^2$. To avoid low-energy Landau poles independent of δg_{hVV}, we consider only $Y \leq 5$ in what follows (see figure 7). This is more "aggressive" than our upper bound on $SU(2)_L$ representations, if we judge by the position of the gauge Landau pole in figure 7. However, when computing Λ_B, we find that gauge Landau poles dominate over other instabilities already for $Y=5$, as discussed in the next section.

As for $\mathrm{SU}(2)_L$, it is not possible to replicate the scaling arguments that led to consider only the lowest color representations, because there is no obvious rescaling of y, y^c that gives a fixed scale of the instability for any value of Y. The only exception arises if Y is so large that the first issue with the theory is the hypercharge Landau pole. In this case the one-loop RGE for the gauge coupling $dg_Y/dt \sim Y^2 g_Y^3$ suggests the rescaling $g'_Y = g_Y/Y$. In this regime there is no gain in going to large Y because $\delta g_{hZZ} \sim g_Y^2 Y^2 (V^3/M^2) \sim (g'_Y)^2 (V^3/M^2)$ and states with large Y are easier to detect, effectively giving a smaller δg_{hZZ} for a fixed scale of the Landau pole. However at the weak scale we can have $y, y^c \gg g_Y$ and for values of $0 < Y \leq 5$ one can be in a regime where the instability of the Higgs quartic or the Landau poles in y, y^c are not appreciably affected by increasing Y (that enters their running at two-loops in association with the comparatively small hypercharge gauge coupling), while δg_{hZZ} increases. For this reason in the following we show results for the scale of new bosons also in models where Y is large.

6 The scale of new bosons

6.1 Definition of Λ_B

We have seen in the previous sections that the models with only new vector-like leptons give the most conservative upper bound on the scale of new bosons Λ_B. Before turning to our results for Λ_B we comment briefly on how to compute it. If Λ_B comes from a loss of perturbativity, we fix it conventionally to be the scale where the coupling hits 4π. This is somewhat arbitrary given that we do not have complete control of the theory already at smaller values of the coupling. However, the running is fast when the Yukawas approach 4π, and changing the upper bound on the coupling does not appreciably affect Λ_B. For example, it was shown in [2] that reducing the threshold to $\sqrt{4\pi}$ does not qualitatively affect the result for Λ_B within the two-loop approximation for the RGEs that we employ also in this work. If a Higgs coupling deviation is measured it will be worth refining these results, but our choice $y \simeq 4\pi$ is good enough for our illustrative purposes.

The case where Λ_B is due to an instability in the Higgs potential is slightly more subtle. In this work we adopt a manifestly gauge invariant criterion to determine the scale where the Higgs potential becomes unstable. We compute at two loops the RGE evolution of the Higgs quartic λ and require a stable theory to satisfy $\lambda(\mu)^{-1} > -14.53 + 0.153 \log[\mathrm{GeV}/\mu]$, for any scale μ [316]. The physical meaning of this criterion has already been discussed quite extensively in the literature (see for instance [316, 317]). The scale Λ_B is then given by $\lambda(\Lambda_B)^{-1} = -14.53 + 0.153 \log[\mathrm{GeV}/\Lambda_B]$.

Intuitively the stability bound is on λ because the effective Higgs potential is well approximated by $V_{\text{eff}}(h) \simeq (\lambda_{\text{eff}}(h)/4) h^4$ when $h \gg v$. One can compute the bounce action for tunnelling from the SM vacuum ($h = v$) to a point h^* in the region $h \gg v$, where $V_{\text{eff}}(v) = V_{\text{eff}}(h^*)$. The bounce action depends on $\lambda(\mu)$ and the inequality that we use corresponds to a tunnelling rate equal to the lifetime of the universe [316]. This criterion is manifestly gauge invariant given that $\lambda(\mu)$ is a measurable, gauge-independent quantity and so is the tunnelling rate. However, this choice does not capture all the corrections to the Higgs effective potential at this order, i.e. $\lambda_{\text{eff}}(\mu) \neq \lambda(\mu)$.

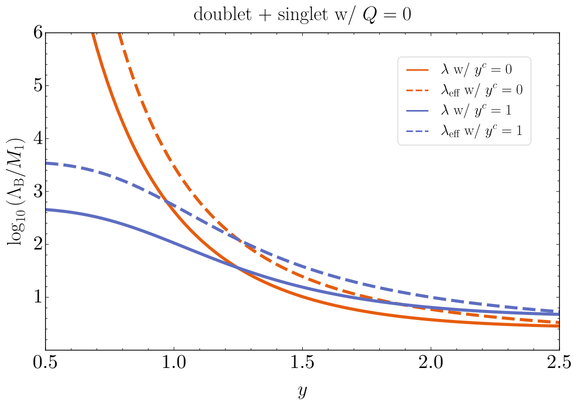

Figure 8. Instability scale of the Higgs potential Λ_B normalized to the lightest new fermion mass M_1 in the *doublet+singlet* theory introduced in section 3.2, as a function of one of its two Yukawa couplings (eq. (3.8)). The solid lines are obtained by running the two-loop RGEs of the Higgs quartic λ. The dashed lines are obtained by extracting the effective Higgs quartic λ_{eff} from the two-loop improved effective potential in Landau gauge.

In figure 8 we show that using $\lambda_{\text{eff}}(\mu)$, as computed from the two-loop improved effective potential[10] in Landau gauge, gives results for Λ_B that can differ from those obtained from $\lambda(\mu)$ by up to a factor of ten for the models considered in this work. However $\lambda_{\text{eff}}(\mu)$ is gauge-dependent. It is a well-known problem that V_{eff} is not gauge invariant [318]. In principle it is possible to extract gauge-invariant quantities from it and the metastability bound that we are interested in should be one of them. However to date we do not know of a way to obtain a gauge-independent result at fixed order in perturbation theory [319, 320]. The gauge dependence arises from electroweak corrections to $\lambda(\mu)$, it is numerically quite mild and often instability scales are quoted in Landau gauge [2, 317]. However, in this work we prefer the theoretically cleaner criterion of bounding $\lambda(\mu)$ and finding a gauge-independent bound.

6.2 Direct and indirect constraints on new fermions

The main goal of the paper is to set a (conservative) upper bound on the scale of new bosons, given an observed deviation in Higgs couplings to WW or ZZ. The lighter the new fermions can be for a fixed Yukawa, the larger the upper bound on Λ_B. In this section we compute current constraints, in the spirit of understanding how far we can go in this direction. Our goal is not to study in detail the constraints on the new fermions, we prefer to make conservative statements to understand how large Λ_B can be. A more thorough study of the bounds on the new fermions will be appropriate if a deviation in hWW or hZZ is discovered.

The strongest candidate for the maximization of Λ_B is the doublet+singlet model discussed in section 3.2, since the masses of its constituents are the hardest to constrain at the LHC. This theory is similar, but not identical, to a Higgsino-Bino system. They

[10]To compute the effective potential we use the same methodology described in detail in the appendix of [2] and obtain the same results.

would have been identical if we had considered the model with a Majorana fermion $N = N^c$ and identified $M_{\widetilde{B}} = M_N$, $M_L = \mu$, $g' v_u = yv$, $g' v_d = y^c v$. However the intuition from SUSY searches that this system is poorly constrained at the LHC holds also in our case.

First of all, it is useful to notice that our coupling deviation is mainly determined by the vector-like mass of the doublet M_L. Taking $M_N = 0$ does not appreciably increase $|\delta\mu_{hVV}|$ compared to $M_N = M_L$. To see this explicitly we integrate out the doublets when the singlets are light. As in section 3.2 we call the doublets ℓ, ℓ^c and the singlets n, n^c. At the weak scale we are left with

$$\mathcal{L} \supset i\left(\frac{|y^c|^2}{|M_L|^2}(Hn)^\dagger \bar{\sigma}^\mu D_\mu (Hn) + \frac{|y|^2}{|M_L|^2}(Hn^c)^\dagger \bar{\sigma}^\mu D_\mu (Hn^c)\right) + \left(\frac{yy^c (Hn)^\dagger H n^c}{M_L} + \text{h.c.}\right) + \mathcal{O}(1/M_L^3)\,. \tag{6.1}$$

From this Lagrangian we can estimate the coupling deviation coming purely from light singlets in the loop. We have

$$\frac{\delta g_{hVV}}{g_{hVV}^{\rm SM}} \sim yy^c |y^c|^2 |y|^2 \frac{v^5}{M_L^5}\,, \tag{6.2}$$

which is very subleading to the $\mathcal{O}(v^2/M_L^2)$ corrections that one obtains by integrating out the doublets at one loop. This simple exercise shows that the limit $M_N \ll M_L$ in eq. (3.8) is not relevant for us, if M_L is constrained to be larger than the weak scale.

This simplifies our analysis and allows us to get some intuition on the collider constraints on the model. Our theory contains two neutral particles with mass $M_{1,2}$ and one charged particle with mass M_L. To better understand the spectrum let us consider a few relevant limits where we can get some analytic intuition. We have discussed how a big hierarchy between M_L and M_N does not increase the coupling deviation. We can therefore start by considering $M_L = M_N$. In this limit we have

$$M_2 > M_L > M_1 \qquad \text{if} \qquad M_L > \frac{yy^c v}{\sqrt{2}(y + y^c)}\,, \tag{6.3}$$

and $M_1 > M_L$ in the opposite case. Even if we take $M_N = 0$ we have a parametrically similar conclusion

$$M_2 > M_L > M_1 \qquad \text{if} \qquad M_L > \frac{yy^c v}{\sqrt{2(y^2 + y^{c2})}}\,. \tag{6.4}$$

Taking $M_{L,N}$ small at fixed y, y^c maximizes the coupling deviation. However these results show that, as we make $M_{L,N}$ smaller, at some point we hit a configuration where our charged particle is the lightest of all.[11] The collider constraints on a stable $Q = 1$ particle are quite stringent. CMS gets $M_{\rm exp} \gtrsim 400$ GeV from DY pair production with $2.5\,\text{fb}^{-1}$ [321]. The ATLAS bound with $36\,\text{fb}^{-1}$ on the charged component of a $\mathrm{SU}(2)_L$ doublet is $M_{\rm exp} \gtrsim 840$ GeV [322].

[11] If we choose to include a relative sign and take for instance $y = -y^c$, the charged state of mass M_L is always the lightest in the spectrum.

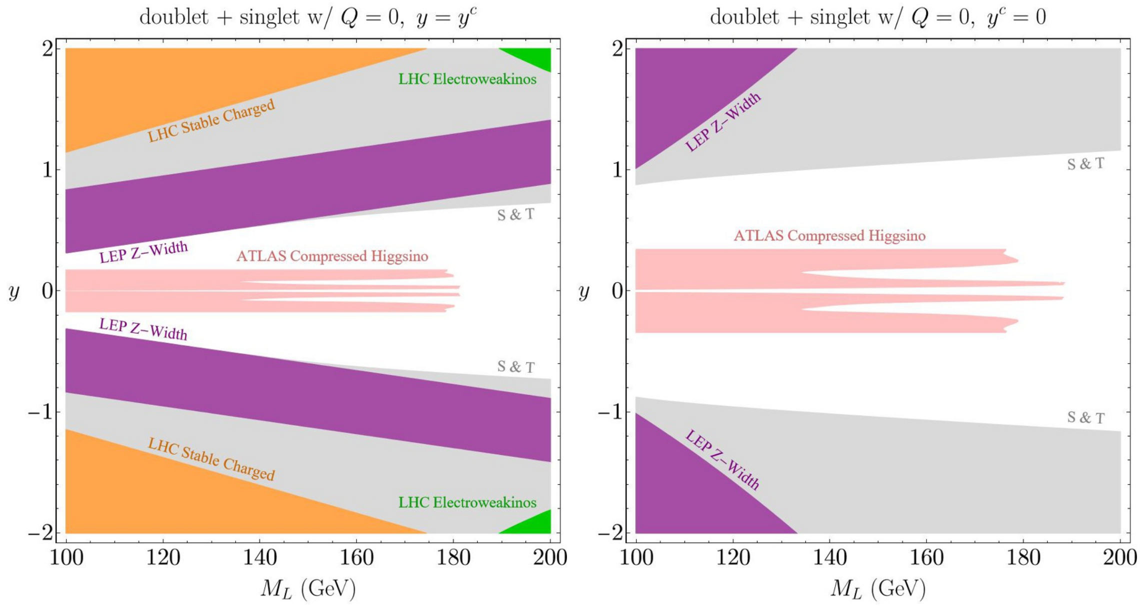

Figure 9. Experimental constraints on the doublet+singlet model from direct searches at colliders and electroweak precision measurements for $y = y^c$ (left panel) and $y^c = 0$ (right panel).

Searches for heavy stable charged particles at the LHC have very low background [321–323] and the same is true for long-lived particles decaying visibly to soft SM particles, so we might try to elude these bounds by adding a small mixing between ℓ, ℓ^c and the SM lepton doublets, just large enough to make the lightest new state decay promptly into SM leptons, but small enough to avoid all indirect constraints from flavor. However, after more than ten years of LHC data analysis, the bounds are stringent also in this case: $M \gtrsim 790\,\text{GeV}$ [324] for a new lepton doublet decaying predominantly to third generation leptons.[12] The small gap at low masses between the LEP bound and this analysis is bridged by CMS and ATLAS searches for charginos that exclude our model for masses around $M_L = 100\,\text{GeV}$ [325–328, 328, 328–337].

In summary if we want M_L as small as possible in order to maximize the coupling deviation we have to take it at least $M_L \gtrsim \mathcal{O}(yv)$. This is not a particularly stringent constraint, as we could take either $y = 0$ or $y^c = 0$ and always have a lightest neutral state in the spectrum (as one can see from eq.s (6.3) and (6.4)). However it is useful to keep it in mind when reading the plots, since the most important LHC bounds at low mass are just the condition $M_L > M_1$.

If we stick to the region of parameter space where the hierarchy between mass eigenstates is $M_2 > M_L > M_1$ the LHC does not strongly constrain the model, since we can have all three masses close enough to make the SM decay products hard to detect, without tuning our parameters. This, together with the relatively small cross sections of our doublet plus singlet model, makes current searches for vector-like leptons or electroweakinos

[12] Since rare events at the LHC are those containing electrons or muons and the τ leptonic branching ratio is smaller than its hadronic one, recasting this analysis would give a stronger constraint on decays to the first two generations.

insensitive to our new particles. In figure 9 we show a summary of existing constraints. The purple bands show where $M_1 < m_Z/2$, at odds with the LEP measurement of the Z-width [35]. The region in orange is where $M_L < M_1$ and we are excluded by stable charged particles searches at the LHC [321–323]. We see that this excluded region sets a bound on the largest coupling deviation that we can achieve, since parametrically it is similar to imposing $M_L \gtrsim yv$. The pink region is excluded by the ATLAS search for compressed Higgsinos [327] and the green one by their more general searches for electroweakinos [325–328, 328, 328–337].

The main qualitative message of figure 9 is that a light doublet+singlet system is still compatible with collider constraints. What really limits the maximal $|\delta\mu_{hVV}|$ are the bounds from electroweak precision measurements on oblique parameters. For $M_1 \geq 100\,\text{GeV}$, these are well captured by the deviations in the Peskin-Takeuchi S and T parameters [338] at $U = 0$. To implement the constraint and check the size of U, we use the following standard definitions [38],

$$\begin{aligned}
T &\equiv \frac{1}{\alpha}\left(\frac{\Pi^{\rm new}_{WW}(0)}{m_W^2} - \frac{\Pi^{\rm new}_{ZZ}(0)}{m_Z^2}\right)\,,\\
S &\equiv \frac{4c_W^2 s_W^2}{\alpha}\left(\frac{\Pi^{\rm new}_{ZZ}(m_Z^2) - \Pi^{\rm new}_{ZZ}(0)}{m_Z^2} - \frac{c_W^2 - s_W^2}{c_W^2 s_W^2}\frac{\Pi^{\rm new}_{\gamma Z}(m_Z^2)}{m_Z^2} - \frac{\Pi^{\rm new}_{\gamma\gamma}(m_Z^2)}{m_Z^2}\right)\,,\\
U &\equiv \frac{4c_W^2 s_W^2}{\alpha}\left(\frac{\Pi^{\rm new}_{WW}(m_W^2) - \Pi^{\rm new}_{WW}(0)}{m_W^2} - \frac{c_W}{s_W}\frac{\Pi^{\rm new}_{\gamma Z}(m_Z^2)}{m_Z^2} - \frac{\Pi^{\rm new}_{\gamma\gamma}(m_Z^2)}{m_Z^2}\right) - S\,,
\end{aligned} \tag{6.5}$$

and compare the calculation with the measured $S - T$ ellipse (95% CL) from the `Gfitter` collaboration [339]. The constraint is shown in gray in figure 9. For smaller masses, $M_1 \leq 100\,\text{GeV}$, the constraints from U and other parameters relevant for light new physics (V, W, X and Y [340, 341]) become important, but do not affect qualitatively the constraints that we show.[13] We leave to future work a more complete calculation of these bounds, in the hope that deviations in hWW and hZZ are discovered.

The situation for higher $\mathrm{SU}(2)_L$ representations or larger Y is different. In this case also LHC constraints have an important role to play. Searches for electroweakinos [325–328, 328, 328–337] can exclude our doublet+triplet model for $M_L \lesssim 240\,\text{GeV}$, for mass splittings as small as $8\,\text{GeV}$ (figure 16 in [327]). Hypercharge assignments that do not allow for a lightest neutral state (i.e. $Y \neq \pm(n-1)/2$ with n the dimension of the largest $\mathrm{SU}(2)_L$ representation) lead to much stronger constraints, comparable or more stringent than those discussed for a stable particle of $Q = 1$ from a $\mathrm{SU}(2)_L$ doublet. For higher $\mathrm{SU}(2)_L$ representations we take conservatively the exclusion on the doublet+triplet model as a benchmark for the lightest mass that we should show in the plots, but show the results also for larger values of M_1.

In practice, in the next section we include the constraints from EWPTs as shaded areas in the plots for Λ_B, showing the result also for regions that are excluded in the simplest models. We implement the collider constraints by limiting the values of M_1 (the lightest

[13]For these small masses, we used the more conservative $S - T$ ellipse (95% CL) from the `Gfitter` collaboration for $U \neq 0$ [342].

mass of the new fermions) that we show in the plots for Λ_B, but we always compute the value of Λ_B for multiple choices of M_1, also much above the most conservatives constraints.

6.3 Results

The HL-LHC is going to measure $\delta\mu_{hVV}$ at the 1.5% level in the WW, ZZ and $\gamma\gamma$ channels [21]. Future lepton colliders (ILC [22], CLIC [23], FCC-*ee* [24], CEPC [25], MuC [26, 27]), in particular FCC-*ee* and the muon collider, have the potential to reach a precision of $\simeq 0.13\%$ [28] comparable to that of LEP on Z couplings. In the plots we show lines corresponding to the 1σ sensitivity in the κ-framework for HL-LHC [21] and in the g^{eff}-framework of [28] for future colliders. These lines are meant mostly to guide the eye, if a deviation is ever found we will conduct a more thorough study (for instance map our results onto higher-dimensional operators and include in the analysis their correlation matrix).

We find that % level deviations measurable at HL-LHC cannot be generated by perturbative theories containing only new fermions. Future lepton colliders can probe deviations small enough to push Λ_B almost to the GUT scale, but only for extremely light new fermions, with masses $\simeq 50\,\text{GeV}$, within reach of HL-LHC. If the new fermions are heavier, we find that new bosons are kinematically within reach of future hadron colliders or at most around $100\,\text{TeV}$ if the new fermions are lighter than $150\,\text{GeV}$. Even if it is implicitly obvious, it is important to explicitly point out that new bosons must appear below Λ_B. For example $\Lambda_B \simeq 10\,M$, with M a typical mass for the new fermions, means that new bosons must already exist at the same scale as the new fermions, otherwise they do not have a chance to cure the instabilities of the running of the fermionic theory.

In the rest of this section we show plots of Λ_B as a function of $\delta\mu_{hVV}$ at fixed values of the lightest new fermions mass M_1. The values of M_1 that we plot reflect the discussion of experimental bounds in the previous section. For example, we mentioned how ATLAS and CMS exclude models almost identical to our doublet+triplet scenario up to $M_1 \simeq 240\,\text{GeV}$, even for mass splittings as small as $8\,\text{GeV}$. In our Λ_B plots for the doublet+triplet model we conservatively allow M_1 to go down to $200\,\text{GeV}$. We do this because a more thorough analysis of experimental constraints might unveil (tuned) regions of parameter space where the lower bounds on M_1 discussed in the previous section are relaxed. We also show larger values of M_1 in line with the bounds.

In the title of each plot of this section we specify the relation between y and y^c used to compute Λ_B. We choose their relative value to maximize the cutoff. In the vast majority of cases this corresponds to $y = \pm y^c$. At very light M_1 and large $\delta\mu$ also the $y^c = 0$ case becomes relevant.

To conclude this section we show the value of Λ_B as a function of deviations in $h\gamma\gamma$. All the fermionic theories that we consider induce a deviation measurable at HL-LHC, with the only exception of the doublet+singlet theory with a neutral singlet. We show these results for completeness, since they might be the first experimental sign of the vector-like leptons that we discuss in this work.

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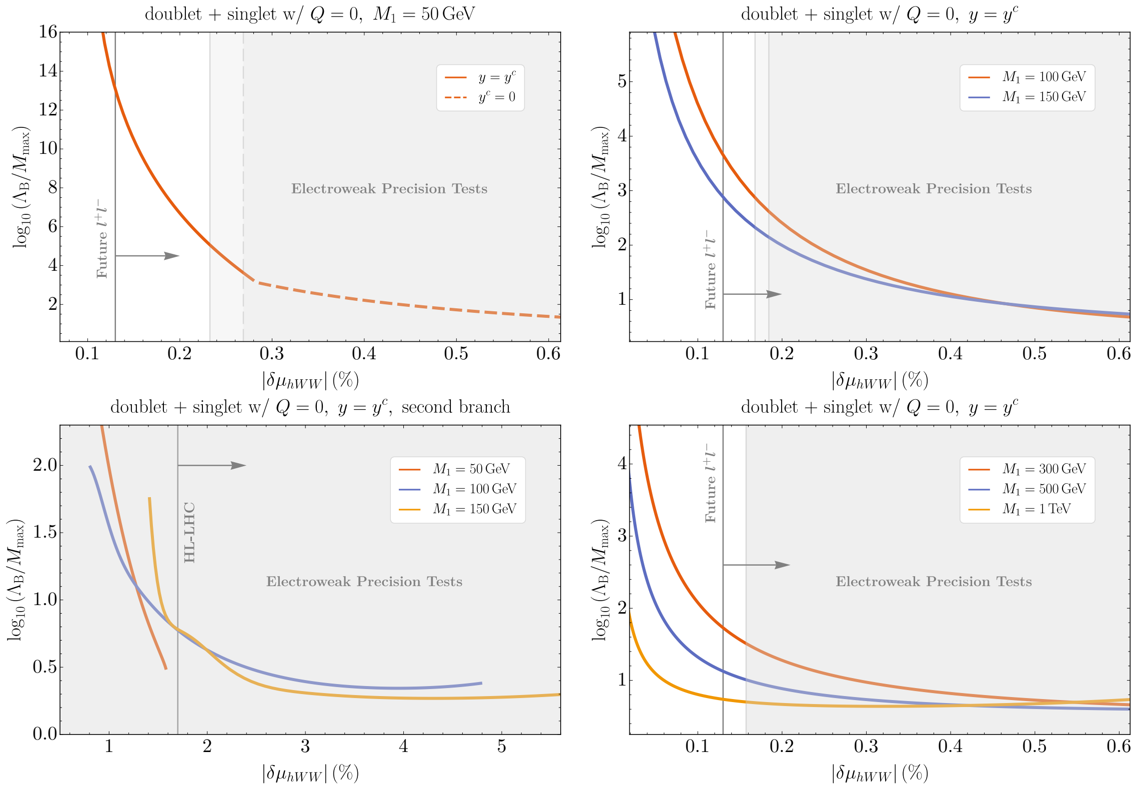

Figure 10. Instability scale of the RGEs Λ_B (i.e. upper bound on the scale of new bosons) as a function of relative hWW coupling deviation in the doublet+singlet model defined in section 3.2. $M_{\max}$ on the y-axis is the largest of the new fermions masses, while M_1 is the smallest one. The smallest values of M_1 in the plots reflect a conservative estimate of collider constraints on the models. The gray shaded areas represent the constraint from EWPTs. In the bottom right panel the constraint is the same for all values of M_1. In the top right panel the EWPT constraint at lowest $\delta\mu$ is on the line with lowest M_1. The constraint gets monotonically weaker at higher masses. The choice $y = y^c$ indicated in the title of each plot maximizes Λ_B, except for the first panel, where $y^c = 0$ gives a larger Λ_B for $\delta\mu_{hWW} \gtrsim 0.3\%$. The charge Q in the title refers to the singlet. The bottom left panel is the only figure in the paper where we picked the solution $M = M_1 - yv/\sqrt{2}$. For larger values of M_1 or smaller values of $\delta\mu$ the only existing solution is $M = M_1 + yv/\sqrt{2}$. We stop the $M_1 = 50\,\text{GeV}$ line at the value of y beyond which $\delta\mu$ starts decreasing. For such small M_1 different terms in the coupling deviation are comparable and start to cancel.

6.3.1 WW and ZZ

In figures 10 and 11 we show the value of Λ_B as a function of the relative hWW coupling deviation. In all cases we find that Λ_B is determined by the instability of the Higgs potential, which occurs before any Landau pole. The only exceptions are models with large hypercharge ($Q = 5$ for the singlet), where we see a scale of the instability independent of $\delta\mu$ at small values of the coupling deviation. In these regions the first problem with the theory is a hypercharge Landau pole.

We also find, regardless of the masses or representations of the new fermions, that coupling deviations measurable at HL-LHC are in tension with EWPTs (area shaded in

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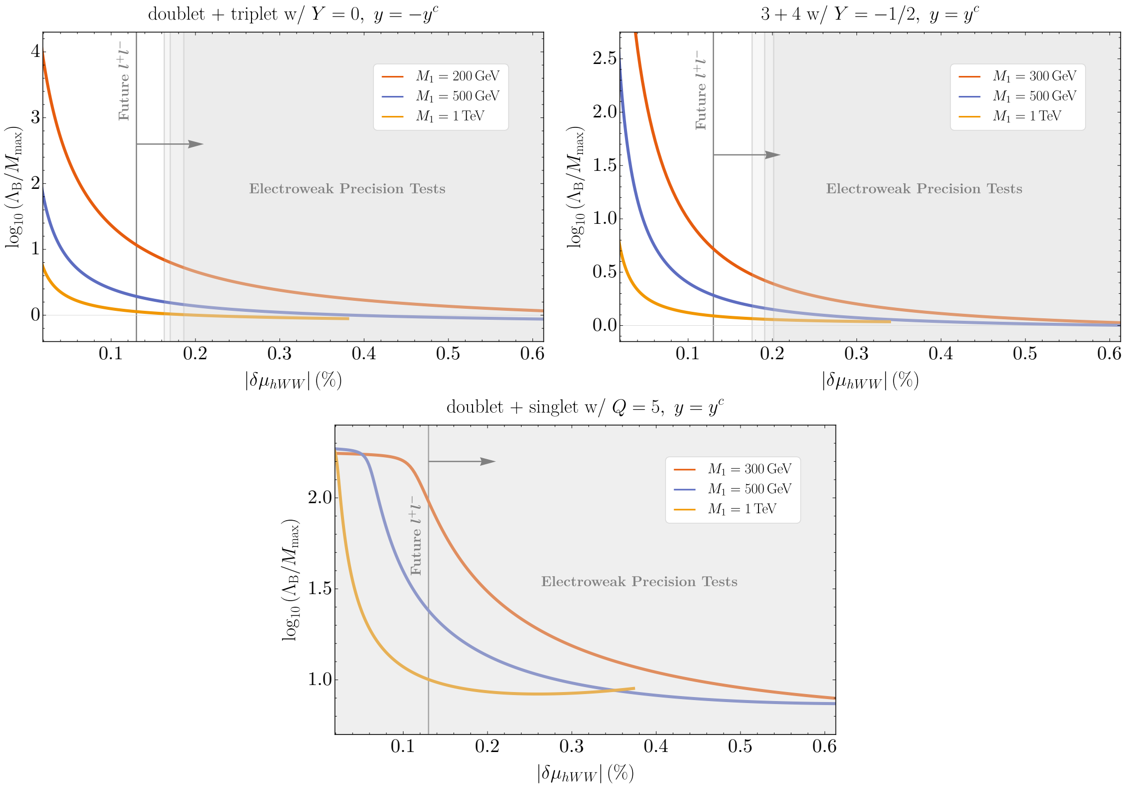

Figure 11. Instability scale of the RGEs Λ_B (i.e. upper bound on the scale of new bosons) as a function of relative hWW coupling deviation in the vector-like lepton models defined in section 3.2. $M_{\max}$ on the y-axis is the largest of the new fermions masses, while M_1 in the legend is the smallest one. The smallest values of M_1 in the plot reflect a conservative estimate of collider constraints on the models. The gray shaded areas represent the constraint from EWPTs. The EWPT constraint at lowest $\delta\mu$ is on the line with lowest M_1. The constraint gets monotonically weaker at higher masses. The choice $y = \pm y^c$ indicated in the title of each plot maximizes Λ_B.

gray). We stop the lines in the plot where we lose perturbative control of the theory (the Higgs quartic becomes rapidly large and negative after the threshold for vacuum decay). Larger values of M_1 correspond to larger Yukawas at fixed $\delta\mu$ and we stop the lines at lower values of $\delta\mu$.

The general message that all figures give is that we lose control of the theory (Λ_B a factor of a few to ten above the heaviest fermion mass $M_{\max}$) long before we can have a deviation large enough for HL-LHC. The only exception is the bottom left panel of figure 10. There we see that for $M_1 \le 150\,\text{GeV}$ we can generate a $\delta\mu$ within reach of HL-LHC, however in this case the new bosons must exist at the same scale as the fermions. This plot is made for $M = \frac{yv}{\sqrt{2}} - M_1$. This solution for M exists only at large enough y and small enough M_1. In all other plots we take $M = \frac{yv}{\sqrt{2}} + M_1$, because the other solution does not exist for Yukawa couplings $\lesssim \mathcal{O}(1)$. In the $M = \frac{yv}{\sqrt{2}} - M_1$ plot, labelled "second branch", we stop the $M_1 = 50\,\text{GeV}$ line at the value of y beyond which $\delta\mu$ starts decreasing. For such small M_1 different terms in the coupling deviation are comparable and start to cancel, giving much smaller values of Λ_B at fixed $\delta\mu$ compared to larger values of M_1.

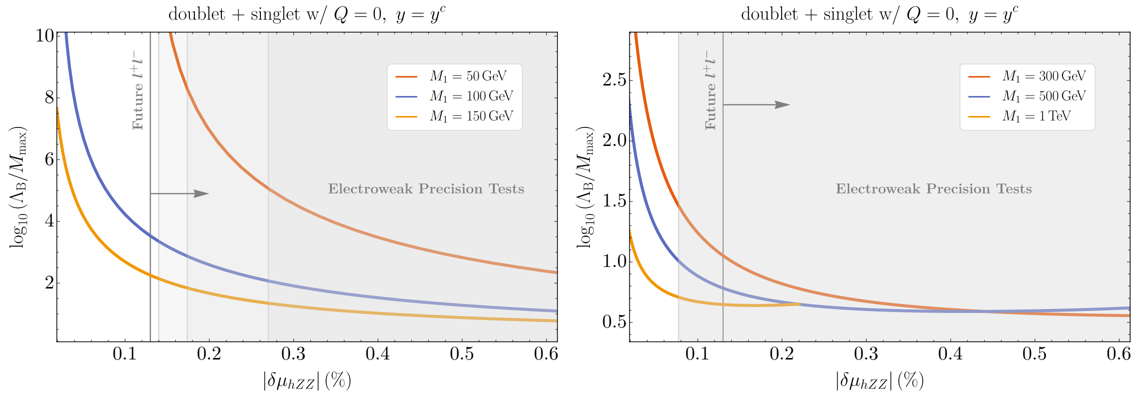

Figure 12. Instability scale of the RGEs Λ_B (i.e. upper bound on the scale of new bosons) as a function of relative hZZ coupling deviation in the doublet+singlet model defined in section 3.2. $M_{\rm max}$ on the y-axis is the largest of the new fermions masses, while M_1 is the smallest. The smallest values of M_1 reflect a conservative estimate of collider constraints on the models. The gray shaded areas represent the constraint from EWPTs. In the right panel the constraint is the same for all values of M_1. In the left panel the EWPT constraint at lowest $\delta\mu$ is on the line with lowest M_1. The constraint gets monotonically weaker at higher masses. The choice $y = y^c$ indicated in the title of each plot maximizes Λ_B.

We conclude that seeing a modification of hWW at HL-LHC requires new bosons at the scale where the deviation is generated. This conclusion is strengthened by the bound from EWPTs (gray shaded area in the figures) that excludes coupling deviations larger than the few permille level. Note that even if this result appears quite strong, it is not guaranteed that the new bosons can be produced at HL-LHC. For example, a strongly coupled ($g_* \simeq 4\pi$) new boson that can affect hWW at tree level, can be roughly as heavy as $M_B \simeq 4\pi v/\sqrt{\delta\mu_{hWW}} \simeq 13\,\text{TeV}\ \times(\sqrt{3\%/\delta\mu_{hWW}})$.

Coupling deviations as small as those detectable at future lepton colliders give in most cases $\Lambda_B \lesssim 100\,\text{TeV}$, implying the existence of new bosons well below this scale, potentially within reach of future hadron colliders. The main exception are light doublet+singlet fermions, where the new bosons can be as heavy as the GUT scale. However in this case we will detect vector-like fermions with masses $M_1 \lesssim 50\,\text{GeV}$ already at HL-LHC.

One last general point common to all Figures is that the EWPTs bound is rather insensitive to M_1 (in the bottom right panel of figure 10 it is even the same for the three values of M_1 in the figure). This is due to the fact that the bound is on almost the same combination of couplings and masses that enter $\delta\mu$, so at fixed $\delta\mu$ it is almost independent of M_1. This is illustrated schematically in figure 2.

It is interesting to compare figure 10, where we show Λ_B for the doublet+singlet model with a neutral singlet, and figure 11, where we show Λ_B for higher $SU(2)_L$ representations and hypercharges. Not surprisingly, scenarios where the new fermions can be lighter (namely the doublet+singlet model) have the largest value of Λ_B for a given coupling deviation. This can be seen by comparing the first three panels of figure 10 with all other subfigures.

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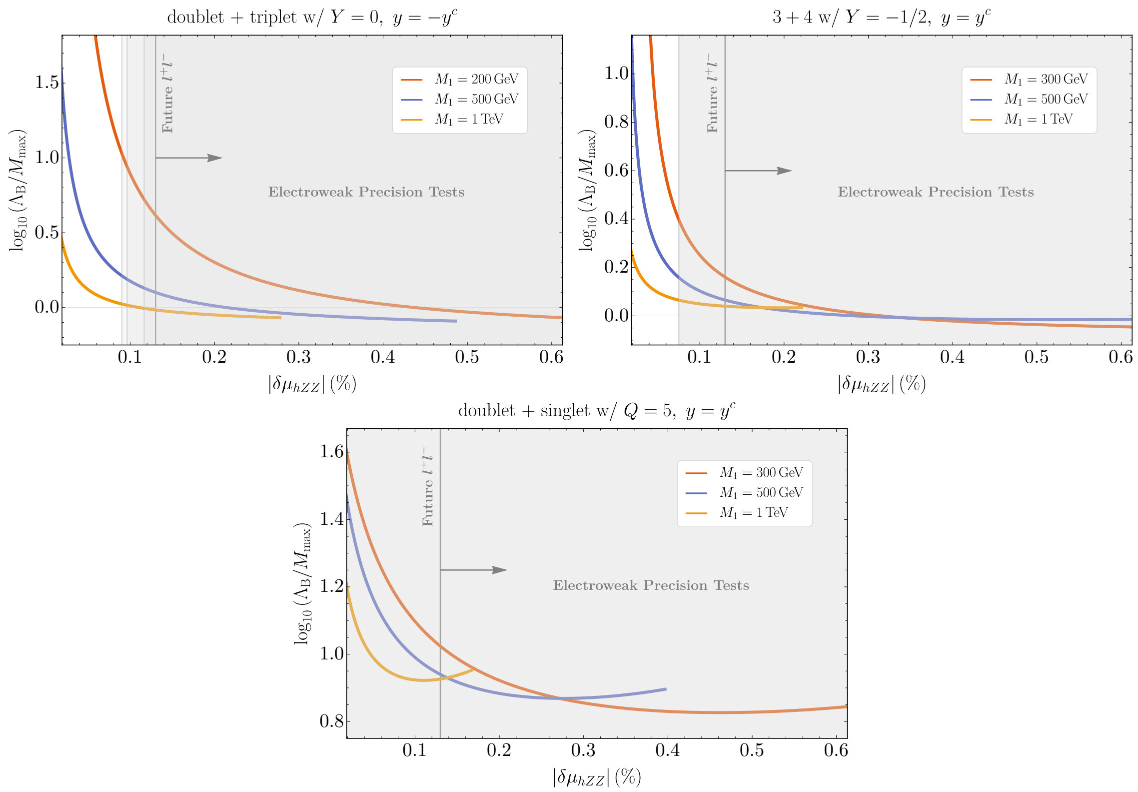

Figure 13. Instability scale of the RGEs Λ_B (i.e. upper bound on the scale of new bosons) as a function of relative hZZ coupling deviation in the vector-like lepton models defined in section 3.2. $M_{\max}$ on the y-axis is the largest of the new fermions masses, while M_1 is the smallest. The smallest values of M_1 reflect a conservative estimate of collider constraints on the models. The gray shaded areas represent the constraint from EWPTs. In the right panel the constraint is the same for all values of M_1. In the left panel the EWPT constraint at lowest $\delta\mu$ is on the line with lowest M_1. The constraint gets monotonically weaker at higher masses. The choice $y = \pm y^c$ indicated in the title of each plot maximizes Λ_B.

What was not completely obvious from our simple arguments in section 5.2 is that at fixed M_1 (lightest new fermion mass) the value of Λ_B does not vary greatly between different representations. This emerges from the comparison of the bottom right panel of figure 10 to the three panels of figure 11. In these subfigures the y^4 term in eq. (5.10) dominates the coupling deviation and we find the instability of the Higgs potential to determine Λ_B in all figures (except at $Q = 5$ and small $\delta\mu$ where the hypercharge Landau pole dominates). The rescaling $y^{(c)} \to n^{1/4} y_n^{(c)}$ makes both the instability of the Higgs potential and the coupling deviation at fixed masses independent of n. We do not show Λ_B for the last few fermionic theories that barely avoid a low energy Landau pole (i.e. a vector-like 5 of $\mathrm{SU}(2)_L$ plus a vector-like 4, and above) since it is even smaller than that shown in figure 11.

To conclude the discussion of hWW it is worth commenting on the regions where the cutoffs slightly increase and where it appears to be independent of $\delta\mu$. The latter case corresponds to a Landau pole in the hypercharge. This shows that already at $Y = 5$ this can be the dominant effect. The increase at large $\delta\mu$ is close to where we lose control of

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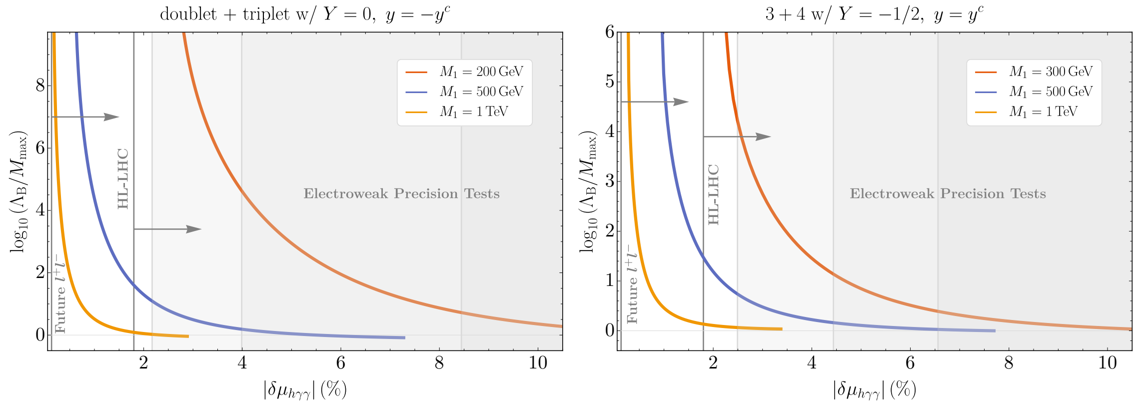

Figure 14. Instability scale of the RGEs Λ_B (i.e. upper bound on the scale of new bosons) as a function of relative $h\gamma\gamma$ coupling deviation in the vector-like lepton models defined in section 3.2. $M_{\max}$ on the y-axis is the largest of the new fermions masses, while M_1 is the smallest. The smallest values of M_1 reflect a conservative estimate of collider constraints on the models. The gray shaded areas represent the constraint from EWPTs. The EWPT constraint at lowest $\delta\mu$ is on the line with lowest M_1. The constraint gets monotonically weaker at higher masses. The choice $y = \pm y^c$ indicated in the title of each plot maximizes Λ_B.

the theory, but it is still a perturbative effect. It arises from the large threshold correction to λ at the matching scale between SM and SM+new fermions, which is proportional to y^4 and positive.

The case of hZZ is illustrated in figures 12 and 13 and is very similar to hWW. All the qualitative statements made for hWW hold also in this case, with slightly lower Λ_B for any given $\delta\mu$. New bosons are responsible for any deviation visible at HL-LHC and even permille level deviations require light new bosons. The behavior of the Λ_B vs $\delta\mu$ curves is explained by the same arguments given for hWW.

6.3.2 $\gamma\gamma$

The case of $h\gamma\gamma$ is much simpler. The result is automatically finite, since the vertex does not exist in the SM at tree-level, and we can write the coupling deviation for general n and Y at large $M_L = M_N \gg v$ in a compact form,

$$\delta\mu_{h\gamma\gamma} = \frac{y y^c v^2}{36 M_L^2 A_{\rm SM}} (-1)^n n \left[(n-1)(n+1) + 4(n+1)Y + 12Y^2 \right], \tag{6.6}$$

with $A_{\rm SM} \simeq 3.3$. Any fermionic theory in the previous subsection, different than the doublet+singlet case ($n = 2$, $Y = -1/2$), gives a deviation also in $h\gamma\gamma$. Since this coupling exist in the SM only at loop level, the relative deviation is much larger than the hWW or hZZ case and potentially visible at HL-LHC. If one believes the fermionic theories in this paper, their first experimental manifestation might be $h\gamma\gamma$, as shown in figure 14. This is the reason why we mention this coupling deviation, but $h\gamma\gamma$ was already discussed in [1, 2] and we do not have anything qualitative to add. As discussed in the first two sections, the fermionic theories in this paper prove by contradiction that observable deviations in hWW

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or hZZ require light new bosons. We do not aim at giving here a detailed phenomenological treatment of vector-like fermion theories besides our model-independent statement on hWW and hZZ.

7 Outlook

We have discussed how observing a deviation in Higgs couplings to WW and ZZ gives information on the scale where new bosons must appear. We considered theories containing only new fermions and showed that a measurable deviation in hWW or hZZ requires large Yukawa couplings that destabilize the Higgs potential. In general, these theories must be completed very close to the scale of the new fermions to avoid a rapid decay of the SM vacuum. Phenomenologically, our most interesting result is that any measurable deviation at HL-LHC requires such a small scale for new bosons that it cannot be generated only by new fermions. Deviations measurable at future lepton colliders either require very light new fermions, with masses $M_\psi \lesssim 150\,\mathrm{GeV}$, or new bosons roughly below 100 TeV.

In this work we have computed a scale Λ_B associated to the instability of the Higgs potential which gives a very conservative upper bound on where new bosons must appear. One natural way to improve this analysis would be to consider explicit new models where the instability is lifted and obtain a more precise determination of the actual masses of the new bosons. It is hard to imagine a way to do this model-independently, but we leave this line of inquiry to other researchers and our future selves, hoping that they can outsmart our present selves.

Another possible future direction consists in refining our treatment of direct and indirect constraints on the new fermions. The scale of new bosons Λ_B is sensitive to the lightest new fermions masses and here we did not go beyond a series of rough, but conservative estimates, thus obtaining an upper bound on Λ_B. We find that it will be interesting to further explore this point in case a credible deviation is measured in hWW or hZZ.

A third possibility for future work is to extend our discussion of the $h\gamma\gamma$ coupling. To a large extent it was treated in [1, 2], but there we did not systematically map all possible representations of the new fermions and we did not get a model-independent upper bound on Λ_B. Note, however, that the results in this paper suggest that the lowest fermions representations (already considered in [1, 2]) will dominate the upper bound.

To conclude, hWW and hZZ couplings are a key observable for the future run of the LHC. Any deviation from the SM prediction signals the presence of light new bosons. If we are lucky, this could either be the first sign of the long awaited symmetry explaining m_h^2 or a definitive sign of unnaturalness.

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A Fermionic low energy theories

In [36] a thorough survey of new heavy fermions that are: 1) experimentally still allowed 2) anomaly-free and 3) can affect Higgs couplings to SM particles was conducted. In this section we review what representations in [36] are relevant for us (i.e. can give an observable deviation in hZZ or hWW).

A.1 One fermion

As discussed in more detail in [36], there are two fermion representations that taken in isolation are consistent with the three conditions stated above, a singlet of all SM gauge groups

$$N = (1,1)_0 \tag{A.1}$$

and a $\mathrm{SU}(2)_L$ triplet

$$\Sigma = (1,3)_0 \,. \tag{A.2}$$

They can both have a Majorana mass and couple to the Higgs via a SM lepton ℓ. The relevant part of the Lagrangian for the singlet reads

$$\mathcal{L}_N = -\frac{M_N}{2} N^2 - y \ell H N \,. \tag{A.3}$$

This gives two neutral mass eigenstates. If $M_N \ll yv$ they are mixed at $\mathcal{O}(1)$ which is excluded by searches for new sterile neutrinos. The opposite limit $M_N \gg yv$ is phenomenologically viable, giving two eigenstates

$$m_{\nu_l} \simeq \frac{y^2 v^2}{2M_N} \,, \qquad m_{\nu_h} \simeq M_N \,, \tag{A.4}$$

mixed at $\mathcal{O}(y^2 v^2/M_N^2)$. Given upper bounds on neutrino masses, the SM-like neutrino must be ν_l, $m_{\nu_l} \lesssim \mathrm{eV}$. The effects of the new state on SM couplings are suppressed by at least m_{ν_l}/v and for our purposes are unobservably small.

A similar reasoning can be followed to conclude that also Σ does not produce observable Higgs coupling deviations, since its neutral component behaves exactly like N

$$\mathcal{L}_\Sigma = -\frac{m_\Sigma}{2}\Sigma^2 - y\ell H \Sigma = -\frac{m_\Sigma}{2}\Sigma^2 - y\nu_l \frac{h+v}{\sqrt{2}}\Sigma^0 + \dots \,. \tag{A.5}$$

Considering the charged components can only strengthen the bounds on m_Σ.

A.2 Two fermions

Any copy of a SM fermion with a vector-like partner satisfies the three conditions stated above. Also any of the exotic leptons (plus a vector-like partner) listed here:

$$\begin{aligned} &\Lambda = (1,2)_{-3/2}, \qquad \Delta = (1,3)_{-1}, \qquad X_T = (3,2)_{7/6}, \qquad Y_B = (3,2)_{-5/6}, \\ &X_Q = (3,3)_{2/3}, \qquad Y_Q = (3,3)_{-1/3}, \end{aligned} \tag{A.6}$$

is a viable possibility. In practice we have listed all fermions that can couple with an existing SM fermion through a Yukawa coupling. Let us first consider the leptons, for concreteness a vector-like partner of the electron described by the two left-handed spinors E, E^c. The relevant part of the Lagrangian reads

$$\mathcal{L}_E = -m_E E E^c - y \ell_\tau H E^c - \mu \tau^c E + \text{h.c.}. \tag{A.7}$$

We coupled E, E^c just to the τ to avoid stringent constraints from flavor changing processes. Different choices are possible, but only strengthen our conclusion below (i.e. that these representations should not be considered in our work). We could immediately discard this option by noting that $H\tau\tau$ and $Z\tau\tau$ are affected at tree-level while the couplings of interest to us are modified at loop level. However let us be more precise. Current constraints on charged particles at the LHC are quite stringent, placing them firmly above the weak scale. Therefore it is sensible to integrate out E and E^c to understand what are the leading effects on SM coupling deviations. At tree level the equations of motion read

$$\begin{aligned} E &= -\frac{y l_\tau H}{m_E} + \frac{\mu^\dagger}{|m_E|^2}(D_\mu^{E^c} \tau^{c\dagger})\bar{\sigma}^\mu , \\ E^c &= -\frac{\mu \tau^c}{m_E} + \frac{y^\dagger}{|m_E|^2}(D_\mu^E l_\tau^\dagger)\bar{\sigma}^\mu . \end{aligned} \tag{A.8}$$

The Lagrangian in eq. (A.7) including all terms up to $\mathcal{O}(1/m_E^2)$ only gives a shift to the τ Yukawa coupling

$$\mathcal{L}_E = \frac{\mu y}{m_E} \ell_\tau H \tau^c + \text{h.c.} + \mathcal{O}(1/m_E^4). \tag{A.9}$$

This contribution to y_τ cannot be observed as a Higgs coupling deviation because the τ mass and its coupling are affected in the same way. However, we are forced to impose

$$\frac{\mu y}{m_E} \lesssim y_\tau^{\text{exp.}} \simeq 0.01 . \tag{A.10}$$

The kinetic terms of the two heavy leptons give a more interesting result

$$\begin{aligned} \mathcal{L}_E^{\text{kin}} &= \frac{i|y|^2}{4|m_E|^2}\left[(H^\dagger D_\mu H)(l_\tau^\dagger \bar{\sigma}^\mu l_\tau) + (H^\dagger D_\mu \vec{\sigma} H)\cdot(l_\tau^\dagger \bar{\sigma}^\mu \vec{\sigma} l_\tau)\right] - y_\tau \frac{|y|^2 |H|^2}{2|m_E|^2} l_\tau H \tau^c \\ &\quad + i\frac{|\mu|^2}{|m_E|^2} \tau^{c\dagger} \bar{\sigma}^\mu D_\mu^{E^c} \tau^c . \end{aligned} \tag{A.11}$$

We can read in appendix B of [2] (eq. B.4) how these operators affect Higgs and Z couplings (we call g_A and g_V the axial and vector coupling of the τ to the Z). The result is

$$\frac{\delta g_{h\tau\tau}}{g_{h\tau\tau}} = \frac{|y|^2 v^2}{2|m_E|^2}, \qquad \frac{\delta g_A}{g_A} = -\frac{|y|^2 v^2}{4|m_E|^2} - \frac{|\mu|^2}{|m_E|^2}, \qquad \frac{\delta g_V}{g_V} = -\frac{|y|^2 v^2}{4|m_E|^2} + \frac{|\mu|^2}{|m_E|^2}. \tag{A.12}$$

The hZZ and hWW coupling deviations arise at loop level and at leading order we have

$$\frac{\delta g_{hVV}}{g_{hVV}} \sim \frac{|y|^2 v^2}{16\pi^2 |m_E|^2} \max[1, y^2/g^2]. \tag{A.13}$$

There are also terms proportional to $y\mu$ or μ^2, but are suppressed by the insertion of the SM τ Yukawa coupling. We see immediately that while we can tune δg_V to be small without affecting δg_{hVV}, we can't have a large δg_{hVV} without an even larger Z-coupling deviation δg_A. The maximal allowed value for δg_{hVV} given a bound on δg_A is obtained at $\mu = 0$

$$\left.\frac{\delta g_{hVV}}{g_{hVV}}\right|_{\max} \sim \frac{1}{\pi^2}\left(\frac{\delta g_A}{g_A}\right)_{\exp}, \tag{A.14}$$

where we took $y \sim 2$, the largest value compatible with not having a Landau pole or an instability right above the new fermions' masses. Z couplings are constrained at the permille level by LEP [41, 343], so this deviation is much smaller than what we can have with the same Yukawa if we introduce 3 or 4 vector-like fermions (see for example section 3 in the main body of the paper), so for a given coupling deviation we get a smaller cutoff and we do not need to consider this case.

B Large $\mathrm{SU}(2)_L$ representations

We can treat the $\mathrm{SU}(2)_L$ representations in the main body of the paper in two equivalent ways. The most common choice in particle physics is to think about a representation of dimension n as a vector with n components

$$L = \begin{pmatrix} L_{(n-1)/2} \\ L_{(n-2)/2} \\ \cdots \\ L_m \\ \cdots \\ L_{-(n-1)/2} \end{pmatrix}, \tag{B.1}$$

labelled by an index m. We can identify each component with the state $|j,m\rangle$ where the equivalent of the total spin is $j = (n-1)/2$. It is then straightforward to write the Lagrangian for the components of L and L^c,

$$\begin{aligned}
\mathcal{L}_{\text{free}} &= \sum_{m=-(n-1)/2}^{(n-1)/2} \left[L_m^\dagger \bar\sigma^\mu \partial_\mu L_m + L_m^{c\dagger}\bar\sigma^\mu \partial_\mu L_m^c - M_L (-1)^{\frac{n-1}{2}-m} L^c_{-m} L_m\right], \\
\mathcal{L}_{\text{gauge}} &= -\frac{g}{\sqrt{2}} \sum_{m=-(n-1)/2}^{(n-1)/2} \left[\sqrt{\frac{n}{2}\left(\frac{n}{2}-1\right) - m(m+1)} L^\dagger_{m+1}\bar\sigma^\mu L_m W^+_\mu + \text{h.c.}\right] \\
&\quad - \sum_{m=-(n-1)/2}^{(n-1)/2} \left[e(Y+m) L_m^\dagger \bar\sigma^\mu L_m A_\mu - \frac{g}{c_W}(m-(Y+m)s_W^2) L_m^\dagger \bar\sigma^\mu L_m Z_\mu\right], \\
\mathcal{L}_{\text{int}} &= -\sum_{m=-(n-2)/2}^{(n-2)/2} \frac{(-1)^{\frac{n-2}{2}-m}}{\sqrt{2}} \left[\sqrt{1-\frac{2m-1}{n-1}}\left(y L_{1/2-m} H^0 N^c_m + y^c L^c_{1/2-m} H^{+*} N_m\right)\right. \\
&\quad \left. + \sqrt{1+\frac{2m+1}{n-1}}\left(y^c L^c_{-1/2-m} H^{0*} N_m + y L_{-1/2-m} H^+ N^c_m\right)\right].
\end{aligned} \tag{B.2}$$

To write $\mathcal{L}_{\rm free}, \mathcal{L}_{\rm gauge}$ and $\mathcal{L}_{\rm int}$ we used the usual Clebsh-Gordan decomposition of the direct product of SU(2) representations and the well-known form of the generators of the group for a representation of dimension n [344]. We have omitted the gauge interactions of L^c because they can easily be deduced from $\mathcal{L}_{\rm gauge}$.

A second option is to treat the representation of dimension n as a symmetric tensor with $n-1$ indexes,[14] each transforming as a fundamental of SU(2)

$$L = L^{i_1 i_2 \dots i_{n-1}}\,. \tag{B.3}$$

This is the most natural way to treat L if one builds up higher dimensional representations from the direct product of lower dimensional ones. Under a SU(2) transformation the tensor L is rotated to L'

$$(L')^{i_1 i_2 \dots i_{n-1}} = U^{i_1 j_1} U^{i_2 j_2} \dots U^{i_{n-1} j_{n-1}} L^{j_1 j_2 \dots j_{n-1}}\,, \qquad U = e^{i\vec{\alpha}\cdot\frac{\vec{\sigma}}{2}}\,. \tag{B.4}$$

If we adopt the convention of lowering the indexes of the fundamental representation using the totally antisymmetric tensor ϵ_{ij}

$$L_i = \epsilon_{ij} L^j\,, \qquad \epsilon = \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}, \tag{B.5}$$

contractions of upper and lower indexes are SU(2)-invariant and we can write $\mathcal{L}_{\rm free}$ and $\mathcal{L}_{\rm int}$ in a compact form

$$\begin{aligned} \mathcal{L}_{\rm free} &= \sum_{(i_1 i_2 \dots i_{n-1})} \Big[(L^\dagger)^{i_1 i_2 \dots i_{n-1}} \overline{\sigma}^\mu \partial_\mu L^{i_1 i_2 \dots i_{n-1}} \\ &\quad + (L^{c\dagger})^{i_1 i_2 \dots i_{n-1}} \overline{\sigma}^\mu \partial_\mu (L^c)^{i_1 i_2 \dots i_{n-1}} - M_L L^c_{i_1 i_2 \dots i_{n-1}} L^{i_1 i_2 \dots i_{n-1}} \Big], \\ \mathcal{L}_{\rm int} &= - \sum_{(i_1 i_2 \dots i_{n-1})} \Big[y L_{i_1 i_2 \dots i_{n-1}} H^{i_1} (N^c)^{i_2 \dots i_{n-1}} + y^c L^c_{i_1 i_2 \dots i_{n-1}} H^\dagger_{i_1} N^{i_2 \dots i_{n-1}} \Big]. \end{aligned} \tag{B.6}$$

Note that the sums extend only over the independent components of the tensors. All tensors are fully symmetric in the exchange of any pair of indexes, so for example we are not including in the sum both $L^{12 i_3 i_4 \dots i_{n-1}}$ and $L^{21 i_3 i_4 \dots i_{n-1}}$ otherwise we would be double-counting.

It is possible to write down also gauge interactions in this notation, but the resulting expressions are cumbersome and not particularly illuminating. It is much better to notice that there is a simple relation between $L^{j_1 j_2 \dots j_{n-1}}$ and our original vector. Let us call $L^{\tilde{k}}$ the component of $L^{j_1 j_2 \dots j_{n-1}}$ with k indexes equal to 1 and $n-k-1$ equal to 2. Its electric charge is

$$Q L^{\tilde{k}} = \left(Y + \frac{k}{2} - \frac{n-1-k}{2} \right) L^{\tilde{k}}\,, \tag{B.7}$$

[14] The dimension d of a symmetric tensor with k indexes running from 1 to N is $d = \begin{pmatrix} N+k-1 \\ k \end{pmatrix}$.

so we can identify $L^{\tilde{k}} = L_{k-(n-1)/2}$, where $L_{k-(n-1)/2}$ is one of the components of the vector in eq. (B.1). We can then write the action of the generators on $L^{i_1 i_2 \dots i_{n-1}}$ starting from the well-known expressions of their action on vector components

$$\begin{aligned} T^{\pm} L_m &= \sqrt{\frac{n}{2}\left(\frac{n-1}{2}\right) - m(m\pm 1)} L_{m\pm 1}\,, \\ T^3 L_m &= m L_m\,. \end{aligned} \tag{B.8}$$

The results in this work were obtained both with the tensor and with the vector notation and verified to be the same.

C Auxiliary functions

The three-body decay widths in the main body of the text and the corresponding coupling deviations have been expressed in terms of the functions listed in this appendix.

C.1 WW

In the case of hWW the functions introduced by the new leptons at $\mathcal{O}(1/M^2)$ in the vector-like mass can be written explicitly:

$$\begin{aligned} R_W(x) &= -180 - \frac{6}{x^2} + 141x + \frac{45}{x} + \left(-36x - \frac{9}{x} + 54\right)\log(x) \\ &\quad + \frac{\left(-720x - \frac{36}{x} + 288\right)\left[\tan^{-1}\left(\frac{2\sqrt{x}-1}{\sqrt{4x-1}}\right) - \cot^{-1}\left(\sqrt{4x-1}\right)\right]}{\sqrt{4x-1}}\,, \end{aligned} \tag{C.1}$$

$$\begin{aligned} P_W^{(1)}(x) &= -320x^{5/2} + 600x^{3/2} - \frac{6}{x^{3/2}} + \left(72x^{5/2} - 252x^{3/2} + 90\sqrt{x} - \frac{12}{\sqrt{x}}\right)\log(x) \\ &\quad + \frac{\left(2016x^{5/2} - 1632x^{3/2} + 456\sqrt{x} - \frac{48}{\sqrt{x}}\right)\left[\tan^{-1}\left(\frac{2\sqrt{x}-1}{\sqrt{4x-1}}\right) - \cot^{-1}\left(\sqrt{4x-1}\right)\right]}{\sqrt{4x-1}} \\ &\quad - 342\sqrt{x} + \frac{68}{\sqrt{x}}\,, \end{aligned} \tag{C.2}$$

$$\begin{aligned} P_W^{(2)}(x) &= -162\sqrt{x} + \frac{68}{\sqrt{x}} + 4x^{5/2} + 96x^{3/2} - \frac{6}{x^{3/2}} + \left(-36x^{3/2} + 54\sqrt{x} - \frac{12}{\sqrt{x}}\right)\log(x) \\ &\quad + \frac{\left(-480x^{3/2} + 312\sqrt{x} - \frac{48}{\sqrt{x}}\right)\left[\tan^{-1}\left(\frac{2\sqrt{x}-1}{\sqrt{4x-1}}\right) - \cot^{-1}\left(\sqrt{4x-1}\right)\right]}{\sqrt{4x-1}}\,, \end{aligned} \tag{C.3}$$

$$\begin{aligned} P_W^{(3)}(x) &= 640x^{5/2} - 918x^{3/2} + \left(-144x^{5/2} + 432x^{3/2} - 72\sqrt{x} + \frac{6}{\sqrt{x}}\right)\log(x) \\ &\quad + \frac{\left(-4032x^{5/2} + 1824x^{3/2} - 336\sqrt{x} + \frac{24}{\sqrt{x}}\right)\left[\tan^{-1}\left(\frac{2\sqrt{x}-1}{\sqrt{4x-1}}\right) - \cot^{-1}\left(\sqrt{4x-1}\right)\right]}{\sqrt{4x-1}} \\ &\quad + 324\sqrt{x} - \frac{46}{\sqrt{x}}\,. \end{aligned} \tag{C.4}$$

C.2 ZZ

In the case of hZZ the dimensionless functions in section 4 are best expressed in terms of the following integrals:

$$f_p = \int_{4\frac{m^2}{m_h^2}}^{(1-z^{1/2})^2} dy \frac{y^p}{(z-y)^2} \sqrt{1-2(z+y)+(z-y)^2}\,, \qquad z = \frac{m_Z^2}{m_h^2}\,. \tag{C.5}$$

For the $h \to Z^* Z \to e^+ e^- Z$ width we have the following auxiliary functions:

$$R_T(z) = \frac{1}{6}\left[f_2 + 2(5z-1)f_1 + (1-z)^2 f_0 + 2\frac{m^2}{m_h^2}(1-z)^2 f_{-1}\right], \tag{C.6}$$

$$P_Z^{(1)}(z) = \frac{1}{6}\left[-f_3 + (3-11z)f_2 + (1-z)(11z-3)f_1 + (1-z)^3 f_0 - 2\frac{m^2}{m_h^2}(1-z)^3 f_{-1}\right], \tag{C.7}$$

$$P_Z^{(2)}(z) = \frac{1}{6}\left[-f_3 + (3+z)f_2 + (z-1)(3+z)f_1 + (1-z)^3 f_0 + 2\frac{m^2}{m_h^2}(1-z)^3 f_{-1}\right], \tag{C.8}$$

$$P_Z^{(3)}(z) = \frac{1}{3}\left[f_3 + (11z-2)f_2 + \left(1-4z+11z^2\right) f_1 + z(1-z)^2 f_0 + 2\frac{m^2}{m_h^2} z(1-z)^2 f_{-1}\right],$$

$$P_T^{(4)}(z) = z\left[-2f_2 + 2f_1 + 2z(z-1)f_0 + 4\frac{m^2}{m_h^2} z(z-1) f_{-1}\right]. \tag{C.9}$$

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