



Higgs physics confronts the M_W anomaly

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ARTICLE INFO

Article history:

Received 20 April 2022

Received in revised form 25 May 2022

Accepted 13 June 2022

Available online 15 June 2022

Editor: G.F. Giudice

ABSTRACT

The recent high-precision measurement of the W mass by the CDF collaboration is in sharp tension with the Standard Model prediction as obtained by the electroweak fit. If confirmed, this finding can only be explained in terms of new physics effects. In this work, we point out a generic connection between the M_W anomaly and Higgs physics observables such as $h \rightarrow \gamma\gamma, Z\gamma$ and the ratio $h \rightarrow ZZ/WW$. Moreover, we systematically classify new physics scenarios which can address the M_W anomaly via a tree-level contribution to the \hat{T} parameter. These include a real scalar triplet, a scalar quadruplet with the same hypercharge of the Higgs doublet, a Z' boson, a vector triplet with unit hypercharge and a vector boson with the gauge quantum numbers of the Higgs doublet. These solutions to the M_W anomaly are characterized by new physics states which are typically too heavy to be discovered in direct searches, but which might leave their imprints in Higgs physics.

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1. Introduction

Electroweak precision observables have played a crucial role to firmly establish the Standard Model (SM) as a quantum field theory and to constrain possible New Physics (NP) extensions. Remarkably, electroweak precision observables were instrumental to indirectly infer the masses of the top quark and the Higgs boson well before their direct detection at Tevatron and LHC, respectively. Nowadays, the SM electroweak fit is performed using as input parameters the fine structure constant α , the muon decay constant G_μ , the Z boson mass M_Z , the strong coupling $\alpha_s(M_Z)$, the top quark mass m_t , the Higgs mass M_h , and the hadronic contribution to the running of α , i.e. $\Delta\alpha_{\text{had}}^5(M_Z)$. In terms of these

parameters, all other observables can be predicted. In particular, the resulting value of the W^\pm boson mass from the electroweak fit is $M_W = 80354.5 \pm 5.7 \text{ MeV}$ [1].

The CDF collaboration has recently published a high-precision measurement of $M_W = 80433.5 \pm 9.4 \text{ MeV}$ [2], whose precision exceeds that of the current PDG world average, $M_W = 80379 \pm 12 \text{ MeV}$ [3], obtained from the combination of all previous measurements from LEP, D0, CDF, and ATLAS. The new CDF value turns out to be considerably larger than the current PDG world average as well as the value previously inferred from the SM electroweak fit [1].

Taking the new CDF result at face value, a few collaborations have already assessed its impact in the global electroweak fit, in the attempt of highlighting the favoured NP scenario to solve this anomaly (see e.g. [4–7]). In particular, it turned out that universal NP models, which are fully described by the fa-

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amous \hat{S} , \hat{T} , U , W , and Y parameters [8–10], provide an overall good quality of the fit. The viable solutions prefer either $\hat{T} \approx 10^{-3}$ and $\hat{S} = U = W = Y = 0$ or highly-correlated positive \hat{S} and \hat{T} parameters of comparable size $\hat{T} \sim \hat{S} \sim 10^{-3}$ and $U = W = Y = 0$. If \hat{S} and \hat{T} are loop-induced, they are of order $(g_{\text{NP}}^4/16\pi^2) \times M_W^2/M_{\text{NP}}^2$ and therefore, weakly-interacting theories require M_{NP} to lie at the electroweak scale to accommodate the M_W anomaly. Such a solution can be hardly reconciled with the direct-search bounds on new particles. Instead, tree-level NP effects –which are equivalent to the effects stemming from a strongly-coupled sector with $g_{\text{NP}} \sim 4\pi$ – can provide the desired values of \hat{S} and \hat{T} even for $M_{\text{NP}} \sim 10$ TeV.

The primary goal of this work is to establish a connection between the NP effects entering M_W and Higgs physics observables. Indeed, since within the SM Effective Field Theory (SMEFT) \hat{S} and \hat{T} receive contributions respectively from the $d=6$ operators $(H^\dagger \tau^a H) W_{\mu\nu}^a B^{\mu\nu}$ and $(H^\dagger D_\mu H)((D_\mu H)^\dagger H)$, it seems rather natural that NP effects in M_W need to be accompanied by modifications of the SM predictions for Higgs decay processes like $h \rightarrow \gamma\gamma, Z\gamma$ and $h \rightarrow ZZ, WW$. In Sect. 2 we quantitatively assess this connection in the context of the SMEFT. Another goal of the present analysis is to systematically classify explicit NP extensions of the SM which can give a sizable contribution to \hat{T} at the tree level. We provide this classification in Sect. 3 and, for those simplified models predicting a positive shift in \hat{T} , we discuss accordingly the correlated signals in Higgs physics. We conclude in Sect. 4 with a summary of our findings.

2. SMEFT approach to the M_W anomaly and Higgs physics

Parametrizing the SMEFT Lagrangian as

$$\mathcal{L}_{\text{SMEFT}} = \mathcal{L}_{\text{SM}} + \sum_i c_i \mathcal{O}_i, \quad (2.1)$$

where we adopt the Warsaw basis [11] and focus in particular on the following subset of operators, which are relevant for electroweak and Higgs physics:

$$\mathcal{O}_{HW} = (H^\dagger H) W_{\mu\nu}^a W^{a\mu\nu}, \quad (2.2)$$

$$\mathcal{O}_{HB} = (H^\dagger H) B_{\mu\nu} B^{\mu\nu}, \quad (2.3)$$

$$\mathcal{O}_{HWB} = (H^\dagger \tau^a H) W_{\mu\nu}^a B^{\mu\nu}, \quad (2.4)$$

$$\mathcal{O}_{HD} = (H^\dagger D_\mu H)((D_\mu H)^\dagger H), \quad (2.5)$$

$$\mathcal{O}_H = (H^\dagger H)^3, \quad (2.6)$$

$$\mathcal{O}_{H\Box} = (H^\dagger H)\Box(H^\dagger H), \quad (2.7)$$

$$\mathcal{O}_{eH} = (H^\dagger H)\bar{\ell}_L e_R H, \quad (2.8)$$

$$\mathcal{O}_{uH} = (H^\dagger H)\bar{q}_L u_R \tilde{H}, \quad (2.9)$$

$$\mathcal{O}_{dH} = (H^\dagger H)\bar{q}_L d_R H, \quad (2.10)$$

with the covariant derivative defined as $D_\mu = \partial_\mu + ig_2 W_\mu^a \tau^a + ig_1 B_\mu Y$. Employing the notation of Refs. [6,10], the leading electroweak oblique corrections are described by¹

$$\hat{S} \equiv \frac{c_W}{s_W} \Pi'(0)_{W_3 B} = \frac{c_W}{s_W} v^2 c_{HWB}, \quad (2.11)$$

$$\hat{T} \equiv \frac{1}{M_W^2} (\Pi_{W_3 W_3}(0) - \Pi_{W^+ W^-}(0)) = -\frac{v^2}{2} c_{HD}, \quad (2.12)$$

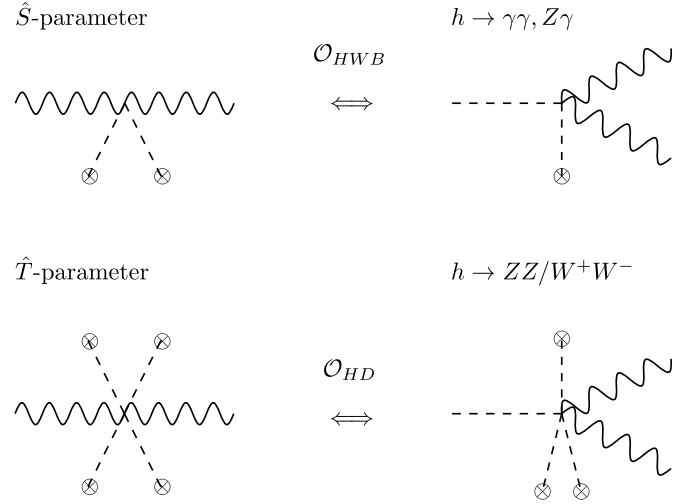


Fig. 1. \hat{S} and \hat{T} vs. Higgs connection.

with $v = 246$ GeV and $s_W \equiv \sin\theta_W$ ($c_W \equiv \cos\theta_W$). We remark that in Eqs. (2.11)–(2.12) we only included so-called “universal” bosonic operators. Upon applying the equations of motion in a given basis, other fermionic operators can contribute as well to the \hat{S} and \hat{T} parameters (see e.g. [12,13]). Concretely, in terms of the Warsaw basis these are four-fermion operators as well as operators of the type $(H^\dagger \overleftrightarrow{D}_\mu H)(\bar{\psi} \gamma^\mu \psi)$. These operators can also lead to contributions to electroweak precision observables beyond the oblique parameters (with the exception of top-quark operators²) and hence are neglected in the present analysis.

The M_W anomaly could be due to a universal new physics correction to \hat{T} [6]

$$\hat{T} \simeq (0.84 \pm 0.14) \times 10^{-3}, \quad (2.13)$$

($c_{HD} = -(0.17 \pm 0.07/\text{TeV}^2)$) as well as a correlated contribution to $\hat{S} \sim 10^{-3}$ ($c_{HWB} \sim (0.07/\text{TeV}^2)$) of the same size of \hat{T} , but compatible with zero [6,7]. The inclusion of higher-order corrections in the momentum expansion of the inverse propagators (Y and W) does not alter significantly the fit [6], while a non-vanishing \hat{U} parameter can also explain by itself the M_W anomaly [7]. However, under the assumption of heavy NP, which is captured by the SMEFT description, the \hat{U} parameter is usually neglected since it arises from $d=8$ operators.

Since the \hat{S} and \hat{T} parameters are obtained by condensing the Higgs fields in \mathcal{O}_{HWB} and \mathcal{O}_{HD} , there is clearly a connection with Higgs physics, as highlighted schematically in Fig. 1.

Writing the SMEFT Lagrangian in the electroweak broken phase as

$$\begin{aligned} \mathcal{L}_{\text{SMEFT}}^{\text{int}} \ni & g_{hWW}^{(1)} h W_\mu^+ W^{-\mu} + g_{hWW}^{(2)} h W_{\mu\nu}^+ W^{-\mu\nu} \\ & + g_{hZZ}^{(1)} h Z_\mu Z^\mu + g_{hZZ}^{(2)} h Z_{\mu\nu} Z^{\mu\nu} + g_{h\gamma\gamma} h F_{\mu\nu} F^{\mu\nu} \\ & + g_{h\gamma Z} h F_{\mu\nu} Z^{\mu\nu} + g_{hhhh} h^4 \\ & + (g_{he} h \bar{e}_L e_R + g_{hu} h \bar{u}_L u_R + g_{hd} h \bar{d}_L d_R + \text{h.c.}) + \dots, \end{aligned} \quad (2.14)$$

one finds at tree level (see e.g. [17])

$$g_{hWW}^{(1)} = \frac{2M_W^2}{v} \left(1 - \frac{v^2}{4} (c_{HD} - 4c_{H\Box}) \right), \quad (2.15)$$

¹ In this notation the S and T parameters of Refs. [8,9] read $S = 4s_W^2 \hat{S}/\alpha$ and $T = \hat{T}/\alpha$.

² Top-quark operators can be (weakly) constrained by top-quark physics [14] and via their loop contributions by electroweak observables [15] and Higgs physics [16].

$$g_{hWW}^{(2)} = 2v c_{HW}, \quad (2.16)$$

$$g_{hZZ}^{(1)} = \frac{M_Z^2}{v} \left(1 + \frac{v^2}{4} (c_{HD} + 4c_{H\Box}) \right), \quad (2.17)$$

$$g_{hZZ}^{(2)} = v \left[\frac{M_W^2}{M_Z^2} c_{HW} + \frac{M_Z^2 - M_W^2}{M_Z^2} c_{HB} + \frac{g_1 g_2}{g_1^2 + g_2^2} c_{HWB} \right], \quad (2.18)$$

$$g_{h\gamma\gamma} = v \left[\frac{M_Z^2 - M_W^2}{M_Z^2} c_{HW} + \frac{M_W^2}{M_Z^2} c_{HB} - \frac{g_1 g_2}{g_1^2 + g_2^2} c_{HWB} \right], \quad (2.19)$$

$$g_{h\gamma Z} = 2v \left[\frac{g_1 g_2}{g_1^2 + g_2^2} c_{HW} - \frac{g_1 g_2}{g_1^2 + g_2^2} c_{HB} + \frac{1}{2} \frac{g_1^2 - g_2^2}{g_1^2 + g_2^2} c_{HWB} \right], \quad (2.20)$$

$$g_{hhh} = -\frac{M_h^2}{2v} \left[1 - \frac{3v^2}{4} (c_{HD} - 4c_{H\Box}) - \frac{2v^4}{M_h^2} c_H \right], \quad (2.21)$$

$$g_{h\psi} = -\frac{m_\psi}{v} \left(1 - \frac{v^2}{4} (c_{HD} - 4c_{H\Box}) \right) + \frac{c_{\psi H} v^2}{\sqrt{2}}, \quad (2.22)$$

with $\psi = e, u, d$. In order to canonically normalize the Higgs kinetic term we adopted the field redefinition

$$h \rightarrow h \left(1 + v^2 (c_{H\Box} - \frac{1}{4} c_{HD}) \left(1 + \frac{h}{v} + \frac{h^2}{3v^2} \right) \right), \quad (2.23)$$

which removes the momentum-dependence in the Higgs self-couplings.

The most relevant Higgs observables, which are affected by the presence of non-zero \hat{S} , $\hat{T} \sim 10^{-3}$, are the Higgs boson decays into vector bosons. Defining the Higgs signal strengths as³

$$\mu_{VV'} \equiv \frac{\Gamma(h \rightarrow VV')}{\Gamma_{\text{SM}}(h \rightarrow VV')}, \quad (2.24)$$

with $V, V' = \gamma, Z, W$, we can compute the corrections arising from the modified Higgs couplings in Eq. (2.14). For the observables related to \hat{S} (c_{HWB}) one finds (see e.g. [18])

$$\begin{aligned} \mu_{\gamma\gamma} &\simeq 1 + \frac{4\pi v^2}{I^\gamma \alpha} (s_W^2 c_{HW} + c_W^2 c_{HB} - s_W c_W c_{HWB}) \\ &\simeq 1 + 0.23 \left(\frac{\hat{S}}{10^{-3}} \right), \end{aligned} \quad (2.25)$$

$$\begin{aligned} \mu_{ZZ} &\simeq 1 + \frac{4\pi v^2}{I^Z \alpha} (s_W c_W c_{HW} - s_W c_W c_{HB} - \frac{1}{2} (c_W^2 - s_W^2) c_{HWB}) \\ &\simeq 1 + 0.084 \left(\frac{\hat{S}}{10^{-3}} \right), \end{aligned} \quad (2.26)$$

where in the last steps we used the SM values $I^\gamma = -1.64$ and $I^Z = -2.84$, and neglected the contribution of c_{HW} and c_{HB} in order to highlight the connection with \hat{S} . Hence, a +20% modification of $\mu_{\gamma\gamma}$ is generically expected for $\hat{S} \sim 10^{-3}$, which is in the ballpark of the present LHC experimental sensitivity at the 10% level [19]. On the other hand, the predicted shift in μ_{ZZ} is presently too small to be detected, although an $\mathcal{O}(10\%)$ sensitivity might be achieved at the HL-LHC [20]. Another observable that is

Table 1

New physics states which can yield a tree-level contribution to \hat{T} via $d \leq 4$ interactions with SM states. Highlighted in pink are the representations predicting a positive shift on \hat{T} . The last column indicates whether a tree-level contribution to \hat{S} is generated (\checkmark) or not (\times).

Field	Spin	SU(3) _C	SU(2) _L	U(1) _Y	sign(\hat{T})	\hat{S}
Δ	0	1	3	0	+	\times
Δ_1	0	1	3	1	-	\times
Θ_1	0	1	4	1/2	+	\times
Θ_3	0	1	4	3/2	-	\times
\mathcal{B}	1	1	1	0	+	\times
\mathcal{B}_1	1	1	1	1	-	\times
\mathcal{W}	1	1	3	0	-	\times
\mathcal{W}_1	1	1	3	1	+	\times
\mathcal{L}	1	1	2	1/2	+/-	\checkmark

directly sensitive to \hat{T} via c_{HD} is the ratio μ_{ZZ}/μ_{WW} [21], which reads

$$\frac{\mu_{ZZ}}{\mu_{WW}} \simeq 1 + 2c_{HD} v^2 \simeq 1 - 0.0034 \left(\frac{\hat{T}}{0.84 \times 10^{-3}} \right). \quad (2.27)$$

However, also in this case the predicted deviation is too small to be presently detected.

3. Heavy new physics: tree-level contributions to \hat{T}

Since the M_W anomaly hints at a sizeable $\hat{T} \sim 10^{-3}$, correlated with an \hat{S} parameter compatible with zero [6,7], we will now focus on heavy NP extensions which can yield a tree-level contribution to \hat{T} via the operator \mathcal{O}_{HD} .⁴

It turns out that such states are either scalars \mathcal{S} or vectors \mathcal{V} , whose quadratic Lagrangian can be written as

$$\mathcal{L}_{\mathcal{S}}^{\text{quad}} = \eta [(D_\mu \mathcal{S})^\dagger D^\mu \mathcal{S} - M_{\mathcal{S}}^2 \mathcal{S}^\dagger \mathcal{S}], \quad (3.1)$$

$$\mathcal{L}_{\mathcal{V}}^{\text{quad}} = \eta [(D_\mu \mathcal{V}_\nu)^\dagger D^\nu \mathcal{V}^\mu - (D_\mu \mathcal{V}_\nu)^\dagger D^\mu \mathcal{V}^\nu + M_{\mathcal{V}}^2 \mathcal{V}_\mu^\dagger \mathcal{V}^\mu], \quad (3.2)$$

with the prefactor $\eta = 1$ ($\eta = 1/2$) for a complex (real) representation. The representations which can generate \mathcal{O}_{HD} , or higher-dimensional variants thereof such as $(H^\dagger H)_{\mathcal{O}_{HD}}$, at tree-level are displayed in Table 1 (for similar classifications see also Refs. [22–25]). In the following, we discuss in detail each simplified model, and for those cases leading to a positive \hat{T} we analyze in turn the correlated signals in Higgs physics.

3.1. $\Delta \sim (1, 3, 0)_{\mathcal{S}}$

From the interaction Lagrangian

$$\mathcal{L}_{\Delta}^{\text{int}} \ni -\kappa_{\Delta} H^\dagger \Delta^a \sigma^a H - \frac{\lambda_{H\Delta}}{2} (H^\dagger H) \Delta^a \Delta^a, \quad (3.3)$$

one obtains

$$c_{HD} = -2 \frac{\kappa_{\Delta}^2}{M_{\Delta}^4}, \quad (3.4)$$

and hence

$$\hat{T} = \frac{\kappa_{\Delta}^2 v^2}{M_{\Delta}^4} = 0.84 \times 10^{-3} \left(\frac{|\kappa_{\Delta}|}{M_{\Delta}} \right)^2 \left(\frac{8.5 \text{ TeV}}{M_{\Delta}} \right)^2, \quad (3.5)$$

which has the correct sign to explain the M_W anomaly.

³ For illustration of the argument, we work under the assumption that the Higgs coupling to gluons is not modified by NP and neglect other production channels than the ones to gluons.

⁴ We note that in a few models discussed below (Δ_1 , \mathcal{B} , \mathcal{W} and \mathcal{L} from Table 1) one also generates fermionic operators that affect \hat{S} and \hat{T} (cf. discussion below Eq. (2.12)). However, these contributions are always proportional to the couplings of the new particle to SM fermions and therefore they can be parametrically suppressed.

When the Δ is not integrated out, the tree-level contribution to \hat{T} can be alternatively understood to arise from the generation of a tree-level vacuum expectation value (VEV) for Δ , that is $\langle \Delta \rangle \equiv v_\Delta = \kappa_\Delta v^2 / (2M_\Delta^2)$. In general, the VEV of a scalar representation $S \sim (1, 2j + 1, y)$ yields [26,27]

$$\hat{T} \simeq 4 \left(\eta [j(j+1) - y^2] - 2y^2 \right) \frac{\langle S \rangle^2}{v^2}, \quad (3.6)$$

with $\langle S \rangle = \alpha v_S$, where $\alpha = 1$ ($\alpha = 1/\sqrt{2}$) for a real (complex) representation and v_S is the VEV of the canonically normalized real scalar component of S . In the case $S = \Delta$ this yields $\hat{T} \simeq \kappa_\Delta^2 v^2 / M_\Delta^4$, as in Eq. (3.5).

The connection between this scalar triplet and electroweak precision measurements was previously considered e.g. in Refs. [28–30]. Note that the perturbativity range of the massive κ_Δ parameter can be obtained by requiring that finite loop corrections to the trilinear scalar vertex $\Delta H^\dagger H$ remain smaller than the tree-level value [31]. This yields $|\kappa_\Delta| / M_\Delta \lesssim 4\pi$ [32]. Hence, a scalar triplet well above the TeV scale and with perturbative couplings can explain the value of \hat{T} while easily evading all direct collider searches. In particular, saturating the perturbativity bound, it turns out that $M_\Delta \lesssim 100$ TeV.

Other coefficients which are unavoidably generated after integrating out Δ are directly correlated with \hat{T} via the coupling κ_Δ :

$$c_H = -4 \frac{\kappa_\Delta^2}{M_\Delta^4} \left(\frac{\lambda_{H\Delta}}{8} - \lambda \right) = -4 \frac{\hat{T}}{v^2} \left(\frac{\lambda_{H\Delta}}{8} - \lambda \right), \quad (3.7)$$

$$c_{H\Box} = \frac{\kappa_\Delta^2}{2M_\Delta^4} = \frac{\hat{T}}{2v^2}, \quad (3.8)$$

$$c_{eH, uH, dH} = \frac{\kappa_\Delta^2 Y_{e,u,d}}{M_\Delta^4} = \frac{\hat{T}}{v^2} Y_{e,u,d}, \quad (3.9)$$

where λ is the SM quartic Higgs coupling and $Y_{e,u,d}$ are SM Yukawas. Hence, in the triplet model a non-zero \hat{T} can be correlated to various Higgs signals (see Eq. (2.14)). The strongest dependence on \hat{T} is via the modification of the trilinear Higgs self-coupling, but one needs to keep in mind that the prospect for its measurement at the HL-LHC is $0.1 < g_{hhh} / g_{hhh}^{\text{SM}} < 2.3$ [33] while the model predicts (setting $\lambda_{H\Delta} = 0$) deviations of $\mathcal{O}(1\%)$ in g_{hhh} (cf. Eq. (2.21)). Similarly, the deviations in the W and Z couplings are out of reach for the LHC.

3.2. $\Delta_1 \sim (1, 3, 1)_S$

From the interaction Lagrangian

$$\mathcal{L}_{\Delta_1}^{\text{int}} \ni -\kappa_{\Delta_1} (\Delta_1^a)^\dagger \tilde{H}^\dagger \sigma^a H + \text{h.c.}, \quad (3.10)$$

one obtains

$$c_{HD} = 4 \frac{|\kappa_{\Delta_1}|^2}{M_{\Delta_1}^4}, \quad (3.11)$$

and hence

$$\hat{T} = -2 \frac{|\kappa_{\Delta_1}|^2 v^2}{M_{\Delta_1}^4}, \quad (3.12)$$

which predicts the wrong sign to explain the new CDF M_W value.

3.3. $\Theta_1 \sim (1, 4, 1/2)_S$

Electroweak quadruplets contribute to \hat{T} via the $d = 8$ operator $(H^\dagger H) \mathcal{O}_{HD}$.⁵ It is hence more practical to directly compute \hat{T} via the VEV contribution. We consider the interaction Lagrangian

$$\mathcal{L}_{\Theta_1}^{\text{int}} \ni M_{\Theta_1}^2 (\Theta_1)_{ijk} (\Theta_1^*)^{ijk} - \lambda_{H3\Theta_1} H^{*i} (\Theta_1)_{ijk} H^{*j} \epsilon^{kl} H_l + \text{h.c.}, \quad (3.13)$$

in a phase convention where $\lambda_{H3\Theta_1}$ is real and we employed a symmetric tensor notation for the quadruplet, with latin indices in $SU(2)_L$ space and $\epsilon = i\sigma^2$. The embedding of the canonically normalized charge eigenstates reads $(\Theta_1)_{111} = \Theta_1^{+++}$, $(\Theta_1)_{112} = \frac{1}{\sqrt{3}} \Theta_1^+$, $(\Theta_1)_{122} = \frac{1}{\sqrt{3}} \Theta_1^0$, $(\Theta_1)_{222} = \Theta_1^-$. In particular, for the VEV of the neutral component one obtains

$$\langle \Theta_1^0 \rangle \equiv \frac{v_{\Theta_1}}{\sqrt{2}} \simeq \frac{\lambda_{H3\Theta_1} v^3}{2\sqrt{6} M_{\Theta_1}^2}, \quad (3.14)$$

and hence, using the VEV formula in Eq. (3.6)

$$\hat{T} \simeq 12 \frac{\langle \Theta_1^0 \rangle^2}{v^2} \simeq \frac{\lambda_{H3\Theta_1}^2 v^4}{2M_{\Theta_1}^4} = 0.84 \times 10^{-3} \lambda_{H3\Theta_1}^2 \left(\frac{1.2 \text{ TeV}}{M_{\Theta_1}} \right)^4, \quad (3.15)$$

which has the correct sign to explain the M_W anomaly. Note that due to the different scaling of the \hat{T} parameter (compared e.g. to the triplet case in Eq. (3.5)) the mass of the quadruplet needs to be around 1 TeV for $\mathcal{O}(1)$ couplings to the Higgs.

At the leading order in the SMEFT, one also generates the Wilson coefficient

$$c_H = \frac{\lambda_{H3\Theta_1}^2}{M_{\Theta_1}^2} = \frac{2M_{\Theta_1}^2}{v^4} \hat{T} = \frac{0.040}{v^2} \left(\frac{M_{\Theta_1}}{1.2 \text{ TeV}} \right)^2 \left(\frac{\hat{T}}{0.84 \times 10^{-3}} \right), \quad (3.16)$$

which is directly correlated with \hat{T} . Substituting the above value of c_H into Eq. (2.21) we then obtain

$$\frac{g_{hhh}}{g_{hhh}^{\text{SM}}} - 1 = -31\% \left(\frac{M_{\Theta_1}}{1.2 \text{ TeV}} \right)^2 \left(\frac{\hat{T}}{0.84 \times 10^{-3}} \right), \quad (3.17)$$

which implies up to an $\mathcal{O}(1)$ variation in the trilinear Higgs self-coupling, depending on the value of M_{Θ_1} , which receives an upper bound from perturbativity (see e.g. [23]). All in all, the reason for such a large effect in the trilinear Higgs self-coupling can be understood from the fact that in the quadruplet case \hat{T} is generated by a $d = 8$ operator, while the contribution to c_H arises at $d = 6$.

3.4. $\Theta_3 \sim (1, 4, 3/2)_S$

We consider the interaction Lagrangian

$$\mathcal{L}_{\Theta_3}^{\text{int}} \ni M_{\Theta_3}^2 (\Theta_3)_{ijk} (\Theta_3^*)^{ijk} - \lambda_{H3\Theta_3} H^{*i} H^{*j} H^{*k} (\Theta_3)_{ijk} + \text{h.c.}, \quad (3.18)$$

in a phase convention where $\lambda_{H3\Theta_3}$ is real. The embedding of the canonically normalized charge eigenstates reads $(\Theta_3)_{111} = \Theta_3^{++++}$, $(\Theta_3)_{112} = \frac{1}{\sqrt{3}} \Theta_3^{+++}$, $(\Theta_3)_{122} = \frac{1}{\sqrt{3}} \Theta_3^{++}$, $(\Theta_3)_{222} = \Theta_3^0$. For the VEV of the neutral component one obtains

⁵ Note that quadruplets do not generate tree-level contributions to the \hat{S} and \hat{U} parameters at $d = 8$ [34].

$$\langle \Theta_3^0 \rangle \equiv \frac{v_{\Theta_3}}{\sqrt{2}} \simeq \frac{\lambda_{H3\Theta_3} v^3}{2\sqrt{2}M_{\Theta_3}^2}, \quad (3.19)$$

and hence, using the VEV formula in Eq. (3.6)

$$\hat{T} \simeq -12 \frac{\langle \Theta_3^0 \rangle^2}{v^2} \simeq -\frac{3\lambda_{H3\Theta_3}^2 v^4}{2M_{\Theta_3}^4}, \quad (3.20)$$

which predicts the wrong sign to solve the M_W anomaly.

3.5. $\mathcal{B} \sim (1, 1, 0)_\gamma$

From the interaction Lagrangian

$$\mathcal{L}_{\mathcal{B}}^{\text{int}} \ni -g_{\mathcal{B}}^H \mathcal{B}^\mu H^\dagger iD_\mu H + \text{h.c.}, \quad (3.21)$$

one obtains

$$c_{HD} = -2 \frac{(\text{Re}(g_{\mathcal{B}}^H))^2}{M_{\mathcal{B}}^2}, \quad (3.22)$$

and hence

$$\hat{T} = \frac{(\text{Re}(g_{\mathcal{B}}^H))^2 v^2}{M_{\mathcal{B}}^2} = 0.84 \times 10^{-3} (\text{Re}(g_{\mathcal{B}}^H))^2 \left(\frac{8.5 \text{ TeV}}{M_{\mathcal{B}}} \right)^2, \quad (3.23)$$

which has the correct sign to explain the M_W anomaly. The possibility of raising the M_W mass via a Z' boson was previously considered e.g. in [6,35]. Note that the Z' phenomenology highly depends on its coupling to SM fermions, which have not been specified here.

3.6. $\mathcal{B}_1 \sim (1, 1, 1)_\gamma$

From the interaction Lagrangian

$$\mathcal{L}_{\mathcal{B}_1}^{\text{int}} \ni -g_{\mathcal{B}_1}^H \mathcal{B}_1^{\mu\dagger} iD_\mu H^\dagger i\sigma^2 H + \text{h.c.}, \quad (3.24)$$

one obtains

$$c_{HD} = \frac{|g_{\mathcal{B}_1}^H|^2}{M_{\mathcal{B}_1}^2}, \quad (3.25)$$

and hence

$$\hat{T} = -\frac{|g_{\mathcal{B}_1}^H|^2 v^2}{2M_{\mathcal{B}_1}^2}, \quad (3.26)$$

which predicts the wrong sign to accommodate the M_W anomaly.

3.7. $\mathcal{W} \sim (1, 3, 0)_\gamma$

From the interaction Lagrangian

$$\mathcal{L}_{\mathcal{W}}^{\text{int}} \ni -\frac{1}{2} g_{\mathcal{W}}^H \mathcal{W}^{\mu a} H^\dagger \sigma^a iD_\mu H + \text{h.c.}, \quad (3.27)$$

one obtains

$$c_{HD} = \frac{(\text{Im}(g_{\mathcal{W}}^H))^2}{2M_{\mathcal{W}}^2}, \quad (3.28)$$

and hence

$$\hat{T} = -\frac{(\text{Im}(g_{\mathcal{W}}^H))^2 v^2}{4M_{\mathcal{W}}^2}, \quad (3.29)$$

which predicts the wrong sign to accommodate the M_W anomaly.

3.8. $\mathcal{W}_1 \sim (1, 3, 1)_\gamma$

From the interaction Lagrangian

$$\mathcal{L}_{\mathcal{W}_1}^{\text{int}} \ni -\frac{1}{2} g_{\mathcal{W}_1}^H \mathcal{W}_1^{\mu a\dagger} iD_\mu H^\dagger i\sigma^a H + \text{h.c.}, \quad (3.30)$$

one obtains

$$c_{HD} = -\frac{|g_{\mathcal{W}_1}^H|^2}{4M_{\mathcal{W}_1}^2}, \quad (3.31)$$

and hence

$$\hat{T} = \frac{|g_{\mathcal{W}_1}^H|^2 v^2}{8M_{\mathcal{W}_1}^2} = 0.84 \times 10^{-3} |g_{\mathcal{W}_1}^H|^2 \left(\frac{3.0 \text{ TeV}}{M_{\mathcal{W}_1}} \right)^2, \quad (3.32)$$

which has the correct sign to explain the M_W anomaly. Note that the interaction in Eq. (3.30) does not necessarily arise from a renormalizable theory and here we have included only $d = 4$ interactions with the SM fields.

3.9. $\mathcal{L} \sim (1, 2, 1/2)_\gamma$

From the interaction Lagrangian

$$\begin{aligned} \mathcal{L}_{\mathcal{L}}^{\text{int}} \ni & -(\gamma_{\mathcal{L}} \mathcal{L}^\dagger D^\mu H + \text{h.c.}) - ig_{\mathcal{L}}^B \mathcal{L}_\mu^\dagger \mathcal{L}_\nu B^{\mu\nu} \\ & - ig_{\mathcal{L}}^W \mathcal{L}_\mu^\dagger \sigma^a \mathcal{L}_\nu W^{a\mu\nu} - h_{\mathcal{L}}^{(2)} \mathcal{L}^\dagger H H^\dagger \mathcal{L} - (h_{\mathcal{L}}^{(3)} (\mathcal{L}^\dagger H)^2 + \text{h.c.}), \end{aligned} \quad (3.33)$$

after an Higgs field redefinition in order to have canonical kinetic terms (see [25]), one obtains

$$c_{HD} = \frac{g_1 g_{\mathcal{L}}^B |\gamma_{\mathcal{L}}|^2}{M_{\mathcal{L}}^4} - \frac{h_{\mathcal{L}}^{(2)} |\gamma_{\mathcal{L}}|^2}{M_{\mathcal{L}}^4} + \frac{2\text{Re}(h_{\mathcal{L}}^{(3)} \gamma_{\mathcal{L}}^{*2})}{M_{\mathcal{L}}^4}, \quad (3.34)$$

and hence

$$\hat{T} = -\frac{g_1 g_{\mathcal{L}}^B |\gamma_{\mathcal{L}}|^2 v^2}{2M_{\mathcal{L}}^4} + \frac{h_{\mathcal{L}}^{(2)} |\gamma_{\mathcal{L}}|^2 v^2}{2M_{\mathcal{L}}^4} - \frac{\text{Re}(h_{\mathcal{L}}^{(3)} \gamma_{\mathcal{L}}^{*2}) v^2}{M_{\mathcal{L}}^4}, \quad (3.35)$$

which can have both signs. Moreover, after integrating out \mathcal{L} , also the Wilson coefficient

$$c_{WB} = -\frac{g_1 g_2 |\gamma_{\mathcal{L}}|^2}{4M_{\mathcal{L}}^4} - \frac{g_2 g_{\mathcal{L}}^B |\gamma_{\mathcal{L}}|^2}{4M_{\mathcal{L}}^4} - \frac{g_1 g_{\mathcal{L}}^W |\gamma_{\mathcal{L}}|^2}{4M_{\mathcal{L}}^4} \quad (3.36)$$

is generated, which implies

$$\hat{S} = \frac{c_W}{s_W} \left[-\frac{g_1 g_2 |\gamma_{\mathcal{L}}|^2 v^2}{4M_{\mathcal{L}}^4} - \frac{g_2 g_{\mathcal{L}}^B |\gamma_{\mathcal{L}}|^2 v^2}{4M_{\mathcal{L}}^4} - \frac{g_1 g_{\mathcal{L}}^W |\gamma_{\mathcal{L}}|^2 v^2}{4M_{\mathcal{L}}^4} \right], \quad (3.37)$$

which is of the same order as the contribution to the \hat{T} parameter. Hence, this model could yield a large effect in the $h \rightarrow \gamma\gamma$ rate (as discussed in Sect. 2).

4. Conclusions

In this work, we have investigated the connection between the M_W anomaly (stemming from the recent M_W measurement by the CDF collaboration [2]) and Higgs physics. This connection is quite natural from the point of view of the SMEFT, since the \hat{S} and \hat{T} parameters arise from $d = 6$ operators containing multiple insertions of the Higgs doublet set on its VEV. Hence, by promoting one of those VEVs to a dynamical field, one automatically predicts modified Higgs signals (cf. Fig. 1). The largest effect turns out to be in $h \rightarrow \gamma\gamma$, which is modified up to +20% for values of $\hat{S} \sim 10^{-3}$

which are generically suggested by the electroweak fit. Modifications of the $h \rightarrow Z\gamma$ rate as well as the ratio $h \rightarrow ZZ/WW$ (which might be used to probe the \hat{T} parameter) are instead too small to be presently detected.

Since a non-zero and positive $\hat{T} \sim 10^{-3}$ is suggested by global electroweak fits, in Sect. 3 we classified SM extensions which predict a positive tree-level shift of \hat{T} . Remarkably, there are only few solutions with NP states that couple directly to the Higgs via $d \leq 4$ interactions: a scalar triplet $\Delta \sim (1, 3, 0)_S$, a scalar quadruplet $\Theta_1 \sim (1, 4, 1/2)_S$, a Z' boson $B \sim (1, 1, 0)_V$, a vector triplet $\mathcal{W}_1 \sim (1, 3, 1)_V$ and a vector boson $\mathcal{L} \sim (1, 2, 1/2)_V$. In all these cases the value of $\hat{T} \sim 10^{-3}$ can be easily explained via a NP state with mass around 10 TeV and $\mathcal{O}(1)$ couplings to the Higgs (barring the quadruplet case which needs to be at the TeV scale and implies as well a large deviation in the trilinear Higgs self-coupling – see Eq. (3.17)).

Although the NP states implied by such scenarios can escape direct detection at particle colliders, they might still leave their imprints via modifications of Higgs signals, which become especially correlated in explicit models.

Declaration of competing interest

The authors declare that they have no known competing financial interests or personal relationships that could have appeared to influence the work reported in this paper.

Acknowledgements

We thank David Marzocca, Alessandro Strumia and Nicolò Toniole for useful discussions.

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