

Universal SMEFT as a tool for discovery of universal new physics at the high luminosity LHC

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We analyze the potential of the Universal SMEFT as a tool to study universal new physics in the neutral and charged Drell-Yan processes at the high-luminosity LHC. In order to do so, we generated pseudodata containing the contributions of beyond the standard model vector bosons to the dilepton production. Our results show that fits with the minimal theoretical bias can be used not only to unveil the existence of new physics but also to accurately extract its properties. Moreover, the results are rather stable with respect to the order of truncation of the effective field theory.

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I. INTRODUCTION

Up to now there has been no direct evidence of new physics (NP) at the CERN Large Hadron Collider (LHC) despite the accumulated statistics. One possibility is that the states associated to the new dynamics are too heavy to produce a resonance signature at the LHC but they are not too far from the kinematic reach. In this scenario it is standard to look for NP signals in the tail of kinematic distributions, with effective field theory (EFT) being the optimum model-independent theoretical tool to perform such analysis.

In this context, dilepton production, i.e., the Drell-Yan (DY) process [1], is one of the most precisely measured observables at the LHC. It can proceed either via neutral current (NC) or charged current (CC)

$$pp \rightarrow \ell^+ \ell^- \quad \text{and} \quad pp \rightarrow \ell^\pm \nu_\ell$$

with $\ell = e, \mu$. Involving only leptons in the final state, provides a clean environment for both experimental studies and theoretical predictions. Thus it constitutes one of the most promising windows for beyond the standard model (BSM) searches. In fact the LHC experimental collaborations have taken advantage of these features to perform precision standard model (SM) tests [2–6] and searches for new resonances [7–10]. In the near future the high-luminosity run of the LHC (HL-LHC) is expected to accumulate 3000 fb^{-1} at 14 TeV, enlarging the parameter space that can be probed.

As mentioned above, a particularly suitable approach to parametrize deviations in kinematic distributions in a model-independent fashion is to employ the EFT framework. Under the minimal assumption that the scalar particle observed in 2012 [11,12] is part of an electroweak doublet, the $SU(2)_L \otimes U(1)_Y$ symmetry can be linearly realized and the resulting EFT is the so-called Standard Model EFT (SMEFT). In this framework, deviations from the SM predictions are parametrized as higher order operators

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{SM}} + \sum_{n>4,j} \frac{f_{n,j}}{\Lambda^{n-4}} \mathcal{O}_{n,j}, \quad (1.1)$$

where Λ is a characteristic high energy scale, $\mathcal{O}_{n,j}$ are higher dimensional operators, and $f_{n,j}$ are their respective Wilson coefficients (WC).

In the bottom-up approach to BSM searches in this context, WCs are treated as being independent in the fit to

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the considered dataset. Open questions in this framework include: Can such an EFT fit with arbitrary WC provide a clear signature of NP? If so, do the extracted values of the WC's capture the properties of the BSM? Do the answers to these two questions depend on the order of the truncation of the EFT series?

These questions are central to the EFT methodology and as such they have been approached in the existing literature. In this respect, the study of correlations in the SMEFT series associated with specific UV completions as well as the challenge of identifying the specific UV completion in SMEFT fits has been previously addressed (see for example [13–25] for an incomplete list). There is also an increasing number of studies of the impact of dimension-eight operators in EFT analyses either stressing the existence of new signal topologies due to dimension-eight operators [26–28] or investigating the changes generated by the inclusion of these operators to fits performed at dimension six, in Drell-Yan [29–36], production of electroweak gauge bosons [37–42], Higgs physics [43–50], and electroweak precision observables [48,51]. Somehow none of these studies directly addressed the question of the discovery of NP in fits with minimal theoretical priors, such as SMEFT fits, which is the focus of our work.

Thus it is with these questions in focus, that we assess in this work the HL-LHC potential to search for NP employing the SMEFT framework. We consider two simplified scenarios in which we extend the SM with the addition of a new $U(1)$ gauge field in one case, and additional $SU(2)$ vector bosons in the other. As is well known, beyond $\mathcal{O}(1/\Lambda^2)$, the SMEFT in its most general form possesses an untreatable number of operators, which precludes a definitive answer to the third question above. The answer becomes accessible, however, in the context of Universal BSM scenarios, that is, when the NP either dominantly couples to the SM bosons and if it couples to fermions, does so via SM currents. Therefore, we will consider two scenarios where the new gauge bosons couple to SM currents, also referred as mirror-hypercharge [or mirror- $U(1)_Y$] and mirror- $SU(2)_L$ models.

Our strategy to quantitatively answer the above questions proceeds as follows. First we simulate pseudodata containing the SM and BSM predictions for the DY observables in the context of the two models for a variety of model parameters (masses and couplings). In addition we also take into account the electroweak precision observables (EWPO) in order to better constrain the parameter space under consideration. Subsequently we perform a fit to the DY pseudodata and EWPO for each model and each set of model parameters using Universal SMEFT (USMEFT) and obtained the best fit and allowed range of all the operator Wilson coefficients included in the fit. In this framework, the signal of NP is, then, associated with the determination of some nonvanishing Wilson coefficients within a particular confidence level. The ability to identify the specific NP is addressed by comparing the extracted ranges of

coefficients from the fit with the predictions from matching the specific model to the USMEFT operators. Furthermore, in order to address the dependence of the answers on the order of the EFT truncation, we perform the analyses including all operators at orders $1/\Lambda^2$ and $1/\Lambda^4$.

The presentation of our work is organized as follows. In Sec. II we summarize the USMEFT operators relevant for the analyses of DY data up to order $1/\Lambda^4$, complemented by Appendix A. The analyzed BSM scenarios are introduced in Sec. III with further details in the Appendix B. The specifics of the framework and assumptions in our analyses are contained in Sec. IV. Our results are presented in detail in Sec. V and we briefly summarize our findings in Sec. VI.

II. OPERATOR BASIS

Within the USMEFT predictions for observables at order $\mathcal{O}(1/\Lambda^4)$ require evaluating the SM contributions, the interference between the $1/\Lambda^2$ amplitude ($\mathcal{M}^{(6)}$) with the SM amplitude, the square of the dimension-six amplitude, as well as the interference of the $1/\Lambda^4$ amplitude with the SM, which we represent as,

$$|M_{\text{SM}}|^2 + \mathcal{M}_{\text{SM}}^* \mathcal{M}^{(6)} + |\mathcal{M}^{(6)}|^2 + \mathcal{M}_{\text{SM}}^* \mathcal{M}^{(6,2)} + \mathcal{M}_{\text{SM}}^* \mathcal{M}^{(8)}. \quad (2.1)$$

$\mathcal{M}^{(8)}$ includes amplitudes with one dimension-eight operator coefficient while $\mathcal{M}^{(6,2)}$ includes the contribution of the insertion of two dimension-six Wilson coefficients in the amplitude.

In order to perform our analyses we must choose a basis of independent operators for the USMEFT. Under the assumption of universality an independent basis can be constructed of *bosonic* operators, that is operators formed exclusively of bosonic fields. As is widely known, operators connected by the use of the classical equations of motion (EOM) of the SM fields lead to the same S -matrix elements [52–55]. Thus, it is possible to trade bosonic operators with higher derivatives for operators involving fermions in the form of combinations of SM currents herein called *fermionic operators*. We refer to the resulting basis of this exchanged as the *rotated basis*. In fact in order to perform the Monte Carlo event simulations it is more convenient to work with the operators in the rotated basis [39,56].

The complete list of universal dimension-six operators is presented in Ref. [57]. Assuming that the Yukawa couplings are negligible as well as requiring C and P conservation one can identify six independent bosonic operators contributing to the weak boson propagators at order $1/\Lambda^2$, hence contributing to EWPO and/or Drell-Yan processes. Two of these bosonic operators can be rotated into fermionic operators by the EOM to form the relevant rotated set of dimension-six operators listed in Table I,

TABLE I. C and P conserving dimension-six operators that contribute to DY and EWPO in the rotated basis for universal theories and their respective Wilson coefficients.

Operator basis		Coefficient
$Q_{\Phi,1}$	$(D_\mu H^\dagger H)(H^\dagger D^\mu H)$	$c_{\Phi,1}$
Q_{WW}	$H^\dagger W_{\mu\nu}^I W^{I,\mu\nu} H$	c_{WW}
Q_{BB}	$H^\dagger B_{\mu\nu} B^{\mu\nu} H$	c_{BB}
Q_{BW}	$H^\dagger B_{\mu\nu} \tau^I W^{I,\mu\nu} H$	c_{BW}
Q_{2JW}	$J_{W\mu}^I J_W^{I\mu}$	c_{2JW}
Q_{2JB}	$J_{B\mu} J_B^\mu$	c_{2JB}

where H stands for the SM Higgs doublet and $W_{\mu\nu}^I$ and $B_{\mu\nu}$ are the $SU(2)_L$ and $U(1)_Y$ field strength tensors, respectively. τ^I are the Pauli matrices. As mentioned above, the fermionic operators involve the SM fermion currents,

$$J_B^\mu = g' \sum_{f \in \{q,l,u,d,e\}} \sum_a Y_f \bar{f}_a \gamma^\mu f_a \quad \text{and} \quad J_W^{I\mu} = \frac{g}{2} \sum_{f \in \{q,l\}} \sum_a \bar{f}_a \gamma^\mu \tau^I f_a, \quad (2.2)$$

with Y_f standing for the fermion f hypercharge, q and l are the quark and lepton doublets, and u , d , and e represent the

fermion singlets. The sum over a is over generations while the $SU(2)_L$ and $U(1)_Y$ gauge couplings are g and g' , respectively.

The full basis of dimension-eight operators for universal theories was presented in Ref. [58] and contains 175 operators of which 86 bosonic operators with higher derivatives can be rotated into fermionic operators by the EOM as shown in Ref. [58]. From those we identify 11 independent bosonic operators contributing to the weak boson propagators of which eight can be rotated by the EOM into combinations involving ten fermionic operators. We list in Table II those independent dimension-eight operators relevant for EWPD and Drell-Yan analyses. The EOMs imply two relations connecting the coefficients of the following operators in the *rotated basis*

$$c_{\psi^2 H^2 D^3}^{(1)} = -\frac{g'}{2} c_{\psi^4 D^2}^{(2)}, \quad c_{\psi^2 H^2 D^3}^{(2)} = -\frac{g}{2} c_{\psi^4 D^2}^{(3)}. \quad (2.3)$$

As discussed in Ref. [36] the EWPO and DY analyses cannot constrain all 17 Wilson coefficients that enter the analyses at order $\mathcal{O}(1/\Lambda^4)$. However, in Ref. [36] we showed that after carefully accounting for both the direct contributions and the indirect effects induced by the finite renormalization of the SM parameters [39] up to $\mathcal{O}(\Lambda^{-4})$, one can construct five effective combinations of Wilson coefficients

TABLE II. C and P conserving USMEFT dimension-eight operators that contribute to DY and EWPO in the rotated basis and their respective Wilson coefficients.

Operator basis		Wilson coefficient
$Q_{H^6}^{(2)}$	$(H^\dagger H)(H^\dagger \tau^I H)(D_\mu H)^\dagger \tau^I D^\mu H$	$c_{H^6}^{(2)}$
$Q_{WBH^4}^{(1)}$	$(H^\dagger H)(H^\dagger \tau^I H) W_{\mu\nu}^I B^{\mu\nu}$	$c_{WBH^4}^{(1)}$
$Q_{W^2 H^4}^{(3)}$	$(H^\dagger \tau^I H)(H^\dagger \tau^J H) W_{\mu\nu}^I W^{J\mu\nu}$	$c_{W^2 H^4}^{(3)}$
$Q_{\psi^4 D^2}^{(2)}$	$D^\alpha J_B^\mu D_\alpha J_{B\mu}$	$c_{\psi^4 D^2}^{(2)}$
$Q_{\psi^2 H^2 D^3}^{(1)}$	$i(D^\mu J_B^\nu + D^\nu J_B^\mu) \times (D_{(\mu} D_{\nu)} H^\dagger H - H^\dagger D_{(\mu} D_{\nu)} H)$	$c_{\psi^2 H^2 D^3}^{(1)}$
$Q_{\psi^4 D^2}^{(3)}$	$D^\alpha J_W^{I\mu} D_\alpha J_{W\mu}^I$	$c_{\psi^4 D^2}^{(3)}$
$Q_{\psi^2 H^2 D^3}^{(2)}$	$i(D^\mu J_W^{I\nu} + D^\nu J_W^{I\mu}) \times (D_{(\mu} D_{\nu)} H^\dagger \tau^I H - H^\dagger \tau^I D_{(\mu} D_{\nu)} H)$	$c_{\psi^2 H^2 D^3}^{(2)}$
$Q_{\psi^2 H^4 D}^{(1)}$	$iJ_B^\mu (H^\dagger \overleftrightarrow{D}_\mu H)(H^\dagger H)$	$c_{\psi^2 H^4 D}^{(1)}$
$Q_{\psi^2 H^4 D}^{(2)}$	$iJ_W^{I\mu} [(H^\dagger \overleftrightarrow{D}_\mu H)(H^\dagger H) + (H^\dagger \overleftrightarrow{D}_\mu H)(H^\dagger \tau^I H)]$	$c_{\psi^2 H^4 D}^{(2)}$
$Q_{\psi^2 H^4 D}^{(4)}$	$\epsilon^{IJK} J_W^{I\mu} (H^\dagger \tau_\mu^J H) D_\mu (H^\dagger \tau^K H)$	$c_{\psi^2 H^4 D}^{(4)}$
$Q_{\psi^4 H^2}^{(4)}$	$J_B^\mu J_{B\mu} (H^\dagger H)$	$c_{\psi^4 H^2}^{(4)}$
$Q_{\psi^4 H^2}^{(5)}$	$J_W^{I\mu} J_{W\mu}^I (H^\dagger H)$	$c_{\psi^4 H^2}^{(5)}$
$Q_{\psi^4 H^2}^{(7)}$	$J_W^{I\mu} J_{B\mu} (H^\dagger \tau^I H)$	$c_{\psi^4 H^2}^{(7)}$

$$\bar{c}_{BW} \equiv c_{BW} \left(1 - \frac{\hat{v}^2}{\Lambda^2} (c_{WW} + c_{BB}) \right) + \frac{1}{2} \left[c_{WBH^4}^{(1)} + \frac{\hat{e}}{2\hat{s}} c_{\psi^2 H^4 D}^{(1)} + \frac{\hat{e}}{\hat{c}} c_{\psi^2 H^4 D}^{(2)} \right] \frac{\hat{v}^2}{\Lambda^2}, \quad (2.4)$$

$$\bar{c}_{\Phi,1} \equiv c_{\Phi,1} + \left[c_{H^6}^{(2)} - \frac{\hat{e}}{\hat{c}} c_{\psi^2 H^4 D}^{(1)} - \frac{\hat{e}}{\hat{s}} (c_{\psi^2 H^4 D}^{(2)} - c_{\psi^2 H^4 D}^{(4)}) \right] \frac{\hat{v}^2}{\Lambda^2}, \quad (2.5)$$

$$\bar{c}_{W^2 H^4}^{(3)} \equiv c_{W^2 H^4}^{(3)} + \frac{\hat{e}}{2\hat{s}} (c_{\psi^2 H^4 D}^{(2)} - c_{\psi^2 H^4 D}^{(4)}) \bar{c}_{2JW}, \quad (2.6)$$

$$\bar{c}_{2JW} \equiv c_{2JW} \left(1 - 2 \frac{\hat{v}^2}{\Lambda^2} c_{WW} \right) + c_{\psi^4 H^2}^{(5)} \frac{\hat{v}^2}{2\Lambda^2} \equiv - \frac{2\hat{s}^2}{\hat{e}^2} \bar{\Delta}_{4F}, \quad (2.7)$$

$$\bar{c}_{2JB} \equiv c_{2JB} \left(1 - 2 \frac{\hat{v}^2}{\Lambda^2} c_{BB} \right) + c_{\psi^4 H^2}^{(4)} \frac{\hat{v}^2}{2\Lambda^2}, \quad (2.8)$$

which together with the coefficients of the dimension-eight operators $c_{\psi^4 H^2}^{(7)}$, $c_{\psi^4 D^2}^{(2)}$, and $c_{\psi^4 D^2}^{(3)}$, can be independently determined by studying the invariant mass distribution of the lepton pairs produced in DY processes.

For the sake of completeness, we summarize the basic elements and the final expressions for the fermion couplings to the electroweak gauge bosons and the four-fermion amplitudes relevant for the EWPO and DY analyses expressed in terms of these eight effective coefficients in Appendix A. In particular, notice that in Eq. (2.7) we have introduced $\bar{\Delta}_{4F}$, which is an effective parameter defined in the literature containing the contribution of fermionic operators to the muon decay width [see Eq. (A1)].

III. BSM SCENARIOS

In order to study the potential of the HL-LHC to unveil signals of universal NP we study two specific models. The first is the “mirror- $U(1)_Y$ model” where a copy of the SM $U(1)_Y$ field is assumed, the second is a mirror- $SU(2)_L$ model where we instead study a copy of the SM $SU(2)_L$ field. In both cases the SM fields have the same charges under the new $U(1)$ ($SU(2)$) gauge field as to the SM gauge field for $U(1)_Y$ ($SU(2)_L$).

In this Section we discuss how the matching up to dimension-eight operators is achieved for both of these models. In order to do so we establish our conventions by writing out the relevant Lagrangians coupling the NP to the SM degrees of freedom. This allows us to, in later Sections, compare the UV behavior of the model where the NP degrees of freedom are allowed to propagate with that of the infrared theory or EFT where the new degrees of freedom have been integrated out.

A. Mirror- $U(1)_Y$ model

We begin by discussing the mirror- $U(1)_Y$ model. In this case we have a duplicate of the SM $U(1)_Y$ field we write as X^μ . For a generic SM field F , the covariant derivative of the field in the UV is given by,

$$\bar{D}_\mu F = (\partial_\mu - ig' Y_F B_\mu - ig' \beta Y_F X_\mu + \dots) F, \quad (3.1)$$

$$= D_\mu F - ig' \beta Y_F X_\mu F. \quad (3.2)$$

In this expression the ellipsis represents the coupling to the other SM gauge bosons. We have defined the covariant derivative in the UV theory to be $\bar{D}$ and that involving only SM gauge bosons to be D . As our goal is to integrate out the X field, it is more convenient to factor the X dependence out of the covariant derivative. Notice that, rather than introducing a separate gauge coupling, we measure the strength of the coupling of X in units of g' with the coupling scaling β , i.e., $g_X = g' \beta$.

For the fermionic sector the explicit X field dependence is given by,

$$i\bar{\psi} \not{D} \psi = i\bar{\psi} \not{D} \psi + \beta X_\mu J_B^\mu, \quad (3.3)$$

where J_B is one of the SM fermionic currents appearing in Eq. (2.2). The fact that the new X field only couples to the SM fermions through this current is precisely the statement that this UV theory is “Universal.” The covariant derivative on the right does not have a bar as it involves only the SM gauge bosons. Performing the same for the Higgs kinetic term yields,

$$(\bar{D}_\mu H)^\dagger (\bar{D}^\mu H) = (D_\mu H)^\dagger (D^\mu H) + ig' \beta Y_H (H^\dagger \vec{D}_\mu H) X^\mu + g'^2 \beta^2 Y_H^2 (H^\dagger H) X_\mu X^\mu \quad (3.4)$$

$$\equiv (D_\mu H)^\dagger (D^\mu H) + g' \beta Y_H H_\mu X^\mu + g'^2 \beta^2 Y_H^2 (H^\dagger H) X_\mu X^\mu. \quad (3.5)$$

The last term, proportional to X^2 , is irrelevant for dimension-six matching and therefore commonly omitted [59,60]. In our case we must retain this term as it is relevant for matching at dimension-eight. Above we have implicitly defined the Higgs current we use for simplifying our expressions below,

$$\mathcal{H}^\mu = i(H^\dagger \vec{D}_\mu H). \quad (3.6)$$

Taking into account the above discussion our full UV Lagrangian is given by,

$$\mathcal{L} = \mathcal{L}_{\text{SM}} - \frac{1}{4} X_{\mu\nu} X^{\mu\nu} + \frac{1}{2} M_1^2 X_\mu X^\mu + \beta X_\mu J_B^\mu + g' \beta Y_H X_\mu \mathcal{H}^\mu + g'^2 \beta^2 (H^\dagger H) X_\mu X^\mu. \quad (3.7)$$

We do not include a kinetic mixing term, $X_{\mu\nu}B^{\mu\nu}$, as in [16]. This is because, in the case of the mirror model where the charges of the SM fields under the new $U(1)$ are the same as under the SM $U(1)$, we can remove the mixing term by a field redefinition.¹ The mass, M_1 , denotes specifically the explicit mass appearing in the mirror- $U(1)_Y$ Lagrangian above. The physical masses in the UV model, used for the full theory simulations, are calculated below the electroweak symmetry breaking scale after diagonalizing mass mixing effects. We present our calculation for the resulting heavy and light neutral gauge boson masses up to order $1/M_1^2$, while the full analytic expression is omitted due to its length,

$$\bar{M}_1^2 = M_1^2 \left(1 + \frac{g'^2 v^2}{4M_1^2} \beta^2 + \dots \right), \quad (3.8)$$

$$\bar{m}_Z^2 = \frac{(g'^2 + g^2)v^2}{4} \left(1 - \frac{g'^2 v^2}{4M_1^2} \beta^2 + \dots \right). \quad (3.9)$$

We perform the tree-level matching up to dimension-eight using the MATCHETE package [61]. The results of the matching were checked by hand by solving the classical EOM for the X field as described in [62]. The dimension-six effective operators that we obtain are

$$\begin{aligned} \mathcal{L}_{\text{SM}} + \mathcal{L}_6 \rightarrow \mathcal{L}_{\text{SM}} - \frac{\beta^2}{2M_1^2} Q_{2JB} + \frac{g'^2}{M_1^2} \beta^2 Y_H^2 (Q_{\Phi,2} - 2Q_{\Phi,1}) \\ - i \frac{g'^2}{M_1^2} \beta^2 Y_H Y_\psi (H^\dagger \overleftrightarrow{D}_\mu H) (\bar{\psi} \gamma_\mu \psi). \end{aligned} \quad (3.10)$$

The last operator needs to be traded for operators in the *rotated* basis. To consistently include effects up to $\mathcal{O}(1/\Lambda^4)$, it is insufficient to use the EOM, and we need to perform field redefinitions [63]. So, we make the following field redefinition:

$$B_\mu \rightarrow B_\mu + \frac{g' \beta^2}{M_1^2} Y_H \mathcal{H}^\mu. \quad (3.11)$$

This redefinition was made by hand and cross checked by manipulating the MATCHETE output in *Mathematica*. We take into account linear redefinitions of the dimension-six Lagrangian and quadratic terms in the field redefinitions of the renormalizable Lagrangian to capture the dimension-eight terms. We find for the dimension-six Lagrangian,

¹See App. A of [20] for details. The field redefinitions that remove the kinetic mixing also shift the heavy-boson mass and the coupling β , with both shifts controlled by the mixing parameter. Since both parameters are treated as free in the present analysis, this has no effect on our results.

$$\begin{aligned} \mathcal{L}_6 = & -\frac{\beta^2}{2M_1^2} Q_{2JB} + \frac{g'^2}{M_1^2} \beta^2 Y_H^2 (2Q_{\Phi,1} - Q_{\Phi,2}) \\ & - \frac{g'}{M_1^2} \beta^2 Y_H \left[2Q_B + \frac{g'}{2} Q_{BB} + \frac{g}{2} Q_{BW} \right]. \end{aligned} \quad (3.12)$$

The operators, $Q_B = i(D_\mu H)^\dagger B^{\mu\nu} (D_\nu H)$ and $Q_{\Phi,2} = \frac{1}{2} \partial_\mu (H^\dagger H) \partial^\mu (H^\dagger H)$, do not contribute to EWPO or the Drell-Yan process at tree level.

The resulting dimension-eight Lagrangian arising from the field redefinitions as well as tree-level matching was then reduced by making extensive use of Matchete's internal integration-by-parts routines in order to align the dimension-eight operators with the purely bosonic basis presented in [58]. These results were then rewritten in the Murphy basis [64] following the relations presented in [58]. The part of the dimension-eight Lagrangian relevant to our analysis of Drell-Yan in the rotated basis is then given by,

$$\begin{aligned} \mathcal{L}_8 = & -\frac{g^2 g'^2 \beta^2}{8M_1^4} (1 + \beta^2) Q_{H^6}^{(2)} - \frac{g g'^2 \beta^2}{8M_1^4} (1 + \beta^2) Q_{\psi^2 H^4 D}^{(2)} \\ & + \frac{g' \beta^2}{4M_1^4} Q_{\psi^2 H^2 D^3}^{(1)} - \frac{\beta^2}{2M_1^4} Q_{\psi^4 D^2}^{(2)} + \frac{g^2 \beta^4}{4M_1^4} Q_{\psi^4 H^2}^{(4)} + \dots \end{aligned} \quad (3.13)$$

The ellipses denotes 19 additional dimension-eight operators that do not contribute to the Drell-Yan observables considered in this work. The full result of the matching can be found in Eq. (B4) in Appendix B. Rewriting these results in terms of our input parameter scheme and relating them to the barred coefficients as defined in Eqs. (2.4)–(2.8) yields the relations found in Table III.

B. Mirror- $SU(2)_L$ Model

Next we consider new heavy gauge bosons, X_μ^a , corresponding to a mirror $SU(2)$ symmetry. The SM fermionic and Higgs doublets are also doublets under this mirror symmetry and we obtain the following UV Lagrangian:

$$\begin{aligned} \mathcal{L} = \mathcal{L}_{\text{SM}} - \frac{1}{4} (X_{\mu\nu}^a)^2 + \frac{1}{2} M_2^2 (X_\mu^a)^2 + \beta X_\mu^a J_W^{a\mu} + \frac{g}{2} \beta X^{a,\mu} \mathcal{H}_\mu^a \\ + \frac{g^2}{4} \beta^2 (H^\dagger H) X^{a,\mu} X_\mu^a, \end{aligned} \quad (3.14)$$

where, again, β measures the strength of the mirror- $SU(2)$ gauge coupling relative to the coupling g of the SM. In this case the mass of the new field is denoted M_2 in order to distinguish it from that of the mirror- $U(1)_Y$ model. Again, in the full UV model this mass is shifted by mixing effects. To order $1/M_2^2$ the masses of the resulting charged and neutral vectors are given by,

TABLE III. Low-energy Wilson coefficients in the rotated basis generated by integrating out the gauge bosons of the mirror- $U(1)_Y$ and mirror- $SU(2)_L$ models that contribute to DY and EWPO.

Effective coefficients	Mirror- $U(1)_Y$			Mirror- $SU(2)_L$		
	Dim-6	Dim-6 squared	Dim-8	Dim-6	Dim-6 squared	Dim-8
$\bar{c}_{BW}$	$-\frac{\hat{e}^2 \beta^2}{4\hat{s}\hat{c}}$	$-\frac{\hat{e}^4 \hat{v}^2 \beta^4}{8\hat{s}_{4W} M_1^2}$	$-\frac{\hat{e}^4 \hat{v}^2 \beta^2}{16 M_1^2 \hat{c}^3 \hat{s}} (1 + \beta^2)$	$-\frac{\hat{e}^2 \beta^2}{4\hat{s}\hat{c}}$	$-\frac{\hat{e}^4 \hat{v}^2 \beta^4}{64\hat{s}^3 \hat{c} M_2^2} (\frac{1}{\hat{c}_{2W}} - 3)$	$-\frac{\hat{e}^4 \hat{v}^2 \beta^2}{16\hat{s}^3 \hat{c} M_2^2}$
$\bar{c}_{\Phi,1}$	$\frac{\hat{e}^2 \beta^2}{2\hat{c}^2}$	$-\frac{\hat{e}^4 \hat{v}^2 \beta^4}{4\hat{c}^4}$	0	0	0	0
$\bar{c}_{2JB}$	$-\frac{\beta^2}{2}$	$-\frac{\hat{e}^2 \hat{v}^2 \beta^4}{4\hat{c}^2 M_1^2}$	$+\frac{\hat{e}^2 \hat{v}^2 \beta^4}{8\hat{c}^2 M_1^2}$	0	0	0
$\bar{c}_{2JW} \equiv -\frac{2\hat{s}^2}{e^2} \bar{\Delta}_{4F}$	0	0	0	$-\frac{\beta^2}{2}$	$-\frac{\hat{e}^2 \hat{v}^2 \beta^4}{4\hat{s}^2 M_2^2}$	$+\frac{\hat{e}^2 \hat{v}^2 \beta^4}{8\hat{s}^2 M_2^2}$
$\bar{c}_{W^2 H^4}^{(3)}$	...	...	$-\frac{\hat{e}^4 \beta^2}{16\hat{s}^2 \hat{c}^2} (1 + \beta^2)$	...	...	0
$c_{\psi^4 H^2}^{(7)}$	...	...	0	...	...	0
$c_{\psi^4 D^2}^{(2)}$	...	...	$-\frac{\beta^2}{2}$	...	...	0
$c_{\psi^4 D^2}^{(3)}$	...	...	0	...	...	$-\frac{\beta^2}{2}$

$$(\bar{M}_2^{\pm,0})^2 = M_2^2 \left(1 + \frac{g^2 v^2}{4M_2^2} \beta^2 + \dots \right), \quad (3.15)$$

$$\bar{m}_W^2 = \frac{g^2 v^2}{4} \left(1 - \frac{g^2 v^2}{4M_2^2} \beta^2 + \dots \right), \quad (3.16)$$

$$\bar{m}_Z^2 = \frac{(g'^2 + g^2) v^2}{4} \left(1 - \frac{g^2 v^2}{4M_2^2} \beta^2 + \dots \right). \quad (3.17)$$

The Higgs current appearing above is defined as,

$$\mathcal{H}_\mu^a = i(H^\dagger \overleftrightarrow{D}_\mu^a H) = i[H^\dagger \tau^a D_\mu H - (D_\mu H)^\dagger \tau^a H], \quad (3.18)$$

while the fermionic current appears in Eq. (2.2).

We integrate out the heavy gauge bosons using MATCHETE as described in the case of the mirror- $U(1)$ model. In order to rewrite the dimension-six operator basis we first make the W field redefinition,

$$W_\mu^a \rightarrow W_\mu^a + \frac{g}{2M_2^2} \beta^2 \mathcal{H}_\mu^a, \quad (3.19)$$

then subsequently redefine the Higgs doublet as,

$$H \rightarrow H - \frac{g^2}{4M_2^2} \beta^2 H |H|^2. \quad (3.20)$$

The resulting dimension-six Lagrangian is given by,

$$\begin{aligned} \mathcal{L}_6 = & \frac{g^2 \beta^2}{M_2^2} \left(\lambda + \frac{\beta^2 g^2 \mu^2}{16M_2^2} \right) Q_{\Phi^6} - \frac{\beta^2}{2M_2^2} Q_{2JW} - \frac{3g^2 \beta^2}{4M_2^2} Q_{\Phi,2} \\ & - \frac{g\beta^2}{M_2^2} Q_W - \frac{g^2}{4M_2^2} \beta^2 Q_{WW} - \frac{g'g}{4M_2^2} \beta^2 Q_{BW}, \end{aligned} \quad (3.21)$$

where λ is the quartic coupling of the Higgs potential. The operators $Q_W = i(D_\mu H)^\dagger \tau^I W^{I\mu\nu} (D_\nu H)$, $Q_{\Phi,2}$, and $Q_{\Phi^6} = (H^\dagger H)^3$ do not contribute to either EWPO or the Drell-Yan process at tree level. The Higgs field redefinition in Eq. (3.20) also results in a shifted quartic Higgs operator, however our analysis is insensitive to this shift.

After rewriting the dimension eight operators in terms of the rotated basis we find the operators relevant to our analysis are,

$$\begin{aligned} \mathcal{L}_8 = & -\frac{g^2 g'^2 \beta^2}{2M_2^4} (1 + \beta^2) Q_{H^6}^{(2)} + \frac{g^3 g' \beta^2}{8M_2^4} Q_{WBH^4}^{(1)} \\ & - \frac{g^2 g' \beta^2}{2M_2^4} (1 + \beta^2) Q_{\psi^2 H^4 D}^{(1)} + \frac{g^3 \beta^4}{4M_2^4} Q_{\psi^2 H^4 D}^{(2)} \\ & + \frac{g^3 \beta^4}{4M_2^4} Q_{\psi^2 H^4 D}^{(4)} + \frac{g\beta^2}{4M_2^4} Q_{\psi^2 H^2 D^3}^{(2)} + \frac{3g^2 \beta^2}{4M_2^4} Q_{\psi^4 H^2}^{(3)} \\ & + \frac{g^2 \beta^4}{4M_2^4} Q_{\psi^4 H^2}^{(5)} - \frac{\beta^2}{2M_2^4} Q_{\psi^4 D^2}^{(3)} + \dots \end{aligned} \quad (3.22)$$

The full matching results including operators that do not affect our analyses (denoted by the ellipsis above) can be found in Appendix B. The results expressed as barred coefficients in terms of the input parameters appear in Table II.

IV. ANALYSIS FRAMEWORK

As a starting step, prior to our HL-LHC studies, we obtain the parameter regions of the BSM models already excluded at 95% confidence level (CL) by the presently available data. Since each of the BSM scenarios considered in this work generates a single oblique parameter: Y for the mirror- $U(1)_Y$ model and W for the mirror- $SU(2)_L$ model,

TABLE IV. Neutral- and charged-current Drell-Yan data to obtain the present USMEFT bounds.

Channel	Distribution	No. bins	Ranges	Dataset	Integrated luminosity
NC	$\frac{d^2\sigma}{dm_{\ell\ell}d y _{\ell\ell}}$	48	$116 \text{ GeV} \leq m_{\ell\ell} \leq 1.5 \text{ TeV}$ $0 \leq y_{\ell\ell} \leq 2.4$	ATLAS 8 TeV	20.3 fb^{-1} [2]
NC	$\frac{dN_{\text{ev}}}{dm_{e^+e^-}}$	20	$250 \text{ GeV} \leq m_{e^+e^-} \leq 5 \text{ TeV}$	ATLAS 13 TeV	139 fb^{-1} [7]
NC	$\frac{dN_{\text{ev}}}{dm_{\mu^+\mu^-}}$	20	$250 \text{ GeV} \leq m_{\mu^+\mu^-} \leq 5 \text{ TeV}$	ATLAS 13 TeV	139 fb^{-1} [7]
NC	$\frac{dN_{\text{ev}}}{dm_{e^+e^-}}$	20	$300 \text{ GeV} \leq m_{e^+e^-} \leq 6 \text{ TeV}$	CMS 13 TeV	137 fb^{-1} [8]
NC	$\frac{dN_{\text{ev}}}{dm_{\mu^+\mu^-}}$	20	$300 \text{ GeV} \leq m_{\mu^+\mu^-} \leq 7 \text{ TeV}$	CMS 13 TeV	137 fb^{-1} [8]
CC	$\frac{d\sigma}{dm_T}$	20	$200 \text{ GeV} \leq m_{T,\ell\nu} \leq 5 \text{ TeV}$	ATLAS 13 TeV	140 fb^{-1} [6]
CC	$\frac{dN}{dm_T}$	20	$440 \text{ GeV} \leq m_{T,e\nu} \leq 7 \text{ TeV}$	CMS 13 TeV	138 fb^{-1} [10]
CC	$\frac{dN}{dm_T}$	20	$600 \text{ GeV} \leq m_{T,\mu\nu} \leq 7 \text{ TeV}$	CMS 13 TeV	138 fb^{-1} [10]

the strongest bounds arise from the high- p_T Drell-Yan tails as emphasized previously in Ref. [65]. Therefore, we perform an analysis of the Run 1 and 2 LHC data in Table IV in the context of these two models. We obtain the following limits on the model parameters:

$$\text{mirror} - U(1)_Y: |\beta| \leq 0.23 \frac{M_1}{\text{TeV}} \quad [95\% \text{CL}], \quad (4.1)$$

$$\text{mirror} - SU(2)_L: |\beta| \leq 0.11 \frac{M_2}{\text{TeV}} \quad [95\% \text{CL}]. \quad (4.2)$$

These current bounds on the models delimit the target region of model parameters for our HL-LHC forecasts. Furthermore, we estimate the reach of the HL-LHC for those models from the on-shell production of the new gauge bosons. This is shown in Fig. 1 where we plot the on-shell production cross section times the leptonic branching ratio into either electron or muon final states as a function of $M = M_{1(2)}$ for the mirror- $U(1)_Y$ (mirror- $SU(2)_L$) for SM couplings $\beta = 1$. Multiplying by the acceptance times efficiencies estimated in Ref. [66] to be in average of the order of 75% (65%) in the NC and 65% in CC one can estimate the number of signal events with 3000 fb^{-1} from the on-shell production of the new gauge bosons. From Ref. [66] we also read a number of 4 (2) background events expected for the NC (CC) channel, which naively requires 18 (13) events for a 5σ discovery. As the dominant β dependence arises from the coupling scaling, the expected number of events will grow as β^2 yielding a 5σ discovery for $\beta \gtrsim 0.1(0.04), 0.25(0.12), 0.6(0.32), 1.6(0.85), 4(2.2)$ for $M_1(M_2) = 4, 5, 6, 7, 8 \text{ TeV}$, respectively. As we will see below with our simulations based on the invariant mass distributions of the full SM + BSM signals we find that the values of β for a given M for which a 5σ excess in the high energy invariant mass bins are a bit higher than those naïve expectations.

As sketched in the introduction, our starting point is the generation of pseudodata for the DY processes following the predictions of the BSM models for grid of values of the model parameters, i.e., the masses and couplings, and confronting it with the expectations within the USMEFT including all the relevant operator coefficients. To this end the BSM model pseudodata as well as the USMEFT predictions were obtained with MadGraph5_aMC@NLO [67] at leading order in QCD and QED. The required UFO files for the two BSM models considered as well as the USMEFT relevant operators were created with the help of FeynRules [68,69]. Parton shower and hadronization was performed using PYTHIA8 [70], and the fast detector simulation was carried out with DELPHES [71]. Jet analyses were performed using FASTJET [72]. Exclusively for the presently available ATLAS NC data [7], the detector response was simulated using Rivet [73,74], with the analysis code provided by the experimental collaboration.

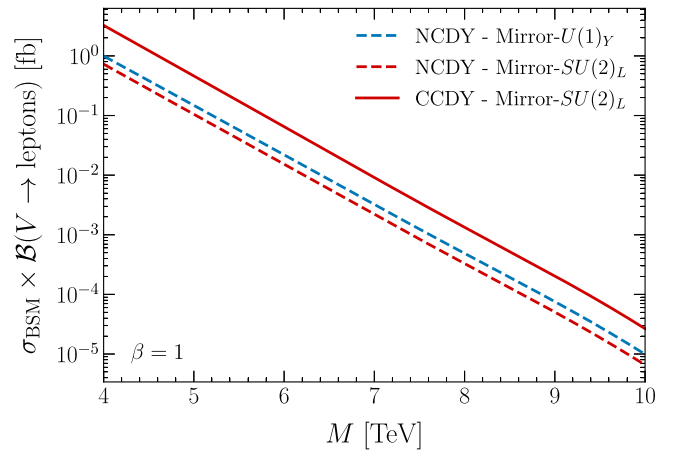


FIG. 1. Production cross section times leptonic branching fraction of the mirror gauge bosons as a function of $M = M_{1(2)}$ for $\beta = 1$.

In the analysis of NC (CC) DY process we employed the invariant (transverse) mass distribution to perform the USMEFT fit to the corresponding pseudodata. The invariant (transverse) mass binning used is from 0.5 (1.0) TeV to 3.5 TeV and an overflow bin for all events with invariant (transverse) masses above 3.5 TeV. Defining as $N_i^{\text{P-D,NC(CC)}}$ the pseudodata in bin i for the NC (CC) process for a given BSM model and parameters, and $N_i^{\text{EFT,NC(CC)}}$ the expected number of events in that same bin for a given value of USMEFT WC, the chi-squared functions for the DY distributions are constructed as

$$\chi_{\text{DY,NC(CC)}}^2(M, \beta | \vec{c}_{\text{NC(CC)}}) = \sum_i \frac{[N_i^{\text{P-D,NC(CC)}}(M, \beta) - N_i^{\text{EFT,NC(CC)}}(\vec{c}_{\text{NC(CC)}})]^2}{N_i^{\text{P-D,NC(CC)}}(M, \beta) + \delta_{\text{syst},i}^2}, \quad (4.3)$$

where

$$N_i^{\text{P-D,NC(CC)}}(M, \beta) = \mathcal{L} \sigma_i^{\text{P-D,NC(CC)}}(M, \beta), \quad (4.4)$$

$$N_i^{\text{EFT,NC(CC)}}(\vec{c}_{\text{NC(CC)}}) = \mathcal{L} \sigma_i^{\text{EFT,NC(CC)}}(\vec{c}_{\text{NC(CC)}}), \quad (4.5)$$

$$\delta_{\text{syst},i} = 0.015 \times m_i [\text{TeV}] \times N_i^{\text{P-D,NC(CC)}}. \quad (4.6)$$

$\mathcal{L}$ denotes the luminosity, which we take to be 3000 fb^{-1} , while $\sigma_i^{\text{P-D,NC(CC)}}(M, \beta)$ stands for the BSM cross section for model parameters (M, β) in that bin and $\sigma_i^{\text{EFT,NC(CC)}}(\vec{c}_{\text{NC(CC)}})$ is the corresponding predicted cross section within USMEFT for WC

$$\vec{c}_{\text{NC}} \equiv (\bar{c}_{BW}, \bar{c}_{\Phi,1}, \bar{\Delta}_{4F}, \bar{c}_{2JB}, c_{\psi^4 D^2}^{(2)}, c_{\psi^4 D^2}^{(3)}, c_{\psi^4 H^2}^{(7)}),$$

and

$$\vec{c}_{\text{CC}} \equiv (\bar{c}_{BW}, \bar{c}_{\Phi,1}, \bar{\Delta}_{4F}, \bar{c}_{W^2 H^4}^{(3)}, c_{\psi^4 D^2}^{(3)}).$$

Both cross sections include the detector effects and event selection cuts. Also following Ref. [66], we have included a systematic uncertainty of 1.5% m_i (TeV) to simulate energy resolution [75]. m_i is the value of the corresponding invariant (transverse) mass in the center of the bin i for all bins but for the overflow bin for which we take as m_i the value of the mass in the lower edge of the bin.

In the EWPO analyses, we include 12 Z-pole observables [76]: Γ_Z , σ_h^0 , $\mathcal{A}_\ell(\tau^{\text{pol}})$, R_ℓ^0 , $\mathcal{A}_\ell(\text{SLD})$, $A_{\text{FB}}^{0,l} R_c^0$, R_b^0 , $\mathcal{A}_c$, $\mathcal{A}_b$, $A_{\text{FB}}^{0,c}$, and $A_{\text{FB}}^{0,b}$ and two W pole observables M_W and Γ_W taken from [76,77]. The SM predictions and their uncertainties due to variations of the SM parameters were extracted from [78]. In order to eliminate the fluctuations associated with the central values of the real EWPO results, in our forecasts, we simulate the central values of EWPOs

following the predictions of the BSM model for the given set of parameters (M, β) . With those we define a function $\chi_{\text{EWPO}}^2(M, \beta | \vec{c}_{\text{EWPO}})$ which depends on the WC,

$$\vec{c}_{\text{EWPO}} \equiv (\bar{c}_{BW}, \bar{c}_{\Phi,1}, \bar{\Delta}_{4F}, \bar{c}_{W^2 H^4}^{(3)}, c_{\psi^4 D^2}^{(2)}, c_{\psi^4 D^2}^{(3)}).$$

In fact, EWPO data is capable of constraining only four combinations of these parameters at order $1/\Lambda^4$ [36], however the combination of the EWPO and DY data resolves the blind directions of the EWPO.

As mentioned in the introduction, we perform our USMEFT fits in four different variants of the analysis:

- (i) In the first case, we truncate the expansion at $\mathcal{O}(\Lambda^{-2})$ by considering only the first two terms in Eq. (2.1). In this scenario referred as $(d=6)$ analysis, we include only the four coefficients of the dimension-six operators and keep their effects in the observables at linear order in the coefficients.
- (ii) In the second analysis, we include all contributions in Eq. (2.1) but the last term. This accounts to consider only the four coefficients of the dimension-six operators and to keep their effects in the observables up to quadratic order in their WC. We refer to this case as $(d=6)^2$ analysis.
- (iii) In the third case, we perform a full $\mathcal{O}(\Lambda^{-4})$ analysis by considering all terms in Eq. (2.1). This case, dubbed $(d=8)$ analysis, contains all the contributions in the previous scenario as well as linear one coming from dimension-eight operators.
- (iv) The fourth analysis is the same as the third but employing the two coefficient combinations in Eq. (5.1). We refer to this case as $(d=8)'$ analysis.

The last ingredient in our studies, which serves the purpose of answering our first question on the HL-LHC potential to uncover NP signals using the EFT framework, we defined for each WC, c_j , the quantity

$$f_{\text{NP}}(c_j) \equiv \frac{|c_j^{\text{best fit}}|}{\Delta c_j} \quad (4.7)$$

where Δc_j is half of the difference between the upper and lower 68% CL limits on c_j . This quantity measures the deviation from the SM prediction ($c_j = 0$) in number of 1σ deviations.

V. RESULTS

Our starting point is the generation of the pseudo-data for the invariant (transverse) mass distribution for DY NC (CC). We generated DY pseudodata for gauge boson masses between 4 and 10 TeV and several values of the coupling scaling β . We show in Fig. 2 examples of these distributions for NC pseudodata for the mirror- $U(1)_Y$ [mirror- $SU(2)_L$] model in the left (right) panel for several values of the model parameters. As a starting step, taking into account that the

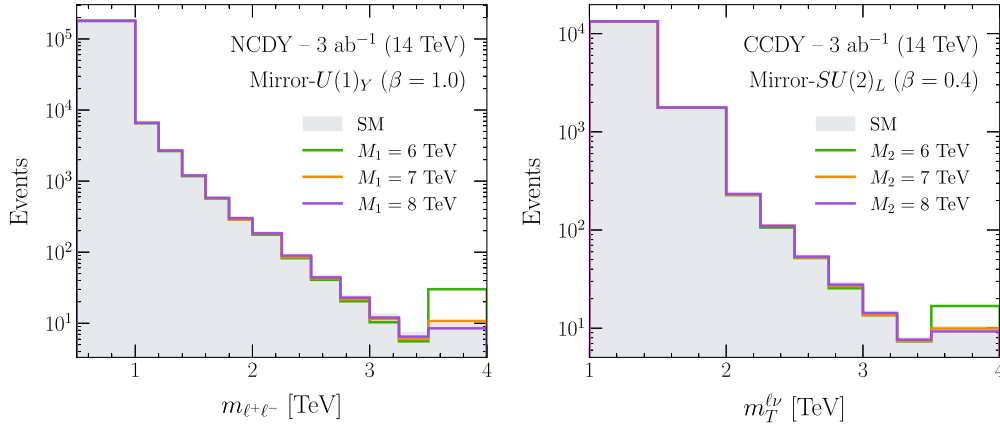


FIG. 2. Left (right) panel displays the expect number of events distribution as a function of the invariant (transverse) mass for the mirror- $U(1)_Y$ (mirror- $SU(2)_L$) extension of the SM. We assumed an integrated luminosity of 3 ab^{-1} .

EFT contribution grows at high invariant (transverse) masses, as illustrated in the figure, we obtained the range of model parameters where a simple counting experiment using the events in the overflow bin is enough to guarantee a 5σ discovery of NP. We do so because, certainly, in this parameter region a new rebinning or a resonance search would be more efficient than an EFT analysis to confirm and identify the NP states.

Next, for each value of the model parameters for which we have generated the pseudodata we performed an EFT fit of the DY pseudodata and EWPO without imposing any assumption on the eight WCs involved. From each of the analyses we find the best fit value and uncertainties of the eight WCs. An example is shown in Fig. 3, where we plot the one-dimension projections of $\Delta\chi^2$ for different orders considered in the analysis for one example of pseudodata generated for each of the models. To guide the eye, we plot as a vertical dotted line the SM value “0” of the coefficients.

Focusing on the result obtained at $(d=6)$ (blue lines) and $(d=6)^2$ (black dashed lines), we see that both analyses lead to similar results indicating that the bounds are dominated by the dimension-six cross section. In particular we see that the coefficients for which the best fit value is most prominently away from zero at this order are $\bar{c}_{2JB}$ for the fit with the mirror- $U(1)_Y$ pseudodata and $\bar{\Delta}_{4F}$ (i.e., $\bar{c}_{2JW}$) for the mirror- $SU(2)_L$ pseudodata. This is expected because these are the coefficients generated by direct matching at tree level (see Table II). Furthermore the four-vertex contact interactions associated with $\bar{c}_{2JB}^{(i)}$ and $\bar{c}_{2JW}$ presents the fastest growth with subprocess center-of-mass energy. The other minor deviations arise from the smaller contributions to indirect finite-renormalization effects, or by correlations between the determined coefficients in the fit. Precisely on that respect, one can see that when including also the dimension-eight operators the sensitivity to the characteristic non-zero value of $\bar{c}_{2JB}$ for the fit with the mirror- $U(1)_Y$ pseudodata is lost: the corresponding $\Delta\chi^2$ is the red dotted line at the bottom

of those panels which is almost flat zero within the plotted range. The reason for this loss of sensitivity is that, as shown in Ref. [36], the analysis at $1/\Lambda^4$ order presents a strong correlation between $\bar{c}_{2JB}$ and $c_{\psi^4 H^2}^{(7)}$ limiting our ability to precisely determine both of the WCs simultaneously. We see this in the left panel of Fig. 4 where we plot their allowed regions which, as seen on the figure, display a very strong correlation.

These two operators enter DY in the four-fermion NC contact amplitudes dominantly in the linear combination $\bar{c}_{2JB} \frac{\hat{v}^2}{\Lambda^2} - \frac{s}{2e} c_{\psi^4 H^2}^{(7)} \frac{\hat{v}^4}{\Lambda^4}$ leading to the very strong correlation observed and the substantial weakening of the bounds on c_{2JB} when compared to the analysis performed at $(d=6)$ or $(d=6)^2$. However, it is important to notice that this correlation appears in the USMEFT prediction irrespective of the pseudodata fitted. This allows to perform our analysis *blindly*, this is, deciding *a priori* to use a different basis for those two coefficients in which one of the new coefficients is chosen to be the combination that dominantly contributes to DY and the other the orthogonal one. To this end we define the “prime” combinations

$$\begin{aligned} \bar{c}'_{2JB} \frac{\hat{v}^2}{\Lambda^2} &= \bar{c}_{2JB} \frac{\hat{v}^2}{\Lambda^2} + 0.33 c_{\psi^4 H^2}^{(7)} \frac{\hat{v}^4}{\Lambda^4}, \\ c_{\psi^4 H^2}^{(7)'} \frac{\hat{v}^4}{\Lambda^4} &= c_{\psi^4 H^2}^{(7)} \frac{\hat{v}^4}{\Lambda^4} - 0.33 \bar{c}_{2JB} \frac{\hat{v}^2}{\Lambda^2}. \end{aligned} \quad (5.1)$$

For all other operator coefficients the prime and unprimed coefficients are the same.

As seen in the right panel in Fig. 4 in these variables the correlation disappears and consequently the coefficient $\bar{c}'_{2JB}$ can be more precisely determined as can be seen also in the one-dimensional projection for this variable for the $(d=8)'$ case in Fig. 3 (full red lines). With all this we find that including the dimension-8 operators, still allows for the identification of the operator coefficients with dominant NP

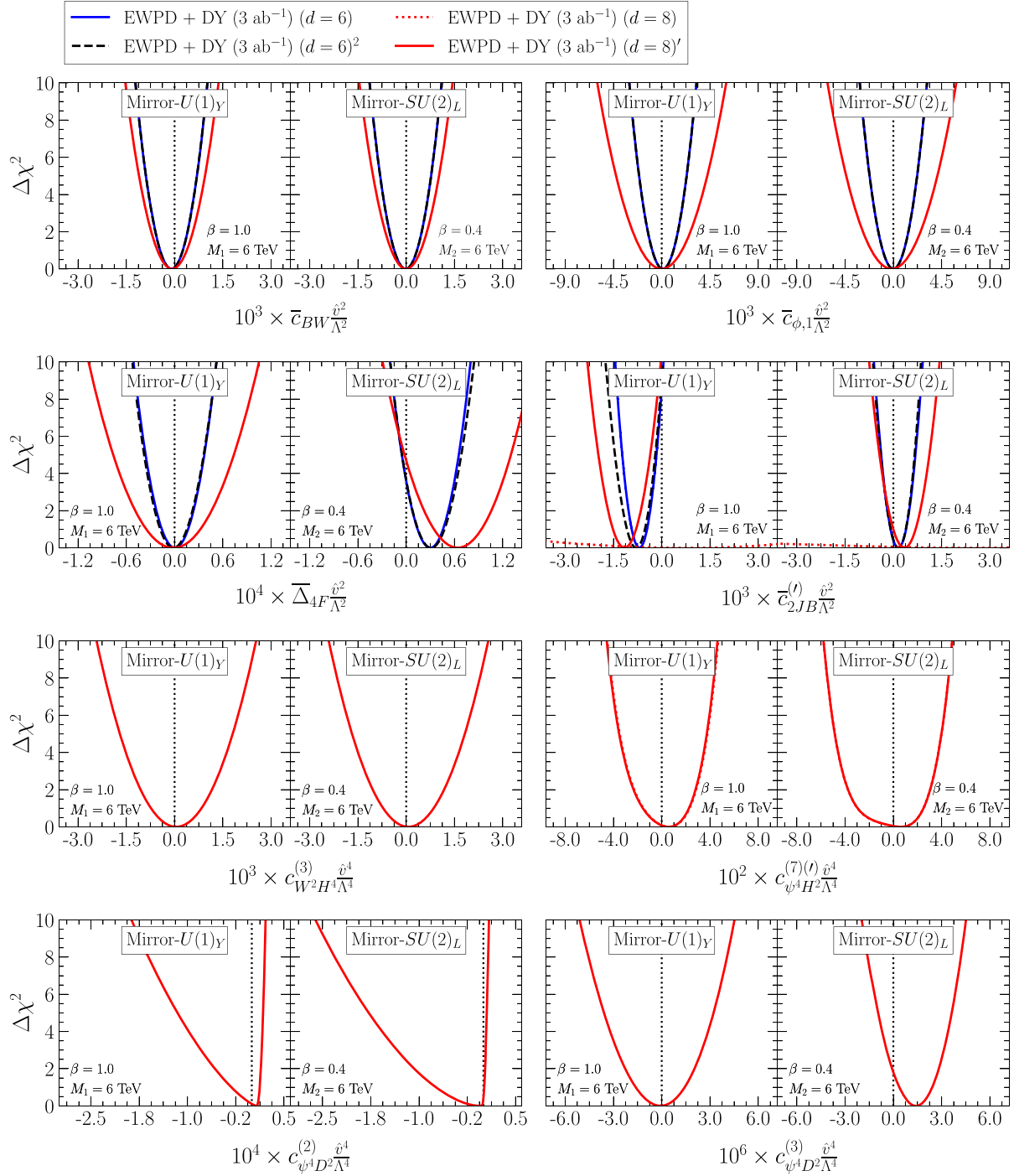


FIG. 3. One-dimensional projections of $\Delta\chi^2$ after marginalization over all other Wilson coefficients for the different analyses performed as labeled in the figure. Notice that for all coefficients but $\bar{c}_{2JB}$ and $c_{\psi^4 H^2}^{(7)}$ the curves for the analysis $(d = 8)$ and $(d = 8)'$ fully overlap so only the full red lines are seen (see text for details).

effects, $\bar{c}'_{2JB}$ for mirror- $U(1)_Y$ pseudodata fits, and $\bar{c}_{2JW}$ for mirror- $SU(2)_L$ pseudodata fits while the precision in the determination of these coefficients is only mildly affected despite the larger parameters space.

From Fig. 3, we also anticipate that, for any point in the model parameter space, those two couplings are the most sensitive to the NP measure f_{NP} . As mentioned above,

this is expected since the four-point contact interactions associated with $\bar{c}_{2JB}^{(\prime)}$ and $\bar{c}_{2JW}$ exhibit the fastest growth with the subprocess center-of-mass energy. In principle, as shown in Table II, $c_{\psi^4 D^2}^{(2)}$ and $c_{\psi^4 D^2}^{(3)}$ are also generated through tree-level matching of direct effects in the mirror- $U(1)_Y$ and mirror- $SU(2)_L$ models at dimension eight. Despite inducing

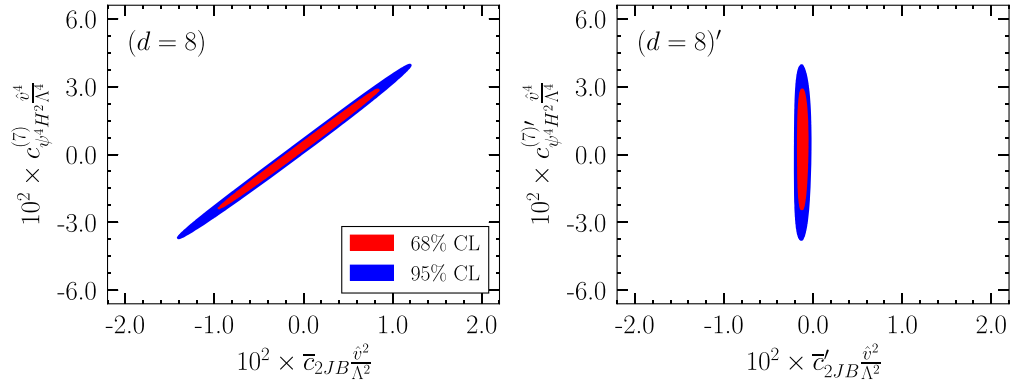


FIG. 4. Left: Allowed regions for the parameters $\bar{c}_{2JB}$ and $c_{\psi^4 H^2}^{(7)}$ for the fit to the pseudodata for mirror- $U(1)_Y$ model with $M_1 = 1$ TeV and $\beta = 1$ at order $(d = 8)$ after marginalization over all other parameters. Right: Same region for the combinations $\bar{c}_{2JB}'$ and $c_{\psi^4 H^2}^{(7)'}$ in Eq. (5.1).

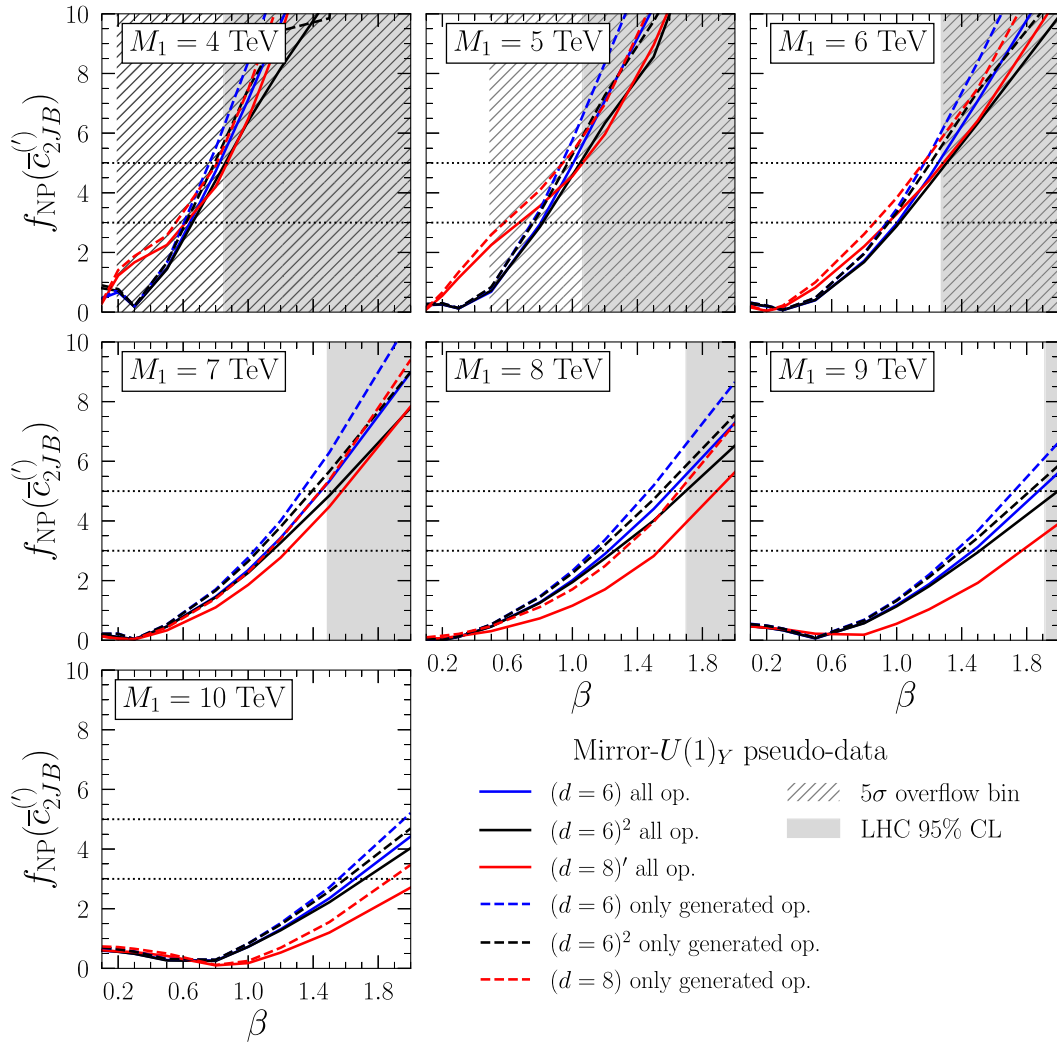


FIG. 5. f_{NP} as a function of β for $c_{2JB}^{(n)}$ obtained with fit to pseudodata generated with the mirror- $U(1)_Y$ model for the masses as labeled in the figure. The full lines correspond to the analysis performed with all the eight relevant operator coefficients, while the dashed lines correspond to the analysis performed including only coefficients of the operators generated in the model (three at $(d = 6)^{(2)}$ and five at $(d = 8)'$). The shaded area is excluded at 95% CL by the present available data while the hashed one is the one where the overflow alone leads to a 5 σ discovery. We present the results for different $1/\Lambda$ orders as indicated. We assumed an integrated luminosity of 3 ab^{-1} .

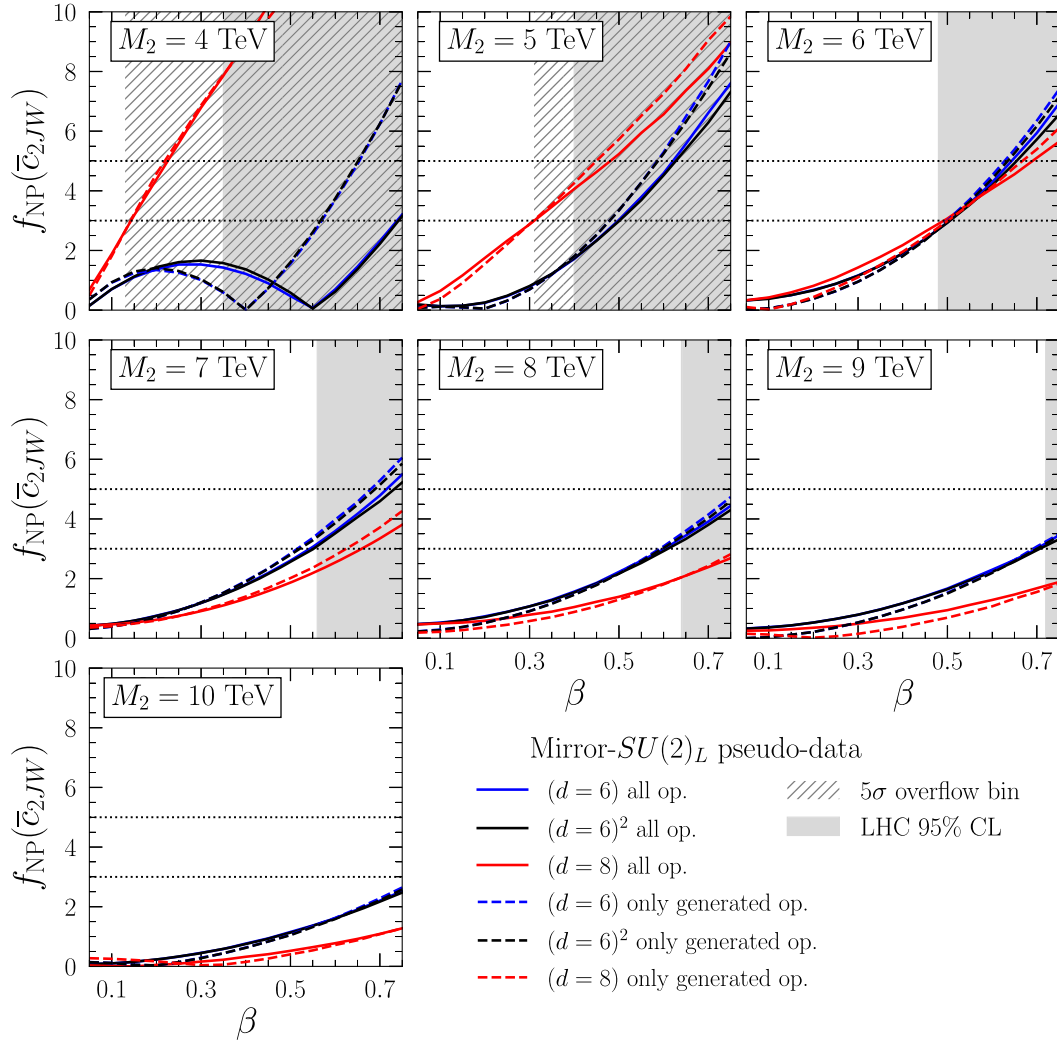


FIG. 6. $f_{\text{NP}}(\bar{c}_{2JW})$ as a function of β obtained with fit to pseudodata generated with the mirror- $SU(2)_L$ model for the masses as labeled in the figure. Conventions are the same as in Fig. 5. In this case the dashed lines correspond to fits including the coefficients of the operators which would be generated in the mirror- $SU(2)_L$ model, (two at $(d=6)^{(2)}$ and three at $(d=8)'$).

energy-growing effects, they are poorly constrained because of the stronger mass suppression.

Figures 5 and 6 depict, respectively, the dependence of $f_{\text{NP}}(\bar{c}_{2JB}^{(i)})$ and $f_{\text{NP}}(\bar{c}_{2JW})$ on β for different values of M_1 and M_2 . First, as expected, the simple analysis of the overflow bin (hashed area) excludes a large fraction of the β range for the smaller mass values (4 and 5 TeV) allowing only small values of β . For the mirror- $U(1)_Y$ case for these lighter masses, the EFT analyses do not lead to any evidence of NP due to the small value of β , while for the mirror- $SU(2)_L$ case the sensitivity is a bit higher but still marginal. In addition, as we can see, for the mirror- $SU(2)_L$ the presently available DY data exclude a larger fraction of the parameter space, since this model contributes to both NC and CC DY processes, especially for masses $M \lesssim 5$ TeV.

The different curves for f_{NP} are obtained with the analyses performed at different orders in the $1/\Lambda$ expansion. Furthermore, to test the robustness of the conclusions we

have performed two variants of the analysis. One in which all the operator coefficients of a given order are included (solid lines), and one in which we artificially fix to zero the coefficients which we know that are not generated by the model (dashed lines). In this case, we only fit three dimension-six coefficients— $\bar{c}_{BW}$, $\bar{c}_{\Phi,1}$, $\bar{c}_{2JB}^{(i)}$ —and two additional dimension-eight coefficients— $c_{\psi^4 D^2}^{(2)}$ and $\bar{c}_{W^2 H^4}^{(3)}$ —in the fits to the mirror- $U(1)_Y$ pseudodata shown in Fig. 5, and two dimension-six coefficients— $\bar{c}_{BW}$, and $\bar{c}_{2JW}$ —and one additional dimension-eight coefficient— $c_{\psi^4 D^2}^{(3)}$ —in the fits to mirror- $SU(2)_L$ pseudodata shown in Fig. 6.

Comparing the results obtained at the different orders in the operator expansion, we see in Fig. 5 that for the fits to mirror- $U(1)_Y$ pseudodata the $d=6$ and $(d=6)^2$ results are rather similar and the $(d=8)'$ analyses reduce f_{NP} by about 10%–20% for higher M_1 masses. Altogether, we find that an EFT 5σ signal (3σ evidence) can be observed

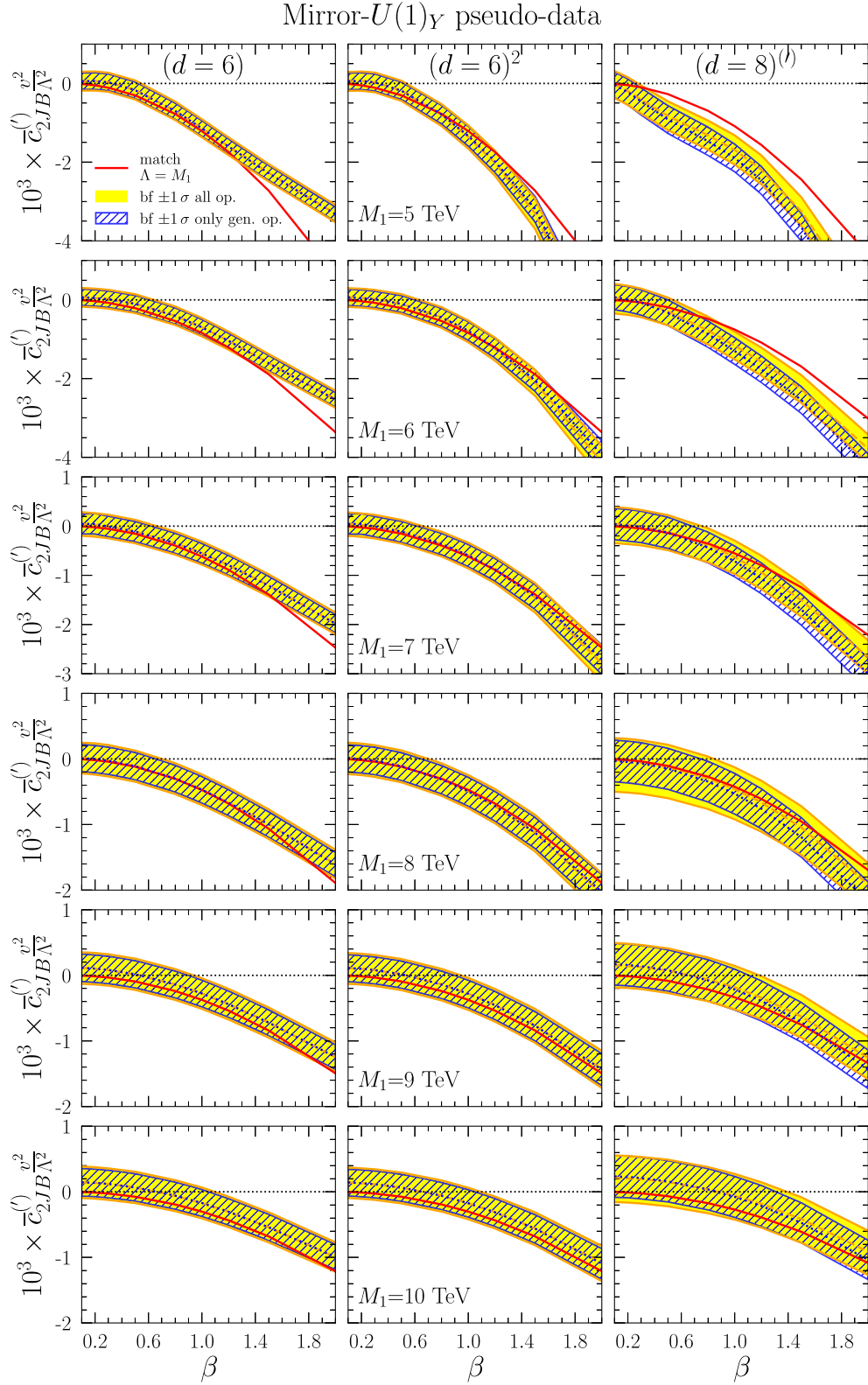


FIG. 7. Allowed one sigma regions for $c_{2JB}^{(l)}$ as a function of β obtained with fit to pseudodata generated with mirror- $U(1)_Y$ model for the masses as labeled in the figure. The three columns correspond to the analysis performed at different orders in the operator expansion. In each panel we show as yellow band the 1σ allowed region around the best fit (orange dashed line) obtained with the analysis including all operators in that order of the expansion. The blue hatched band corresponds to the 1σ allowed region around the best fit (blue dotted line) obtained with the analysis including only the operators, which are generated in the model at that order. The red line is the predicted value by the matching (see Table II). We assumed an integrated luminosity of 3 ab^{-1} .

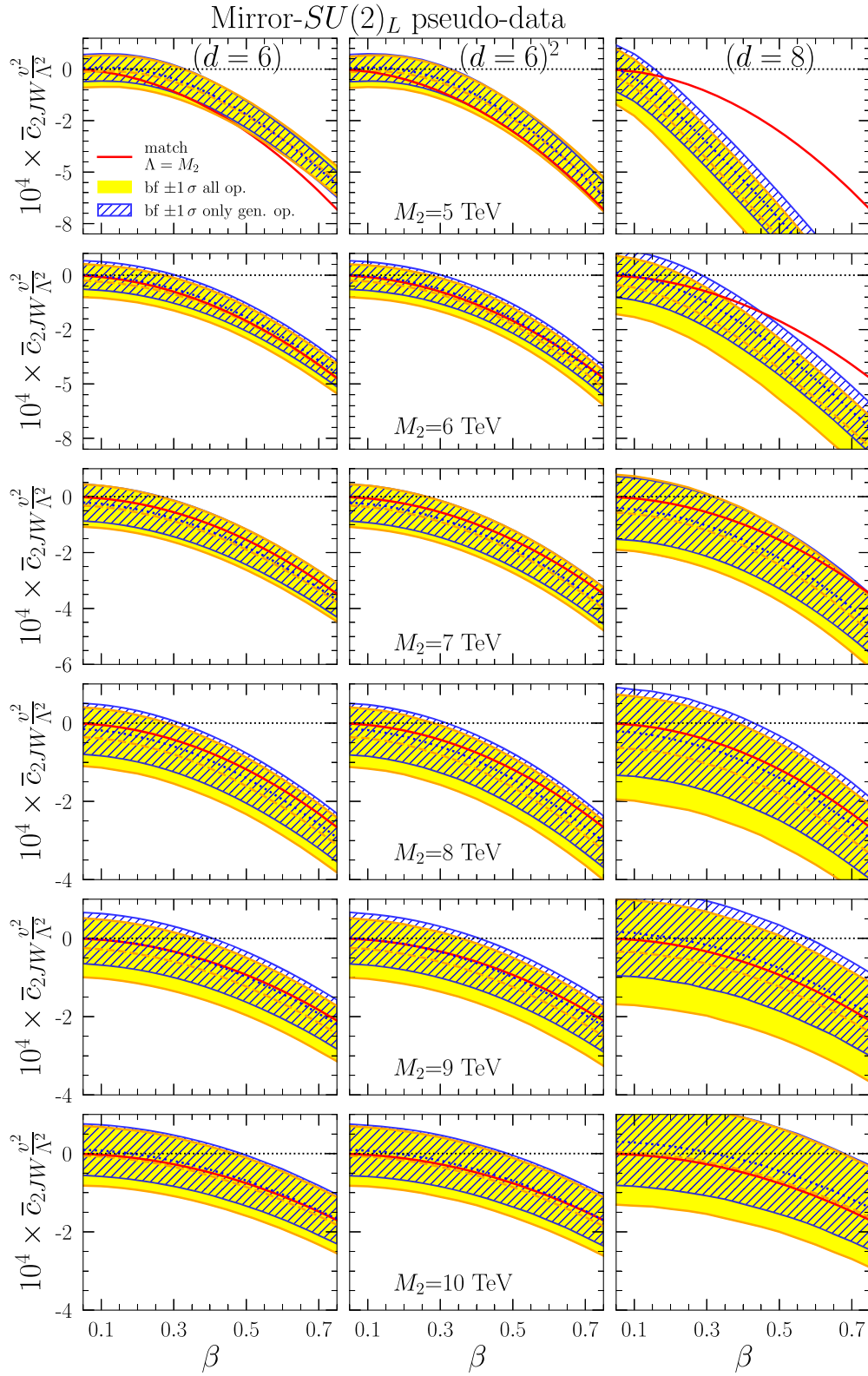


FIG. 8. Allowed one sigma regions for c_{2JW} as a function of β obtained with fit to pseudodata generated with the mirror- $SU(2)_L$ model for the masses as labeled in the figure. Conventions are the same as in Fig. 7.

for M_1 between 6 and 8 (9) TeV and β is in the range 1.3 to 1.9 (1 to 1.8).

Similarly, in the mirror- $SU(2)_L$ scenario, we find that the inclusion of the full dimension-eight contributions reduces the reach for NP due to the slight increase of the 68% allowed range of $\bar{c}_{2JW}$. In this case we find that 5σ discovery using the EFT framework is possible for M_2 ranging from 5 to 10 TeV and β between 0.5 and 1.5 when we consider the $1/\Lambda^4$ contributions. A 3σ evidence of NP is possible for these masses for β ranging from 0.3 to 1.2.

Finally, comparing the solid and dashed lines in both figures we find that in model parameters for which the EFT analysis is consistent, the sensitivity to NP is not improved by artificially selecting only the operators generated by the model.

We now turn to the question of whether the extracted values of the WC's capture the properties of the BSM in each case. To address this issue, we present in Figs. 7 and 8 the 1σ allowed band of $\bar{c}_{2JB}^{(i)}$ extracted from the fits to mirror- $U(1)$ pseudodata and $\bar{c}_{2JW}$ from the fits to mirror- $SU(2)$ pseudodata, respectively. These results are shown as function of the coupling β with the rows corresponding to a given mass and the columns to the different orders considered in the fit. In each panel the red solid line represents the predicted value of the WC from the matching in the model at that order. Furthermore, as for the f_{NP} figures, to test the robustness of the conclusions, we show the results for the complete fits to each order, as well as the result of the “biased” fit in which only the subset of known-to-be-generated operator coefficients are included.

These result for the fits to both models' pseudodata present similar behaviors. First of all, for heavier $M_{1,2}$ masses ($\gtrsim 7$ TeV) the allowed 1σ regions of $\bar{c}_{2JB}^{(i)}$ and $\bar{c}_{2JW}$ contain the *true* value of this WC (i.e., that predicted by the matching). Moreover, the results are stable with respect to the $1/\Lambda$ order indicating that the EFT is describing well the low-energy data. As we consider smaller masses the global fit 1σ bands do not contain the true values for large β , suggesting that the EFT is being employed beyond its applicability range due to the proximity and size of the resonance. The inclusion of higher orders does not mitigate this departure, possibly indicating that the effective extension of the window of applicability of the $1/\Lambda^4$ and higher order contributions is narrow in comparison with that of the $1/\Lambda^2$ terms.

Finally, comparing the blue hatched and solid yellow bands we see that the behavior does not change significantly in the restricted fits. Thus altogether, we find that our results are robust even when blindly fitting the data with minimal theoretical bias.

VI. SUMMARY

The expected accumulated luminosity at the HL-LHC will allow detailed analyses of the high energy tail of kinematic distributions. In the absence of new resonances, the EFT

framework for the analysis of those distributions is a powerful tool to search for signs of NP as well as to constrain possible extensions of the SM. In this work, we study the NC- and CC-DY processes to assess the HL-LHC potential to search for NP using USMEFT as the analysis tool. To this end, we simulated the expected signal of two universal SM extensions: the mirror- $U(1)_Y$ and mirror- $SU(2)_L$ and performed a fit of that pseudodata for the NC and CC DY distributions together with the EWPO with the minimal theoretical priors at different orders in the operator truncation.

We found that the analyses performed at orders dimension-six and dimension-six-squared lead to similar results due to the tight limits on the dimension-six Wilson coefficients. The inclusion of dimension-eight operators in the analyses lead to slightly less precise results due to the increased number of four-fermion operator coefficients.

More qualitatively, our unbiased fits of the pseudodata for both models show that the HL-LHC will be able to lead to a 5σ discovery (3σ evidence) for masses up to $\simeq 7$ (9) TeV provided the coupling is of similar strength to the corresponding the SM $U(1)_Y$ and $SU(2)_L$ couplings. We also analyzed how well a global fit with arbitrary WCs can capture the true value of the BSM theory. Except for the parameter regions where the new resonance could be directly observed, the dominant WC ($\bar{c}_{2JB}^{(i)}$ and $\bar{c}_{2JW}$) true model value is inside the 1σ allowed region for these WCs for a large range of masses and couplings.

We also find that that both the NP reach as well as the identification of the BSM model obtained in the most unbiased fit including all the WCs at a given order of the truncation are very similar to those obtained imposing the relations between WCs generated in each of the models. This further strengthens the robustness of our conclusion: fits with the minimal theoretical bias can be used not only to unveil the existence of new physics but also to accurately extract its properties.

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DATA AVAILABILITY

There are no publicly available research data or software supporting this manuscript. Requests for further information or data should be sent to the authors.

APPENDIX A: EFFECTIVE INTERACTIONS ENTERING OUR ANALYSES

There are six dimension-six and 11 dimension-eight operators entering the Drell-Yan production in USMEFT. In order to perform the Monte Carlo event simulations it is more convenient to work with the operators in the rotated basis including in addition the indirect effects induced by the finite renormalization of the SM parameters [39].

Here, we chose the input parameters to be $\{\hat{\alpha}_{\text{em}}, \hat{G}_F, \hat{M}_Z\}$ and take into account the following three relations to define the renormalized parameters

$$\hat{e} = \sqrt{4\pi\hat{\alpha}_{\text{em}}}, \quad \hat{v}^2 = \frac{1}{\sqrt{2}\hat{G}_F} \quad \text{and} \quad \hat{c}^2\hat{s}^2 = \frac{\pi\hat{\alpha}_{\text{em}}}{\sqrt{2}\hat{G}_F\hat{M}_Z^2},$$

where $\hat{s}$ ($\hat{c}$) is the sine (cosine) of the weak mixing angle $\hat{\theta}$. The contributions of fermionic operators to the muon decay width are written as

$$\left[2\langle H^\dagger H \rangle - \frac{1}{\sqrt{2}\hat{G}_F}\right]_{\text{fermionic}} \equiv \frac{\hat{v}^4}{\Lambda^2}\Delta_{4F} + \frac{\hat{v}^6}{\Lambda^4}\Delta_{4F}^{(8)}, \quad (\text{A1})$$

where Δ_{4F} ($\Delta_{4F}^{(8)}$) contains the dimension-six (-eight) contributions. After finite renormalization of the SM inputs and accounting for the direct contribution from the fermionic dimension-eight operators containing two fermion fields we can parametrize the Z coupling to fermion pairs $\bar{f}f$ as

$$\frac{\hat{e}}{\hat{s}\hat{c}} \left[\hat{g}_{L,R}^f \left(1 + \Delta\bar{g}' + \Delta g'^{\square} + \frac{p^2}{\hat{M}_Z^2} \Delta g' \right) + Q^f \left(\Delta\bar{g} + \Delta g^{\square} + \frac{p^2}{\hat{M}_Z^2} \Delta g' \right) \right], \quad (\text{A2})$$

and the W coupling to left-handed fermions as

$$\frac{1}{\sqrt{2}\hat{s}} \hat{e} \left(1 + \Delta\bar{g}_W + \Delta g_W^{\square} + \frac{p^2}{\hat{M}_W^2} \Delta g'_W \right), \quad (\text{A3})$$

where $\hat{g}_L^f = T_3^f - \hat{s}^2 Q^f$, $\hat{g}_R^f = -\hat{s}^2 Q^f$, T_3^f is the fermion's third component of isospin, and Q^f is its charge. In the expressions above p^2 is the square of the four-momentum of the corresponding gauge boson and the ΔX represent the BSM contributions.

At order $1/\Lambda^4$, the effective couplings $\Delta\bar{g}_{i(W)}$ and $\Delta g'_{i(W)}$ appearing in (A2) and (A3) depend upon the Wilson combinations defined in Eqs. (2.4)–(2.6) as

$$\Delta\bar{g}_1 = -\frac{1}{4}[2\bar{\Delta}_{4F} + \bar{c}_{\Phi,1}] \frac{\hat{v}^2}{\Lambda^2}, \quad (\text{A4})$$

$$\Delta\bar{g}_2 = -\frac{\hat{s}_2}{8\hat{c}_2} [\hat{s}_2(2\bar{\Delta}_{4F} + \bar{c}_{\Phi,1}) + 4\bar{c}_{BW}] \frac{\hat{v}^2}{\Lambda^2}, \quad (\text{A5})$$

$$\Delta\bar{g}_W = -\frac{1}{4\hat{c}_2} [2\hat{s}_2\bar{c}_{BW} + 2\hat{c}^2\bar{\Delta}_{4F} + \hat{c}^2\bar{c}_{\Phi,1}] \frac{\hat{v}^2}{\Lambda^2} - \frac{1}{2}\bar{c}_{W^2H^4}^{(3)} \frac{\hat{v}^4}{\Lambda^4}, \quad (\text{A6})$$

and

$$\Delta g'_1 = \frac{\hat{e}^4}{8\hat{c}^4\hat{s}^4} (\hat{s}^2 c_{\psi^4 D^2}^{(2)} + \hat{c}^2 c_{\psi^4 D^2}^{(3)}) \frac{\hat{v}^4}{\Lambda^4}, \quad (\text{A7})$$

$$\Delta g'_2 = \frac{\hat{e}^4}{8\hat{c}^2\hat{s}^2} (c_{\psi^4 D^2}^{(3)} - c_{\psi^4 D^2}^{(2)}) \frac{\hat{v}^4}{\Lambda^4}, \quad (\text{A8})$$

$$\Delta g'_W = \frac{\hat{e}^4}{8\hat{s}^4} c_{\psi^4 D^2}^{(3)} \frac{\hat{v}^4}{\Lambda^4}, \quad (\text{A9})$$

with $\hat{c}_n = \cos(n\hat{\theta})$ and $\hat{s}_n = \sin(n\hat{\theta})$. The coefficients $\Delta g_{i(W)}^{\square}$ contain the additional contributions, which are quadratic in the dimension-six Wilson coefficients. Their explicit expressions as a function of the effective coefficients can be found in Ref. [36].

In addition to the corrections to the couplings of Z and W bosons to fermion pairs, there are seven contact contributions to four-fermion amplitudes:

- (i) two at dimension six: c_{2JW} and c_{2JB} ,
- (ii) five at dimension eight: $c_{\psi^4 D^2}^{(2)}$, $c_{\psi^4 D^2}^{(3)}$, $c_{\psi^4 H^2}^{(4)}$, $c_{\psi^4 H^2}^{(5)}$, and $c_{\psi^4 H^2}^{(7)}$.

However, in DY, the contribution of c_{2JB} and $c_{\psi^4 H^2}^{(4)}$ (c_{2JW} and $c_{\psi^4 H^2}^{(5)}$) always enter together in the same combination. Furthermore, as before, the dimension-six Wilson coefficients c_{BB} and c_{WW} induce a shift on the gauge coupling constants which results in an associated shift on the four-fermion operator coefficients. Hence, we find that the four-fermion contact amplitudes depend upon the two combinations,

$$\bar{c}_{2JW} \equiv c_{2JW} \left(1 - 2 \frac{\hat{v}^2}{\Lambda^2} c_{WW} \right) + c_{\psi^4 H^2}^{(5)} \frac{\hat{v}^2}{2\Lambda^2} = -\frac{2\hat{s}^2}{\hat{e}^2} \bar{\Delta}_{4F}, \quad (\text{A10})$$

$$\bar{c}_{2JB} \equiv c_{2JB} \left(1 - 2 \frac{\hat{v}^2}{\Lambda^2} c_{BB} \right) + c_{\psi^4 H^2}^{(4)} \frac{\hat{v}^2}{2\Lambda^2}. \quad (\text{A11})$$

In addition $Q_{\psi^4 H^2}^{(7)}$ contributes to a different combination of the momentum independent four-fermion contact amplitudes. Moreover, $Q_{\psi^4 D^2}^{(2)}$ and $Q_{\psi^4 D^2}^{(3)}$ generate distinctive momentum-dependent four-fermion vertices.

APPENDIX B: FULL MATCHING

The full matching results at dimension eight include many operators not relevant to our analysis. We collect the

TABLE V. Full set of dimension-eight operators generated by the mirror- $U(1)_Y$ and mirror- $SU(2)_L$ models. We also include the dimension-six operator $Q_{\Phi,6}$ as discussed in the text.

Bosonic operators	Fermionic operators
$Q_{\Phi,6} = (H^\dagger H)^3$	$Q_{\psi^2 H^5}^{(1)} = (H^\dagger H)^2 (J_H H)$
$Q_{H^8} = (H^\dagger H)^4$	$Q_{\psi^2 H^4 D}^{(1)} = i J_B^\mu (H^\dagger \overleftrightarrow{D}_\mu H) (H^\dagger H)$
$Q_{H^6}^{(1)} = (H^\dagger H)^2 (D_\mu H)^\dagger (D^\mu H)$	$Q_{\psi^2 H^4 D}^{(2)} = i J_W^{I\mu} [(H^\dagger \overleftrightarrow{D}_\mu^I H) (H^\dagger H) + (H^\dagger \overleftrightarrow{D}_\mu H) (H^\dagger \tau^I H)]$
$Q_{H^6}^{(2)} = (H^\dagger H) (H^\dagger \tau^I H) (D_\mu H)^\dagger \tau^I (D^\mu H)$	$Q_{\psi^2 H^4 D}^{(4)} = \epsilon^{IJK} J_W^{I\mu} (H^\dagger \tau_\mu^J H) D_\mu (H^\dagger \tau^K H)$
$Q_{H^4}^{(1)} = (D_\mu H)^\dagger (D_\nu H) (D^\nu H)^\dagger (D^\mu H)$	$Q_{\psi^2 H^2 D^3}^{(1)} = i (D^\mu J_B^\nu + D^\nu J_B^\mu) (D_{(\mu} D_{\nu)} H^\dagger H - H^\dagger D_{(\mu} D_{\nu)} H)$
$Q_{H^4}^{(2)} = (D_\mu H)^\dagger (D_\nu H) (D^\mu H)^\dagger (D^\nu H)$	$Q_{\psi^2 H^2 D^3}^{(2)} = i (D^\mu J_W^{I\nu} + D^\nu J_W^{I\mu}) (D_{(\mu} D_{\nu)} H^\dagger \tau^I H - H^\dagger \tau^I D_{(\mu} D_{\nu)} H)$
$Q_{H^4}^{(3)} = (D_\mu H)^\dagger (D^\mu H) (D^\nu H)^\dagger (D_\nu H)$	$Q_{\psi^2 W H^2 D}^{(11)} = i \epsilon^{IJK} J_W^{I,\nu} (H^\dagger \overleftrightarrow{D}_\mu^J H) W_{\mu\nu}^K$
$Q_{B^2 H^4}^{(1)} = (H^\dagger H)^2 B_{\mu\nu} B^{\mu\nu}$	$Q_{\psi^4 H^2}^{(1)} = (H^\dagger H) J_H^j J_{H,j}^\dagger$
$Q_{W^2 H^4}^{(1)} = (H^\dagger H)^2 W_{\mu\nu}^I W^{I\mu\nu}$	$Q_{\psi^4 H^2}^{(2)} = (H_j J_H^j) (H_k J_H^k)$
$Q_{WBH^4}^{(1)} = (H^\dagger H) (H^\dagger \tau^I H) W_{\mu\nu}^I B^{\mu\nu}$	$Q_{\psi^4 H^2}^{(3)} = (H^{\dagger,k} J_{H,k}^\dagger) (J_H^j H_j)$
$Q_{BH^4 D^2}^{(1)} = (H^\dagger H) (D^\mu H)^\dagger (D^\nu H) B_{\mu\nu}$	$Q_{\psi^4 H^2}^{(4)} = (H^\dagger H) J_B^\mu J_{B\mu}$
$Q_{WH^4 D^2}^{(1)} = (H^\dagger H) (D^\mu H)^\dagger \tau^I (D^\nu H) W_{\mu\nu}^I$	$Q_{\psi^4 H^2}^{(5)} = (H^\dagger H) J_W^{I\mu} J_{W\mu}^I$
	$Q_{\psi^4 D^2}^{(2)} = D^\alpha J_B^\mu D_\alpha J_{B\mu}$
	$Q_{\psi^4 D^2}^{(3)} = D^\alpha J_W^{I\mu} D_\alpha J_{W\mu}^I$
	$Q_{\psi^4 W}^{(1)} = \epsilon^{IJK} W_{\mu\nu}^I J_W^{J\mu} J_W^{K\nu}$

full dimension-eight Lagrangians here. In what follows we have *suppressed the mass scaling* of the operators to improve the presentation. As an example the following term in the Lagrangian for the mirror- $U(1)_Y$ model,

$$\mathcal{L}_8 \subset -\frac{1}{2} \beta^2 Q_{\psi^4 D^2}^{(2)}, \quad (\text{B1})$$

should be understood to scale as,

$$-\frac{1}{2} \beta^2 Q_{\psi^4 D^2}^{(2)} \rightarrow -\frac{1}{2 M_1^2} \beta^2 Q_{\psi^4 D^2}^{(2)}. \quad (\text{B2})$$

The full dimension-eight Lagrangian for the mirror- $U(1)_Y$ model is given by,

$$\begin{aligned} \mathcal{L}_8 = & -\frac{1}{8} v^2 g^2 g'^2 \beta^2 (1 + \beta^2) \lambda_h Q_{\Phi,6} + \frac{1}{4} g^2 g'^2 \beta^2 (1 + \beta^2) \lambda_h Q_{H^8} \\ & -\frac{1}{2} g^2 g'^2 \beta^2 (1 + \beta^2) Q_{H^6}^{(1)} - \frac{1}{8} g^2 g'^2 \beta^2 (1 + \beta^2) Q_{H^6}^{(2)} - \frac{1}{2} g^2 \beta^2 (1 + \beta^2) Q_{H^4}^{(1)} + \frac{1}{2} g^2 \beta^2 (1 + \beta^2) Q_{H^4}^{(2)} \\ & + \frac{1}{16} g^2 g'^2 \beta^2 (1 + \beta^2) Q_{W^2 H^4}^{(1)} - \frac{1}{16} g^4 \beta^2 (1 + \beta^2) Q_{B^2 H^4}^{(1)} + \frac{i}{2} g g'^2 \beta^2 (1 + \beta^2) Q_{W H^4 D^2}^{(1)} - \frac{i}{2} g'^3 \beta^2 (1 + \beta^2) Q_{B H^4 D^2}^{(1)} \\ & + \frac{1}{16} g^2 g'^2 \beta^2 (1 + \beta^2) (Q_{\psi^2 H^5}^{(1)} + Q_{\psi^2 H^5}^{\dagger(1)}) - \frac{1}{8} g g'^2 \beta^2 (1 + \beta^2) Q_{\psi^2 H^4 D}^{(2)} + \frac{1}{4} g' \beta^2 Q_{\psi^2 H^2 D^3}^{(1)} \\ & - \frac{3}{8} g^2 \beta^2 Q_{\psi^4 H^2}^{(1)} + \frac{3}{8} g^2 \beta^2 Q_{\psi^4 H^2}^{(2)} + \frac{3}{8} g^2 \beta^2 Q_{\psi^4 H^2}^{\dagger(2)} - \frac{3}{4} g^2 \beta^2 Q_{\psi^4 H^2}^{(3)} + \frac{1}{4} g^2 \beta^4 Q_{\psi^4 H^2}^{(4)} - \frac{1}{2} \beta^2 Q_{\psi^4 D^2}^{(2)}. \end{aligned} \quad (\text{B3})$$

In the first line we also include a dimension-six operator, $Q_{\Phi,6} = (H^\dagger H)^3$, which results from rewriting the purely bosonic dimension-eight operators in terms of fermionic operators via the EOM following [58], in particular see the top of Appendix A. The full set of operators appearing in this, and the dimension-eight Lagrangian for the mirror- $SU(2)_L$ model can be found in Table V.

For the mirror- $SU(2)_L$ model we find,

$$\begin{aligned}
\mathcal{L}_8 = & -\frac{1}{4}g^2g'^2\beta^2(1+\beta^2)v^2\lambda_h Q_{\Phi,6} + \frac{1}{16}g^2\beta^2(-3g^2\beta^2 + 8g'^2(1+\beta^2))\lambda_h Q_{H^8} \\
& + \frac{1}{16}(9g^4\beta^4 - 12g^2g'^2\beta^2(1+\beta^2))Q_{H^6}^{(1)} - \frac{1}{2}g^2g'^2\beta^2(1+\beta^2)Q_{H^6}^{(2)} \\
& + \frac{1}{2}g^2\beta^2(1+\beta^2)Q_{H^4}^{(1)} + \frac{1}{2}g^2\beta^2(1+\beta^2)Q_{H^4}^{(2)} - g^2\beta^2(1+\beta^2)Q_{H^4}^{(3)} \\
& + \frac{3}{16}g^2g'^2\beta^2(1+\beta^2)Q_{B^2H^4}^{(1)} - \frac{1}{16}g^4\beta^2(1+3\beta^2)Q_{W^2H^4}^{(1)} + \frac{1}{8}g^3g'\beta^2Q_{WBH^4}^{(1)} \\
& + \frac{3i}{2}g^2g'\beta^2(1+\beta^2)Q_{BH^4D^2}^{(1)} - \frac{i}{2}g^3\beta^2(1+3\beta^2)Q_{WH^4D^2}^{(1)} \\
& + \frac{1}{8}g^2g'^2\beta^2(1+\beta^2)(Q_{\psi^2H^5}^{(1)} + Q_{\psi^2H^5}^{\dagger(1)}) - \frac{1}{2}g^2g'\beta^2(1+\beta^2)Q_{\psi^2H^4D}^{(1)} + \frac{1}{4}g^3\beta^4Q_{\psi^2H^4D}^{(2)} \\
& + \frac{1}{4}g^3\beta^4Q_{\psi^2H^4D}^{(4)} + \frac{1}{4}g\beta^2Q_{\psi^2H^2D^3}^{(2)} + g^2\beta^2Q_{\psi^2WH^2D}^{(11)} \\
& - \frac{9}{8}g^2\beta^2Q_{\psi^4H^2}^{(1)} + \frac{3}{8}g^2\beta^2(Q_{\psi^4H^2}^{(2)} + Q_{\psi^4H^2}^{\dagger(2)}) + \frac{3}{4}g^2\beta^2Q_{\psi^4H^2}^{(3)} + \frac{1}{4}g^2\beta^4Q_{\psi^4H^2}^{(5)} - \frac{1}{2}\beta^2Q_{\psi^4D^2}^{(3)} - \frac{1}{2}g\beta^2Q_{\psi^4W}^{(1)}. \quad (B4)
\end{aligned}$$

As in the case of the mirror- $U(1)_Y$ model the dimension-six operator $Q_{\Phi,6}$ is generated in changing from the purely bosonic to fermionic basis. The full set of operators appearing in this Lagrangian can be found in Table V.

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- [1] S. D. Drell and T.-M. Yan, Massive lepton pair production in Hadron-Hadron collisions at high-energies, *Phys. Rev. Lett.* **25**, 316 (1970); *Phys. Rev. Lett.* **25**, 902(E) (1970).
 - [2] G. Aad *et al.* (ATLAS Collaboration), Measurement of the double-differential high-mass Drell-Yan cross section in pp collisions at $\sqrt{s} = 8$ TeV with the ATLAS detector, *J. High Energy Phys.* **08** (2016) 009.
 - [3] V. Khachatryan *et al.* (CMS Collaboration), Measurements of differential and double-differential Drell-Yan cross sections in proton-proton collisions at 8 TeV, *Eur. Phys. J. C* **75**, 147 (2015).
 - [4] G. Walia (CMS Collaboration), Single and double differential Drell-Yan cross section measurements using the CMS Detector, *Springer Proc. Phys.* **203**, 517 (2018).
 - [5] A. Hayrapetyan *et al.* (CMS Collaboration), Measurement of the Drell-Yan forward-backward asymmetry and of the effective leptonic weak mixing angle in proton-proton collisions at $\sqrt{s} = 13$ TeV, *Phys. Lett. B* **866**, 139526 (2025).
 - [6] G. Aad *et al.* (ATLAS Collaboration), Measurement of double-differential charged-current Drell-Yan cross-sections at high transverse masses in pp collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, *J. High Energy Phys.* **07** (2025) 026.
 - [7] G. Aad *et al.* (ATLAS Collaboration), Search for high-mass dilepton resonances using 139 fb $^{-1}$ of pp collision data collected at $\sqrt{s} = 13$ TeV with the ATLAS detector, *Phys. Lett. B* **796**, 68 (2019).
 - [8] A. M. Sirunyan *et al.* (CMS Collaboration), Search for resonant and nonresonant new phenomena in high-mass dilepton final states at $\sqrt{s} = 13$ TeV, *J. High Energy Phys.* **07** (2021) 208.
 - [9] G. Aad *et al.* (ATLAS Collaboration), Search for a heavy charged boson in events with a charged lepton and missing transverse momentum from pp collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, *Phys. Rev. D* **100**, 052013 (2019).
 - [10] A. Tumasyan *et al.* (CMS Collaboration), Search for new physics in the lepton plus missing transverse momentum final state in proton-proton collisions at $\sqrt{s} = 13$ TeV, *J. High Energy Phys.* **07** (2022) 067.
 - [11] G. Aad *et al.* (ATLAS Collaboration), Observation of a new particle in the search for the Standard Model Higgs boson with the ATLAS detector at the LHC, *Phys. Lett. B* **716**, 1 (2012).
 - [12] M. Aaboud *et al.* (ATLAS Collaboration), Search for WW/WZ resonance production in $\ell\nu qq$ final states in pp collisions at $\sqrt{s} = 13$ TeV with the ATLAS detector, *J. High Energy Phys.* **03** (2018) 042.

- [13] S. Dawson, S. Homiller, and S. D. Lane, Putting standard model EFT fits to work, *Phys. Rev. D* **102**, 055012 (2020).
- [14] S. Dawson, S. Homiller, and M. Sullivan, Impact of dimension-eight SMEFT contributions: A case study, *Phys. Rev. D* **104**, 115013 (2021).
- [15] J. Ellis, M. Madigan, K. Mimasu, V. Sanz, and T. You, Top, Higgs, Diboson and electroweak fit to the Standard Model effective field theory, *J. High Energy Phys.* **04** (2021) 279.
- [16] S. Dawson, M. Forslund, and M. Schnubel, SMEFT matching to Z' models at dimension eight, *Phys. Rev. D* **110**, 015002 (2024).
- [17] N. Adhikary, T. Biswas, J. Chakraborty, C. Englert, and M. Spannowsky, Electroweak scalar effects beyond dimension-six in SMEFT, *Phys. Rev. D* **113**, 036003 (2026).
- [18] S. Dawson, D. Fontes, S. Homiller, and M. Sullivan, Role of dimension-eight operators in an EFT for the 2HDM, *Phys. Rev. D* **106**, 055012 (2022).
- [19] J. Ellis, K. Mimasu, and F. Zampedri, Dimension-8 SMEFT analysis of minimal scalar field extensions of the Standard Model, *J. High Energy Phys.* **10** (2023) 051.
- [20] T. Corbett, Top-down and bottom-up: Studying the SMEFT beyond leading order in $1/\Lambda^2$, *SciPost Phys. Core* **7**, 053 (2024).
- [21] I. Brivio, S. Bruggisser, E. Geoffray, W. Killian, M. Krämer, M. Luchmann, T. Plehn, and B. Summ, From models to SMEFT and back?, *SciPost Phys.* **12**, 036 (2022).
- [22] J. ter Hoeve, G. Magni, J. Rojo, A. N. Rossia, and E. Vryonidou, The automation of SMEFT-assisted constraints on UV-complete models, *J. High Energy Phys.* **01** (2024) 179.
- [23] A. Greljo and A. Palavrić, Leading directions in the SMEFT, *J. High Energy Phys.* **09** (2023) 009.
- [24] A. Greljo, A. Palavrić, and A. Smolkovič, Leading directions in the SMEFT: Renormalization effects, *Phys. Rev. D* **109**, 075033 (2024).
- [25] M. Hirsch, L. Mantani, and V. Sanz, Data-driven discovery strategy for Standard Model effective field theory searches, *Phys. Rev. Lett.* **135**, 241801 (2025).
- [26] D. Liu, A. Pomarol, R. Rattazzi, and F. Riva, Patterns of strong coupling for LHC searches, *J. High Energy Phys.* **11** (2016) 141.
- [27] C. Degrande, A basis of dimension-eight operators for anomalous neutral triple gauge boson interactions, *J. High Energy Phys.* **02** (2013) 101.
- [28] S. Alioli, R. Boughezal, E. Mereghetti, and F. Petriello, Novel angular dependence in Drell-Yan lepton production via dimension-8 operators, *Phys. Lett. B* **809**, 135703 (2020).
- [29] S. Alte, M. König, and W. Shepherd, Consistent searches for SMEFT effects in non-resonant Dilepton events, *J. High Energy Phys.* **07** (2018) 144.
- [30] R. Boughezal, E. Mereghetti, and F. Petriello, Dilepton production in the SMEFT at $\mathcal{O}(1/\Lambda^4)$, *Phys. Rev. D* **104**, 095022 (2021).
- [31] R. Boughezal, Y. Huang, and F. Petriello, Exploring the SMEFT at dimension eight with Drell-Yan transverse momentum measurements, *Phys. Rev. D* **106**, 036020 (2022).
- [32] T. Kim and A. Martin, Monolepton production in SMEFT to $\mathcal{O}(1/\Lambda^4)$ and beyond, *J. High Energy Phys.* **09** (2022) 124.
- [33] L. Allwicher, D. A. Faroughy, F. Jaffredo, O. Sumensari, and F. Wilsch, Drell-Yan tails beyond the Standard Model, *J. High Energy Phys.* **03** (2022) 064.
- [34] L. Allwicher, D. A. Faroughy, F. Jaffredo, O. Sumensari, and F. Wilsch, HighPT: A tool for high- p_T Drell-Yan tails beyond the standard model, *Comput. Phys. Commun.* **289**, 108749 (2023).
- [35] L. Allwicher, D. A. Faroughy, M. Martinez, O. Sumensari, and F. Wilsch, On the EFT validity for Drell-Yan tails at the LHC, *Eur. Phys. J. C* **85**, 463 (2025).
- [36] T. Corbett, J. Desai, O. J. P. Eboli, M. C. Gonzalez-Garcia, M. Martinez, and P. Reimitz, Drell-Yan production in universal theories beyond dimension-six SMEFT, *Phys. Rev. D* **112**, 013009 (2025).
- [37] D. Liu and L.-T. Wang, Precision measurement with Diboson at the LHC, [arXiv:1804.08688](https://arxiv.org/abs/1804.08688).
- [38] C. Degrande and H.-L. Li, Impact of dimension-8 SMEFT operators on Diboson productions, *J. High Energy Phys.* **06** (2023) 149.
- [39] T. Corbett, J. Desai, O. J. P. Éboli, M. C. Gonzalez-Garcia, M. Martinez, and P. Reimitz, Impact of dimension-eight SMEFT operators in the electroweak precision observables and triple gauge couplings analysis in universal SMEFT, *Phys. Rev. D* **107**, 115013 (2023).
- [40] A. Martin, A case study of SMEFT $\mathcal{O}(1/\Lambda^4)$ effects in diboson processes: $pp \rightarrow W^\pm(\ell^\pm \nu)\gamma$, *J. High Energy Phys.* **05** (2023) 223.
- [41] D. Gillies, A. Banfi, A. Martin, and M. A. Lim, Dimension-8 operators in WW production via gluon fusion, *J. High Energy Phys.* **06** (2024) 111.
- [42] J. M. Goldberg, H. Liu, and Y. Shadmi, Dimension-8 SMEFT contact-terms for vector-pair production via on-shell Higgsing, *J. High Energy Phys.* **12** (2024) 057.
- [43] C. Hays, A. Martin, V. Sanz, and J. Setford, On the impact of dimension-eight SMEFT operators on Higgs measurements, *J. High Energy Phys.* **02** (2018) 123.
- [44] B. Assi and A. Martin, Energy-enhanced dimension eight SMEFT effects in VBF Higgs production, *J. High Energy Phys.* **02** (2024) 029.
- [45] T. Corbett, A. Martin, and M. Trott, Consistent higher order $\sigma(\mathcal{G}\mathcal{G} \rightarrow h)$, $\Gamma(h \rightarrow \mathcal{G}\mathcal{G})$ and $\Gamma(h \rightarrow \gamma\gamma)$ in geoSMEFT, *J. High Energy Phys.* **12** (2021) 147.
- [46] A. Martin and M. Trott, ggh variations, *Phys. Rev. D* **105**, 076004 (2022).
- [47] T. Corbett, The one-loop tadpole in the geoSMEFT, *SciPost Phys.* **11**, 097 (2021).
- [48] C. Hays, A. Helset, A. Martin, and M. Trott, Exact SMEFT formulation and expansion to $\mathcal{O}(v^4/\Lambda^4)$, *J. High Energy Phys.* **11** (2020) 087.
- [49] M. Flores-Hernandez and A. Martin, Higgs decays to four leptons to $\mathcal{O}(1/\Lambda^4)$ in SMEFT, [arXiv:2602.12326](https://arxiv.org/abs/2602.12326).
- [50] S. Bera, D. Chakraborty, S. Chattopadhyay, and R. S. Gupta, Differential observables for the Higgs-strahlung process to all orders in EFT, [arXiv:2601.08821](https://arxiv.org/abs/2601.08821).
- [51] T. Corbett, A. Helset, A. Martin, and M. Trott, EWPD in the SMEFT to dimension eight, *J. High Energy Phys.* **06** (2021) 076.
- [52] H. D. Politzer, Power corrections at short distances, *Nucl. Phys. B* **172**, 349 (1980).

- [53] H. Georgi, On-shell effective field theory, *Nucl. Phys.* **B361**, 339 (1991).
- [54] C. Arzt, Reduced effective Lagrangians, *Phys. Lett. B* **342**, 189 (1995).
- [55] H. Simma, Equations of motion for effective Lagrangians and penguins in rare B decays, *Z. Phys. C* **61**, 67 (1994).
- [56] I. Brivio, O. J. P. Éboli, M. B. Gavela, M. C. González-Garcia, L. Merlo, and S. Rigolin, Higgs ultraviolet softening, *J. High Energy Phys.* **12** (2014) 004.
- [57] J. D. Wells and Z. Zhang, Effective theories of universal theories, *J. High Energy Phys.* **01** (2015) 123.
- [58] T. Corbett, J. Desai, O. J. P. Eboli, and M. C. Gonzalez-Garcia, Dimension-eight operator basis for universal standard model effective field theory, *Phys. Rev. D* **110**, 033003 (2024).
- [59] M. Madigan, Can PDFs absorb new physics?, *Proc. Sci., EPS-HEP2023* (2024) 479.
- [60] E. Hammou, Z. Kassabov, M. Madigan, M. L. Mangano, L. Mantani, J. Moore, M. M. Alvarado, and M. Ubiali, Hide and seek: How PDFs can conceal new physics, *J. High Energy Phys.* **11** (2023) 090.
- [61] J. Fuentes-Martín, M. König, J. Pagès, A. E. Thomsen, and F. Wilsch, A proof of concept for matchete: an automated tool for matching effective theories, *Eur. Phys. J. C* **83**, 662 (2023).
- [62] B. Henning, X. Lu, and H. Murayama, How to use the Standard Model effective field theory, *J. High Energy Phys.* **01** (2014) 023.
- [63] J. C. Criado and M. Pérez-Victoria, Field redefinitions in effective theories at higher orders, *J. High Energy Phys.* **03** (2018) 038.
- [64] C. W. Murphy, Dimension-8 operators in the Standard Model effective field theory, *J. High Energy Phys.* **10** (2020) 174.
- [65] M. Farina, G. Panico, D. Pappadopulo, J. T. Ruderman, R. Torre, and A. Wulzer, Energy helps accuracy: Electroweak precision tests at hadron colliders, *Phys. Lett. B* **772**, 210 (2017).
- [66] Prospects for searches for heavy Z' and W' bosons in fermionic final states with the ATLAS experiment at the HL-LHC (2018).
- [67] R. Frederix, S. Frixione, V. Hirschi, D. Pagani, H. S. Shao, and M. Zaro, The automation of next-to-leading order electroweak calculations, *J. High Energy Phys.* **07** (2018) 185.
- [68] N. D. Christensen and C. Duhr, FeynRules—Feynman rules made easy, *Comput. Phys. Commun.* **180**, 1614 (2009).
- [69] A. Alloul, N. D. Christensen, C. Degrande, C. Duhr, and B. Fuks, FeynRules 2.0—A complete toolbox for tree-level phenomenology, *Comput. Phys. Commun.* **185**, 2250 (2014).
- [70] T. Sjostrand, S. Mrenna, and P. Z. Skands, A brief introduction to PYTHIA 8.1, *Comput. Phys. Commun.* **178**, 852 (2008).
- [71] J. de Favereau, C. Delaere, P. Demin, A. Giammanco, V. Lemaitre, A. Mertens, and M. Selvaggi (DELPHES 3 Collaboration), DELPHES 3, A modular framework for fast simulation of a generic collider experiment, *J. High Energy Phys.* **02** (2013) 057.
- [72] M. Cacciari, G. P. Salam, and G. Soyez, FastJet user manual, *Eur. Phys. J. C* **72**, 1896 (2012).
- [73] C. Bierlich *et al.*, Robust independent validation of experiment and theory: Rivet version 3, *SciPost Phys.* **8**, 026 (2020).
- [74] A. Buckley, D. Kar, and K. Nordström, Fast simulation of detector effects in Rivet, *SciPost Phys.* **8**, 025 (2020).
- [75] Prospect for a measurement of the Weak Mixing Angle in $pp \rightarrow Z/\gamma^* \rightarrow e^+e^-$ events with the ATLAS detector at the High Luminosity Large Hadron Collider, Report No. ATL-PHYS-PUB-2018-037 (2018), <https://inspirehep.net/files/de6bdcea841dc9e434119338e1dc1cda>.
- [76] S. Schael *et al.* (L3 Collaborations), Precision electroweak measurements on the Z resonance, *Phys. Rep.* **427**, 257 (2006).
- [77] S. Navas *et al.* (Particle Data Group), Review of particle physics, *Phys. Rev. D* **110**, 030001 (2024).
- [78] J. de Blas, M. Pierini, L. Reina, and L. Silvestrini, Impact of the recent measurements of the top-quark and W-boson masses on electroweak precision Fits, *Phys. Rev. Lett.* **129**, 271801 (2022).