PHYSICAL REVIEW D **87,** 055019 (2013)

Flavor with a light dynamical "Higgs particle"

R. Alonso,[1,*] M. B. Gavela,[1,2,†] L. Merlo,[1,2,3,‡] S. Rigolin,[4,§] and J. Yepes[1,∥]

[1]*Departamento de Física Teórica and Instituto de Física Teórica, IFT-UAM/CSIC, Universidad Autónoma de Madrid, Cantoblanco 28049, Madrid, Spain*

[2]*Department of Physics, Theory Division, CERN, CH-1211 Geneva 23, Switzerland*

[3]*TUM Institute for Advanced Study, Technische Universität München, Lichtenbergstrasse 2a, D-85748 Garching, Germany*

[4]*Dipartimento di Fisica "G. Galilei," Università di Padova and INFN, Sezione di Padova, Via Marzolo 8, I-35131 Padua, Italy*

(Received 21 December 2012; published 20 March 2013)

The Higgs-fermion couplings are sensitive probes of possible new physics behind a stable light Higgs particle. It is then essential to identify the flavor pattern of those interactions. We consider the case in which a strong dynamics lies behind a light Higgs and explore the implications within the minimal flavor violation ansatz. The dominant effects on flavor-changing Higgs-fermion couplings stem in this context from operators with mass dimension ≤ 5, and we analyze all relevant chiral operators up to that order, including loop corrections induced by four-dimensional ones. Bounds on the operator coefficients are derived from a plethora of low-energy flavor transitions, providing a guideline on which flavor-changing Higgs interactions may be open to experimental scrutiny. In particular, the coefficient of a genuinely *CP*-odd operator is only softly constrained and therefore its impact is potentially interesting.

DOI: 10.1103/PhysRevD.87.055019 PACS numbers: 12.60.Fr, 11.30.Hv, 12.38.Qk

I. INTRODUCTION

A new resonance at the electroweak (EW) scale has been established at LHC. Both ATLAS and CMS Collaborations have recently presented [1,2] the discovery of an excess of events above the expected Standard Model (SM) background with a local significance of 5σ consistent with the hypothesis of the SM scalar boson [3–5] (so-called "Higgs boson" for short) with mass around 125 GeV.

This resonance is, at the moment, compatible with the SM Higgs interpretation, even if the rate in the diphoton channel, slightly above SM expectations, leaves still open the possibility of nonstandard effects, and furthermore a $\sim 2\sigma$ tension persists between the predictions and measurement of the rate R_b^0 and the forward-backward asymmetry $A_{FB}^{0,b}$, in b-quark production from $e^+ - e^-$ collisions [6,7].

There are essentially two main frameworks that have been proposed over the last decades in order to explain the EW symmetry breaking sector. The first possibility is that the Higgs is a fundamental particle, transforming linearly (as a doublet in the standard minimal picture) under the gauge symmetry group $SU(2)_L \times U(1)_Y$. This line of thought suggests, due to the appearance of the hierarchy problem, to invoke new physics (NP) around the TeV scale in order to definitively stabilize the Higgs (and the EW) mass scale. The minimal supersymmetric standard model and its variations are the best explored options of that kind, and a plethora of supersymmetry partners should populate the scale unveiled by LHC experiments, unless awkward fine-tuning effects take place.

An interesting alternative is that the Higgs dynamics is nonperturbative and associated to a strong interacting force with scale Λ_s, and the gauge symmetry in the scalar sector is nonlinearly realized. In the original "technicolor" formulation [8–10], no physical Higgs particle appears in the low-energy spectrum and only the three would-be-Goldstone bosons responsible for the weak gauge boson masses are retained. The characteristic scale f associated to the Goldstone bosons was identified with the electroweak scale $f = v \equiv 246$ GeV, defined from the W mass $M_W = gv/2$, and respecting $f \geq \Lambda_s/4\pi$ [11]. The smoking gun signature of this technicolor ansatz is the appearance of several vector and fermion resonances at the TeV scale. The discovery of a light Higgs candidate has recently focused the attention on an interesting variant: to consider still a strong dynamics behind the electroweak scalar sector but resulting—in addition—in a composite (instead of elementary) and light Higgs particle. In this scenario, proposed long ago [12–17], the Higgs itself would be one of the Goldstone bosons associated with the strong dynamics at the scale Λ_s, while its mass would result from some explicit breaking of the underlying strong dynamics. It was suggested that this breaking may be caused by the weak gauge interactions or alternatively by nonrenormalizable couplings. These ideas have been revived in recent years and are opportune given the recent experimental data (see for example Ref. [18] for a recent review on the subject). In this class of scenarios, f may lie around the TeV regime, while v is linked to the

*rodrigo.alonso@uam.es
†belen.gavela@uam.es
‡luca.merlo@uam.es
§stefano.rigolin@pd.infn.it
∥ju.yepes@uam.es

electroweak symmetry breaking process and is not identified with f, $v \leq f$. The degree of nonlinearity is then quantified by a new parameter,

$$\xi \equiv \frac{v^2}{f^2}, \tag{1.1}$$

and, for instance, $f \sim v$ characterizes the extreme nonlinear constructions, while $f \gg v$ is typical of scenarios which mimic the linear regime. As a result, for non-negligible ξ there may be corrections to the size of the SM couplings observable at low energies due to NP contributions.

The question we address in this paper is the flavor structure of the NP operator coefficients, when a strong dynamics is assumed at the scale Λ_s and in the presence of a light Higgs particle. In particular, dangerous NP contributions to flavor-changing observables could arise. Indeed, the core of the flavor problem in NP theories consists in explaining the high level of suppression that must be encoded in most of the theories beyond the SM in order to pass flavor-changing neutral current (FCNC) observability tests. Minimal flavor violation (MFV) [19–21] emerged in the last years as one of the most promising working frameworks and it will be used in this work.

Following the MFV ansatz, flavor in the SM *and beyond* is described at low energies uniquely in terms of the known fermion mass hierarchies and mixings. An outcome of the MFV ansatz is that the energy scale of the NP may be as low as few TeV in several distinct contexts [22–25], while in general it should be larger than hundreds of TeV [26]. MFV has been codified as a general framework built upon the flavor symmetry of the kinetic terms [27–34]. For quarks, the flavor group

$$G_f = SU(3)_{Q_L} \times SU(3)_{U_R} \times SU(3)_{D_R} \tag{1.2}$$

defines the non-Abelian transformation properties of the $SU(2)_L$ doublet Q_L and singlets U_R and D_R,

$$Q_L \sim (3,1,1), \quad U_R \sim (1,3,1), \quad D_R \sim (1,1,3). \tag{1.3}$$

To introduce the Yukawa Lagrangian without explicitly breaking G_f, the Yukawa matrices for up (Y_U) and down (Y_D) quarks can be promoted to be spurion fields transforming under the flavor symmetry,

$$Y_U \sim (3, \bar{3}, 1), \qquad Y_D \sim (3, 1, \bar{3}). \tag{1.4}$$

The quark masses and mixings are correctly reproduced once these spurion fields get background values as

$$Y_U = V^\dagger \mathbf{y}_U, \qquad Y_D = \mathbf{y}_D, \tag{1.5}$$

where $\mathbf{y}_{U,D}$ are diagonal matrices whose elements are the Yukawa eigenvalues, and V a unitary matrix that in good approximation coincides with the Cabibbo-Kobayashi-Maskawa (CKM) matrix. These background values break the flavor group G_f, providing contributions to FCNC observables suppressed by specific combinations of quark mass hierarchies and mixing angles. In Ref. [21], the complete basis of gauge-invariant six-dimensional FCNC operators has been constructed for the case of a linearly realized SM Higgs sector, in terms of the SM fields and the Y_U and Y_D spurions. Operators of dimension $d > 6$ are usually neglected due to the additional suppression in terms of the cutoff scale.

The MFV ansatz in the presence on a strong interacting dynamics has been introduced in Ref. [35], where the list of relevant $d = 4$ flavor-changing operators was identified, in the limit in which the Higgs degree of freedom is integrated out. In the nonlinear regime a chiral expansion is pertinent, and this results in a different set of operators at leading order than in the case of the linear regime, as the leading operators in the linear and nonlinear expansion do not match one to one (see for instance the discussion in Ref. [36]). The promotion of the Yukawa matrices to spurions follows the same lines as in the linear regime, though. Indeed, when the SM quarks $\Psi_{L,R}$ couple bilinearly to the strong sector[1]

$$\bar{\Psi}_L Y_\Psi \Psi_R \Theta_s, \tag{1.6}$$

with Θ_s a flavor blind operator in the strong sector, then all flavor information is encoded in Y_Ψ, that, in order to preserve the flavor group G_f, must transform as in Eq. (1.4). Once the spurions have been defined as the only sources of flavor violation (in the SM and beyond), it is possible to build the tower of FCNC operators, invariant under both the gauge and the flavor symmetries. It is customary to define

$$\lambda_F \equiv Y_U Y_U^\dagger + Y_D Y_D^\dagger = V^\dagger \mathbf{y}_U^2 V + \mathbf{y}_D^2, \tag{1.7}$$

which transforms as a (8, 1, 1) under G_f. The only relevant nondiagonal entries are all proportional to the top Yukawa coupling, $(\lambda_F)_{ij} \approx y_t^2 V_{ti}^* V_{tj}$, for $i \neq j$.

Within the spirit of MFV, the flavor structure of all Yukawa terms will be dictated only by its fermion composition; in consequence, the resulting fermion-h couplings get diagonalized together with the fermion mass matrix diagonalization. In other words, flavor-changing couplings require operators of (at least) dimension 5. This property will also apply to the nonlinear analysis below.

In this work we construct the tower of $d \leq 5$ h-fermion flavor-changing operators for a generic strong interacting light Higgs scenario. Which operator basis is chosen in an effective Lagrangian approach is an issue relevant to get the best bounds from a given set of observables, and a

[1] A more complicated case in which the link among the proto-Yukawa interactions and the spurions $Y_{U,D}$ is less direct happens in the context of the partial compositeness [37]. In this case, quarks couple to the strong sector linearly and therefore two Yukawa couplings, Y_{Ψ_L} and Y_{Ψ_R}, for each flavor sector appear. Spurions are then identified with only one of these proto-Yukawa couplings for each flavor sector, with the other assumed flavor diagonal [38]. We will not consider here this possibility.

convenient basis will be used when analyzing flavor, distinct from that applied in Refs. [36,39–41] to analyze the Higgs-gauge sector. A consistent approach requires to revisit as well the $d=4$ flavor-changing operators presented in Ref. [35], by introducing the possibility of a light scalar Higgs, and to consider in addition their main loop-induced effects. In the theoretical discussion we will reconsider the interesting exercise performed in Refs. [40,41] to reach the nonlinear regime from the linear one [39] in the presence of a light Higgs. We will also perform the phenomenological analysis of the strength of the NP fermionic couplings, focusing on the—often stringent—bounds on the operator coefficients that follow from present low-energy measurements on the Higgs-less component of the couplings. This will provide a guideline on which type of flavored Higgs couplings may be at reach at the LHC.

The structure of the paper is the following. Section II describes the framework and it is mainly devoted to the relation between the linear and nonlinear realizations of the electroweak symmetry breaking mechanism with a light scalar Higgs particle. Section III identifies the $d=4$ and $d=5$ flavor-changing couplings. The main phenomenological impact of both $d=4$ and $d=5$ operators is presented in Sec. IV. Finally, we conclude in Sec. V. Technical details on the relation with the strongly-interacting light Higgs (SILH) Lagrangian [39] can be found in Appendix A; the gauge field equations of motion in the presence of flavor-changing contributions are described in Appendix B; the identification of the $d\geq 6$ operators of the linear expansion which correspond to $d=5$ operators of the nonlinear one can be found in Appendix C; the relation between the $d=5$ operator coefficients and the corresponding coefficients in the unitary basis is detailed in Appendix D.

II. THE FRAMEWORK

By "Higgs" we mean here a particle that, at some level, participates in the EW symmetry breaking mechanism, which requires an $SU(2)$ doublet structure. When building up the hybrid situation in which a nonlinear dynamics is assumed but the Higgs is light two strategies are possible: to go from a linear expansion towards a nonlinear one, or conversely to start from the nonlinear realization of the Goldstone boson mechanism and modify it to account for a light Higgs. In general, four (related) scales may be relevant, Λ_s, f, $\langle h\rangle$ and v:

(i) Λ_s is the energy scale of the strong dynamics and the typical size of the mass of the strong scalar and fermionic resonances (in the context of QCD, it corresponds to $\Lambda_{\chi SB}$, the scale of the chiral symmetry breaking [11]).

(ii) f is the characteristic scale associated to the Goldstone bosons that give mass to the gauge bosons and respects $\Lambda_s \leq 4\pi f$ (in the context of QCD, it corresponds to the pion coupling constant f_π).

(iii) $\langle h\rangle$ refers to the order parameter of EW symmetry breaking, around which the physical scalar h oscillates.

(iv) v denotes the EW scale, defined through $M_W = gv/2$. In a general model $\langle h\rangle \neq v$ and this leads to an $\langle h\rangle$ dependence in the low-energy Lagrangian through a generic functional form $\mathcal{F}(h+\langle h\rangle)$.

In nonlinear realizations such as technicolorlike models, it may happen that $\langle h\rangle = v = f$. In the setup considered here with a light h they do not need to coincide, though, although typically a relation links v, $\langle h\rangle$ and f. Thus, a total of three scales will be useful in the analysis, for instance Λ_s, f and v. Without referring to a specific model, one can attempt to describe the NP impact at low energies resorting to an effective Lagrangian approach, with operators made out of SM fields and invariant under the SM gauge symmetry. The transformation properties of the three longitudinal degrees of freedom of the weak gauge bosons can still be described at low energy[2] by a dimensionless unitary matrix transforming as a representation of the global symmetry group:

$$\mathbf{U}(x) = e^{i\sigma_a \pi^a(x)/v}, \qquad \mathbf{U}(x) \to L\mathbf{U}(x)R^\dagger, \tag{2.1}$$

with L, R denoting, respectively, the global transformations $SU(2)_{L,R}$. The adimensionality of $\mathbf{U}(x)$ is the key to understand that the dimension of the leading low-energy operators describing the dynamics of the scalar sector differs for a nonlinear Higgs sector [42–46] and a purely linear regime, as insertions of $\mathbf{U}(x)$ do not exhibit a scale suppression.

It is becoming customary to parametrize the Lagrangian describing a light dynamical Higgs particle h by means of the following ansatz [40,41]:

$$\begin{aligned}\mathcal{L}_h = &\frac{1}{2}(\partial_\mu h)(\partial^\mu h)(1 + c_H\xi\mathcal{F}_H(h)) - V(h)\\ &- \frac{v^2}{4}\,\mathrm{Tr}[\mathbf{V}^\mu\mathbf{V}_\mu]\mathcal{F}_C(h)\\ &+ c_T\frac{v^2}{4}\xi\mathrm{Tr}[\mathbf{T}\mathbf{V}^\mu]\mathrm{Tr}[\mathbf{T}\mathbf{V}_\mu]\mathcal{F}_T(h)\\ &- \left\{\frac{v}{2\sqrt{2}}\bar{Q}_L\mathbf{U}(x)\mathbf{Y}Q_R\mathcal{F}_Y(h) + \mathrm{H.c.}\right\} + \cdots,\end{aligned} \tag{2.2}$$

where dots stand for higher order terms, and $\mathbf{V}_\mu \equiv (\mathbf{D}_\mu\mathbf{U})\mathbf{U}^\dagger$ ($\mathbf{T}\equiv\mathbf{U}\sigma_3\mathbf{U}^\dagger$) is the vector (scalar) chiral field transforming in the adjoint of the gauge group $SU(2)_L$. The covariant derivative reads

[2]Notice that in this low-energy expression for $\mathbf{U}(x)$, the scale associated to the eaten Goldstone Bosons is v and not f. Technically, the scale v appears through a redefinition of the Goldstone Boson fields so as to have canonically normalized kinetic terms.

$$\mathbf{D}_\mu \mathbf{U}(x) \equiv \partial_\mu \mathbf{U}(x) + \frac{ig}{2} W^a_\mu(x)\sigma_a \mathbf{U}(x) - \frac{ig'}{2} B_\mu(x)\mathbf{U}(x)\sigma_3, \tag{2.3}$$

with W^a_μ (B_μ) denoting the $SU(2)_L$ [$U(1)_Y$] gauge bosons and g (g') the corresponding gauge coupling. In these equations, $V(h)$ denotes the effective scalar potential describing the breaking of the electroweak symmetry, the first term in Eq. (2.2) includes the Higgs kinetic term, while the second line describes the W and Z masses and their interactions with h, and the third line shows the usual custodial symmetry breaking term. Finally, restricting our considerations to the quark sector, the fourth line accounts for the Yukawa-like interactions between h and the SM quarks, grouped in doublets of the global symmetry $Q_{L,R}$, with $\mathbf{Y}$ being a 6×6 block diagonal matrix containing the usual Yukawa matrices Y_U and Y_D. The parameters c_H and c_T are model-dependent operator coefficients.

The functions $\mathcal{F}_i(h)$ in Eq. (2.2), as well as other $\mathcal{F}(h)$ functions defined below, encode the generic dependence on the light h particle. Each $\mathcal{F}(h)$ function can be expanded in powers of ξ, $\mathcal{F}(h) = g_0(h, v) + \xi g_1(h, v) + \xi^2 g_2(h, v) + \cdots$, where $g(h, v)$ are model-dependent functions of h. We will not need to enter in their precise dependence in this work; a discussion can be found in Ref. [36] and references therein. We just mention here that in previous literature [40,41] the functional dependence of some of those functions has been expressed as a power series in h/v:

$$\mathcal{F}_C(h) = \left(1 + 2a\frac{h}{v} + b\frac{h^2}{v^2} + \cdots\right),$$
$$\mathcal{F}_Y(h) = \left(1 + c\frac{h}{v} + \cdots\right).$$

The constants a, b and c are model-dependent parameters and encode the dependence on ξ. The a and c_T parameters are constrained from electroweak precision tests: in particular $0.7 \lesssim a \lesssim 1.2$ [47] and $-1.7 \times 10^{-3} < c_T\xi < 1.9 \times 10^{-3}$ [39] at 95% C.L.

The Lagrangian discussed above can be very useful to describe an extended class of Higgs models, ranging from the SM scenario (for $\langle h \rangle = v$, $a = b = c = 1$ and neglecting higher order terms in h), to the technicolorlike ansatz (for $f \sim v$ and omitting all terms in h) and intermediate situations with a light scalar h (in general for $f \neq v$) as in composite or holographic Higgs models [10,12–17,48–50] up to dilatonlike scalar frameworks [51–57]. Note that, although electroweak radiative corrections severely constraint technicolorlike scenarios, in concrete models values of v/f as large as $v/f \sim 0.4$–0.6 are still allowed at the price of acceptable 10% fine-tunings [18,58]. As a result, the study of higher dimension operators is strongly motivated, especially as the limits on ξ are quite model dependent: in the effective Lagrangian approach ξ will be left free $0 < \xi < 1$ while the constraints on custodial breaking effects will be translated into limits on the operator coefficients. For the case of pure gauge and h-gauge couplings, some of the couplings have been explicitly explored in Refs. [39–41] and a complete basis of independent operators up to dimension five has been provided in Ref. [36].

The ξ parameter in Eq. (1.1) defines the degree of nonlinearity of a specific model and in particular $\xi \to 0$ refers to the linear regime, while $\xi \to 1$ to the nonlinear one. For $\xi \ll 1$ the hierarchy between operators mimics that in the linear expansion, where the operators are written in terms of the Higgs doublets H: couplings with a higher number of (physical) Higgs legs are suppressed compared to the SM renormalizable ones, through powers of the high NP scale or, in other words, of ξ [11]. The power of ξ keeps then track of the h dependence of the $d > 4$ operators, where the insertions of h enter only through powers of $(\langle h \rangle + h)/f \simeq \xi^{1/2}(v + h)/v$ and of $\partial_\mu h/f^2$ (see Ref. [36]). In the $\xi \ll 1$ limit, the $\mathcal{F}_i(h)$ functions, appearing in Eq. (2.2) and in the following, would inherit the same universal behavior in powers of $(1 + h/v)$: at order ξ, that is, for couplings that would correspond to $d = 6$ operators of the linear expansion, it follows that

$$\mathcal{F}_i(h) = F(h) \equiv \left(1 + \frac{h}{v}\right)^2. \tag{2.4}$$

An obvious extrapolation applies to the case of couplings weighted by higher powers of ξ, that is, with $d > 6$.

When $\xi \approx 1$ the ξ dependence does not entail a suppression of operators compared to the renormalizable SM operators and the chiral expansion should instead be adopted, although it should be clarified at which level the effective expansion on h/f should stop. Below, the $\mathcal{F}(h)$ functions will be considered completely general polynomial of $\langle h \rangle$ and h (in particular not of derivatives of h) and, when using equations of motion and integration by parts to relate operators, they would be assumed to be redefined when convenient, much as one customarily redefines the constant operator coefficients.

To analyze the passage from the linear to the nonlinear regime, it is an interesting exercise to explore the transition from a $SU(2)_L \times U(1)_Y$ invariant effective Lagrangian in the linear realization of the EW symmetry breaking mechanism to an effective chiral Lagrangian. For instance, in the so-called SILH framework, operators may be written in either the linear [39] [i.e., using the $H(x)$ doublet] or the nonlinear [40,41] (i.e., using the $\mathbf{U}$ matrix and a scalar field h) formalism. We have revisited this procedure in Appendix A.

III. THE FLAVOR SECTOR

The choice of operator basis most suitable when analyzing fermionic couplings is in general one in which fermionic fields participate in the operators.

The flavor-changing sector has not been explicitly taken into consideration in previous analysis of the effective

Lagrangian for a strong interacting light Higgs. Flavor-changing terms do appear in the equations of motion for the gauge field strengths in the presence of the effective operators, and the explicit expressions can be found in Appendix B. However, to include them explicitly would translate into corrections to flavored observables that are quadratic in the effective operator coefficients a_i, and more precisely of the type $\mathcal{O}(a_{FC}\times a_{GH})$, where a_{FC} (a_{GH}) represent the generic flavor-changing (gauge-h) coefficients. Given that a_{GH} are severely constrained by EW data (barring extreme fine-tunings), those quadratic corrections can be disregarded in the rest of the analysis and it is enough to consider the SM equations of motion:

$$(D^\mu W_{\mu\nu})_j = i\frac{g}{4}v^2\mathrm{Tr}[\mathbf{V}_\nu\sigma_j] + \frac{g}{2}\bar{Q}_L\gamma_\nu\sigma_j Q_L, \tag{3.1}$$

$$\partial^\mu B_{\mu\nu} = -i\frac{g'}{4}v^2\mathrm{Tr}[\mathbf{T}\mathbf{V}_\nu] + g'\bar{Q}_L\gamma_\nu\mathbf{h}_L Q_L + g'\bar{Q}_R\gamma_\nu\mathbf{h}_R Q_R. \tag{3.2}$$

In resume, the analysis of the flavor-changing sector can be considered "independent" of that for the gauge-h and flavor-conserving sectors.

A. From $d = 4$ nonlinear operators

With the aid of the (Goldstone) chiral fields $\mathbf{T}$ and $\mathbf{V}_\mu$ it is only possible to write $d = 4$ fermionic operators involving two right-handed or two left-handed fields. In the MFV framework under consideration, only operators built with two left-handed fermions can induce flavor-changing effects at leading order in the spurion expansion. Consequently, terms with two right-handed fermions will not be considered in what follows.

A total of four independent $d = 4$ chiral operators containing left-handed fermion fields can be constructed [35,59–61], namely,

$$\mathcal{O}_1 = \frac{i}{2}\bar{Q}_L\lambda_F\gamma^\mu\{\mathbf{T},\mathbf{V}_\mu\}Q_L, \quad \mathcal{O}_2 = i\bar{Q}_L\lambda_F\gamma^\mu\mathbf{V}_\mu Q_L,$$
$$\mathcal{O}_3 = i\bar{Q}_L\lambda_F\gamma^\mu\mathbf{T}\mathbf{V}_\mu\mathbf{T}Q_L, \quad \mathcal{O}_4 = \frac{1}{2}\bar{Q}_L\lambda_F\gamma^\mu[\mathbf{T},\mathbf{V}_\mu]Q_L. \tag{3.3}$$

Out of these $\mathcal{O}_1 - \mathcal{O}_3$ are CP-even while $\mathcal{O}_4$ is intrinsically CP-odd [35].

Following the discussion in Sec. II, it is pertinent to extend the definition of these chiral couplings in order to include the possibility of a light scalar degree of freedom (related to the EW symmetry breaking), through the ansatz:

$$\mathcal{L}^f_{\chi=4} = \xi\sum_{i=1,2,3}\hat{a}_i\mathcal{O}_i(h) + \xi^2\hat{a}_4\mathcal{O}_4(h), \tag{3.4}$$

where a redefinition by powers of ξ of the operators coefficients defined in Ref. [35] has been implemented, $a_i \equiv \xi\hat{a}_i$ for $i = 1, 2, 3$, while $a_4 \equiv \xi^2\hat{a}_4$. Furthermore

$$\mathcal{O}_i(h) \equiv \mathcal{O}_i\mathcal{F}_i(h), \tag{3.5}$$

where again the functions $\mathcal{F}_i(h)$ contain the dependence on $(h + \langle h\rangle)$. In the present work—restrained to effective couplings of total dimension $d \leq 5$—only terms linear in h should be retained in Eq. (3.4); for the same reason it is neither pertinent to consider couplings containing $\partial_\mu h$ [that is, derivatives of $\mathcal{F}(h)$]. For $\xi \ll 1$, the functions $\mathcal{F}_i(h)$ collapse into combinations of $F(h)$ as defined in Eq. (2.4) for the linear regime:

$$\begin{aligned}
\mathcal{O}_1(h) &\equiv \mathcal{O}_1F(h)(1+\alpha_1\xi F(h)),\\
\mathcal{O}_2(h) &\equiv \mathcal{O}_2F(h)(1+\alpha_2\xi F(h)),\\
\mathcal{O}_3(h) &\equiv \mathcal{O}_3F(h)(1+\alpha_3\xi F(h)),\\
\mathcal{O}_4(h) &\equiv \mathcal{O}_4F^2(h).
\end{aligned} \tag{3.6}$$

The powers of ξ in Eqs. (3.4) and (3.6) facilitate the identification of the lowest dimension at which a "sibling" operator appears in the linear regime. By sibling we mean an operator written in terms of H, that includes the couplings $\mathcal{O}_{1-4}$. For instance, the lowest-dimension siblings of $\mathcal{O}_1$ and $\mathcal{O}_2$ arise at $d = 6$, while that of $\mathcal{O}_4$ appears at $d = 8$ [35]. The case of $\mathcal{O}_3$ is special: indeed, it corresponds to a combination of a $d = 6$ and a $d = 8$ operator of the linear expansion. The parametrization in Eq. (3.6) reflects this correspondence, where for all the operators the contributions from siblings up to $d = 8$ have been accounted for (further contributions will arise considering higher-dimension siblings).

For $\xi \ll 1$ it is consistent to retain just the terms linear in ξ and neglect the contributions from $\mathcal{O}_4(h)$, while it can be shown [35] that $\mathcal{O}_3(h)$ coincides with $-\mathcal{O}_2(h)$ and finally only two linearly independent flavored operators remain [e.g., $\mathcal{O}_1(h)$ and $\mathcal{O}_2(h)$], as previously studied in the literature. On the contrary, in the $\xi \sim 1$ limit all four operators are on the same footing, higher order terms in ξ may contribute, and one recognizes the need of a QCD-like resummation. In particular any chiral operator is made up by an infinite combination of linear ones, an effect represented by the generic $\mathcal{F}_i(h)$ functions, which admit in general an expansion in powers of ξ as discussed previously.

In Ref. [35] we set limits on the coefficients of the operators $\mathcal{O}_1 - \mathcal{O}_4$ from the analysis of $\Delta F = 1$ and $\Delta F = 2$ observables. The inclusion of a light scalar h does not modify the bounds obtained there for the overall coefficients. In fact, the overall operator coefficients in Eq. (3.4) may differ from their Higgsless counterparts in Eq. (3.3) only through a (negligible) loop contribution.

With the inclusion of the light h field, the low-energy effective flavor Lagrangian induced by the SM and the $\mathcal{O}_1(h) - \mathcal{O}_4(h)$ operators in Eq. (3.5) reads, in the unitary gauge [i.e., $\mathbf{U}(x) = \mathbb{1}$] and up to $d = 5$ couplings,

$$
\begin{aligned}
\mathcal{L}^f_{\chi=4} = & -\frac{g}{\sqrt{2}}[W^+_\mu \bar{U}_L \gamma^\mu [a_W(1+\beta_W h/v) \\
& + i a_{CP}(1+\beta_{CP} h/v)](\mathbf{y}_U^2 V + V \mathbf{y}_D^2) D_L + \text{H.c.}] \\
& - \frac{g}{2\cos\theta_W} Z_\mu [a_Z^u \bar{U}_L \gamma^\mu (\mathbf{y}_U^2 + V \mathbf{y}_D^2 V^\dagger) \\
& \times U_L(1+\beta_Z^u h/v) + a_Z^d \bar{D}_L \gamma^\mu (\mathbf{y}_D^2 + V^\dagger \mathbf{y}_U^2 V) \\
& \times D_L(1+\beta_Z^u h/v)],
\end{aligned} \tag{3.7}
$$

where

$$
\begin{aligned}
a_Z^u &\equiv a_1 + a_2 + a_3, \qquad a_Z^d \equiv a_1 - a_2 - a_3, \\
a_W &\equiv a_2 - a_3, \qquad a_{CP} \equiv -a_4.
\end{aligned} \tag{3.8}
$$

The arbitrary coefficients β_i in Eq. (3.7) follow a similar rearrangement to that for a_i in Eq. (3.8), once the $\mathcal{F}(h)$ functions are expanded to first order in h, $\mathcal{F}_i(h) \sim (1 + \beta_i h + \cdots)$; in general each β_i may receive contributions from all orders in ξ for large ξ.

All limits obtained in Ref. [35] for the values of a_Z^d, a_W and a_{CP} resulted from tree-level contributions to observables. It is interesting—and necessary when considering $d = 5$ effective couplings—to analyze as well the possible bounds on the a_i coefficients from their contribution (still disregarding the h insertions) at loop level to radiative processes, such as $b \to s\gamma$ decay. Indeed, the modification of the CKM matrix has a non-negligible impact in the branching ratio of this observable and its precision on both the experimental determination and the theoretical prediction constrains significantly the $a_W - a_{CP}$ parameter space, as we will show in Sec. IV.

Finally, an important difference with strongly interacting heavy Higgs scenarios is the presence at low energies of vertices with additional h external legs, as indicated by Eq. (3.7). This implies interesting phenomenological consequences that will be illustrated later on.

B. From $d = 5$ nonlinear operators

To our knowledge, no discussion of $d = 5$ flavor-changing *chiral* operators has been presented in literature. They may contribute at tree level to relevant flavor-changing observables, as for instance the $b \to s\gamma$ branching ratio. In this subsection all $d = 5$ flavor-changing chiral operators are identified, while interesting phenomenological consequences will be discussed in Sec. IV.

Gauge invariant $d = 5$ operators relevant for flavor must have a bilinear structure in the quark fields of the type $\bar{Q}_L(\ldots)\mathbf{U}(x)Q_R$, where dots stand for objects that transform in the trivial or in the adjoint representation of $SU(2)_L$. Besides the vector and scalar chiral fields $\mathbf{V}_\mu$ and $\mathbf{T}$, they can contain either the rank-2 antisymmetric tensor $\sigma_{\mu\nu}$ or the strength tensors $B_{\mu\nu}$, $W_{\mu\nu}$ and $G_{\mu\nu}$. According to their Lorentz structure, the resulting independent $d = 5$ chiral couplings can be classified in three main groups:

(i) *Dipole-type operators*:

$$
\begin{aligned}
\mathcal{X}_1 &= g' \bar{Q}_L \sigma^{\mu\nu} \mathbf{U} Q_R B_{\mu\nu}, \\
\mathcal{X}_2 &= g' \bar{Q}_L \sigma^{\mu\nu} \mathbf{T}\mathbf{U} Q_R B_{\mu\nu}, \\
\mathcal{X}_3 &= g \bar{Q}_L \sigma^{\mu\nu} \sigma_i \mathbf{U} Q_R W^i_{\mu\nu}, \\
\mathcal{X}_4 &= g \bar{Q}_L \sigma^{\mu\nu} \sigma_i \mathbf{T}\mathbf{U} Q_R W^i_{\mu\nu}, \\
\mathcal{X}_5 &= g_s \bar{Q}_L \sigma^{\mu\nu} \mathbf{U} Q_R G_{\mu\nu}, \\
\mathcal{X}_6 &= g_s \bar{Q}_L \sigma^{\mu\nu} \mathbf{T}\mathbf{U} Q_R G_{\mu\nu}, \\
\mathcal{X}_7 &= g \bar{Q}_L \sigma^{\mu\nu} \mathbf{T} \sigma_i \mathbf{U} Q_R W^i_{\mu\nu}, \\
\mathcal{X}_8 &= g \bar{Q}_L \sigma^{\mu\nu} \mathbf{T} \sigma_i \mathbf{T}\mathbf{U} Q_R W^i_{\mu\nu}.
\end{aligned} \tag{3.9}
$$

(ii) *Operators containing the rank-2 antisymmetric tensor* $\sigma^{\mu\nu}$:

$$
\begin{aligned}
\mathcal{X}_9 &= \bar{Q}_L \sigma^{\mu\nu} [\mathbf{V}_\mu, \mathbf{V}_\nu] \mathbf{U} Q_R, \\
\mathcal{X}_{10} &= \bar{Q}_L \sigma^{\mu\nu} [\mathbf{V}_\mu, \mathbf{V}_\nu] \mathbf{T}\mathbf{U} Q_R, \\
\mathcal{X}_{11} &= \bar{Q}_L \sigma^{\mu\nu} [\mathbf{V}_\mu \mathbf{T}, \mathbf{V}_\nu \mathbf{T}] \mathbf{U} Q_R, \\
\mathcal{X}_{12} &= \bar{Q}_L \sigma^{\mu\nu} [\mathbf{V}_\mu \mathbf{T}, \mathbf{V}_\nu \mathbf{T}] \mathbf{T}\mathbf{U} Q_R.
\end{aligned} \tag{3.10}
$$

(iii) *Other operators containing the chiral vector fields* $\mathbf{V}_\mu$:

$$
\begin{aligned}
\mathcal{X}_{13} &= \bar{Q}_L \mathbf{V}_\mu \mathbf{V}^\mu \mathbf{U} Q_R, \\
\mathcal{X}_{14} &= \bar{Q}_L \mathbf{V}_\mu \mathbf{V}^\mu \mathbf{T}\mathbf{U} Q_R, \\
\mathcal{X}_{15} &= \bar{Q}_L \mathbf{V}_\mu \mathbf{T} \mathbf{V}^\mu \mathbf{U} Q_R, \\
\mathcal{X}_{16} &= \bar{Q}_L \mathbf{V}_\mu \mathbf{T} \mathbf{V}^\mu \mathbf{T}\mathbf{U} Q_R, \\
\mathcal{X}_{17} &= \bar{Q}_L \mathbf{T} \mathbf{V}_\mu \mathbf{T} \mathbf{V}^\mu \mathbf{U} Q_R, \\
\mathcal{X}_{18} &= \bar{Q}_L \mathbf{T} \mathbf{V}_\mu \mathbf{T} \mathbf{V}^\mu \mathbf{T}\mathbf{U} Q_R.
\end{aligned} \tag{3.11}
$$

A fourth group of operators can be constructed from the antisymmetric rank-2 chiral tensor, that transforms in the adjoint of $SU(2)_L$:

$$
\begin{aligned}
\mathbf{V}_{\mu\nu} &\equiv \mathcal{D}_\mu \mathbf{V}_\nu - \mathcal{D}_\nu \mathbf{V}_\mu \\
&= ig \mathbf{W}_{\mu\nu} - i\frac{g}{2} B_{\mu\nu} \mathbf{T} + [\mathbf{V}_\mu, \mathbf{V}_\nu].
\end{aligned} \tag{3.12}
$$

However, the second equality in Eq. (3.12) shows that operators including $\mathbf{V}_{\mu\nu}$ are not linearly independent from those listed in Eqs. (3.9) and (3.10).

The chiral Lagrangian containing the 18 fermionic flavor-changing $d = 5$ operators can thus be written as

$$
\mathcal{L}^f_{\chi=5} = \sum_{i=1}^{18} b_i \frac{\mathcal{X}_i}{\Lambda_s}, \tag{3.13}
$$

where Λ_s is the scale of the strong dynamics and b_i are arbitrary $\mathcal{O}(1)$ coefficients. It is worth to underline that for the analysis of $d = 5$ operators in the nonlinear regime, the relevant scale is Λ_s and not f as for the analysis in the previous section. Indeed, f is associated to light Higgs insertions, while Λ_s refers to the characteristic scale of the strong resonances that, once integrated out, give rise to the operators listed in Eqs. (3.9), (3.10), and (3.11).

A redefinition of the coefficients allows one to make explicit the connection to their lowest-dimension siblings in the linear expansion:

$$\mathcal{L}^f_{\chi=5} = \sqrt{\xi}\sum_{i=1}^{8} \hat{b}_i \frac{\mathcal{X}_i}{\Lambda_s} + \xi\sqrt{\xi}\sum_{i=9}^{18} \hat{b}_i \frac{\mathcal{X}_i}{\Lambda_s}. \qquad (3.14)$$

In the limit of small ξ, $\mathcal{X}_{1-6}$ correspond to $d=6$ operators in the linear expansion, while $\mathcal{X}_7$ and $\mathcal{X}_8$ result from combinations of $d=6$ and $d=8$ siblings. Moreover, $\mathcal{X}_{9-18}$ have linear siblings of $d=8$, but $\mathcal{X}_{17}$ and $\mathcal{X}_{18}$ that are combinations of $d=8$ and $d=10$ operators in the linear regime. The complete list of the linear siblings of the chiral $d=5$ operators can be found in Appendix C.

Because in this work the analysis will be restrained to (at most) $d=5$ couplings, it is not necessary nor pertinent to discuss further the possible extensions $\mathcal{X}_i \to \mathcal{X}_i(h)$ that would include the dependence on a light Higgs through generic $\mathcal{F}_i(h)$ functions.

The phenomenological impact of these contributions can be best identified through the low-energy Lagrangian written in the unitary gauge:

$$\delta\mathcal{L}_{\chi=5} = \delta\mathcal{L}^u_{\chi=5} + \delta\mathcal{L}^d_{\chi=5} + \delta\mathcal{L}^{ud}_{\chi=5}, \qquad (3.15)$$

where

$$\begin{aligned}\delta\mathcal{L}^d_{\chi=5} = {} & \frac{g^2}{4\cos\theta_W^2}\frac{b^d_Z}{\Lambda_s}\bar{D}_L D_R Z_\mu Z^\mu + \frac{g^2}{2}\frac{b^d_W}{\Lambda_s}\bar{D}_L D_R W^+_\mu W^{-\mu} + g^2\frac{c^d_W}{\Lambda_s}\bar{D}_L\sigma^{\mu\nu}D_R W^+_\mu W^-_\nu \\ & + e\frac{d^d_F}{\Lambda_s}\bar{D}_L\sigma^{\mu\nu}D_R F_{\mu\nu} + \frac{g}{2\cos\theta_W}\frac{d^d_Z}{\Lambda_s}\bar{D}_L\sigma^{\mu\nu}D_R Z_{\mu\nu} + g_s\frac{d^d_G}{\Lambda_s}\bar{D}_L\sigma^{\mu\nu}D_R G_{\mu\nu} + \text{H.c.},\end{aligned} \qquad (3.16)$$

$$\begin{aligned}\delta\mathcal{L}^{ud}_{\chi=5} = {} & \frac{g^2}{2\sqrt{2}\cos\theta_W}\left(\frac{b^+_{WZ}}{\Lambda_s}\bar{U}_L D_R W^+_\mu Z^\mu + \frac{b^-_{WZ}}{\Lambda_s}\bar{D}_L U_R W^-_\mu Z^\mu\right) \\ & + \frac{g^2}{2\sqrt{2}\cos\theta_W}\left(\frac{c^+_{WZ}}{\Lambda_s}\bar{U}_L\sigma^{\mu\nu}D_R W^+_\mu Z_\nu + \frac{c^-_{WZ}}{\Lambda_s}\bar{D}_L\sigma^{\mu\nu}U_R W^-_\mu Z_\nu\right) \\ & + \frac{g}{\sqrt{2}}\left(\frac{d^+_W}{\Lambda_s}\bar{U}_L\sigma^{\mu\nu}D_R W^+_{\mu\nu} + \frac{d^-_W}{\Lambda_s}\bar{D}_L\sigma^{\mu\nu}U_R W^-_{\mu\nu}\right) + \text{H.c.},\end{aligned} \qquad (3.17)$$

and analogously for $\delta\mathcal{L}^u_{\chi=5}$ as in $\delta\mathcal{L}^d_{\chi=5}$ interchanging $d \leftrightarrow u$ and $D_{L,R} \leftrightarrow U_{L,R}$. In these equations $W^\pm_{\mu\nu} = \partial_\mu W^\pm_\nu - \partial_\nu W^\pm_\mu \pm ig(W^3_\mu W^\pm_\nu - W^3_\nu W^\pm_\mu)$, while the photon and Z field strengths are defined as $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$ and $Z_{\mu\nu} = \partial_\mu Z_\nu - \partial_\nu Z_\mu$, respectively, and $W^3_\mu = \cos\theta_W Z_\mu + \sin\theta_W A_\mu$. The relations between the coefficients appearing in Eqs. (3.16) and (3.17) and those defined in Eq. (3.13) are reported in Appendix D.

IV. PHENOMENOLOGICAL ANALYSIS

This section first resumes and updates the bounds existing in the literature [35] on the coefficients of the flavor-changing $d=4$ chiral expansion, and then discusses new bounds and other phenomenological considerations with and without a light Higgs:

(i) loop-level impact of fermionic $d=4$ chiral operators ($\mathcal{O}_1$ to $\mathcal{O}_4$) on those same radiative decays;
(ii) tree-level bounds on the fermionic $d=5$ chiral operators $\mathcal{X}_i$, from radiative decays;
(iii) light Higgs to fermions couplings, from operators $\mathcal{O}_1(h)$ to $\mathcal{O}_4(h)$.

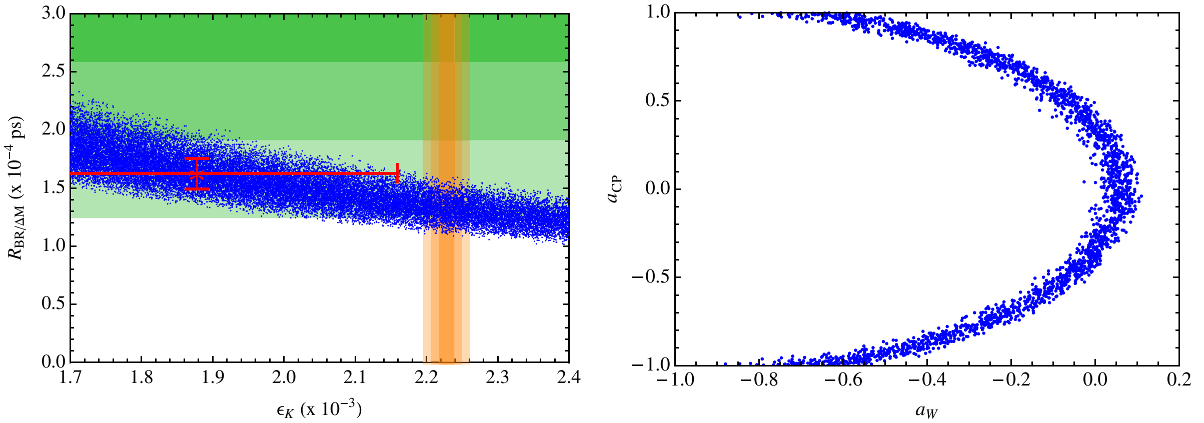
A. $\Delta F = 1$ and $\Delta F = 2$ observables

In Ref. [35], the constraints on the coefficients of the $d=4$ flavor-changing operators of the nonlinear expansion have been analyzed. These bounds resulted from $\Delta F=1$ and $\Delta F=2$ observables and apply straightforwardly to nonlinear regimes with a light h scalar. Operators $\mathcal{O}_1$, $\mathcal{O}_2$ and $\mathcal{O}_3$ induce tree-level contributions to $\Delta F=1$ processes mediated by the Z boson, as can be seen from the Z couplings in the effective Lagrangian in Eq. (3.7), and are severely constrained. Because of the MFV structure of the coefficients, sizable flavor-changing effects may only be expected in the down quark sectors, with data on K and B transitions providing the strongest constraints on a^d_Z,

$$-0.044 < a^d_Z < 0.009 \quad \text{at } 95\% \text{ of C.L.} \qquad (4.1)$$

from $K^+ \to \pi^+\bar{\nu}\nu$, $B \to X_s\ell^+\ell^-$ and $B \to \mu^+\mu^-$ data.

Furthermore, operators $\mathcal{O}_2$, $\mathcal{O}_3$ and $\mathcal{O}_4$ induce corrections to the fermion-W couplings, and thus to the CKM matrix; see Eq. (3.7). This in turn induces modifications [35] on the strength of meson oscillations (at loop level), on $B^+ \to \tau^+\nu$ decay and on the B semileptonic



(a) Correlation plot between ε_K and $R_{BR/\Delta M}$. $a_W, a_{CP} \in [-1, 1]$, $a_Z^d \in [-0.1, 0.1]$

(b) $a_W - a_{CP}$ parameter space for the observables on the left panel inside their 3σ error ranges and $a_Z^d \in [-0.044, 0.009]$.

FIG. 1 (color online). Results for the reference point $(|V_{ub}|, \gamma) = (3.5 \times 10^{-3}, 66°)$. Left panel: In red the SM prediction and its 1σ theoretical error bands for ε_K and $R_{\mathrm{BR}/\Delta M}$ for this reference point; in orange (green) the 1σ, 2σ and 3σ (from the darker to the lighter) experimental error ranges for ε_K ($R_{\mathrm{BR}/\Delta M}$), in blue the correlation between ε_K and $R_{\mathrm{BR}/\Delta M}$ induced by NP contributions. Right panel: Allowed values for a_W and a_{CP} upon the setup of the left panel. See Ref. [35] for further details.

CP asymmetry, among others; more specifically the following processes have been taken into account in Ref. [35]:

(i) *The CP-violating parameter ϵ_K of the $K^0 - \bar{K}^0$ system and the mixing-induced CP asymmetries $S_{\psi K_S}$ and $S_{\psi\phi}$ in the decays $B_d^0 \to \psi K_S$ and $B_s^0 \to \psi\phi$.*—The corrections induced to ϵ_K are proportional to y_t^2, while those to $S_{\psi K_S}$ and $S_{\psi\phi}$ are proportional to y_b^2. Consequently, possible large deviations from the values predicted by the SM are only allowed in the K system.

(ii) *The ratio among the meson mass differences in the B_d and B_s systems, $R_{\Delta M_B} \equiv \Delta M_{B_d}/\Delta M_{B_s}$.*—The NP contributions on the mass differences almost cancel in this ratio and therefore deviations from the SM prediction for this observable are negligible.

(iii) *The ratio among the $B^+ \to \tau^+\nu$ branching ratio and the B_d mass difference, $R_{\mathrm{BR}/\Delta M} \equiv \mathrm{BR}(B^+ \to \tau^+\nu)/\Delta M_{B_d}$.*—This observable is clean from theoretical hadronic uncertainties and the constraints on the NP parameters are therefore potentially strong.

Since only small deviations from the SM prediction for $S_{\psi K_S}$ are allowed, only values close to the exclusive determination for $|V_{ub}|$ are favored [35] for a complete discussion. Moreover, it is possible to constrain the $|V_{ub}| - \gamma$ parameter space, with γ being one of the angles of the unitary triangle, requiring that both $S_{\psi K_S}$ and $R_{\Delta M_B}$ observables are inside the 3σ experimental determination.

Once this reduced parameter space is identified, it is illustrative to choose one of its points as a reference point, in order to present the features of this MFV scenario; for instance for the values $(|V_{ub}|, \gamma) = (3.5 \times 10^{-3}, 66°)$, $S_{\psi K_S}$, $R_{\Delta M_B}$ and $|V_{ub}|$ are all inside their own 1σ values, and the predicted SM values for ϵ_K and $R_{\mathrm{BR}/\Delta M}$ are[3]

$$\epsilon_K = 1.88 \times 10^{-3}, \qquad R_{\mathrm{BR}/\Delta M} = 1.62 \times 10^{-4}. \quad (4.2)$$

The errors on these quantities are $\sim 15\%$ and $\sim 8\%$, respectively, estimated considering the uncertainties on the input parameters and the analysis performed in Ref. [63]. Figure 1 shows the correlation between ϵ_K and $R_{\mathrm{BR}/\Delta M}$ (left panel) and the $a_{CP} - a_W$ parameter space (right panel), requiring that ϵ_K and $R_{\mathrm{BR}/\Delta M}$ lie inside their own 3σ experimental determination. Finally, for those points in the $a_{CP} - a_W$ parameter space that pass all the previous constraints, the predictions for $S_{\psi\phi}$ and the B semileptonic *CP* asymmetry turned out to be close to the SM determination, in agreement with the recent LHCb measurements [64].

In the next subsection, new constraints on the $d = 4$ operator coefficients a_W and a_{CP} will be obtained from their impact at loop level on radiative B decays. The latter data will be also used to constrain the set of $d = 5$ chiral operators coefficients identified in Sec. III B: while they are expected *a priori* to be all of comparable strength, the most powerful experimental constraints should result from the tree-level impact of dipole operators $\mathcal{X}_1$ to $\mathcal{X}_8$, as they include vertices involving just three fields, one of them being a light gauge boson. Photonic penguins and also gluonic penguins and tree-level four-fermion diagrams (through renormalization group mixing effects) will be explored below and contrasted with radiative B decays.

[3]The predicted SM value for ϵ_K differs from that in Ref. [35] due to the new input parameters used: in particular $\hat{B}_K = 0.7643 \pm 0.0097$ has sensibly increased [62].

B. $\bar{B} \to X_s\gamma$ branching ratio

The current experimental value of the $\bar{B} \to X_s\gamma$ branching ratio [65] is

$$\mathrm{Br}(\bar{B} \to X_s\gamma) = (3.55 \pm 0.24 \pm 0.09) \times 10^{-4}, \quad (4.3)$$

for a photon-energy cutoff $E_\gamma > 1.6$ GeV. On the other hand, its next-to-next-to-leading order SM prediction for that same energy cutoff and in the $\bar{B}$-meson rest frame, reads [66–68]

$$\mathrm{Br}\,(\bar{B} \to X_s\gamma) = (3.15 \pm 0.23) \times 10^{-4}. \quad (4.4)$$

The presence of NP can easily modify this prediction, and the precision of both the experimental measure and the SM computation allows one in principle to provide severe bounds on the NP parameters.

The effective Lagrangian relevant for $b \to s\gamma$ decay at the $\mu_b = \mathcal{O}(m_b)$ scale can be written as

$$\mathcal{L}_{\mathrm{eff}} = \frac{4G_F}{\sqrt{2}} V_{ts}^* V_{tb}\Bigg[\sum_{i=1}^{6} C_i(\mu_b)Q_i(\mu_b) + C_{7\gamma}(\mu_b)Q_{7\gamma}(\mu_b) + C_{8G}(\mu_b)Q_{8G}(\mu_b)\Bigg], \quad (4.5)$$

where $Q_{1,2}$, $Q_{3,\ldots,6}$ and $Q_{7\gamma,8G}$ denote the current-current, QCD penguin and magnetic dipole operators, respectively, as it is customary. In this effective Lagrangian, subleading terms proportional to $V_{us}^* V_{ub}$ have been neglected; the same applies to the contributions from the so-called primed operators, similar to those appearing in Eq. (4.5) although with opposite chirality structure, which are suppressed by the m_s/m_b ratio.

The value of the Wilson coefficients $C_i(\mu_b)$ at the scale μ_b is derived applying the QCD renormalization group (RG) analysis to the corresponding Wilson coefficients, evaluated at the effective scale μ of the underlying theory, which is the matching scale linking the effective and full descriptions. For the SM case, this is the electroweak scale $\mu_W = \mathcal{O}(M_W)$. The effects of the RG contributions are in general non-negligible, and indeed the rate of the $b \to s\gamma$ decay in the SM is enhanced by a factor of 2–3 [67] upon the inclusion of these corrections. They originate dominantly from the mixing of charged current-current operators with the dipole operators, and to a smaller extent from the mixing with QCD-penguin operators. These QCD contributions can be formally written as

$$C_i(\mu_b) = U_{ij}(\mu_b, \mu)C_j(\mu), \quad (4.6)$$

where $U_{ij}(\mu_b, \mu)$ are the elements of the RG evolution matrix from the effective scale μ down to μ_b [69].

The expression for the $\bar{B} \to X_s\gamma$ branching ratio is then given as follows:

$$\mathrm{Br}\,(\bar{B} \to X_s\gamma) = R(|C_{7\gamma}(\mu_b)|^2 + N(E_\gamma)), \quad (4.7)$$

where $R = 2.47 \times 10^{-3}$ is simply an overall factor as discussed in Refs. [66,68] and $N(E_\gamma) = (3.6 \pm 0.6) \times 10^{-3}$ is a nonperturbative contribution for the photon-energy cutoff $E_\gamma > 1.6$ GeV. $C_{7\gamma}(\mu_b)$ can be decomposed into SM and NP contributions,

$$C_{7\gamma}(\mu_b) = C_{7\gamma}^{\mathrm{SM}}(\mu_b) + \Delta C_{7\gamma}(\mu_b), \quad (4.8)$$

where, for $\mu_b = 2.5$ GeV, the SM contribution at the next-to-next-to-leading order level, is given by [66–68]

$$C_{7\gamma}^{\mathrm{SM}}(\mu_b) = -0.3523. \quad (4.9)$$

In our context, the NP contributions arise from the non-unitarity of the CKM matrix and the presence of flavor violating Z-fermion couplings (induced by the $d = 4$ chiral operators $\mathcal{O}_{1-4}$ [35]), and from the direct contributions from the $d = 5$ chiral operators $\mathcal{X}_{1-8}$. In the following we will discuss separately these contributions.

1. $d = 4$ contributions

The effective scale of the $d = 4$ chiral operators is $f \geq v$, but no contributions to the Wilson coefficients relevant for $b \to s\gamma$ arise at scales above the electroweak one. As a result, the analysis of these contributions is alike to that in the SM, except for the fact that the NP operators modify the initial conditions at μ_W. The Wilson coefficients at the scale μ_W can be written as

$$C_i(\mu_W) = C_i^{\mathrm{SM}}(\mu_W) + \Delta C_i^{d=4}(\mu_W), \quad (4.10)$$

where the SM coefficients at the LO are given by [70]

$$\begin{aligned} C_2^{\mathrm{SM}}(\mu_W) &= 1, \\ C_{7\gamma}^{\mathrm{SM}}(\mu_W) &= \frac{7x_t - 5x_t^2 - 8x_t^3}{24(x_t - 1)^3} + \frac{-2x_t^2 + 3x_t^3}{4(x_t - 1)^4}\log x_t, \\ C_{8G}^{\mathrm{SM}}(\mu_W) &= \frac{2x_t + 5x_t^2 - x_t^3}{8(x_t - 1)^3} + \frac{-3x_t^2}{4(x_t - 1)^4}\log x_t, \end{aligned} \quad (4.11)$$

with $x_t \equiv m_t^2/M_W^2$.

The NP contributions due to the nonunitarity of the CKM matrix induce modifications in all three Wilson coefficients involved:

$$\begin{aligned} \Delta C_2^{d=4}(\mu_W) &= (a_W - i a_{CP})y_b^2 + (a_W^2 + a_{CP}^2)y_b^2 y_c^2, \\ \Delta C_{7\gamma}^{d=4}(\mu_W) &= (2a_W y_t^2 + (a_W^2 + a_{CP}^2)y_t^4) \times\left(\frac{23}{36} + C_{7\gamma}^{\mathrm{SM}}(\mu_W)\right), \\ \Delta C_{8G}^{d=4}(\mu_W) &= (2a_W y_t^2 + (a_W^2 + a_{CP}^2)y_t^4)\left(\frac{1}{3} + C_{8G}^{\mathrm{SM}}(\mu_W)\right). \end{aligned} \quad (4.12)$$

These terms originate from the corresponding SM diagrams with the exchange of a W boson and are proportional to a_W and a_{CP}: indeed they are due to the modified vertex couplings, both in the tree-level diagram that originates Q_2

TABLE I. The magic numbers for $\Delta C_{7\gamma}(\mu_b)$ defined in Eq. (4.13).

i	1	2	3	4	5	6	7	8
σ_i	$\frac{14}{23}$	$\frac{16}{23}$	$\frac{6}{23}$	$-\frac{12}{23}$	0.4086	-0.4230	-0.8994	0.1456
κ_i	2.2996	-1.0880	$-\frac{3}{7}$	$-\frac{1}{14}$	-0.6494	-0.0380	-0.0185	-0.0057

and in the 1-loop penguin diagrams that give rise to $Q_{7\gamma}$ and Q_{8G}. On the other hand, the new flavor-changing Z-fermion vertices participate in penguin diagrams contributing to the $b \to s\gamma$ decay amplitude, with a Z boson running in the loop [71]. These contributions can be safely neglected, though, because they are proportional to the $a_Z^{u,d}$ parameters, which are already severely constrained from their tree-level impact on other FCNC processes.

Including the QCD RG corrections, the NP contributions at LO to the Wilson coefficients are given by

$$\Delta C_{7\gamma}(\mu_b) = \eta^{\frac{16}{23}}\Delta C_{7\gamma}(\mu_W) + \frac{8}{3}(\eta^{\frac{14}{23}} - \eta^{\frac{16}{23}})\Delta C_{8G}(\mu_W) + \Delta C_2(\mu_W)\sum_{i=1}^{8}\kappa_i\eta^{\sigma_i}, \tag{4.13}$$

with

$$\eta \equiv \frac{\alpha_s(\mu_W)}{\alpha_s(\mu_b)} = 0.45. \tag{4.14}$$

Here κ's and σ's are the magic numbers listed in Table I, while η has been calculated taking $\alpha_s(M_Z = 91.1876\ \mathrm{GeV}) = 0.118$. Because of the simple additive structure of the NP contributions in Eq. (4.10), these magic numbers are the same as in the SM context.

The analysis above allows one to estimate the impact of the experimental value for $\mathrm{BR}(\bar{B} \to X_s\gamma)$ on the NP parameter space of $\mathcal{O}_1 \ldots \mathcal{O}_4$ operators: in Fig. 2 we retake the scatter plot shown in Fig. 1(b), based on the analysis of $\Delta F = 1$ and $\Delta F = 2$ observables for the reference point $(|V_{ub}|, \gamma) = (3.5 \times 10^{-3}, 66°)$, and superimpose the new constraints resulting from the present loop-level impact on $\mathrm{BR}(\bar{B} \to X_s\gamma)$: they are depicted as shadowed (gray) exclusion regions. The figure illustrates that they reduce the available parameter space, eliminating about half of the points previously allowed in the scatter plot of Fig. 1(b).

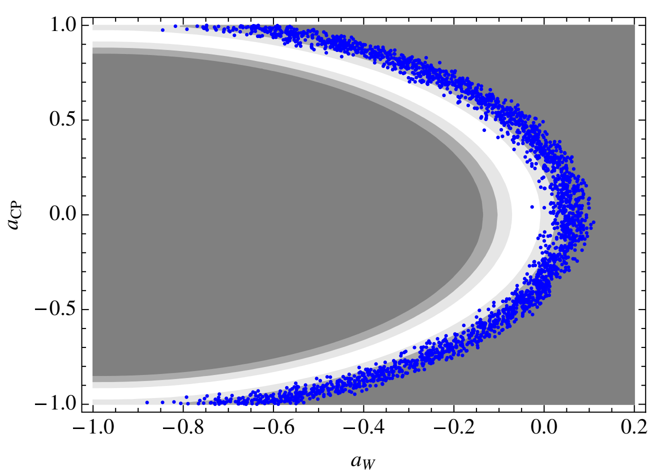


FIG. 2 (color online). $a_W - a_{CP}$ parameter space for ε_K and $\mathrm{BR}(B^+ \to \sigma^+\nu)/\Delta M_{B_d}$ observables inside their 3σ error ranges and $a_Z^d \in [-0.044, 0.009]$ (see Ref. [35] for details). The gray areas correspond to the bounds from the $\mathrm{BR}(\bar{B} \to X_s\gamma)$ at 1σ, 2σ, and 3σ, from the lighter to the darker, respectively.

Figure 2 shows that a_{CP}, the overall coefficient of the genuinely CP-odd coupling $\mathcal{O}_4$, and thus of $\mathcal{O}_4(h)$ in Eq. (3.5), is still loosely constrained by low-energy data. This has an interesting phenomenological consequence on Higgs physics prospects, since it translates into correlated exotic Higgs-fermion couplings, which for instance at leading order in h read

$$\delta\mathcal{L}^h_{\chi=4} \supset a_{CP}\left(1 + \beta_{CP}\frac{h}{v}\right)\mathcal{O}_4. \tag{4.15}$$

For intermediate values of ξ (for which the linear expansion could be an acceptable guideline), the relative weight of the couplings with and without an external Higgs particle reduces to, see Eq. (3.6),

$$\beta_{CP} \sim 4. \tag{4.16}$$

These are encouraging results in the sense of allowing short-term observability. In a conservative perspective, the operator coefficients of the $d = 4$ nonlinear expansion should be expected to be $\mathcal{O}(1)$. Would this be the case, the possibility of NP detection would be delayed until both low-energy flavor experiments and LHC precision on h-fermion couplings nears the $\mathcal{O}(10^{-2})$ level, which for LHC means to reach at least its 3000 fb^{-1} running regime. Notwithstanding this, a steady improvement of the above bounds should be sought.

2. $d = 5$ contributions

For the $d = 5$ chiral operators considered, the effective scale weighting their overall strength is $\Lambda_s \leq 4\pi f$. In the numerical analysis that follows, we will consider for Λ_s the smallest value possible, i.e., $\Lambda_s = 4\pi v$. For this value, the effects due to the $d = 5$ chiral operators are maximized: indeed, for higher scales, the initial conditions for the Wilson coefficients are suppressed with the increasing of Λ_s. This effect is only slightly softened, but not canceled, by the enhancement due to the QCD running from a higher initial scale. For the analytical expressions, we will keep the discussion at a more general level and the high scale will be denoted by μ_s, $\mu_s \gg v$. At this scale the top and W bosons are still dynamical and therefore they do not contribute yet to any Wilson coefficients. The only operators relevant for $b \to s\gamma$ decay and with nonvanishing initial conditions are thus $Q_{7\gamma}$ and Q_{8G}, whose contributions arise from dipole $d = 5$ chiral operator. At the scale μ_s the Wilson coefficients can thus be written as

$$C_i(\mu_s) \equiv C_i^{\rm SM}(\mu_s) + \Delta C_i^{d=5}(\mu_s), \qquad (4.17)$$

where the only nonvanishing contributions are

$$\Delta C_{7\gamma}^{d=5}(\mu_s) = d_F^d \frac{(4\pi)^2 v y_t^2}{\sqrt{2}\mu_s}, \quad \Delta C_{8G}^{d=5}(\mu_s) = d_G^d \frac{(4\pi)^2 v y_t^2}{\sqrt{2}\mu_s}, \qquad (4.18)$$

with d_F^d and d_G^d denoting the relevant photonic and gluonic dipole operator coefficients in Eq. (3.16), respectively.

The QCD RG analysis from μ_s down to μ_b should be performed in two distinct steps:

(i) *A six-flavor RG running from the scale μ_s down to μ_W.*—Focusing on the Wilson coefficients corresponding to the SM and to the $d = 5$ couplings under discussion, at the scale μ_W the coefficients read

$$C_i(\mu_W) \equiv C_i^{\rm SM}(\mu_W) + \Delta C_i^{d=5}(\mu_W), \qquad (4.19)$$

where the only nonvanishing contributions from the set of $d = 5$ flavor-changing fermionic operators are those given by

$$C_{7\gamma}^{d=5}(\mu_W) = \frac{8}{3}(1 - \eta_{\mu_s}^{2/21})\eta_{\mu_s}^{2/3}\Delta C_{8G}^{d=5}(\mu_s) + \eta_{\mu_s}^{16/21}\Delta C_{7\gamma}^{d=5}(\mu_s), \quad C_{8G}^{d=5}(\mu_W) = \eta_{\mu_s}^{2/3}\Delta C_{8G}^{d=5}(\mu_s), \qquad (4.20)$$

with

$$\eta_{\mu_s} \equiv \frac{\alpha_s(\mu_s)}{\alpha_s(\mu_W)}. \qquad (4.21)$$

In the numerical analysis $\eta_{\mu_s} = 0.67$ will be taken.

(ii) *A five-flavor RG running from μ_W down to μ_b.*—This analysis is alike to that described in the previous section, substituting the initial conditions for the Wilson coefficients in Eqs. (4.10), (4.11), and (4.12) for those in Eqs. (4.19), (4.20), and (4.21).

It is interesting to focus on the final numerical result for the BR$(\bar{B} \to X_s\gamma)$, leaving unspecified only the parameters of the $d = 5$ chiral operators $b_{F,G}^d$:

$$\mathrm{BR}\,(b \to s\gamma) = 0.000315 - 0.00175 b_{\rm eff}^d + 0.00247 (b_{\rm eff}^d)^2, \qquad (4.22)$$

where

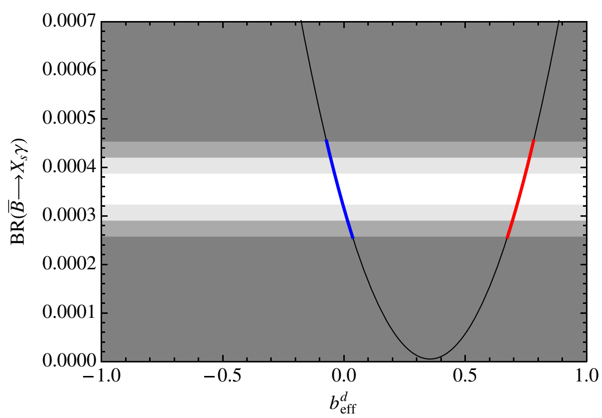
$$b_{\rm eff}^d \equiv 3.8 b_F^d + 1.2 b_G^d. \qquad (4.23)$$

The corresponding plot is shown on the left-hand side of Fig. 3, which depicts the dependence of the branching ratio on $b_{\rm eff}^d$, together with the experimental 3σ regions. Two distinct ranges for $b_{\rm eff}^d$ are allowed:

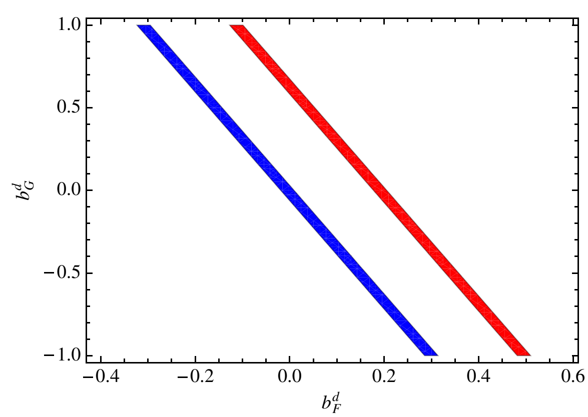
$$-0.07 \lesssim b_{\rm eff}^d \lesssim 0.04 \quad \text{or} \quad 0.67 \lesssim b_{\rm eff}^d \lesssim 0.78. \qquad (4.24)$$

Using the expression for $b_{\rm eff}^d$ in Eq. (4.23), it is possible to translate these bounds onto the $b_F^d - b_G^d$ parameter space, as shown on the right-hand side of Fig. 3. The two narrow bands depict the two allowed regions.

Analogously to the case of $\mathcal{O}_1(h) \ldots \mathcal{O}_4(h)$ operators discussed in the previous subsection, a correlation would hold between a low-energy signal from these $d = 5$ couplings and the detection of exotic fermionic couplings at LHC, upon considering their extension to include h-dependent insertions. Nevertheless, a consistent analysis would require one in this case to consider $d = 6$ couplings



(a) $BR(\bar{B} \to X_s\gamma)$ vs. b_{eff}^d



(b) $b_F^d - b_G^d$ parameter space

FIG. 3 (color online). Left panel: The curve depicts BR$(\bar{B} \to X_s\gamma)$ as a function of $b_{\rm eff}^d$, while the horizontal bands are the experimentally excluded regions at 1σ, 2σ, and 3σ, from the lighter to the darker, respectively. Right panel: The 3σ corresponding allowed $b_F^d - b_G^d$ parameter space is depicted as two separate narrow bands.

of the nonlinear expansion, which are outside the scope of the present work.

V. CONCLUSIONS

The lack of indications of new resonances at LHC data other than a strong candidate to be the SM scalar boson h, together with the alignment of the couplings of the latter with SM expectations, draws a puzzling panorama with respect to the electroweak hierarchy problem. If the experimental pattern persists, either the extended prejudice against fine-tunings of the SM parameters should be dropped, or the new physics scale is still awaiting discovery and may be associated for instance to a dynamical origin of the SM scalar boson. We have focused in this work on possible implications for fermionic couplings of a strong interacting origin of electroweak symmetry breaking dynamics with a light scalar h with mass around 125 GeV, within an effective Lagrangian approach.

The parameter describing the degree of nonlinearity $\xi = (v/f)^2$ must lie in the range $0 < \xi < 1$. For small values, the effective theory converges towards the SM, as the NP contributions can be safely neglected. On the other hand, large values indicate a chiral regime for the dynamics of the Goldstone bosons, which in turn requires one to use a chiral expansion to describe them, combined with appropriate insertions of the light h field.

We identified the flavor-changing operator basis for the nonlinear regime up to $d = 5$. Furthermore, taking into account the QCD RG evolution, the coefficients of these operators have been constrained from a plethora of low-energy transitions. In particular we have analyzed in detail and in depth the constraints resulting from the data on $\bar{B} \to X_s \gamma$ branching ratio. Its impact is important on the global coefficients of the four relevant $d = 4$ flavor-changing chiral couplings at the loop level, and on those of the $d = 5$ dipole operators. The limits obtained constrain in turn the possible fermion-h exotic couplings to be explored at the LHC. A particularly interesting example is that of the intrinsically CP-odd $d = 4$ operator $\mathcal{O}_4$ of the nonlinear expansion, whose coefficient is loosely constrained by data: a correlation is established between the possible signals in low-energy searches of CP violation and anomalous h-fermion couplings at the LHC. Their relative strength is explored for the case of a relatively small ξ. A similar correlation between low-energy flavor searches and LHC signals also follows for all other operators.

ACKNOWLEDGMENTS

We acknowledge partial support by European Union FP7 ITN INVISIBLES (Marie Curie Actions, PITN-GA-2011-289442), CiCYT through Project No. FPA2009-09017, CAM through Project No. HEPHACOS P-ESP-00346, European Union FP7 ITN UNILHC (Marie Curie Actions, PITN-GA-2009-237920), MICINN through Grant No. BES-2010-037869 and the Juan de la Cierva program (JCI-2011-09244), ERC Advance Grant "FLAVOUR" (267104), Technische Universität München—Institute for Advanced Study—funded by the German Excellence Initiative, Italian Ministero dell'Università e della Ricerca Scientifica through the COFIN program (PRIN 2008) and Contracts No. MRTN-CT-2006-035505 and No. PITN-GA-2009-237920 (UNILHC). We thank the Galileo Galilei Institute for Theoretical Physics for the hospitality and the INFN for partial support during the completion of this work. R. A. acknowledges the Harvard Physics department for hospitality during the completion phase of this work. S. R. and J. Y. acknowledge CERN TH department for hospitality during the completion phase of the work.

APPENDIX A: RELATION TO THE SILH BASIS

In this Appendix we revisit the transition from an $SU(2)_L \times U(1)_Y$ invariant effective Lagrangian in the linear realization of the EW symmetry breaking mechanism to an effective chiral Lagrangian, focusing to the so-called SILH framework [39]. The $d = 6$ SILH Lagrangian in Eq. (2.15) of Ref. [39] can be written in terms of $\mathbf{U}$, $\mathbf{V}$, $\mathbf{T}$ and a scalar field h:

$$\begin{aligned}
\mathcal{L}_{\rm SILH} = \xi\Big\{ & \frac{c_H}{2}(\partial_\mu h)(\partial^\mu h)\mathcal{F}(h) + \frac{c_T}{2}\frac{v^2}{4}\,\mathrm{Tr}[\mathbf{T}\mathbf{V}^\mu]\mathrm{Tr}[\mathbf{T}\mathbf{V}_\mu]\mathcal{F}(h)^2 - c_6\lambda\frac{v^4}{8}\mathcal{F}(h)^3 \\
&+ \left(c_y\frac{v}{2\sqrt{2}}\bar{Q}_L\mathbf{U}\,\mathrm{diag}(\mathbf{y_U},\mathbf{y_D})Q_R\mathcal{F}(h)^{3/2} + \mathrm{H.c.}\right) - i\frac{c_W g}{2m_\rho^2}\frac{f^2}{2}(\mathcal{D}_\mu W^{\mu\nu})_i\mathrm{Tr}[\sigma_i\mathbf{V}_\nu]\mathcal{F}(h) \\
&+ i\frac{c_B g'}{2m_\rho^2}\frac{f^2}{2}(\partial_\mu B^{\mu\nu})\mathrm{Tr}[\mathbf{T}\mathbf{V}_\nu]\mathcal{F}(h) + i\frac{c_{HW}g}{16\pi^2}W_i^{\mu\nu}\left(\frac{1}{4}\mathrm{Tr}[\sigma_i\mathbf{V}_\mu\mathbf{V}_\nu]\mathcal{F}(h) - \frac{1}{4}\mathrm{Tr}[\sigma_i\mathbf{V}_\mu]\partial_\nu\mathcal{F}(h)\right) \\
&+ i\frac{c_{HB}g'}{16\pi^2}B^{\mu\nu}\left(\frac{1}{4}\mathrm{Tr}[\mathbf{T}\mathbf{V}_\mu\mathbf{V}_\nu]\mathcal{F}(h) + \frac{1}{4}\mathrm{Tr}[\mathbf{T}\mathbf{V}_\mu]\partial_\nu\mathcal{F}(h)\right) + \frac{c_\gamma g'^2}{16\pi^2}\frac{g^2}{g_\rho^2}\frac{1}{2}B_{\mu\nu}B^{\mu\nu}\mathcal{F}(h) \\
&+ \frac{c_g g_S^2}{16\pi^2}\frac{y_t^2}{g_\rho^2}\frac{1}{2}G^a_{\mu\nu}G^{a\mu\nu}\mathcal{F}(h)\Big\}, \qquad \text{(A1)}
\end{aligned}$$

where the notation of the operator coefficients is as in Ref. [39] and $\mathcal{F}(h) = F(h)$ is the function of the light Higgs fields resulting from the doublet Higgs ansatz as in Eq. (2.4); the Lagrangian above is only complete at leading order for values of $\xi \ll 1$; otherwise other operators of nonlinear parenthood have to be added, as earlier explained.

APPENDIX B: GAUGE FIELD EQUATIONS OF MOTION

When deriving the gauge field SM equations in Eqs. (3.1) and (3.2), all contributions from $d = 4$ operators in $\delta\mathcal{L}_{d=4}$ effective Lagrangians have been neglected, on the assumption that their coefficients are small, typically $a_i < 1$, $i = 1\ldots4$, with typically $a_i \approx 1/(16\pi^2)$. This allows one to trade flavor-conserving currents for gauge terms with derivatives of the gauge field strengths.

Otherwise, for $a_i \sim 1$, taking into account that the gauge sector is already severely modified, and thus keeping only the flavor-changing contributions in $\delta\mathcal{L}^f_{d=4}$, Eqs. (3.1) and (3.2) would get modified to

$$(D^\mu W_{\mu\nu})_j = +i\frac{g}{4}v^2\mathrm{Tr}[\mathbf{V}_\nu\sigma_j] + \frac{g}{2}\bar{Q}_L\gamma_\nu\sigma_j Q_L - \frac{g}{2}\bar{Q}_L\gamma_\nu\lambda_F[(a_2 - a_3)\delta_{jk} - a_4\epsilon_{3jk}]\sigma_k Q_L, \tag{B1}$$

$$\partial^\mu B_{\mu\nu} = -i\frac{g'}{4}v^2\mathrm{Tr}[\mathbf{T}\mathbf{V}_\nu] + g'\bar{Q}_L\gamma_\nu H^L_q Q_L + g'\bar{Q}_L\gamma_\nu H^R_q Q_L - g'\bar{Q}_L\gamma_\nu\lambda_F \times\left[a_1\mathbb{1} + (a_2 + a_3)\frac{\sigma_3}{2}\right]Q_L, \tag{B2}$$

where the right-handed flavor-changing contributions have been also disregarded. However, as gauge-h coefficients are severely constrained by EW precision data (barring extremely fine-tuned regions in the parameter space) the analysis of flavor-changing couplings would get modified only by terms of $O(a_i \times a_{GH})$, $i = 1\ldots4$, being a_{GH} the coefficients in the gauge-h sector, and therefore their impact on the flavor sector is negligible.

APPENDIX C: LINEAR SIBLINGS OF THE $d = 5$ OPERATORS

In this Appendix we connect the operators listed in Eqs. (3.9), (3.10), and (3.11) with those defined in the linear realization, $\mathcal{X}_i \leftrightarrow \sum_j \mathcal{C}_{ij}\mathcal{X}_{Hj}$, where $\mathcal{C}$ is an 18×18 matrix.

The first set of nonlinear operators listed in Eq. (3.9) corresponds to the following eight linear operators containing fermions, the Higgs doublet H, the rank-2 antisymmetric tensor $\sigma^{\mu\nu}$ and the field strengths $B_{\mu\nu}$, $W_{\mu\nu}$ and $G_{\mu\nu}$:

$$\begin{aligned}
&\mathcal{X}_{H1} = g'\bar{Q}_L\sigma^{\mu\nu}HD_RB_{\mu\nu}, \qquad \mathcal{X}_{H2} = g'\bar{Q}_L\sigma^{\mu\nu}\tilde{H}U_RB_{\mu\nu}, \qquad \mathcal{X}_{H3} = g\bar{Q}_L\sigma^{\mu\nu}W_{\mu\nu}HD_R,\\
&\mathcal{X}_{H4} = g\bar{Q}_L\sigma^{\mu\nu}W_{\mu\nu}\tilde{H}U_R, \qquad \mathcal{X}_{H5} = g_s\bar{Q}_L\sigma^{\mu\nu}HD_RG_{\mu\nu}, \qquad \mathcal{X}_{H6} = g_s\bar{Q}_L\sigma^{\mu\nu}\tilde{H}U_RG_{\mu\nu},\\
&\mathcal{X}_{H7} = g\bar{Q}_L\sigma^{\mu\nu}\sigma_iHD_RH^\dagger W_{\mu\nu}\sigma^iH, \qquad \mathcal{X}_{H8} = g\bar{Q}_L\sigma^{\mu\nu}\sigma_i\tilde{H}U_RH^\dagger\sigma^iW_{\mu\nu}H.
\end{aligned} \tag{C1}$$

The operators $\mathcal{X}_{H7,H8}$ have mass dimension $d = 8$, while all the others have (linear) mass dimension $d = 6$. The correspondence among these linear operators and those nonlinear listed in Eq. (3.9) is the following: for $i = 1,\ldots,8$,

$$\mathcal{X}_i \leftrightarrow \sum_{j=1}^{8}\mathcal{C}_{ij}\mathcal{X}_{Hj} \quad \text{with} \quad \mathcal{C} = \frac{\sqrt{2}}{f}\begin{pmatrix}
1 & 1 & 0 & 0 & 0 & 0 & 0 & 0\\
-1 & 1 & 0 & 0 & 0 & 0 & 0 & 0\\
0 & 0 & 1 & 1 & 0 & 0 & 0 & 0\\
0 & 0 & -1 & 1 & 0 & 0 & 0 & 0\\
0 & 0 & 0 & 0 & 1 & 1 & 0 & 0\\
0 & 0 & 0 & 0 & -1 & 1 & 0 & 0\\
0 & 0 & 1 & -1 & 0 & 0 & -4/f^2 & 4/f^2\\
0 & 0 & -1 & -1 & 0 & 0 & 4/f^2 & 4/f^2
\end{pmatrix}. \tag{C2}$$

The second set of nonlinear operators listed in Eq. (3.10) corresponds to the following four linear operators containing fermions, the Higgs doublet H and the rank-2 antisymmetric tensor $\sigma^{\mu\nu}$:

$$\begin{aligned}&\mathcal{X}_{H9}=\bar{Q}_L\sigma^{\mu\nu}HD_R((D_\mu H)^\dagger D_\nu H-(\mu\leftrightarrow\nu)), && \mathcal{X}_{H10}=\bar{Q}_L\sigma^{\mu\nu}\tilde{H}U_R((D_\mu H)^\dagger D_\nu H-(\mu\leftrightarrow\nu)),\\ &\mathcal{X}_{H11}=\bar{Q}_L\sigma_i\sigma^{\mu\nu}HD_R((D_\mu H)^\dagger\sigma^i D_\nu H-(\mu\leftrightarrow\nu)), && \mathcal{X}_{H12}=\bar{Q}_L\sigma_i\sigma^{\mu\nu}\tilde{H}U_R((D_\mu H)^\dagger\sigma^i D_\nu H-(\mu\leftrightarrow\nu)),\end{aligned}\tag{C3}$$

all of them of mass dimension $d=8$. The correspondence among these linear operators and those nonlinear listed in Eq. (3.10) is the following: for $i=9,\ldots,12$,

$$\mathcal{X}_i\leftrightarrow\sum_{j=9}^{12}\mathcal{C}_{ij}\mathcal{X}_{Hj}\quad\text{with}\quad\mathcal{C}=\frac{2\sqrt{2}}{f^3}\begin{pmatrix}0&0&1&1\\0&0&-1&1\\1&-1&0&0\\-1&-1&0&0\end{pmatrix}.\tag{C4}$$

For the third set in Eq. (3.11), we consider the following six linear operators involving fermions and the Higgs doublet H:

$$\begin{aligned}&\mathcal{X}_{H13}=\bar{Q}_LHD_R(D_\mu H)^\dagger D^\mu H, && \mathcal{X}_{H14}=\bar{Q}_L\tilde{H}U_R(D_\mu H)^\dagger D^\mu H,\\ &\mathcal{X}_{H15}=\bar{Q}_L\sigma_iHD_R(D_\mu H)^\dagger\sigma^i D^\mu H, && \mathcal{X}_{H16}=\bar{Q}_L\sigma_i\tilde{H}U_R(D_\mu H)^\dagger\sigma^i D^\mu H,\\ &\mathcal{X}_{H17}=\bar{Q}_LHD_R(D_\mu H)^\dagger HH^\dagger D^\mu H, && \mathcal{X}_{H18}=\bar{Q}_L\tilde{H}U_R(D_\mu H)^\dagger HH^\dagger D^\mu H,\end{aligned}\tag{C5}$$

where the first four operators have mass dimension $d=8$, while the last two have mass dimension $d=10$. It is then possible to establish the following correspondence between these linear operators and those nonlinear listed in Eq. (3.11): for $i=13,\ldots,18$,

$$\mathcal{X}_i\leftrightarrow\sum_{j=13}^{18}\mathcal{C}_{ij}\mathcal{X}_{Hj}\quad\text{with}\quad\mathcal{C}=\frac{2\sqrt{2}}{f^3}\begin{pmatrix}-1&-1&0&0&0&0\\1&-1&0&0&0&0\\0&0&1&1&0&0\\0&0&-1&1&0&0\\2&2&1&-1&-8/f^2&-8/f^2\\-2&2&-1&-1&8/f^2&-8/f^2\end{pmatrix}.\tag{C6}$$

APPENDIX D: $d = 5$ OPERATOR COEFFICIENTS IN THE UNITARY BASIS

In this Appendix, we report the relations between the coefficients appearing in the Lagrangian Eq. (3.16) and the ones defined in Eq. (3.13) for the effective Lagrangian in the unitary basis:

$$\begin{pmatrix}c_W^u\\c_W^d\\c_{WZ}^+\\c_{WZ}^-\\d_F^u\\d_F^d\\d_Z^u\\d_Z^d\\d_W^+\\d_W^-\\d_G^u\\d_G^d\end{pmatrix}=\mathcal{A}\begin{pmatrix}b_1\\\cdots\\b_{12}\end{pmatrix},\qquad\begin{pmatrix}b_Z^u\\b_Z^d\\b_W^u\\b_W^d\\b_{WZ}^+\\b_{WZ}^-\end{pmatrix}=\mathcal{B}\begin{pmatrix}b_{13}\\\cdots\\b_{18}\end{pmatrix},\tag{D1}$$

$$\mathcal{A} = \begin{pmatrix} 0 & 0 & 2i & 2i & 0 & 0 & 2i & 2i & -1 & -1 & 1 & 1 \\ 0 & 0 & -2i & 2i & 0 & 0 & 2i & -2i & 1 & -1 & -1 & 1 \\ 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & 4 & -4 & 0 & 0 \\ 0 & 0 & 0 & 0 & 0 & 0 & 0 & 0 & -4 & -4 & 0 & 0 \\ 1 & 1 & 1 & 1 & 0 & 0 & 1 & 1 & 0 & 0 & 0 & 0 \\ 1 & -1 & -1 & 1 & 0 & 0 & 1 & -1 & 0 & 0 & 0 & 0 \\ -2s_W^2 & -2s_W^2 & 2c_W^2 & 2c_W^2 & 0 & 0 & 2c_W^2 & 2c_W^2 & 0 & 0 & 0 & 0 \\ -2s_W^2 & 2s_W^2 & -2c_W^2 & 2c_W^2 & 0 & 0 & 2c_W^2 & -2c_W^2 & 0 & 0 & 0 & 0 \\ 0 & 0 & 2 & -2 & 0 & 0 & 2 & -2 & 0 & 0 & 0 & 0 \\ 0 & 0 & 2 & 2 & 0 & 0 & -2 & -2 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 1 & 1 & 0 & 0 & 0 & 0 & 0 & 0 \\ 0 & 0 & 0 & 0 & 1 & -1 & 0 & 0 & 0 & 0 & 0 & 0 \end{pmatrix}, \tag{D2}$$

$$\mathcal{B} = \begin{pmatrix} -1 & -1 & -1 & -1 & -1 & -1 \\ -1 & 1 & 1 & -1 & -1 & 1 \\ -1 & -1 & 1 & 1 & 1 & 1 \\ -1 & 1 & -1 & 1 & 1 & -1 \\ 0 & 0 & -2 & 2 & -2 & 2 \\ 0 & 0 & -2 & -2 & 2 & 2 \end{pmatrix}. \tag{D3}$$

[1] G. Aad *et al.* (ATLAS Collaboration), Phys. Lett. B **716**, 1 (2012).
[2] S. Chatrchyan *et al.* (CMS Collaboration), Phys. Lett. B **716**, 30 (2012).
[3] F. Englert and R. Brout, Phys. Rev. Lett. **13**, 321 (1964).
[4] P. W. Higgs, Phys. Lett. **12**, 132 (1964).
[5] P. W. Higgs, Phys. Rev. Lett. **13**, 508 (1964).
[6] O. Eberhardt, G. Herbert, H. Lacker, A. Lenz, A. Menzel, U. Nierste, and M. Wiebusch, Phys. Rev. Lett. **109**, 241802 (2012).
[7] M. Baak, M. Goebel, J. Haller, A. Hoecker, D. Kennedy, R. Kogler, K. Mönig, M. Schott, and J. Stelzer, Eur. Phys. J. C **72**, 2205 (2012).
[8] L. Susskind, Phys. Rev. D **20**, 2619 (1979).
[9] S. Dimopoulos and L. Susskind, Nucl. Phys. **B155**, 237 (1979).
[10] S. Dimopoulos and J. Preskill, Nucl. Phys. **B199**, 206 (1982).
[11] A. Manohar and H. Georgi, Nucl. Phys. **B234**, 189 (1984).
[12] D. B. Kaplan and H. Georgi, Phys. Lett. **136B**, 183 (1984).
[13] D. B. Kaplan, H. Georgi, and S. Dimopoulos, Phys. Lett. **136B**, 187 (1984).
[14] T. Banks, Nucl. Phys. **B243**, 125 (1984).
[15] H. Georgi, D. B. Kaplan, and P. Galison, Phys. Lett. **143B**, 152 (1984).
[16] H. Georgi and D. B. Kaplan, Phys. Lett. **145B**, 216 (1984).
[17] M. J. Dugan, H. Georgi, and D. B. Kaplan, Nucl. Phys. **B254**, 299 (1985).
[18] R. Contino, arXiv:1005.4269.
[19] R. S. Chivukula and H. Georgi, Phys. Lett. B **188**, 99 (1987).
[20] L. J. Hall and L. Randall, Phys. Rev. Lett. **65**, 2939 (1990).
[21] G. D'Ambrosio, G. F. Giudice, G. Isidori, and A. Strumia, Nucl. Phys. **B645**, 155 (2002).
[22] Z. Lalak, S. Pokorski, and G. G. Ross, J. High Energy Phys. 08 (2010) 129.
[23] A. L. Fitzpatrick, G. Perez, and L. Randall, Phys. Rev. Lett. **100**, 171604 (2008).
[24] B. Grinstein, M. Redi, and G. Villadoro, J. High Energy Phys. 11 (2010) 067.
[25] A. J. Buras, M. V. Carlucci, L. Merlo, and E. Stamou, J. High Energy Phys. 03 (2012) 088.
[26] G. Isidori, Y. Nir, and G. Perez, Annu. Rev. Nucl. Part. Sci. **60**, 355 (2010).
[27] V. Cirigliano, B. Grinstein, G. Isidori, and M. B. Wise, Nucl. Phys. **B728**, 121 (2005).
[28] S. Davidson and F. Palorini, Phys. Lett. B **642**, 72 (2006).
[29] A. L. Kagan, G. Perez, T. Volansky, and J. Zupan, Phys. Rev. D **80**, 076002 (2009).
[30] M. Gavela, T. Hambye, D. Hernandez, and P. Hernandez, J. High Energy Phys. 09 (2009) 038.
[31] T. Feldmann, M. Jung, and T. Mannel, Phys. Rev. D **80**, 033003 (2009).

[32] R. Alonso, M. B. Gavela, L. Merlo, and S. Rigolin, J. High Energy Phys. 07 (2011) 012.
[33] R. Alonso, G. Isidori, L. Merlo, L. A. Munoz, and E. Nardi, J. High Energy Phys. 06 (2011) 037.
[34] R. Alonso, M. Gavela, D. Hernandez, and L. Merlo, Phys. Lett. B **715**, 194 (2012).
[35] R. Alonso, M. Gavela, L. Merlo, S. Rigolin, and J. Yepes, J. High Energy Phys. 06 (2012) 076.
[36] R. Alonso, M. Gavela, L. Merlo, S. Rigolin, and J. Yepes, arXiv:1212.3305.
[37] D. B. Kaplan, Nucl. Phys. **B365**, 259 (1991).
[38] M. Redi and A. Weiler, J. High Energy Phys. 11 (2011) 108.
[39] G. F. Giudice, C. Grojean, A. Pomarol, and R. Rattazzi, J. High Energy Phys. 06 (2007) 045.
[40] R. Contino, C. Grojean, M. Moretti, F. Piccinini, and R. Rattazzi, J. High Energy Phys. 05 (2010) 089.
[41] A. Azatov, R. Contino, and J. Galloway, J. High Energy Phys. 04 (2012) 127.
[42] T. Appelquist and C. W. Bernard, Phys. Rev. D **22**, 200 (1980).
[43] A. C. Longhitano, Phys. Rev. D **22**, 1166 (1980).
[44] A. C. Longhitano, Nucl. Phys. **B188**, 118 (1981).
[45] F. Feruglio, Int. J. Mod. Phys. A **08**, 4937 (1993).
[46] T. Appelquist and G.-H. Wu, Phys. Rev. D **48**, 3235 (1993).
[47] A. Azatov and J. Galloway, Int. J. Mod. Phys. A **28**, 1330004 (2013).
[48] K. Agashe, R. Contino, and A. Pomarol, Nucl. Phys. **B719**, 165 (2005).
[49] R. Contino, L. Da Rold, and A. Pomarol, Phys. Rev. D **75**, 055014 (2007).
[50] B. Gripaios, A. Pomarol, F. Riva, and J. Serra, J. High Energy Phys. 04 (2009) 070.
[51] E. Halyo, Mod. Phys. Lett. A **08**, 275 (1993).
[52] W. D. Goldberger, B. Grinstein, and W. Skiba, Phys. Rev. Lett. **100**, 111802 (2008).
[53] L. Vecchi, Phys. Rev. D **82**, 076009 (2010).
[54] B. A. Campbell, J. Ellis, and K. A. Olive, J. High Energy Phys. 03 (2012) 026.
[55] S. Matsuzaki and K. Yamawaki, Phys. Lett. B **719**, 378 (2013).
[56] Z. Chacko, R. Franceschini, and R. K. Mishra, arXiv:1209.3259.
[57] B. Bellazzini, C. Csaki, J. Hubisz, J. Serra, and J. Terning, arXiv:1209.3299 [Eur. Phys. J. C (to be published)].
[58] G. Panico, M. Redi, A. Tesi, and A. Wulzer, arXiv:1210.7114.
[59] T. Appelquist, M. J. Bowick, E. Cohler, and A. I. Hauser, Phys. Rev. D **31**, 1676 (1985).
[60] G. Cvetič and R. Kogerler, Nucl. Phys. **B328**, 342 (1989).
[61] D. Espriu and J. Manzano, Phys. Rev. D **63**, 073008 (2001).
[62] J. Laiho, E. Lunghi, and R. S. Van de Water, Phys. Rev. D **81**, 034503 (2010). Updates available on http://latticeaverages.org/.
[63] J. Brod and M. Gorbahn, Phys. Rev. Lett. **108**, 121801 (2012).
[64] R. Aaij *et al.* (LHCb Collaboration), Phys. Rev. Lett. **108**, 101803 (2012).
[65] K. Nakamura *et al.* (Particle Data Group), J. Phys. G **37**, 075021 (2010).
[66] P. Gambino and M. Misiak, Nucl. Phys. **B611**, 338 (2001).
[67] M. Misiak *et al.*, Phys. Rev. Lett. **98**, 022002 (2007).
[68] M. Misiak and M. Steinhauser, Nucl. Phys. **B764**, 62 (2007).
[69] G. Buchalla, A. J. Buras, and M. E. Lautenbacher, Rev. Mod. Phys. **68**, 1125 (1996).
[70] T. Inami and C. S. Lim, Prog. Theor. Phys. **65**, 297 (1981).
[71] A. J. Buras, L. Merlo, and E. Stamou, J. High Energy Phys. 08 (2011) 124.