A LIGHT DYNAMICAL SCALAR BOSON

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With the discovery of a scalar resonance at ATLAS and CMS, the understanding of the electroweak symmetry breaking origin seems a much closer goal. A strong dynamics at relatively low scales is still a good candidate. In this talk, the complete effective Lagrangian up to $d \leq 5$ will be presented, both for the gauge and the flavour sectors. Interesting features in the flavour phenomenology will be discussed.

1 Framework

With the new resonance at the Electroweak (EW) scale discovered at LHC [1], we can now hope to have hints to understand the origin of the EW symmetry breaking (EWSB) mechanism. The data indicate that the new particle is looking more and more like the Standard Model (SM) scalar boson [2] with mass around 125 GeV, but other possibilities are still viable. In particular, the case of a strong dynamics at the TeV scale Λ_s responsible for the EWSB is still attractive. In the original technicolor ansatz [3], only the three SM would-be-Goldstone bosons (GBs) are retained and are responsible of giving mass to the weak gauge bosons. Given the present discovery, a more attractive scenario is the so-called composite Higgs model, originally introduced in Ref. [4]. In this context, the theory is based on the spontaneous breaking of a large global symmetry, that gives rise to the appearance of several GBs: a realistic model then accounts for three GBs corresponding to the SM ones, and for at least one more GB that takes the role of the scalar boson. The latter is then a composite scalar degree of freedom that arises as massless and gets mass due to an explicit breaking of the symmetry. With respect to the technicolor context, here there are four relevant scales: Λ_s typical of the strong resonances; $f \leq 4\pi\Lambda_s$ that characterises the GBs energy scale; $v = 246$ GeV defined through the W mass, $M_W = gv/2$; and $\langle h \rangle$ that is the vacuum expectation value (VEV) of the scalar particle, providing the EWSB, and that is in general distinct from v. To measure the degree of non-linearity of these schemes, it is customary to introduce the parameter $\xi \equiv (v/f)^2$ that parametrises the impact of the strong dynamics at low-energy.

Without entering into details of a specific model, it is possible to describe the NP effects due to the TeV strong dynamics by making use of an effective Lagrangian approach, dealing with only the SM fields. The SM GBs can be described by a dimensionless unitary matrix:

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

with L, R denoting respectively the $SU(2)_{L,R}$ global transformations of the scalar potential. The adimensionality of $\mathbf{U}(x)$ is the technical key to understand why the dimension of the leading low-energy operators describing the dynamics of the scalar sector differs for a non-linear regime [5] and a purely linear regime [6,7]. In the former, non-renormalisable operators containing extra powers of a light h are weighted by powers of h/f [8], and the GB contributions encoded in $\mathbf{U}(x)$ do not exhibit any scale suppression. In the linear regime, instead, the light h and the three SM GBs are encoded into the scalar doublet H, with mass dimension one: therefore any extra insertion of H is suppressed by a power of the cutoff.

It is becoming customary to parametrise the Lagrangian describing a light dynamical scalar particle h by means of the following ansatz [9] :

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

where dots stand for higher order terms in the (linear) expansion in h/f, 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 $SU(2)_L$. The covariant derivative reads $\mathbf{D}_\mu\mathbf{U}(x) \equiv \partial_\mu \mathbf{U}(x) + ig/2\, W^a_\mu(x)\sigma_a \mathbf{U}(x) - ig'/2\, B_\mu(x)\mathbf{U}(x)\sigma_3$ with W^a_μ (B_μ) denoting the $SU(2)_L$ ($U(1)_Y$) gauge bosons and g (g') the corresponding gauge coupling. In the equations above, $V(h)$ denotes the effective scalar potential describing the breaking of the EW symmetry. The first line in Eq. (1) includes the SMS kinetic term, its scalar potential and the Yukawa-like interactions for quarks, while the second line describes the W and Z masses and their interactions with h, as well as the usual custodial symmetry breaking term labeled by c_T.

The functions $\mathcal{F}_H(h)$, $\mathcal{F}_C(h)$, $\mathcal{F}_T(h)$ and $\mathcal{F}_Y(h)$ above, as well as all $\mathcal{F}(h)$ functions to be used below, encode the generic dependence on $(\langle h\rangle + h)$ and are model-dependent. Each $\mathcal{F}(h)$ function can be expanded in powers of ξ, $\mathcal{F}(h) = g_0(h,v) + \xi g_1(h,v) + \xi^2 g_2(h,v) + \ldots$, where $g(h,v)$ are model-dependent functions of h and of v, once $\langle h\rangle$ is expressed in terms of ξ and v. For not too small ξ the whole series may need to be considered.

The above Lagrangian can be very useful to describe an extended class of "Higgs" models, ranging from the SM scenario with a linear Higgs sector (for $\langle h\rangle = v$, $a = b = c = 1$ and neglecting the higher order terms in h), to the technicolor-like 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/holographic Higgs models [4,8,10] up to dilaton-like scalar frameworks. Note that in concrete models electroweak corrections imply $\xi < 0.2 - 0.4$ [11], but we will leave the ξ parameter free here and account for the constraints on custodial symmetry through limits on the $d = 2$ and higher-dimensional chiral operator coefficients.

In what follows, the complete basis of independent operators up to dimension 5 will be reported, both in the gauge and in the flavour sectors, providing the complete list of interactions of a light h [12,13,14]. This analysis enlarges and completes the operator basis previously considered in Refs. [5,9] and represents a fundamental tool in order to characterise the emerging phenomenology at LHC and investigating on the EWSB origin.

2 The effective Lagrangian in the gauge sector

All CP-even gauge operators appropriate to the non-linear regime will be included in this section, up to mass dimension 5. In the absence of a light h, no pure gauge or gauge-h $d = 5$ operator exists, and it is thus a good guideline to start from the basis of $d = 4$ pure gauge chiral operators and complete it up to $d = 5$ with suitable insertions of h. The connection to the linear regime will be made manifest exploiting the operator dependence on ξ. The Lagrangian can be decomposed as

$$
\begin{aligned}
\mathscr{L}^{d\leq 5}_{gauge-h} =& \mathscr{L}_h - \frac{g_s^2}{4}\, G^a_{\mu\nu}G^{\mu\nu}_a\, \mathcal{F}_G(h) - \frac{g^2}{4}\, W^a_{\mu\nu}W^{\mu\nu}_a\, \mathcal{F}_W(h) - \frac{g'^2}{4}\, B_{\mu\nu}B^{\mu\nu}\, \mathcal{F}_B(h) + \\
&+ \xi \sum_{i=1}^{5} c_i\, \mathcal{P}_i(h) \,+\, \xi^2 \sum_{i=6}^{20} c_i\, \mathcal{P}_i(h) + \xi^3 \sum_{i=21}^{23} c_i\, \mathcal{P}_i(h) + \xi^4\, c_{24}\, \mathcal{P}_{24}(h)\,.
\end{aligned}
\tag{2}
$$

The first line of Eq. (2) contains the kinetic terms for the gauge bosons, with $W_{\mu\nu}$, $B_{\mu\nu}$ and $G_{\mu\nu}$ denoting the $SU(2)_L$, $U(1)_Y$ and $SU(3)_C$ field strengths, respectively. The second line of Eq. (2) contains the following 24 CP-even operators, ordered by their ξ dependence [13]:

$$
\begin{aligned}
&\mathcal{P}_1(h) = g\, g'\, B_{\mu\nu} \mathrm{Tr}\left(\mathbf{T}\, W^{\mu\nu}\right)\, \mathcal{F}_1(h) &&\mathcal{P}_4(h) = i\, g' B_{\mu\nu} \mathrm{Tr}(\mathbf{T}\mathbf{V}^\mu)\, \partial^\nu \mathcal{F}_4(h)\\
&\mathcal{P}_2(h) = i\, g'\, B_{\mu\nu} \mathrm{Tr}\left(\mathbf{T}\,[\mathbf{V}^\mu, \mathbf{V}^\nu]\right)\, \mathcal{F}_2(h) &&\mathcal{P}_5(h) = i\, g\, \mathrm{Tr}(W_{\mu\nu}\mathbf{V}^\mu)\, \partial^\nu \mathcal{F}_5(h)\\
&\mathcal{P}_3(h) = i\, g\, \mathrm{Tr}\left(W_{\mu\nu}\,[\mathbf{V}^\mu, \mathbf{V}^\nu]\right)\, \mathcal{F}_3(h) &&
\end{aligned}
\tag{3}
$$

$$
\begin{aligned}
&\mathcal{P}_6(h) = \left(\mathrm{Tr}\left(\mathbf{V}_\mu\, \mathbf{V}^\mu\right)\right)^2\, \mathcal{F}_6(h) &&\mathcal{P}_{14}(h) = i\, g\, \mathrm{Tr}(\mathbf{T}W_{\mu\nu})\mathrm{Tr}(\mathbf{T}\mathbf{V}^\mu)\, \partial^\nu \mathcal{F}_{14}(h)\\
&\mathcal{P}_7(h) = \left(\mathrm{Tr}\left(\mathbf{V}_\mu\, \mathbf{V}_\nu\right)\right)^2\, \mathcal{F}_7(h) &&\mathcal{P}_{15}(h) = \mathrm{Tr}(\mathbf{T}\,[\mathbf{V}_\mu, \mathbf{V}_\nu])\mathrm{Tr}(\mathbf{T}\mathbf{V}^\mu)\, \partial^\nu \mathcal{F}_{15}(h)\\
&\mathcal{P}_8(h) = g^2\, \left(\mathrm{Tr}\left(\mathbf{T}\, W^{\mu\nu}\right)\right)^2\, \mathcal{F}_8(h) &&\mathcal{P}_{16}(h) = \mathrm{Tr}(\mathbf{V}_\nu\, \mathcal{D}_\mu \mathbf{V}^\mu)\, \partial^\nu \mathcal{F}_{16}(h)\\
&\mathcal{P}_9(h) = i\, g\, \mathrm{Tr}\left(\mathbf{T}\, W_{\mu\nu}\right) \mathrm{Tr}\left(\mathbf{T}\,[\mathbf{V}^\mu, \mathbf{V}^\nu]\right)\, \mathcal{F}_9(h) &&\mathcal{P}_{17}(h) = \mathrm{Tr}(\mathbf{T}\, \mathcal{D}_\mu \mathbf{V}^\mu)\mathrm{Tr}(\mathbf{T}\mathbf{V}_\nu)\, \partial^\nu \mathcal{F}_{17}(h)\\
&\mathcal{P}_{10}(h) = g\, \epsilon^{\mu\nu\rho\lambda} \mathrm{Tr}\left(\mathbf{T}\mathbf{V}_\mu\right) \mathrm{Tr}\left(\mathbf{V}_\nu\, W_{\rho\lambda}\right)\, \mathcal{F}_{10}(h) &&\mathcal{P}_{18}(h) = \mathrm{Tr}\left(\mathbf{V}_\mu\, \mathbf{V}^\mu\right)\, \partial_\nu \partial^\nu \mathcal{F}_{18}(h)\\
&\mathcal{P}_{11}(h) = \mathrm{Tr}\left((\mathcal{D}_\mu \mathbf{V}^\mu)^2\right)\, \mathcal{F}_{11}(h) &&\mathcal{P}_{19}(h) = \mathrm{Tr}\left(\mathbf{V}_\mu\, \mathbf{V}_\nu\right)\, \partial^\mu \mathcal{F}_{19}(h) \partial^\nu \mathcal{F}'_{19}(h)\\
&\mathcal{P}_{12}(h) = \mathrm{Tr}(\mathbf{T}\, \mathcal{D}_\mu \mathbf{V}^\mu)\, \mathrm{Tr}(\mathbf{T}\, \mathcal{D}_\nu \mathbf{V}^\nu)\, \mathcal{F}_{12}(h) &&\mathcal{P}_{20}(h) = \mathrm{Tr}\left(\mathbf{T}\mathbf{V}_\mu\right) \mathrm{Tr}\left(\mathbf{T}\mathbf{V}_\nu\right)\, \partial^\mu \mathcal{F}_{20}(h) \partial^\nu \mathcal{F}'_{20}(h)\\
&\mathcal{P}_{13}(h) = \mathrm{Tr}([\mathbf{T}\,, \mathbf{V}_\nu]\, \mathcal{D}_\mu \mathbf{V}^\mu)\, \mathrm{Tr}(\mathbf{T}\mathbf{V}^\nu) \mathcal{F}_{13}(h) &&
\end{aligned}
\tag{4}
$$

$$
\begin{aligned}
&\mathcal{P}_{21}(h) = \mathrm{Tr}\left(\mathbf{V}_\mu \mathbf{V}^\mu\right) \left(\mathrm{Tr}\left(\mathbf{T}\mathbf{V}_\nu\right)\right)^2 \mathcal{F}_{21}(h) &&\mathcal{P}_{23}(h) = \left(\mathrm{Tr}\left(\mathbf{T}\, \mathbf{V}_\mu\right)\right)^2 \partial_\nu \partial^\nu \mathcal{F}_{23}(h)\\
&\mathcal{P}_{22}(h) = \mathrm{Tr}\left(\mathbf{V}_\mu \mathbf{V}_\nu\right) \mathrm{Tr}\left(\mathbf{T}\mathbf{V}^\mu\right) \mathrm{Tr}\left(\mathbf{T}\mathbf{V}^\nu\right) \mathcal{F}_{22}(h) &&
\end{aligned}
\tag{5}
$$

$$
\mathcal{P}_{24}(h) = \left(\mathrm{Tr}\left(\mathbf{T}\mathbf{V}_\mu\right) \mathrm{Tr}\left(\mathbf{T}\mathbf{V}_\nu\right)\right)^2\, \mathcal{F}_{24}(h)\,.
\tag{6}
$$

The 24 constant parameters c_i are model-dependent coefficients. The powers of ξ, factorized out in the second line of Eq. (2), do not reflect an expansion in ξ, but a reparametrisation that facilitates the tracking to the lowest dimension at which a "sibling" operator appears in the linear expansion. By sibling we mean an operator written in terms of the scalar doublet H, that includes the pure gauge part of the couplings $\mathcal{P}_{1-24}(h)$. It may happen that an operator listed in Eqs. (3)-(6) corresponds to a specific combination of siblings with different dimensions. This is the case, for instance, of $\mathcal{P}_{13}(h)$, whose siblings are of dimension 8 and 10.

For $\xi \ll 1$ the weight of the operators which are accompanied by powers of ξ is scale suppressed compared to that of SM renormalisable couplings. In this limit the Lagrangian above would encode a consistent linear expansion up to $d = 6$ operators, if only the terms of zero and first order in ξ are kept: indeed, operators $\mathcal{P}_6(h)$ to $\mathcal{P}_{24}(h)$ would correspond to $d = 8$ or higher-dimension siblings in the linear expansion. In contrast, in the non-linear regime, that is for $\xi \approx 1$, no such suppression appears and *all* operators in Eqs. (3)-(6) include $d \leq 5$ couplings and should be considered on equal footing. The leading terms of the linear and non-linear expansions do not match.

The different operators defined in Eqs. (3)-(6) correspond to three major categories: pure gauge and gauge-h operators (in blue) which result from a direct extension of the original Appelquist-Longhitano chiral Higgsless basis; operators containing the contraction $\mathcal{D}_\mu \mathbf{V}^\mu$ and no derivatives of $\mathcal{F}(h)$ (in green); operators with one or two derivatives of $\mathcal{F}(h)$ (in red).

3 The effective Lagrangian in the flavour sector

The core of the flavour 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 flavour changing neutral current (FCNC) observability tests. Minimal Flavour Violation (MFV) [15,16] emerged in the last years as one of the most promising working frameworks to this end: the MFV ansatz dictates that flavour in the SM and beyond is described at low-energies uniquely in terms of the known fermion mass hierarchies and mixings. An outcome is that the energy scale of the NP may be as low as few TeV in several distinct contexts [17], while in general it should be larger than hundreds of TeV.

In Ref. [16], the complete basis of gauge-invariant 6-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 cut-off scale.

In the non-linear 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. A total of four independent $d = 4$ chiral operators containing LH fermion fields can be constructed [12,18], namely:

$$\mathscr{L}^f_{\chi=4} = \xi \sum_{i=1,2,3} \hat{a}_i\, \mathcal{O}_i(h) + \xi^2 \hat{a}_4\, \mathcal{O}_4(h) \tag{7}$$

$$\begin{aligned}
&\mathcal{O}_1(h) = \frac{i}{2}\,\bar{Q}_L\,\lambda_F\,\gamma^\mu\,\{\mathbf{T},\mathbf{V}_\mu\}\,Q_L\,\mathcal{F}_1(h)\,, &&\mathcal{O}_2(h) = i\,\bar{Q}_L\,\lambda_F\,\gamma^\mu\,\mathbf{V}_\mu\,Q_L\,\mathcal{F}_2(h)\,,\\
&\mathcal{O}_3(h) = i\,\bar{Q}_L\,\lambda_F\,\gamma^\mu\,\mathbf{T}\,\mathbf{V}_\mu\,\mathbf{T}\,Q_L\,\mathcal{F}_3(h)\,, &&\mathcal{O}_4(h) = \frac{1}{2}\,\bar{Q}_L\,\lambda_F\,\gamma^\mu\,[\mathbf{T},\mathbf{V}_\mu]\,Q_L\,\mathcal{F}_4(h)\,,
\end{aligned} \tag{8}$$

where the parameter λ_F remembers the MFV ansatz,

$$\lambda_F \equiv Y_U\,Y_U^\dagger + Y_D\,Y_D^\dagger = V^\dagger \mathbf{y}_U^2 V + \mathbf{y}_D^2. \tag{9}$$

Out of these $\mathcal{O}_1(h) - \mathcal{O}_3(h)$ are CP-even while $\mathcal{O}_4(h)$ is intrinsically CP-odd [12]. The powers of ξ in Eq. (7) facilitate the identification of the lowest dimension at which a sibling operator appears in the linear regime. The lowest-dimension siblings of $\mathcal{O}_1(h) - \mathcal{O}_3(h)$ arise at $d = 6$, while that of $\mathcal{O}_4$ appears at $d = 8$ [12].

Operators $\mathcal{O}_1(h) - \mathcal{O}_3(h)$ induce tree-level contributions to $\Delta F = 1$ processes mediated by the Z boson and are severely constrained. Due to the MFV structure of the coefficients, sizable flavour-changing effects may only be expected in the down quark sectors, with data on K and B transitions providing the strongest constraints on a_Z^d,

$$-0.044 < a_Z^d < 0.009 \qquad\qquad \text{at } 95\% \text{ of C.L.} \tag{10}$$

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

Furthermore, operators $\mathcal{O}_2(h) - \mathcal{O}_4(h)$ induce corrections to the fermion-W couplings, and thus to the the CKM matrix. This in turn induces modifications [12] on the strength of meson oscillations (at loop level), on $B^+ \to \tau^+\nu$ decay and on the B semileptonic CP-asymmetry, among others; more specifically the following process have been taken into account in Ref. [12]:

- 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 $B_s^0 \to \psi\phi$. Possible large deviations from the values predicted by the SM are only allowed in the K system.

- The ratio among the meson mass differences in the B_d and B_s systems, $R_{\Delta M_B} \equiv \Delta M_{B_d}/\Delta M_{B_s}$. Deviations from the SM prediction for this observable are negligible.

- The ratio among the $B^+ \to \tau^+\nu$ branching ratio and the B_d mass difference, $R_{BR/\Delta M} \equiv BR(B^+ \to \tau^+\nu)/\Delta M_{B_d}$. This observable is clean from theoretical hadronic uncertainties.

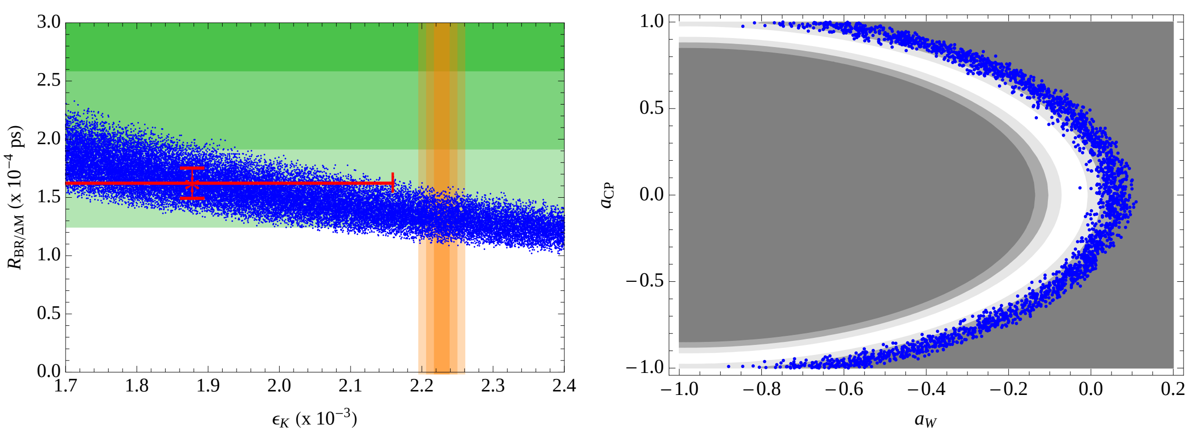
- The $\bar{B} \to X_s\gamma$ branching ratio that benefits of good experimental and theoretical precision.

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 favoured. 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 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 $R_{BR/\Delta M}$ are

$$\epsilon_K = 1.88 \times 10^{-3}\,, \qquad\qquad R_{BR/\Delta M} = 1.62 \times 10^{-4}\,. \tag{11}$$

The errors on these quantities are $\sim 15\%$ and $\sim 8\%$, estimated considering the uncertainties on the input parameters and the analysis performed in Ref. [19]. Fig. 1 shows the correlation between ϵ_K and $R_{BR/\Delta M}$ (left panel) and the $a_{CP} - a_W$ parameter space (right panel), requiring that ϵ_K and $R_{BR/\Delta M}$ lie inside their own 3σ experimental determination. In the latter, the gray areas correspond to the bounds from the $BR(\bar{B} \to X_s\gamma)$. 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.



(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 $a_Z^d \in [-0.044, 0.009]$.

Figure 1: Results for the reference point $(|V_{ub}|, \gamma) = (3.5 \times 10^{-3}, 66°)$. See the text for details.

Fig. 1 on the right shows that a_{CP}, the overall coefficient of the genuinely CP-odd coupling $\mathcal{O}_4(h)$ 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:

$$\delta\mathscr{L}^h_{\chi=4} \supset a_{CP}\left(1 + \beta_{CP}\frac{h}{v}\right)\mathcal{O}_4\ . \tag{12}$$

These are encouraging results in the sense of allowing short-term observability. In a conservative perspective, the operator coefficients of the $d = 4$ non-linear 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 flavour 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.

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