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Higgs Physics, in the SM and Beyond

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Abstract
A lecture on Higgs boson physics to highlight why it is necessary and how it looks like. I review the Standard Model and why a small electroweak scale is our strongest indication for an extended Higgs sector, that can be searched for by a precise study of Higgs properties. To this goal, I discuss effective field theories and how they capture the most relevant effects in large classes of scenarios beyond the Standard Model.

Keywords
Standard model: High-Energy Physics; Higgs Physics.

1 Motivation

1964-1967: A quantum field theoretical (QFT) description of the electroweak (EW) interactions is developed; among a handful of models that can give mass to the $W^\pm$ and Z bosons, one stands out for its predictiveness, simplicity and for seemingly getting as close as possible to a fundamental theory: the Standard Model (SM) of particle physics [1–3]. Its most distinctive prediction is the existence of a resonance, the Higgs boson, whose properties are uniquely fixed by parameters that are already known, except for one, the Higgs boson mass m_h.

2009-2012: the Large Hadron Collider (LHC) is built to collide protons up to 14 TeV energies, and to search explicitly for the Higgs boson, or any alternative source for the EW masses. A discovery is guaranteed by theoretical inconsistency of the EW massive theory above ≈ 3 TeV. And is indeed made on July 4, 2012 [4, 5]. The last parameter of the SM is now measured, $m_h = 125.09$ [6]. With this mass, the Higgs boson properties are just right for a rich experimental program to be carried out, as Higgs decays in a rather equilibrated way to all SM particles. With this mass, we can compute the quantum lifetime of our universe, and find that it is just right to last 13.7 billion years [7].

The most relevant aspect of the Higgs discovery, however, is that it constitutes the last brick of the SM, and this brick is just right, that the theory has in principle a very large range of validity. For this reason, the Higgs discovery interrupts an important trend in particle physics, where well-established fundamental principles were pointing the finger towards guaranteed discoveries.

In these lecture notes I will try to give a feeling of physics before the SM, in the SM, and Beyond the SM, to appreciate the necessity for a Higgs boson, but also its limitations. I think they provide a nice little story to understand why we are interested in studying Higgs physics, and how we are going to do so at the LHC and future colliders. Hundreds of relevant references will be omitted, sorry, and loads of interesting physics will not even be mentioned, sorry - see for instance the complementary reviews Refs. [8–11].

The notes as organized as follows. The QFT of massless and massive spin-1 states, and the necessity for a Higgs mechanism is reviewed in section 2, leading to the SM in section 3. The reader familiar with the SM and interested in the most modern aspects of Higgs physics, relevant for colliders, can skip to section 4, were I discuss the motivations why the SM might not be the end of the story. This section includes a discussion of Effective Field Theories (EFTs) relevant for Higgs physics, ranging from practical aspects relevant for global fits in SM precision tests, to power-counting rules to identify what are the relevant high-energy features that we can test in low-energy experiments.

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2 bSM - before the Standard Model

To appreciate the importance of the Higgs mechanism, and of its SM realization, I must first discuss how the SM looks like without a Higgs boson. In particular I want to discuss the difficulties of providing a QFT description of the massless and massive spin-1 states/resonances observed in Nature: the photon γ and W, Z bosons respectively.

2.1 Gauge Invariance – 4 Legs Good, 2 Legs Better

The observation that physics is invariant under the Poincaré symmetry $\mathcal{P} = \mathbb{R}^4 \ltimes SO(1,3)$ of translations, rotations and boosts, shapes most of our understanding of nature [12]. We realize this symmetry by building objects with well-defined transformation properties and combine them in an invariant way. An important mismatch strikes us at the very start of this program. Physical states $|p\rangle = a^\dagger(\vec{p})|0\rangle$ at finite momentum $\vec{p}$ and spin s, and their scattering amplitudes, transform accordingly to the Little group, the subgroup of Lorentz that leaves the momentum of a particle unchanged (we can think of the momentum as spontaneously breaking the Lorentz group $SO(1,3)$, the Little group is what is left). On the other hand, the fields $\Phi_{\{\mu\}}(x)$ appearing in the Lagrangian in position space, and Feynman diagrams, appear in full representations of $SO(1,3)$ (denoted here with generic indices $\{\mu\}$). The two are however related,

$$\Phi_{\{\mu\}}(x) = \sum_s \int d^3p \, \epsilon_{\{\mu\}}(x; \vec{p}, s) a(\vec{p}, s) \,, \tag{1}$$

with $\epsilon_{\{\mu\}}(x; \vec{p}, s)$ the polarization vectors, that define how representations of Lorentz break into representations of the Little group $LG \subset SO(1,3)$. For $\Lambda^\mu_{\ \nu} \in SO(1,3)$,

$$\Phi_{\{\mu\}}(x) \to D(\Lambda^{-1})\Phi_{\{\mu\}}(\Lambda x) \qquad a(\vec{p}, s) \to \sum_{s'} U^j_{s,s'}(\Lambda^{Little}) a(\Lambda\vec{p}, s), \tag{2}$$

where $D(\Lambda)$ and $U(\Lambda^{Little})$ are representations of the full and Little Lorentz group respectively, that typically differ from one another.

Massive vectors. In this case, the Little group is the group of rotations $SO(3)$. Its representations are well known and classified according to their dimension $2j + 1$, where j is half-integer and refers to the spin of the particle annihilated by $a(\vec{p}, s)$ (and it bounds the sum in Eq. (2) into $s, s' \leq j$). The spin-1 representation, in which we are interested to describe W and Z bosons, is 3-dimensional, corresponding in practice to the two transverse and one longitudinal polarizations of massive vectors.

The smallest full representation of $SO(1,3)$ that can accommodate these 3 states, is the vector field $\Phi_{\{\mu\}} = V^\mu$, with $D(\Lambda) = \Lambda^\mu_{\ \nu}$, that can source 4 degrees of freedom. The remaining one, $4 = 3 \oplus 1$ under $SO(1,3) \to SO(3)$, corresponds to a $j = 1$ state and has a polarization $\epsilon_\mu(p) \propto p_\mu$ in Eq. (1). This is equivalent to the state sourced by the derivative of a scalar $\Phi_{\{\mu\}}(x) = \partial_\mu \phi(x)$ and it does not interest us in the description of spin-1 states. Luckily, it is easy to eliminate it: on the physical states, $p^\mu \epsilon_\mu(p) = 0$ singles out the spin-1 polarizations (for the scalar $p^\mu \epsilon_\mu(p) \propto p^\mu p_\mu = m^2 \neq 0$), which is equivalent to,

$$\langle \psi^{phys'} | \partial_\mu V^\mu | \psi^{phys} \rangle = 0 \tag{3}$$

on physical states, that separates the Hilbert space into two disjoint parts.

The most general Lagrangian (up to dimension-4 and bilinear in V^μ – so as to describe a free field) compatible with the Lorentz transformation of V^μ, can be written as

$$\mathcal{L} = -\frac{1}{4g^2} F_{\mu\nu} F^{\mu\nu} - \frac{\xi}{2}(\partial_\mu V^\mu)^2 - \frac{v^2}{2} V_\mu V^\mu, \tag{4}$$

where we have defined $F_{\mu\nu} \equiv \partial_\mu V_\nu - \partial_\nu V_\mu$ and introduced generic parameters g, v, ξ. Clearly Eq. (4) describes four dynamical degrees of freedom, since it contains the time derivative of all 4 components of V^μ. If we call $\chi = \partial_\mu V^\mu$, the equations of motion for V^μ read

$$-\xi\partial^\nu\chi - v^2 V^\nu = \partial_\mu F^{\mu\nu}/g^2 \qquad \Rightarrow^{\partial^\nu} \qquad -\xi\Box\chi - v^2\chi = 0\,, \tag{5}$$

where in the second equation we have exploited the fact that $F^{\mu\nu}$ is antisymmetric[1] Eq. (5) is interesting because it shows that the fields χ is not sourced by the other components of the vector field, it is a free field. For this reason it is consistent to have it vanish at all times, Eq. (3). In fact, the (Proca) Lagrangian

$$\mathcal{L} = -\frac{1}{4g^2}F_{\mu\nu}F^{\mu\nu} - \frac{v^2}{2}V_\mu V^\mu\,, \tag{6}$$

automatically provides this condition Eq. (3) as a consequence of the equations of motion (Eq. (5) with $\xi = 0$), and can be thought as the correct Lagrangian to describe the dynamics of one massive vector with $m = vg$.

Massless vectors. For massless particles the little group is $ISO(2)$, the isometries of a 2-dimensional plane, and its representations are 2-dimensional[2] and labelled by the helicity of the state: a spin-1 state has 2 degrees of freedom.

The problem is that in this peculiar case, it is not possible to find polarization vectors $\epsilon_\mu(x, \vec{p})$, such that the left-hand and the right hand side of Eq. (1) have the right transformation properties: no 4-vector field can be constructed with annihilation/creation operators of a spin-1 massless particle [12]. If one tries to force so, the resulting monster will do the following under a Lorentz transformation:

$$V_\mu(x) \to \Lambda^\nu{}_\mu V_\nu(\Lambda x) + \partial_\mu\Omega(x, \Lambda), \tag{7}$$

for a function $\Omega(x, \Lambda)$. This second piece in the transformation law for V_μ clearly differentiates it from a Lorentz vector, but its peculiar form suggests that, if a theory can be defined modulo transformations of the type

$$V_\mu(x) \to V_\mu(x) + \partial_\mu\alpha(x) \tag{8}$$

for any function $\alpha(x)$, then Eq. (7) might be concealed with the correct Lorentz transformation for a 4-vector. This is called *gauge* invariance/redundancy and plays a central rôle in our understanding of the fundamental interactions: it is an inevitable consequence of Lorentz invariance and the existence of massless spin-1 states (in this sense, symmetries are a consequence of dynamics, rather than the opposite) and leads to a unique Lagrangian,

$$\mathcal{L} = -\frac{1}{4g^2}F_{\mu\nu}F^{\mu\nu}\,. \tag{9}$$

In fact, this symmetry also accounts for the disappearance of the one degree of freedom w.r.t. the massive case. This type of symmetry is perhaps better referred to as a *redundancy*, since it characterizes a situation in which different mathematical descriptions correspond to the same physical system. In fact, there is even so much of this redundancy (since $\alpha(x)$ is a complete set of functions) that an entire degree of freedom becomes unphysical.

In summary, a massless spin-1 state has 2 degrees of freedom, and its description in terms of a quantum field requires the introduction of gauge invariance. A massive spin-1 state has instead 3 degrees of freedom and gauge invariance is not apparently manifest; its description in terms of a quantum 4-vector Lorentz field requires the additional Lorentz covariant condition Eq. (3).

[1]This argument is not modified if V_μ couples to a conserved current $gV^\mu J_\mu$ in Eq. (4), since this cancels Eq. (5) because $\partial_\nu J^\nu = 0$.

[2]The representations are in fact 1-dimensional, but since parity interchanges helicity $h \to -h$, a manifestly parity preserving description must include multiplets with states of opposite helicity (for $h = 0$ this is trivially satisfied and the representation is in fact 1-dimensional)

2.2 The Higgs Mechanism: $4+1-2=3$

There is an interesting way of writing the Lagrangian for a massive spin-1 field, that is surprisingly equivalent to Eq. (6), but it involves an additional degree of freedom, in the form of a scalar $U(x) = \exp i\phi(x)$, and the Lagrangian

$$\mathcal{L} = -\frac{1}{4g^2}F_{\mu\nu}F^{\mu\nu} + \frac{v^2}{2}(D_\mu U)^2 = -\frac{1}{4g^2}F_{\mu\nu}F^{\mu\nu} - \frac{v^2}{2}(\partial_\mu\phi - V_\mu)^2 \tag{10}$$

where $D_\mu = \partial_\mu - iV_\mu$. This is, in fact, invariant under the symmetry

$$U(x) \to e^{i\alpha(x)}U(x) \qquad \phi(x) \to \phi(x) + \alpha(x) \qquad V_\mu(x) \to V_\mu(x) - \partial_\mu\alpha(x) \tag{11}$$

which includes the gauge invariance Eq. (8) for vector fields, and extends it to the scalar field. The theory described by Eq. (10) must therefore describe $4-2+1=3$ degrees of freedom (gauge invariance removing two degrees of freedom from a 4-vector, and the scalar adds one), equivalently to Eq. (6). That it is equivalent can also be understood by exploiting the gauge invariance and perform a transformation Eq. (11) on Eq. (10) with $\alpha = \phi$: this results in Eq. (6). This equivalence allows us to compare the theories for massive and massless vectors on an equal footing, both of them being based on intrinsic gauge invariance, and differing by the addition of a scalar degree of freedom, with the appropriate transformation properties. This is the essence of the *Higgs mechanism* and, in this form, provides a description of individual massive vectors, associated with Abelian gauge symmetries (such as in massive Quantum Electrodynamics – QED). In particular, the theory described by Eq. (6) or Eq. (10), and their extension to couplings with fermions based on gauge symmetry, can be extrapolated to arbitrary high energy.

Non-Abelian symmetries. We are interested however in providing a description of Nature and of the $W^\pm, Z$ bosons, whose gauge symmetry $SU(2)_L$ is in fact non-abelian. This can be described in a generalization of Eq. (10),

$$\mathcal{L} = -\frac{1}{4g^2}\mathrm{tr}\left[F_{\mu\nu}F^{\mu\nu}\right] + \frac{(v+ah)^2}{2}\mathrm{tr}\left[D_\mu U^\dagger D^\mu U\right] \tag{12}$$

where now $F_{\mu\nu} = \partial_\mu V_\nu - \partial_\nu V_\mu - i[V_\mu, V_\nu]$ for $V_\mu = V_\mu^i\sigma^i$, σ^i the Pauli matrices, and

$$U = e^{i\phi^i\sigma^i} \tag{13}$$

for three scalars ϕ^a, so that the Lagrangian is invariant under $g_L \in SU(2)_L$

$$U \to g_L U\,, \quad V_\mu \to g_L V_\mu g_L^\dagger - i g_L^\dagger \partial_\mu g_L\,. \tag{14}$$

Notice that I've included an additional scalar h in Eq. (12), whose importance will become clear later, but for the moment we can take $a = 0$. Contrary to the abelian case of Eq. (10), that is basically a free theory, Eq. (12) has self-interactions, that contribute for instance to the $2 \to 2$ scattering between four spin-1 states. In this process, a surprising feature appears in the particular channel involving their longitudinal polarization. As shown in the left panel of Fig. 1, this theory predicts an uncontrolled rise with energy in the scattering probability in this channel. In fact, for

$$E \gtrsim \Lambda_{cut} = \frac{4\pi v}{\sqrt{|1-a|}} \tag{15}$$

the theory doesn't make sense anymore as a Lagrangian weakly coupled description of this scattering process (the amplitude calculated with this theory would appear to violate unitarity).

Why does the amplitude grow? Scalar fields that appear in the Lagrangian through the exponential representation Eq. (13) (the non-linear sigma-model), are in fact very special in Nature: their interactions

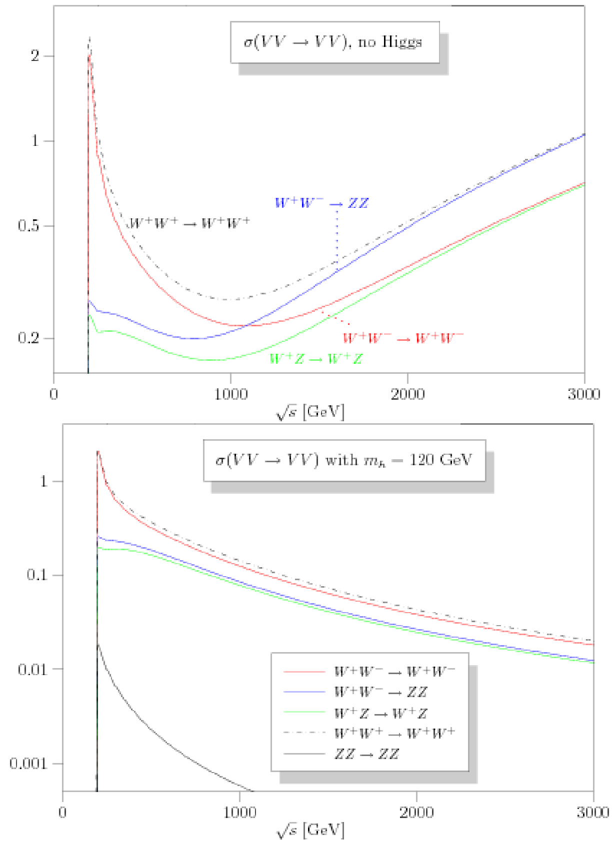

Fig. 1: Taken from [13]. Cross sections (in nanobarns) $\sigma(V_L V_L \to V_L V_L)$ for the longitudinal polarization of vectors in the SM, $V = W^{\pm}, Z$. The LEFT panel shows the energy-growth in the absence of a Higgs boson, while the RIGHT panel includes a Higgs bosons with $m_h = 120$ GeV.

are always associated with derivatives, and this leads to the rapid energy growth observed above. In fact, fields like these are Nambu-Goldstone bosons of spontaneously broken symmetries, and they always correspond to the low energy manifestation of a more complicated microscopic theory. This is were the unphysical behavior of Fig. 1 is bringing us: at high energy the theory of massive vectors does not describe Nature.

a=0 – Pions. A familiar example where we find the necessary scalar Eq. (13), is the pion Lagrangian from quantum chromodynamics (QCD). Here, for small $m_{u,d}$, QCD is almost invariant under the symmetry $\mathcal{G} = SU(2)_L \times SU(2)_R$, independently acting on the left- and right-handed up and down quark doublet

$$\Psi_{L,R} \equiv \begin{pmatrix} u \\ d \end{pmatrix}_{L,R} . \tag{16}$$

At low energy the quarks are no longer the relevant degrees of freedom, they condense at $E \sim \Lambda_{QCD}$ and define a new QCD vacuum, so that, at low energy, only the symmetry $SU(2)_L = SU(2)_R \equiv SU(2)_V = \mathcal{H}$ survives, while independent LR transformations are spontaneously broken. Goldstone theorem predicts three associated Nambu Goldstone bosons (NGB) corresponding to the pions π^a, which are the relevant degrees of freedom at low energy, as opposed to the quarks being the appropriate degrees of freedom at high-energy. These NGBs span the coset manifold $\mathcal{G}/\mathcal{H}$, which can be parametrized by $U = \exp i\pi^a\sigma^a$, transforming as $U \to g_L U g_R$. This is equivalent to Eq. (13), since the quark global $SU(2)_L$ corresponds indeed to the gauged weak interactions, while $U(1)_Y$ can be identified with a subgroup of $SU(2)_R$. So, QCD gives mass to the EW bosons: at $E \lesssim \Lambda_{QCD}$ Eq. (12) (with $a = 0$) describes massive W, Z bosons, while at $E \gtrsim \Lambda_{QCD}$ they appear massless, and the anomalous behavior of Fig. 1 is not realized. Eq. (12) is an effective field theory (EFT) with a finite range of validity. The problem with QCD is that v in Eq. (12) is not a free parameter, but determines also other pion interactions and has been measured $v = f_\pi \approx 130$ MeV, so that $m_{W^\pm} = g\frac{f}{2} \approx 40$ MeV...

This is clearly not what we observe in Nature, but this example remains illustrative as it provides an important proof of principle, that led the physics community to speculate on the existence of a new strong interaction, called Technicolor [14, 15], with new (techni)quarks in addition to the SM ones, such that the magic of QCD is repeated at a higher scale

$$v = f_{TC} = 246\,\text{GeV} \tag{17}$$

thus reproducing the correct mass spectrum for the EW gauge bosons.

a=1 – Higgs. Recall that in Eq. (12) I've included the interaction with an additional real scalar h (with assumed canonical kinetic term). The reason for doing so is that it is easy to see that this gives an additional contribution to the amplitude for $V_L V_L$ scattering, that at thigh-energy has the opposite sign compared to the $a = 0$ one, and leads to Eq. (15). For $a = 1$, the high-energy behavior cancels and Eq. (12) goes from an EFT with a small cutoff to a theory that in principle has an arbitrarily large range of validity $\Lambda_{cut} \to \infty$.

The reason is the following. For $a = 1$ we can write

$$(v + h)U \equiv \Sigma = (\widetilde{H}, H)\,, \quad \Rightarrow \quad \frac{(v+h)^2}{2}\text{tr}\,[D_\mu U^\dagger D^\mu U] = |D^\mu H|^2 \tag{18}$$

where we have rewritten the fields using a different parametrization,

$$H = \exp(i\pi^a\sigma^a/\phi)\begin{pmatrix} 0 \\ \phi \end{pmatrix} = U\begin{pmatrix} 0 \\ \phi \end{pmatrix} = \frac{1}{\sqrt{2}}\begin{pmatrix} \phi_1 + i\phi_2 \\ \phi_3 + i\phi_4 \end{pmatrix}, \tag{19}$$

and $\widetilde{H} = \epsilon H^\dagger$ (ϵ the antisymmetric tensor). This map works only when $|H| \equiv \phi \neq 0$, and is singular in $\phi = 0$. The important aspect is that now H transforms as a fundamental representation **2** of $SU(2)_L$ and his own Lagrangian at dimension ≤ 4 is simply

$$\mathcal{L}_H = \partial_\mu H^\dagger \partial^\mu H - V(H)\,, \qquad V(H) = -m_H^2|H|^2 + \lambda|H|^4\,, \tag{20}$$

which describes four real scalars with non-pathological self interactions, whose scattering amplitudes are well-behaved also at high-energy. This has to be contrasted with the Lagrangian for U only, $v^2\text{tr}\,[\partial_\mu U \partial^\mu U]/2$ that has instead a cutoff at $4\pi v$.

In some sense, we have found a different UV completion for our effective Lagrangian Eq. (12), that simply involves an additional scalar $\phi = v + h$, that however renders the amplitude physical, thanks to its contribution to scattering processes. This is illustrated in the right panel of figure 1. Now, what guarantees $|H| \neq 0$ everywhere so that the field H has a vacuum expectation value $\langle 0|H^\dagger H|0\rangle \neq 0$? This is revealed from the H potential $V(H)$ in Eq. (20). Notice that this is independent of the NGBs, which cancel in $|H|^2$; this is why they are massless NGBs, because they don't appear in the potential. Then we see that for positive $m^2 > 0$ and positive $\lambda > 0$ the potential has a minimum at

$$\langle\phi\rangle \equiv v = \frac{\sqrt{m^2}}{\lambda} \tag{21}$$

and this is the average value of the field everywhere in spacetime. Because of this value, the low-energy Lagrangian has $SU(2)_L \times U(1)_Y$ symmetry realized non-linearly: we say that EW symmetry is

	SM Fields		$SU(3)_C,\, SU(2)_L,\, U(1)_Y$
spin-0	Higgs	H	$(\mathbf{1},\, \mathbf{2}\,,\, -\frac{1}{2})$
spin-1/2	Quarks (×3 families)	$Q_L = (u_L \;\; d_L)$	$(\mathbf{3},\, \mathbf{2}\,,\, \frac{1}{6})$
		$u_R^\dagger$	$(\overline{\mathbf{3}},\, \mathbf{1},\, -\frac{2}{3})$
		$d_R^\dagger$	$(\overline{\mathbf{3}},\, \mathbf{1},\, \frac{1}{3})$
	Leptons (×3 families)	$L_L = (\nu \;\; e_L)$	$(\mathbf{1},\, \mathbf{2}\,,\, -\frac{1}{2})$
		$e_R^\dagger$	$(\mathbf{1},\, \mathbf{1},\, 1)$
spin-1	Gluon	g	$(\mathbf{8},\, \mathbf{1}\,,\, 0)$
	W bosons	$W^\pm \;\; W^0$	$(\mathbf{1},\, \mathbf{3}\,,\, 0)$
	B boson	B^0	$(\mathbf{1},\, \mathbf{1}\,,\, 0)$

Table 1: The SM field content and quantum numbers.

spontaneously broken (EWSB). Now field excitations have to be considered around this minimum, and we see that $\delta\phi \equiv h$ has mass

$$m_h^2 = \lambda v^2 \,. \tag{22}$$

So, this theory gives the 3 massive vectors and one massive scalar, a prediction summarized by Higgs in his original article [2], as *the prediction of incomplete multiplets of scalar and vector bosons*, which granted him and Englert [3] the Nobel prize in 2013.

In this introduction to the Higgs mechanism, I have tried to give a feeling for the necessity of gauge invariance and the Higgs mechanism, and then exposed to examples of the latter. We do not know yet with certitude how the EW symmetry breaking sector looks like, although we already know that it is not of the form of a purely technicolor interaction. Indeed, in 2012, an *incomplete multiplet* – called the Higgs boson – has been discovered at the LHC, suggesting that a field in the linear representation of the EW group provides an appropriate description of nature, at least in the regime of energy tested at the LHC so far.

3 The Standard Model

In a seminal article, Weinberg [1] proposed a model in which fermions and vectors interact with the Higgs field H. He pointed out the well-behaved high-energy limit of amplitudes, saying that *this model may be renormalizable*: this is now the definition of the Standard Model of particle physics. The field content includes gauge bosons $V = G, W, B$ associated with the SM gauge symmetry $SU(3)\times SU(2)_L \times U(1)_Y$, matter fermions $\psi = Q, u, d, L, e$, and of course the Higgs field. The quantum numbers are reported in table 1, while their dynamics is described by a very simple Lagrangian which, at least at first sight, appears to plausibly describe the behavior of elementary particles:

$$\mathcal{L}_{SM} = \sum_\psi i\bar\psi \not{D}\psi + h.c. - \sum_V \frac{1}{4} V_{\mu\nu}V^{\mu\nu} + |D_\mu H|^2 - V(H) + y_{ij}^{D,L} H \bar\psi_i \psi_j + y_{ij}^{U} \tilde{H} \bar\psi_i \psi_j \,, \tag{23}$$

where $D_\mu = (\partial_\mu + ig'YB_\mu + igW_\mu^a\sigma^a + ig_sG_\mu^a\lambda^a)$. Notice that Left- and Right-handed fermion have different quantum numbers (they are referred to as *chiral*), implying that a standard mass term $m(\bar\psi_L\psi_R + \bar\psi_R\psi_L)$ would not respect gauge symmetry. In this sense, the Higgs field becomes necessary also to provide a mass to the fermions: indeed $\bar\psi_L H \psi_R$ is a gauge singlet, and induces a fermion mass after EWSB, $\phi \to h + v$.

3.1 Accidental and Approximate Symmetries

The SM Lagrangian Eq. (23) consists of only relevant and marginal operators (that is, operators with dimension $D \leq 4$). This is its defining feature, that allows it to be a valid theory over many orders of magnitude in distance (in principle) and also that makes it such a predictive theory. Indeed, the limited number of interactions Eq. (23) implies many relations and structures that can be tested. An interesting way of keeping track of these relations, and to understand to what extent these relations are solid, is to identify the symmetries of the SM Lagrangian (some of which might be only approximate).

The SM Lagrangian Eq. (23) is invariant under a $U(1)_B$ symmetry, called *Baryon symmetry* or Baryon number, that acts on quarks and anti-quarks with opposite charge, as well as three $U(1)_{L_i}$ global symmetries that act on the three families of leptons, called *lepton numbers*. These symmetries are *accidental*, in the sense that they follow from the fact that only interactions of dimension ≤ 4 appear in Eq. (23), but they are *exact*. Indeed an operator of dimension-5, $LHHL$, called Weinberg operator, violates the lepton numbers (and gives mass to neutrinos), while at dimension-6, there are operators that violate baryon number. These symmetries imply important predictions in the SM: the absence of proton decay and vanishing neutrino masses.

Many other relations, especially in the context of Higgs physics, can instead be understood in terms of *custodial symmetry* $SU(2)_c$. This is an accidental and approximate symmetry of the Lagrangian. Custodial symmetry is best understood by writing the Higgs field as the 2×2 matrix Σ in Eq. (18). Now, Σ transforms as $\Sigma \to g_L\Sigma$ under $g_L \in SU(2)_L$, but we can conceive an extended (global) transformation $\Sigma \to \Sigma \to g_L \Sigma g_R$ for $g_{L,R} \in SU(2)_{L,R}$. For $g' \to 0$ and vanishing Yukawas, the SM Lagrangian involving the Higgs field can be written as

$$\mathcal{L}_{SM}^{\Sigma} = \mathrm{tr}\,(D_\mu \Sigma^\dagger D^\mu \Sigma) - V(|\Sigma|) \tag{24}$$

and respects this symmetry. The Higgs vev $\langle \Sigma \rangle = \mathrm{diag}(v, v)$ breaks it spontaneously to the diagonal subgroup $SU(2)_c = SU(2)_L = SU(2)_R$. This is called *custodial* symmetry because it implies that the mass of the W and Z bosons be identical. We can now keep track of the parameters that do not respect this symmetry by attributing them spurious transformation properties [16]

$$g' \sim (\mathbf{1_L}, \mathbf{3_R}) \qquad Y^U \sim (\mathbf{1_L}, \mathbf{2_R})\,. \tag{25}$$

This can help to keep track of the size of certain effects in the SM. For instance, the W boson mass matrix $m_W^{ab} W^a W^b$ reduces into $\mathbf{3}\otimes\mathbf{3} = \mathbf{1}\oplus\mathbf{3}\oplus\mathbf{5}$, and the Z/W mass difference $m_Z^2 - m_{W^\pm}^2 = (m_W^2)^{33} - (m_W^2)^{11}$, can only appear in the $\mathbf{5}$ of $SU(2)_L = SU(2)_R = SU(2)_c$. Using Eq. (25) we can see that, in the SM, we can have effects $m_Z^2 - m_{W^\pm}^2 \propto (g')^2$. Indeed, at tree-level,[3]

$$\frac{m_Z^2}{m_W^2} = 1 + \frac{g'^2}{g^2} \tag{27}$$

This is particularly important for BSM physics, where custodial symmetry is no longer accidental and the ratio Eq. (27) can be modified.

3.2 Higgs Physics in the SM

Other *accidental* relations that characterize the SM cannot be attributed directly to symmetries, yet they can be considered on the same footing as defining features of the SM.

[3] At loop-level, effects involving the top-Yukawa become manifest. Our power-counting suggests that these $\propto (Y^t)^4$, but the explicit computation gives, for $m_t \gg m_b$,

$$\Delta\frac{m_Z^2}{m_W^2} = -\frac{3}{2}\frac{(Y^t)^2}{16\pi^2}\cos^2\theta_w\,. \tag{26}$$

This is due to the fact that this contribution is related to an IR effect, regulated by m_t^2 in the propagator, that removes two powers of Y_t from our power-counting. In other words, Y_t cannot be considered a small spurion compatibly with the $m_t \gg m_b$ assumption: in the limit $Y_t \to 0$, the approximated expression Eq. (26) actually becomes infinite, but an exact computation with $m_b, m_t \to 0$ would show the expected behaviour.

The classical Lagrangian Eq. (23) can be expanded as $\phi \to h + v$, and the Goldstone bosons can be absorbed through an $SU(2)_L$ transoformation (unitary gauge).

$$|D_\mu H|^2 \supset (m_W^2 W_\mu^+ W^{-\mu} + \frac{1}{2} m_Z^2 Z_\mu Z^\mu) \left(1 + 2\frac{h}{v} + \frac{h^2}{v^2}\right) \tag{28}$$

$$V(H) \supset -\frac{1}{2} m_h^2 h^2 \left(1 + 2\frac{h}{v} + \frac{h^2}{v^2}\right) \tag{29}$$

$$H\bar{\psi}_i\psi \supset m_\psi \bar{\psi}_i \psi (1 + \frac{h}{v}), \tag{30}$$

where $m_W = gv/2$, $m_Z = m_W/\cos\theta_W$ ($\cos\theta_W = g/\sqrt{g^2 + g'^2}$), $m_h = \lambda v/2$ and $m_\psi = yv/2$ in terms of the Lagrangian parameters. Interestingly, once the masses of the SM particles are measured, Eqs. (28-30) fix uniquely their coupling to physical Higgs bosons, in such a way that it is proportional to their mass. The fact that h interactions are naturally aligned with the fermion masses plays a crucial rôle for the phenomenology of the SM, forbidding to very high accuracy Flavor Changing Neutral Currents (FCNCs).

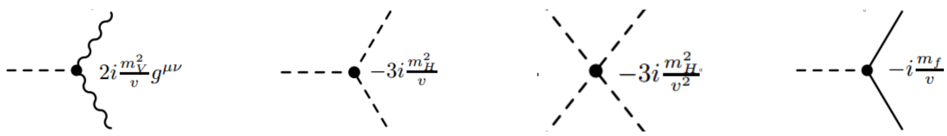

Fig. 2: Tree-level Higgs couplings in the SM.

The SM is a renormalizable theory. In practice this means that *infinite* quantum effects must be unobservable, as they cancel against infinite Lagrangian counter-terms and relate to input (measured) parameters of the theory. On the contrary, *finite* quantum effects are physical and observable: these constitute a robust prediction of the theory and an important test of its structure. For instance, quantum

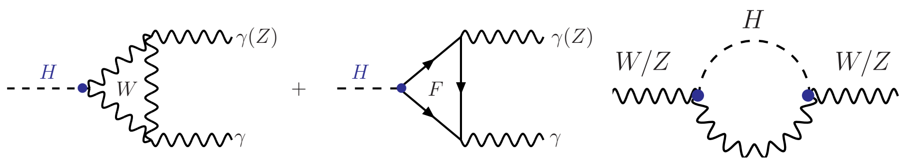

Fig. 3: Some loop effects involving the Higgs in the SM.

effects imply small (finite) departures of the tree-level relations in Fig. 2, that can be ignored for the purpose of this review. They are however important when they are associated with effects that are not present at tree level. I will discuss the most important here, though keep in mind that experiments can be designed to be sensitive also to other effects that I am here neglecting. First of all, the Higgs doesn't couple to photons, to gluons, nor to $Z\gamma$ at tree-level; all these are generated by loop effects, some of which are reported in Fig. 3. The exact expressions can be found e.g. in [11]; what matters to us is that the resulting expressions are completely determined by the tree-level couplings Eqs. (28-30) and are therefore a prediction of the SM. For these reason, knowledge of the SM particle masses gives us a concrete prediction on the physics of the Higgs boson, as shown in Fig. 4 that shows the different branching ratios for Higgs decays in the SM, for different values of m_h (computed before the discovery of the Higgs boson).

It is interesting to point out that loop effects allowed to test the physics of the Higgs boson, even before its discovery. Indeed, the last diagram of Fig. 3 shows a loop effect that contributes a finite amount to processes involving a Z or a W boson. Since H propagate in the loop, these effects are $\propto \log m_h$. Precise measurements of the Z-boson properties at LEP [17], allowed therefore to extract the information of the RH side of Fig. 4, $m_h = 94^{+29}_{-24}$ GeV.

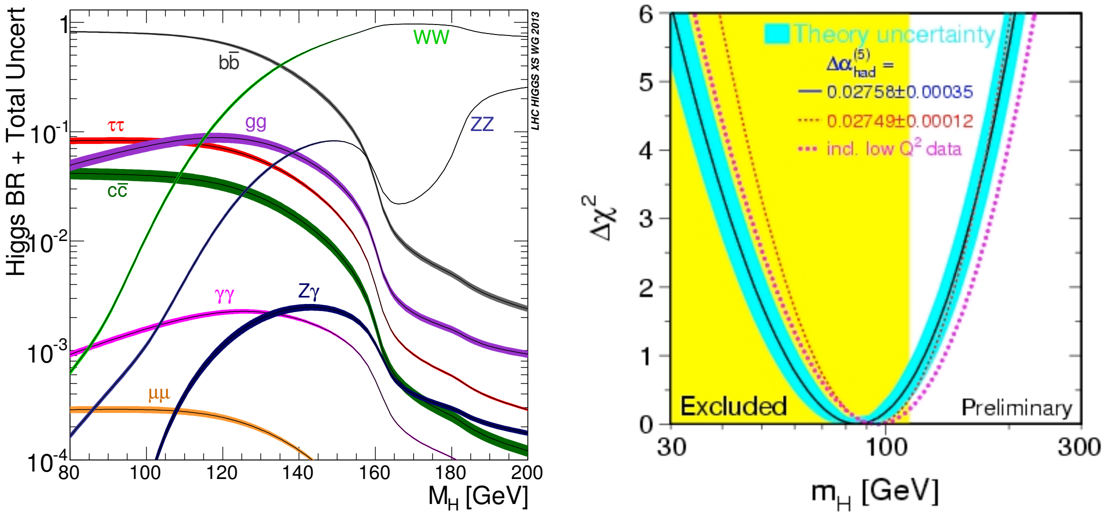

Fig. 4: LEFT: Higgs branching ratios in the SM, for different values of m_h [18]. RIGHT: preferred value of m_h, from a global fit to EW data; the yellow band corresponds to the direct LEP constraints $m_h > 114.4$ GeV [19].

So, when produced, a Higgs boson decays predominantly in b-quarks, W-bosons, gluons, etc. But how is it produced, to begin with? Amusingly, the dominant production mode at the LHC is through a loop effect: the cross-section for gluon fusion $gg \to h$ is large $\sigma_{gg\to h} \approx 44(19)$ pb at 13(8) TeV, simply because the proton content of gluons is very large. This is followed by tree-level, but electroweak, processes: vector boson fusion $qq \to VVqq \to hqq$ has $\sigma_{VBF} \approx 3.7(1.6)$ pb, while vector-associated production $\bar{q}q \to V^* \to Vh$ has $\sigma_{VH} \approx 2.2(1.1)$ pb; yet these are the dominant production modes at e^+e^- colliders. Finally, production in association with top quarks $pp \to t\bar{t}h$ constitute a small fraction of the total cross-section $\sigma_{tth} \approx 0.51(0.13)$ pb.

For what concerns the SM, the only information that was still missing before the LHC was m_h. This could be measured with extreme precision

$$m_h = 125.09 \pm 0.24 \text{ GeV} \quad [6] \tag{31}$$

thanks to the Higgs decaying to $\gamma\gamma$ and ZZ, which have the best mass-resolution (decays to W-bosons are instead penalized by the impossibility to detect neutrinos and hence to reconstruct the invariant mass of the W-pair).

As we will see in the next section, from a BSM point of view, all channels are important, because they allow us to test alternative hypotheses in which the Higgs couplings might depart from the SM ones. In this sense (borrowing Fabiola Gianotti words), *Nature has been kind to us*, because the Higgs mass happens to be such that all decay channels are more or less important. Indeed, had m_h be for instance 170 GeV, we would have been dominated by the WW mode, with no hope to ever observing $\gamma\gamma$ decays.

In the BSM-motivated quest of testing the Higgs couplings, it is important to keep in mind that the LHC is a complicated machine, in which size does not always matter. For instance, the process $gg \to h \to b\bar{b}$ has by far the largest cross-section, but the same is true for the QCD-dominated background, which renders this channel practically invisible. Decays into $\bar{b}b$ have therefore to be tested in the VH production mode, where the leptonic decays of the associated vector allow the event to be detected (or in VBF, where the same rôle is played by the special kinematics of the qq pair).

A convenient way to express the results of this exploration, is portrayed in Fig. 5. Here the Higgs couplings to the SM particles are multiplied by an arbitrary factor μ, the *signal strength*, such that the SM coincides with $\mu = 1$; then the μ_i are fitted as if they were free parameters of the theory. The information contained in these fits, will form the basis for our BSM exploration in the next section.

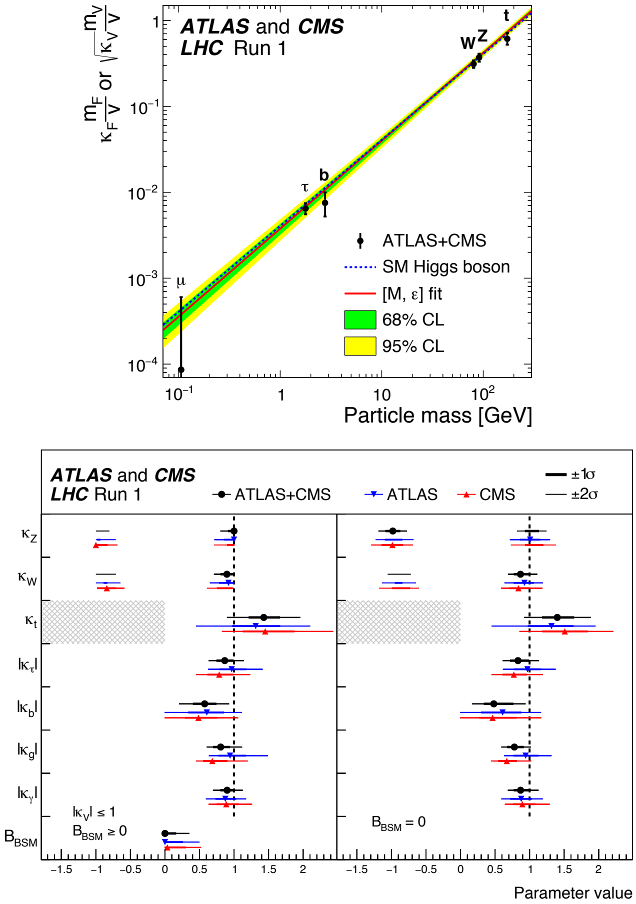

Fig. 5: One of the rare ATLAS+CMS combinations from [20]. LEFT: fit to tree-level Higgs couplings, testing the coupling/mass relation as predicted in the SM, Eqs. (28-30). RIGHT: test of the SM Higgs couplings, performed with a global fit by rescaling each SM tree-level coupling to particle i by a factor κ_i (the "κ" framework) - B_{BSM} denotes a branching ratio into undetected particles.

4 Beyond the Standard Model

I mentioned before that the defining feature of the SM is the possibility, given its matter content, to extrapolate it to very high-energy, thus allowing it to be (very close to) a fundamental theory.[4] First of all, this is just a possibility: nothing forbids the presence of new structure at distances close to the ones probed today. Secondly, we have indications that such structure exists. For instance, electric charge runs towards higher values as the energy is increased, and eventually becomes strong and non-perturbative at $M_{Landau} \approx 10^{275}$ GeV, signaling the existence of new dynamics. Gravity, must ultimately become part of the SM, but it's known to become strong at energies of order $M_{Plank} \approx 10^{19}$GeV. Other in-

[4]More precisely, the SM includes only relevant and marginal operators that represent a finite theory keeping their couplings fixed and sending the cutoff to ∞.

dications include the existence of Dark Matter, the neutrino masses, the necessity for a mechanism for baryo(lepto)genesis, all of which are not accommodated in the SM.

All these reasons motivate searches for physics beyond the SM, in all its possible incarnations. There is however one additional reason, that pushes us to search for BSM physics explicitly *at the LHC* and in particular in *Higgs physics*: Naturalness. This much disputed principle has to do with the hierarchy problem, expressed by Wilson [21] as the impossibility for the existence of light scalar particles that are not associated with the breaking of some symmetry. This argument includes an elementary Higgs scalar, such as the one appearing in the SM, and is a consequence of the fact that the Higgs mass term $|H|^2$ in the Lagrangian has dimension-2 and is therefore a very relevant operator. Relevant operators are UV-sensitive, so that the value of the Higgs mass we observe is related to quantities at high-energy; for instance changing the value of $m_h^2(M_{Plank})$ - the Higgs mass parameter measured at the Planck scale - by a factor of two within the SM, corresponds to changing $m_h^2(\text{TeV})$ by a factor $M_{Plank}^2/\text{TeV}^2 = 10^{32}$! This Wilsonian point of view is the formal implementation of reductionist thinking, that is: low-energy quantities such as $m_h^2(\text{TeV})$ can be computed in terms of more fundamental quantities $m_h^2(M_{Plank})$; in this sense the observed value $m_h^2 \equiv m_h^2(\text{TeV}) \approx (100\text{GeV})^2 \ll 10^{4+32}\text{GeV}^2$ appears to be finely tuned in the UV and this is the essence of the hierarchy problem.[5]

More concretely, in a fundamental theory that predicts m_h^2 in terms of more fundamental parameters, we could integrate $m_h^2(E)$ along the renormalization group flow (denoted here with E, to highlight the physical Wilsonian interpretation of E as the momentum of particles in the loops), to obtain its value at small energy (see e.g. [9]):

$$m_h^2 = \int_0^\infty dE \frac{dm_h^2(E)}{dE} = \int_0^{\Lambda_{UV}} \frac{dm_h^2(E)}{dE} + \underbrace{\int_{\Lambda_{UV}}^\infty dE \frac{dm_h^2(E)}{dE}}_{\equiv m_h^2(\Lambda_{UV})}, \tag{32}$$

where we have separated the contribution from below/above an arbitrarily chosen physical scale Λ_{UV}. The point is that we can compute the first integral in the RH side of Eq. (32) with the observation that at low-energy the SM describes properly Higgs physics; for $\Lambda_{UV} = M_{Plank}$ we obtain a contribution $\Delta m_h^2 = 10^{36}\text{GeV}^2$ as mentioned above, that must be finely cancelled against an (uncorrelated, in the Wilsonian point of view!) contribution from $m_h^2(\Lambda_{UV})$.

Paradoxically, the best way to formulate the hierarchy problem, is in models that solve it, i.e. models that do exhibit additional symmetry, such as composite Higgs models (CHM) or Supersymmetry (SUSY), such that

$$m_h^2(\Lambda_{UV}) = 0\,. \tag{33}$$

This is achieved in CHM because $|H|^2$ is a composite operator in the theory above Λ_{UV} and is in fact irrelevant. In SUSY, instead, non-renormalization theorems imply that $dm_h^2(E)/dE = 0$ and hence Eq. (33). In both cases, however the dominant contribution to m_h^2 comes from the first integral in the RH side of Eq. (32) that is typically a loop factor smaller than Λ_{UV}^2 (in strongly coupled CHM this is possible only if the Higgs is a PNGB of a global SSB [23]). Given the Higgs couplings in the SM, this part can be calculated and one finds that no fine-tuning corresponds to

$$\Lambda_{UV} \lesssim 450\,\text{GeV}\,, \tag{34}$$

that is: new dynamics has to modify Higgs physics at a physical scale accessible to the LHC.

Now, we can tolerate some level of fine-tuning (this can be quantified using the definition of Ref. [24] and increases quadratically with the new physics scale), but these arguments clearly single out different directions for the future of particle physics:

[5]In the SM, this relevant operator is very sensitive to the UV, but in principle one could immagine a modification of the SM where this sensitivity is tamed, because the dimension of the $|H|^2$ operator is smaller. First principles exclude this possibility [22].

– Rethink the Wilsonian approach in the grander scheme of things
– Search directly for the new states at Λ_{UV}
– Test the properties of the Higgs boson, which are expected to depart from the SM ones, signaling a modification of the first integral in Eq. (32)

The first approach has recently provided promising directions that take into account the cosmological history of the universe [25] with or without the inclusion of anthropic arguments [26], and it is not clear how wide the spectrum of possibilities is. The search of direct states in the TeV region constitutes instead the bulk of the LHC search program; its different ramifications depend very much on how we think this different dynamics will be, I refer to reviews on CHM or SUSY.

Instead, testing the properties of the Higgs boson is a well-defined and compact field of research. Indeed, from an experimental point of view, there is only a limited number of observables that can be tested in the context of Higgs physics. Interestingly, also from a theoretical point of view, independently from the details of the UV dynamics, there is only a limited number of ways in which Higgs physics can be modified. This is a consequence, again, of the Wilsonian point of view: integrating out new physics at the scale Λ_{UV} generates a set of local operators, corresponding to an Effective Field Theory (EFT). Given that only the more relevant ones survive at low-energy and given that there is only a finite number of operators of a given relevance, there can be only a finite number of effects that is worth studying in Higgs physics. We discuss this in detail in what follows.

4.1 BSM Higgs Phenomenology

An important thing to keep in mind, when thinking about Higgs phenomenology, is that the structure of the SM is rather unique: the relation between the different couplings, implied mostly by gauge-invariance, allow the theory to be valid to high-scales. So,

Any modification of the SM couplings reduces the cut-off of the theory.

This has to be read together with the definition that

Any theory with a cutoff is an Effective Field Theory,

that can be written as an expansion in local field operators of different dimensionality

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{SM}} + \sum_i \frac{c_i^{(6)}}{M^2}\mathcal{O}_i^{(6)} + \sum_j \frac{c_j^{(8)}}{M^4}\mathcal{O}_j^{(8)} + \cdots, \tag{35}$$

where $c_i^{(D)}$ are called Wilson coefficients and the leading (relevant) contribution $\mathcal{L}_{\text{SM}} \equiv \mathcal{L}_{\leq 4}$ is the one surviving in the limit where M, which is here the mass of new physical states, is taken $\to \infty$.[6]

In other words, what we can learn from testing different aspects of Higgs physics, can be captured in the language of EFTs. For instance, the 'κ' framework [18], where all SM Higgs couplings are rescaled (as in $\kappa_Z = g_{hZZ}/g_{hZZ}^{SM}$ for the $hZ_\mu Z^\mu$ coupling to Z-bosons), lowers the cutoff of the theory for $\kappa \neq 1$, and eventually corresponds to an EFT, despite the fact that no scale appears explicitly in its formulation (we had an example of this for $a \neq 1$ in Eq. (12)).

From a practical, experimental, point of view, there have been different proposals to parametrize the experimental information that can be extracted from measurements of the Higgs properties, but ultimately, the information they carry can be matched to EFTs. Most notably, pseudo-observables - POs

[6] Here I'll discuss only theories that have a well defined decoupling-limit, see however [27] for an interesting case, with a naturally light Higgs, that cannot be captured by such an EFT. Moreover, I refer to Refs. [28, 29] (and to the appendix of [30]) for parameterizations in which EW symmetry is never linearly realized (useful for instance to capture Technicolor theories with an accidentally light Higgs).

- [31, 32] are designed as EFT-inspired expansions around the poles of certain scattering amplitudes involving h; as such they are physical quantities and can be used as an important tool to extract information about EFTs from experiment. They are particularly useful in processes involving on/off-shell EW bosons, such as $h \to V\bar{\psi}\psi$, because their measurement depends on a minimum of theoretical input: precision QCD/EW calculations can then be used to extract from them information about the Wilson coefficients c_i, but since the precision of such computations constantly improve, measurements in terms of POs constitute a durable legacy.

For processes involving QCD states, POs are unfortunately less effective, since the experimental information is encoded in objects -jets - that correspond to a multiplicity of states at the fundamental level. In this case, simplified template or fiducial cross-section measurements [33] try to package the experimental information in a way that has reduced sensitivity to theoretical uncertainties (which are prone to improve when future calculations will be available) and enhanced sensitivity to the effects induced by EFTs.

Having said this, let us discuss what EFTs for Higgs physics really are, were we are interested in theories which have the same field content as the SM (see table 1). EFTs are all about formulating hypotheses about microscopic physics and to follow how these hypotheses can be tested with low-energy experiments. The assumptions I'll make in what follows are:

– There are no states beyond the SM ones at the energies relevant for these experiments, i.e. $E \ll M$. This assumption is motivated by the so far null results of the LHC.
– There is well-defined decoupling limit $M \to \infty$ in which the SM is recovered; in particular the Higgs field behaves (at least approximately) like an $SU(2)_L$ SM linear doublet. This is motivated by the left panel of Fig. 5: experimental data implies that departures from this limit must be small and presumably associated with a small expansion parameter.
– New physics is flavor-universal. This is motivated by the difficulty of accommodating experimental constraints in models with flavor non-universal new physics, but also from simplicity. See Ref. [34] where this assumption is relaxed in the context of Higgs/EW physics. Similarly we assume here that new physics conserves CP.
– We assume, to begin with, that the new physics is such as to generate all allowed operators, so that the leading effects are necessarily given at dimension-6, all higher dimension effects being irrelevant.

This last assumption is the least motivated: typically scenarios of new physics affect different sectors of the SM in different ways, symmetries can forbid some operators in favor of others, and large couplings can enhance effects that would otherwise be small. We will discuss these aspects below, but for the moment this hypotheses provides a conservative way for exploring the new physics landscape without much commitment.

4.1.1 Dimension-6 Effects

An important aspect of EFTs is that some (combinations of) effective couplings are redundant and do not induce any physical effect. Consider an infinitesimal (small ϵ) local transformation

$$\phi(x) \to \phi(x) + \epsilon F(\phi(x), \partial\phi(x)) \tag{36}$$

for a generic function F. Such field redefinition does not change the S-matrix, as long as the redefined field has non-vanishing matrix elements with the states sourced by the original one, but it changes the action by

$$\delta\mathcal{S}[\phi] = \epsilon \int d^4x \frac{\delta\mathcal{S}[\phi]}{\delta\phi(x)} F(\phi(x), \partial\phi(x))\,. \tag{37}$$

This implies that pieces of the action that can be written in the form of Eq. (37), can be eliminated by such field redefinition without changing the physics: they are therefore redundant. An example can clarify this better:

$$\mathcal{S} = \int d^4x \frac{(\partial_\mu \phi)^2}{2} + \frac{c_\phi}{M^2}\phi^3\Box\phi \quad \xrightarrow{\left(\phi \to \phi - c_\phi \frac{\phi^3}{M^2}\right)} \quad \mathcal{S} = \int d^4x \frac{(\partial_\mu \phi)^2}{2} + O\left(\frac{1}{M^4}\right) . \qquad (38)$$

The irrelevant interaction c_ϕ is proportional to the leading-order equations of motion and can be eliminated, up to higher order effects in the $1/M^2$ expansion (which here plays the ϵ role). From a practical point of view, such redundancy can be thought (together with integration by parts) as the freedom of choosing different forms of the Lagrangian (similarly to gauge invariance), to highlight different properties of the theory under scrutiny.

So, under the above assumptions, and focussing on non-redundant sets of operators, there is only a handful of operators that can modify Higgs physics. These are summarized in Table 2 where we exploit the above-mentioned freedom to write them in three different bases, corresponding to SILH [35, 36], Warsaw [37], and BSM primaries/Higgs basis [33,38,39]. Integration by parts and field redefinitions allow to swap the blue operators in the SILH basis with the red ones in the Warsaw basis (plus a redefinition of other Wilson coefficients).

It is sometimes useful to classify these effects according to their transformation properties under the SM $SU(2)_L \times SU(2)_R$ accidental symmetry, which derive from

$SU(2)_L \times SU(2)_R$	Higgs only	Higgs and Derivative
$(\mathbf{1}_L, \mathbf{1}_R)$	$\text{tr}\,(\Sigma^\dagger\Sigma) = H^\dagger H$	$\text{tr}\,(\Sigma^\dagger D_\mu \Sigma) = D_\mu(H^\dagger H)/2$
$(\mathbf{1}_L, \mathbf{3}_R)_{Y=0}$	$\text{tr}\,(\Sigma^\dagger\Sigma\sigma^b) = 0$	$\text{tr}\,(\Sigma^\dagger D_\mu\Sigma\sigma^b) \xrightarrow{b=3} -H^\dagger \overleftrightarrow{D}_\mu H/2$
$(\mathbf{3}_L, \mathbf{1}_R)$	$\text{tr}\,(\Sigma^\dagger\sigma^a\Sigma) = 0$	$\text{tr}\,(\Sigma^\dagger\sigma^a D_\mu\Sigma) = H^\dagger\sigma^a \overleftrightarrow{D}_\mu H/2$
$(\mathbf{3}_L, \mathbf{3}_R)_{Y=0}$	$\text{tr}\,(\Sigma^\dagger\sigma^a\Sigma\sigma^b) \xrightarrow{b=3} -(H^\dagger\sigma^a H)$	$\text{tr}\,(\Sigma^\dagger\sigma^a D_\mu\Sigma\sigma^b) \xrightarrow{b=3} -D_\mu(H^\dagger\sigma^a H)$

(39)

where the $\mathbf{3}_R$ is broken down to its components, since $SU(2)_R$ is not necessarily a BSM symmetry.[7] For instance, in the SILH basis, the operator $\mathcal{O}_T$ is $(\mathbf{1}_L, \mathbf{3}_R) \otimes (\mathbf{1}_L, \mathbf{3}_R) \supset \mathbf{5}_R$ and is indeed associated with maximal custodial symmetry breaking. Similarly, $\mathcal{O}_{H\psi_R}$ and $\mathcal{O}_{H\psi_L}$ are $\sim \mathbf{3}_R$ and also break custodial. On the other hand, using the spurious transformation properties of g' from Eq. (25), $\mathcal{O}_B$ is a singlet, in the sense that it doesn't introduce further custodial symmetry breaking than the SM. All other operators are singlets too. This classification is important because $SU(2)_c$ in the SM implies some relations (see section 3.1) that are well preserved at loop level (as discussed below Eq. (25)); departures from these relations can be well measured and constitute accurate BSM probes.

Clearly each individual operator contributes to different physical processes; for instance $\mathcal{O}_{HQ}$ modifies the Z couplings to left-handed quarks, but also contributes directly to $h \to Z\bar{Q}Q$ decays, etc. This fact complicates a global fit, but also makes its results difficult to present and to interpret, as marginalized constraints are often dominated by the poorest observables.[8] For this reason it is useful to identify the most relevant experiments that can test the operators of Table 2 and organize them according to their precision, and the operators they are sensitive to.

[7] In Eq. (39) we have kept the third one $b = 3$ that is associated with vanishing hypercharge $Y = 0$, but keep in mind that operators, like $\mathcal{O}_R^{ud}$ in the caption of Table 2, can be formed also with the $\pm$ components that give $\tilde{H}^\dagger\sigma^a H$ with hypercharge $Y = \pm 1$

[8] In principle the $n \times n$ correlation matrix relating measurements of n Wilson coefficients carries all the necessary information. In practice, however, results are often presented as marginalized over $n-1$ parameters so that, when the correlation matrix has large off-diagonal components, constraints are dominated by the least sensitive measurements and appear, therefore, artificially weak. These are sometimes called "blind directions".

	SILH	Warsaw	BSM Primaries
LHC Higgs 10%	$\mathcal{O}_6 = \lambda\vert H\vert^6$	$\mathcal{O}_6$	$\Delta\mathcal{L}_{3h}$
	$\mathcal{O}_{BB} = g'^2\vert H\vert^2 B_{\mu\nu}B^{\mu\nu}$	$\mathcal{O}_{BB}$	$\Delta\mathcal{L}^h_{\gamma\gamma}$
	$\mathcal{O}_{GG} = g_s^2\vert H\vert^2 G^A_{\mu\nu}G^{A\mu\nu}$	$\mathcal{O}_{GG}$	$\Delta\mathcal{L}^h_{GG}$
	$\mathcal{O}_{y_\psi} = y_e\vert H\vert^2(\bar\psi_L H\psi_R)_{\psi=u,d,e}$	$\mathcal{O}_{y_\psi}$	$\Delta\mathcal{L}^h_{\psi\psi}$
	$\mathcal{O}_H = \frac{1}{2}(\partial^\mu\vert H\vert^2)^2$	$\mathcal{O}_H$	$\Delta\mathcal{L}^h_{VV}$
	$\mathcal{O}_{HB} = ig'(D^\mu H)^\dagger(D^\nu H)B_{\mu\nu}$	$\mathcal{O}_{WW} = g^2\vert H\vert^2 W^a_{\mu\nu}W^{\mu\nu\, a}$	$\Delta\mathcal{L}^h_{Z\gamma}$
LEPI%	$\mathcal{O}_{HW} = ig(D^\mu H)^\dagger\sigma^a(D^\nu H)W^a_{\mu\nu}$	$\mathcal{O}_{WB} = g'gH^\dagger\sigma^a HW^a_{\mu\nu}B^{\mu\nu}$	$\Delta\mathcal{L}_{\kappa_\gamma}$
	$\mathcal{O}_W = \frac{ig}{2}\left(H^\dagger\sigma^a \overleftrightarrow{D}^\mu H\right)D^\nu W^a_{\mu\nu}$	$\mathcal{O}_{HL} = (iH^\dagger \overleftrightarrow{D}_\mu H)(\bar L_L\gamma^\mu L_L)$	$\Delta\mathcal{L}_{g_1^Z}$
LEPI‰	$\mathcal{O}_B = \frac{ig'}{2}\left(H^\dagger \overleftrightarrow{D}^\mu H\right)\partial^\nu B_{\mu\nu}$	$\mathcal{O}'_{HL} = (iH^\dagger\sigma^a \overleftrightarrow{D}_\mu H)(\bar L_L\sigma^a\gamma^\mu L_L)$	$\Delta\mathcal{L}^Z_{eL}$
	$\mathcal{O}_T = \frac{1}{2}\left(H^\dagger \overleftrightarrow{D}_\mu H\right)^2$	$\mathcal{O}_T$	$\Delta\mathcal{L}^Z_\nu$
	$\mathcal{O}_{H\psi_R} = (iH^\dagger \overleftrightarrow{D}_\mu H)(\bar\psi_R\gamma^\mu\psi_R)_{\psi=u,d,e}$	$\mathcal{O}_{H\psi_R}$	$\Delta\mathcal{L}^Z_{uR,dR,eR}$
	$\mathcal{O}_{HQ} = (iH^\dagger \overleftrightarrow{D}_\mu H)(\bar Q_L\gamma^\mu Q_L)$	$\mathcal{O}_{HQ}$	$\Delta\mathcal{L}^Z_{uL}$
	$\mathcal{O}'_{HQ} = (iH^\dagger\sigma^a \overleftrightarrow{D}_\mu H)(\bar Q_L\sigma^a\gamma^\mu Q_L)$	$\mathcal{O}'_{HQ}$	$\Delta\mathcal{L}^Z_{dL}$

Table 2: CP-even dimension-6 operators that affects Higgs physics. We omit dipole operators for fermions and $\mathcal{O}^{ud}_R = y_u^\dagger y_d(i\tilde H^\dagger \overleftrightarrow{D}_\mu H)(\bar u_R\gamma^\mu d_R)$ since they are suppressed by light fermion Yukawas under the MFV assumption. A complete set of operators can be found in Refs. [16, 37].

Higgs Physics. In principle Higgs physics can test all effects induced by the above operators. However, the sensitivity of present measurements of Higgs properties at the LHC is generally smaller than that of EW tests at LEP. For this reason it is useful to identify *Higgs-only* operators that do not contribute to EW precision observables and are therefore the genuine target of the LHC Higgs program, unconstrained by other experiments. Operators of the form $|H|^2\mathcal{O}_{SM}$, with $\mathcal{O}_{SM}$ a SM operator, only contribute to Higgs physics since, according to the classification Eq. (39), $|H|^2$ is an EW singlet that does not induce EW breaking effects. It is instructive to see why: the operator $\mathcal{O}_{GG}$ for instance appears in the effective Lagrangian (with G in non-canonical normalization) as

$$-\frac{1}{4g_s^2}G^A_{\mu\nu}G^{A\mu\nu} + \frac{c_{GG}}{M^2}|H|^2 G^A_{\mu\nu}G^{A\mu\nu} = -\frac{1}{4}\left(\frac{1}{g_s^2} - 2\frac{v^2}{M^2}\right)G^A_{\mu\nu}G^{A\mu\nu} + \frac{c_{GG}}{2M^2}(2vh + h^2)G^A_{\mu\nu}G^{A\mu\nu}\,.$$

The piece proportional to v^2 can be reabsorbed into a redefinition of g_s, which is an input parameter for the SM and has therefore no physical meaning until measured. In our flavor-universal framework we can think of the SM as having 8 input parameters, e.g. g', g, g_s, m_W, $m_{u,d,e}$ and m_h, implying the existence of 8 Higgs-only operators

$$\text{Higgs-only:} \quad \{\mathcal{O}_{BB}, \mathcal{O}_{WW}, \mathcal{O}_{GG}, \mathcal{O}_H, \mathcal{O}_{y_{u,d,e}}, \mathcal{O}_6\} \tag{40}$$

that are explicitly manifest in the Warsaw basis (upper-center box in Table 2). These can be tested at the LHC principally through the following rates

$$h\to\gamma\gamma, \quad h\to Z\gamma, \quad gg\to h \quad h\to ZZ, WW, \quad gg\to\bar t t h, \quad h\to \bar b b, \quad h\to\bar\tau\tau, \quad gg\to hh \tag{41}$$

which contribute to the results from the right panel of Fig. 5, in addition with constraints from the $h\to Z\gamma$ channel [40] and $pp\to hh$ processes (the latter will test the Higgs cubic interaction that is affected uniquely by $\mathcal{O}_6$, experimental results in this channel are not available yet, see Ref. [41] and references therein).

Notice that Higgs physics has in principle many more observables than the free parameters implied by the EFT, meaning that at this level in the $1/M$ expansion the EFT is in fact predictive and provides relations that can be tested. For instance, the EFT operators Eq. (40) imply that only one operator $\mathcal{O}_H$ modifies both the $hZ_\mu Z^\mu$ and $hW^+_\mu W^{-\mu}$ couplings. Similarly, two operators $\mathcal{O}_{BB}, \mathcal{O}_{WW}$ affect the four $hZ_{\mu\nu}Z^{\mu\nu}$, $hW^+_{\mu\nu}W^{-\mu\nu}$, $hZ_{\mu\nu}A^{\mu\nu}$ and $hA_{\mu\nu}A^{\mu\nu}$ structures. These, and many more, are accidental relations that will be broken by the dimension-8 Lagrangian and they can be thought as the defining features of our assumptions on page 14. These relations can be made more explicit in the mass eigenbasis, writing combination of the 8 operators Eq. (40) as

$$\begin{aligned}
\Delta\mathcal{L}^h_{\gamma\gamma} &= \boldsymbol{\kappa_{\gamma\gamma}} \left(\frac{h}{v}+\frac{h^2}{2v^2}\right)\left[A_{\mu\nu}A^{\mu\nu}+Z_{\mu\nu}Z^{\mu\nu}+2W^+_{\mu\nu}W^{-\mu\nu}\right],\\
\Delta\mathcal{L}^h_{Z\gamma} &= \boldsymbol{\kappa_{Z\gamma}} \left(\frac{h}{v}+\frac{h^2}{2v^2}\right)\left[t_{\theta_W}A_{\mu\nu}Z^{\mu\nu}+\frac{c_{2\theta_W}}{2c^2_{\theta_W}}Z_{\mu\nu}Z^{\mu\nu}+W^+_{\mu\nu}W^{-\mu\nu}\right],\\
\Delta\mathcal{L}^h_{GG} &= \boldsymbol{\kappa_{GG}} \left(\frac{h}{v}+\frac{h^2}{2v^2}\right)G^A_{\mu\nu}G^{A\,\mu\nu}\,,\\
\Delta\mathcal{L}^h_{ff} &= \boldsymbol{\delta g^h_{ff}}\Big(h\bar{f}_L f_R+\text{h.c.}\Big)\left(1+\frac{3h}{2v}+\frac{h^2}{2v^2}\right),\\
\Delta\mathcal{L}_{3h} &= \boldsymbol{\delta g_{3h}}\, h^3\left(1+\frac{3h}{2v}+\frac{3h^2}{4v^2}+\frac{h^3}{8v^3}\right),\\
\Delta\mathcal{L}^h_{VV} &= \boldsymbol{\delta g^h_{VV}}\, 2m_W\left[h\left(W^{+\mu}W^-_\mu+\frac{Z^\mu Z_\mu}{2c^2_{\theta_W}}\right)+\Delta_h\right],
\end{aligned} \tag{42}$$

with Δ_h contributing to processes with at least two physical Higgses. Here we have mede sure that, for instance, only $\Delta\mathcal{L}^h_{\gamma\gamma}$ contributes to the $h\to\gamma\gamma$ rate and that only $\Delta\mathcal{L}_{3h}$ modifies the Higgs cubic. Then, these relations imply predictions: for instance the $hZ_{\mu\nu}Z^{\mu\nu}$ structure receives contributions that are proportional to $\boldsymbol{\kappa_{Z\gamma}}$ and $\boldsymbol{\kappa_{\gamma\gamma}}$ that are both well constrained.

This way of writing the EFT operators has the unique purpose of making manifest the relations between modification to different observables that persist in the dimension-6 Lagrangian. It is a way of writing *observables in terms of observables* that is analogous to defining the SM through the relation between W and Z masses Eq. (27), or the relations between Higgs couplings and masses Eqs. (28-30) that are accidental to the dimension-4 Lagrangian. In practice, these *BSM Primaries* [38], provide a useful way to express the results of a global fit in terms of parameters that are as close as possible to experiments, but at the same time maintain the information about the accidental relations of the dimension-6 Lagrangian. A global fit to Run-1 Higgs data reads [42] (see also [43–49])

$$\boldsymbol{\delta g^h_{VV}} = 1.04\pm 0.03,\quad \boldsymbol{\delta g^h_{tt}} = 1.1^{+0.9}_{-3.0}\quad \boldsymbol{\delta g^h_{bb}} = 1.06^{+0.30}_{-0.23},\quad \boldsymbol{\delta g^h_{\tau\tau}} = 1.04\pm 0.22 \tag{43}$$

$$\boldsymbol{\kappa_{gg}} = 0.0005\pm 0.008,\quad \boldsymbol{\kappa_{\gamma\gamma}} = -0.0003^{+0.0005}_{-0.0007},\quad \boldsymbol{\kappa_{Z\gamma}} = 0.000^{+0.035}_{-0.019}\,. \tag{44}$$

This can be written in terms of constraints on the above operators, implying that $(c_H, c_{y_{u,d,e}})\frac{v^2}{M^2}\sim 10^{-1}$, $c_G\frac{m_W^2}{M^2}\sim 10^{-3}$, while $(c_{WW}, c_{BB})\frac{v^2}{M^2}\sim 10^{-2}$ (notice that $c_{WW}\; c_{BB}$ affect both $\gamma\gamma$ and the poorer $Z\gamma$; it is the latter that dominates this marginalized constraint); see also Fig. 6.

LEPI Electroweak Precision Tests (EWPT). A peculiarity of the Higgs field is that it acquires an expectation value that modifies the symmetry of the vacuum and the propagation of the EW bosons. Physics that modifies the Higgs sector can therefore also contribute to observables in EW physics, through the operators of Table 2 with $\langle H\rangle \to v/\sqrt{2}$.

LEP-I, operating on the Z-pole, provides the most precise measurements in this context, reaching the permille level. From an experimental point of view, we can think of these measurements of $e^+e^-\to$

$Z \to fermions$ as testing the couplings of Z to all the 7 SM (left and right) fermions ν, e_L, e_R, u_L, u_R, d_L, d_R (the Z couplings to these define a set of 7 pseudo-observables for LEP-I). Using (α_{em}, m_W, m_Z) as EW input parameters, it is easy to get convinced that there is no additional experimental information that can be extracted from LEP-I on flavor universal theories; in particular, information on custodial symmetry breaking (often presented as T parameter, that measures departures from the m_Z/m_W mass difference of Eq. (27)) or effects from $Z - B$ mixing (the S parameter [50, 51]) is contained in the Z-coupling measurements.

In the SILH basis, there turn out to be 7 operators contributing to these observables[9], reported in the bottom-left block of Table 2 - more precisely, only the combination $\mathcal{O}_W + \mathcal{O}_B$ affects this type physics, while the orthogonal combination $\mathcal{O}_W - \mathcal{O}_B$ does not,

$$\text{LEP-I}: \quad \{\mathcal{O}_W + \mathcal{O}_B, \mathcal{O}_T, \mathcal{O}_{HQ}, \mathcal{O}'_{HQ}, \mathcal{O}_{Hu}, \mathcal{O}_{Hd}, \mathcal{O}_{He}\} \tag{45}$$

There could have been more operators (as in the Warsaw basis, where there naively appear to be 8 operators contributing), implying that one combination will necessarily remain unconstrained and it will affect the results of global fits with marginalized coefficients (see footnote 8). In this sense, the SILH basis appears as a favorable choice for EWPT. The result of a global fit, restricted to these operators and to LEP-I data [17], implies that the coefficients c_i of the operators of Eq. (45) are constrained at the permille level [52–54],

$$c_i \frac{m_W^2}{M^2} \sim few \times 10^{-3} \tag{46}$$

as illustrated in Fig. 6.

LEP-II vs LHC. We have seen that out of the 17 operators that involve H, Table 2, 7 are constrained by LEP-I measurements and are well described in the SILH basis, while other 8 can be tested with Higgs physics only, and are well described in the Warsaw basis. Two (combinations of) operators remain yet unconstrained. To understand what effects they induce, it is useful to think of the 17 operators in Table 2 as a 17-dimensional sub-manifold in the ∞-dimensional space of all possible observables and single out the 2-dimensional plane that belongs to this 17-dimensional manifold, but does not contribute to any observables measured at LEP-I or to the Higgs-only measurements of the Right panel of Fig. 5. The result of this exercise is a continuation of what we began in Eq. (42), expressing the EFT in the mass eigenstate basis and unitary gauge. We find that [38]

$$\begin{aligned}
\Delta\mathcal{L}^V_{ee} &= \boldsymbol{\delta g^Z_{eR}} \frac{\hat{h}^2}{v^2} Z^\mu \bar{e}_R \gamma_\mu e_R + \boldsymbol{\delta g^Z_{eL}} \frac{\hat{h}^2}{v^2} \left[Z^\mu \bar{e}_L \gamma_\mu e_L - \frac{c_{\theta_W}}{\sqrt{2}} (W^{+\mu} \bar{\nu_L} \gamma_\mu e_L + \text{h.c.}) \right] \\
&\quad + \boldsymbol{\delta g^Z_{\nu L}} \frac{\hat{h}^2}{v^2} \left[Z^\mu \bar{\nu}_L \gamma_\mu \nu_L + \frac{c_{\theta_W}}{\sqrt{2}} (W^{+\mu} \bar{\nu_L} \gamma_\mu e_L + \text{h.c.}) \right] ,
\end{aligned} \tag{47}$$

$$\begin{aligned}
\Delta\mathcal{L}^V_{qq} &= \boldsymbol{\delta g^Z_{uR}} \frac{\hat{h}^2}{v^2} Z^\mu \bar{u}_R \gamma_\mu u_R + \boldsymbol{\delta g^Z_{dR}} \frac{\hat{h}^2}{v^2} Z^\mu \bar{d}_R \gamma_\mu d_R \\
&\quad + \boldsymbol{\delta g^Z_{dL}} \frac{\hat{h}^2}{v^2} \left[Z^\mu \bar{d}_L \gamma_\mu d_L - \frac{c_{\theta_W}}{\sqrt{2}} (W^{+\mu} \bar{u}_L \gamma_\mu d_L + \text{h.c.}) \right] \\
&\quad + \boldsymbol{\delta g^Z_{uL}} \frac{\hat{h}^2}{v^2} \left[Z^\mu \bar{u}_L \gamma_\mu u_L + \frac{c_{\theta_W}}{\sqrt{2}} (W^{+\mu} \bar{u_L} \gamma_\mu d_L + \text{h.c.}) \right]
\end{aligned} \tag{48}$$

$$\Delta\mathcal{L}_{g_1^Z} = \boldsymbol{\delta g_1^Z} \Big[igc_{\theta_W} \Big(Z^\mu (W^{+\nu} W^-_{\mu\nu} - \text{h.c.}) + Z^{\mu\nu} W^+_\mu W^-_\nu \Big) \tag{49}$$

[9] This is true when performing EW fits using (α_{em}, m_W, m_Z) as EW input parameters; if instead one uses (α_{em}, G_F, m_Z) there is an additional non-Higgs operator $(\bar{L}_L \sigma^a \gamma^\mu L_L)(\bar{L}_L \sigma^a \gamma_\mu L_L)$ that enters the fit at dimension-6 level, but also has an additional measurement that constrains it: the muon lifetime is used to extract G_F.

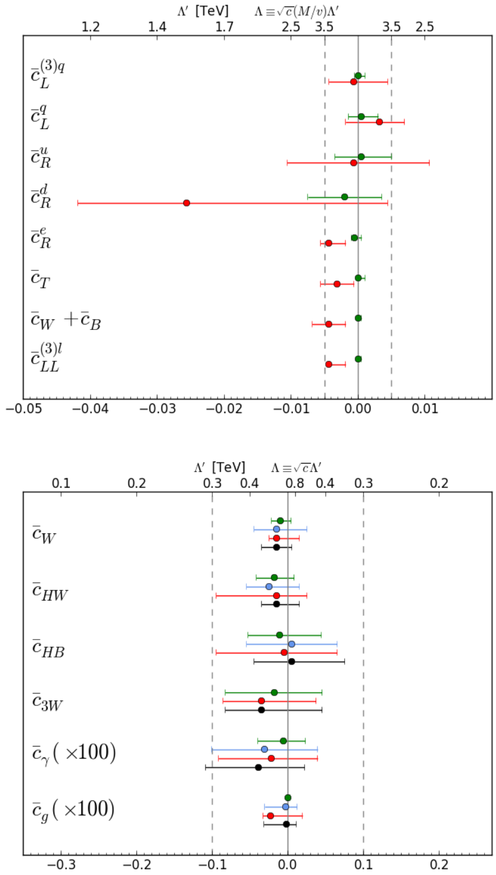

Fig. 6: 95% confidence intervals from a global fit, from Ref. [54]. The green lines denote fits with one coefficient only, while red bars denote fits with multi coefficients, marginalized (in our notation $c^{\psi}_{R,L}$ is $c_{H\psi}$).

$$-2gc^2_{\theta_W}\frac{h}{v}\left(W^-_\mu J^\mu_- + \text{h.c.} + \frac{c_{2\theta_W}}{c^3_{\theta_W}}Z_\mu J^\mu_Z + \frac{2s^2_{\theta_W}}{c_{\theta_W}}Z_\mu J^\mu_{em}\right)\left(1+\frac{h}{2v}\right) + \frac{e^2 v}{2c^2_{\theta_W}}hZ_\mu Z^\mu + \Delta_{h^2}\Bigg]$$

$$\Delta\mathcal{L}_{\kappa_\gamma} = \delta\kappa_\gamma\Bigg[ie\left(1+\frac{h}{v}\right)^2(A_{\mu\nu}-t_{\theta_W}Z_{\mu\nu})W^{+\mu}W^{-\nu} + 2Z_\nu\partial_\mu\left(\frac{h}{v}+\frac{h^2}{2v^2}\right)(t_{\theta_W}A^{\mu\nu}-t^2_{\theta_W}Z^{\mu\nu})$$
$$+\left(\frac{h}{v}+\frac{h^2}{2v^2}\right)\left(t_{\theta_W}Z_{\mu\nu}A^{\mu\nu}+\frac{c_{2\theta_W}}{2c^2_{\theta_W}}Z_{\mu\nu}Z^{\mu\nu}+W^+_{\mu\nu}W^{-\mu\nu}\right)\Bigg], \tag{50}$$

complete Eq. (42) to spam the entire space of dimension-6 operators, in a way that aligns with EW and Higgs observables. Here h corresponds to the physical Higgs and Δ_{h^2} includes interactions with at least two h that are irrelevant for experiments in the near future.

From Eqs. (49,50) we can read that the two remaining directions modify the trilinear gauge couplings (TGCs) between ZW^+W^- and γW^+W^-, as well as the $hV\bar{\psi}\psi$ ($V = W, Z$) amplitude. The reason that both these observables appear simultaneously modified can be traced to the fact that Higgs and the eaten Goldstones belong to the same multiplet, so that some BSM deformations in the Higgs sector can modify also processes with longitudinal W, Z bosons, at least at a fixed order in the $1/M$ expansion. TGCs can be tested in diboson processes at LEP-II or at the LHC, while deformation in $hV\bar{\psi}\psi$ are tested in $pp \to VH$ associated production at the LHC.[10] The constraints from these two search modes

[10] In principle the same amplitude can be tested in $h \to V\bar{\psi}\psi$ decays [55] but, as shown in Fig. 7, constraints from LEP-II

are at present comparable, although the latter are typically extracted from the high-energy regime [56], so that certain conditions regarding the validity of our EFT assumption limit their interpretability [57]. From LEP-II data [58] we read

$$\delta g_{1,Z} = -0.05^{+0.05}_{-0.07}, \quad \delta\kappa_\gamma = 0.05^{+0.04}_{-0.04}\,. \tag{51}$$

These directions Eqs. (49,50) can be translated into combinations of operators in the SILH or Warsaw basis (see Ref. [39] for a gauge-invariant formulation in terms of the above operators), where $\delta g_{1,Z}, \delta\kappa_\gamma \sim c\, m_W^2/M^2$, for c a combination of Wilson coefficients, and are included in Fig. 6.

To conclude, let me reiterate the arguments of this section. Global fits are useful to explore the impact of experimental data on broad BSM hypotheses, that in our case we have defined with the assumptions in page 14 that we hope are able to capture large classes of BSM theories. If the theoretical parameters are not aligned with the experimental observables, the results of a global fit with n parameters, $n-1$ of which are normalized, will be dominated by the poorest experiment. For this reason we have divided experiments with different sensitivity in correlated blocks (LEP-I, LEP-II and Higgs physics) and identified the operators they constrain. In fact, the best way to do this is to align the theoretical parameters to the most precise experiments, as in the BSM primaries/Higgs basis.

The outcome of this discussion is two-fold. First of all, it allowed us to identify the relations that persist in the effective Lagrangian at the level of dimension-6 operators, relations that can be tested or used as a check, if a deviation is found, or can be used to better constraint a given parameter. An example of such relation is reported in Fig. 7, which shows the differential distribution of Higgs decaying to a vector and a pair of fermions, such as in the golden channel. This distribution is affected by many dimension-6 operators, but all of them are bounded by other experiments, either at LEP-I, LEP-II or in Higgs physics; therefore this distribution cannot be arbitrarily modified, and a prediction of the dimension-6 EFT analysis is that, if any deviations are found there, they must be within the blue band, or violate some of our assumptions.

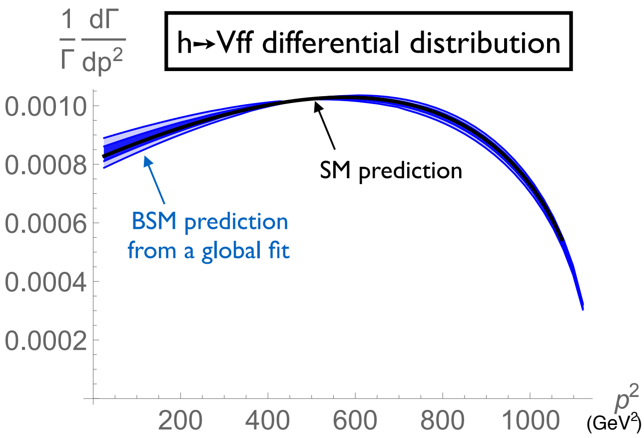

Fig. 7: Normalized differential distribution for Higgs decays in the golden channel. The black line corresponds to the SM prediction, while the blue bands correspond to the (1,2-σ contours) of the BSM-deformation, as allowed by constraints from a global fit on possible dimension-6 modifications, that includes LEP-I and LEP-II data, as well as constraints on $h \to Z\gamma$. From the data in [52] (see also [59]).

are stronger than what can currently be measured in that channel

Secondly, this discussion, shows the most that we can get, in terms of experimental constraints, out of the *most conservative* hypotheses that all operators are democratically generated, and that they can even cancel each other in their contribution to a given observable. Even in this limiting hypothesis we could identify classes of operators that are already very well tested and constrained.

4.2 BSM Perspective

In realistic set-ups, based on physical microscopic hypotheses, only subsets of operators are typically generated, thus providing concrete scenarios that can be tested more precisely, and from which something can be learned. There is no limit in how specific and complicated a BSM model can be, but there is instead a limit in how simple it can be. Here we want to identify the minimal set of ingredients that can characterize a microscopic model and that can, at the same time, represent broad features of more generic BSM scenarios. These ingredients are,

One New Mass Scale M, *One New Coupling* g_*, *Symmetries and Selection Rules*[11]

where we assume that the UV completion admits some perturbative expansion in its couplings. These microscopic properties will be imprinted into the Wilson coefficients of the operators of Table 2. In lack of a specific model in which to compute this UV→IR matching, we can still estimate them through a procedure known as *power-counting* or Naive Dimensional Analysis [60].

The idea is simple. Symmetries or properties of the underlying theory determine if an operator is generated or not. In weakly coupled theories this can simply boil down to whether or not the field mediating an interaction is present at the scale M (an example of this will be given below for SUSY), and can make the difference between large (tree-level) and small (loop-level) effects [61], while in strongly coupled theories a symmetry is typically necessary to generate a selection rule in the IR (for instance in composite Higgs models the Higgs is a psuedo-Goldstone boson with non-linearly realized global symmetry that generates mostly interactions with derivatives ∂H).

The dimension of the operator D determines that its coefficient will scale $\sim 1/M^{D-4}$, in order to make the action dimensionless, as we already know. What is perhaps more exotic is the fact that, counting powers of $\hbar \neq 1$ in the Lagrangian can tell us how many powers of couplings an operator might carry. Indeed, couplings, as well as fields, carry $\hbar$ dimensions [35] (see also Refs. [9, 10]). It is easy to see that, since $[\mathcal{S}] = \hbar$, fields have $[\phi] = \hbar^{1/2}$, while couplings have $[g] = \hbar^{-1/2}$. Any operator in the Lagrangian must have dimension $[c_i \mathcal{O}_i] = \hbar$, and we find that for an operator with n_i fields,

$$c_i \sim (\text{coupling})^{n_i - 2} \tag{52}$$

In what follows we clarify the importance of power-counting through two examples ins SUSY and CH models, the most studied extensions of the Higgs sector that solve the hierarchy problem. Keep in mind that the selection rules discussed here, apply at the matching scale M, but the coefficients will run as they evolve towards lower energy. These effects can be computed entirely within the EFT, see [16, 62–66], and can have some important implications, in particular in situations where a poorly measured Wilson coefficient mixes through renormalization group flow into a very-well measure one.

4.2.1 SUSY

SUSY provides a great example of symmetries and selection rules in action. Superymmetry is compatible only with holomorphic interactions, thus forbidding the up-quark Yukawa $\propto \widetilde{H} = \epsilon H^\dagger$ in the SM Lagrangian Eq. (23). Up-quark masses require therefore the presence of an additional Higgs doublet

[11] Selection rules follow from microscopic symmetries that might or might not be realized (linearly or not) in the IR. Here I tend to refer to selection rules, as those that cannot be identified with symmetries from the low-energy point of view.

with opposite hypercharge $H_2 \in (\mathbf{1}, \mathbf{2}, 1/2)$, which can appear in the Lagrangian without complex conjugation. [12]

Moreover, the symmetry implies that the SM fields are part of super-multiplets containing fields of different spin. Some of these fields can potentially mediate proton decay, so that many SUSY incarnations invoke an accidental symmetry, R-parity, to forbid the relevant interactions responsible for this. Under R, the SM fields that we know, together with the scalar component of H_2, are even, while all other states appearing in the SUSY multiplets are odd. This last observation is uniquely responsible for the EFT structure of SUSY models with R-parity. Indeed, any process between SM (R-even) states, can not be mediated at tree-level by an R-odd resonance, so that the only tree-level BSM effects are those associated with H_2, the only R-even BSM particle. We show the relevant diagram in Fig. 8, where we identify h with the linear combination of H_1, H_2 that is closer to the SM one, and call H the heaviest combination (see Ref. [68]).

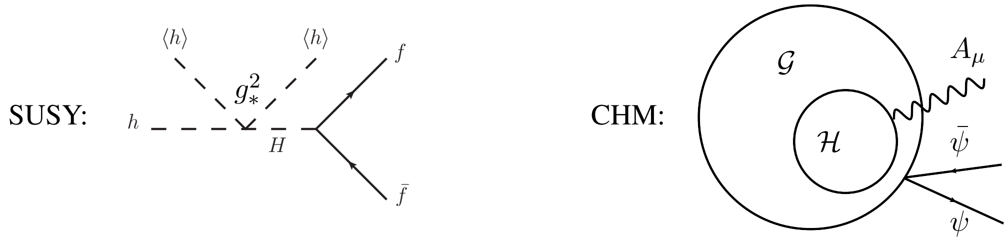

Fig. 8: LEFT: a modification of the Higgs couplings to SM fermion is the largest effect from integrating out BSM physics in SUSY models with R-parity. RIGHT: in CHM a strong sector brakes a global symmetry $\mathcal{G} \to \mathcal{H}$, while the gauge and Yukawa couplings break the symmetries explicitly, but weakly.

So, from the EFT perspective, R-symmetry translates in the *selection rule* that only a very special subset of new interactions is generated sizably at low energy: $\mathcal{O}_{y_u}, \mathcal{O}_{y_d}, \mathcal{O}_{y_e}$. What else can we say about this 'New Physics' sector? Flavor *symmetry* $U(3)^5$, is broken in the SM by the Yukawa couplings; the stringent constraints from flavor physics favor the possibility that the Yukawas be the only source of flavor symmetry breaking also in TeV-scale BSM: a possibility that is known as Minimal Flavor Violation [69] and is indeed realized in these SUSY models (the $H\bar{\psi}\psi$ couplings are aligned with the SM Yukawas). Finally, we can identify the microscopic coupling

$$\mathcal{L}_{UV}^{\text{SUSY}} \supset \frac{g_*^2}{4} h^3 H \tag{53}$$

with the generic g_* introduced on page 21, and m_H with M.

Now that we have identified the *relevant*[13] features of SUSY models, in the generic language of page 21, we can estimate the low-energy EFT:

$$\text{SUSY} \sim \left\{ \begin{array}{l} \text{Selection rule from R-symmetry} \\ \text{Flavor symmetry, broken by } y_\psi \\ \text{New coupling: } g_*^2 \\ \text{Mass scale: } m_H \end{array} \right\} \Rightarrow \mathcal{L}_{eff}^{\text{SUSY}} = \tilde{c}_u \frac{g_*^2}{m_H^2} \mathcal{O}_{y_u} + \tilde{c}_d \frac{g_*^2}{m_H^2} \mathcal{O}_{y_d} + \tilde{c}_e \frac{g_*^2}{m_H^2} \mathcal{O}_{y_e} \tag{54}$$

where we have used the fact that the operators $\mathcal{O}_{y_\psi}$ contain 5 fields so we expect their coefficient $\sim$ $coupling^3$; since they brake flavor symmetries they must involve y_ψ (they are already weighted by one

[12] One of my favorite possibilities is that instead up-type quark Yukawas arise as SUSY breaking effects, while the Higgs doublet, which has quantum numbers $H_2 \in (\mathbf{1}, \mathbf{2}, -1/2)$ be the scalar supersymmetric partner of one of the leptons L, which has the same quantum numbers. In this case h would be the sneutrino [67], implying that we would have already discovered SUSY!

[13] On the most technical meaning of the word: features that are important at low-energy.

power of the Yukawas in Table 2), thus Eq. (52) reads here $coupling^3 \sim y_\psi g_*^2$. An explicit computation reproduces Eq. (54), with

$$\tilde{c}_u = -\cot\beta\,, \quad \tilde{c}_d = \tilde{c}_e = \tan\beta\,. \tag{55}$$

This example show the impressive power of EFTs: 3 parameters are enough to capture the low-energy signatures of all SUSY models with R symmetry and minimal field content. It also shows the importance of being able to identify the relevant microscopic features that shape the Wilson coefficients at the matching scale, as in Eq. (54), for two reasons. Certainly, these power-counting rules provide a useful short-cut for the $BSM \to SM\,EFT$ matching. Most importantly, however, they allow us to identify which are the relevant hypotheses that we are actually testing when we use the EFT to parametrize SM precision tests; all other features of specific BSM models being *de facto* irrelevant. For this reason, we can refer to the assumptions in Eq. (54), as broad, in the sense that they are not specific to a single model.

4.2.2 Composite Higgs Models

Models of Composite Higgs solve the hierarchy problem in two steps. First of all they postulate a compositeness scale $M \sim$ few $\times$ TeV $\ll M_{Planck}$ that is naturally generated, e.g., by dimensional transmutation, like the QCD scale. As explained before, above M, $|H|^2$ is an irrelevant operator and hence small. A more physical picture is simply that the Higgs H is a composite particle that exists only in the low-energy EFT: the contribution to m_h^2 from loops of high virtuality is tamed when the particles in the loop probe Higgs compositeness. This would still imply $m_h \sim M$ naturally. As we will see below, EWPT constrain $M \gtrsim 3$ TeV, thus creating a *little hierarchy problem*, that is solved if the H is also an approximate Nambu Goldstone boson of a spontaneously broken global symmetry $\mathcal{G}/\mathcal{H}$. Of course this is *natural*, as it mimics the pions of QCD, which have been observed in nature.

There are many explicit realizations of these models, but the gross picture is common to all of them. A strong sector confines at the scale M and brakes a symmetry $\mathcal{G} \to \mathcal{H}$, delivering at least 4 massless NGBs at low energy, and nothing else.[14] The $\mathcal{G}/\mathcal{H}$ symmetry is broken explicitly by how this sector couples to the SM: the EW gauge couplings g, g' are associated to the gauging of only a subgroup $SU(2)_L \times U(1) \subset \mathcal{H}$ (therefore breaking H explicitly), and the SM Yukawas always break $\mathcal{G}$ when the SM fermions are coupled to the strong sector. A pictorial representation is given in Fig. 8.

These ingredients are enough to estimate the low-energy EFT of composite Higgs models:

$$\text{CHM} \sim \left\{ \begin{array}{l} \mathcal{G}/\mathcal{H} \text{ symmetry, broken by } g, g', y \\ \text{New coupling in } H\text{-sector: } g_* \\ \text{Resonances mass scale: } M \end{array} \right\} \Rightarrow \mathcal{L}_{eff}^{\text{CHM}} = \frac{M^4}{g_*^2} L\left(\frac{g_* H}{M}, \frac{D_\mu}{M}, \frac{g F_{\mu\nu}}{M^2}, \frac{\lambda_\psi \psi}{M^{3/2}}\right) \tag{56}$$

where $\mathcal{G}/\mathcal{H}$-preserving Higgs interactions must be compatible with the goldstone symmetry and be functions of the CCWZ building blocks [72], that at the leading order read the same for all compact cosets,[15]

$$d_\mu = \partial_\mu H + \cdots\,, \qquad \epsilon_\mu = H \overleftrightarrow{D}_\mu H + \cdots \tag{57}$$

More precisely, the Lagrangian for a Strongly Interacting Light Higgs reads [35],

$$\begin{aligned} \mathcal{L}_{\text{SILH}} \;=\; & \frac{\tilde{c}_H g_*^2}{M^2}\mathcal{O}_H + \frac{\tilde{c}_T g_*^2}{M^2}\mathcal{O}_T - \frac{\tilde{c}_6 g_*^2}{M^2}\mathcal{O}_6 + \left(\frac{\tilde{c}_y g_*^2}{M^2}\mathcal{O}_{y_\psi} + \text{h.c.}\right) + \frac{\tilde{c}_W}{M^2}\mathcal{O}_W + \frac{\tilde{c}_B}{M^2}\mathcal{O}_B \\ & + \frac{\tilde{c}_{HW}}{M^2}\mathcal{O}_{HW} + \frac{\tilde{c}_{HB}}{M^2}\mathcal{O}_{HB} + \frac{\tilde{c}_{BB}}{M^2}\mathcal{O}_{BB} + \frac{\tilde{c}_g}{M^2}\mathcal{O}_{GG}. \end{aligned} \tag{58}$$

[14] Non-minimal models with extended cosets predict more states in the IR, $SO(6)/SO(5)$ for instance includes an additional singlet in the light spectrum, that can in principle play important phenomenological rôles, such as being DM [70,71].

[15] An interesting limiting case is that of $ISO(4)/SO(4)$: here the coset manifold is flat and $\epsilon_\mu = 0$, so that the first strong interactions involving H arise at dimension-8 [30].

This defines the basis the SILH basis previously introduced and explains the normalization of operators $\mathcal{O}_B, \mathcal{O}_{HB} \sim g'$, $\mathcal{O}_W, \mathcal{O}_{HW} \sim g$, while $\mathcal{O}_{BB} \sim g'^2$ and $\mathcal{O}_{GG} \sim {g_s}^2$, appearing in Table 2. Also, $\mathcal{O}_6$ shares the same symmetry as the Higgs quartic $\lambda|H|^2$: this is why we expect $\mathcal{O}_6 \sim \lambda$.

Further UV hypotheses can be easily translated into selection rules for SILH Wilson coefficients. For instance, both coset structures $SU(3)/SU(2) \times U(1)$ and $SO(5)/SO(4)$ are minimal, in the sense they only have 4 NGB degrees of freedom that can be identified with H. Yet they differ in that the latter case $\mathcal{H} = SO(4) \simeq SU(2)_L \times SU(2)_R$ contains custodial symmetry, implying that operators that transform non-trivially under$SU(2)_c$ must be suppressed. In the SILH basis we find,

$$\text{Custodial symmetry} \quad \Rightarrow \quad \tilde{c}_T = 0\,. \tag{59}$$

In many known theories (including weakly coupled 5D theories, and their holographic strongly coupled duals), the dominant effects come from integrating out particles of spin≤ 1 at tree-level, while other effects only arise at loop-level. This hypothesis, called *Minimal Coupling*, implies a further suppression

$$\text{Minimal Coupling} \quad \Rightarrow \quad \tilde{c}_{HW}, \tilde{c}_{HB} \sim \frac{g_*^2}{16\pi^2} \quad \tilde{c}_{BB}, \tilde{c}_{GG} \sim \frac{g_{SM}^2}{16\pi^2} \tag{60}$$

where we have taken into account that, for $\mathcal{O}_{BB}, \mathcal{O}_{GG}$ the couplings to such particles must also break the shift-symmetry and are therefore typically suppressed by a symmetry breaking SM coupling $g_{SM} = g', y_t, \cdots$.

It is interesting to compare the power-counting of Eqs. (58,59,60) with the experimental observations from the previous section. From Fig. 6 we see that the best constraints are on c_{BB} and c_{GG}, that are predicted suppressed by several powers of the weak couplings, and on $c_W + c_B$ that are instead expected $O(1)$, implying

$$M \gtrsim 2.5\,\text{TeV}\,. \tag{61}$$

Despite these stringent constraints, we notice that many Higgs-only operators are g_* enhanced allowing for seizable effects in Higgs physics.[16] A careful study of the Higgs properties will tell us more about these types of models.

We have seen the example of a weakly coupled BSM scenario, where loop-effects were suppressed, so that the dynamical field content and interactions (in this case dictated by R-parity) have an important impact on the low-energy EFT (also the minimal coupling assumption Eq. (60) relied on such weakly coupled picture). In other words, the EFT for weakly coupled BSM models is rather model-dependent.

On the other hand, in strongly coupled BSM scenarios, *everything that is not forbidden is compulsory*, as it might be generated by unsuppressed loop effects involving the strong coupling. In these scenarios, our power-counting rules that identify weak and strong couplings and symmetries, are not only useful, but rather necessary, as the underlying theory is incalculable. In these conditions, broad assumptions about the UV are enough to determine the resulting EFT, independently of the microscopic details.

There are many, more or less specific, assumptions that can characterize physics BSM and that can be captured by power-counting rules like those mentioned above. For instance, New physics can couple to the SM bosons only (this is sometimes called *universal*) [73] or only to the top quark [74], that plays the most important rôle in terms of loop effects to the Higgs mass parameter, or only to the transverse components of vector-bosons [30]. It is important to keep in mind that, when testing a specific property of the Higgs boson, we are specifically looking towards one of these specific BSM directions.

[16]This is in fact a chicken-egg situation: these models were thought in the LEP era with knowledge of EW constraints.

4.2.3 EFT Validity

Our complete discussion so far was based on the existence of a scale separation

$$E \ll M\,. \tag{62}$$

An experiment operating at energy E_{exp}, can provide a constraint (or measurement) on the combination

$$\frac{c_i}{M^2} < \frac{\delta_{exp}}{E_{exp}^2} \tag{63}$$

on the effects of an operator $\mathcal{O}_i$, with precision δ^{exp}. From this we can say that our original assumption Eq. (62) is indeed satisfied if the experimental precision $\delta_{exp} = c_i E_{exp}^2/M^2 \ll c_i$. Vice versa, we can say that our measurement can be consistently interpreted in the EFT context, only in theories with

$$c_i \gg \delta_{exp}\,. \tag{64}$$

We have seen in the above examples that the Wilson coefficients c_i can vary enormously depending from the UV structure: they can be enhanced by a strong coupling or reduced by loop factors, and even vanish. Therefore, it is fair to say that:

There is no model-independent discussion about the EFT validity.

In this context, the power-counting arguments outlined above become particularly useful, because they allow us to identify the broadest features that can make a Wilson coefficient large or small, so that the question of whether the EFT provides a consistent interpretation of our measurement, can be answered in the most generic (less model-dependent) terms.

At the LHC, some experiments (e.g $2 \to 2$ scattering processes) are testing a large range of energies, and the question of whether Eq. (62) is satisfied becomes more subtle, as E_{exp} is in principle unknown. This can be obviated in a number of ways. The most systematic is to perform an additional cut $\sqrt{s} < M_{cut}$ at the level of the analysis on the center-of mass energy of the system.[17] This procedure provides the necessary information on E_{exp} that is now bounded from above by M_{cut} and allows to discuss the EFT validity [57, 77].

In these high-E processes, EFT effects have the common property that they *grow* with some power of the energy, relatively to the SM. On the other hand, the precision of measurements of deviations from the SM, *decreases* with energy, principally due to the rapid fall-off in number of events at high-energy (since parton distribution functions decrease exponentially fast). In some instances this implies that the constraint is dominated by an energy region with little sensitivity, implying that only departures from the SM that exceed the SM itself can be tested,

$$\delta_{exp} = \frac{\delta\sigma}{\sigma_{SM}} \gtrsim 1\,. \tag{65}$$

This is often considered a limitation to test EFTs at hadronic colliders in high-energy processes, as it implies from Eq. (64) that such experiments can not be interpreted in theories with $c_i \lesssim 1$. This corresponds to weakly coupled theories that fill a special place in BSM model-building, as they are calculable and well under control. In this sense $\delta_{exp} \lesssim 1$ can be thought as a target for experiments, that opens the door to interpret their results in a wider and well-motivated context [77]. Yet this doesn't mean that high-energy low-resolution experiments that test EFT are necessarily inconsistent: they provide useful information about strongly coupled theories, where, e.g. $c_i \sim g_*^2 \gg 1$.

[17] In some systems, this might not be known, but a consistent analysis can still be performed along the lines of Refs. [75, 76].

5 Conclusions

I like Higgs physics because it is at the frontier of our exploration of the universe at small distances. As such, it can potentially hide information about structure beyond the SM. The hierarchy problem suggests the existence of such structure, that would imply modification of the Higgs properties. These can be studied in the formalism of EFTs and their power-counting, a dictionary that allows to recognize the relevant ingredients in microscopic theories and read their effects at low-energy.

EFTs can be thought as a structured and well motivated context to perform SM precision tests, the result of which gives us quantitative information on how well we know the SM, in terms of how strong are the constraints on certain classes of theories beyond the SM.

At the same time, EFTs, accompanied with their BSM perspective, provide an important search tool that extends the reach of the LHC beyond its direct reach. This is particularly true for strongly coupled BSM scenarios, that might induce large effects in low-energy processes, despite the scale of new physics being beyond kinematic reach.

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