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Through precision straits to next standard model heights

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Abstract

After the LHC Run 1, the standard model (SM) of particle physics has been completed. Yet, despite its successes, the SM has shortcomings vis-à-vis cosmological and other observations. At the same time, while the LHC restarts for Run 2 at 13 TeV, there is presently a lack of direct evidence for new physics phenomena at the accelerator energy frontier. From this state of affairs arises the need for a consistent theoretical framework in which deviations from the SM predictions can be calculated and compared to precision measurements. Such a framework should be able to comprehensively make use of all measurements in all sectors of particle physics, including LHC Higgs measurements, past electroweak precision data, electric dipole moment, $g-2$, penguins and flavor physics, neutrino scattering, deep inelastic scattering, low-energy e^+e^- scattering, mass measurements, and any search for physics beyond the SM. By simultaneously describing all existing measurements, this framework then becomes an intermediate step, pointing us toward the next SM, and hopefully revealing the underlying symmetries. We review the role that the standard model effective field theory (SMEFT) could play in this context, as a consistent, complete, and calculable generalization of the SM in the absence of light new physics. We discuss the relationship of the SMEFT with the existing kappa-framework for Higgs boson couplings characterization and the use of pseudo-observables, that insulate experimental results from refinements due to ever-improving calculations. The LHC context, as well as that of previous and future accelerators and experiments, is also addressed.

Keywords: Standard Model, Beyond Standard Model, Effective Field Theory, Radiative Corrections, Higgs Physics, Electroweak Precision Data
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1. The Higgs boson

During the LHC Run 1 a new resonance was discovered in 2012 [1, 2]. That resonance, with a mass of 125 ± 0.24 *GeV* [3], is a candidate to be the Higgs boson of the standard model (SM). The spin-0 nature of the resonance is well established [4–6], all the available studies on the couplings of the new resonance conclude it to be compatible with the Higgs boson of the SM within present precision [7, 8], and, as of yet, there is no direct evidence for new physics phenomena beyond the SM (BSM).

Inevitably, after the LHC Run 1 results comes a need for a better understanding of the current "we haven't seen anything (yet)" theoretical *zeitgeist*. Is the SM with a 125 *GeV* Higgs boson the final theory, or indeed can it be? The

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associated problems with the SM are known and include the neutrino masses as well as cosmological evidence for dark matter.

The discovery of a scalar resonance and the absence of direct evidence for new physics forces us to change perspectives and to redefine the problem. In this review, the starting point is to assume quantum field theory (QFT) as the framework with which to study the basic constituents of matter. The parameters of QFT Lagrangians describe the dynamics, something that is at the heart of the needed change of perspective. At LEP, the dynamics were fixed by the SM Lagrangian, with the unknowns being parameters such as the Higgs mass $M_{\rm H}$, the strong coupling constant $\alpha_{\rm s}(M_{\rm Z})$, etc. [9]. In other words, at LEP the SM was the hypothesis and bounds on $M_{\rm H}$ were derived from a comparison with high-precision data. At the LHC, after the 2012 discovery, the unknowns are deviations from the SM, given that the SM is fully specified and constrained by experimental measurements of increasing precision and accuracy. The definition of SM deviations requires a characterization of the underlying dynamics. Whereas (concrete) BSM models represent specific roads toward the Planck scale, it would be of great interest to employ a (more) model-independent approach, a framework that could describe a whole class of paths to the Planck scale.

While studies performed with limited precision may only claim the discovery of a SM-like Higgs boson, as soon as greater precision is available, it may be possible to decipher the nature of the Higgs through the accurate determination of its couplings [10–13].

Given the precision that was expected for LHC Run 1 results, it was natural to begin exploring the couplings using the (original) κ-framework [14, 15]. There is no need to repeat here the main argument, of splitting and scaling different loop contributions in the amplitudes of processes mediated by Higgs bosons. The main shortcoming is that the original κ-framework is only an intuitive language that lacks internal consistency when moving beyond leading order (LO). In parallel, recent years have witnessed an increasing interest in Higgs effective Lagrangians and SM effective field theory (EFT); see in particular Refs. [16–18], Refs. [19–25], Refs. [26, 27], Ref. [28], Ref. [29], Ref. [30], Refs. [31, 32], Ref. [33], Refs. [11, 34, 35] and Refs. [36–41]. EFTs can be used to describe the full set of deviations from the SM and therefore a better name is certainly SMEFT, as used in Ref. [42, 43] and Refs. [19, 44].

It is worth noting that there is no formulation which is completely model-independent and the SMEFT, as any other approach, is based on a given set of (well-defined) assumptions. In full generality we can distinguish a top-down approach (model-dependent) and a bottom-up approach (with fewer assumptions). The top-down approach is based on several steps. First one has to classify BSM models, possibly respecting custodial symmetry and decoupling of high mass states, then the corresponding SMEFT can be constructed, e.g. via a covariant derivative expansion [42]. Once the SMEFT is derived one can construct the corresponding SM deviations, that may be different for each BSM model or class of BSM models. The bottom-up approach starts with the determination of a basis of $\dim = 6$ (or higher) operators and proceeds directly to the classification of SM deviations, possibly respecting the analytic structure of the SM amplitudes. The synthesis is that $\dim = 6$ operators are supposed to arise from a local Lagrangian, containing heavy degrees of freedom decoupled from the presently-probed energy scales. Of course, the correspondence between Lagrangians and effective operators is not bijective because different Lagrangians can give rise to the same operator.

The change of perspective after the LHC Run 1 is equivalent to saying that we have moved from a fully predictive (SM) phase to a "partially predictive (fitting)" one. The predictive phase is defined as follows: in any (strictly) renormalizable theory with n parameters we need to match n data points, and the $(n+1)$-th calculation is a prediction, e.g. as can be done in the SM. In the fitting (partially predictive) phase there will be $(N_6+N_8+\cdots=\infty)$ renormalized Wilson coefficients to be fitted, e.g. by measuring the SM deformations due to a single $\mathscr{O}^{(6)}$ insertion. This represents a departure from the use of a strictly renormalizable theory, with the compromise of gaining, order-by-order, the ability to explore deviations that can only be constrained by fitting to data. As the number of parameters increases it becomes inevitable that only combinations of the parameters can be constrained.

There is a conceptual difference between Higgs physics at the LHC, for which the UV completion is unknown, and other scenarios where EFT techniques are applied and for which there are known UV completions. When the UV completion is known, we consider a theory with both light and heavy particles; the Lagrangian is $\mathscr{L}(m)$ where m is the mass of the heavy degree of freedom. Next, we introduce the corresponding $\mathscr{L}_{\rm eff}$, the effective theory valid up to a scale $\Lambda = m$. We renormalize the two theories, say in the $\overline{\rm MS}$-scheme, taking care that loop-integration and heavy

limit are operations that do not commute, and impose matching conditions among renormalized "light" one-particle irreducible (1PI) Green's functions.

When we compare the present situation with the past an analogy can be drawn. Consider the QED Lagrangian and complement it with dim $= 6$ Fermi operators $\bar{e}_L \gamma^\mu e_L \bar{e}_L \gamma_\mu e_L$, etc.. This EFT can be used to study the muon decay but also $\nu_e(\nu_\mu)e$ scattering in the approximation of zero momentum transfer. Using data on $\sigma_{\bar{\nu}_l e}/\sigma_{\nu_l e}$, one can derive predictions for the Z couplings [45], e.g. for the ratio g_V^e/g_A^e. In principle, one could have realized the possibility of having neutral currents. Understanding that the Yang-Mills theory could match this EFT at very low energy scales took longer [46], and pretending to use this theory to describe the Z-lineshape is not feasible as the Z boson mass is beyond the validity of this EFT.

One could ask: would there be a way to take the Fermi theory and show how this theory would have pointed to massive vector bosons? The answer is yes, due to unitarity violations at large energies (pure S-wave unitarity); for instance, with intermediate vector bosons the $\nu_e e^- \to \nu_e e^-$ scattering is better behaved as it is no longer a pure S-wave process. For $\nu_\mu e^- \to \nu_e \mu^-$ scattering, unitarity applied to the $l = 0$ LO partial wave requires that $E_{cm} < (\pi/2\sqrt{2}G_F)^{1/2} \approx 310\ GeV$. Furthermore, the interaction had a well-known structure, e.g. in neutron decay, muon decay, and neutrino events, that strongly suggested the existence of massive spin-1 particles. In hindsight, the Fermi Lagrangian could have been built from symmetries (of the SM) only, i.e. left-handed leptons are doublets under $SU(2)$ and flavor universality. In that case the Fermi theory contains the only dim $= 6$ operator with a charged current (CC).

Retrospectively one could have written

$$\mathscr{L}_F = G_F\, \overline{\psi}\, \psi\, \overline{\psi}\, \psi = \sum_i \overline{\psi}_p\, O_i\, \psi_n \left[C_i\, \overline{\psi}_e\, O_i\, \psi_\nu + C_i'\, \overline{\psi}_e\, O_i\, \gamma^5\, \psi_\nu \right] + \text{h. c.}\,, \tag{1}$$

where the O_i refer to scalar, ..., tensor structures, and extended it to become

$$\mathscr{L}_F = G_F\, \overline{\psi}\, \psi\, \overline{\psi}\, \psi + a_\Box\, G_F^2\, \overline{\psi}\, \psi \Box\, \overline{\psi}\, \psi + \ldots \tag{2}$$

add counterterms, making it possible for the theory to become predictive at the loop level.

Historically, events went differently: charged currents were measured to be flavor universal, parity violation was discovered, the V-A structure detected, the $SU(2)$ symmetry was postulated, and neutral currents (NC) predicted; finally NCs were discovered and the SM made its success.

The SMEFT used so far is based on several assumptions: one Higgs doublet with a linear representation (for non-linerar see Ref. [47]), no new "light" degrees of freedom and decoupling of heavy degress of freedom, and absence of mass mixing of new heavy scalars with the SM Higgs doublet. Furthermore, most of the approaches present in the literature are LO SMEFT, i.e. they include SM up to next-to-leading order (NLO) and SMEFT "contact" terms. There are two directions for improving upon this scenario: adding dim > 6 operators without touching the SM loops and inserting dim > 4 operators in SM loops. Ideally one should move along the diagonal direction in this space, doing both.

In Ref. [48] it was re-established that a SMEFT can provide an adequate answer for describing SM deviations beyond LO. The direction chosen in Ref. [48] and also in Refs. [44, 49] is to work with the insertion of dim $= 6$ operators. In this construction, the scale Λ that characterizes BSM physics cannot be too small, because dim $= 8$ operators were neglected. But Λ can also not be too large, because dim $= 4$ higher-order corrections may be more important than the dim $= 6$ interference effects. It is worth noting that these statements do not imply an inconsistency of SMEFT. They only mean that higher-dimensional operators and/or higher-order electroweak (EW) corrections (e.g. Ref. [50]) must also be included if one wants to explore larger ranges of Λ. The general SMEFT decomposition of any amplitude, projected into the $\alpha_s = 0$ plane, reads as follows:

$$\mathscr{A} = \sum_{n=N}^{\infty} g^n\, \mathscr{A}_n^{(4)} + \sum_{n=N_6}^{\infty} \sum_{l=0}^{n} \sum_{k=1}^{\infty} g^n\, g_{4+2k}^{l}\, \mathscr{A}_{nlk}^{(4+2k)}\,, \tag{3}$$

where g is the $SU(2)$ coupling constant, and $g_{4+2k} = 1/(\sqrt{2}\,G_{\rm F}\,\Lambda^2)^k$. For each process, N defines the LO for dim $= 4$: e.g. $N = 1$ for $\mathrm{H} \to \mathrm{VV}$ and other tree-level couplings, but $N = 3$ for $\mathrm{H} \to \gamma\gamma$, a loop-induced process at LO. $N_6 = N$ for tree-level processes and $N_6 = N - 2$ for loop-induced processes.

Generally speaking, there is no factorization of SM higher-order terms. Therefore, reweighing the leading-order SM predictions to account for higher-order QCD and EW corrections, assuming factorization from the EFT effects, is not a procedure that produces accurate results (see Ref. [51]). As far as QCD factorization is concerned let us consider the well-known example

$$\mathrm{g}(p_1) + \mathrm{g}(p_2) \to \mathrm{A}(p_a) + \mathrm{B}(p_b) + \mathrm{X} \quad (p_1 = z x_1 P_1,\ p_2 = x_2 P_2)\,, \tag{4}$$

where $(p_a + p_b)^2 = Q^2$, $\tau s = Q^2$, and $z \to 1$ is the soft limit

$$d\,\sigma\left(\tau\,,Q^2\,,\ldots\right) = \int dx_1\,dx_2\,dz\,f_{\rm g}\left(x_1\,,\mu_{\rm F}\right)\,f_{\rm g}\left(x_2\,,\mu_{\rm F}\right)\,\delta\left(\tau - x_1 x_2 z\right)d\,\hat{\sigma}\left(z\,,\alpha_{\rm s}\,,\frac{Q^2}{\mu_{\rm R}^2}\,,\frac{Q^2}{\mu_{\rm F}^2}\ldots\right)\,, \tag{5}$$

where $d\,\hat{\sigma} = d\,\hat{\sigma}^0\,z\,G$ and

$$G^{\rm NLO}\left(z\,,\alpha_{\rm s}\right)\Big|_{\rm soft} = \delta\left(1-z\right) + \frac{\alpha_{\rm s}}{2\,\pi}\left[d_1\,D_1(z) + \left(c_0 + c_1 + \ldots\right)\,\delta\left(1-z\right)\right],\quad D_n(z) = \left[\frac{\ln^n(1-z)}{1-z}\right]_+\,. \tag{6}$$

Non-universal NLO corrections (that are process-dependent) enter through the coefficient c_1 and $D_n(z)$ (plus sub-leading terms,) imply convolution, and dominate the cross-section in the soft limit. For re-evaluation it is important to have the answer in terms of SM deviations: this allows to "reweight" when new (differential) K-factors become available. New input will touch only the dim $= 4$ components.

The rationale in building a QFT of SM deviations is not so much the numerical impact of higher orders (even if some can be sizable) but in promoting a phenomenological tool (the κ-framework) to the full-fledged status of QFT. Another reason for having a complete formalism is to avoid a situation where experimentalists will have to go back and "remove" a provisional formalism from the analysis. To explain SMEFT in a nutshell, consider a process described by a SM amplitude

$$\mathscr{A}_{\rm SM} = \sum_{i=1}^{n} \mathscr{A}_{\rm SM}^{(i)}\,, \tag{7}$$

where the $\mathscr{A}_{\rm SM}^{(i)}$ are gauge-invariant sub-amplitudes. In general, the same process is given by a contact term or a collection of contact terms of dim $= 6$; for instance, direct coupling of H to $\mathrm{VV}(\mathrm{V} = \gamma, \mathrm{Z}, \mathrm{W})$. In order to construct the theory one has to select a set of higher-dimensional operators and start the complete procedure of renormalization. Of course, different sets of operators can be interchanged as long as they are closed under renormalization. It is evident that renormalization is best performed when using the so-called Warsaw basis, see Ref. [52]. Moving from SM to SMEFT we obtain

$$\mathscr{A}_{\rm SMEFT} = \sum_{i=1}^{n} \kappa_i\,\mathscr{A}_{\rm SM}^{(i)} + g_6\,\kappa_{\rm c} + g_6 \sum_{i=1}^{\rm N} a_i\,\mathscr{A}_{\rm nfc}^{(i)}\,, \tag{8}$$

where $g_6^{-1} = \sqrt{2}\,G_{\rm F}\,\Lambda^2$ and $\kappa_i = 1 + g_6\,\Delta\kappa_i$. The last term in Eq. (8) collects all NLO contributions that do not factorize (nfc) and the a_i are Wilson coefficients. The $\Delta\kappa_i$ are linear combinations of the a_i.

We conclude that Eq. (8) gives the correct generalization of the original κ-framework at the price of introducing additional, non-factorizable, terms in the amplitude. In strict LO SMEFT and in the linear realization, only the $\kappa_{\rm c}$ contact term is included with the following drawback: $\kappa_{\rm c}$ is non-zero but $\Delta\kappa_i = 0$. Therefore, when measuring a deviation from the SM prediction we would find a non-zero value for $\kappa_{\rm c}$. However, at NLO the $\Delta\kappa_i$ are non-zero, leading to a (NLO) degeneracy. The interpretation in terms of $\kappa_{\rm c}^{\rm LO}$ or in terms of $\{\kappa_{\rm c}^{\rm NLO}\,,\Delta\kappa_i^{\rm NLO}\}$ is rather different. Indeed, mapping of experimental constraints to Wilson coefficients at LO, or at NLO, should be corrected for if an inferred constraint on a coefficient is to be used in predicting another process. For the $\mathrm{H} \to \gamma\gamma$ decay process, within LO SMEFT we "measure" just $a_{\rm AA} = {\rm s}_\theta^2\,a_{\phi\rm W} + {\rm c}_\theta^2\,a_{\phi\rm B} + {\rm s}_\theta\,{\rm c}_\theta\,a_{\phi\rm WB}$, while at NLO there are contributions proportional to $a_{\rm tW}, a_{\rm tB}, a_{\rm bW}, a_{\rm bB}, a_{\phi\rm W}, a_{\phi\rm B}, a_{\phi\rm WB}$ (with a mixing among $\{a_{\phi\rm W}, a_{\phi\rm B}, a_{\phi\rm WB}\}$). The differences are illustrated in Fig. 1.

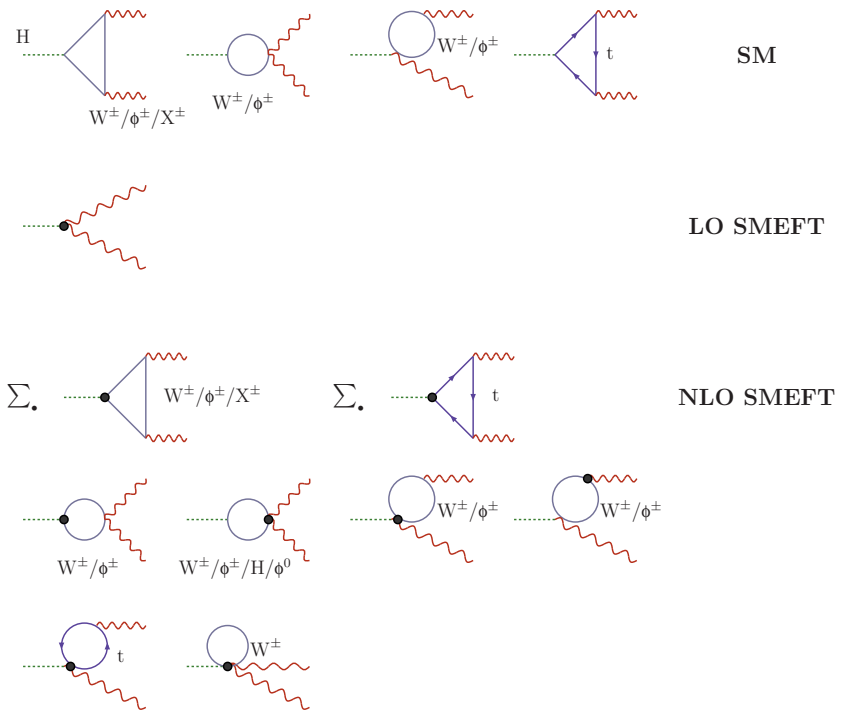

Figure 1: Diagrams contributing to the amplitude for $H \to \gamma\gamma$ in the R_ξ -gauge: SM (first row), LO SMEFT (second row), and NLO SMEFT. Black circles denote the insertion of one dim = 6 operator. $\sum_\bullet$ implies summing over all insertions in the diagram (vertex by vertex). For triangles with internal charge flow $(t, W^\pm, \phi^\pm, X^\pm)$ only the clockwise orientation is shown. Non-equivalent diagrams obtained by the exchange of the two photon lines are not shown. Higgs and photon wave-function factors are not included. The Fadeev-Popov ghost fields are denoted by X.

In NLO SMEFT each κ-parameter has a second index which specifies the corresponding process. One easily discovers that there are correlations among the different κ-parameters and cross-constraints as well: this can be seen by solving the inversion problem ($c_\theta^2 = M_W^2/M_Z^2$):

$$\Delta\kappa_b^{HAZ} - \Delta\kappa_t^{HAZ} - \Delta\kappa_b^{HAA} + \Delta\kappa_t^{HAA} = c_\theta^2\,\Delta\kappa_W^{HAZ} + \left(\frac{3}{2} + 2c_\theta^2\right)\left(\Delta\kappa_t^{HAZ} - \Delta\kappa_t^{HAA}\right) + \left(\frac{1}{2} + 3c_\theta^2\right)\Delta\kappa_W^{HAA} = 0\,. \quad (9)$$

$$a_{t\phi} = \frac{1}{2s_\theta^2}\,a_{\phi D} - 2\,a_{\phi\Box} + \Delta\kappa_t^{HAA}\,, \qquad a_{b\phi} = -\frac{1}{2s_\theta^2}\,a_{\phi D} + 2\,a_{\phi\Box} - \Delta\kappa_b^{HAA}\,,$$
$$a_{\phi\Box} = \frac{1}{4s_\theta^2}\,a_{\phi D} + \frac{1}{2}\,\Delta\kappa_W^{HAA} \qquad 2c_\theta^2\,a_{\phi D} = s_\theta^2\left(\Delta\kappa_b^{HAZ} - \Delta\kappa_b^{HAA}\right). \quad (10)$$

Considering decay processes, $H \to \{F\}$, we define "effective" kappas

$$\mathscr{A}_{SMEFT}^{H\{F\}}(s) = \kappa^{H\{F\}}(s)\,\mathscr{A}_{SM}(s) \qquad \Gamma_H = \sum_{\{F\}} |\,\kappa^{H\{F\}}(M_H^2)\,|^2\,\Gamma_{H\to\{F\}}^{SM}\,. \quad (11)$$

In writing Eq. (11) we have assumed that the Higgs may not decay to new invisible or undetectable particles. Another important point to mention is the dependence of the effective kappas on the scale relevant for the process; that has the consequence that rescaling couplings at the H peak is not the same thing as rescaling them off-peak [34, 53].

Therefore, off-shell measurements are (much) more than consistency checks on $\Gamma_{\rm H}$: observing an excess in the off-shell measurement will be a manifestation of BSM physics, which might or might not need to be in relation with the H width.

It is worth noting that (a priori) discarding subsets of dim $= 6$ operators is not advisable and, as usual, approximations should be the last step in the procedure, after full calculations are performed. The theory of SM deviations, workable to all orders, is still in its infancy but clearly marks the irrelevance of protracted discussion of which SMEFT basis to use; a basis is *by definition* closed under renormalization, and anything that is not a basis, such as many effective Lagrangians, should be viewed with due care. With NLO SMEFT we can study Higgs couplings to very high accuracy and try to understand sources of deviations that may appear in the data from multiple sectors. Potentially, there will be a blurred arrow in the space of Wilson coefficients pointing the way to the UV completion of the SM, and we should simply focus the arrow.

Another important point to mention is the limit where SM deviations are set to zero. In this limit one should recover the most accurate SM predictions. Therefore, it is certainly allowed to decompose an amplitude into form factors multiplying Lorentz structures (as long as the form factors have a known analytical structure) but the limit where "anomalous couplings" vanish (see Eq. 4 of Ref. [54]) should not be interpreted as the LO value of the SM amplitude. This is to say that a measurement that deviated from the LO SM predictions is not necessarily "anomalous". From a constructive point of view, it seems reasonable to require a common language in describing SM deviations, i.e. "HVV anomalous couplings" can be easily incorporated into the more universal notion of SMEFT.

Before the LHC, when there was no measurement of the Higgs boson mass, there were two interesting scenarios in $\mathrm{V_L V_L} \to \mathrm{V_L V_L}$ longitudinal scattering: $M_{\rm W}^2, M_{\rm Z}^2 \ll M_{\rm H}^2 \ll s$ and $M_{\rm W}^2, M_{\rm Z}^2 \ll s \ll M_{\rm H}^2$. With a 125 *GeV* Higgs boson we analyze a new option, $M_{\rm W}^2, M_{\rm Z}^2, M_{\rm H}^2 \ll s$, for which the SM result for the transition matrix is

$$\frac{d}{dt}\,\sigma_{\mathrm{V_L V_L}\to\mathrm{V_L V_L}} = \frac{\left|\mathrm{T}(s,t)\right|^2}{16\,\pi\,s^2}\,, \quad \mathrm{T}^0_{\rm LO} = \frac{1}{16\,\pi\,s}\int_{-s}^{0} dt\,\mathrm{T}_{\rm LO}\,, \quad \mathrm{T}^0_{\rm LO}\left(\mathrm{W_L^+ W_L^-} \to \mathrm{W_L^+ W_L^-}\right) \overset{s\to\infty}{\sim} -\frac{G_{\rm F} M_{\rm H}^2}{4\sqrt{2}\,\pi}\,. \tag{12}$$

As well know, anomalous couplings violate perturbative unitarity. However, one has to be careful in formulating the problem, as the region of interest is $M_{\rm W}^2, M_{\rm Z}^2, M_{\rm H}^2 \ll s \ll \Lambda^2$. In this case, when s approaches Λ^2, the SMEFT must be replaced by its UV completion and it makes no sense to study the limit $s \to \infty$ in SMEFT. However, it is known that heavy degrees of freedom may induce effects of *delayed* unitarity cancellation [55] in the intermediate region. These effects could be detectable in vector boson scattering if there is enough room between $M_{\rm H}$ and the scale of BSM physics. We derive

$$\mathrm{T}^0_{\rm SM+SMEFT} \sim \sum_{n=0}^{2} \mathrm{T}_n\,(G_{\rm F}\,s)^n\,. \tag{13}$$

As expected, the SM part contributes to the constant term, while dim $= 6$ operators have positive powers of s (up to a power of two). The leading behavior is controlled by the $\mathscr{O}_{\phi\rm WB} = \Phi^\dagger\,\tau^a\,\Phi\,\mathrm{F}_a^{\mu\nu}\,\mathrm{F}^0_{\mu\nu}$ operator. Delayed unitarity cancellations might very well be the best window for detecting BSM physics.

2. Not just the LHC Higgs

In our quest for a UV completion of the SM we cannot neglect the sensitivity of electroweak precision data (EWPD). By its general nature, the SMEFT is not confined to describe Higgs couplings and their SM deviations: it can be used to reformulate the constraints from EWPD and to analyze the whole set of processes measurable at LHC and future colliders, such as single and multiple gauge boson production, Drell-Yan physics, associated production of gauge bosons and jets, triple gauge coupling searches, $M_{\rm W}$, asymmetries such as $\mathrm{A_{FB}}$, extraction of $\sin^2\theta_{\rm W}$, etc. Here we present a few examples of EWPD evaluated in NLO SMEFT.

2.1. $\alpha_{\rm QED}$ at the mass of the Z

If we neglect loop-generated (LG) operators [56] in loops, the following result holds for vacuum polarization:

$$\Pi_{\rm AA}^{(\dim=6)}(0) = -8\,({\rm c}_\theta^2/{\rm s}_\theta^2)\,a_{\phi\rm D}\,\Pi_{\rm AA}^{(\dim=4)}(0)\,. \tag{14}$$

One of the key ingredients in computing precision (pseudo-)observables is $\alpha_{\rm QED}$ at the mass of the Z, defined by

$$\alpha(M_{\rm Z}) = \frac{\alpha(0)}{1-\Delta\alpha^{(5)}(M_{\rm Z})-\Delta\alpha_{\rm t}(M_{\rm Z})-\Delta\alpha_{\rm t}^{\alpha\alpha_{\rm s}}(M_{\rm Z})}\,, \tag{15}$$

$$\Delta\alpha^{(5)}(M_{\rm Z}) = \Delta\alpha_{\rm l}(M_{\rm Z})+\Delta\alpha_{\rm had}^{(5)}(M_{\rm Z})\,. \tag{16}$$

The numerical impact of the different corrections is

$$\begin{array}{ll}
\Delta\alpha_{\rm had}^{(5)}(M_{\rm Z}) & = 0.0280398 \\
10^4\times\Delta\alpha_{\rm l}(M_{\rm Z}) & = 0.0314976 \\
10^4\times\Delta\alpha_{\rm t}(M_{\rm Z}) & \approx [-0.62\,,\,-0.55] \\
10^4\times\Delta\alpha_{\rm t}^{\alpha\alpha_{\rm s}}(M_{\rm Z}) & \approx [-0.114\,,\,-0.095]
\end{array}$$

The effect of the SMEFT is equivalent to multiply $\Delta\alpha_{\rm l,t}(M_{\rm Z})$ by $1-\kappa_\alpha$, where

$$\kappa_\alpha = 8\,g_6\,({\rm c}_\theta^2/{\rm s}_\theta^2)\,a_{\phi\rm D} = 0.188\,a_{\phi\rm D} \qquad \text{for } \Lambda = 3\ TeV\,. \tag{17}$$

Therefore, $\mid\kappa_\alpha\,\Delta\alpha_{\rm t}\mid > \Delta\alpha_{\rm l}$ and $\mid\kappa_\alpha\,\Delta\alpha_{\rm t}\mid \approx \mid\Delta\alpha_{\rm t}^{\alpha\alpha_{\rm s}}\mid$.

2.2. The ρ-parameter

Consider the following decomposition of the gauge-boson self-energies (see Ref. [57]):

$$S_{\rm WW} = \frac{g^2}{16\pi^2}\,\Sigma_{\rm WW}\,, \qquad S_{\rm ZZ} = \frac{g^2}{16\pi^2\,{\rm c}_\theta^2}\,\left(\Sigma_{33}-2\,{\rm s}_\theta^2\,\Sigma_{3Q}-{\rm s}_\theta^4\,\Pi_{\rm AA}\,s\right)\,,$$

$$\Sigma_{\rm F} \;=\; \Sigma_{\rm WW}(0)-{\rm Re}\,\Sigma_{33}(M_{\rm Z}^2)+{\rm Re}\,\Sigma_{3Q}(M_{\rm Z}^2) \tag{18}$$

and define $\rho^{-1} = 1+\frac{G_{\rm F}}{2\sqrt{2}\pi^2}\,\Sigma_{\rm F} = 0.99490$; $\Delta\rho$ depends on $a_{\phi\rm D}$, $a_{\phi\Box}$, $a_{\phi\rm f}$, and $a_{\phi\rm f}^{(1,3)}$ (with f = l, u, d,) when considering only potentially-tree-generated (PTG) operators [56]. The leading term, that should not be used for accurate predictions, is

$$\Delta\rho = M_{\rm t}^2\left[\kappa_\rho\,\Delta\rho^{(4)}+g_6\sum_i\Delta\rho_i^{(6)}\,a_i\right], \quad \kappa_\rho = 1+\frac{g_6}{11}\left[\frac{7}{6}\,a_{\phi\rm D}+28\,(a_{\phi\rm q}^{(1)}+a_{\phi\rm q}^{(3)})-20\,a_{\phi\rm t}\right], \tag{19}$$

where $a_i = a_{\phi\rm D}$, $a_{\phi\rm t}$, $a_{\phi\rm q}^{(1,3)}$. The explicit form for $\Delta\rho_i^{(6)}$ will not be reported here.

2.3. The W mass

Working in the α-scheme we can predict $M_{\rm W}$ [57]. The solution is

$$\frac{M_{\rm W}^2}{M_{\rm Z}^2} = \hat{\rm c}_\theta^2+\frac{\alpha}{\pi}\,{\rm Re}\Big\{\left(1-\frac{1}{2}\,g_6\,a_{\phi\rm D}\right)\Delta_{\rm B}^{(4)}(M_{\rm W})+\sum_{\rm gen}\Big[\left(1+4\,g_6\,a_{\phi\rm l}^{(3)}\right)\Delta_{\rm l}^{(4)}(M_{\rm W})$$
$$+\left(1+4\,g_6\,a_{\phi\rm q}^{(3)}\right)\Delta_{\rm q}^{(4)}(M_{\rm W})\Big]+g_6\left[\Delta_{\rm B}^{(6)}(M_{\rm W})+\sum_{\rm gen}\left(\Delta_{\rm l}^{(6)}(M_{\rm W})+\Delta_{\rm q}^{(6)}(M_{\rm W})\right)\right]\Big\}\,, \tag{20}$$

where $\Delta_i^{(4,6)}$ (with $i = \mathrm{l}, \mathrm{q}$, and B) are the dim $= 4, 6$ corrections due to leptons, quarks, and bosons. Furthermore, we have introduced the LO solution (in the α-scheme) for the weak-mixing angle:

$$\hat{s}_\theta^2 = \frac{1}{2}\left[1 - \sqrt{1 - 4\frac{\pi\alpha}{\sqrt{2}G_F M_Z^2}}\right]. \tag{21}$$

The expansion can be improved when working within the SM (dim $= 4$), e.g. by expanding in powers of $\alpha(M_Z)$.

2.4. Dijet data

In the Warsaw basis [52] there are two distinct sets of dim $= 6$ operators: dim $= 6$ four-fermion operators (Tab. 3 in Ref. [52]) and other dim $= 6$ operators (Tab. 2 in Ref. [52]) . The first set is relevant for a) NLO SMEFT predictions involving processes with external fermions (e.g. $\mathrm{H} \to \bar{\mathrm{b}}\mathrm{b}$, $\mathrm{Z} \to \bar{\mathrm{f}}\mathrm{f}$ etc.) and for b) processes dominated by QCD interactions, such as dijet distributions, etc. In the first case, four-fermion operators modify the fermion self-energy, contributing to the fermion mass renormalization and to the fermion wave-function factor, and any $\Phi\bar{\mathrm{f}}\mathrm{f}$ vertex ($\Phi = \mathrm{H}, \mathrm{Z}$ and W). Alternatively, this set of operators is relevant in probing the SM with dijets, e.g. in the study of angular distributions of dijets in the process $\mathrm{pp} \to \mathrm{jj}$, see Ref. [58]. The relevant partonic processes are $\mathrm{uu} \to \mathrm{uu}$, $\mathrm{dd} \to \mathrm{dd}$, and $\mathrm{ud} \to \mathrm{ud}$. It is worth noting that operators such as $(\bar{\mathrm{q}}_\mathrm{L}\,\gamma^\mu\,\mathrm{q}_\mathrm{L})^2$ are PTG and their Wilson coefficients are not necessarily suppressed by the loop factor $1/(16\pi^2)$, which might be important when considering strongly-coupled BSM physics. At the moment, NLO SMEFT predictions for dijet production are not available.

2.5. Flavour physics

Searches for BSM physics in flavour observables have been interpreted in terms of an effective Hamiltonian description using dim $= 6$ operators [59]. The usage of SMEFT, instead of an effective Lagragian or Hamiltonian, would allow to have a consistent and comprehensive way to also incorporate any observed deviations from the SM in flavour physics observables.

2.6. Lepton dipole operators

The Wilson coefficients $a_{\mathrm{lW}}, a_{\mathrm{lB}}$, and $a_{\mathrm{lequ}}^{(3)}$ (see Tab. 2 and Tab. 3 of Ref. [52]) give a LO contribution to remarkably clean windows to BSM physics, namely the $\mu \to \mathrm{e} + \gamma$ decay, the anomalous magnetic moment of the muon, and to the electric dipole moment of the electron (see Ref. [21]) . The NLO calculation is presently not available and will be useful for future precision studies.

2.7. Other examples

Other processes that can be treated within a SMEFT framework are: top pair production [60, 61], neutral triple gauge boson interactions [62], Higgs boson plus jet production [63], and boosted Higgs boson production [64].

2.8. Electroweak precision data

There are several ways to incorporate EWPD. So far, the most common option has been to reduce (a priori) the number of dim $= 6$ operators considered. Open questions regarding this procedure are [65]: should one fit one κ at a time? Should one fit first to the EWPD and then to H observables? A combination of both? The SMEFT is the framework and we are just at the beginning of a new phase that should witness the consolidation of a "common language" between the theory and experimental communities, linking together many different LHC and non-LHC analyses. In any case, it is essential that the derivation of constraints is done in a consistent and basis-independent manner [20]. A recent analysis [66] reaches conclusions that differ from the usual claim, namely that it is not justified to set individual Wilson coefficients to zero in the analysis of LHC data as an attempt to incorporate pre-LHC (EWPD) data.

3. The strategy: from the multi-pole expansion to pseudo-observables to kappas

Concerning the κ-framework, we can say that the κ-parameters are easy to understand in terms of how they change cross sections and partial decay widths. Extending the framework should be seen as expressing the κ parameters in terms of SMEFT coefficients. One question that remains to be answered is the following: could we use and translate part of the LEP language, e.g. that of pseudo-observables (PO), to recast SMEFT parameters into inclusive POs?

What are POs? To be concise we could say that what the experimenters do is to collapse (and/or transform) some "primordial quantities" (such as the number of observed events in some pre-defined set-up) into some "secondary quantities" which we feel that are closer to the theoretical description of the phenomena. How were POs defined at LEP? We will give one example: within the context of the SM, fiducial observables (FO) at LEP are described in terms of some set of amplitudes and cross sections:

$$A_{\rm SM} = A_\gamma + A_{\rm Z} + \text{non-fact.}\,, \quad \sigma\left(\hat{s}\right) = \int dz \mathrm{H}_{\rm in}\left(z,\hat{s}\right) \mathrm{H}_{\rm fin}\left(z,\hat{s}\right) \hat{\sigma}\left(z,\hat{s}\right)\,, \tag{22}$$

where $\mathrm{H}_{\rm in,fin}$ are QED/QCD radiators. Once the amplitude, dressed by the weak loop corrections, is given we use the fact that in the SM there are several effects, such as the imaginary parts or the $\gamma-$Z interference or the pure QED background, that have a negligible influence on the line shape. Therefore, POs are determined by fitting FOs but some ingredients are still taken from the SM, making the model-independent results dependent upon the SM prediction.

In this way, the exact (de-convoluted) cross-section is successively reduced to a Z-resonance. It is a modification of a pure Breit-Wigner resonance because of the s-dependent width:

$$\sigma_{\bar{\rm f}\rm f}\left(s\right) = \sigma_0^{\bar{\rm f}\rm f} \frac{s^2\,\Gamma_{\rm Z}^2}{\left(s-M_{\rm Z}^2\right)^2 + s^2\,\Gamma_{\rm Z}^2/M_{\rm Z}^2} \qquad \sigma_0^{\bar{\rm f}\rm f} = \frac{12\,\pi}{M_{\rm Z}^2}\,\frac{\Gamma_{\rm e}\,\Gamma_{\rm f}}{\Gamma_{\rm Z}^2}\,. \tag{23}$$

The partial widths are computed by including all we know about loop corrections. One needs to specify $M_{\rm Z}$ and the (remaining) relevant SM parameters for the SM-complement. For instance, the explicit formulae for the Z$\bar{\rm f}$f vertex are

$$\rho_{\rm Z}^{\rm f}\,\gamma^\mu\,\left[\left(I_{\rm f}^{(3)} + i\,a_{\rm L}\right)\gamma_+ - 2\,Q_{\rm f}\,\kappa_{\rm Z}^{\rm f}\,\sin^2\theta + i\,a_{\rm Q}\right] = \gamma^\mu\left(\mathscr{G}_{\rm V}^{\rm f} + \mathscr{G}_{\rm A}^{\rm f}\,\gamma^5\right), \tag{24}$$

where $\gamma_+ = 1+\gamma^5$, and $a_{\rm Q,L}$ are the SM imaginary parts. By definition, the total and partial widths of the Z boson include also QED and QCD corrections.

From LEP to LHC, does history repeat itself? Why should it? It should because POs are a platform between realistic observables and theory parameters, allowing experimentalists and theorists to meet half way between; i.e. theorists do not have to run full simulation and reconstruction and experimentalists do not need to fully unfold to model-dependent parameter spaces.

Clearly, the LHC is not LEP and there are many differences. As a consequence, we face new problems, e.g. off-shell LHC physics is not simple and resonant/non-resonant are perfectly tied together, posing severe questions of gauge invariance.

Despite inherent, albeit technical, difficulties the next job for the LHC is the high-precision study of SM-deviations; this will require several steps. For each process, one can write down the SMEFT amplitude, both for resonant and non-resonant parts and compute fiducial observables. Then, express the resonant part as a function of POs without altering the total, something different from the strategy adopted at LEP. The SM non-resonant part also changes and cannot be subtracted. At this point, conventionally-defined POs can be fit to data, and later interpreted in terms of SMEFT Wilson coefficients (or BSM Lagrangian parameters).

In order to define POs at the LHC we need various ingredients [67], e.g. multi-pole expansion (MPE), see Ref. [68], and phase-space factorization. In any process, the residues of the poles (starting from maximal degree) are gauge invariant quantities, see Ref. [69]. The non-resonant part of the amplitude is a gauge-invariant, multivariate, function. That is to say that the residue of the resonant poles can be POs by themselves and expressing them in terms of other

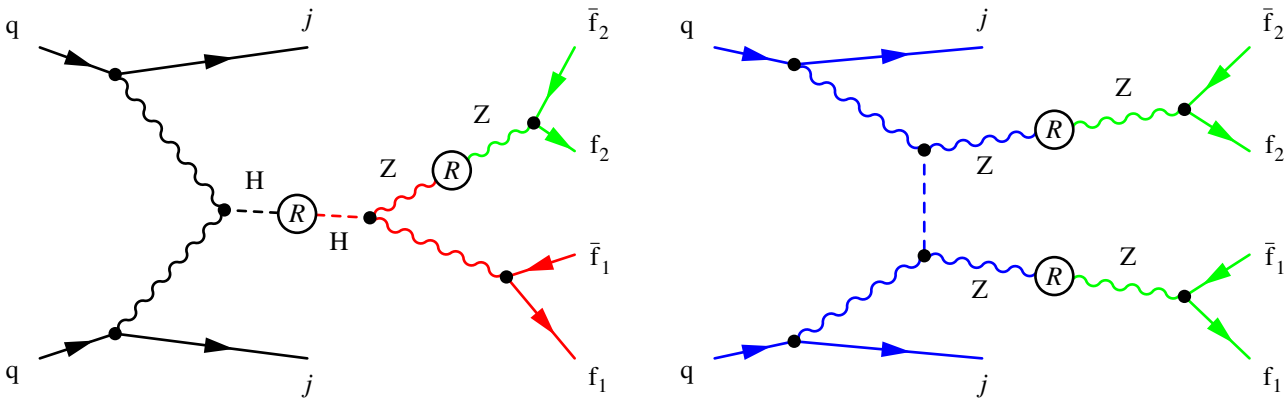

Figure 2: Pseudo-Observables for the triple-resonant (left) and double-resonant (right) parts of the $\mathbf{qq \to \bar{f}_1 f_1 \bar{f}_2 f_2 jj}$ process at NLO. Each color defines one PO and the R blobs correspond to the cuts induced by the (resonant) δ-part of the propagator, shown in Eq. (25).

objects (e.g. SMEFT Wilson coefficients) is an operation the can be postponed to an interpretation step. The end of the chain, when no poles are left, requires an (almost) model-independent SMEFT or model-dependent BSM description; numerically speaking, it depends on the sensitivity of the measurements to the non-resonant part.

The MPE has a dual role: as we mentioned, poles and their residues are intimately related to the gauge-invariant splitting of the amplitude (Nielsen identities); residues of poles (eventually after integration over other variables) can be interpreted as POs, something that requires factorization of the amplitude squared. However, gauge-invariant splitting is not the same as "factorization" of the process into sub-processes; indeed phase-space factorization requires the pole to be inside the physical region

$$
\begin{aligned}
|\Delta|^2 &= \frac{1}{(s-M^2)^2+\Gamma^2 M^2} = \frac{\pi}{M\Gamma}\,\delta\left(s-M^2\right) + \mathrm{PV}\left[\frac{1}{(s-M^2)^2}\right], \\
d\Phi_n\left(P,p_1\ldots p_n\right) &= \frac{1}{2\pi}\,dQ^2\,d\Phi_{n-(j+1)}\left(P,Q,p_{j+1}\ldots p_n\right)\,d\Phi_j\left(Q,p_1\ldots p_j\right).
\end{aligned}
\tag{25}
$$

To "complete" the decay ($d\Phi_j$) we need the δ-function in Eq. (25). We can say that the δ-part of the resonant propagator opens the corresponding line and allows us to define POs. This is not the case for t-channel propagators, which cannot be cut. Consider the process $\mathrm{qq} \to \bar{\mathrm{f}}_1\mathrm{f}_1\bar{\mathrm{f}}_2\mathrm{f}_2 jj$: given the structure of the resonant poles we can define different POs, e.g.

$$
\begin{aligned}
\sigma(\mathrm{qq} \to \bar{\mathrm{f}}_1\mathrm{f}_1\bar{\mathrm{f}}_2\mathrm{f}_2 jj) &\overset{PO}{\longmapsto} \sigma(\mathrm{qq} \to \mathrm{H}jj)\,\mathrm{Br}(\mathrm{H} \to \mathrm{Z}\bar{\mathrm{f}}_1\mathrm{f}_1)\,\mathrm{Br}(\mathrm{Z} \to \bar{\mathrm{f}}_2\mathrm{f}_2), \\
\sigma(\mathrm{qq} \to \bar{\mathrm{f}}_1\mathrm{f}_1\bar{\mathrm{f}}_2\mathrm{f}_2 jj) &\overset{PO}{\longmapsto} \sigma(\mathrm{qq} \to \mathrm{ZZ}jj)\,\mathrm{Br}(\mathrm{Z} \to \bar{\mathrm{f}}_1\mathrm{f}_1)\,\mathrm{Br}(\mathrm{Z} \to \bar{\mathrm{f}}_2\mathrm{f}_2).
\end{aligned}
\tag{26}
$$

These two possibilities are illustrated in Fig. 2. There are fine points to be considered when factorizing a process into "physical" sub-processes. Consider an amplitude that can be factorized as follows:

$$
\mathscr{A} = \mathscr{A}_\mu^{(1)}\,\Delta_{\mu\nu}(p)\,\mathscr{A}_\nu^{(2)}, \tag{27}
$$

where $\Delta_{\mu\nu}$ is the propagator for a spin-1 resonance. We would like to replace

$$
\Delta_{\mu\nu} \to \frac{1}{s-s_\mathrm{c}}\sum_\lambda \varepsilon_\mu(p,\lambda)\,\varepsilon_\nu^*(p,\lambda), \tag{28}
$$

where s_c is the complex pole and ε_μ are the spin-1 polarization vectors. What we obtain is

$$
|\,\mathscr{A}\,|^2 = \frac{1}{|\,s-s_\mathrm{c}\,|^2}\left|\left[\mathscr{A}^{(1)}\cdot\varepsilon\right]\left[\mathscr{A}^{(2)}\cdot\varepsilon^*\right]\right|^2 \tag{29}
$$

However, extracting the δ from the propagator does not necessarily factorize the phase space, i.e. we do not find back the needed form,

$$\sum_{\lambda}\left|\mathscr{A}^{(1)}\cdot\varepsilon(p,\lambda)\right|^2\sum_{\sigma}\left|\mathscr{A}^{(2)}\cdot\varepsilon(p,\sigma)\right|^2. \tag{30}$$

Is there a solution? Yes, if and only if cuts are not introduced. In that case the interference terms between different helicities oscillate over the phase space and drop out, i.e. we achieve factorization, see Refs. [70]. Furthermore, the MPE should be understood as an "asymptotic expansion", see Refs. [71, 72], not as a narrow-width approximation (NWA). The phase space decomposition is obtained by using the two parts in the propagator expansion of Eq. (25): the δ-term is what we need to reconstruct POs, the PV-term gives the remainder, and POs are extracted without making any approximation. It is worth noting that in extracting pseudo-observables, analytic continuation (of on-shell masses into complex poles) is performed only after integrating over residual variables [73].

The MPE returns Green's functions in well-defined kinematic limits, i.e. residues of the poles after extracting the parts which are one-particle-reducible. These residues can then be computed within SMEFT (or any BSM model) and expressed in terms of Wilson coefficients (or BSM Lagrangian parameters).

We can illustrate the MPE-PO connection by using a simple but non-trivial example: the Dalitz decay of the Higgs boson, see Ref. [74]. Consider the process

$$\mathrm{H}(P)\to\bar{\mathrm{f}}(p_1)+\mathrm{f}(p_2)+\gamma(p_3)\,, \tag{31}$$

and introduce invariants $s_{\mathrm{H}}=-P^2$, $s=-(p_1+p_2)^2$, and propagators $\Delta_{\mathrm{A}}(i)=1/s_i$ and $\Delta_{\mathrm{Z}}(i)=1/(s_i-s_{\mathrm{Z}})$. With $s_{\mathrm{H}}=\mu_{\mathrm{H}}^2-i\,\mu_{\mathrm{H}}\,\gamma_{\mathrm{H}}$ we denote the H complex pole, etc.. In the limit of $m_{\mathrm{f}}\to 0$, the total amplitude for process Eq. (31) is given by the sum of three contributions: Z-resonant, A-resonant, and non-resonant,

$$\mathscr{A}\left(\mathrm{H}\to\bar{\mathrm{f}}\mathrm{f}\gamma\right)=\left[\mathrm{A}_{\mathrm{Z}}^{\mu}\left(s_{\mathrm{H}}\,,s\right)\Delta_{\mathrm{Z}}(s)+\mathrm{A}_{\mathrm{A}}^{\mu}\left(s_{\mathrm{H}}\,,s\right)\Delta_{\mathrm{A}}(s)\right]e_{\mu}\left(p_3,l\right)+\mathrm{A}_{\mathrm{NR}}\,, \tag{32}$$

where e_{μ} is the photon polarization vector. The two resonant components are given by

$$\mathrm{A}_{\mathrm{V}}^{\mu}\left(s_{\mathrm{H}}\,,s\right)=\mathscr{T}_{\mathrm{HAV}}\left(s_{\mathrm{H}}\,,s\right)\mathrm{T}_{\nu}^{\mu}\left(q\,,p_3\right)J_{\mathrm{V\,f}}^{\nu}\left(q;p_1,p_2\right)\,, \tag{33}$$

where $J_{\mathrm{V\,f}}^{\mu}$ is the V-fermion (f) current, $\mathrm{T}^{\mu\nu}\left(k_1\,,k_2\right)=k_1\cdot k_2\,\delta^{\mu\nu}-k_1^{\nu}\,k_2^{\mu}$, $q=p_1+p_2$, and $\mathrm{V}=\mathrm{A},\mathrm{Z}$. Having the full amplitude, we start the MPE according to

$$\mathscr{T}_{\mathrm{HAZ}}\left(s_{\mathrm{H}}\,,s\right)=\mathscr{T}_{\mathrm{HAZ}}\left(s_{\mathrm{H}}\,,s_{\mathrm{Z}}\right)+\left(s-s_{\mathrm{Z}}\right)\,\mathscr{T}_{\mathrm{HAZ}}^{(1)}\left(s_{\mathrm{H}}\,,s\right)\qquad\text{etc.}, \tag{34}$$

and obtain the following result:

$$\begin{aligned}\mathscr{A}\left(\mathrm{H}\to\bar{\mathrm{f}}\mathrm{f}\gamma\right)=\mathrm{T}_{\mu\nu}\left(q\,,p_3\right)&\Big[\mathscr{T}_{\mathrm{HAZ}}\left(s_{\mathrm{H}}\,,s_{\mathrm{Z}}\right)\Delta_{\mathrm{Z}}(s)J_{\mathrm{Z\,f}}^{\nu}\left(q;p_1,p_2\right)+\mathscr{T}_{\mathrm{HAA}}\left(s_{\mathrm{H}}\,,0\right)\Delta_{\mathrm{A}}(s)J_{\mathrm{A\,f}}^{\nu}\left(q;p_1,p_2\right)\\&+\mathscr{T}_{\mathrm{HAZ}}^{(1)}\left(s_{\mathrm{H}}\,,s\right)J_{\mathrm{Z\,f}}^{\nu}\left(q;p_1,p_2\right)+\mathscr{T}_{\mathrm{HAA}}^{(1)}\left(s_{\mathrm{H}}\,,s\right)J_{\mathrm{A\,f}}^{\nu}\left(q;p_1,p_2\right)\Big]e^{\mu}\left(p_3,l\right)+\mathrm{A}_{\mathrm{NR}}\end{aligned} \tag{35}$$

It is easy to verify that

$$\sum_{\lambda=0,\pm1}e_{\mu}\left(q\,,\lambda\right)e_{\nu}^{*}\left(q\,,\lambda\right)=\delta_{\mu\nu}+\frac{1}{s}\,q_{\mu}\,q_{\nu}\,,\qquad q\cdot q=-s\,. \tag{36}$$

Consider now the single-resonant, Z, part

$$\mathscr{A}_{\mathrm{SR;AZ}}=\mathscr{T}_{\mathrm{HAZ}}\left(s_{\mathrm{H}}\,,s_{\mathrm{Z}}\right)\mathrm{T}_{\mu\nu}\left(q\,,p_3\right)\Delta_{\mathrm{Z}}(s)J_{\mathrm{Z\,f}}^{\nu}\left(q;p_1,p_2\right)e^{\mu}\left(p_3,l\right) \tag{37}$$

and introduce

$$\mathrm{E}_i(q)=J_{\mathrm{Z\,f}}^{\mu}\left(q;p_1,p_2\right)e_{\mu}^{*}(q,i)\,,\qquad\mathrm{P}_{\mu\,,i}(q)=\mathscr{T}_{\mathrm{HAZ}}\left(s_{\mathrm{H}}\,,s_{\mathrm{Z}}\right)\mathrm{T}_{\mu\nu}\left(q\,,p_3\right)e^{\nu}(q,i)\,. \tag{38}$$

Squaring and summing over spins gives

$$\sum_{\text{spin}} \left|\mathscr{A}_{\text{SR}\,;\,\text{AZ}}\right|^2 = \text{P}_{\mu\,,\,i}(q)\left[\text{P}^{\mu}_{j}(q)\right]^{\dagger}\text{E}_i(q)\,\text{E}^{\dagger}_j(q)\left|\Delta_{\text{Z}}(s)\right|^2, \tag{39}$$

instead of what is expected for a factorized term, namely

$$\sum_{\text{spin}} \left|\mathscr{A}^{\text{fc}}_{\text{SR}\,;\,\text{AZ}}\right|^2 = \frac{1}{3}\sum_{ij}\text{P}_{\mu\,,\,i}(q)\left[\text{P}^{\mu}_{i}(q)\right]^{\dagger}\left|\text{E}_j(q)\right|^2\left|\Delta_{\text{Z}}(s)\right|^2. \tag{40}$$

The result in Eq. (40) is what we need to define the relevant PO, namely $\Gamma\left(\text{H}\to\text{Z}\gamma\right)$. Derivation continues by writing

$$\Gamma_{\text{SR}}\left(\text{H}\to\bar{\text{f}}\text{f}\gamma\right) = \frac{1}{2M_{\text{H}}}\frac{1}{(2\pi)^5}\int d\text{PS}_{1\to3}\sum_{\text{spin}}\left[\left|\mathscr{A}^{\text{fc}}_{\text{SR}\,;\,\text{AZ}}\right|^2 + \Delta\mathscr{A}_{\text{SR}\,;\,\text{AZ}}\right],$$

where $\left|\mathscr{A}_{\text{SR}\,;\,\text{AZ}}\right|^2 = \left|\mathscr{A}^{\text{fc}}_{\text{SR}\,;\,\text{AZ}}\right|^2 + \Delta\mathscr{A}_{\text{SR}\,;\,\text{AZ}}$. Let us now turn to the phase-space integral for the process in Eq. (31). With $P^2 = -M_{\text{H}}^2$, we introduce the Mandelstam variables: $s = -(P-p_3)^2 = -(p_1+p_2)^2$, $t = -(P-p_1)^2 = -(p_2+p_3)^2$, and $u = -(P-p_2)^2 = -(p_1+p_3)^2$, such that $s+t+u = M_{\text{H}}^2$. Then consider the n-dimensional integral

$$\Phi_n(s\,,t) = \int\prod_i d^n p_i\,\theta(p_i^0)\,\delta(p_i^2)\,\delta^n(P-\sum_i p_i)\,\delta((P-p_3)^2+s)\,\delta((p_2+p_3)^2+t)\,, \tag{41}$$

that is related to the phase-space integral by

$$\int d\text{PS}_{1\to3} = \int_{-\infty}^{+\infty} dsdt\Phi_n(s\,,t)\,. \tag{42}$$

Using Eq. (42) as well as $0\le s\le M_{\text{H}}^2$, $0\le t\le M_{\text{H}}^2$, and $0\le s+t\le M_{\text{H}}^2$, we can write

$$\Gamma_{\text{SR}}\left(\text{H}\to\bar{\text{f}}\text{f}\gamma\right) = \frac{1}{2}M_{\text{H}}^5\frac{1}{(2\pi)^5}\int_0^1 dx_s\int_0^{1-x_s} dx_t\,\Phi_4(x_s\,,x_t)\sum_{\text{spin}}\left[\left|\mathscr{A}^{\text{fc}}_{\text{SR}\,;\,\text{AZ}}\right|^2 + \Delta\mathscr{A}_{\text{SR}\,;\,\text{AZ}}\right], \tag{43}$$

where we have introduced scaled variables, $s = x_s M_{\text{H}}^2$, etc.

It is easily seen that $\Delta\mathscr{A}_{\text{SR}\,;\,\text{AZ}}$ vanishes after integration over $0\le x\le 1$, but this is not the case if cuts are introduced. This result is an explicit example of a general proof given in Ref. [70]. We therefore derive the result in the extrapolated scenario. To summarize the steps, we have the following:

The Z single-resonant amplitude. This is given by

$$\begin{aligned}&\mathscr{T}_{\text{HAZ}}\left(s_{\text{H}}\,,s_{\text{Z}}\right)\text{T}_{\mu\nu}\left(q\,,p_3\right)e^{\nu}(p_3\,,l)\,\delta^{\mu}_{\alpha}\,J^{\alpha}_{\text{Zf}}\left(p\,;q,k\right)\Delta_{\text{Z}}(s)\quad\to\\ &\sum_i\mathscr{T}_{\text{HAZ}}\left(s_{\text{H}}\,,s_{\text{Z}}\right)\text{T}_{\mu\nu}\left(q\,,p_3\right)e^{\nu}(p_3\,,l)\,e^{\mu}(q\,,i)\left[e_{\alpha}(q\,,i)\right]^{\dagger}J^{\alpha}_{\text{Zf}}\left(p\,;q,k\right)\Delta_{\text{Z}}(s)\,.\end{aligned} \tag{44}$$

The fully extrapolated scenario. This allows to replace the (squared) S-matrix element with

$$\sum_{ij}\left|\mathscr{T}_{\text{HAZ}}\left(s_{\text{H}}\,,s_{\text{Z}}\right)\text{T}_{\mu\nu}\left(q\,,p_3\right)e^{\nu}(p_3\,,j)\,e^{\mu}(q\,,i)\right|^2\left|\Delta_{\text{Z}}(s)\right|^2\frac{1}{3}\sum_l\left|e_{\alpha}(q\,,l)\,J^{\alpha}_{\text{Zf}}\left(p\,;q,k\right)\right|^2. \tag{45}$$

At this point, if cuts are introduced there is an extra contribution.

The decomposition of the resonant part . We obtain

$$\Gamma_{\rm SR}\left({\rm H}\to\bar{\rm f}{\rm f}\gamma\right)=\frac{1}{16}\,\frac{1}{(2\,\pi)^5}\,\frac{\pi^2}{M_{\rm H}^2}\int_0^1 dx_s\int_0^{1-x_s}dx_t\sum_{ij}\left|{\rm S}_{{\rm H}\to{\rm Z}\gamma}\right|^2\frac{1}{3}\sum_l\left|{\rm S}_{{\rm Z}\to\bar{\rm f}{\rm f}}\right|^2\left|\overline{\Delta}_{\rm Z}(s)\right|^2, \tag{46}$$

where the scaled propagator is $\overline{\Delta}_{\rm Z}(s)=1/(x_s-s_{\rm Z}/M_{\rm H}^2)$. The integrand does not depend on x_t and we can use

$$\int_0^{1-x_s}dx_t=1-x_s\,,\quad \int_0^1 x_s^n\left|\overline{\Delta}_{\rm Z}(s)\right|^2=\frac{\pi}{\mu_{\rm Z}\gamma_{\rm Z}}\,\delta\left(x_s-\frac{\mu_{\rm Z}^2}{M_{\rm H}^2}\right)+\ \text{reg. part} \tag{47}$$

We also introduce

$$\mathrm{F}_{\rm proc}\left(s_{\rm Z}\,,\,s\right)=\sum_{\rm spin}\left|{\rm S}_{\rm proc}\right|^2. \tag{48}$$

The reason for the dependence with s in Eq. (48) is due to kinematical factors. This is to say that the kinematic is real and no approximation is made.

The PO definition. At this point the POs may be defined as

$$\Gamma_{\rm PO}\left({\rm H}\to{\rm Z}\gamma\right)=\frac{1}{16\,\pi}\,\frac{1}{M_{\rm H}}\left(1-\frac{\mu_{\rm Z}^2}{M_{\rm H}^2}\right){\rm F}_{{\rm H}\to{\rm Z}\gamma}\left(s_{\rm Z}\,,\,\mu_{\rm Z}^2\right),\quad \Gamma_{\rm PO}\left({\rm Z}\to\bar{\rm f}{\rm f}\right)=\frac{1}{48\,\pi}\,\frac{1}{\mu_{\rm Z}}\,{\rm F}_{{\rm Z}\to\bar{\rm f}{\rm f}}\left(s_{\rm Z}\,,\,\mu_{\rm Z}^2\right). \tag{49}$$

The final result. The final result can be expressed as

$$\Gamma_{\rm SR}\left({\rm H}\to\bar{\rm f}{\rm f}\gamma\right)=\frac{1}{2}\,\Gamma_{\rm PO}\left({\rm H}\to{\rm Z}\gamma\right)\,\frac{1}{\gamma_{\rm Z}}\,\Gamma_{\rm PO}\left({\rm Z}\to\bar{\rm f}{\rm f}\right)+\ \text{remainder}. \tag{50}$$

In the narrow-width approximation the remainder is neglected; we keep it in our formulation where the goal is to define POs without making approximations. Figure 3 illustrates the MPE of the $\mathrm{H}\to\gamma\bar{\mathrm{f}}\mathrm{f}$ process as described above.

We can repeat the question: what are POs? The conclusion is that residues of resonant poles, κ-parameters, and Wilson coefficients are different layers of POs. The layer closest to theory refers to Wilson coefficients or non-SM parameters in BSM models, such as $\alpha,\beta,M_{\rm sb}$, etc. in two-Higgs-doublet models. There is then a layer using kappas, an intermediate layer defined by residues that is similar to $g^{\rm e}_{\rm V\,A}$ at LEP, and a layer closer to experiment that is similar to $\Gamma(\mathrm{Z}\to\bar{\mathrm{f}}\mathrm{f})$ at LEP. However, while $\Gamma(\mathrm{Z}\to\bar{\mathrm{f}}\mathrm{f})$ can be defined, $\Gamma(\mathrm{H}\to\mathrm{ZZ})$ cannot because not all three particles can be on-shell simultaneously; in other words, POs are defined by convention, but they cannot violate kinematics.

One has to be careful to not confuse the residue of the two Z poles in the $\mathrm{H}\to4\mathrm{f}$ amplitude (needed for a question of gauge invariance) with a partial decay width. The crucial step is shown in Eq. (25): let $s_{\rm V}$ be the complex pole for a particle V, parametrized as $s_{\rm V}=\mu_{\rm V}^2-i\,\mu_{\rm V}\gamma_{\rm V}$. Consider the following integral

$$\mathrm{I}\left(a\,,\,b\,,\,s_{\rm V}\right)=\int_a^b ds\frac{1}{\mid s-s_{\rm V}\mid^2}=\int_a^b ds\frac{1}{(s-\mu_{\rm V}^2)^2+\mu_{\rm V}^2\,\gamma_{\rm V}^2}\,. \tag{51}$$

From Ref. [75] we obtain

$$\begin{aligned}
&\mathrm{I}\left(a\,,\,b\,,\,s_{\rm V}\right)=\frac{\pi}{(\mathrm{A}\,\lambda)^{1/2}}\,\theta(X)\,\theta(1-X)+\mathrm{I}^{\rm reg}\left(a\,,\,b\,,\,s_{\rm V}\right),\\
&\mathrm{I}^{\rm reg}\left(a\,,\,b\,,\,s_{\rm V}\right)=-\frac{1}{2}\sum_{l=1,2}\left|X_l\right|^2\int_0^1 dx\,x^{-1/2}\left(\mathrm{A}X_l^2+\lambda\,x\right)^{-1}=-\frac{1}{\mathrm{A}}\sum_{l=1,2}{}_2\mathrm{F}_1\left(1\,,\,\frac{1}{2}\,;\,\frac{3}{2}\,;\,-\frac{\lambda}{\mathrm{A}X_l^2}\right),\\
&{}_2\mathrm{F}_1\left(1\,,\,\frac{1}{2}\,;\,\frac{3}{2}\,;\,-z^2\right)=\frac{1}{z}\,\arctan z,
\end{aligned} \tag{52}$$

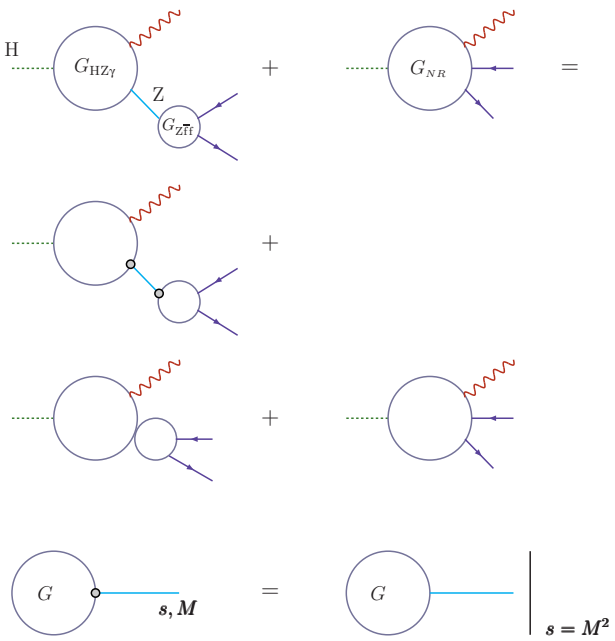

Figure 3: Multi-pole-expansion for $\mathrm{H} \to \gamma\bar{\mathrm{f}}\mathrm{f}$. G stands for Green's function and G_{NR} denotes the non-resonant part of the amplitude. The sum of amplitudes in the second (third) row is gauge-parameter independent. In the last row, an amplitude with an external line of virtuality s and mass M is put on-shell.

with parameters

$$\mathrm{A} = b - a\,, \quad X = \frac{\mu_{\mathrm{V}}^2 - a}{b - a}\,, \quad \lambda = \frac{\mu_{\mathrm{V}}^2 \gamma_{\mathrm{V}}^2}{b - a}\,, \quad X_1 = X_2 = -X\,. \tag{53}$$

In general we have

$$\mathrm{I}_n(a\,, b\,, s_{\mathrm{V}}) = \int_a^b ds \frac{s^n}{|\, s - s_{\mathrm{V}}\,|^2} = \frac{\pi}{(\mathrm{A}\,\lambda)^{1/2}} (\mu_{\mathrm{V}}^2)^n\, \theta(X)\, \theta(1 - X) + \mathrm{I}_n^{\mathrm{reg}}(a\,, b\,, s_{\mathrm{V}})\,. \tag{54}$$

The integral that is needed for the four-body decay,

$$I = \int_0^{M_{\mathrm{H}}^2} ds_1 \int_0^{(M_{\mathrm{H}} - \sqrt{s_1})^2} ds_2\, \lambda^{1/2}\left(M_{\mathrm{H}}^2\,, s_1\,, s_2\right) \left|\Delta_{\mathrm{Z}}(s_1)\, \Delta_{\mathrm{Z}}(s_2)\right|^2\,, \tag{55}$$

can be worked out along the same lines. Factorization of phase-space (the "opening" of a line) requires the identification of "virtuality with mass" (s with μ_{V}^2), which requires $0 \le X \le 1$, i.e. $a \le \mu_{\mathrm{V}}^2 \le b$. Therefore, the natural PO is $\Gamma(\mathrm{H} \to \mathrm{Z} + \bar{\mathrm{f}}\mathrm{f})$.

Some of the POs that were used at LEP have now been calculated at the two-loop level, see Refs. [76–80], and two-loop renormalization of the full SM has been completed in Refs. [75, 81, 82]. While the corresponding theoretical uncertainties are adequate to compare with current precision measurements, significant improvements will be necessary to make full use of the precision foreseen at future facilities, e.g. the FCC-ee. Preliminary studies [83] seem to indicate the need for full two-loop exponentiation for QED ISR, relevant for the measurement itself, and full three-loop EW radiative corrections, relevant for the interpretation. New POs will appear, e.g. σ_{ZH} relying on accurate threshold cross section measurements sensitive to loop corrections.

4. What to fit

In the linear realization of SMEFT, a subset of dim = 6 operators involves a Higgs doublet Φ that contains both the physical Higgs field H and its vacuum expectation value, v. When using a dim = 6 operator there is a term coming from the replacement of Φ with v (not with H) and one gets a shift in dim = 4 operators, i.e. kinetic and mass terms. Normalization of these terms must always be the canonical one, i.e. the one appearing in the SM Lagrangian. This means that one has to redefine all fields and parameters, including the ghost sector, even before starting the actual calculation of observables. Furthermore, this set of redefinitions affects the sources and this must be taken into account when building S-matrix elements out of Green's functions, for details see Ref. [48]. These extra terms are essential in defining the SMEFT content of all POs.

A question that is often raised concerns the "optimal" parametrization of the dim = 6 basis. Clearly all bases are equivalent and there is no obstacle in "extracting" (Wilson) coefficients as defined in a particular basis. However, certain linear combinations of Wilson coefficients in one basis become a single Wilson coefficient in another basis and a mapping of this type, that puts coefficients and (pseudo-)observables in a one-to-one correspondence, may seem appropriate when considering LO constraints from EWPD [37]. But even at this level one should be careful, since Wilson coefficients mix under renormalization. Furthermore, it could be sensible to start with fits at the level of POs (or kappas), as usually done in flavor physics, instead of directly on Wilson coefficients.

Based on these considerations there are suggestions on separating weakly- and strongly-constrained combinations of Wilson coefficients, possibly disregarding the latter. However, this is currently done in the lowest order implementation of the experimental constraints and there is already strong evidence that NLO SMEFT provides non-negligible corrections, which are relevant for per-mille/few percent constraints. For a given observable $\mathscr{O}$ one can compute the deviation $\mathscr{O}_{\rm SMEFT}/\mathscr{O}_{\rm SM} - 1$ and the corresponding probability distribution function (pdf) with the result that the LO pdf differs from the NLO pdf at the level required by the projected accuracy.

The problem is as follows: we have a basis where constraints are not in one-to-one correspondence with the POs, which is ideal for the NLO extension (renormalization, mixing of Wilson coefficients, etc.) and a basis where LO implementation of constraints is automatic but (much) less suitable for NLO extension. The obvious solution is to perform the full NLO SMEFT analysis in the basis which is well-suited (therefore reducing the SMEFT theoretical uncertainty), not only for Higgs physics but also for EWPD, and only after that one identifies weakly- and strongly-constrained combinations.

To summarize, mapping of experimental constraints to Wilson coefficients at LO, and at NLO, should be corrected for if an inferred coefficient is to be used in another process. Any LO analysis will miss contributions from the running of Wilson coefficients (renormalization group) and from finite "non-factorizable" terms that are not negligible: as stressed in Sect. 2.4 of Ref. [66] this source of uncertainty (pertinent to the LO studies) is not currently included in the fits. The best way to improve the uncertainty due to missing higher orders is to move a step forward in the perturbative expansion (both in g and in Λ).

5. Effective solutions

There is now convincing evidence from the LHC results that one should use more theoretical tools; not only consider more theories (i.e. specific models) but also make use of EFTs. Both have specific functions, and both are required [84].

The Euler-Heisenberg theory of photon-photon scattering and the Fermi theory of weak interactions are prototypical examples of EFTs. In both theories only are "relevant" fields are considered and other fields are hidden. Both theories are valid only up to a scale Λ, e.g. $E_\gamma \ll m_{\rm e}$, and unitarity is violated at large scales by Fermi theory. Both theories are non-renormalizable, are based on certain symmetries, and provided stepping stones for scientific advancement.

Also gravity is amenable to an EFT description, see Refs. [85, 86]. This allows to predict the effect of quantum physics on the gravitational interaction of two heavy masses. However, such an EFT would only be valid for "ordinary"

distances (where the curvature is small) and far away from singularities. Of course, for such a description to be relevant would require there being no new physics all the way up to the Planck scale.

If one has a full BSM model it is not necessary to use SMEFT to describe BSM (Higgs) physics because one can always compute anything from the full BSM model. However, is very convenient to use SMEFT because that forces us to concentrate on universal aspects of SM deviations. Of course, the SMEFT modifies the high energy behavior of any UV completion and the effective theory is only a valid description of the physics at energies below the scale of new physics. Essentially, interpreting data via Wilson coefficients may allow to discern the UV completion of the SM. In general, we should strive to devise a more fundamental description, since the idea of an ultimate theory has a powerful aesthetic pull. However, do we have that theory? We have models, mostly "ad hoc" models that cannot be the "fundamental theory" and that are sometimes introduced to "cure" a specific set of experimental results. Without entering a detailed discussion, even if we assume that a "fundamental" theory does exist, e.g. superstrings, presently we cannot test its "resolved" regime, i.e. phenomena at a very large scale.

This is why SMEFT in the bottom-up approach is so useful: we do not know what the tower of UV completions is (or if it exists at all) but we can formulate the SMEFT and perform calculations with it without needing to know what happens at arbitrarily high scales.

6. Through the precision straits

Directly marrying Wilson coefficients and precision data to quantify deviations from the SM is one option, but not necessarily the most convenient. Another way of taking the next step is based on POs, any quantity that is connected to "data" by a set of well-defined assumptions. Properly defining POs requires care, since we cannot randomly isolate portions of an observable and break gauge invariance. That is why the MPE is useful, as it provides (gauge parameter independent) subsets of the amplitudes. Of course, residues of resonant poles can be computed within the SMEFT in terms of Wilson coefficients, but that is more akin to an interpretation step. Instead, we argue for defining POs closer to the experimental observations by defining POs similar to those at LEP; the LEP POs have stood the test of time and still today accurately encode the information contents of the data they were derived from. Introducing such POs and splitting one observable into products of POs related to "sub-processes" requires factorization of the full phase-space and to either make no cuts (implying an extrapolation) or compute (and include) correction terms.

To conclude, the journey to the next standard model may require crossing narrow straits of precision physics. If that is what nature has in store for us, we must equip ourselves with both a range of concrete BSM models as well as a general SMEFT. Both will be indispensable tools in navigating an ocean of future experimental results. The LEP experience has proven that those results can stand the test of time when expressed in terms of POs. And as long as POs are well defined and calculations are performed in a general and coherent way, nature can be systematically probed and our knowledge of it improved.

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