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Effective Lagrangians from functional matching

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We briefly review a variant of functional matching to derive an Effective Field Theory (EFT) for heavy particles at the one-loop level in the top-down approach. The method integrates out heavy fields that correspond to mass eigenstates, i.e. after removing mixing effects by diagonalizing mass matrices. Tree- and loop-level effects are separated by employing the background-field method, hard and soft modes are separated with the use of the expansion by regions. The method is exemplified for the Higgs Singlet Extension of the Standard Model where the mass $M_{\rm H}$ of the additional Higgs boson is considered large, and the Higgs mixing angle α is assumed to scale like $1/M_{\rm H}$, in order to guarantee decoupling in the large-$M_{\rm H}$ limit. Our calculation is agnostic w.r.t. the type (SMEFT vs. HEFT) of the emerging EFT. Eventually the emerging EFT Lagrangian can be transformed into SMEFT form, but only at the cost of introducing fermionic EFT operators, although no such operators are directly generated upon solving the functional integral over the heavy Higgs field.

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1. Introduction

Effective Field Theories (EFTs), such as Standard Model (SM) Effective Field Theory (SMEFT) or Higgs Effective Field Theory (HEFT), provide an excellent framework to address physics beyond the Standard Model (BSM) induced by new heavy particles. In the top-down approach, EFT Lagrangians of underlying UV-complete models are derived by integrating out the heavy fields, and the EFT couplings (Wilson coefficients) can be confronted with constraints from experimental data.

In the two foundational papers [1, 2] we have proposed a variant of *functional matching* where the heavy fields are integrated out directly in the path integral. The method, which employs the background-field method and the expansion by regions, is to some extent a further development of the method proposed in Ref. [3, 4], but also shares features with similar proposals in the literature [5–7]. Specifically, we have integrated out the heavy Higgs field H (with mass $M_{\rm H}$) in a Higgs Singlet Extension of the SM (SESM) under the assumption that the Higgs mixing angle α is suppressed like $1/M_{\rm H}$ in the limit $M_{\rm H} \to \infty$, in order to guarantee the decoupling of H in this limit. Fermions are treated in the massless approximation. Since the SM-like Higgs doublet is represented non-linearly, our calculation is agnostic w.r.t. the important question whether the resulting EFT is of SMEFT or HEFT type. In the the considered large-$M_{\rm H}$ limit, we find (and prove) that the SMEFT framework with only bosonic operators is not sufficient, although there is no coupling to massless fermions in the BSM sector of the SESM. The EFT Lagrangian can, however, be brought into SMEFT form upon introducing fermionic SMEFT operators. EFTs for singlet-Higgs extensions with heavy Higgs bosons have been discussed in the literature quite extensively in different variants of the model and the inspected large-mass limits, with Refs. [8–10] close to our setup and Refs. [7, 11, 12] at least related; more details on the relation to our work are given in Ref. [2].

In contrast to most discussions of EFT matching procedures, we are very explicit with respect to the procedure of renormalization, which is essential when confronting EFT predictions with results of the full underlying model. Our calculation is validated by comparing our EFT predictions with full SESM higher-order predictions for a representative set of electroweak precision observables.

2. The Higgs singlet extension of the SM with a heavy Higgs boson

Before we go over the individual steps of the functional matching method, we briefly introduce the salient features of the SESM, entirely following the conventions and notation of Refs. [1, 2], which are in turn adopted from the formulations of the SM and SESM given in Refs. [13] and [14], respectively. The SESM extends the SM by a real singlet scalar field $\sigma = v_1 + h_1$ with vacuum expectation value (vev) v_1 and Higgs excitation h_1, which interacts with SM particles via the portal provided by the Higgs potential (featuring a $\mathbb{Z}_2$ symmetry under $\sigma \to -\sigma$)

$$V^{\rm SESM} = -\frac{1}{2}\mu_2^2 \,\mathrm{tr}\left[\Phi^\dagger\Phi\right] + \frac{1}{16}\lambda_2\Big(\,\mathrm{tr}\left[\Phi^\dagger\Phi\right]\Big)^2 - \mu_1^2\sigma^2 + \lambda_1\sigma^4 + \frac{1}{2}\lambda_{12}\sigma^2\,\mathrm{tr}\left[\Phi^\dagger\Phi\right], \qquad (1)$$

where Φ represents the SM-like SU(2) doublet in matrix notation. We consistently parametrize Φ non-linearly,

$$\Phi = \frac{1}{\sqrt{2}}\,(v_2 + h_2)\,U\,, \qquad U = \exp\!\left(2\mathrm{i}\frac{\varphi}{v_2}\right), \qquad \varphi = \frac{1}{2}\varphi_a\tau_a\,, \qquad \tau_a = \text{Pauli matrices}\,, \qquad (2)$$

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where v_2 quantifies its vev, h_2 is the physical Higgs component, which is a gauge singlet, and φ_a are Goldstone-boson fields. The two singlet Higgs fields h_1, h_2 mix to fields h and H corresponding to mass eigenstates (with masses M_h and M_H) via a rotation with angle α ($s_\alpha \equiv \sin\alpha$, $c_\alpha \equiv \cos\alpha$),

$$\begin{pmatrix} H \\ h \end{pmatrix} = \begin{pmatrix} c_\alpha & s_\alpha \\ -s_\alpha & c_\alpha \end{pmatrix} \begin{pmatrix} h_1 \\ h_2 \end{pmatrix}. \tag{3}$$

Instead of using the five original parameters μ_1^2, μ_2^2, λ_1, λ_2, and λ_{12} of the Higgs potential, we reparametrize the Higgs sector in terms of parameters that are closer to phenomenology: the two Higgs masses M_h, M_H, the mixing angle α, the vev v_2, which is closely related to the W-boson mass M_W, and one of the quartic couplings, for which we take λ_{12}. To control the large-mass limit, we introduce the parameter $\zeta \equiv M_\mathrm{H}/M_\mathrm{h}$ and assume the scaling $M_\mathrm{H} \sim \zeta$, $\alpha \sim s_\alpha \sim \zeta^{-1}$, while $\lambda_{12} \sim 1$ and all SM-like parameters are kept fixed for $\zeta \to \infty$.

3. Recap of the functional matching method

a) Background-field method (BFM) and non-linear Higgs realization

Generically, the BFM splits all fields $\phi \to \tilde{\phi} = \hat{\phi} + \phi$ into background parts $\hat{\phi}$, which are classical (in the sense that they do not appear in loops), and quantum parts ϕ, which correspond to quantum fluctuations and are the integration variables of the path integral. Diagrammatically, the fields $\hat{\phi}$ deliver the tree-like lines and and the fields ϕ the lines within loops in Feynman graphs. Since all vertices that are part of a one-loop subdiagram involve exactly two quantum fields, at the one-loop level the relevant part of the Lagrangian is quadratic in quantum fields.

In the non-linear realization of the Higgs doublet Φ, it is convenient to work with the multiplicative split $U \to \tilde{U} = \hat{U}U$ for the matrix U of the Goldstone-boson fields [3, 4]. Owing to the gauge invariance of the background-field effective action, it is possible to absorb all background Goldstone-boson fields into the background gauge fields by a straightforward Stückelberg transformation, i.e. we effectively adopt the unitary background gauge by setting $\hat{U} \to \mathbb{1}$, which reduces the amount of algebraic work in the subsequent steps considerably.

b) Separation of hard and soft field modes

Our implementation of the method of regions [15] for large-mass expansions of Feynman diagrams splits both quantum and background fields into soft (s) and hard (h) field modes. In the SESM with a heavy H boson, the original field H is split into $\hat{H}_s + \hat{H}_h + H_s + H_h$. This separation disentangles EFT contributions that correspond to higher-order effects from hard interactions ending up in Wilson coefficients of effective operators on one side and effective operators entering loops built from soft fields on the other.

The splitting of one-loop diagrams into two integration domains of small and large momenta can be realized as a splitting of the path integral into two functional integrals extending over soft and hard field modes, $\int \mathcal{D}H \to \int \mathcal{D}H_s \int \mathcal{D}H_h$. While the use of the method of regions was not yet part of of the procedure of Ref. [3, 4], where only hard field modes appeared in the calculation of the effective Lagrangian, it is an integral part of most of the functional methods for integrating out heavy fields proposed in the literature [1, 2, 5–7] (for more references, see e.g. Ref. [2]).

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c) Equations of motion (EOMs) for the soft modes of the heavy fields and renormalization

By design, the EFT is valid for energies where for any considered process the particle momenta p_i on tree-level lines of Feynman graphs are much smaller than the heavy masses, i.e. $|p_i^\mu| \ll M_\mathrm{H}$ in our case. In other words, background fields only possess soft modes, and the heavy fields carrying these soft momenta do not represent dynamical degrees of freedom in the EFT. The EOMs for the soft modes of the heavy field can be solved up to the relevant order in the large-mass expansion, and this solution can be used to eliminate them from the Lagrangian.

This procedure produces the EFT Lagrangian for the SESM in the large-M_H limit valid at tree level. We obtain

$$\mathcal{L}_\mathrm{eff}^\mathrm{tree} = \mathcal{L}_\mathrm{SM}^\mathrm{tree} - \frac{g_2^2 s_\alpha^2}{8v_2}\tilde{h}(v_2+\tilde{h})^2\big(\tilde{C}_\mu^a\big)^2 - \frac{s_\alpha^2}{8v_2^2}\tilde{h}^2\Box\tilde{h}^2 + \frac{M_\mathrm{h}^2 s_\alpha^2}{8v_2^3}\tilde{h}^3(6v_2^2+7\tilde{h}v_2+2\tilde{h}^2) + O\big(\zeta^{-4}\big)\,. \quad (4)$$

Here, g_2 denotes the SU(2) gauge coupling, and the vector fields $\tilde{C}_\mu^a$ are the combination of SU(2) gauge fields W_μ^a and the U(1) gauge field B_μ corresponding to the massive $\mathrm{W}^\pm$ and Z bosons,

$$C_\mu = \frac{\tau_a}{2}C_\mu^a = W_\mu + \frac{s_\mathrm{w}}{c_\mathrm{w}}B_\mu\frac{\tau_3}{2} = \frac{1}{2}\left(W_\mu^1\tau_1 + W_\mu^2\tau_2 + \frac{1}{c_\mathrm{w}}Z_\mu\tau_3\right). \quad (5)$$

By virtue of the EOM of the light Higgs field, the tree-level effective Lagrangian $\mathcal{L}_\mathrm{eff}^\mathrm{tree}$ of Eq. (4) is equivalent to $\mathcal{L}_\mathrm{eff}^\mathrm{tree} = \mathcal{L}_\mathrm{SM}^\mathrm{tree} - \frac{s_\alpha^2}{2v_2}O_{\Phi\Box}^\mathrm{SMEFT}$, where $O_{\Phi\Box}^\mathrm{SMEFT} = \frac{1}{2}\mathrm{tr}\{(\Phi^\dagger\Phi)\Box(\Phi^\dagger\Phi)\}$ is one of the usual SMEFT operators of the Warsaw basis [16].

The same procedure can be applied to the counterterm Lagrangian of the full SESM. Since renormalization constants are already one-loop effects, only background fields are relevant in this context. This step requires a proper power-counting of all parameters and fields in the heavy-mass limit. As emphasized in Ref. [1], in order to obtain a consistent effective Lagrangian the large-mass expansion must be carefully performed taking into account that the full-theory renormalization constants may have a different scaling behaviour than the corresponding renormalized quantities. "Good schemes" should respect the decoupling behaviour of the theory at leading mass order also at higher orders in the EFT expansion. In this context, the choice of a proper renormalization of the Higgs vevs, i.e. a proper tadpole scheme, is crucial in $\overline{\mathrm{MS}}$ renormalization of full-theory parameters [1].

d) Integrating out the hard modes of the heavy quantum fields in the path integral

Next, we deal with the hard modes of the heavy quantum fields. Since the part of the Lagrangian that is relevant at one-loop order is quadratic in the quantum fields, the path integral over the hard field modes of the heavy quantum fields is of Gaussian type and can be done analytically. The major complication in this step is the fact that there are also terms that are linear in the heavy quantum fields. These terms can be removed by a field redefinition of the hard quantum field modes in a fully algorithmic manner. This means that the resulting part of the Lagrangian quadratic in the hard modes of the heavy quantum fields can be directly identified based again on a simple power-counting argument and the large-mass expansion can be directly carried out via a Neumann series [1, 2, 5–7], which establishes another technical improvement over the procedure described in Ref. [3, 4].

Schematically, integrating out the hard quantum field H_h and performing the large-mass expansion, turns the Lagrangian $\mathcal{L}^{\text{1-loop}}(\hat{\phi},\phi,H_h)$ that is bilinear in quantum fields into the effective

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Lagrangian $\mathcal{L}_{\text{eff}}^{\text{1-loop}}(\hat{\phi},\phi)$

$$\int \mathcal{D}\phi \int \mathcal{D}H_h \, \exp\left\{\mathrm{i}\int \mathrm{d}^D x \, \mathcal{L}^{\text{1-loop}}(\hat{\phi},\phi,H_h)\right\} \to \int \mathcal{D}\phi \, \exp\left\{\mathrm{i}\int \mathrm{d}^D x \, \mathcal{L}_{\text{eff}}^{\text{1-loop}}(\hat{\phi},\phi)\right\}. \quad (6)$$

Note that (after the field redefinition) the functional integration over the hard modes of the light quantum fields (ϕ_h) does not contribute to $\mathcal{L}_{\text{eff}}^{\text{1-loop}}$ in dimensional regularization, because those corrections are related to scaleless one-loop diagrams.

e) Final form of the effective Lagrangian

The effective Lagrangian resulting from the previous steps only involves soft background and quantum fields, but none of the modes of the heavy fields. The effects of the latter are absorbed into the Wilson coefficients of the effective operators.

Finally, redefinitions of the light quantum fields and IBP in the effective Lagrangian may be used to eliminate redundant operators that only influence off-shell Green functions, but no physical scattering amplitudes. This step is, in particular, necessary to bring the effective Lagrangian into canonical form, i.e. to SMEFT or (more generally) HEFT form.

4. Structure and validation of the resulting EFT Lagrangian for the SESM

In Ref. [2] we have given two forms for the resulting EFT Lagrangian:[1] the first form contains only bosonic operators apart from the SM parts, but necessarily includes EFT operators of non-SMEFT form; the second form contains only SMEFT operators, but involves also fermionic operators that are not directly generated by integrating out the heavy H field. The fermionic SMEFT operators in the latter case result from the EOMs of the gauge-boson fields, which mix bosonic and fermionic operators. From this formal perspective the occurrence of fermionic SMEFT operators is not surprising. Nevertheless, we believe it is worth emphasizing in the context of data analyses within SMEFT that BSM effects may show up in fermionic operators even if the BSM sector does not directly couple to fermions, as it is the case in the SESM with massless fermions.

Form 1 of the one-loop EFT Lagrangian is given in the form

$$\delta\mathcal{L}_{\text{eff}}^{\text{1-loop,BSM}} = \sum_{i\in\text{bos.SMEFT}} C_i^{\text{SMEFT}} O_i^{\text{SMEFT}} + \sum_{n=1}^{8} C_n^{\cancel{\text{SMEFT}}} O_n^{\cancel{\text{SMEFT}}}, \quad (7)$$

where the sum in the first SMEFT-like part runs over the bosonic SMEFT operators O_Φ^{SMEFT}, $O_{\Phi\Box}^{\text{SMEFT}}$, $O_{\Phi D}^{\text{SMEFT}}$, $O_{\Phi W}^{\text{SMEFT}}$, $O_{\Phi B}^{\text{SMEFT}}$, $O_{\Phi WB}^{\text{SMEFT}}$ as defined in Sec. 2.4 of Ref. [13]. The corresponding Wilson coefficients look like

$$C_{\Phi WB}^{\text{SMEFT}} = -\frac{e^2 s_\alpha^2 I_{20}}{16\pi^2 D(D-2) s_{\text{w}} c_{\text{w}} v_2^2}, \qquad I_{20} = \frac{1}{\epsilon} - \gamma_{\text{E}} + \ln(4\pi) + \ln\left(\frac{\mu^2}{M_{\text{H}}^2}\right) + O(\epsilon), \quad (8)$$

with the two-point vacuum integral $I_{20} \equiv B_0(0, M_{\text{H}}^2, M_{\text{H}}^2)$ in $D = 4 - 2\epsilon$ space-time dimensions, the electromagnetic charge e, s_{w} and c_{w} denoting sine and cosine of the weak mixing angle,

[1] In the first preprint version of Ref. [2] only the first form had been given. The possibility to transform it to SMEFT form was pointed out to us by Gerhard Buchalla, who is gratefully acknowledged for this.

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respectively, and the unphysical (matching) scale μ in dimensional regularization. Only the Wilson coefficient $C_{\Phi\Box}^{\text{SMEFT}}$ involves a contribution from BSM renormalization constants (specifically from δs_α), because only the EFT operator $O_{\Phi\Box}^{\text{SMEFT}}$ appears at tree level. As renormalization scheme for s_α, we admit any scheme in which the renormalization constant δs_α shows the same scaling in ζ as s_α at lowest order — a requirement that is, e.g., met by the on-shell scheme suggested in Ref. [17].[2] In unitary gauge, the non-SMEFT operators exemplarily read

$$\begin{aligned} O_1^{\cancel{\text{SMEFT}}} &= (v_2+h)^2 (D_B^\mu C_\mu)^a (D_B^\nu C_\nu)^a\,, \quad (D_B^\mu C_\nu)^a = \partial^\mu C_\nu^a + g_1 \epsilon^{3ab} B^\mu C_\nu^b\,, \quad \dots \\ O_8^{\cancel{\text{SMEFT}}} &= (v_2+h)^2 \left[(C_\mu^a)^2 (C_\nu^b)^2 - C_\mu^a C^{\mu,b} C_\nu^a C^{\nu,b} \right]. \end{aligned} \tag{9}$$

The corresponding Wilson coefficients look very much like $C_{\Phi WB}^{\text{SMEFT}}$ given above. While the SMEFT Wilson coefficients C_i^{SMEFT} do not only contain the suppression factor s_α^2, but also factors like $\lambda_{12}^2 v_2^2/M_{\text{H}}^2$, the non-SMEFT coefficients $C_n^{\cancel{\text{SMEFT}}}$ are all proportional to s_α^2, showing that this part is entirely due to mixing effects.

Form 2 of the one-loop EFT Lagrangian reads

$$\delta\mathcal{L}_{\text{eff}}^{\text{1-loop,BSM}} = \sum_{i\in\text{SMEFT}} \hat{C}_i\, O_i^{\text{SMEFT}}\,, \tag{10}$$

where the sum now also includes the fermionic SMEFT operators $O_{\Phi F}^{(1),\text{SMEFT}}$, $O_{\Phi F}^{(3),\text{SMEFT}}$, and $O_{\Phi f}^{\text{SMEFT}}$ with left-handed fermion doublets $F = L, Q$ and right-handed fermion singlets $f = l, u, d$ [13]. Four-fermion operators do not occur. The Wilson coefficients in Eq. (10) read

$$\begin{aligned} \hat{C}_\Phi &= \frac{(D^2+6D+20) M_{\text{h}}^4 s_\alpha^2 I_{20}}{2D(D^2-4)\pi^2 v_2^6} - \frac{(D^2+4) e^2 M_{\text{h}}^2 s_\alpha^2 I_{20}}{4D(D^2-4)\pi^2 s_{\text{w}}^2 v_2^4} + \frac{(D-4)(M_{\text{H}}^2 s_\alpha^2 - 2\lambda_{12} v_2^2)^3 I_{20}}{48 M_{\text{H}}^2 \pi^2 v_2^6} \\ &\quad + \frac{(D-4)(D+2) M_{\text{h}}^2 s_\alpha^2 (M_{\text{H}}^2 s_\alpha^2 - 2\lambda_{12} v_2^2) I_{20}}{4D(D-2)\pi^2 v_2^6}\,, \\ \hat{C}_{\Phi\Box} &= -\frac{s_\alpha \delta s_\alpha}{v_2^2} - \frac{(D^2+4) e^2 s_\alpha^2 (3-2s_{\text{w}}^2) I_{20}}{16D(D^2-4)\pi^2 s_{\text{w}}^2 c_{\text{w}}^2 v_2^2} + \frac{(M_{\text{H}}^2 s_\alpha^2 - 2\lambda_{12} v_2^2) I_{20}}{192 D(D-2) M_{\text{H}}^2 \pi^2 v_2^4} \\ &\quad \times \left[(D^3 - 18D^2 + 128D - 96) M_{\text{H}}^2 s_\alpha^2 - 2D(D^2 - 42D + 44)\lambda_{12} v_2^2 \right], \\ \hat{C}_{\Phi D} &= -\frac{(D^2+4) e^2 s_\alpha^2 I_{20}}{4D(D^2-4)\pi^2 c_{\text{w}}^2 v_2^2}\,, \qquad \hat{C}_{\Phi W} = -\frac{(D-4) e^2 s_\alpha^2 I_{20}}{8D(D^2-4)\pi^2 s_{\text{w}}^2 v_2^2}\,, \\ \hat{C}_{\Phi B} &= \frac{s_{\text{w}}^2}{c_{\text{w}}^2} \hat{C}_{\Phi W}\,, \qquad \hat{C}_{\Phi WB} = \frac{2 s_{\text{w}}}{c_{\text{w}}} \hat{C}_{\Phi W}\,, \qquad \hat{C}_W = 0\,, \\ \hat{C}_{\Phi F}^{(3)} &= \frac{(D-6) e^2 s_\alpha^2 I_{20}}{16D(D^2-4)\pi^2 s_{\text{w}}^2 v_2^2}\,, \qquad \hat{C}_{\Phi F}^{(1)} = \frac{s_{\text{w}}^2}{c_{\text{w}}^2} \hat{C}_{\Phi F}^{(3)} Y_{\text{w},F}\,, \qquad \hat{C}_{\Phi f} = \frac{s_{\text{w}}^2}{c_{\text{w}}^2} \hat{C}_{\Phi F}^{(3)} Y_{\text{w},f}\,. \end{aligned} \tag{11}$$

Note that the Wilson coefficients $\hat{C}_i$ of the bosonic operators do not exactly coincide with the C_i^{SMEFT} of the first form of the EFT Lagrangian, but their structure is very similar. All non-zero fermionic

[2] In the on-shell scheme used in Ref. [1] for s_α the renormalization constant δs_α contains hard and soft contributions and depends on the tadpole scheme. This entails differences whether the EFT calculation of observables is carried out in the linear or non-linear Higgs realization. Since we have employed the *Parameter Renormalized Tadpole Scheme (PRTS)* (see, e.g., Ref. [13]) in the EFT construction in the non-linear realization, any direct application of our EFT using the linear Higgs realization implies the *Gauge-Invariant Vacuum expection value Scheme (GIVS)* [18**?**] for the tadpole treatment if δs_α is kept unchanged (see Ref. [2] for further details).

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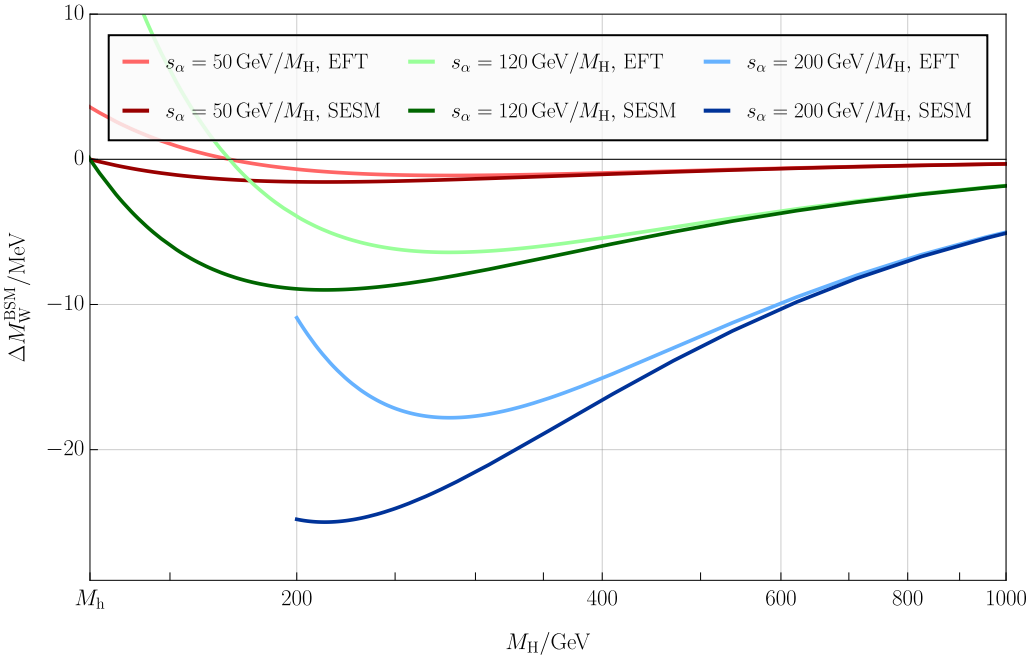

Figure 1: BSM correction $\Delta M_\mathrm{W}^\mathrm{BSM}$ to the W-boson mass, shown for the full SESM and the EFT approximation. The product $s_\alpha M_\mathrm{H}$ is kept fixed in each curve. (Taken from Ref. [2].)

Wilson coefficients are proportional to s_α^2, and the ones of $O_{\Phi F}^{(1),\mathrm{SMEFT}}$ and $O_{\Phi f}^{\mathrm{SMEFT}}$ receive factors of the weak hypercharges $Y_{F/f}$ of the respective fermions.

We have validated both forms of the effective Lagrangian at next-to-leading order in the coupling expansion by verifying that the difference between EFT and full-theory predictions for several observables vanishes faster than $1/M_\mathrm{H}^2$ in the large-M_H limit. Specifically, we examined the W-boson mass M_W derived from muon decay, some W/Z-boson decay widths, the effective weak mixing angle, and the leptonic four-body Higgs decay $\mathrm{h} \to \mathrm{WW} \to \nu_\mathrm{e}\mathrm{e}^+\mu^-\bar{\nu}_\mu$ in Ref. [2]. In Fig. 1 we exemplarily show the BSM contribution to M_W to illustrate this fact.

5. Conclusions

We briefly summarize the salient features and particular strengths of our method [1, 2] to derive the effective Lagrangian at the one-loop level, some of them common to the approaches of Refs. [5–7]:

(i) a clear separation of tree-level and loop effects of the heavy field modes by employing the BFM;

(ii) the possibility to fix the (background) gauge in intermediate steps of the calculation and to restore gauge invariance of the effective Lagrangian at the end;

(iii) transparency in the sense that at each stage of the calculation it is possible to identify the origin of all contributions to the effective Lagrangian in terms of (classes of) Feynman diagrams;

(iv) flexibility due to the fact that no ansatz is made for the effective Lagrangian. Whether the emerging EFT is of SMEFT or HEFT type is part of the result.

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(v) An automation of the method is possible, since it is fully algorithmic. In principle, given a BSM Lagrangian, a proper definition of the large-mass limit with a corresponding power-counting scheme, and some details on the renormalization of the BSM sector, the actual determination of the effective Lagrangian at the one-loop level can be carried out by computer algebra.

As an illustrative example we have derived the EFT Lagrangian to order $O(1/M_{\rm H}^2)$ for the SESM with massless fermions and a heavy Higgs boson of mass $M_{\rm H}$ in the "weak-coupling limit" where the Higgs mixing angle α is suppressed with $1/M_{\rm H}$. We have written the emerging EFT Lagrangian in two alternative forms: one which is of SMEFT form, but involves both bosonic and fermionic SMEFT operators, and another one that involves only bosonic EFT operators, albeit not all of SMEFT type.

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