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CHALLENGES OF EFT AT THE LHC*

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I will discuss the difficulties of putting constraints on the Effective Field Theories (EFT) at the Large Hadron Collider (LHC). In particular, I will analyse the generic properties of the $2 \to 2$ scattering in the presence of the higher dimensional operators indicating that some of the beyond Standard Model effects are parametrically smaller than naive expectations.

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1. Introduction

One of the major goals of the future LHC program is to perform a careful analysis of the properties of the Standard Model (SM) particles. The interactions between SM fields are generically modified in the presence of the new Beyond Standard Model (BSM) particles. If there is a mass gap between the SM and the new physics (NP) states, all the effects of the heavy fields can be parametrized in terms of the higher dimensional operators and the Lagrangian at the electroweak scale will be given by[1]

$$\mathcal{L} = \mathcal{L}^{\mathrm{SM}} + \mathcal{L}^6 + \mathcal{L}^8 + \cdots\,, \qquad \mathcal{L}^D = \textstyle\sum_i c_i^{(D)} \mathcal{O}_i^{(D)}\,, \\ c_i^{(D)} \sim \frac{1}{\Lambda^{D-4}}\,. \tag{1}$$

Usually, we expect the dimension-six terms to capture the leading effects of the new physics corrections to the SM interaction what motivates the truncation of the series in Eq. (1) keeping only the first two terms. The leading BSM effect in $1/\Lambda^2$ expansion will come from the interference between the SM and the BSM diagrams and will be proportional to $\propto 1/\Lambda^2$. The rest of the proceeding will be dedicated to the studies and constraints on this interference term.

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[1] I assume the baryon and lepton number conservation.

2. SM amplitudes

At the LHC, the energy scale of the collision energy is not fixed and one can exploit the relative energy growth of the BSM non-resonant contribution in order to put the stronger bound on the Wilson coefficients. In the limit of $E \gg m_{\mathrm{W}}$, we can treat the SM particles as massless fields. In this limit, it is well-known that the SM amplitudes follow the helicity selection rules, which forces the total helicity of any amplitude in the $2 \to 2$ processes to satisfy[2]

$$h\left(A_4^{\mathrm{SM}}\right) \equiv \sum_i h_i = 0\,, \tag{2}$$

with the exceptions of the four-fermion amplitudes generated by the scalar exchange where the total helicity can have the values $H = \pm 2$ (see Table I for the values of the possible helicities).

TABLE I

Helicities of the four-point amplitudes in the SM and BSM.

A_4	$\lvert h\left(A_4^{\mathrm{SM}}\right)\rvert$	$\lvert h\left(A_4^{\mathrm{BSM}}\right)\rvert$
$VVVV$	0	4,2
$VV\phi\phi$	0	2
$VV\psi\psi$	0	2
$V\psi\psi\phi$	0	2
$\psi\psi\psi\psi$	2,0	2,0
$\psi\psi\phi\phi$	0	0
$\phi\phi\phi\phi$	0	0

3. BSM amplitudes

Let us look now at the amplitudes generated by the dimension-six operators. For simplicity, we will start with the massless case corresponding to the high-energy limit. Since we are interested in the analysis of the $2 \to 2$ scattering, the discussion becomes clearer in the basis where there are no bivalent operators and the number of trivalent operators is reduced to minimum, which happens to be the case in the so-called Warsaw basis [4]. Another simplification which we can make is to use a supersymmetric notation for the field strength tensor

$$F_{\mu\nu}\sigma^{\mu}_{\alpha\dot{\alpha}}\sigma^{\nu}_{\beta\dot{\beta}} \equiv F_{\alpha\beta}\bar{\epsilon}_{\dot{\alpha}\dot{\beta}} + \bar{F}_{\dot{\alpha}\dot{\beta}}\epsilon_{\alpha\beta}\,, \tag{3}$$

[2] We assume all the particles to be incoming.

where $F, (\bar{F})$ transform as (1,0)((0,1)) representations of the Lorentz group, projecting the helicity $+(-)$ states of the vector bosons. Generically, we can characterize [3] any amplitude in terms of the holomorphic and anti-holomorphic weights

$$w(A) = n(A) - h(A)\,, \qquad \bar{w}(A) = n(A) + h(A)\,, \tag{4}$$

where $n(A)$ is the number of legs and $h(A)$ is the total helicity. Then, we can relate the operator with the helicity of the amplitude it generates [3] by generalizing Eq. (4)

$$w(\mathcal{O}) = \min_A \{w(A)\}\,, \qquad \bar{w}(\mathcal{O}) = \min_A \{\bar{w}(A)\}\,, \tag{5}$$

where the minimization is taken over all the amplitudes with only one operator insertion. In practice, one can show the weight is fixed by the amplitude with the smallest number of the legs [3]. Then, for the dimension-six operators, the weights are given in Table II, and we can trivially see that the helicity of the amplitude generated by the operator $\mathcal{O}$ will be constrained

$$\bar{w}(\mathcal{O}) - n \leq h\left(A_n^{\mathcal{O}}\right) \leq n - w(\mathcal{O})\,. \tag{6}$$

Then by inspecting the values of the possible helicities of the BSM and SM amplitudes, we can conclude that for the $2 \to 2$ processes, the interference term will always vanish if there is at least one transverse vector boson.

TABLE II

Holomorphic and anti-holmorphic weights for the dimension-six operators. The minimal number of legs and the corresponding helicity is also indicated. The last column shows an estimate for the Wilson coefficient based on the "one scale one coupling" power counting.

$\mathcal{O}_i$	$n_{\min}$	$h_{\min}$	$(w, \bar{w})$	c_i
F^3	3	3	(0,6)	g_*/Λ^2
$F^2\phi^2$, $F\psi^2\phi$, ψ^4	4	2	(2,6)	g_*^2/Λ^2
$\psi^2\bar{\psi}^2$, $\psi\bar{\psi}\phi^2 D$, $\phi^4 D^2$	4	0	(4,4)	g_*^2/Λ^2
$\psi^2\phi^3$	5	1	(4,6)	g_*^3/Λ^2
ϕ^6	6	0	(6,6)	g_*^4/Λ^2

4. Phenomenological implications

So far, we have been considering the $2 \to 2$ scattering in the very high energy ($E \gg m_{\rm w}$) limit. However, very similar methods can be used to classify non-zero mass effects. In the high-energy limit, the corrections to the results in Table I should be controlled by the small parameters

$$\varepsilon_V \equiv \frac{m_V}{E}\,, \qquad \varepsilon_\psi \equiv \frac{m_\psi}{E}\,. \tag{7}$$

However, every extra power of ϵ_i corresponds to the Higgs vev insertion, see the diagrams in Fig. 1. But the Higgs vev flips the total fermion and

Fig. 1. Helicity flip due to the Higgs vev.

the vector boson helicity by factor of ± 1. Thus, we can predict the leading energy behaviour for all the helicity amplitudes for the SM and BSM, see Table III.

TABLE III

Scaling with energy for various SM and BSM amplitudes.

Channel	SM	BSM_6/E^2
$+ + + +$	ε_V^4	ε_V^0
$+ + + -$	ε_V^2	ε_V^0
$+ + - -$	ε_V^0	ε_V^2
$+\frac{1}{2}\ -\frac{1}{2}\ + +$	ε_V^2	ε_V^0
$+\frac{1}{2}\ -\frac{1}{2}\ + -$	ε_V^0	ε_V^2
$+\frac{1}{2}\ -\frac{1}{2}\ 0\ +$	ε_V^1	ε_V^1
$+\frac{1}{2}\ -\frac{1}{2}\ 0\ 0$	ε_V^0	ε_V^0
$0 + + +$	ε_V^3	ε_V^1
$0 + + -$	ε_V^1	ε_V^1
$0\ 0 + +$	ε_V^2	ε_V^0
$0\ 0 + -$	ε_V^0	ε_V^2
$0\ 0\ 0 +$	ε_V^1	ε_V^1
$0\ 0\ 0\ 0$	ε_V^0	ε_V^0

Now, after we know the scaling of all the amplitudes with energy, we can proceed to the discussion of the phenomenological implications. One of the most important questions of any EFT analysis is the region of the validity, *i.e.* where the effects of the dimension-eight operators can be safely ignored.

To make the discussion more precise, we will assume the "one scale one coupling" power counting for all the Wilson coefficients of the dimension-six operators (see Table II). Let us consider the scattering of the longitudinal vector bosons. Then, the differential cross section can be parametrized as:

$$\sigma_{\mathrm{L}} \sim \frac{g_{\mathrm{SM}}^4}{E^2}\left[1+\underbrace{\frac{g_*^2}{g_{\mathrm{SM}}^2}\frac{E^2}{\Lambda^2}}_{\mathrm{BSM}_6\times\mathrm{SM}}+\underbrace{\frac{g_*^4}{g_{\mathrm{SM}}^4}\frac{E^4}{\Lambda^4}}_{\mathrm{BSM}_6^2}+\underbrace{\frac{g_*^2}{g_{\mathrm{SM}}^2}\frac{E^4}{\Lambda^4}}_{\mathrm{BSM}_8\times\mathrm{SM}}+\dots\right], \tag{8}$$

where we have explicitly indicated all the various contributions to the cross section: interference with dimension-six operators, dimension-six operators squared, interference with dimension-eight, *etc.* The results are presented in Fig. 2, where we have explicitly indicated the origins of the leading and subleading contributions. We can clearly see that the EFT analysis with the dimension-six operators *only* is always consistent and the contribution of the dimension-eight operators is subleading till the scale of the cut-off, where any EFT description becomes invalid anyway. Now, let us consider the scattering of the two transverse vector bosons:

$$\sigma_{\mathrm{T}} \sim \frac{g_{\mathrm{SM}}^4}{E^2}\left[1+\overbrace{\frac{g_*}{g_{\mathrm{SM}}}\frac{m_{\mathrm{W}}^2}{\Lambda^2}}^{\mathrm{BSM}_6\times\mathrm{SM}}+\overbrace{\frac{g_*^2}{g_{\mathrm{SM}}^2}\frac{E^4}{\Lambda^4}}^{\mathrm{BSM}_6^2}+\overbrace{\frac{g_*^2}{g_{\mathrm{SM}}^2}\frac{E^4}{\Lambda^4}}^{\mathrm{BSM}_8\times\mathrm{SM}}+\overbrace{\frac{g_*^4}{g_{\mathrm{SM}}^4}\frac{E^8}{\Lambda^8}}^{\mathrm{BSM}_8^2}+\dots\right]. \tag{9}$$

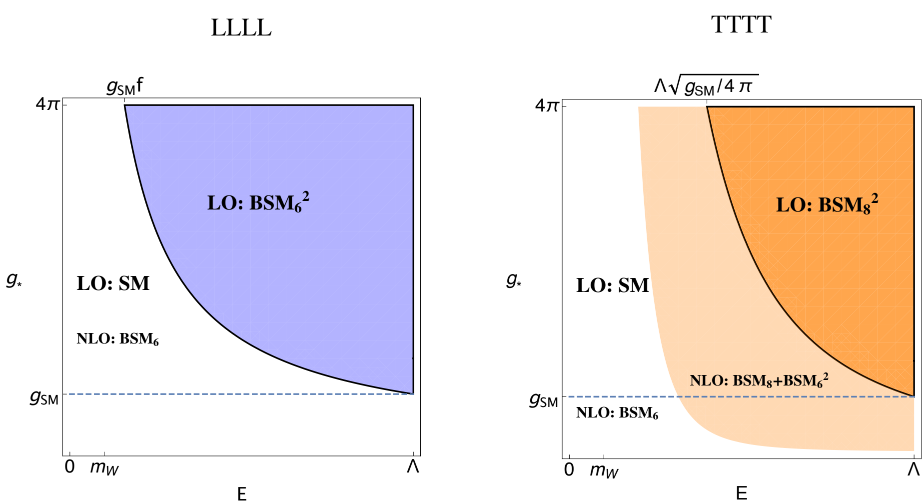

Fig. 2. EFT validity range: origin of the leading and subleading contributions to the cross sections are indicated.

In this case, the interference term has an additional suppression $\sim m_W^2/E^2$ compared to the previous example. This makes the EFT truncation at the level of the dimension-six operators valid only in the low energy and small coupling limit, once we are off this region, the contributions from the dimension-eight operators become essential.

5. Outlook

We have analysed the generic properties of the amplitudes in SM and its extensions by the dimension-six operators. We have shown that the interference between the SM and the BSM amplitudes, in the presence of the transverse vector bosons, is always suppressed compared to the naive expectations. This leads to the additional difficulties in measuring the Wilson coefficients of the corresponding operators at the LHC. At the same time, our results indicate that the $2 \to 3$ processes do not suffer from the above-mentioned suppression and can be used [2, 5] as an additional measurement which can become competitive and complementary to the $2 \to 2$ processes [5].

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