Hidden Symmetries and HEFT

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based on arXiv:1610.00150

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Outlook

Probing the Higgs potential
The model
Functional Symmetries and Renormalization
Expanding around the p.c. renormalizable theory:
the privileged operator $\partial_\mu(\Phi^\dagger\Phi)\partial^\mu(\Phi^\dagger\Phi)$
Applications: anomalous trilinear Higgs coupling
Conclusions

Probing the Higgs potential

The discovery at LHC of a particle compatible
with the Standard Model Higgs boson in 2012
has been an unprecedented breakthrough in the study of
the electroweak symmetry breaking mechanism.

Parameterization of Beyond-the-Standard-Model (BSM) physics
at the LHC is usually carried out
within the Effective Field Theory approach.

Recent summary:

Handbook of LHC Higgs cross sections:
4. Deciphering the nature of the Higgs sector
arXiv:1610.07922

Higgs Effective Field Theories

Add operators of higher dimension to the SM Lagrangian compatible with the symmetries of the theory

$$\mathcal{L}_{\text{eff}} = \mathcal{L}_{\text{SM}} + \sum_i \frac{c_i^{(5)}}{\Lambda} \mathcal{O}_i^{(5)} + \sum_i \frac{c_i^{(6)}}{\Lambda^2} \mathcal{O}_i^{(6)} + \sum_i \frac{c_i^{(7)}}{\Lambda^3} \mathcal{O}_i^{(7)} + \sum_i \frac{c_i^{(8)}}{\Lambda^4} \mathcal{O}_i^{(8)} + \cdots$$

c are the Wilson coefficients, Λ is some large energy scale

Dimension-six bosonic operators

Table 97: Bosonic D=6 operators in the SILH basis.

Bosonic CP-even		Bosonic CP-odd	
O_H	$\frac{1}{2v^2}\left[\partial_\mu(H^\dagger H)\right]^2$		
O_T	$\frac{1}{2v^2}\left(H^\dagger \overleftrightarrow{D}_\mu H\right)^2$		
O_6	$-\frac{\lambda}{v^2}(H^\dagger H)^3$		
O_g	$\frac{g_s^2}{m_W^2} H^\dagger H\, G^a_{\mu\nu} G^a_{\mu\nu}$	$\widetilde{O}_g$	$\frac{g_s^2}{m_W^2} H^\dagger H\, \widetilde{G}^a_{\mu\nu} G^a_{\mu\nu}$
O_γ	$\frac{g'^2}{m_W^2} H^\dagger H\, B_{\mu\nu} B_{\mu\nu}$	$\widetilde{O}_\gamma$	$\frac{g'^2}{m_W^2} H^\dagger H\, \widetilde{B}_{\mu\nu} B_{\mu\nu}$
O_W	$\frac{ig}{2m_W^2}\left(H^\dagger \sigma^i \overleftrightarrow{D}_\mu H\right) D_\nu W^i_{\mu\nu}$		
O_B	$\frac{ig'}{2m_W^2}\left(H^\dagger \overleftrightarrow{D}_\mu H\right) \partial_\nu B_{\mu\nu}$		
O_{HW}	$\frac{ig}{m_W^2}\left(D_\mu H^\dagger \sigma^i D_\nu H\right) W^i_{\mu\nu}$	$\widetilde{O}_{HW}$	$\frac{ig}{m_W^2}\left(D_\mu H^\dagger \sigma^i D_\nu H\right) \widetilde{W}^i_{\mu\nu}$
O_{HB}	$\frac{ig'}{m_W^2}\left(D_\mu H^\dagger D_\nu H\right) B_{\mu\nu}$	$\widetilde{O}_{HB}$	$\frac{ig}{m_W^2}\left(D_\mu H^\dagger D_\nu H\right) \widetilde{B}_{\mu\nu}$
O_{2W}	$\frac{1}{m_W^2} D_\mu W^i_{\mu\nu} D_\rho W^i_{\rho\nu}$		
O_{2B}	$\frac{1}{m_W^2} \partial_\mu B_{\mu\nu} \partial_\rho B_{\rho\nu}$		
O_{2G}	$\frac{1}{m_W^2} D_\mu G^a_{\mu\nu} D_\rho G^a_{\rho\nu}$		
O_{3W}	$\frac{g^3}{m_W^2} \epsilon^{ijk} W^i_{\mu\nu} W^j_{\nu\rho} W^k_{\rho\mu}$	$\widetilde{O}_{3W}$	$\frac{g^3}{m_W^2} \epsilon^{ijk} \widetilde{W}^i_{\mu\nu} W^j_{\nu\rho} W^k_{\rho\mu}$
O_{3G}	$\frac{g_s^3}{m_W^2} f^{abc} G^a_{\mu\nu} G^b_{\nu\rho} G^c_{\rho\mu}$	$\widetilde{O}_{3G}$	$\frac{g_s^3}{m_W^2} f^{abc} \widetilde{G}^a_{\mu\nu} G^b_{\nu\rho} G^c_{\rho\mu}$

From arXiv:1610.07922

UV properties of the HEFT

Power-counting renormalizability is lost
Physical Unitarity (cancellation of ghost states)
guaranteed by BRST symmetry & Slavnov-Taylor identities
Froissart bound usually not respected

In general all possible terms allowed by symmetry
must be included in an EFT approach

J.Gomis, S.Weinberg, Are nonrenormalizable gauge theories renormalizable?
Nucl.Phys. B469 (1996) 473-487

One-loop anomalous dimensions in the HEFT

However a *tour de force* computation
of one-loop anomalous dimensions in the HEFT
has revealed surprising cancellations.

R.Alonso, E.Jenkins, A.Manohar, M.Trott
arXiv:1308.2627 , arXiv:1310.4838 , arXiv:1312.2014 , arXiv:1409.0868

Not all mixings in principle allowed by the symmetries
do indeed arise at one loop level

Holomorphy

Holomorphic part of the Lagrangian: built out of

$$X^+, R, \bar{L}$$

Anti-holomorphic part of the Lagrangian: built out of

$$X^-, \bar{R}, L$$

The remnant is non-holomorphic

Basic idea: holomorphic operators
do not mix with anti-holomorphic and non-holomorphic operators

True at the one-loop level (up to some breaking proportional to Yukawa couplings)
on the S-matrix elements (not off-shell)

One-loop Anomalous Dimensions for a gauge theory

	$(w,\bar{w})$	F^3 $(0,6)$	$F^2\phi^2$ $(2,6)$	$F\psi^2\phi$ $(2,6)$	ψ^4 $(2,6)$	$\psi^2\phi^3$ $(4,6)$	$\bar{F}^3$ $(6,0)$	$\bar{F}^2\phi^2$ $(6,2)$	$\bar{F}\bar{\psi}^2\phi$ $(6,2)$	$\bar{\psi}^4$ $(6,2)$	$\bar{\psi}^2\phi^3$ $(6,4)$	$\bar{\psi}^2\psi^2$ $(4,4)$	$\bar{\psi}\psi\phi^2 D$ $(4,4)$	$\phi^4 D^2$ $(4,4)$	ϕ^6 $(6,6)$
F^3	$(0,6)$			$\times$	$\times$	$\times$			$\times$	$\times$	$\times$	$\times$	$\times$	$\times$	$\times$
$F^2\phi^2$	$(2,6)$				$\times$	$\times$				$\times$	$\times$	$\times$			$\times$
$F\psi^2\phi$	$(2,6)$									$\times$				$\times$	$\times$
ψ^4	$(2,6)$	$\times$	$\times$			$\times$	$\times$	$\times$	$\times$	$\times$	$\times$	y^2		$\times$	$\times$
$\psi^2\phi^3$	$(4,6)$	$\times^*$									y^2				$\times$
$\bar{F}^3$	$(6,0)$			$\times$	$\times$	$\times$			$\times$	$\times$	$\times$	$\times$	$\times$	$\times$	$\times$
$\bar{F}^2\phi^2$	$(6,2)$				$\times$	$\times$				$\times$	$\times$	$\times$			$\times$
$\bar{F}\bar{\psi}^2\phi$	$(6,2)$				$\times$									$\times$	$\times$
$\bar{\psi}^4$	$(6,2)$	$\times$	$\times$	$\times$	$\times$	$\times$	$\times$	$\times$			$\times$	$\bar{y}^2$		$\times$	$\times$
$\bar{\psi}^2\phi^3$	$(6,4)$					$\bar{y}^2$	$\times^*$								$\times$
$\bar{\psi}^2\psi^2$	$(4,4)$		$\times$		$\bar{y}^2$	$\times$		$\times$		y^2	$\times$			$\times$	$\times$
$\bar{\psi}\psi\phi^2 D$	$(4,4)$														$\times$
$\phi^4 D^2$	$(4,4)$				$\times$					$\times$		$\times$			$\times$
ϕ^6	$(6,6)$	$\times^*$		$\times$	$\times$		$\times^*$		$\times$	$\times$		$\times$			

Table II. Anomalous dimension matrix for dimension six operators in a general quantum field theory. The shaded entries vanish by our non-renormalization theorems, in full agreement with [2]. Here y^2 and $\bar{y}^2$ label entries that are non-zero due to non-holomorphic Yukawa couplings, $\times$ labels entries that vanish because there are no diagrams [13], and $\times^*$ labels entries that vanish by a combination of counterterm analysis and our non-renormalization theorems.

From C.Cheung and C.Shen, arXiv: 1505.01844

UV properties of HEFT dimension-six operators from derivative interactions

We restrict ourselves to the Higgs potential sector.
In the first approximation we do not couple the scalars with the gauge bosons and the fermions

$$S = \int d^4x \left[\frac{1}{2}\partial^\mu \sigma \partial_\mu \sigma + \frac{1}{2}\partial^\mu \phi_a \partial_\mu \phi_a - \frac{M^2}{2} X_2^2 + \frac{1}{v}(X_1 + X_2)\Box\left(\frac{1}{2}\sigma^2 + v\sigma + \frac{1}{2}\phi_a^2 - vX_2\right)\right].$$

$\Phi = \frac{1}{\sqrt{2}}\begin{pmatrix} i\phi_1 + \phi_2 \\ \sigma + v - i\phi_3 \end{pmatrix}$ SU(2) doublet

X_2 SU(2) singlet

Equations of motion for the auxiliary field

$$\frac{\delta S}{\delta X_1} = \frac{1}{v}\Box\Big(\frac{1}{2}\sigma^2 + v\sigma + \frac{1}{2}\phi_a^2 - vX_2\Big)$$

Going on-shell with the auxiliary field
and neglecting zero modes of the Laplacian
we obtain

$$X_2 = \frac{1}{2v}\sigma^2 + \sigma + \frac{1}{2v}\phi_a^2$$
$$= \Phi^\dagger\Phi - \frac{v^2}{2}$$

On-shell equivalence with the standard formulation

By substituting the e.o.m. solution for the singlet field in the classical action one gets back the ordinary quartic potential

$$S|_{\text{on-shell}} = \int d^4x \left[\frac{1}{2}\partial^\mu\sigma\partial_\mu\sigma + \frac{1}{2}\partial^\mu\phi_a\partial_\mu\phi_a - \frac{1}{2}\frac{M^2}{v^2}\left(\frac{1}{2}\sigma^2 + v\sigma + \frac{1}{2}\phi_a^2\right)^2\right]$$

with coupling constant

$$\lambda = \frac{1}{2}\frac{M^2}{v^2}$$

right sign of the quartic potential needed to ensure stability from the sign of the mass term, in turn fixed by the requirement of the absence of tachyons

Dangereous interactions for renormalizability

The model contains derivative interactions of the schematic form

$$\chi \Box \chi^2$$

i.e. an operator of dimension 5.

Renormalizability?

More symmetries needed

BRST implementation of the on-shell constraint

Off-shell there is one more scalar field X_1.
What about this field? Physical or unphysical?

BRST symmetry (it does not originate from gauge invariance)

$$sX_1 = vc\,, \qquad sc = 0\,, \qquad s\sigma = s\phi_a = sX_2 = 0\,,$$

$$s\bar{c} = \frac{1}{2}\sigma^2 + v\sigma + \frac{1}{2}\phi_a^2 - vX_2\,.$$

Ghost action

$$S_{ghost} = -\int d^4x\, \bar{c}\Box c\,.$$

Invariance under the nilpotent BRST symmetry formally associated with a $U(1)_{constr}$ group

A.Q., Phys.Rev. D73 (2006) 065024

Full tree-level vertex functional

$$\Gamma^{(0)} = \int d^4x \Big[\frac{1}{2}\partial^\mu\sigma\partial_\mu\sigma + \frac{1}{2}\partial^\mu\phi_a\partial_\mu\phi_a - \frac{M^2}{2}X_2^2$$
$$- \bar{c}\Box c + \frac{1}{v}(X_1 + X_2)\Box\Big(\frac{1}{2}\sigma^2 + v\sigma + \frac{1}{2}\phi_a^2 - vX_2\Big)$$
$$+ \bar{c}^*\Big(\frac{1}{2}\sigma^2 + v\sigma + \frac{1}{2}\phi_a^2 - vX_2\Big)\Big].$$

Antifield

Slavnov-Taylor identity

$$\int d^4x \left(vc\frac{\delta\Gamma}{\delta X_1} + \frac{\delta\Gamma}{\delta\bar{c}^*}\frac{\delta\Gamma}{\delta\bar{c}}\right) = 0$$

Propagators

The quadratic part is diagonalized by $\sigma = \sigma' + X_1 + X_2$

$$\Delta_{\sigma'\sigma'} = \frac{i}{p^2}\,, \quad \Delta_{\phi_a\phi_b} = \frac{i\delta_{ab}}{p^2}\,, \quad \Delta_{\bar{c}c} = \frac{i}{p^2}$$

$$\Delta_{X_1X_1} = -\frac{i}{p^2}\,, \quad \Delta_{X_2X_2} = \frac{i}{p^2 - M^2}\,.$$

The derivative interaction only depends on $X = X_1 + X_2$,
whose propagator has an improved UV behaviour

$$\Delta_{XX} = \frac{iM^2}{p^2(p^2 - M^2)}$$

Thus the derivative interaction
is harmless and p.c. renormalizability still holds

Physical states

They are described by the standard BRST construction of the physical Hilbert space in terms of the asymptotic charge Q

$$\mathcal{H}_{phys} = \ker Q / \mathrm{Im}\, Q$$

The scalar singlet is physical:

$$X_1, \sigma' \notin \mathcal{H}_{phys}$$

$$X_2 \in \mathcal{H}_{phys}$$

$$c = \{Q, X_1\}\,, \qquad [Q, \bar{c}] = \sigma' + X_1$$

The only physical degrees of freedom are the scalar X_2 and the ϕ_a components of the SU(2) doublet

Functional identities & Renormalization

Ghost and antighost equation

$$\frac{\delta\Gamma}{\delta\bar{c}} = -\Box c\,, \frac{\delta\Gamma}{\delta c} = \Box\bar{c}$$

X_1 equation

$$\frac{\delta\Gamma}{\delta X_1} = \frac{1}{v}\Box\frac{\delta\Gamma}{\delta\bar{c}^*}$$

X_2 equation & shift symmetry $\delta X_1(x) = \alpha(x)\,, \quad \delta X_2(x) = -\alpha(x)$

$$\frac{\delta\Gamma}{\delta X_2(x)} = \frac{1}{v}\Box\frac{\delta\Gamma}{\delta\bar{c}^*} + \Box(X_1 + X_2) - M^2 X_2 - v\bar{c}^*$$

Global SU(2) invariance

$$\int d^4x \Big[-\frac{1}{2}\alpha_a\phi_a\frac{\delta\Gamma}{\delta\sigma} + \Big(\frac{1}{2}(\sigma+v)\alpha_a + \frac{1}{2}\epsilon_{abc}\phi_b\alpha_c\Big)\frac{\delta\Gamma}{\delta\phi_a}\Big] = 0\,.$$

Stability of the theory

The ghosts remain free fields also at the quantum level.

The X_1 and X_2 equations fix the amplitudes
involving at least one X_1 and X_2 fields in terms of amplitudes without.

The latter have a good UV behaviour

Global SU(2) invariance is a symmetry of the theory

UV indices

Field/ext. source	UV index
$\bar{c}^*$	2
σ'	1
ϕ_a	1

There is no power-counting with positive UV index on the physical singlet and the auxiliary field X_ι, however they are controlled by the functional identities in terms of ancestor amplitudes with external legs with positive UV dimensions

Counterterms for ancestor amplitudes

$$\mathcal{L}_{ct} = \mathcal{L}_{ct,1} + \mathcal{L}_{ct,2}$$

$$\begin{aligned}\mathcal{L}_{ct,1} = &- \mathcal{Z}\Big(\frac{1}{2}\partial^\mu\sigma\partial_\mu\sigma + \frac{1}{2}\partial^\mu\phi_a\partial_\mu\phi_a\Big) - \mathcal{M}\Big(\frac{1}{2}\sigma^2 + v\sigma + \frac{1}{2}\phi_a^2\Big) \\ &- \mathcal{G}\Big(\frac{1}{2}\sigma^2 + v\sigma + \frac{1}{2}\phi_a^2\Big)^2 - \mathcal{R}_1\bar{c}^*\Big(\frac{1}{2}\sigma^2 + v\sigma + \frac{1}{2}\phi_a^2\Big).\end{aligned}$$

$$\mathcal{L}_{ct,2} = -\mathcal{R}_2\bar{c}^* - \frac{1}{2}\mathcal{R}_3(\bar{c}^*)^2$$

SM with the Derivative Representation of the Higgs potential

The coupling to the gauge fields and the fermions happens through the doublet

$$\partial_\mu \Phi^\dagger \partial^\mu \Phi \to D_\mu \Phi^\dagger D^\mu \Phi$$

$$-G_Y \bar{\psi}_L \Phi \psi_R - \tilde{G}_Y \bar{\psi}_L \psi_R \Phi^C + \text{h.c.}$$

The mass term for X_2, the (X_1+X_2)-dependent sector as well as the ghosts and antighosts associated with the SU(2) constraint are left unchanged.

Custodial Symmetry

Still true in the derivative formulation
of the Higgs potential

$$\Omega = (\Phi^C, \Phi)$$

$$\frac{1}{2}\sigma^2 + v\sigma + \frac{1}{2}\phi_a^2 = \frac{1}{2}\mathrm{Tr}\ (\Omega^\dagger\Omega)$$

By requiring X_1, X_2 to be invariant under global $SU(2)_L$ x $SU(2)_R$
we recover the usual custodial symmetry

$$\Omega \to L\Omega R^\dagger$$

in the limit where the hypercharge coupling vanishes.

SU(2) x U(1) x $U(1)_{constr}$ BRST symmetry

The full BRST symmetry of the theory is

$$s = s_0 + s_1$$

s_0 is the SU(2)xU(1) BRST symmetry of the SM
s_1 is the $U(1)_{constr}$ BRST symmetry

$$s_0\chi = 0 \qquad \text{for } \chi \in \{X_2, X_1, \bar{c}, c\}$$

$$s_1\chi = 0 \qquad \text{on gauge fields, fermions and the } \Phi \text{ doublet}$$

$$s^2 = 0\,, \qquad s_0^2 = s_1^2 = 0\,, \qquad \{s_0, s_1\} = 0$$

The z-model

The X_2 equation is not the most general functional symmetry holding true for the vertex functional.
The breaking term on the R.H.S. of the shift symmetry stays linear in the quantum fields even if one adds a kinetic term for the scalar singlet

$$\int d^4x \, \frac{z}{2} \partial^\mu X_2 \partial_\mu X_2$$

Upon integration over the auxiliary field this is equivalent to the addition of the dimension-six operator

$$\int d^4x \, \frac{z}{v^2} \partial_\mu \Phi^\dagger \Phi \partial^\mu \Phi^\dagger \Phi$$

Modified identities and propagators

The X_2 equation changes

$$\frac{\delta\Gamma}{\delta X_2(x)} = \frac{1}{v}\Box\frac{\delta\Gamma}{\delta\bar{c}^*} + \Box(X_1 + (1-z)X_2) - M^2 X_2 - v\bar{c}^*$$

The X_2 propagator becomes

$$\Delta_{X_2 X_2} = \frac{i}{(1+z)p^2 - M^2}$$

and p.c. renormalizability is lost, as a consequence of the worsened UV behaviour of the X field

$$\Delta_{XX} = \frac{i(-zp^2 + M^2)}{p^2[(1+z)p^2 - M^2]}$$

z-dependence of the vertex functional

$$\frac{\partial \Gamma}{\partial z} = \int d^4x \, \frac{\delta \Gamma}{\delta R(x)}$$

$R(x)$ coupled to $-\frac{1}{2} X_2 \Box X_2$

Thus we can expand the non p.c.-renormalizable theory around z=0 in terms of insertion of the kinetic operator for X_2 at zero momentum

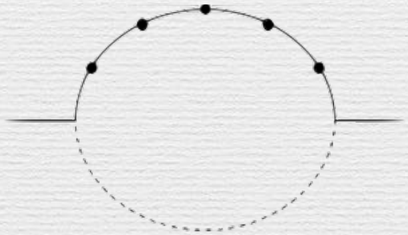


FIG. 1: Mass insertions (the dots) on a X_2 propagator (drawn as a solid line)

z-dependence from tree-level Feynman rules

Rescale the X_2 lines by a factor $1/(1+z)$
and redefine the mass through

$$M^2 \to M^2/(1+z)$$

Side remark:
No power-counting renormalizability,
however no freedom to add finite counterterms
vanishing at z=0
(otherwise further sources of z-dependence would be added)

Application: anomalous trilinear Higgs coupling

G.Degrassi, P.Giardino, F.Maltoni, D.Pagani
arXiv:1607.04251

$$V_{H^3} = \lambda_3 v H^3$$

induced e.g. by an anomalous dim. 6 operator $(\Phi^\dagger\Phi)^3$

Mixing with $\partial_\mu(\Phi^\dagger\Phi)\partial^\mu(\Phi^\dagger\Phi)$?

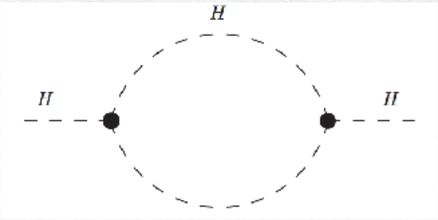
Gauge invariance?

NLO corrections

According to arXiv:1607.04251 there are two types of corrections:

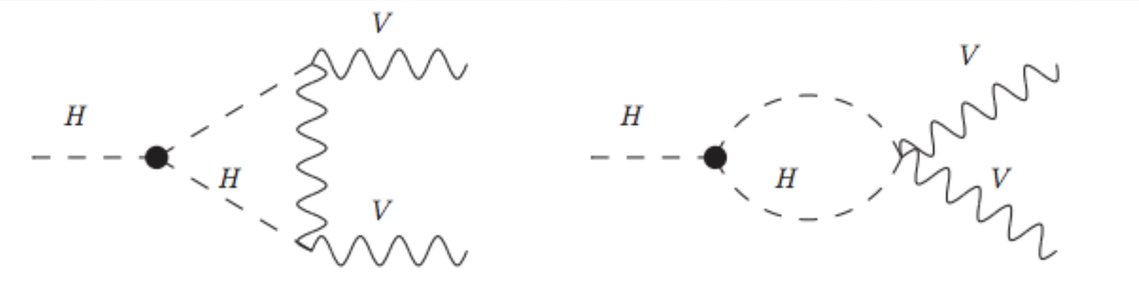
$$\Sigma_{NLO} = Z_H \Sigma_{LO}(1 + \kappa_\lambda C_1)$$

Z_H is related to the wave function correction associated with the diagram

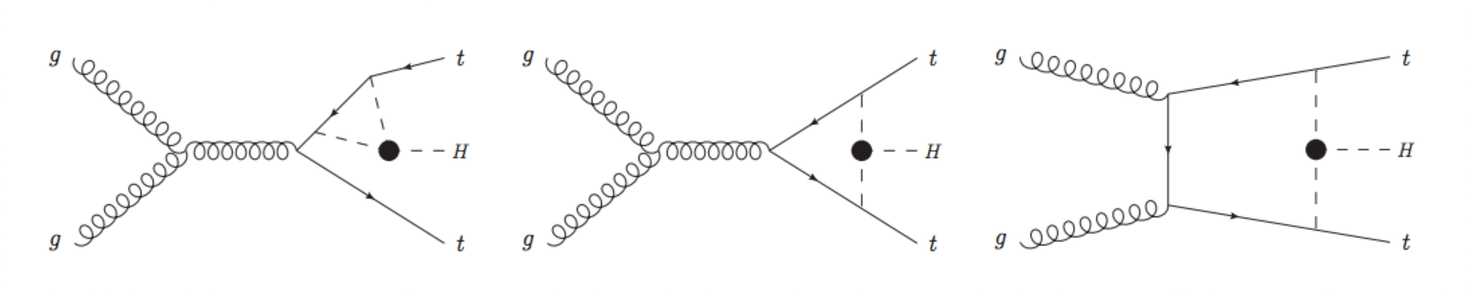


C_1 is a process-dependent coefficient arising from the interference from one-loop virtual EW corrections

Examples of EW vertex corrections



Processes involving massive vector bosons



Contributions to $t\bar{t}H$

The X2-equation in the presence of a cubic interaction

$$g_6 X_2^3$$

One needs one more external source L
to define the X2-equation in the presence of a cubic interaction

$$\int d^4x \, L X_2^2$$

The X2-equation becomes

$$\frac{\delta\Gamma}{\delta X_2} = \frac{1}{v}\Box\frac{\delta\Gamma}{\delta\bar{c}^*} + 3g_6\frac{\delta\Gamma}{\delta L} + \Box(X_2 + 1(1-z)X_2) - M^2 X_2 - v\bar{c}^* + 2LX_2$$

valid to all orders in the loop expansion

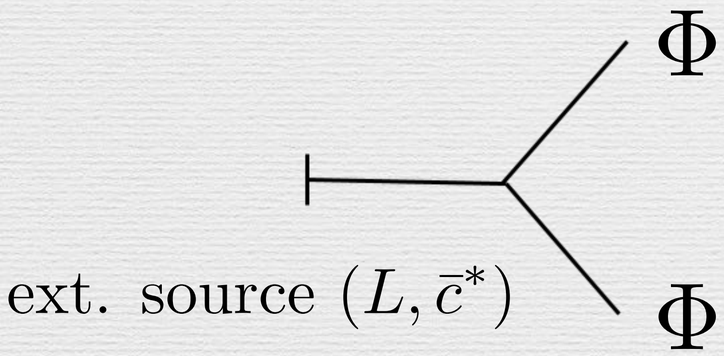
The two-point function

$$\Gamma^{(j)}_{X_2(-p)X_2(p)} = \frac{1}{v^2} p^4 \Gamma^{(j)}_{\bar{c}^*(-p)\bar{c}^*(p)} - \frac{3g_6}{v} p^2 \Gamma^{(j)}_{L(-p)\bar{c}^*(p)} - \frac{3g_6}{v} p^2 \Gamma^{(j)}_{L(p)\bar{c}^*(-p)} + 9g_6^2 \Gamma^{(j)}_{L(-p)L(p)}$$

$\sim_{j=1} (\partial_\mu X_2 \partial^\mu X_2)^2$

$\sim_{j=1} (\partial_\mu X_2)^2$

$\sim_{j=1} X_2^2$



At one loop level (j=1) the amplitudes involving the external sources on the r.h.s. are logarithmically divergent

One generates divergences for monomials involving X2 and up to four derivatives

Basis dependence

The form of the X2-equation depends on the basis.
It simplifies if one uses

$$X = X_1 + X_2, X_2, \sigma$$

$$\frac{\delta\Gamma}{\delta X_2} = 3g_6 \frac{\delta\Gamma}{\delta L} - \Box X - M^2 X_2 - z\Box X_2 - v\bar{c}^*$$

so that one has for the two-point function

$$\Gamma^{(j)}_{X_2(-p)X_2(p)} = 9g_6^2 \Gamma^{(j)}_{L(-p)L(p)} \sim_{j=1} X_2^2$$

However one needs to consider mixed propagators

$$\Delta_{XX_2}, \Delta_{\sigma X_2}$$

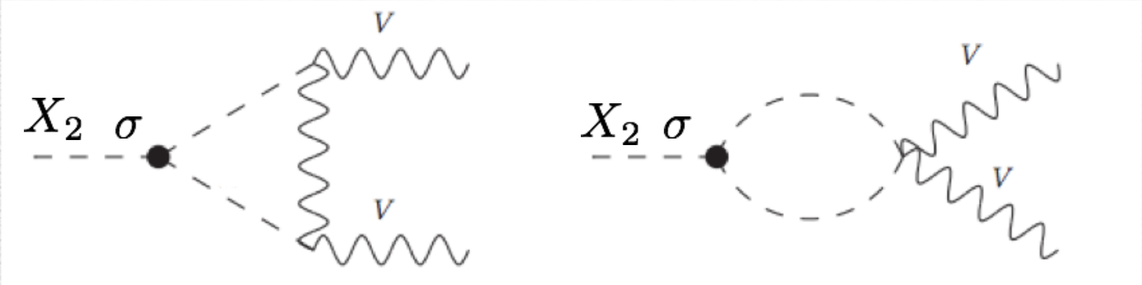
and mixed 1-PI Green's functions $\Gamma^{(j)}_{\sigma X_2}, \Gamma^{(j)}_{XX_2}$
yielding back the divergences described before

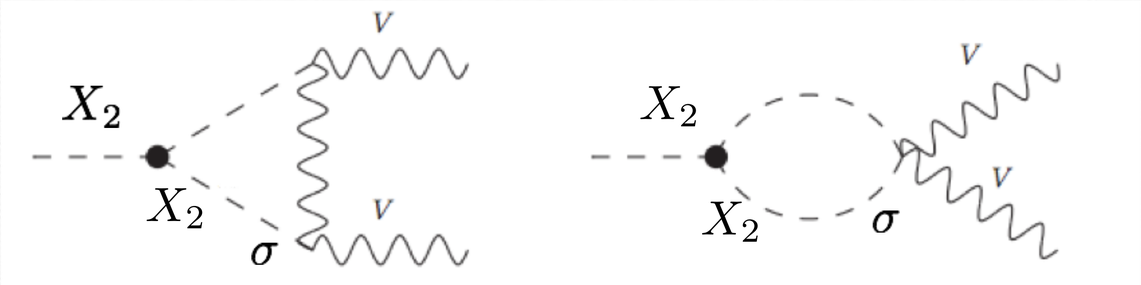
Vertex functions

Anomalous coupling associated with $g_6 X_2^3$

SM contribution controlled by the derivative interaction and the sigma-dependent part of the action

At the one-loop level diagrammatic (gauge-invariant) separation of the SM and BSM part in each process





Sample SM and BSM diagrams

UV divergences in vertex functions

Again controlled by the X2-equation

More complicated relations
involving external sources insertions

The latter have better UV properties
than the field X2

Valid off-shell, to any order in the loop expansion

Outlook

- A particular dimension-six operator has special UV properties in the HEFT extensions of the SM
- The latter properties are transparent in the derivative Higgs interaction formalism (application to the study of anomalous interactions)
- The non-power-counting renormalizable theory is constructed as an expansion around the p.c. renormalizable theory at z=0