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Off-shell renormalization of spontaneously broken effective gauge theories

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The off-shell renormalization of spontaneously broken effective gauge theories is carried out following the Algebraic Renormalization approach. We provide a constructive way to obtain the canonical transformations with respect to the Batalin-Vilkovisky bracket, associated with the gauge group of the theory, that generalize the familiar linear field redefinitions arising in power-counting renormalizable theories. By taking into account such generalized field redefinitions one can consistently extract the β-functions of the higher dimensional operators order by order in the loop expansion. The gauge invariance of the ultraviolet coefficients for physical operators is explicitly checked at one loop order.

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1. Introduction

Lack of discovery of new signals of physics Beyond the Standard Model (BSM) at the current LHC experimental program has triggered in recent years a huge amount of work on the theory side. In particular a prominent role has been played by the so-called Standard Model effective field theories (SMEFTs), i.e. models whose interactions can expanded on a basis of higher dimensional operators suppressed by some large energy scale Λ (for recent reviews see, e.g., [1, 2]). The classical Lagrangian of these models takes the form

$$\mathscr{L}_{SMEFT} = \mathscr{L}_{SM} + \sum_i \frac{c_i^{[5]}}{\Lambda} \mathscr{O}_i^{[5]} + \sum_i \frac{c_i^{[6]}}{\Lambda^2} \mathscr{O}_i^{[6]} + \cdots \tag{1.1}$$

where $\mathscr{L}_{SM}$ denotes the SM Lagrangian and $\mathscr{O}_i^{[k]}$ is a dimension k gauge invariant operator with coefficient $c_i^{[k]}$. Usually in the literature operators up to dimension 6 are considered [3, 4, 5, 6].

Due to the presence of higher-dimensional interaction vertices generated by $\mathscr{O}_i^{[k]}$, SMEFTs are no more power-counting renormalizable, hence an infinite number of counterterms is required to subtract the ultraviolet (UV) divergences.

In a seminal paper [7] Weinberg and Gomis have provided a complete characterization of these counterterms in terms of the quantum symmetries holding true for the model (1.1).

Let us denote collectively by $\Phi = \{\Phi^I\}$ the fields of the theory and by $\Phi^* = \{\Phi^*_I\}$ the corresponding antifields. For the models we are interested in, Φ^* are the external sources coupled at the classical level to the BRST transformations $s\Phi$ of the quantized fields Φ.

One can then introduce the Batalin-Vilkovisky (BV) [8, 9] bracket[1]

$$(X,Y) = \int d^4x \sum_I \Big[\frac{\delta_r X}{\delta \Phi^I} \frac{\delta_l Y}{\delta \Phi^*_I} - \frac{\delta_r X}{\delta \Phi^*_I} \frac{\delta_l Y}{\delta \Phi^I} \Big] . \tag{1.2}$$

Φ and Φ^* can be thought of as conjugate variables since they obey the fundamental BV bracket

$$(\Phi^I, \Phi^*_J) = \delta^I_J ; \qquad (\Phi^I, \Phi^J) = (\Phi^*_I, \Phi^*_J) = 0 . \tag{1.3}$$

A field-antifield redefinition $\Phi' = \Phi'(\Phi, \Phi^*), \Phi^{*'} = \Phi^{*'}(\Phi, \Phi^*)$ is said to be canonical if it preserves the BV bracket, namely

$$(\Phi^{'I}, \Phi^{*'}_J) = \delta^I_J , \qquad (\Phi^{'I}, \Phi^{'J}) = (\Phi^{*'}_I, \Phi^{*'}_J) = 0 . \tag{1.4}$$

Let us denote by $\Gamma^{(0)}$ the tree-level vertex functional obtained from the tree-level action

$$S_0 = \int d^4x \, \mathscr{L}_{SMEFT} + \sum_I \int d^4x \, \Phi^*_I s\Phi^I \tag{1.5}$$

after carrying out the gauge-fixing through a suitable canonical transformation [10]. The full vertex functional $\Gamma = \Gamma^{(0)} + \sum_{j=1}^\infty \hbar^j \Gamma^{(j)}$ then obeys the BV master equation

$$(\Gamma, \Gamma) = 0 , \tag{1.6}$$

[1] The subscripts l, r denote the left- and right- functional derivatives respectively

that supersedes gauge invariance at the quantum level [10].

In [7] it has been proven in full generality that the model can indeed be renormalized in the modern sense: all divergences can be removed order by order in the loop expansion by adding a suitable set of counterterms (possibly infinitely many) in such a way to preserve the BV master equation (1.6).

More specifically, the renormalized action S can be written in an expansion in $\hbar$ as

$$S_0 = S + \hbar\Delta_1 + \frac{1}{2}\hbar^2\Delta_2 + \cdots \,, \tag{1.7}$$

where Δ_j contain the gauge-invariant operators arising at order j in the loop expansion. S_0 in Eq. (1.7) is not enough to remove all the UV divergences of the SMEFTs. One has also to take into account the generalization of the familiar linear field redefinitions of power-counting renormalizable theories [11]. This can be done by an appropriate field-antifield redefintion $\Phi \to \Phi'(\Phi,\Phi^*)$, $\Phi^* \to \Phi^{*'}(\Phi,\Phi^*)$ that respects the fundamental BV bracket and is thus induced by a canonical transformation with generator $F(t) = \hbar t F_1 + \frac{1}{2}\hbar^2 t^2 F_2 + \cdots$ [7].

Then, the transformed bare action finally takes the form

$$S_0' = S + \hbar[\Delta_1 + (F_1,S)] + \frac{1}{2}\hbar^2\Big[\Delta_2 + 2(F_1,\Delta_1) + (F_2,S) + (F_1,(F_1,S)))\Big] + \cdots \tag{1.8}$$

and is sufficient to remove all the UV divergences of the theory [7]. Notice that no restriction to on-shell quantities, customary in the EFT approach, is required.

The issue is then to work out a procedure to determine the Δ_i's and the F_i's in Eq. (1.8). This task has been solved in a series of papers where we have carried out the off-shell renormalization of spontaneously broken effective gauge theories by using aconstructive approach to Algebraic Renormalization [12, 13, 14, 15, 16, 17, 18].

This boils down to solve the BV master equation (equivalently, the Slavnov-Taylor (ST) identity, as is also called in the literature [19, 20]) for effective gauge theories that bear a spontaneously broken phase.

This can be achieved as follows. Let us consider first the Abelian Higgs-Kibble model [21] supplemented by a typical dimension 6 derivative operator

$$\frac{g}{\Lambda^2}\phi^\dagger\phi(D^\mu\phi)(D_\mu\phi)\,, \tag{1.9}$$

$\phi = \frac{1}{\sqrt{2}}(\sigma + v + i\chi)$ being the Higgs scalar with vacuum expectation value v; g the coupling constant of the non power-counting renormalizable operator; and $D_\mu = \partial_\mu - ieA_\mu$ the covariant derivative with respect to the Abelian gauge field A_μ.

The operator (1.9) contains a vertex of the form

$$\sigma^2\partial^\mu\sigma\partial_\mu\sigma\,,$$

that gives rise to UV divergent diagrams with an arbitrary number of external legs already at one loop, as shown in Figure 1.

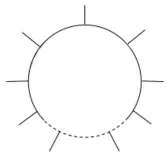

Figure 1: A UV divergent one-loop diagram with an infinite number of external legs. Derivatives from the vertex $\sigma^2\partial^\mu\sigma\partial_\mu\sigma$ act on the internal propagators.

Nevertheless, there are relations between the UV divergent parts of the one-loop amplitudes; but they become transparent only after one uses as a dynamical variable to describe the physical component of the scalar field ϕ the gauge invariant combination

$$X_2 \sim \frac{1}{v}\Big(\phi^\dagger\phi - \frac{v^2}{2}\Big). \tag{1.10}$$

This happens because the dependence of the vertex functional Γ on the field X_2 can be strongly restricted by a set of functional identities that involve external sources with a better UV behaviour, thus reducing the number of independent UV divergences [22, 23, 24, 25].

We name the formulation of the model based on the X_2-field the X-theory. The main features of the X-theory are summarized in Sect. 2. In particular it is relatively straightforward to obtain the canonical transformation F that controls the generalized field redefinitions (GFRs) leading to Eq. (1.8). Once the GFRs are known, one can determine the gauge invariant operators entering in the Δ_j's in Eq. (1.8) and explicitly check that their coefficients do not depend on the gauge.

After completing the renormalization of the X-theory one eventually goes on-shell with the X_2-field and its companion X_1, the field that plays the role of the Lagrange multiplier enforcing the constraint (1.10). It can be proven that one gets back the original amplitudes in the ϕ-formalism [25]. The latter are however decomposed into different pieces corresponding to different sectors of the starting X-theory that exhibit in a more transparent way the relations between the UV divergences of the non-power-counting renormalizable theory. For instance one can prove that an arbitrary derivative-free analytic scalar potential only depending on the combination (1.10) can be renormalized in terms of 11 independent parameters [23].

In the Abelian case the functional identities underlying the X-formalism provide an alternative explanation of the cancellation patterns that have been observed in the literature and traced back to holomorphicity [26, 27] or remnants of supersymmetry [28].

We have carried out our computations at one loop order with the full dependence on the higher dimensional coupling g, i.e. we are not limited to the usual linear approximation in the couplings. Moreover the procedure is based on a parameterization of the gauge-invariant operators in terms of the so-called contractible pairs [29] that lends itself to a direct generalization to a non-Abelian gauge group. In particular the solutions to the ST identity in the non-Abelian case have been discussed in [30].

2. The X-theory

The tree-level vertex functional in the X-formalism can be written as [25]

$$\begin{aligned}\Gamma^{(0)} = \int d^4x \Big[& -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} + (D^\mu\phi)^\dagger(D_\mu\phi) - \frac{M^2-m^2}{2}X_2^2 - \frac{m^2}{2v^2}\Big(\phi^\dagger\phi - \frac{v^2}{2}\Big)^2 \\ & -\bar{c}(\Box+m^2)c + \frac{1}{v}(X_1+X_2)(\Box+m^2)\Big(\phi^\dagger\phi - \frac{v^2}{2} - vX_2\Big) \\ & + \frac{gv}{\Lambda^2}X_2(D^\mu\phi)^\dagger(D_\mu\phi) + T_1(D^\mu\phi)^\dagger(D_\mu\phi) \\ & + \frac{\xi b^2}{2} - b\Big(\partial A + \xi ev\chi\Big) + \bar{\omega}\Big(\Box\omega + \xi e^2 v(\sigma+v)\omega\Big) \\ & + \bar{c}^*\Big(\phi^\dagger\phi - \frac{v^2}{2} - vX_2\Big) + \sigma^*(-e\omega\chi) + \chi^* e\omega(\sigma+v)\Big]. \end{aligned} \tag{2.1}$$

The first line is the action of the Abelian Higgs-Kibble model in the X-formalism. The gauge field A_μ has mass $M_A = ev$. As explained above, one also adds a singlet field X_2 that provides a gauge-invariant parameterization of the physical scalar mode. When going on-shell with the field X_1, that plays the role of a Lagrange multiplier, one recovers the constraint[2] $X_2 \sim \frac{1}{v}(\phi^\dagger\phi - v^2/2)$. Inserting the latter back into the first line of Eq. (2.1), the m^2-term cancels out and one is left with the usual Higgs quartic potential with coefficient $\sim M^2/2v^2$.

Thus Green's functions in the target theory have to be m^2-independent, a fact that can be used as a very strong check of the computations, since m^2 appears both in Feynman amplitudes and in the associated invariants.

The decoupling of the unphysical Lagrange multiplier X_1 is guaranteed by the invariance of the vertex functional under the following constraint BRST differential [25, 31]

$$\mathcal{s}X_1 = vc; \quad \mathcal{s}\phi = \mathcal{s}X_2 = \mathcal{s}c = 0; \quad \mathcal{s}\bar{c} = \phi^\dagger\phi - \frac{v^2}{2} - vX_2\,. \tag{2.3}$$

The constraint ghost and antighost $c, \bar{c}$ remain free.

The BRST symmetry associated with the U(1) gauge invariance reads

$$\begin{aligned} sA_\mu = \partial_\mu\omega; \quad s\omega = 0; \quad s\bar{\omega} = b; \quad sb = 0; \quad s\phi = ie\omega\phi\,, \\ s\sigma = -e\omega\chi; \quad s\chi = e\omega(\sigma+v)\,. \end{aligned} \tag{2.4}$$

The gauge-fixing in the fourth line of Eq. (2.1) is carried out *à la* BRST by introducing the Nakanishi-Lautrup field b and the pair of ghost, antighost fields $\omega, \bar{\omega}$. The last line of Eq. (2.1) includes the couplings with the antifields σ^*, χ^* associated to the non linear BRST transformations

[2] Going on-shell with X_1 yields the condition

$$(\Box+m^2)\Big(\phi^\dagger\phi - \frac{v^2}{2} - vX_2\Big) = 0\,, \tag{2.2}$$

so that the most general solution is $X_2 = \frac{1}{v}\Big(\phi^\dagger\phi - \frac{v^2}{2}\Big) + \eta$, η being a scalar field of mass m. However it can be proven that in perturbation theory the correlators of the field η with any gauge-invariant operator are zero [25], and so one can safely set $\eta = 0$.

$s\sigma, s\chi$. We remark that one can avoid the introduction of the antifield for the gauge connection A_μ as well as for the antighost $\bar{\omega}$ since their BRST transformation is linear in the quantum fields and thus it is not subject to an independent renormalization. The external source $\bar{c}^*$ is coupled to the BRST transformation $\mathcal{s}\bar{c}$ in Eq. (2.3).

Finally in the third line of Eq. (2.1) we have introduced the higher dimensional operator[3]

$$\frac{gv}{\Lambda^2} X_2 (D^\mu \phi)^\dagger (D_\mu \phi) \tag{2.5}$$

that reduces, by going on-shell with X_1 and thus implementing the constraint $X_2 \sim \frac{1}{v}\left(\phi^\dagger\phi - \frac{v^2}{2}\right)$, to $\frac{g}{\Lambda^2}\left(\phi^\dagger\phi - \frac{v^2}{2}\right)(D^\mu\phi)^\dagger(D_\mu\phi)$.

The dependence on $X_{1,2}$ is governed by the $X_{1,2}$-equations of motion

$$\frac{\delta\Gamma}{\delta X_1} = \frac{1}{v}(\square + m^2)\frac{\delta\Gamma}{\delta\bar{c}^*}, \tag{2.6}$$

$$\frac{\delta\Gamma}{\delta X_2} = \frac{1}{v}(\square + m^2)\frac{\delta\Gamma}{\delta\bar{c}^*} + \frac{gv}{\Lambda^2}\frac{\delta\Gamma}{\delta T_1} - (\square + m^2)X_1 - (\square + M^2)X_2 - v\bar{c}^* . \tag{2.7}$$

The X_2-equation requires in particular the introduction of the source T_1 in the third line of Eq. (2.1). Eqs.(2.6) and (2.7) fix 1-PI amplitudes with at least one $X_{1,2}$-external leg in terms of 1-PI amplitudes with the insertion of somehow better behaved external sources: $\bar{c}^*$ has UV degree 2 (at $T_1 = 0$), and thus only dimension 6 amplitudes with the insertion of up to three $\bar{c}^*$'s can be UV divergent (at one-loop level). On the other hand, amplitudes with insertion of the source T_1 are easy to resum, as we will show on the example of the amplitudes fixing the GFRs.

It is noteworthy that Eqs.(2.6) and (2.7) can be solved algebraically to all orders in the loop expansion by the following substitutions:

$$\bar{\mathcal{c}}^* = \bar{c}^* + \frac{1}{v}(\square + m^2)(X_1 + X_2); \qquad \mathcal{T}_1 = T_1 + \frac{gv}{\Lambda^2}X_2. \tag{2.8}$$

The vertex functional Γ obeys the ST identity (equivalently the BV bracket) for the model at hand [24, 25]:

$$\mathcal{S}(\Gamma) = \int d^4x \left[\partial_\mu\omega\frac{\delta\Gamma}{\delta A_\mu} + \frac{\delta\Gamma}{\delta\sigma^*}\frac{\delta\Gamma}{\delta\sigma} + \frac{\delta\Gamma}{\delta\chi^*}\frac{\delta\Gamma}{\delta\chi} + b\frac{\delta\Gamma}{\delta\bar{\omega}}\right] = 0. \tag{2.9}$$

At one loop order both $\Gamma^{(1)}$ and its UV divergent part $\overline{\Gamma}^{(1)}$ satisfies the linearized ST identity

$$\begin{aligned}\mathcal{S}_0(\overline{\Gamma}^{(1)}) = &\int \mathrm{d}^4x \left[\partial_\mu\omega\frac{\delta\overline{\Gamma}^{(1)}}{\delta A_\mu} + e\omega(\sigma + v)\frac{\delta\overline{\Gamma}^{(1)}}{\delta\chi} - e\omega\chi\frac{\delta\overline{\Gamma}^{(1)}}{\delta\sigma} + b\frac{\delta\overline{\Gamma}^{(1)}}{\delta\bar{\omega}}\right.\\ &\left.+\frac{\delta\Gamma^{(0)}}{\delta\sigma}\frac{\delta\overline{\Gamma}^{(1)}}{\delta\sigma^*} + \frac{\delta\Gamma^{(0)}}{\delta\chi}\frac{\delta\overline{\Gamma}^{(1)}}{\delta\chi^*}\right]\\ &= s\overline{\Gamma}^{(1)} + \int \mathrm{d}^4x \left[\frac{\delta\Gamma^{(0)}}{\delta\sigma}\frac{\delta\overline{\Gamma}^{(1)}}{\delta\sigma^*} + \frac{\delta\Gamma^{(0)}}{\delta\chi}\frac{\delta\overline{\Gamma}^{(1)}}{\delta\chi^*}\right],\end{aligned} \tag{2.10}$$

[3] w.r.t. the conventions of [24, 25] the coupling g has been rescaled by a factor $\frac{v}{\Lambda}$ in order to get a pre-factor $1/\Lambda^2$ for dimension 6 operators, as is customary in the literature.

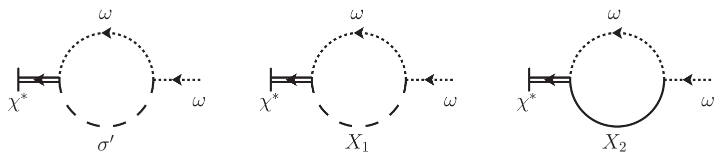

Figure 2: Diagrams contributing to the UV divergent 1-PI amplitude $\Gamma^{(1)}_{\chi^*\omega}$.

whose most general solution can be written as [24, 25]

$$\overline{\Gamma}^{(1)} = \overline{\mathscr{I}}^{(1)}_{\rm gi} + \mathcal{S}_0(\overline{Y}^{(1)}). \tag{2.11}$$

with $\overline{\mathscr{I}}^{(1)}_{\rm gi}$ a gauge-invariant functional [to be identified with $-\Delta_1$ in Eq. (1.8)] and a so-called cohomologically trivial piece $\mathcal{S}_0(\overline{Y}^{(1)})$ [to be identified with $-(F_1, S)$ in Eq. (1.8)].

3. Generalized Field Redefinitions

The functional $\overline{Y}^{(1)}$ can be easily evaluated by looking at the antifield-dependent 1-PI amplitudes. Let us consider what happens in the Feynman gauge. One starts with the sector at $T_1 = 0$ and finds that there is just one UV divergent amplitude, namely[4] $\Gamma^{(1)}_{\chi^*\omega}$. The Feynman diagrams contributing to this amplitudes are depicted in Fig. 2. We use dimensional regularization with $\varepsilon = 4 - D$. Repeated insertions of T_1 on the scalar lines can be straightforwardly resummed to give the result (in Feynman gauge $\xi = 1$)

$$\overline{Y}^{(1)}\Big|_{\xi=1} \supset \rho_1 \mathcal{S}_0 \int d^4x \, \frac{1}{1+T_1}(\sigma^*\sigma + \chi^*\chi), \qquad \rho_1 = \frac{M_A^2}{8\pi^2 v^2}\frac{1}{\varepsilon}. \tag{3.1}$$

4. X-theory renormalization

Once the functional $\overline{Y}^{(1)}$ has been determined and consequently the GFRs are known, one must work out the coefficients of the invariants parameterizing the gauge-invariant sector of the one-loop solution to the ST identity, namely $\overline{\mathscr{I}}^{(1)}_{\rm gi}$. We can restrict ourselves to those operators contributing under the mapping to operators of dimension up to six in the ϕ-theory (i.e., the original theory obtained by going on-shell with $X_{1,2}$).

The list of invariants has been obtained in [24]. There are: 12 linearly independent invariants in the sector with pure external sources and their ordinary derivatives; 13 linearly independent invariants in the mixed fields-external sources sector; and 10 invariants depending only on the fields. Their coefficients can be fixed by projecting the invariants on a basis of Lorentz-covariant monomials in the fields, the external sources and their derivatives and by matching the UV divergences arising in the corresponding 1-PI amplitudes.

A technical but rather important point of the projection is the use of the so-called contractible pair basis [29], pairing together the symmetrized derivatives of the gauge connections with the derivatives of the ghost fields in the so-called BRST doublets [29, 32], that decouple from the gauge-invariant sector.

[4] Subscripts denote functional differentiation with respect to (w.r.t.) the arguments

5. Mapping

Once the renormalization of the X-theory is completed, one can go on shell with $X_{1,2}$ and recover the results on the original ϕ-theory one is interested in. The procedure is purely algebraic.

As already discussed, the X_1-equation of motion enforces the constraint $X_2 = \frac{1}{v}\Big(\phi^\dagger\phi - \frac{v^2}{2}\Big)$. Once one takes into account this constraint, the X_2-equation of motion yields

$$(\square + m^2)(X_1 + X_2) = -(M^2 - m^2)X_2 + \frac{gv}{\Lambda^2}(D^\mu\phi)^\dagger D_\mu\phi - v\bar{c}^* . \tag{5.1}$$

At one loop accuracy only tree-level equations of motions are needed. By substituting the above expressions for $X_{1,2}$ into the replacement rules (2.8) we arrive at the explicit form of the mapping transformation (at zero external sources):

$$\bar{c}^* \to -\frac{(M^2-m^2)}{v^2}\Big(\phi^\dagger\phi - \frac{v^2}{2}\Big) + \frac{g}{\Lambda^2}(D^\mu\phi)^\dagger D_\mu\phi ; \qquad \mathscr{T}_1 \to \frac{g}{\Lambda^2}\Big(\phi^\dagger\phi - \frac{v^2}{2}\Big). \tag{5.2}$$

Since the right-hand side of the above equation contains operators of dimension at least 2, in order to obtain target operators of up to dimension 6 it is clear that we need to consider amplitudes with up to 3 external sources $\bar{c}^*$ and T_1. They are included in the list detailed in Sect. 4.

Let us summarize the main results of this analysis.

- Under the mapping the cohomologically trivial invariant in Eq. (3.1) induces a generalized field redefinition given by:

$$\sigma \to \sigma + \frac{\rho_1}{1+\frac{g}{\Lambda^2}\Big(\phi^\dagger\phi - \frac{v^2}{2}\Big)}\sigma; \qquad \chi \to \chi + \frac{\rho_1}{1+\frac{g}{\Lambda^2}\Big(\phi^\dagger\phi - \frac{v^2}{2}\Big)}\chi. \tag{5.3}$$

  Hence we see that field redefinitions are not even polynomial even at one loop order and they contribute already at the linearized order in the coupling g. This is a crucial remark: if the effects of these non-linear field redefinitions are not properly taken into account, the renormalization of the coefficients of gauge-invariant operators cannot be correctly carried out.

- One can check by explicit computations that the coefficients of the gauge invariant operators are gauge-independent (as they should) only if the contributions from the GFRs to the UV divergent amplitudes are considered.

Let us illustrate how one can determine the UV coefficient of a dimension 6 operator by this technique. Suppose we are interested in the operator[5]

$$\mathscr{O}^{[6]}_{10} = \int d^4x \Big(\phi^\dagger\phi - \frac{v^2}{2}\Big)F^2_{\mu\nu} .$$

We denote by $\widetilde{\lambda}_{10}$ the coefficient of this operator. The tilde refers to the fact that it is an operator in the ϕ-theory. $\widetilde{\lambda}_{10}$ is determined under the mapping as follows

$$\begin{aligned}\widetilde{\lambda}_{10} &= -\frac{M^2-m^2}{v^2}\theta_9 + \frac{g}{\Lambda^2}\theta_{10} + \lambda_{10} \\ &= -\frac{\mu^{-\varepsilon}}{32\pi^2}\frac{g^2M_A^2}{\Lambda^4}\frac{1}{\varepsilon} ,\end{aligned} \tag{5.4}$$

[5] The subscript refers to the numbering used in [24]

where θ_9, θ_{10} are the coefficients of the X-theory operators $\int d^4x\, \bar{c}^* F_{\mu\nu}^2$ and $\int d^4x\, T_1 F_{\mu\nu}^2$ respectively (that do contribute to $\mathscr{O}_{10}^{[6]}$ by the action of the mapping (5.2) on the external sources $\bar{c}^*, T_1$). λ_{10} is the coefficient of the operator $\mathscr{O}_{10}^{[6]}$ in the X-theory.

One can see from Eq. (5.4) that i) the coefficient does not depend on ξ, i.e. it is gauge-invariant, as it should; ii) the dependence on the parameter m^2 has disappeared, with cancellations arising from the m^2-dependence of the coefficients θ_9, θ_{10} and λ_{10} and of the projection equation in the first line of Eq. (5.4).

It is then straightforward to determine the β-function of the operator $\mathscr{O}_{10}^{[6]}$ in the one-loop approximation:

$$\beta_{10} = (4\pi)^2 \frac{\mathrm{d}}{d\log\mu} \widetilde{\lambda}_{10}\,. \tag{5.5}$$

By taking into account only terms dependent on the higher dimensional coupling g we obtain

$$\beta_{10} \supset -\frac{g^2 M_A^2}{2\Lambda^4}\,. \tag{5.6}$$

Notice that this result goes beyond the linearized approximation in the higher dimensional couplings usually considered in the literature.

6. Conclusions

We have carried out the complete off-shell renormalization of the Abelian Higgs-Kibble model in the presence of a dimension 6 derivative-dependent operator within the Algebraic Renormalization approach. Cohomological tools play a crucial role in this class of models since they allow to disentangle non-linear generalized field redefinitions that do affect the Green's function of the theory. Several regularities in the UV behaviour of amplitudes that are not apparent in the conventional formalism can be made explicit by means of the X-formalism, where the physical scalar degree of freedom is described by the gauge invariant variable X_2. These regularities are encoded in certain functional identities obeyed by the vertex functional. The study of the renormalization of the Abelian Higgs-Kibble model supplemented by all parity-preserving dimension 6 operators is currently under way. This should be considered as a last step before tackling the analysis of the physically relevant case of the $\mathrm{SU}(2) \times \mathrm{U}(1)$ gauge group.

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