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Symmetry properties and renormalizability of massive gauge theories in Delbourgo-Jarvis gauge *

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Abstract

The massive Curci-Ferrari model in the nonlinear Delbourgo-Jarvis gauge can be renormalized analogously to Yang-Mills theories in linear gauges if besides the (anti-) Slavnov-Taylor and (anti-) Delduc-Sorella identities also the equations of motion for the auxiliary field and the (anti-)ghost fields as well as the so-called local Lee identity is required.

1 Massive gauge theory: the Curci-Ferrari model

In the algebraic BRST approach to gauge theories [1] renormalizability is proved by showing that the symmetries of the theory have no anomalies, i.e. that the related cohomologies are trivial. As is well known the anomaly coefficients cannot be determined algebraically. But, possibly the absence of anomalies may be shown after enlarging the symmetries of the theory in an appropriate manner. For the Landau gauge this is obtained by requiring invariance under BRST and anti-BRST transformations s and $\bar{s}$, respectively, as well as invariance under $sp(1)$ -transformations generated by the FP operator, the DJNO and the anti-DJNO operator q, d and $\bar{d}$, respectively.

If, however, massive gauge fields, e.g. with mass $m(s) = (1 - s)m$, are to be considered, as it is necessary at least intermediately (e.g. for Yang-Mills theory with instanton background) in the framework of the BPHZL-renormalization [2] to allow

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for an IR-regularization, then only nonlinear gauges are permissible¹. With respect to this the Curci-Ferrari model [3] in the (nonlinear) Delbourgo-Jarvis gauge, being the nonsingular analogue of the Landau gauge, will be considered.

In its most symmetric form the model is defined by the following (classical) action (for the explicit form see (11)):

$$S(g, m; \xi, \rho, \sigma) = -(2g)^{-2} \int d^4x \text{Tr}(F_{\mu\nu}(\mathcal{Q})F^{\mu\nu}(\mathcal{Q})) \quad (1)$$

$$+ (s_m \bar{s}_m + m^2) \int d^4x \text{Tr} \left(\frac{1}{2} \mathcal{Q}_\mu \mathcal{Q}^\mu + \xi \bar{\mathcal{C}} \mathcal{C} - G_\mu \mathcal{Q}^\mu + U \bar{\mathcal{C}} - \bar{U} \mathcal{C} + \frac{1}{2} \sigma G_\mu G^\mu + \rho \bar{U} U \right)$$

with

$$\mathcal{Q}_\mu = Q_\mu + \sigma G_\mu, \quad \mathcal{C} = C + \rho U, \quad \bar{\mathcal{C}} = \bar{C} + \rho \bar{U}, \quad \mathcal{B} = B + (\zeta - \bar{\zeta}) \rho S, \quad (2)$$

where $Q_\mu, C, \bar{C}$ and B are the gauge, ghost, antighost and Nakanishi-Lautrup field, respectively, and $G_\mu, U, \bar{U}$ and S are associated sources having the same quantum numbers as these fields, respectively; ξ is the gauge parameter, ζ an arbitrary real parameter ($\bar{\zeta} = 1 - \zeta$) reflecting the freedom in defining B according to $\mathcal{B} = \bar{\zeta} s_m \bar{\mathcal{C}} - \zeta \bar{s}_m \mathcal{C}$, whereas σ and ρ are to be introduced since the fields and their associated sources mix under renormalization.

The mass-dependent (anti-) BRST operators s_m (and $\bar{s}_m$) are defined through their action on the fields and the sources as well as the gauge parameter according to:

$$s_m \mathcal{Q}_\mu = D_\mu(\mathcal{Q}) \mathcal{C}, \quad s_m \mathcal{B} = \zeta [\mathcal{B} + \bar{\zeta} \{ \mathcal{C}, \bar{\mathcal{C}} \}, \mathcal{C}] - m^2 \mathcal{C},$$

$$s_m \mathcal{C} = -(1/2) \{ \mathcal{C}, \mathcal{C} \}, \quad s_m \bar{\mathcal{C}} = \mathcal{B} - \zeta \{ \mathcal{C}, \bar{\mathcal{C}} \}, \quad (3)$$

$$s_m G_\mu = P_\mu, \quad s_m P_\mu = 0, \quad s_m U = R, \quad s_m R = 0,$$

$$s_m \bar{U} = S, \quad s_m S = -m^2 U, \quad s_m \bar{P}_\mu = -m^2 G_\mu, \quad s_m \bar{R} = -2m^2 \bar{U}, \quad (4)$$

$$s_m \xi = \chi, \quad s_m \chi = 0, \quad s_m \bar{\chi} = -m^2 \xi. \quad (5)$$

The corresponding relations defining the action of the anti-BRST operator $\bar{s}_m$ are obtained by the following conjugation $\mathbf{C}$ which leaves the classical action (1) invariant, and obviously is an extension of the ghost-antighost conjugation:

$$C \rightarrow \bar{C}, \quad U \rightarrow \bar{U}, \quad P_\mu \rightarrow \bar{P}_\mu, \quad R \rightarrow \bar{R}, \quad S \rightarrow -S, \quad \zeta \rightarrow \bar{\zeta}, \quad \chi \rightarrow \bar{\chi},$$

$$\bar{C} \rightarrow -C, \quad \bar{U} \rightarrow -U, \quad \bar{P}_\mu \rightarrow -P_\mu, \quad \bar{R} \rightarrow R, \quad (S \equiv \bar{S}), \quad \bar{\zeta} \rightarrow \zeta, \quad \bar{\chi} \rightarrow \chi, \quad (6)$$

$$s_m \rightarrow \bar{s}_m, \quad \bar{s}_m \rightarrow -s_m, \quad \mathbf{d} \rightarrow \bar{\mathbf{d}}, \quad \bar{\mathbf{d}} \rightarrow \mathbf{d}, \quad \mathbf{q} \rightarrow -\mathbf{q}; \quad (7)$$

¹The well known fact that massive gauge theories, if renormalizable, violate unitarity does not matter here since after carrying out all necessary subtractions the limit $s \rightarrow 1$ is taken.

Q_μ, G_μ and B as well as ξ are $\mathbf{C}$ -even. The additional sources $\bar{P}_\mu$ and $\bar{R}$ together with S and $\bar{U}$ couple to the nonlinear parts of the BRST transformations (3) of the fields (see (11)) and analogously for the $\mathbf{C}$ -conjugate fields and sources; the source G_μ couples to $s_m \bar{s}_m Q_\mu$.

Let us now introduce the Delbourgo-Jarvis-Nakanishi-Ojima(DJNO) transformations [4] defined by:

$$dQ_\mu = 0, \quad d\bar{C} = C, \quad dC = 0, \quad dB = (\zeta - \bar{\zeta})C^2, \quad (8)$$

$$\begin{aligned} d\bar{P}_\mu &= P_\mu, \quad dP_\mu = 0, \quad d\bar{U} = U, \quad dU = 0, \\ dS &= R, \quad dR = 0, \quad d\bar{R} = 2S, \quad dG_\mu = 0, \end{aligned} \quad (9)$$

$$d\xi = 0, \quad d\bar{\chi} = \chi, \quad d\chi = 0. \quad (10)$$

Again the anti-DJNO transformations are obtained by $\mathbf{C}$ -conjugation.

This finishes the definition of our model whose action has the following explicit form:

$$\begin{aligned} S(g, m; \xi, \rho, \sigma) = & \quad (11) \\ - \frac{1}{4}g^{-2} \int d^4x \text{Tr}(F_{\mu\nu}(Q)F^{\mu\nu}(Q)) & + m^2 \int d^4x \text{Tr}\left(\frac{1}{2}Q_\mu Q^\mu + \xi \bar{C}C\right) \\ + \int d^4x \text{Tr}(\partial_\mu B Q^\mu & + \zeta \partial_\mu C D^\mu(Q)\bar{C} - \bar{\zeta} \partial_\mu \bar{C} D^\mu(Q)C) \\ + \xi \int d^4x \text{Tr}(m^2 \bar{C}C & + \frac{1}{2}(B - \zeta\{C, \bar{C}\})^2 + \frac{1}{2}(B + \bar{\zeta}\{\bar{C}, C\})^2) \\ + \int d^4x \text{Tr}(-\xi m^2 \bar{C}C & + \chi \bar{C}(B + (\bar{\zeta} - \zeta)\bar{C}C) - \bar{\chi}C(B - (\zeta - \bar{\zeta})C\bar{C})) \\ + \int d^4x \text{Tr}(-G_\mu D^\mu(Q)B & + (\bar{P}_\mu - \bar{\zeta}[G_\mu, \bar{C}])D^\mu(Q)C - (P_\mu - \zeta[G_\mu, C])D^\mu(Q)\bar{C}) \\ + \int d^4x \text{Tr}(-S\{C, \bar{C}\} & + \bar{R}C^2 + R\bar{C}^2 + \bar{U}[B + \bar{\zeta}\{C, \bar{C}\}, C] - U[B - \zeta\{C, \bar{C}\}, \bar{C}]) \\ + \sigma \int d^4x \text{Tr}(P_\mu \bar{P}^\mu - \frac{1}{2}m^2 G_\mu G^\mu) & + \rho \int d^4x \text{Tr}(R\bar{R} - S^2). \end{aligned}$$

Let us now state the (classical) symmetries of the Curci-Ferrari model given by (11). The superalgebra generated by $\{s_m, \bar{s}_m; d, \bar{d}, q\}$ is easily verified to be the algebra $\text{osp}(2,1)$ [5]. It is given by the following $\mathbf{q}$ -graded commutation relations ²:

$$s_m^2 = -m^2 d, \quad \{s_m, \bar{s}_m\} = m^2 q, \quad \bar{s}_m^2 = -m^2 \bar{d}, \quad (12)$$

$$[d, s_m] = 0, \quad [\bar{d}, s_m] = -\bar{s}_m, \quad [q, s_m] = s_m, \quad (13)$$

$$[\bar{d}, \bar{s}_m] = 0, \quad [d, \bar{s}_m] = s_m, \quad [q, \bar{s}_m] = -\bar{s}_m, \quad (14)$$

$$[q, d] = 2d, \quad [q, \bar{d}] = -2\bar{d}, \quad [d, \bar{d}] = -q. \quad (15)$$

²This algebraic structure is independent on defining the (anti-) BRST and (anti-) DJNO operations with or without the gauge parameter dependent extension by (5) and (10).

Now, using these relations it is easy to check that the action $S(g, m; \xi, \rho, \sigma)$ is invariant under (anti-) BRST as well as (anti-) DJNO transformations:

$$\begin{aligned} s_m S(g, m; \xi, \rho, \sigma) &= 0, & d S(g, m; \xi, \rho, \sigma) &= 0, \\ \bar{s}_m S(g, m; \xi, \rho, \sigma) &= 0, & \bar{d} S(g, m; \xi, \rho, \sigma) &= 0; \end{aligned} \quad (16)$$

invariance under q is trivial since $S(g, m; \xi, \rho, \sigma)$ has ghost number zero by definition. Furthermore, the (anti-) DJNO symmetry of the action according to (12) is a consequence of the (anti-) BRST invariance!

Some further remarks are in order. At first it is evident that in the limit $m^2 = 0$ the usual algebra of nilpotent (anti-) BRST operators $s^2 = 0$, $\{s, \bar{s}\} = 0$, $\bar{s}^2 = 0$ appears, and that the $sp(1)$ -algebra (15) of $d, \bar{d}, q$, being now an independent requirement, decouples. Of course, a massless Yang-Mills theory in nonlinear gauge obtains ³. Taking additionally the limit $\xi \rightarrow 0$ we finally come back to the Landau gauge.

Secondly, it should be emphasized that the mass is no independent parameter of the theory. More exactly there exists no symmetric invariant which is related to the mass since the Slavnov-Taylor identities explicitly depend on m .

Thirdly, in the action the auxiliary field B occurs quadratically (as in linear gauges), and through $(\zeta - \bar{\zeta})B(\bar{C}C)$ it couples to the (anti-) ghost fields, which also have quartic selfcoupling; therefore B does not decouple from the dynamics. And finally, as is seen from (11) most of the sources also occur bilinear (being coupled by σ and ρ). However, these complications may be surmounted as will be shown in the next two sections.

2 Ward identities for the vertex functional Γ

In order to construct the renormalized theory associated to the classical action (1) the symmetries of that action have to be expressed by corresponding Ward identities of the 1PI functional which hold, first of all, at lowest perturbative order, i.e. for $\Gamma^0(g, m; \xi, \rho, \sigma)$, and then these identities are to be required to hold for the renormalized 1PI functional $\Gamma(g, m; \xi, \rho, \sigma)$. Since these identities are purely algebraic requirements they may be formulated for a generic functional which also will be denoted by Γ :

$$\int d^4x \mathbf{W} \Gamma = 0 \quad (\text{rigid Ward identity}), \quad (17)$$

$$\mathbf{S}_m(\Gamma) = 0 \quad (\text{Slavnov - Taylor identity}), \quad (18)$$

$$\mathbf{D}(\Gamma) = 0 \quad (\text{Delduc - Sorella identity}), \quad (19)$$

$$\mathbf{T}_\zeta(\Gamma) = 0 \quad (\text{Bonora - Pasti - Tonin identity}), \quad (20)$$

and, in addition, the $\mathbf{C}$ -conjugated second and third identity; the first and the last identity (fixing the definition of $\mathcal{B}$) are $\mathbf{C}$ -even. The above operations are defined

³As has been shown by Curci and Ferrari [3] in the massless case two independent gauge parameters $\xi, \bar{\xi}$ may be introduced. But then (anti-) DJNO invariance cannot be maintained.

in the following manner:

$$\mathbf{W} = \sum [\Phi, \delta/\delta\Phi],$$

where the sum goes over all fields and sources, generically denoted by Φ , with the commutator or anticommutator depending on even or odd $\mathbf{q}$ -grading of Φ , respectively. The (bilinear) Slavnov-Taylor (ST) and the Delduc-Sorella (DS) [6] operations as usual are given by

$$\mathbf{S}_m(\Gamma) = \int d^4x \sum \text{Tr} \left((\mathbf{s}_m \Phi) \frac{\delta\Gamma}{\delta\Phi} \right) \quad \text{and} \quad \mathbf{D}(\Gamma) = \int d^4x \sum \text{Tr} \left((\mathbf{d}\Phi) \frac{\delta\Gamma}{\delta\Phi} \right),$$

respectively, where the nonlinear parts of $\mathbf{s}_m \Phi$ and $\mathbf{d}\Phi$ are to be replaced by the functional derivative of Γ with respect to the corresponding source $\bar{P}_\mu, \bar{R}, S$ or $\bar{U}$ (e.g. $\zeta [\mathcal{B} + \bar{\zeta} \{ \mathcal{C}, \bar{\mathcal{C}} \}, \mathcal{C}]$ by $\zeta \delta\Gamma/\delta\bar{U}$). The Bonora-Pasti-Tonin (BPT) operation [7] is defined by

$$\mathbf{T}_\zeta(\Gamma) = \frac{\partial\Gamma}{\partial\zeta} - \int d^4x \text{Tr} \left(\frac{\delta\Gamma}{\delta S} \frac{\delta\Gamma}{\delta B} \right).$$

For linear gauges the action S , if appropriately chosen, coincides with Γ^0 , but in the case of nonlinear gauges this will not be true. Namely, the ST operation applied to $S(g, m; \xi, \rho, \sigma)$ leads to the expression

$$\mathbf{S}_m(S) = 2\zeta\bar{\zeta}\rho \int d^4x \text{Tr} \left(\frac{\delta S}{\delta B} \frac{\delta S}{\delta \bar{C}} \right) \neq 0.$$

Therefore, S has to be changed in such a manner that this unwanted term is cancelled. The explicit determination is very cumbersome, it leads to the following result:

$$\Gamma^0(g, m; \xi, \rho, \sigma) = S(g, m; \xi, \rho, \sigma) + \frac{\zeta\bar{\zeta}\rho}{1 - 4\zeta\bar{\zeta}\rho\xi} \int d^4x \text{Tr} \left(\frac{\delta S}{\delta B} \frac{\delta S}{\delta B} \right). \quad (21)$$

As may be seen by an explicit computation this expression fulfills the identities (17) – (20) and therefore constitutes the correct starting point for perturbative renormalization.

Here, it should be remarked that $S(g, m; \xi, 0, 0)$ also fulfils the identities (17) – (20) and, in fact, this restricted action was the expression we started from. But, as is easily seen, it is not stable against small perturbations, i.e. not all the field monomials Δ introduced in such a way that $S + \epsilon\Delta$ fulfills the identities (17) – (20) up to order $O(\epsilon^2)$ can be obtained from S by a variation of the parameters of the restricted action. Therefore this action has to be extended by introducing the ρ - and σ -dependent combinations (2) of the fields and the associated sources.

3 Gauge and consistency conditions

The auxiliary field B , despite of the fact that it does not decouple from the dynamics, is not a primary one. However, its dynamics can be constrained to be the same as for the classical action by the following gauge condition (in fact, it is the renormalized equation of motion):

$$\tau \delta\Gamma/\delta B = \mathbf{B}_\xi(\Gamma) + \partial_\mu \mathbf{B}_\sigma^\mu + \mathbf{B}, \quad (22)$$

with $\tau = (1 - \rho\xi)$ and

$$\begin{aligned} \mathbf{B}_\xi(\Gamma) &= \xi(2B + (\zeta - \bar{\zeta}) \delta\Gamma/\delta S) + \chi(\bar{C} + \rho\bar{U}) - \bar{\chi}(C + \rho U), \\ \mathbf{B}_\sigma^\mu &= (1 - \sigma)G^\mu - Q^\mu, \\ \mathbf{B} &= [Q_\mu, G^\mu] + \{U, \bar{C}\} - \{\bar{U}, C\}. \end{aligned}$$

Here, we emphasise that, contrary to the common belief, (22) is a well posed condition since its nonlinear content is respected by the functional derivation of Γ with respect to the source S .

Applying now $\delta/\delta B$ on the ST identity an equation completely analogous to (22) obtains for $\delta\Gamma/\delta\bar{C}$ which is nothing else than the antighost equation of motion:

$$-\tau \delta\Gamma/\delta\bar{C} = \mathbf{C}_\xi(\Gamma) + \partial^\mu \mathbf{C}_\sigma^\mu(\Gamma) + \mathbf{C}(\Gamma), \quad (23)$$

with

$$\begin{aligned} \mathbf{C}_\xi(\Gamma) &= \xi\delta\Gamma/\delta\bar{U} - 2m^2\xi C - m^2\xi(C + \rho U) \\ &\quad - \bar{\chi}(\rho R - \delta\Gamma/\delta\bar{R}) + \chi(B + \bar{\zeta}\rho\delta\Gamma/\delta B - \rho S - \bar{\zeta}\delta\Gamma/\delta S), \\ \mathbf{C}_\sigma^\mu(\Gamma) &= (1 - \sigma)P^\mu - \delta\Gamma/\delta\bar{P}_\mu, \\ \mathbf{C}(\Gamma) &= [Q_\mu, P^\mu] - [\bar{C}, R] + [C, S] - [U, B] - [G_\mu, \delta\Gamma/\delta\bar{P}_\mu] \\ &\quad - [\bar{U}, \delta\Gamma/\delta\bar{R}] - \zeta[U, \delta\Gamma/\delta S] + \zeta[C, \delta\Gamma/\delta B]; \end{aligned}$$

an analogous condition for the ghost equation of motion obtains by C-conjugation (the explicit expressions will not be given here).

If now $\delta/\delta\bar{C}$ is applied on the anti-ST identity we finally get a further independent condition which we called local Lee identity:

$$\mathbf{L}_\sigma \Gamma = m^2 \partial^\mu Q_\mu \quad \text{with} \quad \mathbf{L}_\sigma \equiv \partial_\mu \left((1 - \sigma) \frac{\delta}{\delta Q_\mu} + \frac{\delta}{\delta G_\mu} \right) + \mathbf{W}. \quad (24)$$

Of course, by integration the rigid Ward identity (17) – (20) obtains. Surprisingly, this identity is independent of ξ , and also of ζ and ρ . It expresses almost the (original) local gauge symmetry of the classical action to hold for Γ (up to the mass term in (24)).

Together with the identities (17) – (20) the conditions (22) – (24) constitute a set of algebraic relations which express the symmetries to be fulfilled by the vertex functional. Additional consistency conditions do not appear since the (linearized) ST- and DS-operations fulfil (anti-) commutation relations analogous to (12) – (15); the BPT-identity is not independent, but follows by the help of $\partial\mathbf{S}_m(\Gamma)/\partial\zeta = 0$.

4 Concluding remarks

In fact, the algebraic relations for Γ , introduced in the preceding Sections, define a renormalizable theory as has been shown, together with a detailed presentation of the results given here, in a separate paper [8]. This proof is given for the nonextended case using only the relations (3), (4) and (8), (9), as well as in the extended case where also the (anti-) BRST and (anti-) DJNO transformations (5) and (10) of the gauge parameter ξ are taken into account. In the latter case it is possible to study, according to the procedure introduced by Piguet and Sibold [9], the gauge (in)dependence of the theory. Especially, we were able to show that the extended (anti-) ST and (anti-) DJ identities are free of anomalies. This overcomes a drawback of the consideration by Bauer, Ottitsch, Piguet and Schweda [10], who did not take into account the anti-BRST and anti-DJNO transformations and therefore were not able to show complete freedom of anomalies: an (unphysical) anomaly of the DS identity remained.

A further result concerns the number of independent z -factors of the renormalized theory. If no gauge condition is required and therefore also the related consistency conditions are absent, 7 z -factors related to the independent parts of the classical action (1) appear, namely

$$z_Q, \quad z_B, \quad z_G; \quad z_g, \quad z_\xi, \quad z_\rho \quad \text{and} \quad z_\sigma,$$

the others being related according to

$$\begin{aligned} z_m^2 &= z_Q^{-1} z_G^{-1}, & z_C &= z_{\bar{C}} = z_m^{-1} z_B, & z_S &= z_R = z_{\bar{R}} = z_B^{-1}, \\ z_P &= z_{\bar{P}} = z_m^{-1} z_Q^{-1}, & z_U &= z_{\bar{U}} = z_m^{-1} z_B^{-1}. \end{aligned}$$

If however the gauge condition (22) together with the consistency conditions (23) and (24) is imposed 4 additional relations obtain

$$z_G = 1, \quad z_\xi = z_Q z_B^{-1}, \quad z_\sigma = z_Q + (1 - z_Q)/\sigma, \quad z_\rho = z_B^2 + z_Q^{-1} z_B (1 - z_Q z_B)/\rho\sigma,$$

and therefore only the following three z -factors are independent:

$$z_Q, \quad z_B; \quad z_g$$

(in complete analogy to the case of linear gauges).

This seemingly contradicts a recent result of de Boer, Skenderis, van Nieuwenhuizen and Waldron [11] which, disregarding anti-BRST symmetry and requiring no gauge condition, states that for the Curci-Ferrari model 5 independent z -factors appear. This, however, can be explained also within our approach: having no anti-BRST symmetry it is only necessary to introduce the external sources $\bar{P}_\mu$ and $\bar{R}$. But then no mixing of the fields according to (2) appears and therefore the parameters ρ and σ are absent reducing the number of independent parameters from 7 to 5 only.

Finally, let us remark that this method of enlarging the symmetries can be extended to gauge theories with generic background configurations, too; also the case of special (e.g. instanton) backgrounds may be considered. Thereby a result of Rouet [12] has been improved concerning the dependence of the Greens functions on the background A_μ (these results will be given elsewhere).

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