TWO-DIMENSIONAL INTERACTIONS IN A CLASS OF TENSOR GAUGE FIELDS FROM LOCAL BRST COHOMOLOGY

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Lagrangian interactions in a class of two-dimensional tensor gauge field theory are derived by means of deforming the solution to the master equation with specific cohomological techniques. Both the gauge transformations and their algebra are deformed. The gauge algebra of the coupled model is open.

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1. Introduction

The key point in the development of the BRST symmetry is represented by its reconstruction on cohomological grounds [1]. This cohomological approach added to the BRST sets extremely powerful capabilities, such as the investigation of perturbative renormalization [2, 3], the anomaly-tracking mechanism [3,4], the simultaneous study of local and rigid invariances of a given theory [5], as well as the reformulation of the construction of consistent interactions in gauge theories [6] in terms of the deformation theory [7], or, actually, in terms of the deformation of the solution to the master equation. A wide class of models has been studied within the background of the deformation of the master equation [8].

In this paper we solve the problem of constructing all consistent Lagrangian interactions in a special class of two-dimensional tensor gauge field theories from the deformation of the solution to the master equation. We

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begin with a free two-dimensional theory that involves a vector field and a two-tensor field with no special symmetry. The gauge transformations of the free action are Abelian and irreducible, such that we deal with a Lagrangian model of Cauchy order two. Consequently, the free antifield-BRST symmetry decomposes into a sum between the Koszul–Tate differential and the exterior derivative along the gauge orbits only, $s = \delta + \gamma$. Next, we deform the solution to the master equation of the free model. The first-order deformation belongs to the zeroth order cohomological space of the BRST differential modulo the exterior spacetime derivative, $H^0(s|d)$. Its computation proceeds by expanding the cocycles according to the antighost number, and by further using the cohomological spaces $H(\gamma)$ and $H_2(\delta|d)$. We completely determine the first-order deformation, which stops at antighost number two and contains some antisymmetric functions that involve only the undifferentiated vector field. Its consistency requires that the above mentioned antisymmetric functions fulfill a certain identity, which is shown to possess solutions. Under these circumstances, we can set all the deformations of order higher than one to be equal to zero. From the deformed solution to the master equation that is consistent to all orders in the coupling constant we extract the Lagrangian action of the interacting model, together with its gauge transformations and accompanying gauge algebra. We find that the gauge transformations are modified with respect to the initial ones, while their algebra is open.

2. The BRST symmetry of the free model

We begin with the Lagrangian action in two spacetime dimensions

$$S_0\left[A_{\mu\nu}, B_\mu\right] = \int d^2x \varepsilon^{\mu\nu} A_{\mu\lambda} \partial_\nu B^\lambda, \tag{2.1}$$

where the fields are bosonic, and $A_{\mu\nu}$ is a two-tensor field with no specific symmetry. The notation $\varepsilon^{\mu\nu}$ signifies the two-dimensional antisymmetric symbol, with $\varepsilon^{01} = +1$. We work with the flat Minkowskian metric $g_{\mu\nu}$ of signature $(1, -1)$. Action (2.1) is invariant under the gauge transformations

$$\delta_\epsilon A_{\mu\nu} = \partial_\mu \epsilon_\nu, \ \delta_\epsilon B_\mu = 0\,, \tag{2.2}$$

with ϵ_ν the bosonic vector gauge parameter.

In order to construct the BRST symmetry of this free model, we introduce the field/ghost, respectively, antifield spectra

$$\Phi^\Delta = \left(A_{\mu\nu}, B_\mu, \eta_\mu\right), \ \Phi^*_\Delta = \left(A^{*\mu\nu}, B^{*\mu}, \eta^{*\mu}\right), \tag{2.3}$$

where η_μ represents the fermionic ghost associated with the vector gauge parameter ϵ_μ. The above gauge transformations are Abelian and irreducible.

Consequently, the BRST differential s reduces to the sum between the Koszul–Tate differential δ and the exterior longitudinal derivative γ only

$$s = \delta + \gamma \,, \tag{2.4}$$

that are, respectively, graded in terms of the antighost number (antigh) and the pure ghost number (pgh). While the Koszul–Tate differential ($\delta^2 = 0$, $\mathrm{antigh}\,(\delta) = -1$, $\mathrm{pgh}\,(\delta) = 0$) realizes a resolution of smooth functions defined on the stationary surface of field equations, the exterior longitudinal derivative ($\mathrm{pgh}\,(\gamma) = 1$, $\mathrm{antigh}\,(\gamma) = 0$) anticommutes with δ and turns out to be a true differential in the particular case of the model under study $\left(\gamma^2 = 0\right)$. Its cohomological space at pure ghost number zero computed in the homology of δ, $H^0\left(\gamma | H_*\left(\delta\right)\right)$, is given by the algebra of Lagrangian physical observables, and is in the meantime isomorphic to the zeroth order cohomological space of s, $H^0\left(s\right)$, that contains the so-called BRST observables. The degrees (antigh) and (pgh) of the BRST generators (2.3) are given by

$$\begin{aligned} \mathrm{pgh}\,(A_{\mu\nu}) &= 0\,, \qquad \mathrm{pgh}\,(B_\mu) = 0\,, \\ \mathrm{pgh}\,(\eta_\mu) &= 1\,, \qquad \mathrm{pgh}\,(\Phi^*_\Delta) = 0\,, && (2.5) \\ \mathrm{antigh}\left(\Phi^\Delta\right) &= 0\,, \qquad \mathrm{antigh}\,(A^{*\mu\nu}) = 1\,, && (2.6) \\ \mathrm{antigh}\,(B^{*\mu}) &= 1\,, \qquad \mathrm{antigh}\,(\eta^{*\mu}) = 2\,, && (2.7) \end{aligned}$$

while the actions of δ and γ read as

$$\delta A_{\mu\nu} = 0\,, \quad \delta B_\mu = 0\,, \quad \delta\eta_\mu = 0\,, \tag{2.8}$$

$$\begin{aligned} \delta A^{*\mu\nu} &= -\varepsilon^{\mu\lambda}\partial_\lambda B^\nu, \\ \delta B^{*\mu} &= \varepsilon_{\alpha\beta}\partial^\beta A^{\alpha\mu}, \\ \delta\eta^{*\mu} &= -\partial_\nu A^{*\nu\lambda}, \end{aligned} \tag{2.9}$$

$$\begin{aligned} \gamma A_{\mu\nu} &= \partial_\mu\eta_\nu\,, \quad \gamma B_\mu = 0, \quad \gamma\eta_\mu = 0\,, && (2.10) \\ \gamma A^{*\mu\nu} &= 0\,, \qquad \gamma B^{*\mu} = 0\,, \quad \gamma\eta^{*\mu} = 0\,. && (2.11) \end{aligned}$$

The overall degree of the BRST complex is named ghost number (gh) and is defined like the difference between the pure ghost number and the antighost number, such that $\mathrm{gh}\,(s) = 1$. The BRST symmetry of the free model is canonically generated in the structure of antibracket $(,)$ via a generator $\overset{(0)}{S}$, namely, $s\cdot = \left(\cdot, \overset{(0)}{S}\right)$. This structure is defined by decreeing the fields/ghosts conjugated in the antibracket with the corresponding antifields. The canonical generator is bosonic, its ghost number is equal to zero, and represents the solution to the master equation $\left(\overset{(0)}{S},\overset{(0)}{S}\right)=0$. Its expression is found by means of developing $\overset{(0)}{S}$ according to the antighost number. In the case of our free model, it takes the simple form

$$\overset{(0)}{S}=S_0+\int d^2xA^{*\mu\nu}\partial_\mu\eta_\nu\,, \tag{2.12}$$

and contains only components of antighost number equal to zero and one.

3. The deformation procedure

A consistent deformation of the free action (2.1) and of its gauge invariances (2.2) defines a deformation of the corresponding solution to the master equation that preserves both the master equation and the field/antifield spectra. So, if $S_0\left[A_{\mu\nu},B_\mu\right]+g\int d^2xa_0+O\left(g^2\right)$ stands for a consistent deformation of the free action, with deformed gauge transformations $\bar{\delta}_\epsilon A_{\mu\nu}=\partial_\mu\epsilon_\nu+g\beta_{\mu\nu}+O\left(g^2\right)$, $\bar{\delta}_\epsilon B_\mu=g\beta_\mu+O\left(g^2\right)$, then the deformed solution to the master equation

$$S=\overset{(0)}{S}+g\int d^2xa+O\left(g^2\right) \tag{3.1}$$

satisfies $(S,S)=0$, where $a=a_0+A^{*\mu\nu}\bar{\beta}_{\mu\nu}+B^{*\mu}\bar{\beta}_\mu+$ 'more' (g is the so-called deformation parameter or coupling constant). The terms $\bar{\beta}_{\mu\nu}$, $\bar{\beta}_\mu$ are obtained by replacing the gauge parameters ϵ_μ with the fermionic ghosts η_μ in the functions $\beta_{\mu\nu}$ and β_μ.

The master equation $(S,S)=0$ holds to order g if and only if

$$sa=\partial_\mu j^\mu, \tag{3.2}$$

for some local j^μ. This means that the nontrivial first-order consistent interactions belong to $H^0\left(s|d\right)$, where d is the exterior space-time derivative. In the case where a is a coboundary modulo d ($a=s\lambda+\partial_\mu b^\mu$), then the deformation is trivial (it can be eliminated by a redefinition of the fields). In order to investigate the solution to (3.2), we develop a according to the antighost number

$$a=a_0+a_1+\ldots+a_J\,,\quad \mathrm{antigh}\left(a_k\right)=k\,, \tag{3.3}$$

where the last term can be assumed to be annihilated by γ, $\gamma a_J=0$. Thus, we need to know the cohomology of γ, $H\left(\gamma\right)$, in order to determine the terms of the highest antighost number in a. From (2.10)–(2.11) it is simple to see that the cohomology of γ is generated by $\partial_{[\mu} A_{\nu]\lambda}$, B_μ, $A^{*\mu\nu}$, $B^{*\mu}$, $\eta^{*\mu}$ together with their spacetime derivatives, as well as by the undifferentiated ghosts η^μ. (The derivatives of the ghosts are discarded from $H(\gamma)$ since they are γ-exact, as results from the first formula in (2.10).) If we denote by $e^M(\eta^\mu)$ a basis in the space of the polynomials in the ghosts, it follows that the general solution to the equation $\gamma a_J = 0$ takes the form

$$a_J = \alpha_J \left(\left[\partial_{[\mu} A_{\nu]\lambda}\right], [B_\mu], [A^{*\mu\nu}], [B^{*\mu}], [\eta^{*\mu}] \right) e^J(\eta^\mu), \tag{3.4}$$

where $\mathrm{antigh}(\alpha_J) = J$ and $\mathrm{pgh}\left(e^J(\eta^\mu)\right) = J$. In (3.4) the notation $f([q])$ signifies that f depends on q and its derivatives up to a finite order.

By projecting the equation (3.2) on antighost number $(J-1)$, we obtain

$$\delta a_J + \gamma a_{J-1} = \partial_\mu m^\mu. \tag{3.5}$$

Inserting (3.4) into (3.5), it follows that the existence of a_{J-1} requires that α_J pertains to $H_J(\delta|d)$, *i.e.*,

$$\delta \omega_J = \partial_\mu n^\mu. \tag{3.6}$$

On the other hand, since the free model under consideration is of Cauchy order two, from [9] it follows that $H_k(\delta|d) = 0$ for $k > 2$, so we can assume that the development (3.3) stops at antighost number two

$$a = a_0 + a_1 + a_2, \tag{3.7}$$

where a_2 is of the form (3.4) with $J = 2$, and α_2 belongs to $H_2(\delta|d)$. On the one hand, the most general representative of $H_2(\delta|d)$ is

$$\alpha_2 = \frac{\delta W}{\delta B^\lambda}\eta^{*\lambda} + \frac{1}{2}\varepsilon_{\mu\nu}\frac{\delta^2 W}{\delta B^\lambda \delta B^\rho} A^{*\mu\lambda} A^{*\nu\rho}, \tag{3.8}$$

where $W\left(B^\lambda\right)$ is an arbitrary function that involves only the undifferentiated vector fields. On the other hand, the elements of pure ghost number two of a basis in the ghosts are of the type

$$\eta^\alpha \eta^\beta. \tag{3.9}$$

Combining (3.8) with (3.9), we find that the last component in (3.7) becomes

$$a_2 = \frac{1}{2}\left(\frac{\delta W_{\alpha\beta}}{\delta B^\lambda}\eta^{*\lambda} + \frac{1}{2}\varepsilon_{\mu\nu}\frac{\delta^2 W_{\alpha\beta}}{\delta B^\lambda \delta B^\rho} A^{*\mu\lambda} A^{*\nu\rho} \right) \eta^\alpha \eta^\beta, \tag{3.10}$$

where the functions $W_{\alpha\beta}$ are antisymmetric, $W_{\alpha\beta} = -W_{\beta\alpha}$, due to the anticommutation of the ghosts. Then, the equation (3.5) for $J = 2$ takes the form $\delta a_2 + \gamma a_1 = \partial_\mu m^\mu$. Taking into account (3.10), we determine the piece of antighost number one from the first-order deformation like

$$a_1 = -\frac{\delta W_{\alpha\beta}}{\delta B^\lambda} A^{*\mu\lambda} A_\mu{}^\beta \eta^\alpha - W_{\alpha\beta} B^{*\alpha} \eta^\beta . \tag{3.11}$$

Further, (3.2) projected on antighost number zero produces the equation $\delta a_1 + \gamma a_0 = \partial_\mu k^\mu$, whose solution reads as

$$a_0 = -\frac{1}{2} \varepsilon^{\mu\nu} W_{\alpha\beta} A_\mu{}^\beta A_\nu{}^\alpha . \tag{3.12}$$

Putting together the relations (3.10–3.12), we can write down the first-order deformation of the solution to the master equation corresponding to the free model under study as

$$\begin{aligned} S_1 \;=\; & \int d^2x \left(-\frac{1}{2} \varepsilon^{\mu\nu} W_{\alpha\beta} A_\mu{}^\beta A_\nu{}^\alpha - \frac{\delta W_{\alpha\beta}}{\delta B^\lambda} A^{*\mu\lambda} A_\mu{}^\beta \eta^\alpha - W_{\alpha\beta} B^{*\alpha} \eta^\beta \right. \\ & \left. + \frac{1}{2} \left(\frac{\delta W_{\alpha\beta}}{\delta B^\lambda} \eta^{*\lambda} + \frac{1}{2} \varepsilon_{\mu\nu} \frac{\delta^2 W_{\alpha\beta}}{\delta B^\lambda \delta B^\rho} A^{*\mu\lambda} A^{*\nu\rho} \right) \eta^\alpha \eta^\beta \right) . \end{aligned} \tag{3.13}$$

Next, we investigate the higher-order deformations. If we make the notations $S_2 = \int d^2x b$ and $(S_1, S_1) = \int d^2x \Delta$, the second-order deformation is subject to the equation (written in local form)

$$\Delta = -2sb + \partial_\mu \theta^\mu . \tag{3.14}$$

With the help of (3.13), by direct computation we arrive at

$$\begin{aligned} (S_1, S_1) \;=\; & \int d^2x \left(m_{\beta\lambda\rho} \left(\varepsilon^{\mu\nu} A_\mu{}^\beta A_\nu{}^\rho \eta^\lambda + B^{*\lambda} \eta^\rho \eta^\beta \right) \right. \\ & - \frac{\delta m_{\beta\lambda\rho}}{\delta B^\nu} \left(A^{*\mu\nu} A_\mu{}^\beta \eta^\lambda \eta^\rho + \frac{1}{3} \eta^{*\nu} \eta^\beta \eta^\lambda \eta^\rho \right) \\ & \left. - \frac{1}{6} \varepsilon_{\mu\sigma} \frac{\delta^2 m_{\beta\lambda\rho}}{\delta B^\nu \delta B^\tau} A^{*\mu\nu} A^{*\sigma\tau} \eta^\beta \eta^\lambda \eta^\rho \right) , \end{aligned} \tag{3.15}$$

where we employed the notation

$$m_{\beta\lambda\rho} = W_{\alpha[\beta} \frac{\delta W_{\lambda\rho]}}{\delta B_\alpha} . \tag{3.16}$$

Since none of the terms in (3.15) can be written like in the right hand-side of (3.14), the consistency of the first-order deformation requires that

$$m_{\beta\lambda\rho} = 0 . \tag{3.17}$$

The solution to (3.17) reads as

$$W_{\alpha\beta}\left(B^{\mu}\right)=c_{1}\varepsilon_{\alpha\beta}W\left(B^{\mu}\right)+c_{2}\left(\varepsilon_{\alpha\lambda}B^{\lambda}B_{\beta}-\varepsilon_{\beta\lambda}B^{\lambda}B_{\alpha}\right), \tag{3.18}$$

where $W\left(B^{\mu}\right)$ is a scalar function depending on the undifferentiated vector field B^{μ}, while c_1 and c_2 are some constants. Then, the solution to (3.14) can be taken of the form $b=0$, which further yields $S_2=0$. Accordingly, we find that all the equations that stipulate the higher-order deformations of the solution to the master equation are satisfied for $S_k=0$, $k>2$.

4. Identification of the interacting theory

So far, we have constructed the complete deformed solution to the master equation for the model under discussion, which is consistent to all orders in the coupling constant, of the type

$$\begin{aligned}S\ =\ &\int d^{2}x\left(\varepsilon^{\mu\nu}A_{\mu\lambda}\left(\partial_{\nu}B^{\lambda}+\frac{1}{2}gW^{\lambda\rho}A_{\nu\rho}\right)\right.\\&+A^{*\mu\nu}\left(\partial_{\mu}\eta_{\nu}-g\frac{\delta W_{\alpha\beta}}{\delta B^{\nu}}A_{\mu}{}^{\beta}\eta^{\alpha}\right)-gB^{*\alpha}W_{\alpha\beta}\eta^{\beta}\\&\left.+\frac{1}{2}g\left(\frac{\delta W_{\alpha\beta}}{\delta B^{\lambda}}\eta^{*\lambda}+\frac{1}{2}\varepsilon_{\mu\nu}\frac{\delta^{2}W_{\alpha\beta}}{\delta B^{\lambda}\delta B^{\rho}}A^{*\mu\lambda}A^{*\nu\rho}\right)\eta^{\alpha}\eta^{\beta}\right),\end{aligned} \tag{4.1}$$

where $W_{\alpha\beta}\left(B^{\mu}\right)$ is given by (3.18). With the help of (4.1), we identify the resulting interacting theory and its gauge structure. Thus, the component of antighost number zero is nothing but the Lagrangian action of the coupled theory

$$\bar{S}_{0}\left[A_{\mu\nu},B_{\mu}\right]=\int d^{2}x\left(\varepsilon^{\mu\nu}A_{\mu\lambda}\left(\partial_{\nu}B^{\lambda}+\frac{1}{2}gW^{\lambda\rho}A_{\nu\rho}\right)\right). \tag{4.2}$$

The terms linear in the antifields of the original fields provide the gauge transformations of the action (4.2) like

$$\bar{\delta}_{\epsilon}A_{\mu\nu}=\partial_{\mu}\epsilon_{\nu}-g\frac{\delta W_{\alpha\beta}}{\delta B^{\nu}}A_{\mu}{}^{\beta}\epsilon^{\alpha},\quad \bar{\delta}_{\epsilon}B_{\mu}=-gW_{\mu\alpha}\epsilon^{\alpha}, \tag{4.3}$$

such that the gauge generators in De Witt condensed notations are expressed by

$$Z_{\mu\nu\alpha}^{(A)}=g_{\nu\alpha}\partial_{\mu}-g\frac{\delta W_{\alpha\beta}}{\delta B^{\nu}}A_{\mu}{}^{\beta},\quad Z_{\mu\alpha}^{(B)}=-gW_{\mu\alpha}. \tag{4.4}$$

Regarding the antighost number two pieces, they are of two kinds: ones are linear in the antifields of the ghosts, while the others are quadratic in the antifields of the fields. This means that the deformed gauge algebra of the interacting theory is non-Abelian, and, more precisely, open. The commutators among the gauge transformations (4.3) take the concrete form (again in De Witt notations)

$$\begin{aligned} Z^{(A)}_{\rho\lambda\alpha}\frac{\delta Z^{(A)}_{\mu\nu\beta}}{\delta A_{\rho\lambda}} & -Z^{(A)}_{\rho\lambda\beta}\frac{\delta Z^{(A)}_{\mu\nu\alpha}}{\delta A_{\rho\lambda}}+Z^{(B)}_{\rho\alpha}\frac{\delta Z^{(A)}_{\mu\nu\beta}}{\delta B_{\rho}} \\ &= g\frac{\delta W_{\alpha\beta}}{\delta B_{\gamma}}Z^{(A)}_{\mu\nu\gamma}+g\varepsilon_{\mu\rho}\frac{\delta^2 W_{\alpha\beta}}{\delta B^{\nu}\delta B^{\lambda}}\frac{\delta\bar{S}_0}{\delta A_{\rho\lambda}}\,, \end{aligned} \tag{4.5}$$

$$Z^{(B)}_{\rho\alpha}\frac{\delta Z^{(B)}_{\mu\beta}}{\delta B_{\rho}}-Z^{(B)}_{\rho\beta}\frac{\delta Z^{(B)}_{\mu\alpha}}{\delta B_{\rho}}=g\frac{\delta W_{\alpha\beta}}{\delta B_{\gamma}}Z^{(B)}_{\mu\gamma}. \tag{4.6}$$

Obviously, the gauge transformations (4.3) are irreducible, just like those of the free model.

5. Conclusion

To conclude, in this paper we have shown that there exist consistent and nontrivial Lagrangian interactions which can be introduced in a special class of two-dimensional tensor gauge field theories. Starting with the BRST differential of the free theory, we fully compute the first-order deformation, which contains some antisymmetric functions of the undifferentiated vector field. Next, we investigate its consistency, and evaluate the appearance of higher-order deformations. It turns out that the consistency of the deformation procedure restricts the antisymmetric functions to fulfill a certain identity, in which case all the higher-order deformations can be taken to vanish. As a result, we derive a coupled model with modified gauge transformations, endowed with an open gauge algebra.

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