Remarks on the Curci-Ferrari model

Peter M. Lavrov

Tomsk State Pedagogical University,
Kievskaya St. 60, 634061 Tomsk, Russia

Abstract

Dependence of Green’s functions for the Curci-Ferrari model on the parameter resembling the gauge parameter in massless Yang-Mills theories is investigated. It is shown that the generating functional of vertex functions (effective action) depends on this parameter on-shell.

Keywords: Gauge theories, BRST symmetry, massive Yang-Mills fields

PACS: 04.60.Gw, 11.30.Pb

1E-mail: lavrov@tspu.edu.ru
1 Introduction

Recently, it was claimed [1] that the Curci-Ferrari (CF) model [2] can be presented as a unitary and renormalizable model for massive Yang-Mills fields without Higgs fields. From the beginning it was well known that the CF model obeys the property of renormalizability [3, 4] and the action of this model is invariant under modified BRST and anti-BRST transformations. In massless limit, the action of the CF model reduces to the Faddeev-Popov (FP) action [5] constructed in a one parameter linear gauge. The FP action is invariant under the BRST transformations [6, 7] as well as under the anti-BRST transformations in special gauges [8, 9]. The BRST symmetry plays a fundamental role in quantum theory of gauge fields [10]. Note, for example, that breaking of BRST symmetry as it occurs in Yang-Mills theories when one takes into account the Gribov horizon [11, 12, 13] leads to the gauge dependence of effective action in gauge theories on-shell [14, 15]. In turn it means inconsistency for physical interpretation of results obtained within this approach. In Yang-Mills theories both the BRST and anti-BRST transformations are nilpotent. Nilpotency of the BRST symmetry allows to formulate suitable conditions (the so-called Kugo-Ojima criterion) for a physical state space providing unitarity of S-matrix in non-abelian gauge theories [16]. In contrast to this case, the modified BRST and anti-BRST transformations in the CF model are not nilpotent. Namely, this fact was considered as a reason for violation of unitarity in this theory for a long time [3, 17, 4]. Reformulation of the CF model proposed in [1] is connected with using local non-linear transformations of massive vector fields and rewritten the CF action in terms of new variables to obtain a model for massive Yang-Mills fields without Higgs fields. The statement about unitarity of S-matrix for this model contradicts with previous conclusions about non-unitarity of the CF model [3, 17, 4] and sounds rather strange from the point of view of the equivalence theorem [18, 19] because two theories under consideration are connected through a change of variables which satisfies conditions of the theorem. Now it is clear that the unitarity problem for the CF model [2] and the model of massive Yang-Mills fields without Higgs fields [1] needs in further investigations.

In present paper the dependence of Green’s functions for the CF model on a parameter resembling the gauge parameter in massless Yang-Mills theories is investigated. It is shown that the generating functional of vertex functions (effective action) does depend on this parameter even on-shell.

The paper is organized as follows. In Section 2, the CF model is considered. In Section 3, dependence of Green’s functions for the CF model on the parameter \( \beta \) is studied. Finally, Section 4 gives concluding remarks.

We employ the condensed notation of DeWitt [20]. Derivatives with respect to sources are taken from the left, while those with respect to fields are taken from the right. Left derivatives with respect to fields are labeled by a subscript \( l \).
2 The Curci-Ferrari model

Consider a massive extension of the massless Yang-Mills theory proposed by Curci and Ferrari [2]. The CF model is described by the action

\[ S = S_{YM} + S_{gf} + S_m, \]  

(2.1)

where \( S_{YM} \) is the Yang-Mills action of fields \( A^a_\mu \), which take values in the adjoint representation of the Lie algebra \( su(N) \) so that, \( a = 1, \ldots, N^2 - 1, \)

\[ S_{YM} = -\frac{1}{4} F^{a}_{\mu\nu} F^{a}_{\mu\nu} \quad \text{with} \quad F^{a}_{\mu\nu} = \partial_\mu A^a_\nu - \partial_\nu A^a_\mu + f^{abc} A^b_\mu A^c_\nu, \]  

(2.2)

and \( \mu, \nu = 0, 1, \ldots, D-1, \) the Minkowski space has signature \((+, -, \ldots, -)\), and \( f^{abc} \) denote the (totally antisymmetric) structure constants of \( su(N) \), the symbol \( \int d^D x \) is omitted. The action \( S_{gf} \) has the form

\[ S_{gf} = B^a \partial^\mu A^a_\mu + \bar{C}^a \partial^\mu D^a_{\mu b} C^b + \frac{\beta}{4} B^a B^a + \frac{\beta}{4} \tilde{B}^a \tilde{B}^a, \]  

(2.3)

where

\[ \tilde{B}^a = -B^a + N^a, \quad N^a = N^a(C, \bar{C}) = f^{abc} \bar{C}^b C^a, \quad D^a_{\mu b} = \delta^{ab} \partial_\mu + f^{acb} A^c_\mu \]  

(2.4)

and \( \beta \) is a parameter of the model. The action \( S_m \) contains a mass \( m \) for the vector fields \( A^a_\mu \) and the ghosts \( C^a \) and antighosts \( \bar{C}^a \)

\[ S_m = \frac{1}{2} m^2 A^a_\mu A^a_\mu + m^2 \bar{C}^a C^a. \]  

(2.5)

Here the notations \( B^a \) for bosonic auxiliary fields were used. In massless limit they are identified with the Nakanishi - Lautrup fields.

Note that \( S_{YM} + S_{gf} \) can be presented as the action constructed by the rules of Faddeev-Popov quantization [5], \( S_{FP} \), in one-parameter linear gauge \( \chi^a \)

\[ \chi^a = \partial^\mu A^a_\mu + \frac{\beta}{2} B^a \]  

(2.6)

and modified by the additional term \( S_{ad} \)

\[ S_{YM} + S_{gf} = S_{FP} + S_{ad}, \]  

(2.7)

where

\[ S_{ad} = \frac{\beta}{4} N^a N^a - \frac{\beta}{2} B^a N^a. \]  

(2.8)

The action (2.2) is invariant under the gauge transformations

\[ \delta A^a_\mu = D^{ab}_{\mu} \xi^b, \]  

(2.9)
where $\xi^a = \xi^a(x)$ are arbitrary functions of space-time coordinates. In turn, the actions $S_{FP}$ and $S_{ad}$ are invariant under BRST transformation [6, 7]

$$
\begin{align*}
\delta_B A^a_\mu &= D^a_{\mu b} C^b \theta, \\
\delta_B C^a &= \frac{1}{2} f^{abc} C^b C^c \theta, \\
\delta_B \bar{C}^a &= B^a \theta, \\
\delta_B B^a &= 0,
\end{align*}
$$

(2.10)

where $\theta$ is a constant Grassmann parameter. Moreover, these actions are invariant under the anti-BRST transformation [8, 9]

$$
\begin{align*}
\bar{\delta}_B A^a_\mu &= D^a_{\mu b} \bar{C}^b \bar{\theta}, \\
\bar{\delta}_B C^a &= \left(-B^a + f^{abc} \bar{C}^b C^c\right) \bar{\theta}, \\
\bar{\delta}_B \bar{C}^a &= \frac{1}{2} f^{abc} \bar{C}^b \bar{C}^c \theta, \\
\bar{\delta}_B B^a &= - f^{abc} \bar{C}^b B^c \theta,
\end{align*}
$$

(2.11)

with $\bar{\theta}$ being a constant Grassmann parameter [2]. The action of the CF model is not invariant under the BRST transformation because of $\delta_B S_m \neq 0$ but it is invariant under the modified BRST transformation $\delta_{mB} S = 0$ [2], where

$$
\begin{align*}
\delta_{mB} A^a_\mu &= D_{\mu}^{ab} C^b \theta, \\
\delta_{mB} C^a &= \frac{1}{2} f^{abc} C^b C^c \theta, \\
\delta_{mB} \bar{C}^a &= B^a \theta, \\
\delta_{mB} B^a &= m^2 C^a \theta,
\end{align*}
$$

(2.12)

as well as under the modified anti-BRST transformation $\bar{\delta}_{mB} S = 0$ [2]. In what follows the explicit form of the modified anti-BRST transformation will not be essential, and we omit it.

Note only that existence of anti-BRST symmetry for Yang-Mills theories in the gauge (2.6) is not specific property of these theories in special gauges. For any classical gauge theory in any admissible gauge one can construct a quantum version respecting both the BRST and anti-BRST symmetries [21, 22, 23, 24]. In contrast to the usual BRST (or anti-BRST) transformation, the modified BRST (or modified anti-BRST) transformation is not nilpotent. It was a reason to claim violation of unitarity for the CF model [3, 17, 4].

Returning to the CF model it needs definitely to say that from the beginning it should be considered as a non-gauge model in contrast to the Faddeev-Popov action $S_{FP}$ constructed for Yang-Mills action $S_{YM}$ which is invariant under gauge transformations $\delta A^a_\mu = D^a_{\mu b} \xi^b$. In sector of vector fields $A^a_\mu$ the action of CF model, $S_{mYM}$,

$$
S_{mYM} = -\frac{1}{4} F^a_{\mu\nu} F^{\mu\nu a} + \frac{1}{2} m^2 A^a_\mu A^a_\mu
$$
is not gauge invariant at all. In particular, there is no reason to refer $\beta$ as the gauge parameter. It is a parameter of the theory with initial classical non-degenerated action $S$ (2.1), for which the physical space contains particles corresponding to the massive vector fields $A^a_{\mu}$ and scalar anticommuting fields $C^a, \bar{C}^a$. This point of view will be supported in the next Section by investigation of the dependence of Green’s functions on this parameter.

3 Dependence of Green’s functions on parameter $\beta$

In this section we will study dependence of Green’s functions for the CF model (2.1)-(2.5) on parameter $\beta$. We start with the vacuum functional $Z_\beta$ for the CF model explicitly indicating dependence on $\beta$

$$Z_\beta = \int D\phi \exp \left( \frac{i}{\hbar} S \right),$$  \hspace{1cm} (3.1)

where $\phi$ denotes the set of all fields of the theory under consideration,

$$\phi^i = (A^a_{\mu}, \bar{C}^a, C^a, B^a).$$  \hspace{1cm} (3.2)

Let $Z_{\beta+\delta\beta}$ be the vacuum functional corresponding to small variation of the parameter $\beta$: $\beta \rightarrow \beta + \delta\beta$. It leads to variation of the CF action (2.1): $S \rightarrow S + \delta\beta S$, where

$$\delta\beta S = \left( \frac{1}{4} B^a B^a + \frac{1}{4} \tilde{B}^a \tilde{B}^a + m^2 \bar{C}^a C^a \right) \delta\beta.$$  \hspace{1cm} (3.3)

Then we have

$$Z_{\beta+\delta\beta} = \int D\phi \exp \left( \frac{i}{\hbar} \left[ S + \delta\beta S \right] \right).$$  \hspace{1cm} (3.4)

From (3.3) and (3.4) it follows the equation

$$\frac{\partial Z_\beta}{\partial \beta} = \frac{i}{4\hbar} < B^a B^a > + \frac{i}{4\hbar} < \tilde{B}^a \tilde{B}^a > + \frac{i}{\hbar} m^2 < \bar{C}^a C^a >$$  \hspace{1cm} (3.5)

where $< \cdots >$ means a vacuum expectation value of corresponding quantities, for example,

$$< \bar{C}^a C^a > = \int D\phi \ \bar{C}^a C^a \exp \left( \frac{i}{\hbar} S \right).$$  \hspace{1cm} (3.6)

Now let us use the invariance of $S$ (2.1) under the modified BRST transformation (2.12) to investigate the functional $Z_{\beta+\delta\beta}$. To this end, in the functional integral (3.4) we can make a change of variables being given by Eqs. (2.12) with some functional $\Lambda = \Lambda(\phi)$ instead of the constant Grassmann odd parameter $\theta$. It is clear that the CF action (2.1) is invariant under such a change of variables. If we restrict ourself to the first order in $\Lambda(\phi)$ and $\delta\beta$ then there appears contribution only coming from the integration measure

$$Z_{\beta+\delta\beta} = \int D\phi \ exp \left( \frac{i}{\hbar} \left[ S + \delta\beta S + \delta M \right] \right),$$  \hspace{1cm} (3.7)
where

$$\delta M = -i\hbar \left( \frac{\delta \Lambda(\phi)}{\delta A^a_\mu} D^a_\mu C^b - \frac{1}{2} \frac{\delta \Lambda(\phi)}{\delta C^a} f^{abc} C^b C^c - \frac{\delta \Lambda(\phi)}{\delta B^a} B^a - m^2 \frac{\delta \Lambda(\phi)}{\delta B^a} C^a \right).$$  \ (3.8)$$

Choosing the functional $\Lambda(\phi)$ as

$$\Lambda(\phi) = \frac{i}{2\hbar} \left( \bar{C}^a B^a - \frac{1}{2} f^{abc} \bar{C}^a \bar{C}^b \bar{C}^c \right) \delta \beta$$  \ (3.9)$$

we find that

$$\delta \beta S + \delta M = \frac{1}{2} m^2 \bar{C}^a C^a \delta \beta.$$  \ (3.10)$$

In massless limit, the vacuum functional $Z_\beta$ does not depend on the parameter $\beta$. It is no wonder that there is no dependence on this parameter because in this limit the CF action reduces to the FP action for massless Yang-Mills when $\beta$ plays a role of gauge parameter and nilpotency of the BRST transformations is restored. If $m \neq 0$ then there is an essential dependence of vacuum functional on this parameter and $\beta$ becomes a physical parameter defining, for example, a mass, $m_c$, of scalar anticommuting fields $C^a$ and $\bar{C}^a$ in the form $m^2_c = \beta m^2$ because the equations of motion read

$$(\Box + m^2_c) C^a + \cdots = 0,$$

where $\Box = \partial^\mu \partial^\mu$ and the dots mean terms which are non-linear in $\phi^i$. Similar equations hold for fields $\bar{C}^a$. Unfortunately, we cannot use the relation (3.10) to find the representation of dependence of $Z_\beta$ on $\beta$

$$\frac{\partial Z_\beta}{\partial \beta} = \frac{i}{2\hbar} m^2 < \bar{C}^a C^a >$$

as one might think considering (3.7) and (3.10). In the case $m^2 \neq 0$ the dependence of $Z_\beta$ on $\beta$ becomes essential and the change of variables

$$\phi^i \to \phi'^i = \phi^i + \frac{i}{\hbar} \Lambda(\phi) R^i(\phi), \quad \Lambda(\phi) = \frac{i}{\hbar} \bar{\Lambda}(\phi)$$

used in (3.8) and (3.9) is beyond the strong definition of functional integral within loop expansions (in $\hbar$) [25]. Here the condensed notations $\delta_{mB} \phi^i = R^i(\phi) \theta$ for the modified BRST transformation (2.12) were used. Note that such kind of transformations serves as a tool to prove the gauge independence of vacuum functional (and physical quantities) in Yang-Mills theories as well as in general gauge theories [26]. In the case of gauge theories, it does not lead to conflicts if one considers physical quantities because they are gauge invariant ones and the change of variables touches a modification of gauge fixing functional only.

We can investigate dependence of Green’s functions on parameter $\beta$ for the CF model as well. The generating functional of Green’s functions, $Z_\beta(J)$, is written in the form

$$Z_\beta(J) = \int D\phi \exp \left( \frac{i}{\hbar} [S(\phi) + J_i \phi^i] \right),$$  \ (3.11)$$
where the action \( S \) is defined through relations (2.1), the set of fields \( \phi^i \) is given in (3.2), and \( J_i = (j^a_\mu, K^a, \tilde{K}^a, L^a) \) are usual sources to fields \( \phi^i \) with relevant distributions of Grassmann and ghost parities. Let us consider the CF model corresponding a small variation of parameter \( \beta \) (\( \beta \to \beta + \delta \beta \)). Then the generating functional for Green’s functions is

\[
Z_{\beta+\delta\beta}(J) = \int D\phi \exp \left( \frac{i}{\hbar} [S(\phi) + \delta \beta S + J_\mu \phi^\mu] \right) \tag{3.12}
\]

where \( \delta \beta S \) is defined in (3.3). As a result we obtain the equation

\[
\frac{\partial Z_{\beta}(J)}{\partial \beta} = \int \frac{i}{\hbar} \left[ \left( \frac{\delta}{\delta \hat{K}} \right)^2 - \frac{\hbar}{2i} \frac{\delta^2}{\delta \hat{L}^a \delta \hat{L}^a} \right] Z_{\beta}(J), \tag{3.13}
\]

describing the dependence of Green’s functions on the parameter \( \beta \). In terms of the generating functional of connected Green’s functions, \( W_\beta(J) = \hbar/i \ln Z_{\beta}(J) \), the equation (3.13) takes the form

\[
\frac{\partial W_\beta(J)}{\partial \beta} = \frac{1}{2} \left( \frac{\delta W_\beta}{\delta \hat{L}^a} \frac{\delta W_\beta}{\delta \hat{L}^a} + \frac{\hbar}{i} \frac{\delta^2 W_\beta}{\delta \hat{L}^a \delta \hat{L}^a} \right) - \\
- \frac{1}{2} \left( \frac{\delta W_\beta}{\delta \hat{L}^a} + \frac{\hbar}{i} \frac{\delta}{\delta \hat{L}^a} \right) N^a \left( \frac{\delta W_\beta}{\delta \hat{K}} + \frac{\hbar}{i} \frac{\delta}{\delta \hat{K}} \right) + \\
+ \frac{1}{4} N^a \left( \frac{\delta W_\beta}{\delta \hat{K}} + \frac{\hbar}{i} \frac{\delta}{\delta \hat{K}} \right) N^a \left( \frac{\delta W_\beta}{\delta \hat{K}} + \frac{\hbar}{i} \frac{\delta}{\delta \hat{K}} \right) + \\
+ m^2 \left( \frac{\delta W_\beta}{\delta \hat{K}^a} \frac{\delta W_\beta}{\delta \hat{K}^a} + \frac{\hbar}{i} \frac{\delta^2 W_\beta}{\delta \hat{K}^a \delta \hat{K}^a} \right). \tag{3.14}
\]

Introducing the generating functional of vertex functions (effective action), \( \Gamma_{\beta}(\phi) \), being defined through the Legendre transformation of \( W_{\beta}(J) \),

\[
\Gamma_{\beta}(\phi) = W_{\beta}(J) - J_i \phi^i, \quad \phi^i = \frac{\delta W_{\beta}}{\delta J_i}, \quad \frac{\delta \Gamma_{\beta}(\phi)}{\delta \phi^i} = - J_i, \tag{3.15}
\]

the equation corresponding to (3.14) has the form

\[
\frac{\partial \Gamma_{\beta}(\phi)}{\partial \beta} = \frac{1}{2} \hat{B}^a \hat{B}^a - \frac{1}{2} \hat{B}^a N^a (\hat{C}, \hat{\tilde{C}}) + \frac{1}{4} N^a (\hat{C}, \hat{\tilde{C}}) N^a (\hat{C}, \hat{\tilde{C}}) + m^2 \hat{C}^a \hat{\tilde{C}}^a, \tag{3.16}
\]

where the notations

\[
\hat{\phi}^i = \phi^i + i \hbar (\Gamma^{-1})_{ij} \frac{\delta}{\delta \phi^j}, \quad (\Gamma')_{ij} = \delta_{ij} \frac{\delta \Gamma_{\beta}}{\delta \phi^i} \quad (\Gamma''_{ij})_{k} = \delta_{kj}, \tag{3.17}
\]

were used. We see that the dependence of effective action on the parameter \( \beta \) does not disappear even on-shell defined by the equations of motion of \( \Gamma_{\beta}(\phi) \) that confirms a physical character of the parameter \( \beta \). In tree approximation \( \Gamma_{\beta} = S, \hat{\phi}^i = \hat{\phi}^i \), and from (3.10) it follows (3.3).
4 Discussion

We investigated the dependence of Green’s functions for the CF model on the parameter resembling the gauge parameter in massless Yang-Mills fields. In particular, it was shown that the effective action for this model depends on this parameter on-shell. It allowed to consider this parameter as a physical one which can be associated with definition of mass for scalar anticommuting fields of the CF model. It was found that violation of nilpotency of the BRST symmetry can be interpreted as a source for appearance of an additional physical parameter in comparison with a gauge theory for which the full configuration space has the same structure. This situation is quite similar to that in the Gribov-Zwanziger theory \[11, 12, 13\] when violation of the BRST symmetry of the Gribov-Zwanziger action was interpreted as a source for the Gribov parameter to be a physical parameter \[27\].

It was pointed out that from the beginning the CF model should be considered as non-degenerated system of massive vector fields and massive scalar anticommuting fields. From this point of view the analysis of unitarity given in \[1\] looks like incomplete because the physical state space should include particles corresponding to massive scalar anticommuting fields as real ones. In turn, presence of these particles in physical state space does not give a chance for the CF model to be unitary because of the breakdown of norm-positivity \[10, 16\].

Acknowledgments

The author thanks I.L. Buchbinder and I.V. Tyutin for useful discussions of this paper. The work is supported by the LRSS grant 224.2012.2 as well as by the RFBR grant 12-02-00121 and the RFBR-Ukraine grant 11-02-90445.

References

[1] K.-I. Kondo, A unitary and renormalizable model for massive Yang-Mills fields without Higgs fields, \[arXiv:1202.4162\] [hep-th].
[2] G. Curci and R. Ferrari, On a class of Lagrangian models for massive and massless Yang-Mills fields, Nuovo Cim. A32 (1976) 151.
[3] G. Curci and R. Ferrari, The unitarity problem and the zero-mass limit for a model of massive Yang-Mills theory, Nuovo Cim. A35 (1976) 1.
[4] J.de Boer, K. Skenderis, P. van Nieuwenhuzen and A. Waldron, On the renormalizability and unitarity of the Curci-Ferrari model for massive vector bosons, Phys. Lett. B367 (1996) 175.
[5] L.D. Faddeev and V.N. Popov, Feynman diagrams for the Yang-Mills field, Phys. Lett. B25 (1967) 29.
[6] C. Becchi, A. Rouet and R. Stora, Renormalization of the abelian Higgs-Kibble model, Commun. Math. Phys. 42 (1975) 127;
[7] I.V. Tyutin, *Gauge invariance in field theory and statistical physics in operator formalism*, Lebedev Inst. preprint N 39 (1975), arXiv:0812.0580.

[8] G. Curci and R. Ferrari, *Slavnov transformations and supersymmetry*, Phys. Lett. B63 (1976) 91.

[9] I. Ojima, *Another BRS transformation*, Prog. Theor. Phys. 64 (1979) 625.

[10] N. Nakanishi and I. Ojima, *Covariant operator formalism of gauge theories and quantum gravity*, World Scientific, Singapore, 1990.

[11] V.N. Gribov, *Quantization of Nonabelian Gauge Theories*, Nucl.Phys. B139 (1978) 1.

[12] D. Zwanziger, *Action from the Gribov horizon*, Nucl. Phys. B321 (1989) 591.

[13] D. Zwanziger, *Local and renormalizable action from the Gribov horizon*, Nucl. Phys. B323 (1989) 513.

[14] P. Lavrov, O. Lechtenfeld and A. Reshetnyak, *Is soft breaking of BRST symmetry consistent?*, JHEP 1110 (2011) 043; arXiv:1108.4820 [hep-th].

[15] P. M. Lavrov, O. V. Radchenko and A. A. Reshetnyak, *Soft breaking of BRST symmetry and gauge dependence*, Mod. Phys. Lett. A27 (2012) 1250067; arXiv:1201.4720 [hep-th].

[16] T. Kugo and I. Ojima, *Local covariant operator formalism of non-abelian gauge theories and quark confinement problem*, Progr. Theor. Phys. Suppl. 66 (1979) 1.

[17] I. Ojima, *Comments on massive and massless Yang-Mills Lagrangian with quartic coupling of Faddeev-Popov ghosts*, Z. Phys. C13 (1982) 173.

[18] R.E. Kallosh and I.V. Tyutin, *The equivalence theorem and gauge invariance in renormalizable theories*, Sov. J. Nucl. Phys. 17 (1973) 98.

[19] I.V. Tyutin, *Once again on the equivalence theorem*, Phys. Atom. Nucl. 65 (2002) 194.

[20] B.S. DeWitt, *Dynamical Theory of Groups and Fields*, Gordon and Breach, New York, 1965.

[21] I.A. Batalin, P.M. Lavrov and I.V. Tyutin, *Covariant quantization of gauge theories in the framework of extended BRST symmetry*, J. Math. Phys. 31 (1991) 1487.

[22] I.A. Batalin, P.M. Lavrov and I.V. Tyutin, *An Sp(2)-covariant quantization of gauge theories with linearly dependent generators*, J. Math. Phys. 32 (1991) 532.

[23] I.A. Batalin, P.M. Lavrov and I.V. Tyutin, *Remarks on the Sp(2) - covariant quantization of gauge theories*, J. Math. Phys. 32 (1991) 2513.

[24] P.M. Lavrov, *Sp(2) renormalization*, Nucl. Phys. B849 (2011) 503.

[25] A.A. Slavnov, *Continual integral in perturbation theory*, Theor. Math. Fiz. 22 (1975) 177.

[26] I.A. Batalin and G.A. Vilkovisky, *Gauge algebra and quantization*, Phys. Lett. 102B (1981) 27; *Quantization of gauge theories with linearly dependent generators*, Phys. Rev. D28 (1983) 2567.

[27] N. Vandersickel and D. Zwanziger, *The Gribov problem and QCD dynamics*, arXiv:1202.1491 [hep-th].