ФИЗИКА ЭЛЕМЕНТАРНЫХ ЧАСТИЦ И АТОМНОГО ЯДРА. ТЕОРИЯ

ON COMPOSITE FIELDS APPROACH
TO GRIBOV COPIES ELIMINATION
IN YANG–MILLS THEORIES

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We suggest a way of introduction of the Gribov–Zwanziger horizon functional H for Yang–Mills theories by means of composite fields technique, as $\sigma(\phi) = H$. A different form of the same horizon functional in gauges χ and χ' is taken into account via (gauged) field-dependent BRST transformations connecting quantum Yang–Mills actions in these gauges. We introduce generating functionals of Green's functions with composite fields and derive the corresponding Ward identities. A study of gauge dependence shows that the effective action in Yang–Mills theories with the composite field H does not depend on the gauge on the extremals determined by the Yang–Mills fields ϕ alone.

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INTRODUCTION

The BRST symmetry concept, expressing a gauge invariance via a one-parameter global supersymmetry [1], appears as a defining tool for quantum description of the gauge theory, because of all known fundamental interactions are described in terms of Yang–Mills type theories [2], provides the success of perturbative calculations at high energy as well as strong evidence that the interactions of quarks and gluons are correctly described by QCD.

The problem of gauge-fixing with the help of a differential condition (like Coulomb or Landau gauges) as shown by V. Gribov [3] and developed by I. Singer [4] for Yang–Mills (YM) theories cannot be correctly realized within Faddeev–Popov (FP) procedure [5] even perturbatively for a whole spectrum of the momenta distribution for the gauge fields, A_μ^a, in the deep IR region, due to an infinitely large number of discrete gauge copies, appearing outside of so-called first Gribov region, $\Omega(A)$. The resolution of this problem can be realized by an addition to the quantum action, constructed by FP receipt, of the special horizon functional, $H(A)$ [6] known as the Gribov–Zwanziger (GZ) theory, which was constructed only for the Landau gauge and is not BRST invariant.

Until recently, the study of the complete GZ theory has been carried out almost entirely within the Landau gauge (see, e.g., [7] and references therein), whereas a restriction to $\Omega(A)$

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in the path integral has been made for YM theories in the approximation being quadratic in the fields and using the covariant [8] and maximal Abelian gauges [9]. At the same time, there is a significant arbitrariness in the choice of admissible gauges, related, in part, to the choice of a reference frame, see, e.g., [10]. Thus, it is well known Green's functions depend on the choice of a gauge; however, this dependence has a special structure, such that it should be absent in physical quantities like the S-matrix. The contemporary study of gauge dependence and unitarity in the Lorentz-covariant quantum description of gauge models is based on BRST symmetry. Therefore, any violation of the BRST invariance may lead to gauge-dependent and to non-unitary S-matrix. A consideration of the first problem within so-called *soft BRST symmetry breaking* concept for YM and general gauge theories [11] within BV quantization scheme [13] revealed a requirement on special transformation of a gauge variation, δM, for BRST symmetry breaking term M in order to provide gauge independence of the effective action (EA) on mass shell. Recently, in [14] it was shown that this requirement on δM is always fulfilled within class of general gauge theory with *soft breaking of BRST symmetry* being based on the concept of field-dependent BRST symmetry transformations introduced within YM theories in [15] to relate quantum FP action given in a fixed gauge to the one in different gauge, and used to determine the GZ horizon functional within GZ theory in R_ξ-family of the gauges. This result solves the above problem of gauge independence for the gauge models with soft BRST symmetry breaking, if for the model given in any fixed gauge reference frame in addition to the quantum BV action the form of the functional M is also changed in another gauge by the rule, firstly, given in [11]. However, an unitarity, analyzed for YM theories in [12], met the obstacle for the resulting quantum theory within such an introduction of the soft BRST symmetry breaking term. We intend to introduce another way to involve GZ horizon functional $H(A)$ into path integral for YM theory, as the composite field [16, 17].

The purpose of the paper is, first, to consider an introduction of soft BRST symmetry breaking terms into FP quantum action by means of composite fields, second, to obtain Ward identities for YM theory with composite fields and study the gauge-dependence problem.

The paper is organized as follows. In Sec. 1 we remind on the key issues of the soft BRST symmetry breaking in YM theories and derive the Ward identity for the EA, which provides a basic result involving the variation of the EA under a variation of the gauge fermion. Section 2 is devoted to application of the above result to GZ theory. The use of the composite field technique for an incorporation of the Gribov horizon functional $H(A)$ and derivation of the Ward identities together with describing gauge-dependence problem is considered in Sec. 3.

1. SOFT BRST SYMMETRY BREAKING IN YM THEORIES

An extended by antifields, ϕ_A^*, generating functional of Green's functions, $Z_M = Z_M(J, \phi^*)$, for YM theory with soft BRST symmetry breaking term, $M = M(\phi, \phi^*)$, $[M(\phi, 0) = m(\phi)]$ is determined following [11] by the path integral depending on sources J_A to the total set of the fields, ϕ^A, with classical gauge A_μ^a, ghost C^a, antighost $\overline{C}^a$, Nakanishi–Lautrup B^a fields, for $a = 1, \ldots, \dim SU(N)$, $\mu = 0, 1, \ldots, D-1$ in the condensed DeWitt's and Ref. [11] notations as follows:

$$Z_M = \int d\phi \, \exp\left\{\frac{\imath}{\hbar}\big(S_0(A) + s\psi(\phi) + \phi_A^* s\phi^A + M + J_A\phi^A\big)\right\}. \tag{1}$$

Here the classical action, $S_0(A)$, to be invariant under gauge transformations, $\delta A_\mu^a = D_\mu^{ab} \xi^b$, with YM covariant derivative, D_μ^{ab}, completely antisymmetric $SU(N)$ structure constants f^{abc} and arbitrary functions ξ^b on Minkowski space-time $R^{1,D-1}$, nilpotent Slavnov variation, s, gauge fermion, ψ, are given by the relations

$$S_0(A) = -\frac{1}{4} \int d^D x F_{\mu\nu}^a F^{\mu\nu a} \quad \text{for} \quad [D_\mu^{ab}, F_{\mu\nu}^a] = [\delta^{ab}\partial_\mu + f^{acb}A_\mu^c, \partial_{[\mu}A_{\nu]}^a + f^{abc}A_\mu^b A_\nu^c], \quad (2)$$

$$sF(\phi, \phi^*) = \frac{\delta F}{\delta \phi^A} s\phi^A, \quad s\phi^A = (D_\mu^{ab}C^b, \frac{1}{2}f^{abc}C^b C^c, B^a, 0), \quad (3)$$

$$\psi(\phi) = \overline{C}^a \chi^a(A, B) \quad \text{for the gauge} \quad \chi^a(A, B) = \partial^\mu A_\mu^a + \frac{\xi}{2}B^a = 0. \quad (4)$$

In terms of the operator s BRST non-invariance of the bosonic functionals, M, m means, $(sM, sm) \neq (0,0)$. For vanishing M we deal with usual path integral, $Z \equiv Z_0(J, \phi^*)$.

The Ward identities for Z_M and for EA (generating functional of vertex Green's functions), $\Gamma_M = \Gamma_M(\phi, \phi^*)$, introduced via Legendre transform of $\frac{\hbar}{\imath} \ln Z_M$ with respect to J_A: $\Gamma_M = \frac{\hbar}{\imath} \ln Z_M - J\phi$ for $\phi = \frac{\hbar}{\imath}(\delta \ln Z_M)/(\delta J)$, have the form

$$\left(J_A + M_A \left(\frac{\hbar}{i} \frac{\delta}{\delta J}, \phi^* \right) \right) \left(\frac{\hbar}{i} \frac{\delta}{\delta \phi_A^*} - M^{A*} \left(\frac{\hbar}{i} \frac{\delta}{\delta J}, \phi^* \right) \right) Z_M = 0, \quad (5)$$

$$\frac{\delta \Gamma_M}{\delta \phi^A} \frac{\delta \Gamma_M}{\delta \phi_A^*} = \frac{\delta \Gamma_M}{\delta \phi^A} \widehat{M}^{A*} + \widehat{M}_A \frac{\delta \Gamma_M}{\delta \phi_A^*} - \widehat{M}_A \widehat{M}^{A*}, \quad (6)$$

where the notations $(M_A, M^{A*}) \left(\frac{\hbar}{i} \frac{\delta}{\delta J}, \phi^* \right) \equiv \left(\frac{\delta M}{\delta \phi^A}, \frac{\delta M}{\delta \phi_A^*} \right) \Big|_{\phi \to \frac{\hbar}{i} \frac{\delta}{\delta J}}$ and $\left(\widehat{M}_A, \widehat{M}^{A*} \right) \equiv (M_A, M^{A*}) \Big|_{\phi \to \widehat{\phi}}$ have been used in accordance with [11, 14] with operator-valued fields

$$\widehat{\phi}^A = \phi^A + \imath\hbar (\Gamma_M''^{-1})^{AB} \frac{\delta_l}{\delta \phi^B} \text{ with } (\Gamma_M''^{-1})^{AC}(\Gamma_M'')_{CB} = \delta_B^A, (\Gamma_M'')_{AB} = \frac{\delta_l}{\delta \phi^A}\left(\frac{\delta \Gamma_M}{\delta \phi^B} \right). \quad (7)$$

Note, in obtaining (5) and (6) we have not utilized the BRST symmetry breaking equation [11], $M_A, M^{A*} = 0$. In turn, the result of gauge dependence research for Z_M, Γ_M can be presented in the form, for $\delta\psi = \xi/2\overline{C}^a B^a$,

$$\delta Z_M = \frac{\imath}{\hbar} \left[\left(J_A + M_A \left(\frac{\hbar}{\imath} \frac{\delta}{\delta J}, \phi^* \right) \right) \times \right.$$
$$\left. \times \left(\frac{\delta}{\delta \phi_A^*} - \frac{\imath}{\hbar} M^{A*} \left(\frac{\hbar}{\imath} \frac{\delta}{\delta J}, \phi^* \right) \right) \delta\psi \left(\frac{\hbar}{\imath} \frac{\delta}{\delta J} \right) + \delta M \left(\frac{\hbar}{\imath} \frac{\delta}{\delta J}, \phi^* \right) \right] Z_M, \quad (8)$$

$$\delta \Gamma_M = \frac{\delta \Gamma_M}{\delta \phi^A} \widehat{F}_M^A \widehat{\delta\Psi} - \widehat{M}_A \widehat{F}_M^A \widehat{\delta\psi} + \widehat{\delta M} \quad (9)$$

$$\text{for} \quad \widehat{F}_M^A \equiv -\frac{\delta}{\delta \phi_A^*} + (-1)^{\varepsilon_B(\varepsilon_A+1)}(\Gamma_M''^{-1})^{BC}\left(\frac{\delta_l}{\delta \phi^C} \frac{\delta \Gamma_M}{\delta \phi_A^*} \right) \frac{\delta_l}{\delta \phi^B} \quad \text{and} \quad \varepsilon_A \equiv \varepsilon(\phi^A). \quad (10)$$

On the extremals, $\left(J_A, \delta\Gamma_M \delta\phi^A\right) = 0$ respectively for Z_M, Γ_M the corresponding variations are vanishing, $\left(\delta Z_M\big|_{J=0}, \delta\Gamma_M\big|_{\Gamma_{MA}=0}\right) = 0$, provided that the gauge variation $\psi \to \psi + \delta\psi$ affects not only the BRST exact part of the action, changed on the term, $s\delta\psi$, but also the functional, M, $\delta M = -(sM)\dfrac{\imath}{\hbar}\delta\psi$, which was shown in [14]. Now, we are able to apply these results for special choice of BRST symmetry breaking term, when $M = H(A)$.

2. GAUGE INDEPENDENCE FOR GRIBOV–ZWANZIGER THEORY

In case of GZ theory the Gribov horizon functional $H(A)$ in the Landau gauge, $\psi_0(\phi)$ determined by Eq. (4) for $\xi = 0$ [6], and its Slavnov variation in non-local formulation read

$$H(A) = \gamma^2 \left(f^{abc}A_\mu^b (K^{-1})^{ad} f^{dec} A^{e\mu} + D(N^2-1)\right) \quad \text{for} \quad (K^{-1})^{ad}(K)^{db} = \delta^{ab}, \quad (11)$$

$$sH = \gamma^2 f^{abc} f^{cde} \left[2D_\mu^{bq} C^q (K^{-1})^{ad} - f^{mpn} A_\mu^b (K^{-1})^{am} K^{pq} C^q (K^{-1})^{nd}\right] A^{e\mu}, \quad (12)$$

with a thermodynamic Gribov parameter, γ, to be determined in self-consistent way from the gap equation, $\dfrac{\partial}{\partial\gamma}\left(\dfrac{\hbar}{i}\ln Z_H(0,0)\right) = 0$. The Ward identities for $Z_H(J,0)$ and for EA, $\Gamma_H(\phi,0)$, are easily obtained from Eqs. (5) and (6) as follows:

$$J_A\langle s\phi^A\rangle + \langle H_{\mu a} sA^{\mu a}\rangle = 0, \quad \text{where} \quad \langle\mathcal{O}\rangle = Z_H^{-1}(J,0)\mathcal{O}\left(\frac{\hbar}{i}\frac{\delta}{\delta J}\right)Z_H(J,0), \quad (13)$$

$$\frac{\delta\Gamma_H(\phi,0)}{\delta\phi^A}s\phi^A(\widehat{\phi}) = \widehat{H}_{\mu a}sA^{\mu a}(\widehat{\phi}), \quad \text{where} \quad H_{\mu a}(A) = \frac{\delta H}{\delta A^{\mu a}}, \quad (14)$$

with the use of the Sec. 2 notations.

Equations (13) and (14), together with the representations (8)–(10), permit one to get the variations for $Z_H(J,0)$ and $\Gamma_H(\phi,0)$ as

$$\delta Z_H = \left(\frac{\imath}{\hbar}\right)^2 \left(J_A\langle(s\phi^A)\delta\psi\rangle + \langle H_{\mu a}(sA^{\mu a})\delta\psi\rangle + \frac{\hbar}{\imath}\langle\delta H\rangle\right)Z_H, \quad (15)$$

$$\delta\Gamma_H = \frac{\delta\Gamma_H}{\delta\phi^A}\widehat{F}_H^A\langle\delta\Psi\rangle\big|_{\phi^*=0} - \widehat{H}_{\mu a}\widehat{F}_H^{\mu a}\widehat{\delta\psi}\big|_{\phi^*=0} + \widehat{\delta H}. \quad (16)$$

Because the gauge variation, $\delta G(A)$, is induced by the variation, $\delta\psi$, of any functional, $G(A)$, given only on YM fields configuration space may be presented by means of the gauge transformation with parameters ξ^b constructing from $\delta\psi$ as

$$G(A) \to G(A) + \delta G(A) = G(A) + G_{\mu a}D^{\mu ab}\xi^b \quad \text{for} \quad \xi^b = -\frac{\imath}{\hbar}C^b\delta\psi, \quad (17)$$

which was, in fact, shown, firstly, with the help of field-dependent BRST transformations in [15], it is obvious from (15) and (16) that on mass-shell we have $\delta Z_H\big|_{J=0} = 0$, $\delta\Gamma_H\big|_{\Gamma_{HA}=0} = 0$. As byproduct, the representation for $H(A)$ in new gauge reference frame, $\psi + \delta\psi$, not necessarily related by infinitesimal gauge variation follows from (17) for $G = H$.

3. A COMPOSITE FIELD REPRESENTATION FOR
THE GRIBOV HORIZON FUNCTIONAL

Since unitarity cannot be verified explicitly, due to BRST non-invariance of the action with $M(\phi, \phi^*)$ $[H(A)]$ term for the model with Z_M (1), in particular, for GZ model with non-local $H(A)$ (11)[1], we will treat this term as composite field, $\sigma(A) = H(A)$. Doing so, we define generating functional of Green's functions with composite fields, by the relation $Z_{L\sigma} = Z_M$ as

$$Z(J, \phi^*, L) = Z_{L\sigma}(J, \phi^*) =$$
$$= \int d\phi \, \exp\left\{ \frac{\imath}{\hbar} \left(S_0(A) + s\psi + \phi_A^* s\phi^A + L_m \sigma^m(\phi) + J_A \phi^A \right) \right\}, \quad (18)$$

with sources L_m, $\varepsilon(L_m) = 0$, for σ^m. Restricted by the case $m = 1$, we introduce EA, $\Gamma(\phi, \phi^*, \Sigma) \equiv \Gamma_\Sigma$, with composite field via Legendre transform of $\ln Z_{L\sigma}$ w.r.t. J, L by the formulae following [16]:

$$\Gamma_\Sigma = \frac{\hbar}{\imath} \ln Z_{L\sigma} - J\phi - L(\Sigma + \sigma(\phi)) \text{ for } \phi^A = \frac{\hbar}{\imath} \frac{\delta \ln Z_{L\sigma}}{\delta J_A} \text{ and } \Sigma = \frac{\hbar}{\imath} \frac{\delta \ln Z_{L\sigma}}{\delta \sigma} - \sigma(\phi), \quad (19)$$

where

$$(J_A, L) = \left(-\frac{\delta \Gamma_\Sigma}{\delta \phi^A} + \frac{\delta \Gamma_\Sigma}{\delta \Sigma} \frac{\delta \sigma(\phi)}{\delta \phi^A}, -\frac{\delta \Gamma_\Sigma}{\delta \Sigma} \right) \equiv \mathcal{N}_\alpha \text{ and } \Phi^\alpha \equiv (\phi^A, \Sigma). \quad (20)$$

Note, first, that tree approximation for Γ_Σ in loop expansion, $\Gamma_\Sigma = \sum_n \hbar \Gamma_\Sigma^{(n)}$, coincides with FP action, $\Gamma_\Sigma^{(0)} = S_0 + s\psi + \phi^* s\phi$, providing an influence of the composite field beyond the tree level. Second, the boundary conditions for $L = 0$: $Z(J, \phi^*, 0) = Z_0(J, \phi^*)$, and for $L = 1$: $Z(J, \phi^*, 1) = Z_H(J, \phi^*)$, are fulfilled.

The Ward identities for $Z_{L\sigma}(J, 0)$ and Γ_Σ take the form being inherited from (13) and (14),

$$J_A \langle s\phi^A \rangle_L + L \langle H_{\mu a} sA^{\mu a} \rangle_L = 0 \text{ for } \langle \mathcal{O} \rangle_L = Z_{L\sigma}^{-1}(J, 0) \mathcal{O}\left(\frac{\hbar}{\imath} \frac{\delta}{\delta J} \right) Z_{L\sigma}(J, 0), \quad (21)$$

$$\frac{\delta \Gamma(\phi, 0, \Sigma)}{\delta \phi^A} s\phi^A(\widehat{\phi}) = -\frac{\delta \Gamma(\phi, 0, \Sigma)}{\delta \Sigma} \left(\widehat{H}_{\mu a} - H_{\mu a}(\phi) \right) sA^{\mu a}(\widehat{\phi}). \quad (22)$$

Again with the use of (21) and (22) the study of gauge dependence for $Z_{L\sigma}(J, 0)$ and Γ_Σ looks like

$$\delta Z_{L\sigma} = \left(\frac{\imath}{\hbar} \right)^2 \left(J_A \langle (s\phi^A) \delta\psi \rangle_L + L \left\{ \langle H_{\mu a}(sA^{\mu a}) \delta\psi \rangle_L + \frac{\hbar}{\imath} \langle \delta H \rangle_L \right\} \right) Z_{L\sigma}, \quad (23)$$

$$\delta \Gamma_\Sigma = \frac{\delta \Gamma_\Sigma}{\delta \phi^A} \widehat{F}_{L\sigma}^A \langle \delta\Psi \rangle \big|_{\phi^*=0} - \frac{\delta \Gamma_\Sigma}{\delta \Sigma} \left\{ \widehat{\delta H} - \left(\widehat{H}_{\mu a} - H_{\mu a}(A) \right) \widehat{F}_{L\sigma}^{\mu a} \widehat{\delta\psi} \big|_{\phi^*=0} \right\} \quad (24)$$

for $\widehat{F}_{L\sigma}^{\mu a} \equiv -\frac{\delta}{\delta A_{\mu a}^*} + (-1)^{\varepsilon_\beta} (K''^{-1})^{\beta\alpha} \left(\frac{\delta_l}{\delta \Phi^\alpha} \frac{\delta \Gamma_\Sigma}{\delta A_{\mu a}^*} \right) \frac{\delta_l}{\delta \Phi^\beta} \text{ and } (K'')_{\beta\alpha} = \frac{\delta \mathcal{N}_\alpha}{\delta \Phi^\beta}, \quad (25)$

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[1]For local representation of Gribov horizon functional in different gauges and study of the gauge dependence problem considered in Sec. 2, see [14, 15].

where we have used Eqs. (18)–(20) and Ward identities (21) and (22). Note that in the above representations for the variations $\delta Z_{L\sigma}$, $\delta\Gamma_\Sigma$ we have taken into account the variation of the composite field itself (cf. with [17]).

Resuming, we state that on mass-shell determined by the surfaces $J_A = 0$ and $(\delta\Gamma_\Sigma)/(\delta\phi^A) = 0$ and due to the rules of transforming of $H(A)$ under gauge variation $\delta\psi$, the generating functional $Z_{L\sigma}(J,0)$ and EA $\Gamma_\Sigma(\phi,0\Sigma)$, respectively, do not depend on the choice of the gauge for any corresponding values of the source L and additional extremal $(\delta\Gamma_\Sigma)/(\delta\Sigma)$, providing basis for consistency of this approach.

Concluding, we may use the latter concept of consideration of Gribov horizon functional as the composite field to study unitarity of the GZ theory starting from some fixed gauge.

REFERENCES

1. *Becchi C., Rouet A., Stora R.* Renormalization of the Abelian Higgs–Kibble Model // Commun. Math. Phys. 1975. V. 42. P. 127–162;
   Tyutin I. V. Gauge Invariance in Field Theory and Statistical Physics in Operator Formalism. Lebedev Inst. Preprint. 1975. No. 39; arXiv:0812.0580[hep-th].

2. *Weinberg S.* The Quantum Theory of Fields. V. II. Cambridge: Cambridge Univ. Press, 1996.

3. *Gribov V. N.* Quantization of Non-Abelian Gauge Theories // Nucl. Phys. B. 1978. V. 139. P. 1–19.

4. *Singer I. M.* Some Remarks on the Gribov Ambiguity // Commun. Math. Phys. 1978. V. 60. P. 7–12.

5. *Faddeev L. D., Popov V. N.* Feynman Diagrams for the Yang–Mills Field // Phys. Lett. B. 1967. V. 25. P. 29–30.

6. *Zwanziger D.* Local and Renormalizable Action from the Gribov Horizon // Nucl. Phys. B. 1989. V. 323. P. 513–544.

7. *Capri M. A. L. et al.* A Remark on the BRST Symmetry in the Gribov–Zwanziger Theory // Phys. Rev. D. 2010. V. 82. P. 105019; arXiv:1009.4135 [hep-th].

8. *Sobreiro R. F., Sorella S. P.* A Study of the Gribov Copies in Linear Covariant Gauges in Euclidean Yang–Mills Theories // JHEP. 2005. V. 0506. 054; arXiv:hep-th/0506165.

9. *Dudal D. et al.* Gribov Ambiguities in the Maximal Abelian Gauge // Braz. J. Phys. 2007. V. 37. P. 320–325; arXiv:1210.5651[hep-th].

10. *DeWitt B. S.* Dynamical Theory of Groups and Fields. N.Y.: Gordon and Breach, 1965.

11. *Lavrov P., Lechtenfeld O., Reshetnyak A.* Is Soft Breaking of BRST Symmetry Consistent? // JHEP. 2011. V. 1110. 043; arXiv:1108.4820 [hep-th].

12. *Kugo T., Ojima I.* Local Covariant Operator Formalism of Non-Abelian Gauge Theories and Quark Confinement Problem // Progr. Theor. Phys. Suppl. 1979. V. 66. P. 1–130.

13. *Batalin I. A., Vilkovisky G. A.* Gauge Algebra and Quantization // Phys. Lett. B. 1981. V. 102. P. 27–31.

14. *Reshetnyak A.* On Gauge Independence for Gauge Models with Soft Breaking of BRST Symmetry. arXiv:1312.2092[hep-th].

15. *Lavrov P., Lechtenfeld O.* Gribov Horizon beyond the Landau Gauge // Phys. Lett. B. 2013. V. 725. P. 386–388.

16. *Cornwall J. M., Jackiw R., Tomboulis E.* Effective Action for Composite Operators // Phys. Rev. D. 1974. V. 10. P. 2428–2445.

17. *Lavrov P. M., Odintsov S. D.* The Gauge Dependence of the Effective Action of Composite Fields in General Gauge Theories // Intern. J. Mod. Phys. 1989. V. 4. P. 5205–5212.