A simplified version of Higher Covariant Derivative regularization.

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Abstract

A simplified version of Higher Covariant Derivative regularization for Yang-Mills theory is constructed. This may make Higher Covariant Derivative method more attractive for practical calculations.

1 Introduction.

The construction of invariant regularizations for gauge theories is very important both for the practical calculations and for the investigation of interesting nonperturbative phenomena. The most popular scheme as yet is a dimensional regularization [14]. The drawback of this regularization is an absence of obvious generalization to nonperturbative approaches and inapplicability to chiral and supersymmetric models. The lattice formulation [15] preserves gauge invariance and has nonperturbative meaning, but it also does not seem to be applicable for the regularization of chiral, topological and supersymmetric theories.

The regularization of gauge theories by higher covariant derivatives [2, 3] and gauge invariant Pauli-Villars regulators [4] is an alternative nonperturbative regularization scheme which has an advantage of being applicable to chiral and supersymmetric models. In Ref.[1] we constructed an unambiguous regularized functional which avoided the objections raised in Refs. [6, 7, 8].

In this paper we obtain a more compact and convenient expression for the regularized Lagrangian which is simpler and requires much less computational efforts than the one in Ref.[1]. This may make the Higher Covariant Derivative method more attractive for the practical calculations in models where dimensional regularization is not applicable.

2 Obtaining the regularized functional.

One of the ways to regularize the theory is to modify the propagators by introducing into Lagrangian higher derivative terms. However this procedure breaks gauge invariance. To preserve the symmetry one can add into the Yang-Mills(YM) Lagrangian the terms containing higher covariant derivatives [2, 3].
We choose the regularized Lagrangian in the form proposed in [1]:

\[ \mathcal{L}_{YM}(x) \rightarrow \mathcal{L}^\Lambda(x) = \mathcal{L}_{YM}(x) + \frac{1}{4\Lambda^4} \left[ \frac{\delta S_0}{\delta A^a_\mu(x)} \right]^2 \]  

(1)

where:

\[ \mathcal{L}_{YM} = \frac{1}{8} \text{tr} F^2_{\mu\nu} \quad ; \quad S_0 = \int (\nabla^2 F^a_{\mu\nu})^2 dx \]  

(2)

Here \( F^a_{\mu\nu} \) is the usual curvature tensor and \( \nabla \) is the covariant derivative:

\[ F^a_{\mu\nu} = \partial_\nu A^a_\mu - \partial_\mu A^a_\nu + t^{abc} A^b_\mu A^c_\nu \]  

(3)

\[ \nabla^a_\mu = \partial_\mu \delta^a_\mu + t^{abc} A^c_\mu \]  

(4)

Then the generating functional for the case of the Lorentz gauge can be written as follows:

\[ Z[J] = \int \exp \left\{ i \int \left[ \mathcal{L}_{YM}(x) + J^a_\mu(x) A^a_\mu(x) - \frac{1}{2} (h^a_\rho(x))^2 \right] dx + \frac{i}{4\Lambda^4} \int \left[ \frac{\delta S_0}{\delta A^a_\rho(x)} \right]^2 dx \} \delta(\partial_\mu A_\mu) \det M(A) \prod_x DA_\mu \]  

(5)

where \( M^{ab} = \partial_\mu \nabla^a_\mu \). It is more convenient for us to transform this functional to another form. Firstly along the lines of Ref.[1] we bring in the integration over the auxiliary fields \( h_\rho \):

\[ Z[J] = \int \exp \left\{ i \int \left[ \mathcal{L}_{YM}(x) + J^a_\mu(x) A^a_\mu(x) - \frac{1}{2} (h^a_\rho(x))^2 \right] dx + \frac{i}{2\Lambda^5} \int h^a_\mu(x) \frac{\delta S_0}{\delta A^a_\rho(x)} dx \} \delta(\partial_\mu A_\mu) \det M(A) \prod_x Dh_\rho DA_\mu \]  

(6)

To separate the integration over covariant transversal and longitudinal parts of \( h_\rho \) we multiply the functional (6) by "unity":

\[ \text{det} \nabla^2 \int \prod_x \delta(\nabla_\mu (h_\mu + \nabla_\mu u)) Du(x) = 1 \]  

(7)

and change variables:

\[ h_\mu \rightarrow h_\mu - \nabla_\mu u \]  

(8)

Then we get:

\[ Z[J] = \int \exp \left\{ i \int \left[ \mathcal{L}_{YM}(x) + J^a_\mu(x) A^a_\mu(x) - \frac{1}{2} (h^a_\rho(x) - \nabla_\rho u^a_\rho(x))^2 \right] dx + \right. \]  

\[ + \left. \frac{i}{2\Lambda^5} \int (h^a_\rho(x) - \nabla_\rho u^a_\rho(x)) \frac{\delta S_0}{\delta A^a_\rho(x)} dx \right\} \times \]  

\[ \times \delta(\partial_\mu A_\mu) \delta(\nabla_\mu h_\mu) \det M(A) \det \nabla^2 \prod_x Du Dh_\rho DA_\mu \]  

(9)
Due to the gauge invariance of $S_0$:

$$
\int \nabla_\mu u^\mu(x) \frac{\delta S_0}{\delta A^\mu_\rho(x)} dx = 0 \tag{10}
$$

only covariant transversal part of $h_\mu$ gives nontrivial contribution to eq. (3). On the surface $\nabla_\mu h_\mu = 0$ we have:

$$
\int h^\mu_\rho(x) \nabla_\mu u^\mu(x) dx = 0 \tag{11}
$$

Using Eqs. (11,12) we can omit these terms from the exponent of equation (3). Integrating over $u(x)$ we get:

$$
Z[J] = \int \exp \left\{ i \int \left[ \mathcal{L}_{YM}(x) + J^\mu_\mu(x) A^\mu_\rho(x) - \frac{1}{2} (h^\mu_\rho(x))^2 + \frac{1}{2 \Lambda^5} \int h^\mu_\rho(x) \frac{\delta S_0}{\delta A^\mu_\rho(x)} dx \right] \right\} \times \\
\times \delta(\nabla_\mu h_\mu) \delta(\partial_\rho A_\rho) \det \mathcal{M}(\mathcal{A}) \det \frac{i}{2} \nabla^2 \prod_x Dh_\mu DA_\rho \tag{12}
$$

The free propagators generated by the exponent in Eq. (12) have the following UV behaviour: $A_\mu A_\nu \sim k^{-12}; h_\rho h_\sigma \sim k^{-10}; A_\mu h_\rho \sim k^{-6}$. Using the functional (12) one can see that the divergence index of arbitrary diagram is equal to:

$$
\omega_1 = 4 - 4(I - 1) - 4N_\mathcal{A} - 6L_\mathcal{A} - 4L_h - E_\mathcal{A} - E_h - 3E_{gh} - \frac{7}{2} E_\delta \tag{13}
$$

where $I$ is the number of loops, $L_\mathcal{A}, L_h$ are the numbers of internal and $E_\mathcal{A}, E_h$ are the numbers of external lines of transversal parts of the fields $\mathcal{A}$ and $h$ correspondingly, $N_\mathcal{A}$ is the number of vertices generated by $\mathcal{L}_{YM}$, $E_{gh}$ is the number of external lines of the fields representing $\det \mathcal{M}(\mathcal{A})$ and $\det \frac{i}{2} \nabla^2$, $E_\delta = E_\mathcal{A}^{long} + E_\lambda$ is the number of external lines representing longitudinal part of field $h$ and auxiliary field $\lambda$ arising in the representation of delta-function:

$$
\delta(\nabla_\mu h_\mu) = \int \exp \left\{ \int \lambda \nabla_\mu h_\mu dx \right\} \prod_x \lambda \tag{14}
$$

($E_{gh}$ and $E_\delta$ are even numbers). One can see that if $E_h > 0$, the diagram is convergent as it includes at least one internal line of $\mathcal{A}$ field. The only divergent graphs are the one loop diagrams with $L_\mathcal{A} = L_h = N_\mathcal{A} = 0$, $E_h = E_{gh} = E_\delta = 0$ and $E_\mathcal{A} = 2, 3, 4$. Therefore the sum of all one-loop divergent diagrams of the functional (12) with external gauge field lines $\mathcal{A}$ can be presented as follows:

$$
Z_{div}[\mathcal{A}] = \int \exp \left\{ \frac{i}{2 \Lambda^5} \int h^\mu_\rho(x) \frac{\delta^2 S_0}{\delta \nabla_\mu h_\rho(x) \delta \nabla_\mu h_\rho(y)} q^\mu_\rho(y) dy + \ldots \right\} \delta(\nabla_\mu h_\mu) \delta(\partial_\rho q_\rho) \\
\prod_x Dh_\mu Dq_\mu \det \mathcal{M}(\mathcal{A}) \det \frac{i}{2} \nabla^2 \tag{15}
$$

Here $\ldots$ denotes the terms which provide the infrared convergence of the integral (13) and have no influence on its ultraviolet behavior.

Although higher loop diagrams acquire by power counting a negative superficial divergent dimension, the divergencies of one loop diagrams are not smoothed and these diagrams require some additional regularization.
It was proposed in the paper [4] (see also [5]) that such a regularization may be provided by a modified Pauli-Villars (PV) procedure.

In this article we introduce the PV interaction which satisfies the following conditions:

A) it is gauge invariant;

B) it completely decouples from the physical fields in the limit when the masses of PV fields go to infinity;

C) it exactly compensates the remaining divergencies of the functional (12) and the cancelation of divergent contributions holds for individual diagrams (i.e. the divergent diagrams of the functional (13) and PV interaction have the same structure and one needs no auxiliary regulator).

D) when integrated over the fields $A_{\mu}$ it does not produce any new divergent subgraphs in the multiloop diagrams (this is known as the problem of overlapping divergencies, see for example [3, 10]).

The PV functional with such properties and external gauge field lines $A$ looks as follows:

$$I_{PV}(A) = \int \exp \left\{ \frac{i}{2\Lambda^5} \int \frac{\delta^2 S_0}{\delta A_\mu^a(x) \delta \bar{A}_\mu^b(y)} (B_\mu - \nabla_\mu \nabla^{-2} \nabla_\nu B_\nu)^b_y dxy + iM \int \delta_\mu^a(x) (B_\mu^{tr} + \nabla_\mu \psi)^a_x dx \right\} \delta(\partial_\rho \bar{B}_\rho) \prod_x DB_\mu DB_\mu^{tr} det^{-1} \mathcal{M}(A)$$

$$\prod_j det \frac{1}{2}(\nabla^2 - M_j^2)$$  \hspace{1cm} (16)

Here $\bar{B}_\mu$, $B_\mu$ are anticommuting PV fields, and PV conditions hold:

$$\sum_j c_j = -1 \hspace{0.5cm}; \hspace{0.5cm} \sum_j c_j M_j^2 = 0 \hspace{1cm} (17)$$

Let us check that the functional (16) satisfies all the four conditions imposed above.

A) To demonstrate the gauge invariance of the functional (16) let us make the change of variables:

$$B_\mu = B_\mu^{tr} + \nabla_\mu \psi$$  \hspace{1cm} (18)

where:

$$B_\mu^{tr} = B_\mu - \nabla_\mu \nabla^{-2} \nabla_\nu B_\nu \hspace{0.5cm}; \hspace{0.5cm} \psi = \nabla^{-2} \nabla_\nu B_\nu \hspace{0.5cm}; \hspace{0.5cm} \nabla_\mu B_\mu^{tr} = 0$$  \hspace{1cm} (19)

and $B_\mu^{tr}, \psi$ are anticommuting fields. The Jacobian of the transformation (18) is equal to $det^{-\frac{1}{2}} \nabla^2$ (the inverse sign of the power of determinant is due to the Grassmannian nature of variables). Then the functional (16) acquires the form:

$$I_{PV}(A) = \int \exp \left\{ \frac{i}{2\Lambda^5} \int \frac{\delta^2 S_0}{\delta A_\mu^a(x) \delta \bar{A}_\mu^b(y)} (B_\mu^{tr})^b_y dxy + iM \int \delta_\mu^a(x) (B_\mu^{tr} + \nabla_\mu \psi)^a_x dx \right\} \delta(\partial_\rho \bar{B}_\rho) \prod_x DB_\mu DB_\mu^{tr} D\psi \times$$

$$\times det^{-1} \mathcal{M}(A) det^{-\frac{1}{2}} \nabla^2 \prod_j det \frac{1}{2}(\nabla^2 - M_j^2)$$  \hspace{1cm} (20)
One sees that the only \( \psi \)-dependent term in the exponent of the expression (20) is equal to zero on the surface \( \nabla \mu B^\mu = 0 \):

\[
\int B_\mu^a(x) \nabla^\mu \psi^a(x) dx = 0
\]  

(21)

This allows us to rewrite the functional (20) in the form:

\[
I_{PV}(A) = \int I^1(A, B^{tr}) I^2(A, B^{tr}) \prod_x DB^{tr}
\]  

(22)

where:

\[
I^1(A, B^{tr}) = \int \exp \left\{ i \int \frac{\delta^2 S_0}{\delta A^a_\mu(x) \delta A^a_\mu(y)} (B^{tr}_{\mu})^a_b dx dy + iM \int B^{tr}_\mu(x) (B^{tr}_{\mu})^a_b dx \right\} \prod_x DB^{tr}_\mu \delta \left( \nabla^\mu B^{tr}_\mu \right) \prod \det^{-\frac{1}{2}} \det \prod \left( \nabla^2 - M^2 \right)
\]  

(23)

and

\[
I^2(A, B^{tr}) = \det^{-1} \mathcal{M}(A) \int \delta(\partial_\rho B^{tr}_\rho + \mathcal{M} \psi) \prod_x D\psi
\]  

(24)

The functional \( I^1(A, B^{tr}) \) is invariant under simultaneous transformations:

\[
\begin{align*}
A_\mu &\to A_\mu + [A_\mu, \epsilon] + \partial_\mu \epsilon \\
B^{tr}_\mu &\to B^{tr}_\mu + [B^{tr}_\mu, \epsilon] \\
B^{tr}_\mu &\to B^{tr}_\mu + [B^{tr}_\mu, \epsilon]
\end{align*}
\]  

(25)

The functional \( I^2(A, B^{tr}) \) is also gauge invariant because after the integration over \( \psi \) it is equal to nonessential constant and actually it does not depend on \( A \) and \( B^{tr} \) (remember that \( \psi \) is Grassmannian field and integration over \( \psi \) in (24) produces \( \det \mathcal{M}(A) \)). That proves the gauge invariance of the functional (14).

The demonstration of the gauge invariance of the PV interaction (16) is the most essential point of this paper. It allows us to avoid a rather complicated two-step procedure which we used in Ref.[1] and to make much simpler the resulting expression for the regularized functional.

B) The basic requirement for any regularization is that the terms of the effective action which in a certain order of perturbation theory are finite in the original theory, should recover their exact finite values after the removal of regulating mass parameters (see Ref.[11]). In particular, the regularized partition function must converge to the original one on a formal level. Here we are going to demonstrate that in the limit \( M \to \infty, M_j \to \infty \) the PV fields of the functional (16) decouple from the physical fields and contribute only to local counterterms. Indeed, rescaling the fields in the expression (16)

\[
\begin{align*}
B^{tr}_\mu &\to \frac{1}{\sqrt{M}} B^{tr}_\mu \\
B^{tr}_\mu &\to \frac{1}{\sqrt{M}} B^{tr}_\mu
\end{align*}
\]  

(26)

we get:

\[
\lim_{M \to \infty, M_j \to \infty} I_{PV}(A) = \int \exp \left\{ i \int B^{tr}_\mu(x) B^{tr}_\mu(x) dx \right\} \delta(\nabla \mu B^{tr}_\mu) \delta(\partial_\mu B^{tr}_\mu) \prod_x DB^{tr}_\mu DB^{tr}_\mu \times \det^{-1} \mathcal{M}(A)
\]  

(27)
Making again the change of variables (18) for both fields $B_\mu, B_\rho$ (the Jacobian of this transformation is equal to $det^{-1}\nabla^2$) we see that:

$$\lim_{M \to \infty; M_j \to \infty} I_{PV}(A) = \int \exp\left\{ i \int \left[ (\overline{B}_\mu^{tr} + \nabla_\mu \overline{\psi})^a (B_\mu^{tr} + \nabla_\mu \psi)^a \right] dx \right\} \times$$

$$\times \delta(\nabla^2 \overline{\psi}) \delta(\partial_\mu B_\rho^{tr} + M_\psi) \prod_x D\overline{B}_\mu^{tr} DB_\mu^{tr} D\psi D\psi det^{-1} \mathcal{M}(A) det^{-1}\nabla^2 =$$

$$= \int \exp\left\{ i \int \overline{B}_\mu^{tr}(x) B_\mu^{tr}(x) dx \right\} \prod_x D\overline{B}_\mu^{tr} DB_\mu^{tr} = const$$

(28)

So the PV interaction (16) completely decouples from the physical fields in the limit $M \to \infty, M_j \to \infty$.

C) Now we shall show that the PV functional (16) exactly compensates the remaining one-loop divergencies of the functional (12). One can see that due to the gauge invariance of $S_0$ the nonlocal term in the exponent of the expression (16) can be rewritten in the form:

$$\int \int \overline{B}_\rho^a(x) \frac{\delta^2 S_0}{\delta A_\rho^a(x) \delta A_\mu^b(y)} (\nabla_\mu \nabla^{-2} \nabla_\nu B_\nu^b) dx dy = \int \overline{B}_\rho^a(x) \frac{\delta S_0}{\delta A_\rho^a(x)} \rho g_{abc} (\nabla^{-2} \nabla_\nu B_\nu^c) dx$$

(29)

and generates only convergent diagrams, as some derivatives in the r.h.s. of Eq.(29) act on external fields $A_\mu$. Therefore the divergent diagrams generated by Eqs.(15) and (16) have the same structure (if $\overline{B}_\mu$ corresponds to $h_\mu$ and $B_\rho$ corresponds to $q_\rho$) and their sum is finite. This fact proves that the PV functional (16) compensates the remaining one-loop divergencies of the functional (12).

D) At this point we shall show that when integrated over the fields $A_\mu$ the PV functional (16) does not produce any new divergent subgraphs in multiloop diagrams and the generating functional (12) regularized by the PV interaction (16) is finite not only at one-loop level, but also for any number of loops. Let us write this regularized functional, which is the final output of our paper:

$$Z[J] = \int \exp\left\{ i \int \left[ \mathcal{L}_{YM}(x) + J_\rho^a(x) A_\rho^a(x) - \frac{1}{2} (h_\rho^a(x))^2 - \frac{1}{2\Lambda^3} h_\rho^a(x) \frac{\delta S_0}{\delta A_\rho^a(x)} \right] dx +$$

$$+ \frac{i}{2\Lambda^3} \int \overline{B}_\rho^a(x) \frac{\delta^2 S_0}{\delta A_\rho^a(x) \delta A_\mu^b(y)} (B_\mu - \nabla_\mu \nabla^{-2} \nabla_\nu B_\nu^b) dx dy + iM \int \overline{B}_\rho^a(x) B_\mu^a(x) dx \right\} \times$$

$$\times \delta(\nabla_\mu h_\mu) \delta(\nabla_\mu \overline{B}_\mu) \delta(\partial_\mu A_\rho) \delta(\partial_\rho B_\mu) det^{\frac{1}{2}} \nabla^2 \prod_j det^{\frac{1}{2}} (\nabla^2 - M_j^2) \prod_x D\overline{B}_\mu DB_\rho Dh_\mu DA_\rho$$

(30)

One can see that the divergency index of arbitrary diagram generated by the functional (30) is much the same as in expression (13) except for the contribution of external lines of PV fields:

$$\omega_2 = \omega_1 - E_B - E_{\overline{B}} - 3E_{gh}^{PV} - \frac{7}{2} E_{\delta}^{PV}$$

(31)

where $E_B = E_{\overline{B}}$ are the numbers of external lines of transversal parts of the fields $B$ and $\overline{B}$ correspondingly, $E_{gh}^{PV}$ is the number of external lines of the fields representing...
\[ \det \mathcal{F} = \frac{1}{2} \left( \nabla^2 - M_\lambda^2 \right), \quad E_{PV} = E_{PV}^{long} + E_{PV}^{\lambda} \]

is the number of external lines representing longitudinal part of field \( \overline{B} \) and auxiliary field \( \lambda^{PV} \) arising in the representation of delta-function:

\[ \delta(\nabla_\mu \overline{B}_\mu) = \int \exp \left\{ \int \lambda^{PV} \nabla_\mu \overline{B}_\mu dx \right\} \prod_x \lambda^{PV} \]  

(\( E_{PV}^{\lambda} \) and \( E_{PV}^{\delta} \) are even numbers). One can see that if \( E_{B} > 0 \), the diagram is convergent, as it includes at least one internal line of \( A \) field. So for the diagram to be divergent the following conditions must hold: \( E_{B} = E_{\overline{B}} = 0 \), \( E_{PV} = E_{PV}^{\lambda} = 0 \) and the number of loops \( I = 1 \). As was shown above, all divergencies in such one-loop diagrams generated by the functional (30) are compensated because of the application of PV procedure.

The absence of new divergent diagrams generated by the PV interaction when integrated over \( A_\mu \) is due to the fact that the propagators of the fields \( A_\mu \) decrease for large momenta faster than the propagators of PV fields. Such a solution of the problem of overlapping divergencies was proposed in Ref. [5] and realized in Ref. [1]. This way of solving the problem imposes certain restrictions on the form of higher covariant derivative term in the intermediate Lagrangian (1). As was pointed out in Ref. [1], if one is interested in calculation of only one-loop diagrams, one can use a much more simple intermediate Lagrangian:

\[ \mathcal{L}_{YM}(x) \rightarrow \mathcal{L}^\Lambda(x) = \mathcal{L}_{YM}(x) + \frac{1}{8\Lambda^4} \text{tr}(\nabla^2 F_{\mu\nu})^2 \] 

(Qua the examples of such calculations see Refs. [6, 9].) However starting from Lagrangian (33), one can not introduce the ”correct” PV interaction which does not generate additional divergent subgraphs in multiloop diagrams (overlapping divergencies) when integrated over \( A_\mu \) (for details see Ref. [1]). So if one is interested in generating functional which is finite in any order of perturbation theory and also has nonperturbative meaning, one should use the functional (30).

Let us make a conclusion to this section. We constructed regularized generating functional (30) which is ”correct” in a sense that:

A) It satisfies the correct Slavnov-Taylor identities (see Refs. [12, 13]), as we used only gauge invariant regularizing terms.

B) In the limit \( M \rightarrow \infty \); \( M_j \rightarrow \infty \) and then \( \Lambda \rightarrow \infty \) all unphysical excitations decouple from the physical fields and contribute only to local counterterms.

C and D) All the diagrams generated by the functional (30) are finite as the integral over anticommuting fields \( \overline{B}, B \) subtracts the divergent one-loop diagrams which arise due to integration over \( A, h \) and no divergent subgraphs in multiloop diagrams (overlapping divergencies) are present.

At the same time the functional (30) is simpler that the one obtained in Ref. [1] and requires much less computational effort, so it can be more attractive for using in practical calculations.

Finally we notice that the nonlocal term (29) in the expression (30) does not contribute in the limit \( M \rightarrow \infty \). For finite \( M \) we have shown it produces finite diagrams, and in the limit \( M \rightarrow \infty \) its contribution disappears. Being interested finally in the limit \( M \rightarrow \infty \) we can omit this term in the Eq. (30). It simplifies the final expression for the regularized functional and make the effective action local. Omitting this term we break the gauge invariance for finite \( M \). The Slavnov-Taylor identities will be violated by finite terms of order \( O(M^{-1}) \). These terms are harmless as they have no influence on the counterterms and disappear in the limit \( M \rightarrow \infty \).
3 Discussion.

In this paper we constructed a consistent invariant regularization of gauge models within the framework of Higher Covariant Derivative method. The obtained formulation is considerably simpler and requires much less computational effort than the one in Ref.[1]. This may make the regularization of gauge theories by higher covariant derivatives and gauge invariant Pauli-Villars regulators much more attractive for using in practical calculations in models where dimensional regularization is not applicable.

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