The Yang-Mills theory as a massless limit of a massive nonabelian gauge model.

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Abstract
A gauge invariant infrared regularization of the Yang-Mills theory applicable beyond perturbation theory is constructed.

1 Introduction
It is well known that the Yang-Mills theory cannot be obtained as a limit, when $m \to 0$, of the massive Yang-Mills model. The longitudinal polarization of the massive vector field does not decouple in the zero mass limit and gives a nonzero contribution [1]. It was shown recently [2], that a massive nonabelian gauge model, which includes additional excitations, may be used as a gauge invariant infrared regularization of the Yang-Mills theory. However the procedure proposed in [2] may be used only in the framework of perturbation theory with respect to the coupling constant. On the other hand in the presence of infrared singularities a perturbation theory is not applicable, although the regularization, proposed in [2] is still useful for calculation of Green functions of gauge invariant operators.

In the present paper we propose a gauge invariant nonabelian model which may be used as an infrared regularization of the Yang-Mills theory both in the framework of the perturbation theory and beyond it.

2 The massive nonabelian gauge invariant model.
It was discussed in the papers ([3], [4],[5]) that impossibility to select a unique gauge beyond the perturbation theory is not the intrinsic property of the Yang-Mills model, but is related to its particular formulation. Adding new excitations which decouple asymptotically it is possible to quantize nonabelian gauge models in a manifestly Lorentz invariant way both in perturbation theory and beyond it.

Having that in mind we propose to use for the gauge invariant infrared regularization of the Yang-Mills theory the following Lagrangian

$$L = -\frac{1}{4} F_{\mu\nu}^a F_{\mu\nu}^a - m^{-2} (D^2 \phi)^*(D^2 \phi) + (D_\mu \phi)^*(D_\mu \phi) + (D_\mu b)^*(D_\mu b)$$

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\[ + \alpha^2 (D_\mu \tilde{\phi})^* (D_\mu \tilde{\phi}) - \alpha^2 m^2 (b^* e + e^* b) \]  

(1)

where \( \phi \) is a two component complex doublet, and

\[ \tilde{\phi} = \phi - \hat{\mu}; \quad \hat{\mu} = (0, \mu \sqrt{2} g^{-1}) \]

(2)

\( \mu \) is an arbitrary constant. \( D_\mu \) denotes the usual covariant derivative. To save the place we consider here the group \( SU(2) \).

In the following we shall use the parametrization of \( \phi \) in terms of Hermitean components

\[ \phi = \left( \frac{i \phi^1 + \phi^2}{\sqrt{2}} (1 + \frac{g}{2 \mu} \phi^0), \phi^0 - i \frac{\phi^3}{\sqrt{2}} (1 + \frac{g}{2 \mu} \phi^0) \right) \]

(3)

The complex anticommuting scalar fields \( b, e \) will be parametrized as follows

\[ b = \left( \frac{ib^1 + b^2}{\sqrt{2}}, \frac{b^0 - ib^3}{\sqrt{2}} \right) (1 + \frac{g}{2 \mu} \phi^0) \]

\[ e = \left( \frac{ie^1 + e^2}{\sqrt{2}}, \frac{e^3}{\sqrt{2}} \right) \]

(4)

where the components \( e^\alpha \) are Hermitean, and \( b^\alpha \) are antihermitean. This particular parametrization of the classical fields is used as we want to get rid off the ambiguity in choosing the gauge for quantization completely.

The Lagrangian (1) is obviously invariant with respect to "shifted" gauge transformations

\[ A_\mu^a \to A_\mu^a + \partial_\mu \eta^a - g \epsilon^{abc} A_\mu^b \eta^c \]

\[ \phi^a \to \phi^a + \frac{g}{2} \epsilon^{abc} \phi^b \eta^c + \eta^a \mu + \frac{g^2}{4 \mu} \phi^a \phi^b \eta^b \]

\[ \phi^0 \to \phi^0 - \frac{g}{2} (\phi^a \eta^a)(1 + \frac{g}{2 \mu} \phi^0) \]

\[ b^a \to b^a + \frac{g}{2} \epsilon^{abc} b^b \eta^c + \frac{g}{2} b^a \eta^a + \frac{g^2}{4 \mu} b^a \phi^b \eta^b \]

\[ e^a \to e^a + \frac{g}{2} \epsilon^{abc} e^b \eta^c + \frac{g}{2} e^a \eta^a \]

\[ b^0 \to b^0 - \frac{g}{2} b^a \eta^a + \frac{g^2}{4 \mu} (\phi^a \eta^a) b^0 \]

\[ e^0 \to e^0 - \frac{g}{2} e^a \eta^a \].

(5)

This Lagrangian is also invariant with respect to the supersymmetry transformations

\[ \phi \to \phi - b \epsilon \]

\[ e \to e - \frac{D^2 (\phi - \hat{\mu})}{m^2} \epsilon \]

\[ b \to b \]

(6)

where \( \epsilon \) is a constant Hermitean anticommuting parameter. This symmetry plays a crucial role in the proof of decoupling of unphysical excitations. It holds for any \( \alpha \), but for \( \alpha = 0 \) these transformations are also nilpotent. Note that for further discussion we need only the existence of the conserved charge \( Q \) and nilpotency of the asymptotic charge \( Q_0 \), as
the physical spectrum is determined by the asymptotic dynamics. (We use the standard
scattering problem, assuming that the asymptotical spectrum coincides with the spectrum
of the free Hamiltonian). In the case under consideration the nilpotency of the asymptotic
charge requires $\alpha = 0$, and the massive theory with $\alpha \neq 0$ is gauge invariant but not
unitary. It may seem strange as usually the gauge invariance is a sufficient condition of
unitarity, because one can pass freely from a renormalizable gauge to the unitary one,
where the spectrum includes only physical excitations. In the present case there is no
"unitary" gauge. Even in the gauge $\phi^a = 0$, there are unphysical excitations.

For gauge transformations (5) the gauge $\phi^a = 0$ is admissible both in perturbation
theory and beyond it. Indeed, if $\phi^a = 0$, then under the gauge transformations (5) the
variables $\phi^a$ become

$$\delta \phi^a = \mu \eta^a$$

and the condition $\phi^a = 0$ implies that $\eta^a = 0$. In the gauge $\phi^a = 0$ the variables $\phi^a$ are
parameterized as usual

$$\phi = (0, \phi^0 \sqrt{2}) + \hat{\mu}$$

It is also obvious that for $\alpha \neq 0$ the Lagrangian (1) describes a massive vector field and
does not produce infrared singularities.

In terms of shifted variables the Lagrangian (1) looks as follows

$$L = -\frac{1}{4} F_{\mu \nu}^a F^{a}_{\mu \nu} - m^{-2} (D^2 \phi)^* (D^2 \phi) + m^{-2} (D^2 \phi)^* (D^2 \hat{\mu})$$
$$+ m^{-2} (D^2 \hat{\mu})^* (D^2 \phi) - m^{-2} (D^2 \hat{\mu})^* (D^2 \hat{\mu}) + (D_\mu \hat{\epsilon})^* (D_\mu \hat{\epsilon})$$
$$+ (D_\mu \hat{\epsilon})^* (D_\mu \phi) + \alpha^2 (D_\mu \phi)^* (D_\mu \phi) - \alpha^2 (D_\mu \phi)^* (D_\mu \hat{\mu})$$
$$- \alpha^2 (D_\mu \hat{\mu})^* (D_\mu \phi) + \alpha^2 (D_\mu \hat{\mu})^* (D_\mu \hat{\mu}) - \alpha^2 m^2 (b^* e + e^* b)$$ (9)

and the term

$$\alpha^2 (D_\mu \hat{\mu})^* (D_\mu \hat{\mu}) = \frac{\alpha^2 \mu^2}{2} A^2_\mu$$ (10)

produces the mass for the vector field.

The term

$$m^{-2} (D^2 \hat{\mu})^* (D^2 \hat{\mu}) = \frac{\mu^2}{2m^2} [ (\partial_\mu A_\mu)^2 + \frac{g^2}{2} (A^2)^2 ]$$ (11)

makes the theory renormalizable for any $\alpha$. To avoid complications due to the presence
of the Yang-Mills dipole ghosts at $\alpha = 0$ we put $\mu^2 = m^2$. The effective Lagrangian in the
gauge $\phi^a = 0$ may be written as follows

$$L_{\text{ef}} = L + \lambda^a \phi^a - \mu \bar{c}^a c^a$$ (12)

where $L$ is the Lagrangian (11) and $\bar{c}^a, c^a$ are nondynamical ghost fields.

Invariance of the Lagrangian (11) with respect to the gauge transformation (5) and the
supersymmetry transformations (6) makes the effective Lagrangian invariant with respect
to the simultaneous BRST transformations corresponding to (5) and the supersymmetry
transformations (9). If $s_1$ is the nilpotent operator, corresponding to the simultaneous
BRST and supersymmetry transformations and the transformations of the fields $\lambda, \bar{c}, c$
are

$$s_1 \lambda^a = 0; \quad s_1 c^a = -\frac{g}{2} \epsilon_{abc} \bar{c}^b c^c; \quad s_1 \bar{c}^a = \lambda^a$$ (13)
the effective Lagrangian may be also written in the form

\[ L_{\text{ef}} = L + s_1 \bar{c}^a \phi^a = L(x) + \lambda^a \phi^a - \bar{c}^a (\mu c^a - b^a) \]  

(14)

As it was indicated in the paper [5], one can integrate over \( \bar{c}, c \) in the path integral determining expectation value of any operator corresponding to observable. It leads to the change \( c^a = b^a \mu^{-1} \). After such integration the effective Lagrangian becomes invariant with respect to the transformations which are the sum of the BRST transformations and the supersymmetry transformations (6) with \( c^a = b^a \mu^{-1} \):

\[
\begin{align*}
\delta A^a_\mu & = D_\mu b^a \mu^{-1} \epsilon \\
\delta \phi^a & = 0 \\
\delta \phi^0 & = -b^0 (1 + \frac{g}{2 \mu}) \epsilon \\
\delta e^a & = \left( \frac{g}{2 \mu} \varepsilon^{abc} b^c + \frac{ge^0 b^a}{2 \mu} + i \frac{D^2 (\phi)^a}{\mu^2} \right) \epsilon \\
\delta e^0 & = -\frac{ge^a b^a}{2 \mu} - \frac{D^2 (\phi^0)}{\mu^2} \epsilon \\
\delta b^a & = \frac{g}{2 \mu} \varepsilon^{abc} b^c \\
\delta b^0 & = 0
\end{align*}
\]  

(15)

As the transformation (15) preserves the condition \( \phi^a = 0 \), we omitted in this transformation all the terms proportional to \( \phi^a \). For the asymptotic Hamiltonian these transformations acquire the form

\[
\begin{align*}
\delta A^a_\mu & = \partial_\mu b^a \mu^{-1} \epsilon \\
\delta \phi^a & = 0 \\
\delta \phi^0 & = -b^0 \epsilon \\
\delta e^a & = \partial_\mu A^a_\nu \mu^{-1} \\
\delta e^0 & = -\partial^2 \phi^0 \mu^{-2} \\
\delta b^a & = 0 \\
\delta b^0 & = 0
\end{align*}
\]  

(16)

According to the Neuther theorem the invariance with respect to the transformations (15) generates a conserved charge \( Q \), and the physical asymptotic states may be chosen to satisfy the equation

\[ \hat{Q}\psi |_{\text{as}} = 0 \]  

(17)

where

\[
\begin{aligned}
Q_0 = \int d^3 x [ & (\partial_0 A^a_i - \partial_i A^a_0) \mu^{-1} \partial_\mu b^a - \mu^{-1} \partial_\nu A^a_\nu \partial_0 b^a + \mu^{-2} \partial^2 (\partial_0 \phi^0) b^0 - \mu^{-2} \partial_0 b^0 \partial^2 (\phi^0) \\
& - \mu a^2 b^a A^a_0 ]
\end{aligned}
\]  

(18)

Due to the conservation of the Neuther charge this condition is invariant with respect to dynamics.
Using Ostrogradsky canonical variables for higher derivative systems one can rewrite the eq. (18) in the form

\[ Q_0 = \int d^3x \left[ -(\partial_i p_i^a b^a \mu^{-1} + \mu a^2 b^a A_0^a) + \mu^{-1} p_0^a \partial_0 b - (p_1 - \alpha^2 \varphi_2^0) b + p_2 \partial_0 b \right] \quad (19) \]

Here

\[ H = p_i^a \dot{A}_i^a + p_0^b b + p_e \dot{e} + p_1 \varphi_2 + p_2 \partial_0 \varphi_2 - L \]
\[ p_0^a = \dot{A}_0^a - \partial_0 A_i^a; \quad p_0^b = -\partial_0 A_i^b; \quad p_0^e = \dot{e}^a; \quad p_0^b = \dot{b}^a \]
\[ \varphi_1^0; \quad \varphi_2 = \dot{\varphi}_0^0; \quad p_2 = -\mu^{-2} \partial^2 \varphi_0^0; \]
\[ p_1 = \mu^{-2} \partial^2 \varphi_0^0 + \alpha^2 \dot{\varphi}_0^0; \quad b^0 = b, e^0 = e, p_b = e, p_e = b \quad (20) \]

In these notations the free Hamiltonian looks as follows

\[ H_0 = \frac{p_i^2}{2} - \frac{p_0^2}{2} - \partial_i p_i^a A_0^a + p_0^b \partial_i A_i^b + \frac{1}{4}(F_{ij}^a)^2 + \frac{1}{2} b^a \partial_i e^a + \]
\[ + p_b p_e + \partial_i b \partial_i e + p_1 \varphi_2 - \frac{\mu^2}{2} p_2^2 + p_2 \Delta \varphi_1 \]
\[ - \frac{\alpha^2 \mu^2}{2} A_0^2 + \frac{\alpha^2 \mu^2}{2} A_i^2 + \alpha^2 \mu^2 b e - \frac{\alpha^2}{2} \varphi_0^2 + \frac{\alpha^2}{2} \partial_i \varphi_1 \partial_i \phi_1 \quad (21) \]

We want to prove that the Lagrangian (14) really describes the infrared regularization of the Yang-Mills theory. That means for \( \alpha \neq 0 \) it corresponds to a massive gauge invariant theory and in the limit \( \alpha = 0 \) it describes the usual three dimensionally transversal excitations of the Yang-Mills field. Of course for \( \alpha \neq 0 \) the spectrum includes also some unphysical excitations.

In the limit \( \alpha = 0 \) only the first two lines of the eq. (21) survive. They contain the terms depending only on the fields \( A_0, A_i \) and corresponding canonical momenta which coincide with the usual Yang-Mills Hamiltonian in the diagonal Feynman gauge and the fields \( \phi_0, b_0, e_0 \). The fields \( b_a, e_a \) play the role of the Faddeev-Popov ghosts. By the usual arguments the longitudinal and temporal components of the Yang-Mills field as well as the fields \( b_a, e_a \) decouple, and the physical states may include only transversal components of the Yang-Mills field and variables corresponding to the fields \( \phi_0^0, b_0^0, e_0^0 \). Below we shall show that the fields \( \phi_0^0, b_0^0, e_0^0 \) also decouple.

It follows from the eqs. (20) that the asymptotical momenta \( p_1, p_2 \) satisfy the free field equations

\[ \partial^2 p_{1,2} = 0; \quad (22) \]

Therefore these momenta allow the standard expansion

\[ p_{1,2}(x) = (2\pi)^{-\frac{3}{2}} \int d^3k \frac{i\sqrt{\omega}}{2} (a_{p_{1,2}}^+(k) \exp\{ikx\} - a_{p_{1,2}}^-(k) \exp\{-ikx\}) \quad (23) \]

\[ k_0 = \omega = \sqrt{k^2} \]

A similar expansion may be written for the asymptotic fields \( b, e, \dot{b}, \dot{e} \)

\[ b(e)(x) = (2\pi)^{-\frac{3}{2}} \int d^3k \frac{1}{\sqrt{2\omega}} (b(e)^+(k) \exp\{ikx\} + b(e)^-(k) \exp\{-ikx\}) \]

\[ b(e)(x) = (2\pi)^{-\frac{3}{2}} \int d^3k \frac{1}{\sqrt{2\omega}} (b(e)^+(k) \exp\{ikx\} + b(e)^-(k) \exp\{-ikx\}) \]
\[ \dot{b}(\dot{e})(x) = (2\pi)^{-\frac{3}{2}} \int d^3k \frac{1}{\sqrt{2\omega}} (b(e)^+(k) \exp\{ikx\} - b(e)^-(k) \exp\{-ikx\}) \]

\[ k_0 = \omega = \sqrt{k^2} \]  

(24)

The canonical variables describing the fields \( \phi_0 \) in general are not oscillatory and cannot be presented by the operators in the usual Fock space.

The asymptotic fields \( \varphi_{1,2}, p_{1,2} \) satisfy the equations of motion which follow from the Hamiltonian \( (21) \):

\[ \dot{\varphi}_1 = \varphi_2; \quad \dot{\varphi}_2 + \mu^2 p_2 - \Delta \varphi_1 = 0; \quad -\dot{p}_2 = p_1; \quad -\dot{p}_1 = \Delta p_2 \]  

(25)

In accordance with these equations the part of the asymptotic BRST charge depending on variables \( p_{1,2} \)

\[ \tilde{Q}_0 = \int d^3x (p_2 \dot{b} + p_1 b) \]  

(26)

is conserved and does not depend on time. Therefore one can put \( x_0 \) in the variables \( p_{1,2}, b, \dot{b} \) equal to zero.

Using the equations \( (23, 25) \) we get

\[ p_1(x, 0) = (2\pi)^{-\frac{3}{2}} \int d^3k \frac{1}{\sqrt{2\omega}} (a_1^+(k) \exp\{-ikx\} - a_1^-(k) \exp\{ikx\}) \]

\[ p_2(x, 0) = -(2\pi)^{-\frac{3}{2}} \int d^3k \frac{1}{2\sqrt{\omega}} (a_1^+(k) \exp\{-ikx\} + a_1^-(k) \exp\{ikx\}) \]  

(27)

\[ b(e)(x, 0) = (2\pi)^{-\frac{3}{2}} \int d^3k \frac{1}{\sqrt{2\omega}} (b(e)^+(k, 0) \exp\{-ikx\} + b(e)^-(k) \exp\{ikx\}) \]

\[ \dot{b}(e)(x, 0) = i(2\pi)^{-\frac{3}{2}} \int d^3k \frac{\sqrt{\omega}}{\sqrt{2^3}} (b(e)^+(k) \exp\{-ikx\} - b(e)^-(k) \exp\{ikx\}) \]  

(28)

Let us introduce the following combinations of the operators \( \dot{p}_{1,2}, \dot{\varphi}_{1,2}, \dot{b}, \dot{e} \), which satisfy the commutation relations of creation and annihilation operators

\[ \frac{\hat{p}_2(k)\omega(k) + i\hat{p}_1(k)}{\sqrt{2\omega}} = -\hat{a}^+(k); \quad \frac{\hat{\varphi}_1(k)\omega(k) + i\hat{\varphi}_2(k)}{\sqrt{2\omega}} = \hat{a}^-(k) \]

\[ \frac{\hat{p}_2(k)\omega(k) - i\hat{p}_1(k)}{\sqrt{2\omega}} = -\hat{a}^-(k); \quad \frac{\hat{\varphi}_1(k)\omega(k) - i\hat{\varphi}_2(k)}{\sqrt{2\omega}} = \hat{a}^+(k) \]

\[ \frac{\hat{b}(k)\omega(k) + i\hat{b}(k)}{\sqrt{2\omega}} = -\hat{b}^+(k); \quad \frac{\hat{e}(k)\omega(k) + i\hat{e}(k)}{\sqrt{2\omega}} = \hat{e}^-(k) \]

\[ \frac{\hat{b}(k)\omega(k) - i\hat{b}(k)}{\sqrt{2\omega}} = -\hat{b}^-(k); \quad \frac{\hat{e}(k)\omega(k) - i\hat{e}(k)}{\sqrt{2\omega}} = \hat{e}^+(k) \]  

(29)

Note that noncommuting pairs are formed by the operators \( \hat{a}, \hat{\varphi} \) and \( \hat{b}, \hat{e} \).

For oscillatory variables \( p_{1,2}, b, e \) this definition coincides with the standard one.

In terms of these operators the asymptotic BRST charge may be written as follows

\[ Q_0 = i \int d^3k (\hat{a}^+(k)\hat{b}^-(k) - \hat{b}^+(k)\hat{a}^-(k)) \]  

(30)
I wish to emphasize that the states generated by the operators $\hat{a}^+$ do not belong to the Fock space. However this fact is irrelevant for the physical interpretation as the part of the Hamiltonian $\hat{H}$, which depends on the variables $b, e, \varphi_1, \varphi_2$ and conjugated momenta, is BRST exact: it may be presented as the anticommutator of the BRST charge $\hat{Q}_0$ with some operator $\hat{A}$

$$\hat{H}_0 = [\hat{Q}_0, \hat{A}]_+; \quad \hat{A} = \int d^3x(\dot{\varphi}_2 \dot{e} - \frac{\mu^2}{2} \hat{p}_2 \dot{e} + \Delta \hat{\varphi}_1 \dot{e})$$  

(31)

It follows that the part of the Hamiltonian $\hat{H}$, which depends on $\varphi_{1,2}, p_{1,2}, b, e$ does not contribute to the expectation value calculated with the help of the physical states, annihilated by $Q_0$. (For similar construction see [6]). Therefore it is irrelevant for the energy of any physical state, and we can define the vacuum as the vector annihilated by the operators $\hat{a}^-, \hat{a}^+, \hat{b}^-, \hat{e}^-$. 

We may introduce the operator $\hat{K}$ by the formula

$$\hat{K} = \int d^3x(\dot{\varphi}_1(x,0)\dot{e}(x,0) - \varphi_2(x,0)\dot{e}(x,0)) = -i \int d^3k(\hat{a}^+(k)\hat{e}^-(k) - \hat{e}^+(k)\hat{a}^-(-k))$$  

(32)

The anticommutator of $\hat{Q}_0$ and $\hat{K}$ is proportional to the number of unphysical modes generated by the operators $\hat{a}^+, \hat{a}^-, \hat{b}^+, \hat{e}^+$. By the usual arguments any vector satisfying the eq.(17) can be presented in the form

$$|\psi >_{phys} = |\psi >_{tr} + \hat{Q}_0 |\chi >$$  

(33)

where $|\psi >_{tr}$ contains only excitations corresponding to the transversal modes of the Yang-Mills field, and expectation value of any observable, calculated with the help of the vectors $|\psi >_{phys}$ coincides with the expectation value calculated with the help of $|\psi >_{tr}$. 

It completes the proof of the fact that for $\alpha = 0$ the Lagrangian (14) produces the same expectation values for all observables calculated with the help of of the physical vectors, satisfying the eq.(17) as the standard Yang-Mills theory. At the same time for $\alpha \neq 0$ it describes a massive gauge invariant theory, which does not have infrared singularities.

### 3 Discussion

In this paper we showed that the Lagrangian (10) may be uniquely quantized irrespectively of using the perturbation theory with respect to the coupling constant. For any value of the coupling constant it allows a unique canonical quantization in the gauge $\phi^a = 0$. 

For $\alpha \neq 0$ it does not produce infrared singularities and is gauge invariant. Introducing some ultraviolet gauge invariant regularization, for example dimensional or higher covariant derivatives, one get the gauge invariant regularization which makes all the quantities finite.

For $\alpha = 0$ it gives for the expectation values of observable operators, calculated with the help of the physical vectors, annihilated by the BRST charge the same result as the usual Yang-Mills Lagrangian. Of course, if the calculations are performed in the framework of perturbation theory, for $\alpha = 0$ the infrared singularities reappear. However this Lagrangian may serve as a starting point for nonperturbative calculations, in particular for calculations explaining the phenomenon of quark confinement.

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