New Algebra of Local Symmetries for Regge Limit of Yang–Mills Theories

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Abstract

Local effective action is derived to describe Regge asymptotic of Yang–Mills theories. Local symmetries of the effective action originating from the gauge symmetry of the underlying Yang–Mills theory are studied. Multicomponent effective action is introduced to express the symmetry transformations as field transformations. The algebra of these symmetries is decomposed onto a semi-direct sum of commutative algebras and four copies of the gauge algebra of the underlying Yang–Mills theory. Possibility of existence of solitons corresponding to the commutative subalgebra of the symmetry algebra is mentioned.

Since the discovery of the asymptotic freedom in Yang-Mills theories [1], it is theoretically understood that one and the same quantum field theory may appear as describing the dynamics of different sets of degrees of freedom: what appears at long distances as chiral dynamics of meson fields [2] reveals the presence of quarks and gluons at short distances [3].

In this talk, we present a theoretical construction which points out the presence of a new distinctive description of the dynamics of Yang-Mills fields. Phenomenologically, this description, which we call *Effective Regge QCD*, should reveal itself in high-energy hadron scattering (for example, at LHC) when the transverse momenta of the produced particles are much smaller than the collision energy but, at the same time, much larger than the hadron mass scale of 1 GeV.

We need Effective Regge QCD in particular after the recent calculation [4], where next-to-leading energy logarithms of perturbative QCD were calculated and analyzed. The result of this analysis may be interpreted as follows:

whatever small the gauge coupling would be provided by the asymptotic freedom, colliding quarks and gluons become strongly interacting at sufficiently high collision energy because of the presence of the energy logarithms. Thus, *perturbative* QCD becomes inadequate when the collision energy becomes high enough while the transverse momentum remains fixed (these kinematic conditions are usually called the Regge limit).

At the same time, we do not expect hadrons to be the adequate degrees of freedom to describe the production of particles with transverse momenta much exceeding 1 GeV, because observation of such particles should resolve quarks and gluons inside the colliding hadrons.

We suggest to look for a formulation of Yang-Mills dynamics which would conceal the strongly interacting quarks and gluons belonging to the *different* colliding hadrons into new degrees of freedom. Each of this new "particles" should contain excitations arriving from both spatial infinities along the collision axis to participate in the collision. As such, the new degrees of freedom are not localized particle-like objects but rather the elementary scattering processes. The adequate degrees of freedom should be weakly interacting in the Regge limit.

The first step in finding the new degrees of freedom is to construct the effective theory in terms of quarks and gluons. It should be a simplified version of Yang-Mills theory. The simplification should not spoil the Regge limit of the underlying Yang-Mills theory. Various approaches to such a construction along with theoretical and phenomenological motivations see in [5, 6, 7, 8].

Here we develop further the approach of [8]. Its core is to consider excitations of the gauge field moving fast rightward along a space axis z as excitations of a separate field A^R of the effective theory. Analogously, an independent field A^L is introduced to describe excitations of the underlying theory which are fast left-movers. It was noted in [8] that there exists a finite effective action describing the dynamics of these fields in the limit of infinite energy of the relative motion of the right- and left-movers. It was also noted in [8] that the effective action admits a multiplicative infrared renormalization. Corresponding renormgroup reproduces in the case of gluodynamics the known Regge trajectory of the Reggeized gluon.

To move beyond perturbation theory in the resummation of the energy logarithms of gauge theories, it is crucial to understand the symmetry behind the multiplicative renormalizability established in [8] in the one-loop approximation. Below we construct an algebra of local symmetries of an effective Regge theory. We will show that this algebra contains Abelian subalgebras. Usually, the presence of Abelian local symmetries correspond to the presence of monopole-like classical solutions of the field equations (see, for example, [9]). These solitons of the Effective Regge QCD are our candidates for the new degrees of freedom needed to describe high-energy hadron collisions.

First, we derive a local effective action. The effective action of [8] is nonlocal and is a result of integrating out some fields of the local effective theory we introduce below. It is obtained from the underlying Yang–Mills theory in two steps. (In what follows we assume the gauge group of the underlying theory to be $SU(N)$. Generalization to arbitrary gauge group is straightforward. We also reduce our consideration to a pure Yang–Mills theory without matter fields.)

The first step connects the underlying Yang–Mills theory with an extended theory which is equivalent to the Yang–Mills theory but has an extended field content. Namely, the extended theory is a theory of four vector gauge fields—S(oft), H(ard), R(ight), and L(eft). Its action is

$$S_{ext}(\{A^\Gamma\}) = S_{YM}(\sum_\Gamma P_\Gamma A^\Gamma), \tag{1}$$

where $\{P_\Gamma\}$, $\Gamma = \{S, H, L, R\}$ is a complete set of orthogonal projectors acting on the fields. For any such a set, vacuum expectations of products of the fields $P_\Gamma A^\Gamma$ calculated with the action (1) coincide with the vacuum expectations of the corresponding products of $P_\Gamma A$, where A is the gauge field, calculated in the framework of the underlying Yang–Mills theory:

$$\langle \Pi_i P_{\Gamma_i} A^{\Gamma_i} \rangle^{S_{ext}} = \langle \Pi_i P_{\Gamma_i} A \rangle^{S_{YM}}.$$

The projectors of the particular set we need act on a Fourier transformed field $\Phi(k)$ as follows:

$$\begin{aligned} P_S\Phi(k) &= \theta(|k_+| < \mu)\theta(|k_-| < \mu)\Phi(k), \\ P_H\Phi(k) &= \theta(|k_+| > \mu)\theta(|k_-| > \mu)\Phi(k), \\ P_R\Phi(k) &= \theta(|k_+| > \mu)\theta(|k_-| < \mu)\Phi(k), \\ P_L\Phi(k) &= \theta(|k_+| < \mu)\theta(|k_-| > \mu)\Phi(k), \end{aligned} \tag{2}$$

where $\theta(condition)$ equals to unit if *condition* is fulfilled and vanishes otherwise. μ above is an arbitrary scale which will not enter any physical prediction by renormgroup invariance, and $k_\pm$ are the light-cone components of the momentum, $k_\pm = k_0 \pm k_z$. Now the labels Γ make sense: S(oft) field describes the gluon modes of both light cone components of momentum small, etc.

The last, second step in the derivation of the effective action is to drop out some vertexes of the extended action (1) along with projectors. This is justified if one considers the large (= exceeding scale μ) components of momenta as very large and the small (= less than μ) components as very small. Obviously, the projectors are equal to unit and some vertices are forbidden by momentum conservation for such momenta. The resulting action, $S_{eff}(A^\Gamma)$, allows one to reproduce correlators of the underlying Yang–Mills theory in the Regge limit:

$$\lim_{\lambda\to\infty} \langle \Pi_i (R_{\Gamma_i}(\lambda) A)\rangle^{S_{YM}} = \lim_{\lambda\to\infty} \langle \Pi_i (R_{\Gamma_i}(\lambda) A^{\Gamma_i})\rangle^{S_{eff}}, \tag{3}$$

where the Regge limit is realized by the rescaling operators R_Γ ($\Gamma = S$, R, L, H). For example, left rescaling $R_L(\lambda)$ is just a Lorentz boost making $(-)$-components of momenta large. It acts on vector fields as follows:

$$\begin{aligned}
R_L(\lambda) A_+(x^+, x^-, x^i) &= \lambda A_+(\lambda x^+, \frac{x^-}{\lambda}, x^i),\\
R_L(\lambda) A_-(x^+, x^-, x^i) &= \frac{1}{\lambda} A_-(\lambda x^+, \frac{x^-}{\lambda}, x^i),\\
R_L(\lambda) A_i(x^+, x^-, x^i) &= A_i(\lambda x^+, \frac{x^-}{\lambda}, x^i).
\end{aligned} \tag{4}$$

Here and below the Lorentz indices denoted by small Latin characters label the transverse components of Lorentz tensors: $i = 1, 2$. λ scales as invariant energy of the relative motion of right- and left-movers; $x^\pm = \frac{x^0 \pm x^3}{2}$ are longitudinal components of the space-time coordinates (the normalization implies that $kx = k_+x^+ + k_-x^- + k_ix^i$). The right rescaling, $R_R(\lambda)$, acts as a boost in the opposite direction: $R_R(\lambda) = R_L(\frac{1}{\lambda})$. The remaining rescalings, H and S, should make both longitudinal momentum components large and small respectively. The scales for the field magnitudes are adjusted to ensure the finiteness of the classical equations for the rescaled fields of the effective theory in the limit $\lambda \to \infty$. They are as follows:

$$\begin{aligned}
&R_H(\lambda) A_\pm(x^\pm, x^i) = \frac{1}{\lambda} A_\pm(\frac{x^\pm}{\lambda}, x^i), \quad R_H(\lambda) A_i(x^\pm, x^i) = A_i(\tfrac{x^\pm}{\lambda}, x^i),\\
&R_S(\lambda) A_\pm(x^\pm, x^i) = \lambda A_\pm(\lambda x^\pm, x^i), \quad R_S(\lambda) A_i(x^\pm, x^i) = \lambda^2 A_i(\lambda x^\pm, x^i).
\end{aligned} \tag{5}$$

To conclude the description of the relation between the effective theory and the underlying Yang–Mills theory, we note that one should not forget altogether the cutoffs introduced by projectors in (2) and use them to define functional integral implicit in the right-hand-side of equation (3). We note here that it is the usual practice to ignore cutoffs and to consider various effective theories as the local ones until some divergence does not enforce the opposite.

Now we wright down S_{eff} explicitly. Recall that it is obtained by substituting $A = A^S + A^R + A^L + A^H$ (projectors ignored) into $S_{YM}(A)$ and by dropping out vertexes forbidden by momentum conservation. To make this dropping out, mark each monomial of fields entering $S_{YM}(\sum_\Gamma A^\Gamma)$ by two indices, (ω_L, ω_R). The indices are taking the following values: $\omega_{L,R} = 0, 1, 2$ depending on the number of legs injecting large left or, correspondingly, right $(=(-)$ or $(+))$ momentum component in the vertex. For example, vertex $A^H A^S$ mark by $(\omega_L = 1, \omega_R = 1)$. If the number of legs injecting large momentum exceeds two, put the value of corresponding index to 2 (2 will stand for "more than one leg"). Dropping out vertexes forbidden by momentum conservation amounts to dropping out the vertexes marked by unit value of any index (large momentum is injected along a single leg and has no way out). Thus, S_{eff} is a sum of four components each of which is a collection of vertexes marked by the indices of the same value:

$$S_{eff}(\{A^\Gamma\}) = \sum_{\omega_L=0,2} \sum_{\omega_R=0,2} S_{(\omega_R,\omega_L)}(\{A^\Gamma\}). \tag{6}$$

In what follows we denote $S_{(0,0)}$ as $S_{S(oft)}$. It has no large longitudinal momentum components entering its vertexes and depends on A^S only. Analogously, $S_{(2,0)} = S_{L(eft)}$ (depends on A^L and A^S only), $S_{(0,2)} = S_{R(ight)}$ (depends on A^R and A^S only), and $S_{(2,2)} = S_{H(ard)}$ since it involves large values of both longitudinal momentum components and depends on all the fields. The explicit expressions for the components of the action are as follows:

$$\begin{aligned} S_H(\{A^\Gamma\}) = {} & S_{YM}(\sum_\Gamma A^\Gamma) + S_{YM}(A^S) - \\ & S_{YM}(A^R + A^S) - S_{YM}(A^L + A^S) + \\ & D^\nu(A^R + A^S)(A^H + A^L)^\mu F_{\nu\mu}(A^R + A^S) + \\ & D^\nu(A^L + A^S)(A^H + A^R)^\mu F_{\nu\mu}(A^L + A^S) - \\ & ([A^{L\nu}, A^{R\mu}] + D^\nu(A^S)(A^{L\mu} + A^{R\mu} + A^{H\mu}))F_{\nu\mu}(A^S) - \end{aligned}$$

$$
\begin{aligned}
&\frac{1}{2}(D^{\mu}(A^S)A^{L\nu} - D^{\nu}(A^S)A^{L\mu})(D_{\mu}(A^S)A^R_{\nu} - D_{\nu}(A^S)A^R_{\mu}),\\
S_L(A^L, A^S) &= S_{YM}(A^L + A^S) - S_{YM}(A^S) + D^{\nu}(A^S)A^{L\mu}F_{\nu\mu}(A^S),\\
S_R(A^R, A^S) &= S_{YM}(A^R + A^S) - S_{YM}(A^S) + D^{\nu}(A^S)A^{R\mu}F_{\nu\mu}(A^S),\\
S_S(A^S) &= S_{YM}(A^S).
\end{aligned} \tag{7}
$$

Here $S_{YM}(A) = \frac{1}{2}F_{\mu\nu}(A)F^{\mu\nu}(A)$ (with the trace over color indices and integration over space-time implicit), ν and μ are Lorentz indices, $F_{\mu\nu}(A) = \partial_\mu A_\nu - \partial_\nu A_\mu - [A_\mu, A_\nu]$, and $D^\nu(A)\Phi = \partial^\nu\Phi - [A^\nu, \Phi]$ (it is assumed that D's act only on their nearest right neighbors in the above formula).

Now we turn to derivation of local symmetries of effective action $S_{eff} = S_H + S_L + S_R + S_S$. To this end note that gauge symmetry of S_{YM} states that some local functional vanishes. It is obtained by an infinitesimal gauge transformation of the action and depends on gauge field A and on field α parameterizing the transformation. In searching the form this symmetry takes after replacing the Yang–Mills action by the effective one, it is natural to substitute all the fields entering the functional which vanishes due to the gauge invariance by the sum of four independent fields corresponding to disparate values of longitudinal momenta and drop out the vertexes forbidden by momentum conservation. In particular, the single field α parameterizing the gauge transformation of the Yang–Mills action is to be substituted by the sum: $\alpha = \alpha^H + \alpha^L + \alpha^R + \alpha^S$. In this way one obtains that some functional of fields A^Γ, α^Γ vanishes: $V(\{A^\Gamma, \alpha^\Gamma\}) = 0$. Consequently, it looks plausible that there are at least four independent local symmetries of the effective action corresponding to independent fields α^Γ. The only problem in proving this is to understand how the vanishing functional is connected to the effective action.

It can be noticed that the obtained equality can be written as follows:

$$
V = \sum_{\Gamma} V_\Gamma = 0, \tag{8}
$$

where each V_Γ is a field variation of S_Γ:

$$
V_\Gamma = D_\Gamma S_\Gamma. \tag{9}
$$

It is crucial that variations D_Γ are *different* for different Γ's. Thus, to describe the symmetry as a symmetry with respect to variations of fields we need to consider action S_{eff} as a multicomponent (vector) quantity: $S_{eff} = \{S_\Gamma\}$. As the procedure of breaking $S_{eff} = \sum_\Gamma S_\Gamma$ into peaces S_Γ is uniquely determined above for any polynomial of fields respecting momentum conservation, there is no ambiguity in this representation of the action as a multicomponent quantity. To have a possibility of applying the transformation several times consecutively, we need to break each of the variations V_Γ into a sum of functionals by the rule defined above: $V_\Gamma = \sum_{\Gamma'} V_{\Gamma'\Gamma}$. It corresponds to breaking each variation D_Γ into a sum: $D_\Gamma = \sum_{\Gamma'} D_{\Gamma'\Gamma}$. Thus, we arrive at the following representation of the gauge transformation of any multicomponent functional $\{S_\Gamma\}$:

$$(DS)_\Gamma = D_{\Gamma\Gamma'} S_{\Gamma'}, \tag{10}$$

where summation over Γ' is implied. Each component of matrix D is a field variation. It is homogeneously and linearly dependent of four gauge parameters $\{\alpha^\Gamma\}$: $D = \sum_\Gamma D^\Gamma(\alpha^\Gamma)$, where each D^Γ depends only on α^Γ. Each D^Γ is a symmetry of the effective action:

$$D^\Gamma(\alpha) S_{eff} = \sum_{\Gamma_1} D^\Gamma_{\Gamma_1\Gamma_2}(\alpha) S_{\Gamma_2} = 0, \tag{11}$$

where summation over Γ_2 is implied and α is a field taking values in the algebra of gauge symmetries of the underlying Yang–Mills theory.

Now we present an explicit form of matrixes D^Γ as a decomposition over the following basis of matrixes four by four:

$$(e^{\Gamma_1\Gamma_2})_{\Gamma_1'\Gamma_2'} = \delta^{\Gamma_1}_{\Gamma_1'} \delta^{\Gamma_2}_{\Gamma_2'}. \tag{12}$$

In terms of $e^{\Gamma_1\Gamma_2}$ variations D^Γ look as follows:

$$\begin{aligned}
D^S(\alpha) &= (-[A^H\delta_H,\alpha] - [A^L\delta_L,\alpha] - [A^R\delta_R,\alpha] + D(A^S)\alpha\delta_S)e^{HH} + \\
&\quad (-[A^L\delta_L,\alpha] + D(A^S)\alpha\delta_S)e^{LL} + \\
&\quad (-[A^R\delta_R,\alpha] + D(A^S)\alpha\delta_S)e^{RR} + (D(A^S)\alpha\delta_S)e^{SS}, \\
D^R(\alpha) &= (-[(A^H + A^L)\delta_H,\alpha] + D(A^R)\alpha\delta_R)e^{HH} + (D(A^R)\alpha\delta_R)e^{RR} + \\
&\quad (-[A^H\delta_L,\alpha] - [A^R\delta_S,\alpha])e^{HL} + (-[A^R\delta_S,\alpha])e^{RS}, \\
D^H(\alpha) &= (D(A^H)\alpha\delta_H)e^{HH} + (-[A^H\delta_S,\alpha])e^{HS} + \\
&\quad (-[(A^H + A^L)\delta_R,\alpha])e^{HR} + (-[(A^H + A^R)\delta_L,\alpha])e^{HL}.
\end{aligned} \tag{13}$$

Here condensed notations for variations of fields are used: For example, $(D(A)\alpha\delta_S)A^S_\mu = D_\mu(A)\alpha = \partial_\mu\alpha - [A_\mu,\alpha]$ and $[A\delta_L,\alpha]A^L_\mu = [A_\mu,\alpha]$. The missing expression for D^L is obtained from the above expression for D^R by the following substitutions of indices: $R \to L\,, L \to R$. We checked by explicit calculation that S_{eff} defined in (6),(7) is invariant under action of the above D^Γ, i.e., equation (11) does hold.

The next step is to calculate commutators of D^Γ's. It turns out that some of the commutators produce new local symmetries of the effective action. These symmetries can be alternatively obtained in the way similar to the one used above in the derivation of D^Γ. To this end one should start with the symmetry transformation of the underlying Yang–Mills theory generated by a commutator $[\alpha, \beta]$. Because of the presence of two fields, α and β, the procedure of multiplication of fields and subsequent dropping out of forbidden vertexes differs from the one we considered above with the single gauge parameter α. It is evident from this derivation that further multiplication of fields parameterizing the transformation in the underlying Yang–Mills theory does not produce new local symmetries of the effective theory. This conjecture can be proved by direct calculation of double commutators of D^Γ's: we checked that double commutators can be linearly expressed through D^Γ and $[D^{\Gamma_1}, D^{\Gamma_2}]$. This left us with a closed algebra consisting of nine generators. They are four D^Γ's defined in (13) and five additional generators defined as follows:

$$D^{HH}([\beta,\alpha]) = [D^H(\alpha), D^H(\beta)],\ D^{RR}([\beta,\alpha]) = [D^R(\alpha), D^R(\beta)],$$
$$D^{LL}([\beta,\alpha]) = [D^L(\alpha), D^L(\beta)],$$
$$D^{HR}([\beta,\alpha]) = [D^H(\alpha), D^R(\beta)],\ D^{HL}([\beta,\alpha]) = [D^H(\alpha), D^L(\beta)]. \tag{14}$$

To present explicitly the total algebra of commutators, we need a new notation for the superscripts on the above nine generators: the new labeling includes two indexes, ω_L and ω_R, each index runs from 0 to 2. The values of the new indexes are calculated by the values of the old ones as follows: Each letter H of the old superscripts gives unit contribution to both ω_L and ω_R, each L gives unit contribution to ω_L, each R gives unit contribution to ω_R, and any S give no contribution to any ω. For example, $D^{HH} = D^{(2,2)}$ (ω_L is first from the left), $D^{HR} = D^{(1,2)}$, while $D^S = D^{(0,0)}$. In this notations the algebra of local symmetries of the effective Regge theory looks as follows:

$$[D^{(\omega_L,\omega_R)}(\alpha), D^{(\omega'_L,\omega'_R)}(\beta)] = D^{(\omega_L+\omega'_L,\omega_R+\omega'_R)}([\beta,\alpha]), \tag{15}$$

where if $\omega_\Gamma + \omega'_{\Gamma'} > 2$, the sum is replaced by 2, and α, β belong to the algebra of the underlying Yang–Mills theory.

Now we decompose algebra (15) into a semi-direct sum of $su(N)$ (which is the gauge algebra of the underlying Yang–Mills theory) and commutative subalgeras. To this end, note that algebra (15) coincides with the gauge algebra of Yang–Mills theory which gauge group has the following algebra:

$$\mathcal{A} = su(N) \times \mathcal{A}_L \times \mathcal{A}_R, \tag{16}$$

where each $\mathcal{A}_\Gamma$ is a commutative algebra generated by three generating elements e_Γ^ω $(\omega = 0, 1, 2)$ with the following table of products:

$$e_\Gamma^{\omega_1} e_\Gamma^{\omega_2} = e_\Gamma^{\omega_1 + \omega_2}, \tag{17}$$

where the sum of ω's is understood in the same way as in (15). This table of products is simplified by going over to a new set of generating elements: $e_\Gamma = e_\Gamma^0$ (the unit element of $\mathcal{A}_\Gamma$), $\xi_\Gamma = e_\Gamma^1 - e_\Gamma^2$, and $P_\Gamma = e_\Gamma^2$. Now (17) takes the following form:

$$e_\Gamma e_\Gamma = e_\Gamma,\ e_\Gamma \xi_\Gamma = \xi_\Gamma,\ e_\Gamma P_\Gamma = P_\Gamma,\ \xi_\Gamma \xi_\Gamma = 0,\ \xi_\Gamma P_\Gamma = 0,\ P_\Gamma P_\Gamma = P_\Gamma. \tag{18}$$

From table (18) it is evident that sets $su(N) \times P_L \times P_R$, $su(N) \times \xi_L \times \xi_R$, $su(N) \times P_L \times \xi_R$, and $su(N) \times \xi_L \times P_R$ generate ideal $\mathcal{I}$ of algebra $\mathcal{A}$. This ideal can be decomposed into a direct sum of subalgebras corresponding to the above sets. Quotient algebra $\mathcal{A}/\mathcal{I}$ has in turn an ideal generated by $su(N) \times P_L \times e_R$, $su(N) \times e_L \times P_R$, $su(N) \times \xi_L \times e_R$, and $su(N) \times e_L \times \xi_R$. Thus, we have the following decomposition of $\mathcal{A}$:

$$\mathcal{A} = [su_S(N) \uplus (su_L(N) \oplus su_R(N) \oplus A_L \oplus A_R)] \uplus (su_H(N) \oplus A_H \oplus A_{HL} \oplus A_{HR}), \tag{19}$$

where $\uplus$ denotes semi-direct sum of algebras, $\oplus$ denotes direct sum of algebras, and

$$\begin{gathered} su_S(N) = su(N) \times e_L \times e_R,\ su_L(N) = su(N) \times P_L \times e_R, \\ su_R(N) = su(N) \times e_L \times P_R,\ su_H(N) = su(N) \times P_L \times P_R, \end{gathered} \tag{20}$$

while the Abelian subalgebras A_L, A_R, A_H, A_{HL} and A_{HR} are as follows:

$$\begin{gathered} A_L = su(N) \times \xi_L \times e_R,\ A_R = su(N) \times e_L \times \xi_R, \\ A_H = su(N) \times \xi_L \times \xi_R,\ A_{HL} = su(N) \times \xi_L \times P_R, \\ A_{HR} = su(N) \times P_L \times \xi_R. \end{gathered} \tag{21}$$

Dimension of each Abelian algebra entering the decomposition (19) coincides with the one of the gauge algebra of the underlying Yang–Mills theory (N^2-1 for $su(N)$).

The fact that there are Abelian algebras in the decomposition (19) makes it plausible that corresponding group is not a simply connected space. This in turn makes plausible the existence of solitons in the effective Regge theory. This guess finds some grounds in a direct analysis of the classical field equations of the effective Regge theory which will be published elsewhere.

In conclusion, effective action to describe Regge limit of Yang–Mills theories is introduced, its local symmetries are determined, and a decomposition of the algebra of symmetries into a semi-direct sum of Abelian algebras and four copies of the gauge algebra of the underlying Yang–Mills theory is found. The local symmetries are defined as transformations of any translationally invariant field functional but cannot be defined as field transformations. However, the latter is possible by the expense of introducing a notion of multicomponent action. A conjecture is risen on the presence of solitons in the effective Regge Yang–Mills theory.

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