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Infrared Abelization of Yang-Mills Theory via Abelian Higgs Variables*

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Abstract A calculation procedure for effective Abelian-Higgs-like action is presented via Faddeev-Niemi decomposition of $SU(2)$ gauge field. A natural gauge fixing is explicitly identified in the decomposition and intrinsic relation between Abelian projection and the symmetry-breaking of the Yang-Mills dynamics in infrared limit is illustrated. A London like equation is derived for chromo-electronic field.

Key words Yang-Mills field decomposition, Abelization, London like equation

1 Introduction

Since quantum chromodynamics(QCD) has been recognized as the standard theory of strong interactions, a main challenge has been to interpret the permanent confinement from first principles. It seems that the nonperturbative dynamics of infrared QCD can merely be approached effectively by lattice QCD[1, 2] yet, though many other approaches (such as QCD sum rules[3]) have been proposed, frequently employing assumptions and approximations which remain to be justified. In the framework of pure Yang-Mills(YM) theory an appealing proposal was made by Faddeev and Niemi[4] so as to separate collective infrared variables from gauge field variables by decomposing gauge connection into an Abelian part, an unit color-vector $\boldsymbol{n}$ as well as dual variables, which manifests the pure YM theory as an effective Abelian theory with duality structure between chromo-electric and chromo-magnetic field. This structure is hoped being consistent with the dual-superconductor picture[5, 6] of the confinement via the dual Meissner effect due to the supposed monopole condensation, in which the chromo-electric field between two colored sources is squeezed into a fluxtubes(vortices) and the later permanently confines colored sources(quarks) within hadrons. As shown by lattice simulations[7, 8], dual dynamics of the infrared QCD dominates in the Maximal Abelian gauge, where monopole degrees of freedom seems form a condensate responsible for confinement. Recently, lattice simulations for center vortices and monopoles(see Ref. [1] for a review) revive the interests of analytical analysis[9—11] of nonperturbative QCD. In this Letter we present a calculation procedure for effective Abelian-Higgs action based on the Faddeev-Niemi decomposition of $SU(2)$ gauge field. A natural gauge fixing is explicitly identified and the intrinsic relation between Abelian projection and the symmetry-breaking of the infrared YM dynamics is illustrated. A London like equation is also derived in terms of chromo-electronic field.

2 $SU(2)$ field decomposition and Abelian projection

It seems that the nonperturbative dynamics of

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infrared QCD can merely be approached effectively by lattice QCD[1, 2] yet, though many other approaches (such as QCD sum rules[3]) have been proposed, frequently employing assumptions and approximations which remain to be justified. According to the original idea of 't Hooft[12], Abelian projection is realized by fixing the non-Abelian part of the gauge ambiguity, breaking full gauge symmetry into that of maximal Abelian subgroup. The singularities in gauge condition lead to difference between two group manifolds and were interpreted as magnetic monopoles in the electrodynamics with residual Abelian symmetry.

We begin with $SU(2)$ YM theory where connection $\boldsymbol{A}_\mu = A_\mu^a \tau^a$ $(\tau^a = \sigma^a/2)$ describes 6 transverse ultraviolet degrees of freedom. We use inner product $\tau^a \cdot \tau^b \equiv 2\mathrm{Tr}(\tau^a\tau^b) = \delta^{ab}$, $\boldsymbol{A}\cdot\boldsymbol{B} \equiv A^a B^a$, and across product $\boldsymbol{A}\times\boldsymbol{B} = -\mathrm{i}[\boldsymbol{A},\boldsymbol{B}]$ for short. To parameterize $\boldsymbol{A}_\mu$ in the spirit of Abelian projection, we invoke the infrared 'magnetic' variable $\boldsymbol{n}(=n^a\tau^a)$, an unit vector in internal (color) space[13]. This vector naturally provides an preferred direction, breaking $SU(2)$ to $U(1)$ and leaving residual $U(1)$ symmetry (rotation around $\boldsymbol{n}$) intact. Here, the rotation around $\boldsymbol{n}$ by angle α has the form of $U(\alpha) = \cos(\alpha/2) + \mathrm{i}\boldsymbol{n}\cdot\boldsymbol{\sigma}\sin(\alpha/2) = \mathrm{e}^{\mathrm{i}\alpha\boldsymbol{n}\cdot\boldsymbol{\sigma}/2}$.

Making across product of $D_\mu\boldsymbol{n} - \partial_\mu\boldsymbol{n} = g\boldsymbol{A}_\mu\times\boldsymbol{n}$ with $\boldsymbol{n}$, where g is coupling constant, one gets

$$\boldsymbol{A}_\mu = A_\mu\boldsymbol{n} + g^{-1}\partial_\mu\boldsymbol{n}\times\boldsymbol{n} + \boldsymbol{b}_\mu \tag{1}$$

where $A_\mu \equiv \boldsymbol{A}_\mu\cdot\boldsymbol{n}$ transforms as an Abelian connection $(A_\mu \to A_\mu + \partial_\mu\alpha/g)$ for $U(1)$ rotation $U(\alpha)$ and $\boldsymbol{b}_\mu = g^{-1}\boldsymbol{n}\times D_\mu(\boldsymbol{A}_\mu)\boldsymbol{n}$ is $SU(2)$ covariant. Here, the first part $A_\mu\boldsymbol{n}$ in (1) is valued in $H = U(1)$ while the second and third terms, both of which are orthogonal to $\boldsymbol{n}$, is valued in orbit $SU(2)/H$. We note that (1) is true variable change[14] if one impose two constraints on $\boldsymbol{b}_\mu$. We also note that the fact that $\boldsymbol{A}_\mu$ does not depend upon the second term in Eq. (1) implies it has intrinsic structure. To find all relevant variables, we further decompose $\boldsymbol{b}_\mu$ in terms of $\boldsymbol{n}$. Observed that the orbit space $SU(2)/H$ can be spanned by basis $\partial_\mu\boldsymbol{n}$ and $\partial_\mu\boldsymbol{n}\times\boldsymbol{n}$, one can re-parameterize $\boldsymbol{b}_\mu$ as

$$\boldsymbol{b}_\mu = g^{-1}\rho\partial_\mu\boldsymbol{n} + g^{-1}\sigma\partial_\mu\boldsymbol{n}\times\boldsymbol{n}. \tag{2}$$

Here, the scalar fields ρ and σ can be combined to define a complex field $\phi = \rho + \mathrm{i}\sigma$. Substituting (2) into (1) we get the Faddeev-Niemi ansatz[4] for $SU(2)$ connection

$$\boldsymbol{A}_\mu = A_\mu\boldsymbol{n} + g^{-1}\partial_\mu\boldsymbol{n}\times\boldsymbol{n} + g^{-1}\rho\partial_\mu\boldsymbol{n} + g^{-1}\sigma\partial_\mu\boldsymbol{n}\times\boldsymbol{n}. \tag{3}$$

The transformation role of ρ and σ can be given by covariance of $\boldsymbol{b}_\mu$ under the rotation $U(\alpha)$. Noticing that $[\partial_\mu\boldsymbol{n}, \mathrm{e}^{-\mathrm{i}\alpha\boldsymbol{n}}] = \alpha\partial_\mu\boldsymbol{n}\times\boldsymbol{n}$, $[\partial_\mu\boldsymbol{n}\times\boldsymbol{n}, \mathrm{e}^{-\mathrm{i}\alpha\boldsymbol{n}}] = -\alpha\partial_\mu\boldsymbol{n}$, one finds

$$\begin{aligned}\boldsymbol{b}_\mu^U &= g^{-1}\mathrm{e}^{\mathrm{i}\alpha\boldsymbol{n}}(\rho\partial_\mu\boldsymbol{n} + \sigma\partial_\mu\boldsymbol{n}\times\boldsymbol{n})\mathrm{e}^{-\mathrm{i}\alpha\boldsymbol{n}} = \\ &\quad g^{-1}(\rho - \alpha\sigma)\partial_\mu\boldsymbol{n} + g^{-1}(\sigma + \alpha\rho)\partial_\mu\boldsymbol{n}\times\boldsymbol{n},\end{aligned}$$

which implies $\delta\rho = -\alpha\sigma$ and $\delta\sigma = \alpha\rho$, or

$$\delta(\rho + \mathrm{i}\sigma) = \mathrm{i}\alpha(\rho + \mathrm{i}\sigma).$$

Thus, the covariance of $\boldsymbol{b}_\mu$ yields

$$\boldsymbol{\phi} \to \phi\mathrm{e}^{\mathrm{i}\alpha}. \tag{4}$$

For this reason, the complex variables ϕ can indeed form a charged scalar field with $U(1)$ symmetry. However, the second term in Eq. (3) is not $U(1)$ covariant.

Since the connection $\boldsymbol{A}_\mu$ has 12 field components while the right hand side of Eq. (3) has 8 degrees of freedom, corresponding to 4 components of A_μ, 2 independent components of n^a and 2 components (ρ, σ), the new variables (A_μ, n^a, ϕ) are short of 4 degrees for Eq. (3). As far as on-shell degrees of freedom, A_μ has 2 transverse polarization components and this results in $2+2+2$ variables in Eq. (3), corresponding to 6 on-shell polarization components of $\boldsymbol{A}_\mu$.

Notice that Eq. (3) eliminates 4 degrees and it does not obey full gauge transformation law, for instance, under gauge rotation $\delta\boldsymbol{n} = \boldsymbol{n}\times\boldsymbol{\varepsilon}$, where ε is an iso-vector with the direction different from $\boldsymbol{n}$, we know that Eq. (3) corresponds to the singular gauge in which $A_\mu\boldsymbol{n}$ is diagonalized variables and the number of other off-diagonal terms in Eq. (3) is reduced. Therefore, one can identify Eq. (3) as a realization of the Abelian projection. One can also see that $\boldsymbol{n}$-field appears as a topological degree of freedom purely associated with orientation of $\boldsymbol{n}$.

3 Abelian Higgs action in terms of collective dual variables

We assume that order-disorder transition do occur in the infrared regime of YM theory as dual-superconductor picture of confinement suggested. Then, the effective model of QCD in confining phase should be described by the condensate of magnetic monopoles pairs and Abelized fields. As a relativistic generalization of effective superconductor model, the Abelian Higgs model has long been proposed to describe the confining phase of QCD, in which the string-like singularities provide the confining forces between field sources[5, 6].

To see the dual structure of infrared YM theory, we use the fact of asymptotic freedom of QCD[15]: coupling constant g tends to infinity in large-distance (infrared) limit. With Eq. (3), one finds

$$G_{\mu\nu}\cdot\boldsymbol{n}=F_{\mu\nu}+H_{\mu\nu}+g^{-1}\boldsymbol{n}\cdot(D_\mu\boldsymbol{n}\times D_\nu\boldsymbol{n}),$$

where $F_{\mu\nu}\equiv\partial_\mu A_v-\partial_v A_\mu$ and $H_{\mu\nu}\equiv -g^{-1}\boldsymbol{n}\cdot(\partial_\mu\boldsymbol{n}\times\partial_\nu\boldsymbol{n})$ stands for the chromo-electric and chromo-magnetic field strengths, respectively. One can identify magnetic potential C_μ by $H_{\mu\nu}\equiv\partial_\mu C_\nu-\partial_\nu C_\mu$. As a physical field, the total field exhibiting 'electronic-magnetic' duality in color space should be gauge-invariant and such a field can be given by the 't Hooft tensor[16]

$$f_{\mu\nu}=G_{\mu\nu}\cdot\boldsymbol{n}-\boldsymbol{n}\cdot(D_\mu\boldsymbol{n}\times D_\nu\boldsymbol{n})=F_{\mu\nu}+H_{\mu\nu}.$$

To find the relation of Abelian projection with symmetry breaking, we fix $\boldsymbol{n}$ to $\boldsymbol{n}_0=(\sin\gamma\cos\beta,\ \sin\gamma\sin\beta,\ \cos\gamma)$. This yields $C_\mu=g^{-1}(\cos\gamma\,\partial_\mu\beta+\partial_\mu\alpha)$, which, together with chromo-electric potential A_μ, shares only $U(1)$ symmetry ($C_\mu\to C_\mu+g^{-1}\partial_\mu\alpha$) of the partial symmetry $\mathrm{e}^{-\sigma^3\alpha}$ of $U=\mathrm{e}^{-\sigma^3\alpha}\mathrm{e}^{-\sigma^2\gamma}\mathrm{e}^{-\sigma^3\beta}$ (we assume $\boldsymbol{n}_0$ is along σ^3), whereas they do not under $\mathrm{e}^{-\sigma^2\gamma}\mathrm{e}^{-\sigma^3\beta}$. Therefore, Abelian projection responds, in $SU(2)$ case, to assigning specific direction field $\boldsymbol{n}(x)$ at each spacetime point x in Eq. (3).

By ignoring the $1/g^2$ terms(taking IR limit), one finds that Abelian component is dominant: $\boldsymbol{G}_{\mu\nu}=(F_{\mu\nu}+H_{\mu\nu})\boldsymbol{n}$, which leads to an Abelized theory with two dual $U(1)$ field ($A_\mu\leftrightarrow C_\mu$)

$$\mathscr{L}_{dual}=\frac{1}{4}(F_{\mu\nu}+H_{\mu\nu})^2. \tag{5}$$

This means, the infrared theory become the diagonal one with ϕ decoupled, consistent with the idea of Abelian Projection. Breaking of $SU(2)\to U(1)$, or fixing of the direction of $\boldsymbol{n}$ (quantum operator) at each point x makes it acquire nonvanishing vacuum expectation value(vev.), that is, $\langle n^a(x)\rangle=n^a(x)$ (c-number field), and

$$\langle\partial^\mu n^a(x)\,\partial_\nu n^a(x)\rangle=\delta^\mu_\nu\langle(\partial n^a)^2\rangle=-\delta^\mu_\nu m^2, \tag{6}$$

in which m is a mass scale and the minus sign comes from the fact $\partial_\mu\boldsymbol{n}$ is space-like for our static case. Here, δ^μ_ν arises from the requirement of Lorentz invariance of vev. and m^2 is due to that $\partial_\mu n^a$ may has a normalized factor (length in isospace) and it has dimension of [Mass]. Thus, the S^2 symmetry (rotation of $\boldsymbol{n}$) of the theory was broken by the QCD vacuum so that

$$\begin{aligned}\langle C^2_\mu\rangle &= g^{-2}\langle(\partial\boldsymbol{n})^2\rangle=-g^{-2}m^2,\\ \langle H_{\mu\nu}\rangle &= 0,\end{aligned}$$

since H is anti-symmetric. Clearly, all vev. of the field components with Lorentz indices or color indices explicitly, such as $\langle A_\mu\rangle$, $\langle C_\mu\rangle$ and $\langle\partial_\mu n^a\rangle$, are zero due to its Lorentz invariance and gauge invariance.

Including all off-diagonal components in Equ. (3) one finds

$$\begin{aligned}\boldsymbol{G}_{\mu\nu} &= \boldsymbol{n}[F_{\mu\nu}+(1-\rho^2-\sigma^2-2\sigma)H_{\mu\nu}]+\\ &\quad (g^{-1}\nabla_\mu\rho+2A_\mu)\partial_\nu\boldsymbol{n}-(g^{-1}\nabla_\nu\rho+2A_\nu)\partial_\mu\boldsymbol{n}+\\ &\quad g^{-1}\nabla_\mu\sigma\partial_\nu\boldsymbol{n}\times\boldsymbol{n}-g^{-1}\nabla_\nu\sigma\partial_\mu\boldsymbol{n}\times\boldsymbol{n},\end{aligned} \tag{7}$$

where $n_{\mu\nu}=\delta_{\mu\nu}(\partial_\rho\boldsymbol{n})^2-\partial_\mu\boldsymbol{n}\cdot\partial_\nu\boldsymbol{n}$, $\nabla_\mu\rho=\partial_\mu\rho+gA_\mu\sigma$ and $\nabla_\mu\sigma=\partial_\mu\sigma-gA_\mu\rho$. This enable us to define a $U(1)$ covariant derivative

$$\nabla_\mu\phi=\nabla_\mu\rho+\mathrm{i}\nabla_\mu\sigma=(\partial_\mu-\mathrm{i}gA_\mu)\phi,$$

With Eq. (7), one gets

$$\begin{aligned}\mathscr{L}_{dual} &= -\frac{1}{4}\{F^2_{\mu\nu}+(1-\rho^2-\sigma^2-2\sigma)^2H^2_{\mu\nu}+\\ &\quad 2(1-\rho^2-\sigma^2-2\sigma)F_{\mu\nu}H^{\mu\nu}+\\ &\quad \frac{2n_{\mu\nu}}{g^2}(\nabla^\mu\rho+2gA^\mu)(\nabla^\nu\rho+2gA^\nu)+\\ &\quad \frac{2n_{\mu\nu}}{g^2}\nabla^\mu\sigma\nabla^\nu\sigma\}.\end{aligned} \tag{8}$$

Here, we looks ϕ as collective charged field of monopole pairs. The effective Lagrangian, by taking QCD vacuum as a condensate of monopole pairs and averaging Eq. (8) over $\boldsymbol{n}$ with

$$\langle n_{\nu}^{\mu}\rangle=\delta_{\nu}^{\mu}\langle(\partial n^{a})^{2}\rangle=-\delta_{\nu}^{\mu}m^{2},$$

can be given by the Abelian Higgs like model

$$\mathscr{L}^{\rm eff} = -\frac{1}{4}F_{\mu\nu}^{2}+\frac{m^{2}}{2g^{2}}|\nabla_{\mu}\phi|^{2}+2m^{2}A_{\mu}^{2}+ \frac{2m^{2}}{g}A^{\mu}{\rm Re}(\nabla_{\mu}\phi)-V(\phi), \tag{9}$$

where

$$V(\phi)=\frac{\lambda}{4}(|\phi|^{2}-1+2{\rm Im}\phi)^{2},$$

where $\lambda\equiv\langle H_{\mu\nu}^{2}\rangle$ is positive scale and with dimension 4 and it can be shown that

$$\lambda=2m^{4}/g^{2}.$$

The verification can be done via the gauge and Lorentz invariance of the vev. In fact, from Eq. (6), one has

$$\begin{aligned}\langle H_{\mu\nu}^{2}\rangle &= g^{-2}\varepsilon_{abc}\varepsilon_{mkl}\langle n^{a}\rangle\langle n^{m}\rangle\langle\partial_{\mu}n^{b}\partial^{\mu}n^{k}\rangle\langle\partial_{\nu}n^{c}\partial^{\nu}n^{l}\rangle= \\ & g^{-2}\varepsilon_{abc}\varepsilon_{mkl}n^{a}n^{m}\delta^{bk}\langle(\partial\boldsymbol{n})^{2}\rangle\delta^{cl}\langle(\partial\boldsymbol{n})^{2}\rangle= \\ & g^{-2}\varepsilon_{abc}\varepsilon_{mbc}n^{a}n^{m}m^{4}=2!g^{-2}m^{4}.\end{aligned}$$

The effective Lagrangian is then

$$\mathscr{L}^{\rm eff} = -\frac{1}{4}F_{\mu\nu}^{2}+\frac{m^{2}}{2g^{2}}|\nabla_{\mu}\phi|^{2}+2m^{2}(1+{\rm Im}\phi){\rm A}_{\mu}^{2}+ \frac{2m^{2}}{g}A^{\mu}{\rm Re}(\partial_{\mu}\phi)-V(\phi). \tag{10}$$

We see here that, as a consequence of $\boldsymbol{n}$-field condensation, Abelian gluon field A_μ has acquired a mass $\sim m$. Accordingly, we have massive gluon without invoking the Higgs-like spontaneous symmetry-breaking(SSB) mechanism as in superconductor or dual superconductor picture for confinement. In contrast, it is due to SSB of a color direction field $\boldsymbol{n}(x)$ in the nontrivial QCD vacuum, provided that magnetic-pairs Bose condensed. The latter has recently confirmed by lattice simulation[8] in maximally Abelian gauge. For $SU(2)$ theory, change of variables Eq. (3) is equivalent to maximally Abelian gauge fixing, since $U(1)$ is maximal Abelian subgroup of $SU(2)$. Then, one has desired mass generation of gluon field for infrared QCD, the mechanism of which is not the Higgs one as commonly expected as in dual superconductor picture[5].

We note that the marginal terms have not included in Eq. (9) since new variables Eq. (3) are on-shell degrees of freedom. The marginal-term inclusion can be done by using off-shell field decomposition[17] and then calculating the effective action through quantum partition functional $Z\sim\int[{\rm d}n^{a}]{\rm e}^{-{\rm i}S}$. Toward the leading infrared term, however, our model is sufficient for the effective description of low-energy YM theory Eq. (9).

London like equation

To see the relation between the Abelian projection and mass generation of Abelian gluon field, we need dual Meissner effect as a possible signal of the monopole condensation. Here, we show that the effective model of QCD can yield, in long-distance limit, a London like equation for Abelized fields. Varying Eq. (10) gives

$$\begin{aligned}&\partial_{\mu}F^{\mu\nu}=\frac{m^{2}}{2g^{2}}[{\rm i}\phi^{*}\overleftrightarrow{\partial}^{\nu}\phi-2{\rm Re}\partial^{\nu}\phi]-m^{2}|\phi+2|^{2}A^{\nu},\\&\nabla_{\mu}\nabla^{\mu}\phi=-\frac{\partial V(\phi)}{\partial\phi^{*}}-\frac{m^{2}}{g^{2}}\partial_{\mu}A^{\mu}+{\rm i}m^{2}A_{\mu}^{2},\ \text{and c.c.,}\end{aligned} \tag{11}$$

where

$$\frac{\partial V(\phi)}{\partial\phi^{*}}=\frac{m^{2}}{g^{2}}[|\phi|^{2}-1-2{\rm Im}\phi](\phi+i).$$

We see that the chromo-electric field strongly coupled with the charged scalar field ϕ with coupling g while ϕ is weakly coupled to itself in the effective dynamics.

Taking the $g\to\infty$ limit in Eq. (11) and using Lorenz gauge for A^μ, we find

$$\begin{aligned}\phi\approx\phi_{0} &= -{\rm i}m^{2}/g^{2},\ \text{and c.c.,}\\ \partial_{\mu}F^{\mu\nu} &= j^{\nu}=-m_{V}^{2}A^{\nu},\end{aligned} \tag{12}$$

in which the second equation in Eq. (12) takes the form of London's equation. Here

$$m_{V}=m(4+m^{4}/g^{4})^{1/2}\approx2m,$$

is the mass scale responsible for dual Meissner effect and its inverse $\lambda_{L}=1/m_{V}$ determines the transverse dimensions of the chromo-electric field A_μ penetrating into the vacuum condensate. As in superconductor, Eq. (12) implies that chromo-electric field decays as

$$A_\mu(d) = A_\mu(0)\exp(-d/\lambda_L)$$

as they depart from the singular vortex tube(string), where d stands for the distance away from string. This is consistent with the dual superconductor picture[5].

We note that the London like equation similar to Eq. (12) was also derived by Dzhunushaliev[9] via ordered Abelian components assumption and Abelian projection. It is pointed out here, however, that the generation of the vector field mass m_V is owing to the quantum fluctuation of the spacial variation of the direction $\boldsymbol{n}$. We also note that the uniform assumption for scalar field ϕ only follows in infrared(long-distance) limit.

In conclusion, we presented a calculation procedure for effective Abelian-Higgs-like action based on the Faddeev-Niemi decomposition of $SU(2)$ gauge field. A natural gauge fixing is explicitly identified in this decomposition and the intrinsic relation between Abelian projection and the symmetry-breaking of the YM dynamics is illustrated in infrared limit. A London like equation is derived for chromo-electronic field with the mass generation being mainly due to quantum fluctuation of direction $\boldsymbol{n}$. This enhances the dual superconductor picture as the possible mechanism of quark confinement.

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基于阿贝尔黑格斯变量的杨–米尔斯理论的红外阿贝尔化*

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摘要 通过$SU(2)$规范场的法捷耶夫-Niemi分解给出了有效阿贝尔–黑格斯型作用量的一个计算方法. 具体指出了该分解中所用的自然规范固定以及阿贝尔投射与杨–米尔斯理论的红外动力学之间的内在关系. 推导出了色电场的一个伦敦型方程.

关键词 杨–米尔斯场分解 阿贝尔化 伦敦型方程

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