NONPERTURBATIVE GLUON MAGNETIC MASS IN THREE-DIMENSIONAL GLUODYNAMICS

By E. Casado*

Instituto de Matemática, Cibernética y Computación, La Habana

and O. K. Kalashnikov

P. N. Lebedev Physical Institute, USSR Academy of Sciences, Moscow

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The infrared behavior of the Euclidean three-dimensional gluodynamics in axial gauge is investigated. A nonperturbative expression for the three gluon vertex is constructed and with its aid the Schwinger-Dyson equation for the gluon polarization tensor is solved in the infrared limit. A nonzero gluon magnetic mass proportional to the square of the effective coupling constant $\tilde{g}^2 = g^2T$ is found.

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1. Introduction

Four-dimensional Quantum Chromodynamics at finite temperature (QCD_4) has been studied in the last years with great interest and success [1–4]. These investigations provide a good basis to construct from the first principles a reliable quantitative theory of a quark-gluon matter which is supported by the present experiments in most of the important points. Today a special interest is focused on the study of the QCD phase diagrams [5, 6] and on the predictions of the experimental signals [3, 7] inherent to the different phase states of this matter. The information obtained must clarify the statistical peculiarities of modern chromodynamics and refine their predictions related to neutron stars and the evolution of the early Universe. However, in spite of the great advances achieved, the infrared behavior of QCD_4 remains an open problem. In particular, the question about the existence of a nonzero gluon magnetic mass [8] and its dependence on the temperature and the coupling constant do not have a conclusive answer.

Two essentially different nonperturbative infrared limits at $T \neq 0$ for the transversal part of the QCD_4 polarization tensor (the limit of Π_{ij}-components) are known today:

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* Present address: Instituto Superior de Ciencias y Tecnologia Nucleares. Quinta de los Molinos, Salvador Allende y Luaces, Habana 4, Cuba.

the logarithm one [9] and the asymptotic behavior of the finite mass type [8, 10–12]. At present, however, there is no reliable way to choose between them, although the hypothesis of a finite screening of the gluomagnetic forces at $T \neq 0$ seems to be preferable. The known nonperturbative schemes (intended for studying the infrared limit of the QCD_4 polarization tensor) face actual difficulties (when the closed approximations have to be built for vertices [13]) and all the results obtained need an additional support. At the same time, the three-dimensional version of the chromodynamics (QCD_3 as the Euclidean theory in quantum field limit [14]) with the dimensional coupling constant $\tilde{g}^2 = g^2 T$ seems very useful to solve this problem. The infrared limit of QCD_3 must be qualitatively the same as the infrared limit of the temperature QCD_4 but the three-dimensional chromodynamics is more simple (the superrenormalized theory) and one may hope to obtain essential progress in this way. Moreover, all quarks may be omitted in this case since their contributions do not change the infrared behavior of the temperature chromodynamics [15].

The use of the nonperturbative calculation schemes leads to the intensive exploitation of the axial gauge

$$n_\mu A_\mu = 0 \tag{1}$$

for which the Schwinger-Dyson equation and the Slavnov-Taylor identities take their simplest forms, due to the absence of ghosts. Two main gauges are selected according to the choice of the gauge vector n_μ: the temporal axial gauge where $n = (\mathbf{0}, 1)$ and coincides with the standard vector u_μ in QCD_4, and the spatial-like gauge for which $n = (\mathbf{n}, 0)$. The latter gauge is usually applied for QCD_3 and reduces to the three-dimensional condition $n_i A_i = 0$.

In the present paper the effective three-dimensional QCD with the gauge $\mathbf{n} \cdot \mathbf{A} = 0$ is studied. We use a nonperturbative method [14, 16] of solving the Schwinger-Dyson equation for the gluon polarization tensor which exploits the special Ansatz proposed by us for the three-gluon vertex. The case of a finite screening of the gluomagnetic forces is only considered and an appropriate gluon magnetic mass is calculated. The result obtained shows that besides the trivial solution $m_{\rm mag} = 0$, there is another nonzero solution which is proportional to the square of the dimensional coupling constant $\tilde{g}^2$ of the three dimensional theory (here $\tilde{g}^2 = g^2 T$).

The content of the paper is as follows. In Section 2 we briefly outline the main features of the Euclidean QCD_3 and sketch the nonperturbative approach we use. In the next Section we discuss the properties of our Ansatz for the three-gluon vertex Γ_3, which plays an essential role in the nonperturbative scheme. This expression for Γ_3 is used in Section 4, where the equation for the gluon magnetic mass is obtained and solved. Finally, in the last Section, we discuss our results.

2. *The nonperturbative approach*

Our goal is to investigate the infrared limit of the finite-temperature gluodynamics (QCD_4) through the study of the effective three-dimensional Yang-Mills Euclidean theory [14] with the dimensional coupling constant $\tilde{g}^2 = g^2 T$.

We use the axial gauge of a special kind [17]

$$n_i A_i = 0, \quad i = 1, 2, 3 \tag{2}$$

which is more suitable for QCD_3 if one exploits the nonperturbative methods. The properties of this effective theory were explained in [14], together with the basic aspects of a nonperturbative approach used. In the following we only outline the main features of interest.

The gluon propagator $D_{ij}(p)$ and the "inverse" gluon propagator $\tilde{D}_{ij}^{-1}(p)$ are connected as follows

$$D_{ik}(p)\tilde{D}_{kj}^{-1}(p) = \delta_{ij}-(p_i n_j)/(np), \tag{3}$$

where all indices belong to the 3-dimensional space. Here $D_{ij}(p)$ is orthogonal to the gauge vector n_i

$$n_i D_{ij}(p) = 0 \tag{4}$$

and $\tilde{D}_{ij}^{-1}(p)$ is transverse to p_i

$$p_i \tilde{D}_{ij}^{-1}(p) = 0. \tag{5}$$

The gluon polarization operator $\Pi_{ij}(p)$ is defined by the standard expression

$$\tilde{D}_{ij}^{-1}(p) = \tilde{D}_{ij}^{-1}(p)^{(0)}+\Pi_{ij}(p), \tag{6}$$

where

$$\tilde{D}_{ij}^{-1}(p)^{(0)} = p^2\delta_{ij}-p_i p_j. \tag{7}$$

Since $\Pi_{ij}(p)$ is transverse, its more general tensor structure is given by the following expression

$$\Pi_{ij}(p) = (\delta_{ij}-p_i p_j/p^2)p^2\Pi_1(p)$$
$$+\left(\delta_{ij}-\frac{p_i n_j+n_i p_j}{(np)}+\frac{n_i n_j p^2}{(np)^2}\right)(np)^2\Pi_2(p), \tag{8}$$

where the structure functions $\Pi_1(p)$ and $\Pi_2(p)$ depend on the two independent variables $|p|$ and $(\hat{p}\cdot\hat{n})$. Here $(\hat{p}\cdot\hat{n}) = p\cdot n/|p|\,|n|$.

We will study the nonperturbative behavior of $\Pi_{ij}(p)$ only in the infrared limit $|p| \to 0$, and we will calculate the gluon magnetic mass following its usual definition

$$m_{\text{mag}}^2 = \lim_{|p|\to 0} \tfrac{1}{2}\Pi_{ii}(p). \tag{9}$$

We assume that the infrared limit of $\Pi_{ij}(p)$ when $|p| \to 0$ is independent of $\boldsymbol{n}$. (We will check the validity of this assumption at the end of the calculations, when we show that a solution of such kind does exist.) Under this assumption one finds that

$$\lim_{|p|\to 0} p^2\Pi_2(p) = 0 \tag{10}$$

and due to this fact all calculations become rather simple. We also note that $\Pi_1(p)$ does not depend now on $(\hat{p} \cdot \hat{n})$.

Consequently, $\Pi_{ij}(p)$ in the infrared limit $|p| \to 0$ reduces to the following expression

$$\Pi_{ij}(|p| \to 0) = (\delta_{ij} - p_i p_j / p^2)\Pi_1(|p|)p^2. \tag{11}$$

The gluon propagator, as it follows from (11), has the same tensor structure as the free propagator and depends only on one structure function $\Pi_1(|p|)$

$$D_{ij}(|p| \to 0) = \frac{1}{p^2 + p^2\Pi_1(|p|)} \times \left(\delta_{ij} - \frac{n_i p_j + n_j p_i}{(np)} + \frac{p_i p_j n^2}{(np)^2}\right). \tag{12}$$

It is worth noting that our assumption is analogous to the one made by Baker et al. [18] in their study of the gluon propagator in 4-dimensional gluodynamics at $T = 0$.

In order to determine the nonperturbative behavior of $\Pi_1(|p|)$, we use the Schwinger--Dyson equation for the gluon polarization tensor (see, for instance, Ref. [14]) contracted with $n_i n_j$

$$\delta^{ab} n_i n_j \Pi_{ij}(p) = -\frac{n_i n_j}{2}\left[\int \frac{d^3 r}{(2\pi)^3} \Gamma^{(0)abcc}_{ijkl} D_{kl}(r) + \int \frac{d^3 r}{(2\pi)^3} \Gamma^{(0)acd}_{iks}(p, q, r) D_{km}(q) D_{sn}(r) \Gamma^{bcd}_{jmn}(-p, -q, -r)\right]. \tag{13}$$

Equation (13) is exact within QCD_3 for the gauge accepted, because after contraction with $n_i n_j$ the higher-order exact graphs in which the four-gluon vertex Γ_4 is present reduce to zero. No other equation is needed for our aim, and it is important to note that only the exact three-gluon vertex Γ_3 determines Eq. (13).

The closed equation for $\Pi_1(|p|)$ is obtained from (13) if the exact three-gluon vertex Γ_3 (in the infrared domain only) is expressed in terms of Π_1.

The crucial point for the nonperturbative approach suggested is to construct the Ansatz for $\Gamma_3(p, q, r)$ that correctly approximates all the infrared peculiarities of the exact vertex when one of its momenta goes to zero. We are going to discuss now this point in detail.

3. *The three-gluon vertex*

The constructed three-gluon vertex $\Gamma_3(p, q, r)$ must be expressed only through the structure functions Π_j — which define expression (8) — and must satisfy a number of exact properties of the real vertex. Particularly, it must obey the Bose symmetry condition

$$\Gamma(p, q, r)^{abc}_{ijk} = \Gamma(q, p, r)^{bac}_{jik} = \Gamma(r, q, p)^{cba}_{kji} = \Gamma(p, r, q)^{acb}_{ikj} \tag{14}$$

as well as the exact Slavnov-Taylor identity

$$p_i \Gamma(p, q, r)^{abc}_{ijk} = i\tilde{g} f^{abc}[\tilde{D}^{-1}_{jk}(q) - \tilde{D}^{-1}_{jk}(r)] \tag{15}$$

and must also have some correct limits.

The problem is that the vertex $\Gamma^{(L)}(p, q, r)_{ijk}^{abc}$ satisfying these properties is still too arbitrary. Any function $\Gamma^{(T)}(p, q, r)_{ijk}^{abc}$ which satisfies (14) and the usual transverse conditions

$$p_i\Gamma^{(T)}(p, q, r)_{ijk}^{abc} = q_j\Gamma^{(T)}(p, q, r)_{ijk}^{abc} = r_k\Gamma^{(T)}(p, q, r)_{ijk}^{abc} = 0 \tag{16}$$

is not fixed by Eq. (15), and the new function

$$\Gamma(p, q, r)_{ijk}^{abc} = \Gamma^{(L)}(p, q, r)_{ijk}^{abc} + \Gamma^{(T)}(p, q, r)_{ijk}^{abc} \tag{17}$$

also satisfies (14) and (15). Thus, we need to complement Properties (14) and (15) with another requirement in order to specify the form of the Ansatz for Γ_3.

The one-loop perturbative calculations [19] point out that the standard relation for Π_{ij} which is known as the differential Slavnov-Taylor identity

$$\Gamma(p, -p, 0)_{ijk}^{abc} = -i\tilde{g}f^{abc}\frac{\partial \tilde{D}_{ij}^{-1}(p)}{\partial p_k} \tag{18}$$

must hold. But for the nonperturbative calculations within QCD_3 this requirement is not acceptable, because in QCD_3 it is exactly equivalent to impose that

$$\lim_{|p| \to 0} \Pi_{ij}(p) = 0. \tag{19}$$

However, it is precisely the validity of (19) what we want to investigate, and we expect that it does not take place in general.

The fact that requirement (18) is equivalent to condition (19) was verified in the one--loop nonperturbative approximation [20]. But it is a specific of QCD_3 that in the exact theory, to require (18) is also the same as to demand (19). Our proof of this is analogous to the one made by other authors in four-dimensional gluodynamics at zero temperature [21], and it is important to stress that only exact properties of QCD_3 are used here.

From the Slavnov-Taylor identity (15) once the momentum r is equalled to zero we obtain the more simple expression of this identity

$$p_i\Gamma(p, -p, 0)_{ijk}^{abc} = i\tilde{g}f^{abc}[\tilde{D}_{jk}^{-1}(p) - \tilde{D}_{jk}^{-1}(0)], \tag{20}$$

where the existence of a vertex limit is proposed in the usual way, without singularities. If for QCD_3 expression (18) takes place, this means that identity (20) reduces to the following equality

$$-i\tilde{g}f^{abc}p_i\frac{\partial \tilde{D}_{ij}^{-1}(p)}{\partial p_k} = i\tilde{g}f^{abc}[\tilde{D}_{jk}^{-1}(p) - \tilde{D}_{jk}^{-1}(0)] \tag{21}$$

which can be rewritten as

$$\tilde{D}_{jk}^{-1}(0) = \tilde{D}_{jk}^{-1}(p) + p_i\frac{\partial \tilde{D}_{ij}^{-1}(p)}{\partial p_k}. \tag{22}$$

The last expression is equivalent to the important condition

$$\tilde{D}_{jk}^{-1}(0) = \frac{\partial}{\partial p_k}(p_i \tilde{D}_{ij}^{-1}(p)) \tag{23}$$

from which we obtain the nonstandard result

$$\tilde{D}_{jk}^{-1}(0) = 0 \tag{24}$$

taking into acount the transversality (5) of $\tilde{D}_{ij}^{-1}(p)$. Then, having in mind Eqs (6) and (7), we must conclude that

$$\Pi_{jk}(0) = 0 \tag{25}$$

and therefore, that no finite gluon magnetic mass exists if the differential Slavnov-Taylor identity holds.

The equivalence found above is a very important result. Now, the additional requirement we shall impose on Γ_3 is not that (18) holds for any Π_{ij}, but instead, that our Ansatz for $\Gamma_3(p, q, r)$ depends on the structure functions of $\Pi_{ij}(p)$ in such a way that

$$\lim_{|r|\to 0} \Gamma(p, q, r)_{ijk}^{abc} = -i\tilde{g}f^{abc}\frac{\partial \tilde{D}_{ij}^{-1}(p)}{\partial p_k} \tag{26}$$

only when

$$\lim_{|r|\to 0} \Pi_{ij}(r) = 0.$$

The vertex which satisfies requirements (14), (15) and (26), without any assumption on the behavior of Π_1 and Π_2 can be constructed. But, in the present paper, we are not considering the general case for Γ_3, since we only exploit this vertex to close Equation (13) in which we have already adopted the approximate expression for $\Pi_{ij}(p)$.

Therefore, in order to be consistent with the approximation made, we use here the more simple Ansatz for Γ_3

$$\Gamma(p, q, r)_{ijk}^{abc} = \Gamma^{(\mathrm{L})}(p, q, r)_{ijk}^{abc} + \Gamma^{(\mathrm{T})}(p, q, r)_{ijk}^{abc}, \tag{27}$$

where $\Gamma^{(\mathrm{L})}$ is built in a standard manner

$$\Gamma^{(\mathrm{L})}(p, q, r)_{ijk}^{abc} = -i\tilde{g}f^{abc}$$
$$\times\Bigg(\delta_{ij}\{p_k[1+\Pi_1(|p|)] - q_k[1+\Pi_1(|q|)]\}$$
$$+[q_i p_j - \delta_{ij}(p\cdot q)]\frac{\Pi_1(|p|)-\Pi_1(|q|)}{p^2-q^2}(p_k - q_k)\Bigg) + \text{cyclic symmetric terms} \tag{28}$$

and $\Gamma^{(T)}$ is introduced in order to satisfy requirement (26)

$$\Gamma^{(T)}(p, q, r)_{ijk}^{abc} = -i\tilde{g}f^{abc}\frac{2}{p^2+q^2+r^2}$$

$$\times[\Pi_1(|p|)+\Pi_1(|q|)+\Pi_1(|r|)]$$

$$\times\left[(q\cdot rp_k-p\cdot rq_k)\delta_{ij}+\frac{r_ip_jq_k-q_ip_kr_j}{3}\right]+\text{cyclic symmetric terms.} \tag{29}$$

The proposed vertex is free of kinematic singularities and reproduces the exact bare vertex when $\Pi_1 = 0$.

The main assumption within this nonperturbative approach is that the constructed Ansatz (27)–(29) for Γ_3 coincides with the exact vertex in the infrared limit, and due to this fact our master equation (13) can be solved exactly.

We now return to Equation (13).

4. *The self-consistent equation and its solutions*

Substituting our Ansatz for Γ_3 in Equation (13) and taking within it the limit $|p| \to 0$ we obtain a closed equation for $\Pi_1(|p|)$, because in this limit the gluon propagators are approximated by only one function according to Eq. (12).

In order to solve the equation found we consider the simplest Ansatz for $\Pi_1(|p|)$

$$\Pi_1(|p|) = \frac{M^2}{p^2}, \tag{30}$$

where M^2 is a positive constant directly connected with the gluon magnetic mass. After assumption (30) is taken into account all algebra is easily performed and the resulting equation for M^2 has the following form

$$M^2[1-(\hat{n}\cdot\hat{p})^2] = \frac{\tilde{g}^2N}{n^2}$$

$$\times\int\frac{d^3r}{(2\pi)^3}n_k\frac{\partial}{\partial r_k}\left\{\frac{(n\cdot r)}{r^2+M^2}\left[1+\frac{1}{(\hat{n}\cdot\hat{r})^2}\right]\right\}$$

$$+\tilde{g}^2NM^2\int\frac{d^3r}{(2\pi)^3}\frac{1}{(r^2+M^2)^2}\left[2(\hat{n}\cdot\hat{p})^2+2(\hat{r}\cdot\hat{p})^2\right.$$

$$+3(\hat{n}\cdot\hat{r})^2-5(\hat{n}\cdot\hat{p})(\hat{r}\cdot\hat{p})(\hat{n}\cdot\hat{r})-\frac{2(\hat{n}\cdot\hat{p})(\hat{r}\cdot\hat{p})}{(\hat{n}\cdot\hat{r})}$$

$$\left.+\frac{1-2(\hat{n}\cdot\hat{p})^2}{(\hat{n}\cdot\hat{r})^2}+\frac{(\hat{n}\cdot\hat{p})(\hat{r}\cdot\hat{p})}{(\hat{n}\cdot\hat{r})^3}\right], \tag{31}$$

where N is the dimension of the SU(N) group considered (here $N = 3$). It is important to note that calculation of all integrals within Eq. (31) is straightforward.

The first integral in the right-hand side of Eq. (31) equals zero. This fact can be verified by taking the derivative and evaluating the resulting integral with the use of both, dimensional regularization [22] and the principal value prescription [17] for the poles $(n \cdot r)^{-m}$

$$(n \cdot r)^{-m} = \lim_{\varepsilon \to 0} \tfrac{1}{2} [(n \cdot r + i\varepsilon)^{-m} + (n \cdot r - i\varepsilon)^{-m}]. \tag{32}$$

The second integral in the right-hand side of Eq. (31) is free of infrared and ultraviolet divergences. The only singularities present are the axial gauge poles, which must be handled by means of prescription (32).

Performing all the integrations, the right-hand side of Eq. (31) gives the expression

$$\frac{\tilde{g}^2 N}{12\pi} [1 - (\hat{n} \cdot \hat{p})^2] \, |M| \tag{33}$$

whose $(\hat{n} \cdot \hat{p})$ dependence exactly cancels the term $[1 - (\hat{n} \cdot \hat{p})^2]$ in the left-hand side of Eq. (31). This agrees with our initial hypothesis about the independence of the infrared limit of Π_{ij} on $\boldsymbol{n}$.

The equation obtained for M has the form

$$M^2 = \frac{\tilde{g}^2 N}{12\pi} |M| \tag{34}$$

and besides the trivial solution

$$|M| = 0 \tag{35}$$

contains the other one

$$|M| = \frac{\tilde{g}^2 N}{12\pi} \tag{36}$$

which corresponds to the nonzero gluomagnetic mass.

Solution (35) reproduces the situation known from the lowest order of perturbation theory [23]. In this case $\Pi_{ij}(|p| = 0) = 0$, and the vertex $\Gamma_3(p, -p, 0)$ satisfies (18).

Solution (36) corresponds to a new situation where $\Pi_{ij}(|p| \to 0) \neq 0$ and Γ_3 does not satisfy (18). In this case $|M|$ represents a dynamically generated mass which regularizes the infrared domain of the three-dimensional theory.

5. *Conclusions*

In the present paper we have obtained the nonperturbative infrared behavior of the gluon polarization tensor which demonstrates the finite dynamical screening of gluomagnetic forces when $|p| \to 0$. The basic equation (34) found for the gluon magnetic mass possesses two solutions which are gauge invariant within the class of axial gauges $n_i A_i = 0$, but have different physical senses. Solution (35) repeats the standard perturbative result and does not seem to reflect the correct infrared behavior of QCD_3. The other solution has a completely nonperturbative nature and represents the possible mechanism for the dynamical screening of gluomagnetic forces. For QCD_4 Solution (36) (which is qualitatively the same as the one obtained in Ref. [24]) leads to the analytical behavior for the gluon magnetic mass

$$m_{\text{mag}} = \frac{N}{12\pi} g^2 T \tag{37}$$

which has been discussed by many authors (see, for instance, Refs. [8, 10]) but is not supported (if the QCD_3 results are rejected) by direct calculations. Moreover, calculations performed within QCD_4 [6, 12] have given the different result

$$m_{\text{mag}}^2 = a^2 g^4 T^2 \ln(1 + k^2/g^2) \tag{38}$$

for the gluon magnetic mass, which is probably more inherent to the real gluodynamics.

Unfortunately, the present state of knowledge makes it impossible to choose between these limits, since there are no arguments to insist that the infrared limit of QCD_3 is exactly equal (not only qualitatively) to the appropriate infrared limit of QCD_4. According to this fact no disagreements exist, and we consider result (37) found for the gluon magnetic mass within QCD_3 as a good infrared limit for the QCD_4 polarization tensor. Particularly, this limit coincides (up to the numerical factor) with the behavior of m_{mag} found in Ref. [10].

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