Electric and magnetic screenings of gluons in a model with dimension-2 gluon condensate

Fukun Xu† and Mei Huang†,‡

1 Institute of High Energy Physics, Chinese Academy of Sciences, Beijing, China
2 Theoretical Physics Center for Science Facilities, Chinese Academy of Sciences, Beijing, China

(Dated: November 23, 2011)

Electric and magnetic screenings of the thermal gluons are studied by using the background expansion method in a gluodynamic model with dimension-2 gluon condensate. At low temperature, the electric and magnetic gluons are degenerate. With the increasing of temperature, it is found that the electric and magnetic gluons start to split at certain temperature $T_0$. The electric screening mass changes rapidly with temperature when $T > T_0$, and the Polyakov loop expectation value rises sharply around $T_0$ from zero in the vacuum to a value around 0.8 at high temperature. This suggests that the color electric deconfinement phase transition is driven by electric gluons. It is also observed that the magnetic screening mass keeps almost the same as its vacuum value, which manifests that the magnetic gluons remains confined. Both the screening masses and the Polyakov loop results are qualitatively in agreement with the Lattice calculations.

I. INTRODUCTION

QCD vacuum is characterized by spontaneous chiral symmetry breaking and color confinement. It is expected that chiral symmetry can be restored and color degrees of freedom can be freed at high temperature and/or density.

The spontaneous chiral symmetry breaking is well understood by the dimension-3 quark condensate $\langle \bar{q}q \rangle$ in the vacuum, which is the order parameter in the chiral limit when the current quark mass is zero $m = 0$, and the chiral restoration is characterized by the vanishing of quark condensate.

The mechanism of confinement still remains as a challenge. The confinement is normally taken as the color singlet nature of the spectrum. However, the color singlet spectrum nature is not unique for QCD, but also holds for gauge-Higgs theories in which the gauge group is spontaneously broken. From the specific feature of QCD dynamics, the Regge trajectories of hadrons indicate the string-picture of hadrons, and the confinement can be described by the string picture of hadrons or the linear potential between two quarks at large distances, i.e. $V_{QQ}(R) = \sigma R$ with $\sigma$ the string tension. There have been great efforts in understanding the emergence of string-like object, e.g. the Abrikosov flux tubes [2], the dual superconductor scenario induced by monopole condensation [3], and the center vortices [4]. In the limit of infinite heavy current quark mass, the flux tube never breaks, and it corresponds to the scenario of “permanent confinement”. From the symmetry point of view, when the current quark mass goes to infinity $m \to \infty$, QCD becomes pure gauge SU(3) theory, which is center symmetric in the vacuum. The non-vanishing string tension corresponds to the area law for the Wilson loop, vanishing Polyakov lines, perimeter-law for the ‘t Hooft loops or the area-law falloff for the vortex free energy [5]. The deconfinement phase transition referring to the “permanent confinement” is characterized by the breaking of center symmetry, and the usually used order parameter is the Polyakov loop expectation value $\langle L \rangle$ [6].

There have been also great efforts in understanding confinement and deconfinement from low-energy Gluodynamics. Varies of vacuum condensates provide important information to understand the non-perturbative dynamics of QCD. For example, the gauge invariant dimension-4 gluon condensate $\langle g^2 G^2 \rangle$ has been widely investigated in both QCD sum rules and lattice calculations [7,9], and the non-vanishing value of the condensate does not signal the breaking of any symmetry directly, but rather the non-perturbative dynamics of strongly interacting gluon fields. In last decade, there have been growing interests in dimension-2 gluon condensates $\langle g^2 A^2 \rangle$ in SU($N_c$) gauge theory [10–20], with the local dimension-2 operator

$$A^2(x) = \sum_{a=1}^{N_c^2-1} \sum_{\mu=1}^{4} A^a_\mu(x) A^a_\mu(x).$$

The dimension-2 gluon condensate breaks the property of gauge invariance, and it has been investigated in varies of gauges. For example, the dimension-2 operator $A^2$ gets a special meaning in the Landau gauge [15,18], in which the condensate is at an extremum and plays as a saddle point on its gauge orbit, and a BRST-invariant mixed gluon-ghost condensate has been introduced in [16]. Though it is not gauge invariant, the growing interests in the dimension-2 gluon condensate lies in that it is related to the production of the dynamical gluon mass, and the possible connection between the minimal value of the $< A^2 >_{\text{min}}$
and the topological defects (e.g., the magnetic monopoles \[13\]). Furthermore, the dimension-2 gluon condensate has a more close relation with confinement, the dimension-2 gluon condensate yields the UV corrections \(A^2/Q^2\) in the QCD running coupling constant \(\alpha_s(Q^2)\), which leads to the linear potential \(\sigma R\) at short distances with \(\sigma_s \approx g_R^2 < A_0^2 >\).

It is of great interest to investigate the behavior of the dimension-2 gluon condensate at finite temperature and its role in the deconfinement phase transition. At zero temperature case the space-time space is symmetric under the \(O(4)\) rotation, i.e., all Lorentz components of the gauge field \(A_\mu\) contribute equally to the vacuum. In the finite temperature, it is more appropriate to divide the gauge boson into time-like (electric) and space-like (magnetic) components \[21, 22\]. This can be viewed as the di-electric theories \[23–33\]. Significant evidence shows that gluon confinement is not a perturbative effect is one of the main features of the quark-gluon plasma (QGP) and has been widely investigated in lattice and effective theories \[21\].

On the other hand, the color screening effect is one of the main features of the quark-gluon plasma (QGP) and has been widely investigated in lattice and effective theories \[21\]. Significant evidence shows that gluon confinement is not affected by a small (physical) number of light quarks \[27, 32\] and the nonperturbative features of QCD are most probably generated in the gauge sector. It is therefore reasonable to study the behavior of screening of gluons at finite temperature. Lattice result shows that the QCD coupling constant strength near the critical temperature \(T_c\) is still of the order of one \[28\], and the perturbation theory cannot be applied in this region. Especially in the regime right above the critical temperature, the nonperturbative effects are supposed to be important.

Therefore, in this work we extend the pure gluodynamic model with dimension-2 gluon condensate in the vacuum \[34\], and estimating the electric as well as magnetic screening masses of gluons at finite temperature. We also investigate the contribution of the dimension-2 gluon condensate to the deconfinement phase transition.

This paper is structured as follows. In Sec. II we introduce the pure gluodynamic model with dimension-2 gluon condensate in the vacuum \[34\], and give the numerical results of the electric and magnetic screening masses as well as the Polyakov loop expectation value in Sec. IV and give the summary in Sec. V.

### II. THE GLUODYNAMIC MODEL WITH DIMENSION-2 GLUON CONDENSATE

In this section, we follow Ref. \[34\] to introduce the Celenza-Shakin model which gives the effective action for pure gluon system with dimension-2 gluon condensate. As an overall notation the paper is in the framework of Euclidean space.

The pure gluon part of QCD Lagrangian is described by

\[
\mathcal{L}_G = -\frac{1}{4} G^\mu_\nu G^\mu_\nu, \tag{2}
\]

with

\[
G^\mu_\nu = \partial_\mu A^\nu_\alpha - \partial_\nu A^\mu_\alpha + gf^{abc} A^\mu_\alpha A^\nu_\beta. \tag{3}
\]

Motivated by the Nambu–Jona-Lasinio model with quark-antiquark condensate in the vacuum, which is similar to the BCS pairing condensation in the superconductor, Celenza-Shakin proposed the "pairing" of two gluons condensates in the vacuum in Ref. \[34\]. The gluon field can be decomposed into a condensate field \(\tilde{A}^a_\mu\) and a fluctuating field \(\tilde{\phi}^a_\mu\) \[34, 35\] as,

\[
A^\mu_\alpha(x) := \tilde{A}^a_\mu + \tilde{\phi}^a_\mu(x), \tag{4}
\]

where \(\tilde{A}^a_\mu\) is macroscopically occupied and independent of \(x\), which carries zero vacuum expectation value, i.e. \(\langle\text{vac}|\tilde{A}^a_\mu|\text{vac}\rangle = 0\).

The Fourier transformation of Eq. (4) has the form of

\[
A^\mu_\alpha(k) := A^{a\mu}_\alpha(k = 0) + \tilde{\phi}^{a\mu}_\alpha(k) \equiv \bar{A}^{a\mu}_\alpha + \tilde{\phi}^{a\mu}_\alpha(k), \tag{5}
\]

where the background \(\bar{A}^{a\mu}_\alpha\) carries only zero momentum mode, and for simplicity we assume it to be a constant.

By using the expansion Eq. (5), the gluon part of the QCD Lagrangian becomes

\[
\mathcal{L}_G = -\frac{1}{4} [g^\mu_\nu g^\rho_\sigma + 2gf^{abc}g^\mu_\nu(a^{abc}_\mu b^{abc}_\rho + b^{abc}_\mu a^{abc}_\rho + b^{abc}_\rho a^{abc}_\mu)] + 2g^2 f^{abc} f^{cde} \tilde{\phi}^{a\mu}_\rho b^{b\nu}_\alpha (\tilde{\phi}^{c\mu}_\nu d^{c\rho}_\alpha + \tilde{\phi}^{d\rho}_\nu c^{d\nu}_\alpha + \tilde{\phi}^{c\nu}_\rho b^{c\nu}_\alpha) + g^2 f^{abc} f^{cde} \tilde{\phi}^{a\mu}_\rho b^{b\nu}_\alpha (\tilde{\phi}^{c\mu}_\nu d^{c\rho}_\alpha + \tilde{\phi}^{d\rho}_\nu c^{d\nu}_\alpha + \tilde{\phi}^{c\nu}_\rho b^{c\nu}_\alpha) + g^2 f^{abc} f^{cde} \tilde{\phi}^{a\mu}_\rho b^{b\nu}_\alpha (\tilde{\phi}^{c\mu}_\nu d^{c\rho}_\alpha + \tilde{\phi}^{d\rho}_\nu c^{d\nu}_\alpha + \tilde{\phi}^{c\nu}_\rho b^{c\nu}_\alpha), \tag{6}
\]
As a further assumption one can treat $\hat{A}^a_\mu$ as a classical variable:

$$\hat{A}^a_\mu := \phi_0 \hat{\eta}^a_\mu,$$

where $\phi_0$ is constant and $\hat{\eta}^a_\mu$ is a vacuum vector. The vector $\hat{\eta}^a_\mu$ has the following properties:

$$\hat{\eta} \equiv \frac{\eta}{|\eta|}, \quad \eta^a_\mu \equiv (\eta^a_\mu, \hat{\eta}^a_\mu), \quad (\hat{\eta}^a_\mu, \hat{\eta}^a_\nu) = 1, \quad \eta^2 = \eta^a_\mu \eta^a_\mu = 32.$$  

The averaging procedure for an operator $O[\hat{\eta}]$ may be written as

$$\langle O[\hat{\eta}] \rangle_{\eta} = \frac{\int \prod \eta \; \delta(\hat{\eta} \cdot \hat{\eta} - 1) O[\hat{\eta}]}{\int \prod \eta \; \delta(\hat{\eta} \cdot \hat{\eta} - 1)}.$$  

Now that the field $\eta^a_\mu$ plays as the vacuum degree of freedom, then one can consider the expectation value this averaging as the vacuum expectation i.e.

$$\langle \text{vac} | O[A^a_\mu] | \text{vac} \rangle \equiv \langle \text{vac} | O[\hat{A}^a_\mu] | \text{vac} \rangle \equiv \langle O[\hat{\eta}^a_\mu] \rangle_{\eta}.$$  

After taking the expecting value in terms of $\eta^a_\mu$, one gets

$$\langle \hat{A}^a_\mu \hat{A}^b_\nu \rangle_{\eta} = \frac{\delta^{ab} \delta_{\mu \nu}}{8} \phi_0^2, \quad \langle \hat{A}^a_\mu \hat{A}^a_\nu \rangle_{\eta} = \phi_0^2.$$  

Actually it is the nonzero expectation value of the double combination $A^2$ plays as an order parameter representing the existence of condensate but not the gauge field $A^a_\mu$ as one spontaneously has the constraint of

$$\langle O[(\hat{A}^a_\mu)^{\text{odd}}] \rangle_{\eta} = 0.$$  

Then the Lagrangian after this background expansion becomes

$$\langle \mathcal{L} \rangle_{\eta} = -\frac{1}{4} (GG)_{\eta} = -\frac{1}{4} \left[ \mathcal{G} + 2m_E^2 \mathcal{A}^2 + 4b \phi_0^4 \right],$$  

with

$$m_E^2 = \frac{9}{32} g^2 \phi_0^2, \quad b = \frac{9}{136} g^2.$$  

The gluon gets mass because of the existence of nonperturbative dimension-2 gluon condensate. We note that the dimension-four gluon condensate $\langle g^2 A^2 \rangle_{\eta}$ is proportional to dimension-2 gluon condensate $\langle g^2 \rangle_{\eta}^2$.

### III. ELECTRIC AND MAGNETIC SCREENING AT FINITE TEMPERATURE

We now use the Lagrangian in Eq. (10) as the effective model of pure gluon system. At finite temperature, the temporal and spatial direction of the gluon field is in general different, i.e.

$$\mathcal{A} := (\mathcal{A}_4, \mathcal{A})$$

and the Lagrangian can be written as

$$\langle \mathcal{L} \rangle_{\eta} = -\frac{1}{4} (GG)_{\eta} = -\frac{1}{4} \left[ \mathcal{G} + 2(m_E^2 \mathcal{A}_4^2 + m_M^2 \mathcal{A}^2) + 4b \phi_0^4 \right],$$  

In the zero temperature limit, one has $m_E^2 = m_M^2 = m_E^2$.

By adding the gauge-fixing term in Lagrangian i.e. $\mathcal{L}_{\text{fix}} = -\frac{1}{2} (\partial_\mu \mathcal{A}_\mu)^2$, one can solve the gluon propagator of the fluctuating field $\mathcal{A}_\mu$ from the equation of

$$\left[ K^2 \delta_{\mu \nu} - (1 - 1/\xi) K_\mu K_\nu + m_E^2 \delta_{\mu 4} + m_M^2 \delta_{\mu j} \right] : D_{\nu \nu}(K) = \delta_{\mu \sigma}.$$  

(14)
In the following numerical calculation, we choose the dimension-four gluon condensate value of $\langle g^2 F^2 \rangle = (0.009 \pm 0.006) \times 4\pi^2\text{GeV}^4$ [42]. We take the value of dimension-four gluon condensate as $\langle g^2 A^2 \rangle = 1.16\text{GeV}^2$ and the gluon mass $m_g = 571\text{MeV}$.

For the calculation of the momentum integral, we employ a soft-cutoff function (for example see [43]), which takes the form of

$$f(K) = e^{-\Lambda K^2} \equiv e^{-\Lambda (\omega+q^2)}.$$  

(20)

In the following numerical calculation, we choose $\Lambda = 0.3\text{[GeV]}^{-1}$. 

The gluon propagator has the form of

$$D_{\mu\nu}(K) = \frac{P_{\mu\nu}}{K^2 + m_M^2} + \frac{K^2 P_{\mu\nu}}{K^2 (K^2 + m_E^2) - K^2 (m_M^2 - m_M^2) + \xi (k^2 m_E^2 + k^2 m_M^2 + m_M^2 m_E^2)}.$$  

(15)

In the limit of $\xi \to \infty$, i.e. in the unitary gauge, the gluon propagator takes the form of

$$D_{\mu\nu}(K) = \frac{1}{K^2 + m_E^2} \left( \delta_{\mu\nu} - \frac{k_\mu k_\nu}{k^2} \right) + \frac{1}{k^2 m_M^2 + K_1^2 m_E^2 + m_E^2 m_M^2} \left( \delta_{44} m_M^2 + K_4 K_4 + m_E^2 k_M k_V \right).$$  

(16)

In the zero temperature limit ($m_E^2 = m_M^2 = m_E^2$) it becomes a simple form

$$D_{\mu\nu}(K) = \frac{1}{K^2 + m_M^2} \left( \delta_{\mu\nu} - \frac{K_4 K_4}{m_M^2} \right).$$

The screening masses are defined as the gluon self-energy tensor $\Pi_{\mu\nu}^{\text{ab}}(p_4, p)$ at the static limit ($p_4 = 0$, $p \to 0$) [36, 37], and the electric and magnetic screening masses take the following expressions:

$$m_E^2 \delta_{44} \delta^{ab} = -\Pi_{44}^{\text{ab}}(0, p \to 0), \quad m_M^2 \delta_{ij} \delta^{ab} = -\Pi_{ij}^{\text{ab}}(0, p \to 0).$$  

(17)

Here the gluon self-energy tensor is with full propagator so that it contains both the perturbative and the nonperturbative contributions of the interaction of the gauge field. As it was pointed by some authors that, the above definition does not yield a gauge invariant definition of the screening masses in a strict sense.

On the other hand, we suggest a nonperturbative iterative relation of gluon mass similar to Dyson-Schwinger method [38], i.e. the value of screening mass especially at finite temperature is decided by the gluon self-energy Fig.1.

The direct calculation by using propagator Eq.16 gives

$$\Gamma_{G_4}^{\text{ab}}(P = 0) = -g^2 N_c \delta_{44} \delta^{ab} T \sum_n \int \frac{d^3 k}{(2\pi)^3} \frac{(2 - \omega^2_2 + k^2 + m_M^2)}{(K^2 + m_M^2)^2} \frac{k^2 m_M^2 - \omega_2^2 m_E^2 + m_E^2 m_M^2}{(k^2 m_M^2 + \omega_2^2 m_E^2 + m_E^2 m_M^2)^2}.$$  

(18)

$$\Gamma_{G_{ij}}^{\text{ab}}(P = 0) = -g^2 N_c \delta_{ij} \delta^{ab} T \sum_n \int \frac{d^3 k}{(2\pi)^3} \frac{(2 - \omega_2^2 + k^2/3 + m_M^2)}{(K^2 + m_M^2)^2} \frac{k^2 m_M^2/3 + \omega_2^2 m_E^2 + m_E^2 m_M^2}{(k^2 m_M^2 + \omega_2^2 m_E^2 + m_E^2 m_M^2)^2}.$$  

(19)

Then one immediately gets the results with parameters given.

IV. RESULTS AND DISCUSSIONS

We firstly investigate the thermal behavior of electric and magnetic screening masses by using the definition Eq.17 and the electric and magnetic gluon self-energy in Eqs.15 and 19.

In our model, there are two input parameters, i.e. the dimension-four gluon condensate $g^2 G^2$ or the nonperturbative coupling constant $g$, and the momentum cutoff parameter $\Lambda$ at zero temperature. For simplicity we assume that the coupling constant $g$ and cutoff parameter $\Lambda$ remain constants even at finite temperature. The value of dimension-four gluon condensate at zero temperature are derived both in QCD sum-rules (lower range of the interval) [7, 39] and in lattice (higher range of the interval) [40, 41]. Different authors give different results but an acceptable candidate is $\langle g^2 G^2 \rangle = (0.009 \pm 0.006) \times 4\pi^2\text{GeV}^4$. We take the value of dimension-four gluon condensate as $\langle g^2 G^2 \rangle = 0.009 \times 4\pi^2\text{GeV}^4$, which corresponds to the dimension-2 gluon condensate $\langle g^2 A^2 \rangle = 1.16\text{GeV}^2$ and the gluon mass $m_g = 571\text{MeV}$.

For the calculation of the momentum integral, we employ a soft-cutoff function (for example see [43]), which takes the form of

$$f(K) = e^{-\Lambda K^2} \equiv e^{-\Lambda (\omega+q^2)}.$$  

(20)
A. The electric and magnetic screening masses

The electric and magnetic screening masses as functions of the temperature are shown in Fig. 2, the solid line and the dashed-dotted line are for the electric and magnetic part, respectively.

![Figure 2: The electric and magnetic screening masses as functions of the temperature.](image)

It is found that both electric and magnetic screening masses are degenerate and remain unchanged at low temperature, and the electric and magnetic components start to split at the temperature $T_0 = 150$ MeV. In the temperature region $T > T_0$, the electric screening mass rise rapidly with the increase of temperature, however, the magnetic screening mass of the gluons remains almost the same as its vacuum value.

In order to compare with the lattice data in Ref. [30], we divide the screening masses by the temperature. We also assume the critical temperature $T_c = T_0 = 150$ MeV, where $m_E$ and $m_M$ start to split. (The exact value of $T_c$ is not important here, and will not affect the qualitative property of the ratio of the screening mass over the temperature.) Fig. 3 shows the ratios of $m_E/T$ and $m_M/T$ as functions of $T/T_c$ and compare with the lattice data in Ref. [30]. The solid line and the dashed-dotted line are for the electric and magnetic part, respectively.

![Figure 3: The ratios of the screening masses $m_E/T$ and $m_M/T$ as functions of $T/T_c$. The lattice data are taken from Ref. [30].](image)

It is found that the ratio of the electric screening mass over temperature $m_E/T$ is around $\sim 1.8$ in the region of $2 < T/T_c < 5$, which is qualitatively in agreement with the lattice result $m_E/T \sim 2.3$. The ratio of the magnetic screening mass over temperature $m_M/T$ is around 1 in the region of $2 < T/T_c < 5$, which is almost the same as the lattice result $m_M/T \sim 1$. It is worthy of
mentioning that \( m_E/T > m_M/T \) in the temperature region of \( T/T_c < 3 \) cannot be explained by using the perturbative scaling \( m_E \sim gT \) and \( m_M \sim g^2T \), because of the coupling constant \( g(T) > 1 \) in this region.

**B. Gauge dependence investigation**

![Figure 4: The gauge dependence of the screening mass as functions of the temperature.](image)

We have shown the screening masses by using the gluon propagator Eq. (15) in the unitary gauge, i.e. \( \xi \to \infty \). By fixing the model parameters \( \langle g^2G^2 \rangle = 0.009 \times 4\pi^2\text{GeV}^4 \) and \( \Lambda = 0.3\text{[GeV]}^{-1} \), in Fig. 4 we show the screening masses in different gauges, the solid lines are for the Landau gauge \( \xi = 0 \), the dash-dotted lines are for the Feynman gauge \( \xi = 1 \), and the dotted lines are for the unitary gauge \( \xi \to \infty \). It is found that below \( T = 500\text{MeV} \), the screening masses are independent on different gauges. The gauge dependence starts to show up when \( T > 500\text{MeV} \), the electric screening mass is more sensitive to the gauge fixing than the magnetic screening mass. In the temperature region we are interested in, both electric and magnetic screening masses are not sensitive to the gauge fixing.

**C. The Polyakov loop expectation value**

The deconfinement phase transition is characterized by the Polyakov-loop expectation value. The Polyakov-loop is defined as

\[
L(x) = \mathcal{P}\exp[i\int_0^\beta d\tau A_4(x, \tau)].
\]

(21)

In order to investigate the relation between the dimension-2 gluon condensate and the deconfinement phase transition, it is necessary to calculate the Polyakov-loop expectation value. By using perturbative expansion [44], it has been observed in Ref. [45] that the Polyakov loop expectation value is associated with the electric dimension-2 gluon condensate by the following relation:

\[
< L >= \exp[-\frac{g^2 < A_4^2 >}{4N_cT^2}].
\]

(22)

In our model, the electric dimension-2 gluon condensate has a simple relation with the electric screening mass square, i.e. \( < A_4^2 >= m_E^2 \).

We show the Polyakov loop expectation value as a function of \( T/T_c \) in Fig. 5 and compare the results with lattice data in Ref. [46]. It is found that the Polyakov loop expectation value is zero in the vacuum and low temperature region, it starts to rise at around \( 0.5T_c \), then rise sharply to a value of 0.8 at high temperature. We have taken \( T_c = T_0 = 150\text{MeV} \), where the electric and magnetic gluons start to split. It is worthy of mentioning that the susceptibility of the Polyakov loop expectation value indeed gives the critical temperature at around \( T_c = T_0 \). Our simple model indicates that the color electric deconfinement phase transition is driven by the electric gluons, and the nonperturbative dimension-2 gluon condensate plays an important role, it still gives at least 80% contribution to the Polyakov loop expectation value even at temperature region \( T > 3T_c \).
V. CONCLUSIONS

We have investigated the electric and magnetic screenings of the thermal gluons in a gluodynamic model with dimension-2 gluon condensate in zero momentum, which spontaneously generates the effect dynamical gluon mass in the vacuum.

It is found that the electric and magnetic gluons are degenerate at low temperature. With the increasing of temperature, the electric and magnetic gluons start to split at a certain temperature around $T_0 = 150$ MeV. The electric screening mass changes rapidly with temperature at $T > T_0$, and the Polyakov loop expectation value rises sharply around $T_0$ from zero in the vacuum to a value around 0.8 at high temperature. This suggests that the color electric deconfinement phase transition is driven by electric gluons. It is also observed that the magnetic screening mass keeps almost the same as its vacuum value, which manifests that the magnetic gluons remains confined. Both the screening masses and the Polyakov loop results are qualitatively in agreement with the Lattice calculations.

The Polyakov loop expectation value in this work is calculated by using the perturbation expansion, a more convenient way to derive the Polyakov loop expectation value is by using AdS/CFT method in the 5D holographic model, e.g. in Ref. [47] with a dimension-2 dilaton field background. It is worthy of mentioning that the dimension-2 dilaton field corresponds to a dimension-2 gluon condensate operator, and in Ref. [47], the Polyakov loop expectation value at finite temperature agrees well with the lattice data [46].

The model we used in this paper is quite simple, but it captures some important feature of gluon dynamics in the vacuum as well as in at finite temperature. We can conclude that the dimension-2 gluon condensate plays an essential role both in confinement as well as in deconfinement phase transition.

Acknowledgments: The authors thank valuable discussions with H. Chen, T. Hatsuda, T. Mukherjee, N. Su, Q.S. Yan. The work of M.H. is supported by CAS program "Outstanding young scientists abroad brought-in", CAS key project KJX2-EW-N01, NSFC under the number of 10735040 and 10875134, and K.C.Wong Education Foundation, Hong Kong.

[1] Y. Nambu, Phys. Rev. 117, 648-663 (1960).
[2] Y. Nambu, in Symmetries and Quark Models, ed. R. Chand, Gordon and Breach (1970); Y. Nambu, Phys. Rev. D 10, 4262 (1974); H. B. Nielsen, submitted to the 15th International Conference on High Energy Physics, Kiev (1970); H. B. Nielsen and P. Olesen, Nucl. Phys. B 61, 45 (1973); L. Susskind, Nuovo Cim. 69A (1970) 457.
[3] ’t Hooft, Nucl.Phys.B190.455(1981); Mandelstam, Phys.Rep.C23.245(1976).
[4] G. ’t Hooft, Nucl. Phys. B 138, 1 (1978).
[5] J. Greensite, Eur. Phys. J. ST 140, 1 (2007).
[6] A. M. Polyakov, Phys. Lett. B 72, 477 (1978).
[7] M. A. Shifman, A. I. Vainshtein, V. I. Zakharov, Nucl. Phys. B147, 385-447 (1979); B147, 448-518 (1979).
