Thermal aspects of dual QCD

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Abstract: The dual QCD version based on the magnetic symmetry structure of non-Abelian gauge theories, has been analyzed for its dual dynamics and thermal behaviour of QCD vacuum. During the thermal evolution, in addition to the restoration of magnetic symmetry at high temperatures, the gradual weakening of the confining force and the thermal evaporation of the magnetic glueballs, has been shown to indicate a smooth transition of hadronic system to the deconfined phase via a weakly bound QGP phase.

Though QCD is a commonly accepted gauge theory of strong interactions [1], the formulations of its dual version has equally important place mainly because of its highly non-perturbative nature in low energy regime [2, 3] and the associated complex phase structure. Furthermore, there has been a renewed interest [2–4] in the study of the hadronic matter under unusual conditions of high temperatures and densities mainly to establish and analyze the existence of the QGP and the detailed phase structure of QCD which is important from the point of view of not only high-energy heavy-ion collision events in modern high energy physics labs but from the evolutionary scenario of the universe also. In view of these facts, in the present paper, we have attempted analyze mainly the thermal response of a infrared effective magnetic symmetry based dual version of QCD [5].

In a $(4 + n)$-dimensional metric manifold (g_{AB}), introducing the magnetic symmetry as an additional isometry, the Killing condition for the Lie derivative along the magnetic vector leading to a gauge covariant magnetic symmetry condition, $D_\mu \hat{m} = 0$ ($\hat{m}$ being a scalar multiplet belonging to the adjoint representation of the gauge group G) leads to an exact solution for the gauge fields (with $G \equiv SU(2)$) and its little group $H \equiv U(1)$) as,

$$\mathbf{W}_\mu = A_\mu \hat{m} - g^{-1} (\hat{m} \times \partial_\mu \hat{m}), \quad (1)$$

where the second term on r.h.s. is of topological origin since the multiplet $\hat{m}$ may be viewed to define the homotopy of the mapping $\Pi_2(S^2)$ as, $\hat{m} : S^2_R \rightarrow S^2 = SU(2)/U(1)$, which identifies the monopoles of the non-Abelian gauge symmetry. The dual structure of and the dual dynamics between colour isocharges and topological charges becomes evident when we express the gauge potential (1) and corresponding gauge fields in magnetic gauge ($\hat{m} \xrightarrow{U} \xi_3 = (0,\ 0,\ 1)^T$) as, $\mathbf{W}_\mu \xrightarrow{U} (A_\mu + B_\mu)\hat{\xi}_3$ and $\mathbf{G}_{\mu\nu} = [\mathbf{W}_{\nu,\mu} - \mathbf{W}_{\mu,\nu} + g\,\mathbf{W}_\mu \times \mathbf{W}_\nu] \xrightarrow{U} (F_{\mu\nu} + B^{(d)}_{\mu\nu})\hat{\xi}_3$. The corresponding Lagrangian with built-in dual structure, in quenched approximation, may then be expressed as,

$$\pounds^{(m)}_d = -\frac{1}{4} B^2_{\mu\nu} + \left| \left[\partial_\mu + i\,4\pi\, g^{-1} B^{(d)}_\mu\right] \phi \right|^2 - 3\lambda\alpha_s^{-2}(\phi^*\phi - \phi_0^2)^2, \quad (2)$$

$\phi(x)$ being the complex scalar monopole field and $\alpha_s = g^2/4\pi$. The quadratic effective potential introduced here induces the dynamical breaking of magnetic symmetry which, in turn, leads the magnetic condensation of QCD vacuum and confines the color isocharges through dual Meissner effect. The associated flux-tube structure of dual QCD then leads to the inter-quark potential given by,

$$V(R) = C^2 \left[-\frac{\alpha_s}{16} \frac{\exp(-2m_B R)}{R} + \frac{\alpha_s m_B^2}{2} R \right], \quad (3)$$

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where, $m_B = (8\pi\alpha_s^{-1})^{\frac{1}{2}}\phi_0$ (the mass of magnetic glueballs appearing as the vector mode of magnetically condensed QCD vacuum) and the linear part is indicative of the confining features of dual QCD vacuum through its characteristic flux-tube structure.

In order to further explore the response of the multi-flux based dual QCD vacuum under extreme thermal conditions, let us start with the partition functional for the present dual QCD in thermal equilibrium at a constant temperature $T(\equiv \beta^{-1})$, given by,

$$Z[J] = \int D[\phi]D[B_\mu^{(d)}]\exp(i\int d^4x(\pounds - J|\phi|^2)). \quad (4)$$

Using the mean-field approach for separating the monopole field fluctuations ($\tilde{\phi}$) from its mean value (ϕ) as, $\phi \longrightarrow (\phi + \tilde{\phi})\exp(i\xi)$, identifying the quadratic source which retains the symmetry of the classical potential as, $J = -6\lambda\alpha_s^{-2}(\phi^2 - \phi_0^2)$ and using the functional Legendre transformations to calculate the effective action as, $S_{eff} = -i\ln Z[J] + \int d^4xJ\phi^2$, the thermodynamical potential may be computed using the Dolan and Jakiw approach [6] in the following form,

$$V(\phi,T) = 3\lambda\alpha_s^{-2}(\phi^2 - \phi_0^2)^2 - \frac{7}{90}\pi^2T^4 + \left(\frac{4\pi\alpha_s + \lambda}{2\alpha_s^2}\right)T^2\phi^2. \quad (5)$$

The minimization of $V(\phi,T)$ w.r.t. ϕ then leads to the value of T-dependent monopole condensate and magnetic glueball mass as,

$$<\phi>_0^{(T)} = \sqrt{\phi_0^2 - \left(\frac{4\pi\alpha_s + \lambda}{\lambda}\right)\frac{T^2}{12}},$$

$$m_B^{(T)} = \sqrt{8\pi\alpha_s^{-1}}\left[\phi_0^2 - \left(\frac{4\pi\alpha_s + \lambda}{\lambda}\right)\frac{T^2}{12}\right]^{\frac{1}{2}},$$

for $T < T_c$ and zero otherwise. It then leads to the disappearance of the QCD monopole condensate and consequent evaporation of magnetic glueballs indicating the restoration of magnetic symmetry and the deconfinement of quarks around a critical temperature given by,

$$T_c = 2\phi_0\sqrt{\frac{3}{4\pi\alpha_s + 1}}, \quad (6)$$

which leads to its values as 268 MeV, 221 MeV and 174 MeV for the values strong coupling (α_s) as 0.22, 0.47 and 0.96 respectively. Furthermore, the relative magnitude of the interquark confinement force at non-zero temperatures may also be obtained as,

$$\frac{F_c^T}{F_c^0} = \frac{(m_B^{(T)})^2}{(m_B^0)^2} = \left[1 - \left(\frac{T}{T_c}\right)^2\right], \quad (7)$$

which leads to a continuous decay of the confining force and the string tension at finite temperatures and its disappearance at $T_c$ultimately which implies a second-order phase transition. The Magnetic glueballs then evaporate and most of the monopole condensate turns into thermal monopoles [7]. The weakening of confining force is then expected to trigger the melting of flux-tubes leading to a weakly bound (QGP) phase before transiting to the completely deconfined phase.

References

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