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Magnetic Symmetry, Regge Trajectories, and the Linear Confinement in Dual QCD

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Using the magnetic symmetry structure of non-Abelian gauge theories, we analyze the flux tube formulation and its implications on the hadronic Regge trajectories and the confinement of color isocharges in magnetically condensed (with as well as without the electric excitations) QCD vacuum. Starting with the fiber bundle structure of QCD, the dual potentials are used to construct the QCD Lagrangian which has been shown to develop a unique flux tube configuration in its dynamically broken phase. The vector mass mode of the condensed vacuum has been shown to play a leading role in flux tube energy and other confinement parameters. Using the flux tube energy and the angular momentum, the Regge trajectories for hadrons have been obtained and the linear confining properties of dual QCD have been established. The dyonic flux tube structure of the condensed QCD vacuum has been obtained by inducing the electric excitation of QCD monopoles and the confining nature along with the linearity of Regge trajectories in dyonically condensed QCD vacuum are shown to remain intact. Implications of the modification in Regge slope parameter, on improving the confining properties of dual QCD vacuum are also discussed.

KEY WORDS: magnetic symmetry; monopoles/dyons; dual QCD; Regge trajectories; colour confinement.
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1. INTRODUCTION

In the history of the diversified ways of the study of the fundamental building blocks of matter and their interactions, the quark model of Gell-Mann and Zweig (Gell-Mann, 1962, 1964; Zweig) occupies a special position. Though, all hadrons are composed of quarks with three color degrees of freedom as seen by deep inelastic scattering experiments, the colored quarks are not the part of the physical spectrum and are permanently confined to the interior of hadrons. Quantum Chromodynamics (QCD) is a most popular modern gauge theory dealing with the

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strong interactions of hadrons (Gross and Wilczek, 1973; Politzer, 1973) and the confinement problem of QCD is the most challenging problem of the present day theoretical physics. This is mainly because of the highly nonperturbative nature of QCD in low-energy region where it essentially becomes a strong coupling gauge theory. Consequently, the infrared sector of QCD needs to be explored with some nontrivial techniques and, therefore, some effective models (Baker *et al.*, 1991; Pandey and Chandola, 2000; Suganuma *et al.*, 1998; Suzuki, 1988; 't Hooft, 1978) have been developed. In this connection, the dual formulation of QCD has played an important role in the explanation of the confinement mechanism and other low energy properties of hadrons by incorporating a direct analogy of the nature of QCD vacuum with that of the conventional superconductivity (Mandelstam, 1976, 1980; Nambu, 1974; Nielsen and Olesen, 1973). In such a dual picture, the colored monopole condensation plays the role analogous to the cooper-pair condensation in the conventional superconductivity. The idea has got further boost by 't Hooft's Abelian projection technique ('t Hooft, 1981, 1975) which indicates that the topological objects (viz, monopoles and dyons) (Julia and Zee, 1975; Prasad and Sommerfield, 1975) are the essential ingredients of the gauge theories to explore the superconducting nature of the QCD vacuum. Further, the lines of color electric flux in QCD are believed to be aligned to form a thin flux tube connecting the opposite color electric charges having a definite spin, flavor and momenta (Glendenning and Matsui, 1983; Parisi, 1975; Wyld, 1976; Wyld and Cutter, 1976). Recent lattice gauge theory (LGT) calculations (Creutz, 1983; Diacomo *et al.*, 2002; Rothe, 1992; Suzuki and Yotsuyanagi, 1990) also confirm the flux tube structure of the hadrons as well as the essential role of the topological objects in explaining the confinement mechanism. With the introduction of such topological objects in the dual QCD vacuum, the dual dynamics between color isocharges and topological charges can be best ascribed by magnetic symmetry (Cho, 1980) which paves the way to formulate the QCD as a viable dual gauge theory. Further, in view of the results of Witten and Schierholz (Schierholz, 1995; Witten, 1983) that the monopoles are necessarily dyons and the recent interest (Akhmedov, 1998; Gonzalez-Arroyo and Simonov, 1996; Simonov, 1996) in the models of QCD vacuum involving non-Abelian dyons, it is necessary to study the dyonic excitations also in the context of non-Abelian QCD formulations. In the present study, using the magnetic symmetry structure of QCD vacuum, we have investigated the flux tube formulation from the energy point of view and analyzed its implications over hadronic Regge trajectories and the confinement of color for the cases when the QCD vacuum undergoes the state of either pure monopole or dyonic condensation. Using the global topology of the gauge fields, the SU(2) chromodynamic Lagrangian has been constructed in terms of the dual potentials which has been shown to lead to the magnetically condensed state for the QCD vacuum in its dynamically broken phase. Flux tube structure of the resulting dual superconducting QCD vacuum has been analyzed and the confinement parameters are expressed in terms of the vector mass mode of the condensed vacuum. Deriving the flux tube energy, the Regge trajectories for hadrons have been obtained and the Regge slope parameter has been shown to depend on the glueball masses and the strong coupling in dual QCD. The electric excitation of condensed monopoles in dual QCD vacuum has been shown to lead it to the dyonic flux tube structure where the dual superconducting and the confining nature of dual QCD vacuum are shown to still remain intact. The linearity of the Regge trajectories in dyonically condensed vacuum is also maintained. However, the Regge slope parameter is modified as a result of the increase in dyonic flux tube energy. Its implications on improving the confining properties of dual QCD vacuum are discussed.

2. MAGNETIC SYMMETRY STRUCTURE AND SU(2) QCD LAGRANGIAN

In order to analyze the QCD as an infrared effective dual gauge theory, let us first briefly review the magnetic symmetry and the topological structure associated with the color gauge theory. The magnetic symmetry (Cho, 1980; Pandey *et al.*, 2001; Nandan *et al.*, 2002) is defined as an additional isometry of the internal fiber space in terms of a scalar multiplet described by a set of the self-consistent Killing vector fields belonging to the adjoint representation of a gauge group (G). The $(4 + n)$ dimensional unified space may be identified as a principal fiber bundle $P(M, G)$ over space–time if we identify the quotient space (P/G) as the base manifold (M) with the canonical projection $\Pi : P \rightarrow M$. For the simplest case of the gauge group $G = SU(2)$ with it's little group $H = U(1)$, the magnetic symmetry, which restricts the connection to those whose holonomy bundle becomes a reduced bundle, may be introduced in terms of a gauge-covariant condition given by

$$D_\mu \hat{m} = (\partial_\mu + g\mathbf{W}_\mu \times)\hat{m} = 0, \tag{1}$$

where $\hat{m}$ is a scalar multiplet which belongs to the adjoint representation of the gauge group G. The exact solution of Eq. (1) leads to the gauge potential $\mathbf{W}_\mu$ for $SU(2)$ gauge symmetry in the following form:

$$\mathbf{W}_\mu = A_\mu \hat{m} - g^{-1}(\hat{m} \times \partial_\mu \hat{m}), \tag{2}$$

where $\hat{m}.\mathbf{W}_\mu \equiv A_\mu$ is the color electric potential unrestricted by magnetic symmetry, while the second term on right hand side is completely determined by the magnetic symmetry and is topological in origin. The magnetic symmetry may, therefore, be used to describe the topological structure of the gauge theory and the scalar multiplet $\hat{m}$ may be viewed to define the homotopy of mapping $\Pi_2(S^2)$ as, $\hat{m} : {S_R}^2 \rightarrow S^2 = SU(2)/U(1)$. The associated field strength is then given by

$$\mathbf{G}_{\mu\nu} = \mathbf{W}_{\nu,\mu} - \mathbf{W}_{\mu,\nu} + g\mathbf{W}_\mu \times \mathbf{W}_\nu \equiv (F_{\mu\nu} + B_{\mu\nu})\hat{m}, \tag{3}$$

where $F_{\mu\nu} = A_{\nu,\mu} - A_{\mu,\nu}$ and $B_{\mu\nu} = -g^{-1}\hat{m}.(\partial_\mu \hat{m} \times \partial_\nu \hat{m})$. The dual symmetric separation of the gauge fields may be used to bring the topological structure of the theory into dynamics explicitly using the magnetic gauge obtained by rotating $\hat{m}$ to a prefixed space–time independent direction in isospace as, $\hat{m} \xrightarrow{U} \hat{\xi}_3 = (0, 0, 1)^T$. Using the simple parameterization, $\hat{m} = (\sin\alpha\cos\beta, \sin\alpha\sin\beta, \cos\alpha)^T$, and choosing, $U = \exp(-\alpha t_2 - \beta t_3)$, the gauge potential is obtained as

$$\mathbf{W}_\mu \xrightarrow{U} \mathbf{W}'_\mu = (A_\mu + B_\mu)\hat{\xi}_3, \tag{4}$$

where $B_\mu = g^{-1}\cos\alpha\partial_\mu\beta$. The associated gauge field strength in magnetic gauge is then given by

$$\mathbf{G}_{\mu\nu} \xrightarrow{U} \mathbf{G}'_{\mu\nu} \equiv (F_{\mu\nu} + B_{\mu\nu})\hat{\xi}_3, \tag{5}$$

which has the topological contribution as given by

$$B_{\mu\nu} = -g^{-1}\sin\alpha(\partial_\mu\alpha\partial_\nu\beta - \partial_\nu\alpha, \partial_\mu\beta) \equiv B_{\nu,\mu} - B_{\mu,\nu}. \tag{6}$$

Since the magnetic charge is of topological in origin, it represents the monopole field and therefore, in magnetic gauge, the topological properties of the magnetic symmetry are brought into dynamics explicitly. In order to study the physical implications of the associated dual potentials, we start with the gauge-invariant dual QCD Lagrangian for SU(2) gauge group with a quark doublet source ($\psi(x)$), given by

$$\pounds = -\frac{1}{4}|\mathbf{G}_{\mu\nu}|^2 + \bar{\psi}(x)i\gamma^\mu D_\mu\psi(x) - m_0\bar{\psi}(x)\psi(x). \tag{7}$$

In the Lagrangian given by Eq. (7), the topological objects appear as the point-like singular objects and not as a regular field. However, in order to avoid such undesirable features in the theory, one may use the dual magnetic potential $B_\mu^{(d)}$ and a complex scalar field (ϕ) at the same time for the topological object. A correct field-theoretical analysis of such a nontrivial QCD vacuum may then be obtained by SU(2) gauge-invariant Lagrangian which in the quenched approximation is given as

$$\pounds = -\frac{1}{4}\left|B_{\mu\nu}^{(d)}\right|^2 + \left|\left(\partial_\mu + i4\pi g^{-1}B_\mu^{(d)}\phi\right)\right|^2 V(\phi^*\phi), \tag{8}$$

where the field tensor $B_{\mu\nu}^{(d)} = \frac{1}{2}\epsilon_{\mu\nu\sigma\rho}B^{\sigma\rho} = B_{\nu,\mu}^{(d)} - B_{\mu,\nu}^{(d)}$. The Lagrangian (8) exactly coincides with that of the Ginzburg–Landau theory of superconductivity and may therefore be identified to generate the dynamical breaking of the magnetic symmetry of QCD vacuum through the effective potential $V(\phi^*\phi)$. The dynamical breaking of magnetic symmetry by the effective potential ultimately leads to the magnetic condensation of QCD vacuum which manifests itself in terms of the appearance of two (vector and scalar) mass modes of the condensed vacuum.

Such an effective potential is fixed by using the single-loop expansion technique (Coleman and Weinberg, 1973) along with the requirement of ultraviolet finiteness and infrared instability and is obtained in the following form:

$$V(\phi^*\phi) = \frac{24\pi^2}{g^4}\left[\phi_0^4 + (\phi^*\phi)^2\left(2\ln\frac{\phi^*\phi}{\phi_0^2} - 1\right)\right], \tag{9}$$

where ϕ_0 represents the vacuum expectation value of the monopole field. The field equations associated with the Lagrangian (8) may then also be derived in the following form:

$$\left(\partial^\mu - \frac{i4\pi}{g}B^{(d)\mu}\right)\left(\partial_\mu + \frac{i4\pi}{g}B_\mu^{(d)}\right)\phi + \frac{24\pi^2}{g^4}\left[4\phi\phi^*\ln\frac{\phi\phi^*}{\phi_0^2}\right]\phi = 0, \tag{10}$$

$$B_{\mu\nu,}^{(d)\nu} + \frac{i4\pi}{g}(\phi^*\partial_\mu\phi - \phi\partial_\mu\phi^*) - \frac{32\pi^2}{g^2}B_\mu^{(d)}\phi\phi^* = 0. \tag{11}$$

These equations govern the dynamics of QCD vacuum in its dynamically broken phase and lead to a definite flux tube structure to the dual QCD vacuum which may be shown to impart it the appropriate (linear) confining properties.

3. MAGNETIC CONDENSATION AND REGGE TRAJECTORIES IN DUAL QCD

In order to explain the nature of the magnetically condensed dual QCD vacuum, let us first analyze the field equations associated with the Lagrangian (8) under the cylindrical symmetry. Taking the flux tube orientation along the z-axis and keeping in view the uniqueness of the function $\phi(x)$, let us use the following cylindrically symmetric ansatz

$$\mathbf{B}^{(d)} = -\hat{\theta}B(\rho), \quad B_0^{(d)} = 0, \tag{12}$$

and

$$\phi(\rho) = \exp(i\,n\theta)\chi(\rho). \tag{13}$$

The coupled nonlinear field equations associated with the Lagrangian (8) in the static case may then be derived in the following form:

$$\frac{1}{\rho}\frac{d}{d\rho}\left(\rho\frac{d\chi}{d\rho}\right) - \left(\frac{n}{\rho} + \frac{4\pi}{g}B\right)^2\chi - \frac{24\pi^2}{g^4}\left(4\chi^2\ln\frac{\chi^2}{\phi_0^2}\right)\chi = 0, \tag{14}$$

$$\frac{d}{d\rho}\left[\frac{1}{\rho}\frac{d}{d\rho}(\rho B(\rho))\right] - \frac{8\pi}{g}\left(\frac{n}{\rho} + \frac{4\pi}{g}B\right)\chi^2 = 0, \tag{15}$$

where the primes refer to the differentiation with respect to the variable ρ. The appropriate boundary conditions for the finite energy vortex solutions are

$$B(\rho) = -\frac{ng}{4\pi\rho},\ B_0(\rho) = 0 \quad \text{and} \quad \chi(\rho) = \phi_0 \quad (f \text{ or } \rho \xrightarrow{\lim} \infty), \tag{16}$$

$$\chi(\rho) = B(\rho) = B_0(\rho) = 0 \quad (f \text{ or } \rho \xrightarrow{\lim} 0). \tag{17}$$

For analyzing the energy contribution associated with the dual QCD field configuration governed by the Eqs. (14) and (15), the energy per unit length of the flux tube may be calculated in the following form:

$$E_{(m)}(B,\chi) = 2\pi \int_0^\infty \rho d\rho \left[\left(\frac{1}{2\rho^2}(\rho B)'^2 + \chi'^2 \right) + \left(\frac{n}{\rho} + \frac{4\pi}{g} B \right)^2 \chi^2 \right.$$
$$\left. + \frac{24\pi^2}{g^4} \left(\phi_0^4 + \chi^4 \left(2 \ln \frac{\chi^2}{\phi_0^2} - 1 \right) \right) \right]. \tag{18}$$

The energy contribution, given by Eq. (18), gets minimized in its natural way with the boundary conditions given by Eq. (16) for the large-distances which enforces monopole condensation in the dual QCD vacuum and consequently the localization of the color electric flux in the form of thin flux tubes. For further simplification of the field Eqs. (14) and (15) and the flux tube energy in the the magnetically condensed QCD vacuum, let us introduce some new functions as given below.

$$r = \frac{8\pi}{g^2}\sqrt{3}\phi_0\rho, \tag{19}$$

$$K(r) = \frac{4\pi}{g}\rho B(\rho), \tag{20}$$

$$H(r) = \frac{\chi(\rho)}{\phi_0}, \tag{21}$$

which lead to the field equations or static configuration in the following form:

$$H'' + \frac{H'}{r} - \frac{(n+K)^2}{r} H - H^3 \ln H = 0, \tag{22}$$

$$K'' - \frac{K'}{r} - \Gamma(n+K)H^2 = 0, \tag{23}$$

where $\Gamma = \frac{2\pi}{3}\alpha_s$ and $\alpha_s = \frac{g^2}{4\pi}$ is the strong coupling constant. In the absence of any exact solution, the solutions of these equations in the asymptotic limit may be obtained by using the boundary conditions given by Eq. (16) in the following form:

$$H(r) = 1 - AK_0(r), \tag{24}$$

$$K(r) = -n + BrK_1\left(\sqrt{\frac{2\pi\alpha_s}{3}}\gamma\right), \tag{25}$$

where K_0 and K_1 are the modified Bessel's functions and A and B are the integration constants. Using Eq. (19), the solution given by Eq. (25) may be reexpressed as

$$K(\rho) = -n + C\rho^{\frac{1}{2}}\exp(-m_B\rho), \tag{26}$$

where C is a constant and $m_B = 4\sqrt{2}\pi g^{-1}\phi_0$ is the vector mass mode, the inverse of which leads to the penetration depth ($\lambda_{\text{QCD}}^{(d)} = m_B^{-1}$) of the color electric flux emanating from the color isocharges. The whole QCD vacuum, as a result of the dynamical breaking of magnetic symmetry, then acquires the flux tube structure which imparts it the perfect dual superconducting nature. The magnitude of the associated dual Meissner effect, responsible for the phase transition from normal to flux tube phase, is thus determined by the vector mass mode of the condensed vacuum. The energy content of the flux tube given by Eq. (18) may further be expressed in terms of functions defined by Eqs. (19)–(21) in the following form:

$$E_{(m)}(K, H) = I_{(m)}\phi_0^2 \equiv I_{(m)}\frac{\alpha_s}{8\pi}m_B^2, \tag{27}$$

where

$$I_{(m)} = \int 2\pi r\,dr\left[\frac{(K')^2}{r^2} + (H')^2 + \frac{1}{r^2}(K+n)^2H^2 + \frac{3}{16\pi\alpha_s}(1 + H^4(4\ln H - 1))\right]. \tag{28}$$

For the further analysis of the confinement mechanism in terms of the Regge trajectories and the associated confinement potential, let us use the expression (27) to obtain the contributions to the classical mass and the angular momentum of the color isocharges (quarks) as massless and spinless particles sitting at the opposite ends (r_1 and r_2) of the flux tube (with its length, $R = r_1 - r_2$), where each point at a distance R_0 from the center of the tube has a local velocity equal to $v = \frac{2R_0}{R}$ in natural system of units (Chew and Frautschi, 1961; Gasiorowicz and Rosner, 1981; Regge, 1959). With these considerations, the classical mass $M_{(m)}$ (relativistically) of the flux tube may be derived in the form given as

$$M_{(m)} = 2\int_0^{\frac{R}{2}} dR_0\frac{E_{(m)}(K, H)}{(1 - v^2)^{\frac{1}{2}}} = I_{(m)}\frac{\alpha_s}{16}m_B^2R. \tag{29}$$

Further, the total angular momentum of the flux tube is the angular momentum (orbital) without any spin contribution from quarks (as the oscillations caused by the monopole field in dual QCD vacuum are taken to be very small), and is given by

$$J^{(m)} = 2\int_0^{\frac{R}{2}} dR_0\frac{E_{(m)}(K, H)v}{(1 - v^2)^{\frac{1}{2}}} = \frac{\alpha_s}{64}m_B^2R^2. \tag{30}$$

Comparing Eqs. (29) and (30), we get the inter-relationship between the total angular momentum and the classical mass of the flux tube in the following form:

$$J^{(m)} = \frac{4}{I_{(m)}\alpha_s m_B^2} M_{(m)}^2. \tag{31}$$

It leads to the linear relationship between $J^{(m)}$ and $M_{(m)}^2$ in the present flux tube model and such trajectories followed by mesons may be identified as the Regge trajectories if we identify

$$\frac{4}{I_{(m)}\alpha_s m_B^2} M_{(m)}^2 = \alpha', \tag{32}$$

as the Regge slope parameter in the present dual QCD model. It demonstrates that the linear pattern of Regge trajectories in dual QCD collectively depends on the glueball masses and strong coupling which incidently has the running (increasing) behavior in low energy region. For the fixed glueball masses, it also predicts the appearance of the trajectories of lower slopes for the couplings in deep infrared sector which is expected to incorporate the mesons containing heavy quarks only. Further, the linearity of the relation (31) also confirms the linear confining nature of the associated confinement potential which depends on the dual gluon mass in a natural way.

4. DYONIC EXCITATIONS AND THE CONFINING STRUCTURE OF DUAL QCD

In view of the importance of the dyonic objects in current non-Abelian gauge theories (Akamedov, 1998; Gonzakz-Arroyo and Simonov, 1996; Julia and Zee, 1975; Prasad and Sommerfield, 1975; Schierholtz, 1995; Simonov, 1996; Witten, 1983) including QCD, let us extend the analysis presented in the previous section to the case where the topological object (monopole) carries an additional degree of freedom in terms of the nonvanishing temporal part of the gauge field. Retaining the temporal part $B_0^{(d)}$ in Eq. (12), in fact, excites the electric degree of freedom of the fundamental monopole which imparts it a definite nonzero electric charge and transforms to a dyonic object. Consequently, the dynamical breaking of the magnetic symmetry leads to the dyonic condensation of QCD vacuum which has definite bearing on the process of (color) confinement. Using $B_0^{(d)} = B_0(\rho)$ along with the ansatz given by Eqs. (12) and (13) in the previous section, the net electric charge associated with the resulting dyonic excitations is given by

$$Q_e = \int dS^i B_{0i} = 2{Q_m}^2 \int d^2x\, B_0 \chi^2, \tag{33}$$

where $Q_m = \frac{4\pi}{g}$ is the quantized magnetic charge. Further, the additional field equation, in addition to those given by Eq. (14) (with an additional term

$(4\pi g^{-1} B_0)^2 \chi)$ and (15), for nonvanishing B_0 is then obtained by using the Eqs. (10)–(13) in the following form:

$$\frac{1}{\rho}\frac{d}{d\rho}\left(\rho\frac{dB_0}{d\rho}\right) - \frac{32\pi^2}{g^2} B_0 \chi^2 = 0, \tag{34}$$

where the appropriate boundary condition for B_0 is given by

$$B_0(\rho) = 0, \quad \text{for} \quad \rho \to 0 \quad \text{or} \quad \infty. \tag{35}$$

In addition to the Nielsen–Olesen ansatz given by the Eqs. (20) and (21), taking the following ansatz for $B_0(\rho)$ in terms of a function $J(r)$,

$$B_0(\rho) = \frac{2\sqrt{3}}{g}\phi_0 J(r), \tag{36}$$

the field Eq. (34) may be expressed in a more simple form as

$$J'' + \frac{J'}{r} - \Gamma J H^2 = 0, \tag{37}$$

where the function $J(r)$ satisfies the conditions, $J(r) \to 0$ or ∞ for the stability reasons. This equation along with those given by Eqs. (22) and (23) then governs dynamics of the dyonically condensed QCD vacuum and they ultimately leads to the dyonic flux tube structure to the dual QCD vacuum in its dynamically broken phase. With these considerations, the energy per unit length of the dyonic flux tube in the condensed QCD vacuum may be derived in the following form:

$$E_{(D)}(K, H, J) = E_{(m)}(K, H) + E_{(e)}(J, H), \tag{38}$$

where $E_{(m)}(K, H)$ is given by Eq. (26) and $E_{(e)}(J, H)$ is given by

$$E_{(e)}(J, H) = 2\pi\phi_0^2 \int_0^\infty r\, dr\, (2J'^2 + J^2 H^2), \tag{39}$$

which is the additional contribution as a result of the excitation of the electric degree of freedom of the fundamental monopole. Using the Eq. (37), it may be further expressed as

$$\begin{aligned} E_{(D)}(K, H, J) &= [I_{(m)}(K, H) + I_{(e)}(J, H)]\phi_0^2 \\ &= [I_{(m)}(K, H) + I_{(e)}(J, H)]\frac{\alpha_s}{8\pi}\left(m_B^{(D)}\right)^2, \end{aligned} \tag{40}$$

where

$$I_{(e)} = 2\pi \int_0^\infty r\, dr\, (2J'^2 + J^2 H^2). \tag{41}$$

The mass $m_B^{(D)}$ here refers to the dyonic glueball mass which differs from its pure magnetic counterpart as the ϕ_0 now represents the VEV of the dyonic field. These considerations, in the present case, thus lead to the contribution to classical mass and the angular momentum of the dyonic flux tube with massless and spineless quarks at its opposite ends, in the following form:

$$M_{(D)} = [I_{(m)} + I_{(e)}]\frac{\alpha_s}{16}\left(m_B^{(D)}\right)^2 R, \tag{42}$$

$$J^{(D)} = \frac{\alpha_s}{64}\left(m_B^{(D)}\right)^2 R^2, \tag{43}$$

which leads to a linear relationship as

$$J^{(D)} = 4\frac{[I_{(m)} + I_{(e)}]^{-1}}{\alpha_s\left(m_B^D\right)^2} M_{(D)}^2. \tag{44}$$

It represents the Regge trajectories if the coefficient,

$$4\frac{[I_{(m)} + I_{(e)}]^{-1}}{\alpha_s\left(m_B^D\right)^2} M_{(D)}^2 = \beta, \tag{45}$$

is identified as the Regge slops parameter (RSP) in a conventional way. At first instance, it demonstrates the overall decreases in the RSP ($\beta < \alpha'$) as a result of the dyonic excitations over its pure magnetic counterpart. Consequently, when the dyonic glueball masses are not much different than the magnetic glueball masses, the dyonic excitation of the condensed QCD vacuum is expected to produce the confining potential so as to favor the mesons with heavy quark flavors. On the other hand, for maintaining both the RSP (α' and β) to its well-known value, i.e.,

$$\alpha' = \beta = 0.93\,\text{GeV}^{-2},$$

the Eq. (45) indicates that the change of the slope in dyonic case must be then compensated by an decrease in dyonic glueball mass which is then marked by an overall decrease in VEV of the dyonic field. Under such circumstances, the confinement parameters (flux penetration, coherence etc.) of the dyonically condensed QCD vacuum show an enhancement (as they are inversely proportional to the glueball masses) such that a more strong confinement is reflected with the dyonic excitations in condensed QCD vacuum. It demonstrates that the dyonic condensation of QCD vacuum is potentially capable of predicting new hadron trajectories on one hand and put a step forward to provide a mechanism for the absolute confinement of color in QCD on other.

5. CONCLUSIONS

Using the magnetic symmetry structure of QCD vacuum, the flux tube formulation has been developed without and with the excitation of the electric degree of freedom of the fundamental monopole. The dynamical breaking of the magnetic symmetry has been shown to leave the QCD vacuum in magnetically condensed dual superconducting state which imparts it the properties necessary for confining the color electric sources. In fact, in the strong-coupling limit, where the magnetic symmetry is dynamically broken, the confinement forces are set in and the system passes into the flux tube phase. It indicates that the quark pairs in such superconducting dual QCD vacuum joined by the flux filament of purely magnetic or dyonic nature, acquire a stable configuration (as there is no way for the flux to leak away) which has the behavior similar to that of the linearly rising hadron trajectories. The field equations given by Eqs. (22) and (23) and (37) along with (22) and (23) govern the flux tube structure for the cases of pure monopole and dyonic condensation respectively in dual QCD vacuum. Such flux tubes and the resulting confining nature of the dual QCD vacuum is identified by the appearance of the vector mass mode of the condensed vacuum which reflects itself in the form of the magnetic and dyonic glueballs in the two cases respectively. The flux tube energy derived in terms of the Eq. (27) for pure monopole condensed QCD vacuum case, leads to a linear relationship between angular momentum and the classical mass square as given by Eq. (31) which identifies the Regge slope parameter in terms of the parameters (α_s and m_B) of dual QCD vacuum as expressed by Eq. (32). It confirms the linear potential for the confinement of color isocharges in the theory. For the known light quark meson Regge trajectories, the value of slope parameter ($\approx 0.93\ \mathrm{GeV}^{-2}$) may be used to compute the glueball masses leading to the linear confining structure to dual QCD vacuum. On the other hand, the revival of the electric excitation of the fundamental monopole is shown to lead to the dyonically condensed QCD vacuum, the flux tube structure of which is governed by the additional Eq. (37) in association with Eqs. (22) and (23). The dyonic flux tube energy expression derived in terms of the Eqs. (38) and (40) shows an increase by an additional factor given by Eq. (39) when compared to its pure monopole counterpart. It leads to the slight decrease in Regge slope parameter expressed by Eq. (45) as a result of the electric excitation of the monopole. The effect is, in fact, just opposite to the case of the velocity dependent potential (which is important for mesons containing light quarks) where the slope of Regge tarjectories tends to increase. The increase in flux tube energy for the dyonically condensed QCD vacuum and the resulting decrease in the dyonic glueball masses for the fixed RSP, however tend to enhance the color electric flux penetration and condensed vacuum coherence parameters with the increase in confining potential strength in comparison to the pure monopole condensation of QCD vacuum. Further, the real analysis of the Regge slope parameter in dual QCD need to include the dynamical quarks which involve the contribution from the associated velocity dependent potentials also and shall be dealt separately.

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