No confinement without Coulomb confinement

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We compare the physical potential $V_D(R)$ of an external quark-antiquark pair in the representation $D$ of $SU(N)$, to the color-Coulomb potential $V_{coul}(R)$ which is the instantaneous part of the 44-component of the gluon propagator in Coulomb gauge $D_{44}(\vec{x},t) = V_{coul}(|\vec{x}|)\delta(t) + \text{(non-instantaneous)}$. We show that if $V_D(R)$ is confining, $\lim_{R \to \infty} V_D(R) = +\infty$, then the inequality $V_D(R) \leq -C_D V_{coul}(R)$ holds asymptotically at large $R$, where $C_D > 0$ is the Casimir in the representation $D$. This implies that $-V_{coul}(R)$ is also confining.
1. Introduction

The problem of confinement of color charge has been with us for a long time. There are many approaches to this problem such as dual Meissner effect by monopole condensation [1], effective string theory [2], center dominance [3], and color-Coulomb potential [4]. One seeks to choose variables so that the most important degrees of freedom have a simple expression. For this purpose gauge fixing can be a useful technique.

Confinement is most commonly characterized by the behavior of \( V_D(R) \), the gauge-invariant potential energy between an external quark pair at separation \( R \) in the representation \( D \) of the gauge structure group SU(N). It may be found from a rectangular Wilson loop \( W_D(R, T) \) in the representation \( D \), but for our purposes it is more convenient to obtain it from the correlator, \( \langle P_D(\vec{x}) P^*_D(\vec{y}) \rangle \), of a pair of Polyakov or thermal Wilson loops at \( \vec{x} \) and \( \vec{y} \) in the representation \( D \), on a Euclidean lattice of period \( T \) in the 4-direction. The Polyakov loop is the lattice analog of the continuum expression

\[
P_D(\vec{x}) = \text{tr} \left[ P \exp \left( \int_0^T A_{D,4}(\vec{x}, t) dt \right) \right],
\]

where \( A_{D,\mu} \equiv A_\mu t_D^a \), and the \( t_D^a \) satisfy the Lie algebra commutation relations \([t_D^a, t_D^b] = f^{abc} t_D^c\) in the representation \( D \). As discussed recently [2], this correlator has the expansion

\[
\langle P_D(\vec{x}) P^*_D(\vec{y}) \rangle = \sum_{n=0}^{\infty} \exp(-E_n,\vec{x},\vec{y}T), \tag{1.1}
\]

where the \( E_n,\vec{x},\vec{y} \) are the eigenvalues, \( H \Psi_n = E_n,\vec{x},\vec{y} \Psi_n \), of the lattice QCD hamiltonian \( H \), specified below, that includes an external quark and anti-quark at \( \vec{x} \) and \( \vec{y} \) in the representation \( D \). In the large-\( T \) limit, the sum is dominated by the first term, with lowest energy eigenvalue, \( E_{0,\vec{x},\vec{y}} \). It is rotationally symmetric in the continuum limit, and we identify the physical quark-antiquark potential \( V_D(R) \) with this energy eigenvalue, after separation of divergences,

\[
E_{0,|\vec{x}_1-\vec{x}_2|} = E_0 + \Delta + V_D(|\vec{x} - \vec{y}|). \tag{1.2}
\]

Here \( E_0 \) is the energy of the vacuum state in the absence of an external quark pair, and \( \Delta \) is the diverging self-energy of the external quarks. According to the Wilson confinement criterion, which is expected to hold in pure gluodynamics without dynamical quarks, but with external quarks in the fundamental representation, \( V_F(R) \) diverges linearly at large \( R \), \( V_F(R) \sim \sigma R \) where \( \sigma \) is the conventional string tension.
In the Coulomb gauge, there is a simple scenario [4] that attributes confinement of color charge to the long range of the color-Coulomb potential, $V_{\text{coul}}(R)$. This quantity characterizes the instantaneous part of the 44-component of the gluon propagator $\langle A_4^a(\vec{x},t)A_4^b(0,0)\rangle = D_{44}(\vec{x},t)\delta^{ab} = V_{\text{coul}}(|\vec{x}|)\delta(t)\delta^{ab} + \text{(non-instantaneous)}$. Since $A_4$ couples universally to the color-charge, this can account for confinement of color-charge, provided that $V_{\text{coul}}(R)$ is indeed long range. A remarkable feature of the Coulomb gauge in QCD, a property not shared by any Lorentz gauge, is that $A_4 = g_0 A_4^{(0)} = g_r A_4^{(r)}$ is a renormalization-group invariant [4], [5]. Here $g_0$ and $A_4^{(0)}$, and $g_r$ and $A_4^{(r)}$ are, respectively, the unrenormalized and renormalized charges and perturbative gauge connections. This means that $D_{44}$, and hence also its instantaneous part $V_{\text{coul}}(R)$, is independent of both the cut-off $\Lambda$ and the renormalization mass $\mu$. This property allows the fundamental QCD quantity, $V_{\text{coul}}(R)$, the instantaneous part of the gluon propagator, to be identified with the phenomenological QCD potential [6] and [7]. Its fourier transform, $\tilde{V}_{\text{coul}}(|\vec{k}|)$, provides a convenient definition of the running coupling constant, $\alpha_s(|\vec{k}|/\Lambda_{\text{coul}}) = g_{\text{coul}}^2(|\vec{k}|)/(4\pi) = \tilde{k}^2\tilde{V}_{\text{coul}}(|\vec{k}|)/(4\pi x_0)$, where $x_0 = 12N/(11N - N_f)$, $\Lambda_{\text{coul}}$ is a finite QCD mass scale, and $N_f$ is the number of quark flavors [8]. The result obtained here means that if the Wilson criterion for confinement is satisfied, then, with this definition, the running coupling constant $\alpha_s(|\vec{k}|/\Lambda_{\text{coul}})$ diverges in the infrared like $1/\tilde{k}^2$, a clear manifestation of infrared slavery.

We shall show that a necessary condition for confinement according to the Wilson criterion is that the instantaneous color-Coulomb potential be confining. In symbols: if $\lim_{R \to \infty} V_D(R) = +\infty$, then $V_D(R) \leq -C_D V_{\text{coul}}(R)$ holds asymptotically at large $R$. The minus sign occurs because antiquark has opposite charge to quark. Here $C_D = -\sum_a (t_D^a)^2 > 0$ is the (positive) value of the Casimir invariant, and in the fundamental representation $C_F = (N^2 - 1)/(2N)$. A considerable simplification is hereby achieved because $V_D(R)$ is defined by means of a path-ordered exponential that involves gluon n-point functions of all orders, whereas $V_{\text{coul}}(R)$ is defined in terms of the gluon 2-point function, $D_{44}$. We also note the striking numerical result [9], [10] that $V_D(R)$ exhibits Casimir scaling, $V_D(R) = (C_D/C_F)V_F(R)$ quite accurately at least in a rather large range of $r$ and 8 representations $D$. This suggests that the above bound may be saturated in this range, for this would explain Casimir scaling, that is not easy to understand otherwise [11]. The result also makes it imperative to extend present programs to calculate $V_{\text{coul}}(R)$ numerically [12], and analytically from first principles [7], and to derive phenomenological quantities from it [6].
2. Lattice Coulomb-gauge QCD hamiltonian

The energy \( E_0, |\mathbf{x}_1 - \mathbf{x}_2| \) is of course gauge invariant, and the lattice QCD hamiltonian \( H \) may be chosen in any gauge. Its most familiar form is in the temporal gauge \( U_4 = 1 \), corresponding to \( A_4 = 0 \),

\[
H_{\text{temp}} = g_0^2(2a)^{-1} \sum_{\mathbf{x},i} \mathcal{E}_{\mathbf{x},i}^2 + 2(g_0^2a)^{-1} \sum_p \text{Re Tr} U_p,
\]

where \( \sum_p \) is the sum over all spatial plaquettes \( p \) (on a single time-slice). Here \( \mathcal{E}_{\mathbf{x},i} \) is the color-electric field operator that satisfies

\[
[\mathcal{E}_{\mathbf{x},i}^a, U_{\mathbf{y},j}^b] = i\delta^{ab}\delta_{\mathbf{x},\mathbf{y}}\delta_{ij} \text{ and } [\mathcal{E}_{\mathbf{x},i}^a, \mathcal{E}_{\mathbf{y},j}^b] = -if^{abc}\delta_{\mathbf{x},\mathbf{y}}\delta_{ij}\mathcal{E}_{\mathbf{x},i}^c. \]

We place an external quark at \( \mathbf{x}_1 \) in the representation \( D \), and an external anti-quark at \( \mathbf{x}_2 \) in the representation \( D^* \), with color vectors that act on the first and second indices of the wave-functional \( \Psi_{\alpha\beta}(U) \) according to \( (\lambda^a_1 \Psi)_{\alpha\beta} = (\lambda^a_D)_{\alpha\gamma}\Psi_{\gamma\beta} \) and \( (\lambda^a_2 \Psi)_{\alpha\beta} = -((\lambda^a_D)^*_{\beta\gamma}\Psi_{\alpha\gamma}) \), where \( \lambda^a_D = it^a_D \). In the temporal gauge, the color charges of the external quarks do not appear in the hamiltonian \( H_{\text{temp}} \), but rather in the subsidiary condition \( G^a(\mathbf{x})\Psi = 0 \). This is an expression of Gauss’s law, for \( G^a(\mathbf{x}) \) is a precise lattice analog of the continuum Gauss’s law operator \( G^a(\mathbf{x}) = -\left(\vec{D} \cdot \vec{E}\right)^a(\mathbf{x}) + \rho^a_{\text{qu}}(\mathbf{x}) \), where \( E_i^a(\mathbf{x}) = i\delta_{\mathbf{x},\mathbf{y}} \vec{\delta} A^a_i(\mathbf{x}) \), \( D^{ac}_i = \delta^{ac}\partial_i + f^{abc} A^b_i \) is the gauge-covariant derivative, and \( \rho^a_{\text{qu}}(\mathbf{x}) = \lambda^a_1\delta(\mathbf{x} - \mathbf{x}_1) + \lambda^a_2\delta(\mathbf{x} - \mathbf{x}_2) \) is the color-charge density of the external quarks. In the temporal gauge \( G^a(\mathbf{x}) \) is the generator of 3-dimensionally local gauge transformations of the quark and gluon variables, satisfying \([G^a(\mathbf{x}), G^b(\mathbf{y})] = i\delta(\mathbf{x} - \mathbf{y})f^{abc}G^c(\mathbf{x})\), and the subsidiary condition is the statement of gauge invariance of the wave functional.

One would expect that Gauss’s law is essential for confinement, and the lattice Coulomb hamiltonian \( H_{\text{coul}} \) [13] may be derived from \( H_{\text{temp}} \) by solving Gauss’s law as subsidiary condition [14]. For our purposes the resulting lattice Coulomb hamiltonian has the same structure to the continuum Coulomb hamiltonian [15]. To simplify the exposition, we shall use continuum language, but it is understood that this is short-hand for the correct lattice kinematics, and divergences are controlled by use of the lattice Coulomb hamiltonian, as will be made clear.

To get to the Coulomb gauge from the temporal gauge, one integrates out the gauge degrees of freedom using the Faddeev-Popov formula in all gauge-invariant matrix elements. In particular for the hamiltonian, one obtains \( H_{\text{coul}} \), defined by its matrix elements

\[
(\Psi_1, H_{\text{coul}}\Psi_2) = \int_{\Lambda} dA^{\text{tr}} \det M (1/2) \int d^3x \left[ g_0^2(E^a_i\Psi_1)^*E^a_i\Psi_2 + g_0^{-2}\Psi_1^*\vec{B}^2\Psi_2 \right],
\]
there the wave-functionals $\Psi_{\alpha\beta}(A^{\text{tr}})$ depend only on 3-dimensionally transverse continuum configurations $\partial_i A_{i}^{\text{tr}} = 0$, and a contraction on color indices is understood. The color-magnetic field is given by $B_i^a = \partial_2 A_3^{\text{tr},a} - \partial_3 A_2^{\text{tr},a} + f^{abc} A_2^{\text{tr},b} A_3^{\text{tr},c}$, etc., and the color-electric field by $E_i^a = E_i^{\text{tr},a} - \partial_i \phi^a$, where $E_i^{\text{tr},a} = i \frac{\delta}{\delta A_i^{\text{tr},a}}$, and $\phi^a(\vec{x})$ is the color-Coulomb potential operator. In this matrix element, $\phi^a(\vec{x})$ acts directly on the wave functional $\Psi$. The definition of $H_{\text{coul}}$ is completed by specifying that $\phi^a(\vec{x}) \Psi \equiv (M^{-1} \rho_{\text{phys}})^a(\vec{x}) \Psi$, which expresses $\phi^a(\vec{x}) \Psi$ in terms $\rho_{\text{qu}}$ and transverse gluon variables only. This is the solution of the subsidiary condition $G^a(\vec{x}) \Psi = 0$, or $M^{ac}(A^{\text{tr}}) \phi^c \Psi = \rho_{\text{phys}}^a \Psi$. Here $M^{ac}(A^{\text{tr}}) \equiv -D_i^{ac}(A^{\text{tr}}) \partial_i = -\partial_i D_i^{ac}(A^{\text{tr}}) = -\partial^2 \delta^{ac} - f^{abc} A_i^{\text{tr},b} \partial_i$ is the 3-dimensional Faddeev-Popov operator, and $\rho_{\text{phys}}^a \equiv -f^{abc} A_i^{\text{tr},b} E_i^{\text{tr},c} + \rho_{\text{qu}}^a$ is the color-charge density of the external quarks plus the color-charge of the dynamical gluon degrees of freedom only. The associated color charge, $Q^a = \int d^3 x \, \rho_{\text{phys}}^a(\vec{x})$ may be identified with the physical color charge, for it generates global gauge transformations on all variables $[Q^a, A_i^{\text{tr},b}] = if^{abc} A_i^{\text{tr},c}$ etc., and satisfies $[Q^a, Q^b] = if^{abc} Q^c$. The second term in $M$ is characteristic of non-Abelian gauge theory. It is responsible for anti-screening because, for typical configurations, this term produces a small denominator in $M^{-1}$. The subscript $\Lambda$ on the integral $\int_\Lambda dA^{\text{tr}}$ means that a region that includes only one Gribov copy is integrated over. This may be chosen as in the minimal Coulomb gauge, but the proof does not depend on the particular way this is chosen.

3. Bound on $V_D(R)$ from trial wave function

The energy $E_{|\vec{x}_1 - \vec{x}_2|} \equiv \langle \Psi, H_{\text{coul}} \Psi \rangle$ of any trial wave function $\Psi$ provides an upper bound on the ground-state energy, $E_0_{|\vec{x}_1 - \vec{x}_2|} \leq E_{|\vec{x}_1 - \vec{x}_2|}$. As trial function we take the product wave function, $\Psi_{\alpha\beta}(A^{\text{tr}}) = N^{-1/2}_{\text{D}} \delta_{\alpha\beta} \Phi_0(A^{\text{tr}})$. Here $\Phi_0(A^{\text{tr}})$ is the exact wave functional of the vacuum state in the absence of external quarks, and $N^{-1/2}_{\text{D}} \delta_{\alpha\beta}$, where $N_{\text{D}}$ is the dimension of representation $D$, is the external quark-pair state of total color-charge zero, $(\lambda_1 + \lambda_2)^a \Psi = 0$. The Coulomb hamiltonian has the decomposition $H_{\text{coul}} = H_{\text{gl}} + H_{\text{gl,qu}} + H_{\text{qu,qu}}$, that follows from the decomposition of the color-charge density $\rho_{\text{phys}}^a = -f^{abc} A_i^{\text{tr},b} E_i^{\text{tr},c} + \rho_{\text{qu}}^a$ in $\phi \Psi = M^{-1} \rho_{\text{phys}} \Psi$. Here $H_{\text{gl}}$ is the Coulomb hamiltonian in the absence of external quarks, $H_{\text{gl,qu}}$ is linear in $\rho_{\text{qu}}$, and $H_{\text{qu,qu}} = (1/2) \int d^3 x \, (\partial_i M^{-1} \rho_{\text{qu}})^2(\vec{x})$. In the last expression there is no ordering problem because $\rho_{\text{qu}}$ commutes with $A^{\text{tr}}$. By definition of $\Phi_0$ we have $H_{\text{gl}} \Phi_0 = E_0 \Phi_0$, where
$E_0$ is the vacuum energy in the absence of external quarks. From $(\Psi, H_{\text{gl}}\Psi) = E_0$, and $(\Psi, H_{\text{gl,qu}}\Psi) = 0$, we get for the trial energy,

$$E_{|\vec{x}_1 - \vec{x}_2|} = E_0 + \Delta' - C_D(\Phi_0, [M^{-1}(-\partial^2)M^{-1}]^a_{\vec{x}_1, \vec{x}_2} \Phi_0) \quad (3.1)$$

(no sum on $a$), where we have used $\langle \lambda^n_1 \lambda^n_2 \rangle = (N^2 - 1)^{-1} \delta^{ab} \langle \lambda^n_a \lambda^n_b \rangle = -(N^2 - 1)^{-1} \delta^{ab} C_D$. Here $\Delta'$ is another self energy of the external quarks that is independent of $\vec{x}_1$ and $\vec{x}_2$. The inequality $E_{0,|\vec{x}_1 - \vec{x}_2|} \leq E_{|\vec{x}_1 - \vec{x}_2|}$ reads $\Delta(\Lambda) + V_D(R, \Lambda) \leq \Delta'(\Lambda) - C_D(\Phi_0, [M^{-1}(-\partial^2)M^{-1}]^a_{\vec{x}_1, \vec{x}_2} \Phi_0)$. We have cancelled the vacuum energy $E_0$ that diverges with the volume of space and, having done so, we may take the volume of space to infinity, keeping the ultraviolet cut-off $\Lambda = a^{-1}$ in place, where $a$ is the lattice spacing. Here $\Delta(\Lambda)$ and $\Delta'(\Lambda)$ are self-energies of the external quarks. The formula [8] for the color-Coulomb potential, $V_{\text{coul}}(|\vec{x}_1 - \vec{x}_2|) \delta^{ab} = (\Phi_0, [M^{-1}(-\partial^2)M^{-1}]^a_{\vec{x}_1, \vec{x}_2} \Phi_0)$, allows us to write the inequality as $\Delta(\Lambda) + V_D(R, \Lambda) \leq \Delta'(\Lambda) - C_D V_{\text{coul}}(R, \Lambda)$.

We have inserted a dependence on the cut-off $\Lambda$ in $V_D(R, \Lambda)$ and $V_{\text{coul}}(R, \Lambda)$, because these are lattice quantities that depend on the lattice spacing, $a = \Lambda^{-1}$. However, having separated out the self-energies, both the quark potential $V_D(R, \Lambda)$ and the lattice color-Coulomb potential $V_{\text{coul}}(R, \Lambda)$ have finite, $\Lambda$-independent continuum limits, $\lim_{\Lambda \to \infty} V_D(R, \Lambda) = V_D(R)$ and $\lim_{\Lambda \to \infty} V_{\text{coul}}(R, \Lambda) = V_{\text{coul}}(R)$. For $V_D(R)$ is a physical energy, and $V_{\text{coul}}(R)$ is a renormalization-group invariant, as noted in the Introduction. If $V_D(R, \Lambda)$ is confining, $\lim_{R \to \infty} V_D(R, \Lambda) = +\infty$, then, for sufficiently large $R$, the self-energies $\Delta(\Lambda)$ and $\Delta'(\Lambda)$ are negligible compared to $V_D(R, \Lambda)$, and the inequality $V_D(R, \Lambda) \leq - C_D V_{\text{coul}}(R, \Lambda)$ holds asymptotically at large $R$, for finite cut-off $\Lambda$. This bound also holds in the continuum limit, because dimensional and renormalization-group considerations tell us that the terms that vanish as $\Lambda \to \infty$ are of relative order $1/(\Lambda R)^n$, where $n$ is positive, so they also vanish asymptotically at large $R$.\footnote{This condition is necessary, as shown by the following counter-example. Take $V_D(R, \Lambda) = \sigma R$, and $V_{\text{coul}}(R, \Lambda) = c/R + m^4 R^2/\Lambda$, with self-energies $\Delta(\Lambda) = a \Lambda$, and $\Delta'(\Lambda) = (a + 1) \Lambda$. The inequality at finite $\Lambda$ reads $\sigma R \leq \Lambda + c/R + m^4 R^2/\Lambda = 2m^2 R + c/R + (\Lambda - m^2 R)^2/\Lambda$. It is satisfied for all finite $\Lambda$ and $R$, provided that $\sigma \leq 2m^2$ and $c \geq 0$. But the continuum limit of $V_{\text{coul}}(R)$ is $c/R$, and $\sigma R < c/R$ is not satisfied at large $R$.} We conclude that, if $V_D(R)$ is confining, $\lim_{R \to \infty} V_D(R) = +\infty$, then in the continuum limit, the inequality, $V_D(R) \leq - C_D V_{\text{coul}}(R)$, holds asymptotically at large $R$, as asserted.
4. Conclusion

The bound implies that if the potential between external quarks in the fundamental representation increases linearly at large $R$, $V_F(r) \sim \sigma R$, where $\sigma$ is the standard string tension, then the color-Coulomb potential $V_{\text{coul}}(R)$ increases at least linearly at large $R$, and moreover if its increase is also linear, $-V_{\text{coul}}(R) \sim \sigma_{\text{coul}} R$, as has been conjectured [4], where $\sigma_{\text{coul}}$ is a string tension that characterizes $V_{\text{coul}}(R)$, then the conventional string tension satisfies the bound $\sigma \leq (N^2 - 1)/(2N)\sigma_{\text{coul}}$.

What has been learned about QCD dynamics? We have found that if the Wilson confinement criterion holds, then $V_{\text{coul}}(R)$, the instantaneous part of the gluon propagator $D_{44}$ in Coulomb gauge, is confining. Moreover, from $V_{\text{coul}}(|\vec{x} - \vec{y}|) = \langle [M^{-1}(-\partial^2)M^{-1}]_{\vec{x},\vec{y}} \rangle$, this can happen only if the Faddeev-Popov or ghost Green function, $[M^{-1}(A^{tr})]_{\vec{x},\vec{y}}$, is long-range for configurations $A^{tr}$ that dominate the functional integral. This confirms the confinement scenario originally proposed by Gribov [16], and advocated by the author [4]. [The scenario reads, in brief, that in the minimal Coulomb gauge, configurations are restricted to the Gribov region, where the Faddeev-Popov operator is positive, $M(A^{tr}) > 0$. The boundary of this region occurs where the lowest eigenvalue of $M(A^{tr})$ vanishes, $\lambda_0(A^{tr}) = 0$. Moreover the dimension $n$ of configuration space is very large, being of the order of the volume $V$ of the lattice. Entropy favors a population highly concentrated close to this boundary, where $\lambda_0(A^{tr})$ is small, for the same reason that, in a space of very high dimension $n$, the density of a ball $r < r_0$ is very sharply peaked near $r_0$, being given by $r^{n-1}dr$. Consequently, for the configurations that dominate the functional integral, $M^{-1}(A^{tr})$ is enhanced, and thus also $V_{\text{coul}}(R)$.]}

To simplify the exposition, we considered gluodynamics without dynamical quarks, but the proof holds if they are included. However if dynamical quarks are present in the fundamental representation $F$, as occurs in nature, then the physical potential $V_F(R)$ between external quarks does not manifest confinement. For at some radius $R_b$ the string breaks by polarization of sea quarks from the vacuum, and for $R > R_b$, $V_F(R)$ represents a residual potential between a pair of mesons, analogous to the van der Waals potential. In this case the bound obtained here does not imply that $V_{\text{coul}}(R)$ is confining. Nevertheless, according to the confinement scenario in Coulomb gauge, $V_{\text{coul}}(R)$ is a fundamental quantity that remains linearly rising even when $V_F(R)$ is not. It is precisely the linear rise of $V_{\text{coul}}(R)$ that causes string-breaking, by making it energetically preferable to polarize sea-quarks from the vacuum.
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References

[1] S. Mandelstam, Phys. Rep. 23C, 245 (1976).
[2] M. Lüscher and P. Weisz, Quark confinement and the bosonic string, hep-lat/0207003.
[3] L. Del Debbio, M. Faber, J. Giedt, J. Greensite, and S. Olejnik Detection of Center Vortices in the Lattice Yang-Mills Vacuum, Phys. Rev. D58 (1998) 094501.
[4] D. Zwanziger, Renormalization in the Coulomb gauge and order parameter for confinement in QCD, Nucl. Phys. B 518 (1998) 237.
[5] L. Baulieu, D. Zwanziger, Renormalizable Non-Covariant Gauges and Coulomb Gauge Limit, Nucl.Phys. B 548 (1999) 527, hep-th/9807024.
[6] Adam Szczepaniak, Eric S. Swanson, Chueng-Ryong Ji, Stephen R. Cotanch, Glueball Spectroscopy in a Relativistic Many-Body Approach to Hadron Structure, Phys. Rev. Lett. 76 (1996) 2011-2014.
[7] A. P. Szczepaniak, E. S. Swanson, Coulomb gauge QCD, confinement and the constituent representation, Phys.Rev. D65:025012 (2002).
[8] Attilio Cucchieri, Daniel Zwanziger, Renormalization-group calculation of the color-Coulomb, Phys. Rev. D65 (2001) 014002.
[9] G. S. Bali, Casimir scaling of SU(3) static monopoles, hep-lat/0006022.
[10] S. Deldar, Static SU(3) potentials for sources in various representations, hep-lat/9911008.
[11] V. I. Shevchenko, Yu. A. Simonov, Casimir scaling as a test of QCD vacuum, hep-ph/0001299.
[12] Attilio Cucchieri, Daniel Zwanziger, Gluon propagator and confinement scenario in Coulomb gauge, hep-lat/0209068.
[13] D. Zwanziger, Lattice Coulomb hamiltonian and static color-Coulomb field, Nucl. Phys. B 485 (1997) 185.
[14] D. Zwanziger, Continuum and Lattice Coulomb-gauge hamiltonian, in Cambridge 1997, Confinement, duality, and non-perturbative aspects of QCD, P. van Baal, Ed. hep-th/9710157.
[15] N. Christ and T. D. Lee, Phys. Rev. D22 (1980) 939.
[16] V.N. Gribov, Nucl. Phys. B 139 (1978) 1.