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Quark-quark interaction with correction from nonperturbative QCD

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The leading nonperturbative correction to the gluon propagator is derived. The Breit-Fermi-type one-gluon exchange potential with the above-mentioned nonperturbative correction is formulated in the nonrelativistic scheme. The result shows that, except for the regular color Coulomb, color electric, and color magnetic terms, both linear confinement and r^3 deconfinement terms can automatically be obtained.

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I. INTRODUCTION

Since QCD was introduced by particle physicists, the assumption that the baryon consists of quarks and gluons has widely been accepted. It is obvious that the traditional baryon-baryon interaction based on the one-boson-exchange theory, various form factors, and the phenomenological repulsive core should be explained in this more fundamental point of view. In the past 10 years, research into this aspect has extensively been carried out. These investigations can roughly be cast into two groups. One of them is based on the quark potential model [1–5].

Although most of the research into the quark potential model is in the nonrelativistic scheme, the results turn out quite well, especially for the hadron spectrum [6] and baryon-baryon scattering [4,5]. There are several reasons. The effect of c.m. motion can be correctly treated. The Breit-Fermi term in the Hamiltonian contains part of the relativistic treatment. Particularly, the strong-coupling constant α_s, which is slightly greater than 1, may include further relativistic and other nonperturbative effects.

However, something is still not clear. In the Hamiltonian of the hadron, one has to write the confinement term phenomenologically. In the baryon-baryon scattering case, the situation is even complicated. It is common knowledge that at long distances the interaction between two hadrons is governed by the pion-exchange mechanism. From the view point of field theory, the effective pion field or Skyrme field [7–10] is an effective theory where no quark and gluon degrees of freedom show up. Because of the finite size of the hadron, at short or even medium distance, two hadrons have more overlap, and quarks and gluons should present their contributions. In fact, there are lots of discussion about this subject in particle physics, the chiral bag model being just one of them [11]. Therefore, one should derive an effective baryon-baryon potential which is applicable not only for the large but also medium and short separations between two hadrons.

For the long- and medium-range baryon-baryon interactions, Yu, Zhang, and Shen [12] used the mechanism of quark-antiquark pair creation via one-gluon exchange and obtained effective one-boson-exchange potentials, which are very close to those obtained from meson-exchange theory, and the effective two-pion-exchange potential, which is weaker than the phenomenological σ-meson-exchange potential. For the short-range part, Oka and Yazaki [1] and Faessler *et al.* [2] gave the repulsive core from the one-gluon-exchange model. But all these results are based on the operator derived in the perturbative QCD scheme which is only applicable for the case of asymptotic freedom. It is clear that there should be some "medium- and long-distance QCD" effects. They may play an important role, especially in the scattering process. Unfortunately, these effects, at present, are not fully understood. There is not any satisfactory way to solve this problem. However, one still can put in some nonperturbative effects phenomenologically [13]. Many authors have studied the full propagators of the gluon and quark from QCD sum rules [14–16].

Our strategy is that beside the potential terms caused by the perturbative QCD, we bring the nonperturbative effect, namely, the nonzero quark and gluon condensates, into the interaction between quarks; then we can obtain an effective potential which contains not only the pertur-

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bative contribution, but also the nonperturbative effect.

As the first attempt, we just take the terms in the first order next to the perturbative leading order; in other words, just as in Ref. [16], there is only one condensate, no matter whether it is a gluon condensate or a quark condensate, in the full gluon propagator. In the next section, a brief formulation is given, and the results are discussed in Sec. III.

II. BRIEF FORMULATION

Years ago, Shifman, Vainstein, and Zakharov (SVZ) [13] introduced the nonzero quark condensation $m_q\langle 0|q\bar{q}|0\rangle$ and gluon condensation $(\alpha_s/\pi)\langle 0|G^a_{\mu\nu}G^{a\mu\nu}|0\rangle$, where $\langle 0|\cdots|0\rangle$ denotes the vacuum expectation value of the operator. In perturbative field theory, these condensations are zero, whereas because of a complicated QCD vacuum structure, SVZ phenomenologically introduced the nonzero vacuum condensate as parameters. The vacuum expectation value $\langle q\bar{q}\rangle$ means that a pair of quarks with zero energy and momentum condenses into vacuum, whereas $\langle GG\rangle$ has the same meaning for a pair of gluons. Since SVZ's work, many possible applications of QCD sum rules have been suggested [17].

Our opinion is that such condensations might give a substantial contribution to the effective medium-range potential.

First, we ignore loop calculations and assume that the one-gluon exchange between two quarks can give general properties; i.e., we do not consider ladder diagrams. Moreover, as the gluon travels along a longer distance, there may be more $q\bar{q}$ or GG pairs produced. We presume that one-pair production is more important than that with more condensating pairs, because it at least reflects a first-order approximation. In fact, we consider a correlator

$$\langle 0|J^a_\mu(x)J^a_\nu(0)|0\rangle \ , \tag{1}$$

where $J^a_\mu(x)=\bar{q}_i(x)\gamma^\mu(\lambda^a_{ij}/2)q_j(x)$. Graphically, we show several diagrams in Fig. 1.

In this work we omit all one-loop calculations, such as the quark self-energy and gluon vacuum polarization, as well as other vertex corrections. It is easy to see that Fig. 1(a) is the contribution from perturbative QCD, i.e., a Coulomb-type potential; Fig. 1(b) does not give a contribution, because the momentum which flows into or out of the vacuum must be zero; and Fig. 1(f) must contribute nothing, because the vacuum cannot possess any quantum number such as color. Then the remaining thing to be done is to calculate contributions from Figs. 1(c)–1(e). These calculations are straightforward, but very tedious. We use the SVZ method throughout the calculation.

In our calculation we use the fixed-point gauge

$$Z^\mu A_\mu(Z)=0 \ . \tag{2}$$

Then

$$A_\mu(Z)=\tfrac{1}{2}Z^\alpha G_{\alpha\mu}(0)+\tfrac{1}{3}Z^\alpha Z^\beta\nabla_\alpha G_{\beta\mu}(0)+\tfrac{1}{8}Z^\alpha Z^\beta Z^\gamma\nabla_\alpha\nabla_\beta G_{\gamma\mu}(0)+\tfrac{1}{30}Z^\alpha Z^\beta Z^\gamma Z^\delta\nabla_\alpha\nabla_\beta\nabla_\gamma G_{\delta\mu}+\cdots \ . \tag{3}$$

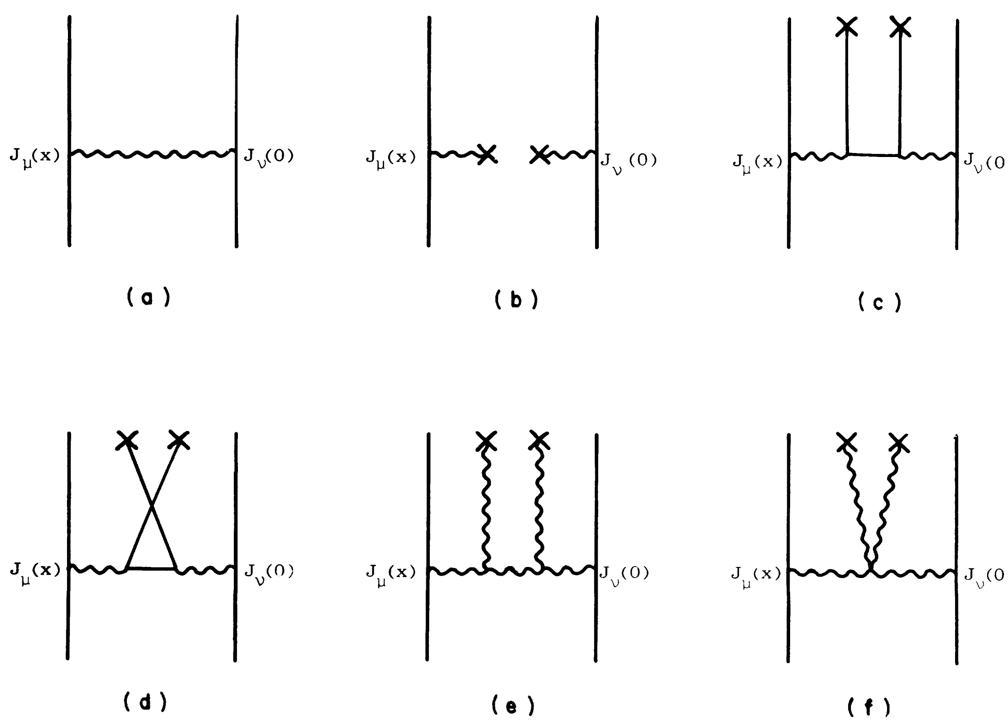

FIG. 1. One-gluon-exchange diagrams with one-pair $q\bar{q}$ or GG production.

In Eq. (3) we take the first term only; that is, we assume that the distance $|Z|$ is sufficiently small (the momentum transfer is large enough) so that we can cut the series at the second term safely. The obtained gluon propagator in momentum space which is directly related to the correlator reads as

$$\int d^4x\, e^{iqx}\langle 0|TA_\mu^a(x)A_\nu^b(0)|0\rangle = -i\frac{g_{\mu\nu}-q_\mu q_\nu/q^2}{q^2}F(\mathbf{q}^2)\delta_{ab}\ , \quad (4)$$

where the additional momentum-dependent factor $F(\mathbf{q}^2)$ is

$$F(\mathbf{q}^2)=1+\frac{1}{3}g^2\sum_{\beta=u,d,s}\frac{m_\beta\langle\psi_\beta\bar{\psi}_\beta\rangle}{q^2(q^2-m_\beta^2)}+\frac{9}{32}g^2\langle G^2\rangle\frac{1}{q^4}\ , \quad (5)$$

where $\langle\psi_\beta\bar{\psi}_\beta\rangle$ denotes the quark condensate with flavor β and $\langle G^2\rangle$ represents the gluon condensate. This full propagator of the gluon includes both perturbative and nonperturbative contributions to first order. It is qualitatively consistent with Larsson's result [14], but with different numerical coefficients as a result of the different methods used.

Next, we use this full propagator of the gluon instead of the normal perturbative one-gluon propagator to derive the Breit-Fermi interaction.

The M matrix can be written as

$$M_{fi}=\left[-\frac{g^2}{4}\right](\lambda_1^a\lambda_2^a)\left\{[\bar{u}_1(\mathbf{p}_1')\gamma^0u_1(\mathbf{p}_1)\frac{i}{\mathbf{q}^2}F(\mathbf{q}^2)\bar{u}_2(\mathbf{p}_2')\gamma^ju_2(\mathbf{p}_2)] + \left[\bar{u}_1(\mathbf{p}_1')\gamma^iu_1(\mathbf{p}_1)\frac{i}{q^2}\left[\delta_{ij}-\frac{q_iq_j}{\mathbf{q}^2}\right]F(\mathbf{q}^2)\bar{u}_2(\mathbf{p}_2')\gamma^ju_2(\mathbf{p}_2)\right]\right\}\ , \quad (6)$$

with

$$u_\lambda(\mathbf{p})=\left[\frac{E(\mathbf{p})+m}{2E(\mathbf{p})}\right]^{1/2}\begin{bmatrix}\xi_\lambda\\ \dfrac{\boldsymbol{\sigma}\cdot\mathbf{p}}{E(\mathbf{p})+m}\xi_\lambda\end{bmatrix}\ . \quad (7)$$

Considering all terms up to order $1/c^2$, M_{fi} can be rewritten as

$$\begin{aligned}M_{fi}=&-\frac{i}{4}g^2(\lambda_1^a\lambda_2^a)F(\mathbf{q}^2)\xi_{\lambda_1}^+\xi_{\lambda_2}^+\left[\frac{1}{\mathbf{q}^2}-\frac{1}{8m_1^2}-\frac{1}{8m_2^2}-\frac{\boldsymbol{\sigma}_1\cdot\boldsymbol{\sigma}_2}{4m_1m_2}-\frac{\mathbf{p}_1\cdot\mathbf{p}_2}{m_1m_2\mathbf{q}^2}+\frac{(\boldsymbol{\sigma}_1\cdot\mathbf{q})(\boldsymbol{\sigma}_2\cdot\mathbf{q})}{4m_1m_2\mathbf{q}^2}+\frac{(\mathbf{p}_1\cdot\mathbf{q})(\mathbf{p}_2\cdot\mathbf{q})}{m_1m_2\mathbf{q}^4}\right.\\&+i\left[\frac{1}{2m_1m_2\mathbf{q}^2}[-\boldsymbol{\sigma}_1\cdot(\mathbf{q}\times\mathbf{p}_2)+\boldsymbol{\sigma}_2\cdot(\mathbf{q}\times\mathbf{p}_1)]\right.\\&\left.\left.+\frac{1}{4m_1^2\mathbf{q}^2}\boldsymbol{\sigma}_1\cdot(\mathbf{q}\times\mathbf{p}_1)-\frac{1}{4m_2^2\mathbf{q}^2}\boldsymbol{\sigma}_2\cdot(\mathbf{q}\times\mathbf{p}_2)\right]\right]\xi_{\lambda_2}\xi_{\lambda_1}\\=&-i\xi_{\lambda_1}^+\xi_{\lambda_2}^+U(\mathbf{p}_1,\mathbf{p}_2,\mathbf{q})\xi_{\lambda_2}\xi_{\lambda_1}\ .\end{aligned}$$

Therefore the potential between two quarks in momentum space can be obtained in the following form:

$$\begin{aligned}U(\mathbf{p}_1,\mathbf{p}_2,\mathbf{q})=&\frac{g^2}{4}(\lambda_1^a\lambda_2^a)\left[1+\frac{g^2}{3}\sum_{\beta=u,d,s}\frac{m_\beta\langle\psi_\beta\bar{\psi}_\beta\rangle}{q^2(q^2-m_\beta^2)}+\frac{9}{32}g^2\langle G^2\rangle\frac{1}{q^4}\right]\\&\times\left\{\frac{1}{\mathbf{q}^2}-\frac{1}{8m_1^2}-\frac{1}{8m_2^2}-\frac{\boldsymbol{\sigma}_1\cdot\boldsymbol{\sigma}_2}{4m_1m_2}-\frac{\mathbf{p}_1\cdot\mathbf{p}_2}{m_1m_2\mathbf{q}^2}+\frac{(\boldsymbol{\sigma}_1\cdot\mathbf{q})(\boldsymbol{\sigma}_2\cdot\mathbf{q})}{4m_1m_2\mathbf{q}^2}+\frac{(\mathbf{p}_1\cdot\mathbf{q})(\mathbf{p}_2\cdot\mathbf{q})}{m_1m_2\mathbf{q}^4}\right.\\&\left.+i\left[\left[\frac{\boldsymbol{\sigma}_1}{4m_1^2}+\frac{\boldsymbol{\sigma}_2}{2m_1m_2}\right]\cdot\frac{\mathbf{q}\times\mathbf{p}_1}{\mathbf{q}^2}-\left[\frac{\boldsymbol{\sigma}_2}{4m_2^2}+\frac{\boldsymbol{\sigma}_1}{2m_1m_2}\right]\cdot\frac{\mathbf{q}\times\mathbf{p}_2}{\mathbf{q}^2}\right]\right\}\ . \quad (8)\end{aligned}$$

Furthermore, by making a Fourier transformation, the q-q potential in the coordinate space can be read as

$$U(\mathbf{p}_1,\mathbf{p}_2,\mathbf{r})=\int U(\mathbf{p}_1,\mathbf{p}_2,\mathbf{q})e^{i\mathbf{q}\cdot\mathbf{r}}\frac{d\mathbf{q}}{(2\pi)^3}$$

$$\begin{aligned}
&=\frac{g^2}{4\pi}\frac{\lambda_1^a\lambda_2^a}{4}\left\{\left[1-g^2\sum_\beta\frac{\langle\psi_\beta\bar\psi_\beta\rangle}{3m_\beta}\left[\frac{1}{8m_1^2}+\frac{1}{8m_2^2}+\frac{\sigma_1\cdot\sigma_2}{6m_1m_2}+\frac{1}{m_\beta^2}\right]\right]\frac{1}{r}-\pi\left[\frac{1}{2m_1^2}+\frac{1}{2m_2^2}+\frac{2\sigma_1\cdot\sigma_2}{3m_1m_2}\right]\delta(\mathbf{r})\right.\\
&\quad-g^2\left[\sum_\beta\frac{\langle\psi_\beta\bar\psi_\beta\rangle}{6m_\beta}-\frac{9}{64}\langle G^2\rangle\left[\frac{1}{8m_1^2}+\frac{1}{8m_2^2}+\frac{\sigma_1\cdot\sigma_2}{8m_1m_2}\right]\right]r+g^2\frac{3}{256}\langle G^2\rangle r^3\\
&\quad+g^2\sum_\beta\frac{\langle\psi_\beta\bar\psi_\beta\rangle}{3}\left[\frac{1}{8m_1^2}+\frac{1}{8m_2^2}+\frac{\sigma_1\cdot\sigma_2}{4m_1m_2}+\frac{1}{m_\beta^2}\right]\frac{e^{-m_\beta r}}{m_\beta r}-\frac{1}{2m_1m_2}\left[\mathbf{p}_1\cdot\mathbf{p}_2+\frac{(\mathbf{p}_1\cdot\mathbf{r})(\mathbf{p}_2\cdot\mathbf{r})}{r^2}\right]\frac{1}{r}\\
&\quad+g^2\sum_\beta\frac{\langle\psi_\beta\bar\psi_\beta\rangle}{8m_1m_2m_\beta}\left[\mathbf{p}_1\cdot\mathbf{p}_2-\frac{(\mathbf{p}_1\cdot\mathbf{r})(\mathbf{p}_2\cdot\mathbf{r})}{r^2}\right]r-g^2\frac{3\langle G^2\rangle}{1536m_1m_2}\left[5(\mathbf{p}_1\cdot\mathbf{p}_2)-\frac{3(\mathbf{p}_1\cdot\mathbf{r})(\mathbf{p}_2\cdot\mathbf{r})}{r^2}\right]r^3\\
&\quad-\frac{1}{4m_1m_2}S_{12}\frac{1}{r^3}-g^2\sum_\beta\frac{\langle\psi_\beta\bar\psi_\beta\rangle}{72m_1m_2m_\beta}S_{12}\frac{1}{r}-g^2\frac{9}{1024m_1m_2}\langle G^2\rangle S_{12}r\\
&\quad-\left[\left[\frac{\sigma_1}{4m_1^2}+\frac{\sigma_2}{2m_1m_2}\right]\cdot(\mathbf{r}\times\mathbf{p}_1)-\left[\frac{\sigma_2}{4m_2^2}+\frac{\sigma_1}{2m_1m_2}\right]\cdot(\mathbf{r}\times\mathbf{p}_2)\right]\frac{1}{r^3}\\
&\quad-g^2\sum_\beta\frac{\langle\psi_\beta\bar\psi_\beta\rangle}{6m_\beta}\left[\left[\frac{\sigma_1}{4m_1^2}+\frac{\sigma_2}{2m_1m_2}\right]\cdot(\mathbf{r}\times\mathbf{p}_1)-\left[\frac{\sigma_2}{4m_2^2}+\frac{\sigma_1}{2m_1m_2}\right]\cdot(\mathbf{r}\times\mathbf{p}_2)\right]\frac{1}{r}\\
&\quad\left.+g^2\frac{9}{256}\langle G^2\rangle\left[\left[\frac{\sigma_1}{4m_1^2}+\frac{\sigma_2}{2m_1m_2}\right]\cdot(\mathbf{r}\times\mathbf{p}_1)-\left[\frac{\sigma_2}{4m_2^2}+\frac{\sigma_1}{2m_1m_2}\right]\cdot(\mathbf{r}\times\mathbf{p}_2)\right]r\right\},
\end{aligned}\tag{9}$$

with

$$S_{12}=\frac{3(\sigma_1\cdot\mathbf{r})(\sigma_2\cdot\mathbf{r})}{r^2}-\sigma_1\cdot\sigma_2\ .\tag{10}$$

Before approaching the final expression $U(\mathbf{p}_1,\mathbf{p}_2,\mathbf{r})$, there are two points that should be mentioned.

First, our derivation is valid when the momentum transfer q is large enough or the interacting distance is small enough. Thus one can only take the first term in Eq. (3) without breaking the feasibility of this approximation and ensure the integrand $U(\mathbf{p}_1,\mathbf{p}_2,\mathbf{q})$ in Eq. (9) is reliable. In order to obtain a reasonable $U(\mathbf{p}_1,\mathbf{p}_2,\mathbf{r})$, Eq. (9) should be rewritten as

$$U(\mathbf{p}_1,\mathbf{p}_2,\mathbf{r})=\int_0^\delta\frac{d\mathbf{q}}{(2\pi)^3}e^{i\mathbf{q}\cdot\mathbf{r}}f(q)+\int_\delta^\infty\frac{d\mathbf{q}}{(2\pi)^3}e^{i\mathbf{q}\cdot\mathbf{r}}U(\mathbf{p}_1,\mathbf{p}_2,\mathbf{q})\ ,\tag{11}$$

where δ is a cutoff parameter and $f(\mathbf{q})$ is another function related to the long-distance (large r) behavior. One easy way to define $f(\mathbf{q})$ is by extrapolating $U(\mathbf{p}_1,\mathbf{p}_2,\mathbf{q})$ to this area and multiplying a weight factor such as $e^{-\Lambda^2/q^2}$, where the adjustable parameter Λ is also related to δ. That makes the analytic integration of Eq. (11) different. There could be some ways to solve it. For example, Huang *et al.* modified Richardson's potential by introducing a cutoff at a small value of q and giving an asymptotic behavior Kr as $r\to\infty\,(q<\varepsilon)$ [19]. Our way to handle this problem is by introducing a factor in coordinate space. We adopt this factor as that obtained from the calculation of lattice gauge theory [20], with which a better long-distance behavior of the potential can be shown. Moreover, the authors of Ref. [19] also replaced the linear term Kr by $Kr[(1-e^{-\mu r})/\mu r]$. It is easy to see that, as $r\to 0$, this term becomes linear in r, but as $r\to\infty$, it tends to constant K/μ. Similar to their treatment, we modify our long-distance term in the Fourier transformation by the same factor; then we have

$$V(r)=\frac{1-e^{-\mu r}}{\mu r}\int_0^\infty\frac{dq}{(2\pi)^3}U(\mathbf{p}_1,\mathbf{p}_2,\mathbf{q})\ ,\tag{12}$$

where μ is a parameter which replaces δ or Λ mentioned above and can be determined by fitting the data. Thus the integrand $U(\mathbf{p}_1,\mathbf{p}_2,\mathbf{q})$ is extrapolated to the small-q range, and the additional factor takes care of the caused modification. In fact, we merely hope our result is valid at short and medium distances and the resultant scalar potential might play the role of the confinement. Since, in our ansatz, the effect at long distance, where the effect of the confinement is also ambiguous, is taken into account by the well-known effective pion field, it is not necessary to worry about the confinement effect at a range longer than few femtometers.

The second problem is caused by the approximation, where all loop contributions are ignored because of loop suppression. In fact, a new term $(1/Q^2)\ln Q^2$ in momentum space can be given by the gluon vacuum polarization. As a consequence, the Coulomb potential can be slightly changed. We omit it at the present stage.

Now we reach the final form of the effective potential between quarks:

$$V(\mathbf{r})=U(\mathbf{p}_1,\mathbf{p}_2,\mathbf{r})\left[\frac{1-e^{-\mu r}}{\mu r}\right] . \tag{13}$$

III. DISCUSSION

In Eq. (9) the summation is taken over all light quarks u, d, and s, and the masses of constituent quarks are taken as $m_u=m_d=350$ MeV and $m_s=550$ MeV. The phenomenological vacuum expectation values of $(\psi\bar{\psi})$ and $(G_{\mu\nu}^2)$ are given as

$$\langle 0|u\bar{u}|0\rangle=\langle 0|d\bar{d}|0\rangle =1.3\langle 0|s\bar{s}|0\rangle=(0.25\ \mathrm{GeV})^3$$

and $\langle 0|(\alpha_s/\pi)G_{\mu\nu}^2|0\rangle=0.012\ \mathrm{GeV}^4$. The effective coupling constant $\alpha_s=g^2/4\pi$ in Eq. (9) is adopted as $\alpha_s=0.5$.

There are lots of interesting features in the resultant potential $V(\mathbf{r})$. First of all, there is a linear term which can play a role of the regular confinement potential; meanwhile, a r^3 term with a sign opposite that of the coefficient to the linear term would give the deconfinement effect. Moreover, a Yukawa-type term automatically appears. It may somewhat provide an interaction at longer distance, such as part of the effects brought about by the pseudoscalar meson exchange. This is easy to understand. In our framework, since the underlying mechanism is the current-current correlator, one may obtain an effective potential, which governs the scattering process, in the medium and short ranges.

Similarly, the spin-orbital and tensor interactions are also modified by the additional terms generated by the quark and gluon condensates. With these modifications the theoretical baryon spectrum and scattering phase shifts would be varied.

Here we continue our discussion on confinement-related terms. The r^3 term comes from the gluon propagator, where a gluon pair is produced. This indicates that the condensate can modify the potential via the full gluon propagator. Because of the appearance of the r^3 term, which provides the deconfinement effect, the behavior of the confinement potential is somewhat different from the commonly used linear confinement potential. This deviation will definitely cause some changes in the spectrum and scattering phase shifts. Does the r^3 term agree with our hypothesis, or does it mean that the quark can be defined through tunneling the potential barrier? The answer is negative. It just looks like the picture in the old string theory. Two quarks (or quark-antiquark) antiquark) are linked by a string. When they are pulled away from each other,the tension on the string becomes greater and greater until the string is broken. Then, at each broken end, a quark (or antiquark) is generated to keep the color-singlet feature of the system. Thus, at large r, the r^3 term becomes very important, and at the very large r, our hypothesis may fail. However, an equivalent to the ladder approximation in field theory, if we consider higher-order nonperturbative corrections, i.e., take into account more quark and gluon condensates, the $r^5, r^7, \ldots$ terms may appear with different signs, and the shape of the confinement potential may converge to that given phenomenologically by the lattice gauge calculation [20]. It is noteworthy that by using the confinement potential in Ref. [20] or a confinement potential in the error function form, Yang, Deng, and Zhang [6] calculated the hadron spectrum and greatly improved the energy level for $N=2$. This indicates that the deconfinement term does exist, and the value of the total confinement potential, at least, should not be increased with increased r. Thus the above-mentioned tunneling problem might be solved.

It is natural to have higher-order terms of r or to consider more condensates. Because all $(q\bar{q})$ and (GG) condensates mix and some overlap, where q (or $\bar{q}$) or G lines cross each other adding contributions, the calculation becomes more complicated. For example, in the ladder approximation, both horizontal and vertical ladders in the same order must be taken into account at the same time. In our opinion, first, one should find out how the lowest order of nonperturbative QCD offers its contribution to the potential. If the result is encouraging, one should work further on higher-order corrections.

Finally, we mention that all parameters in $V(\mathbf{r})$ should be fixed in the spectrum and scattering calculations. These results will appear in a future paper where the hadron spectrum and NN scattering phase-shift calculations are performed.

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