PHYSICAL REVIEW C **76**, 025201 (2007)

Soliton in the global color model with a sophisticated effective gluon propagator

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(Received 8 January 2007; revised manuscript received 7 June 2007; published 2 August 2007)

With a sophisticated effective gluon propagator, Maris-Tandy model, we solve the Dyson-Schwinger equation to get the quark propagator and then study the soliton solution in the global color model (GCM). Along the constraints on the parameters fitted to the pion decay constant, we take several sets of parameters and find that some of the properties of soliton can be produced in the GCM soliton model with a special choice of parameters. We also discuss the influences of the parameters and the ultraviolet perturbative term on the property of the soliton. We find that the interaction among quarks is the one with self-adjusting characteristic and only the fine-tuned interaction can generate an appropriate solition, but not that much stronger attraction produces more stable soliton.

DOI: 10.1103/PhysRevC.76.025201 PACS number(s): 12.39.−x, 03.75.Lm, 12.40.Yx, 11.10.Lm

I. INTRODUCTION

There have been lots of phenomenological models to try to give the structure of baryons. It has been shown that soliton models, including the two most successful kinds, the Skyrme model (or the chiral soliton model, see for example Ref. [1]) and the chiral quark soliton model (or Nambu-Jona-Lasinio (NJL) soliton model, see for example Refs. [2–4]), can give a vivid picture of baryons. However, the Skyrme model has no valence quark degree of freedom, while the chiral quark soliton model based on the NJL model keeps the information of the valence quark degree of freedom but only takes into account the local interaction between quarks. On the other hand, with the phenomenological success, people search for the quantum chromodynamics (QCD) foundation of these models. With a point interaction in the coordinate space to the model of the gluon propagator, one can derive the local NJL model. Nevertheless, the realistic low energy effective interaction in QCD is not the point one. Diakonov *et al*. [4] have first derived the nonlocal NJL interaction model from the dilute instanton model [5] and the global color model (GCM) [6] was also developed to implement the nonlocality of the low energy physics.

In the GCM, the gauge symmetry was discarded and then QCD reduced to a finite-range current-current interaction theory. One can get a quark-meson interaction model or a quark-diquark interaction model after bosonization [6] of the current-current interaction. Then baryons can be modeled as solitons with quark moving in the background of chiral meson fields [7,8]. With the GCM soliton, one can also discuss the dynamical confinement and, in particular, the effect of the nonlocality. Furthermore, the dependence of baryon properties on hadron matter medium has also been studied in the GCM [9–14].

To carry out the calculation of the GCM soliton, one needs a dressed quark propagator. It has been shown that such a dressed propagator can be provided by the Dyson-Schwinger (DS) equations [6]. The research of DS equations has been developed with an extremely complicated technique (for reviews, see for example, Refs. [15–19], and references therein). In most recent years, great progress has been made in systematic studies on quark-quark interaction (see for example Refs. [19–28]). One should, in principle, take the full gluon and quark propagators with clear analytical structure (c.f. given in Refs. [19–22]) to study baryon structure. However, the interaction vertex has not yet been determined well. For simplicity and practical calculation, it is still powerful to solve the quark DS equation in the rainbow approximation with an input of phenomenological (effective) gluon propagator at present stage. And it has been shown that such a handling can analytically continue the quark propagator and study its property in the time-like region [29], which is necessary in studying hadron structure in the frameworks of the Bethe-Salpeter equation, Faddeev equations, and soliton equations. There have now been two kinds of effective (phenomenological) gluon propagators, the infrared divergent one (e.g., [21,30]) and the infrared regular one (e.g., [26,29,31–34]). It has been known that the Maris-Tandy model involves both the infrared enhancement and the ultraviolet perturbative term and can describe the pion properties quite well [32]. It is then regarded as a sophisticated and successful effective gluon propagator. In this paper we will study the properties of the soliton in the GCM with the effective gluon propagator in the Maris-Tandy model and try to search for some relation between the effective gluon propagator and some properties of the soliton in the framework of the GCM soliton model.

The paper is organized as follows. In Sec. II, we describe briefly the GCM formalism and the soliton in GCM. In Sec. III, we give the numerical results and discussions. Finally we summarize this work and give some comments in Sec. IV.

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II. BRIEF DESCRIPTION OF THE SOLITON MODEL IN THE GCM

The global color model (GCM) is constructed by the the generation functional in Euclidean space as [6]

$$Z = \int \mathcal{D}\bar{q}(x)\mathcal{D}q(y)\exp\left[-\int d^4x\bar{q}(x)\partial\!\!\!/ q(x) - \frac{g^2}{2}\int\int d^4xd^4y\bar{q}(x)\gamma_\mu\frac{\lambda^a}{2}q(x) \times D(x-y)\bar{q}(y)\gamma_\mu\frac{\lambda^a}{2}q(y)\right]. \quad (1)$$

This functional is invariant under global color $SU(3)$ transformation rather than the gauge color $SU(3)$ transformation. $D(x-y)$ is the effective gluon propagator, which is a parametrized function to model the low energy property and dynamics such as those of the hadrons. It has been shown that the infrared enhancement of $D(x-y)$ is closely related to the chiral symmetry breaking (see for example Refs. [15,35,36]).

After bosonization [6], one can get the action of the bilocal fields $B^\theta(x,y)$

$$S[B^\theta(x,y)] = -Tr\ln\left[\gamma\cdot\partial\delta(x-y) + \frac{M^\theta}{2}B^\theta(x,y)\right] + \int d^4xd^4y\frac{B^\theta(x,y)B^\theta(y,x)}{2g^2D(x-y)}. \quad (2)$$

The bilocal fields $B^\theta(x,y)$ have the same quantum numbers with the mesons and hence people identify the fluctuation above the vacuum configuration as mesons. One can determine the vacuum configuration through the saddle point condition $\frac{\delta S}{\delta B_0^\theta}=0$, which induces a truncated Dyson-Schwinger equation with rainbow approximation, which reads

$$\Sigma(p) = g^2\int\frac{d^4q}{(2\pi)^4}D(p-q)\frac{\lambda^a}{2}\gamma_\mu\frac{1}{i\gamma\cdot q+\Sigma(q)}\gamma_\mu\frac{\lambda^a}{2}, \quad (3)$$

where $D(k)$ is the effective gluon propagator in momentum space [i.e., the Fourier transformation of $D(x)$]. The self-energy can usually be decomposed as

$$\Sigma(p) = i\gamma\cdot p[A(p^2)-1] + B(p^2) = G^{-1}(p) - i\gamma\cdot p, \quad (4)$$

where $G^{-1}(p)$ is the inverse of the dressed quark propagator. Equation (3) is then in fact two coupled integral equations of $A(p^2)$ and $B(p^2)$.

The dynamical chiral symmetry breaking is due to the Nambu solution $B(p^2)\neq 0$, which causes quark a dynamical mass $B(p^2)/A(p^2)$. One can model the low energy property through a certain form of $A(p^2)$ and $B(p^2)$ determined by solving Eq. (3) with an effective gluon propagator put forward phenomenologically or that derived from lattice QCD or some other approaches such as the instanton model [5].

To extract the information of a specific baryon from the Lagrangian, one usually starts from the baryon correlation function [37], which in the case of nucleon is

$$\Pi_N(T) = \frac{1}{Z}\int\mathcal{D}\bar{q}\mathcal{D}q\mathcal{D}B^\theta J_N\left(\vec{y},\frac{T}{2}\right) \times J_N^+\left(\vec{x},-\frac{T}{2}\right)e^{-\int d^4x\mathcal{L}}, \quad (5)$$

where

$$J_N(\vec{x},t) = \frac{1}{N_c!}\varepsilon^{\beta_1\cdots\beta_{N_c}}\Gamma^{\{f\}}_{JJ_3,TT_3}q_{\beta_1,f_1}(\vec{x},t)\cdots q_{\beta_{N_c},f_{N_c}}(\vec{x},t) \quad (6)$$

is the nucleon current. $\Gamma^{\{f\}}_{JJ_3,TT_3}$ is a matrix, β_i is color index, and f_i denotes the flavor and spin structure. In the Euclidian space, the asymptotic behavior of the nucleon correlation function is

$$\lim_{T\to+\infty}\Pi_N(T)\sim e^{-M_NT}, \quad (7)$$

and then one can get the nucleon mass M_N.

As people known, it is very hard to calculate the above correlation function with the full time and space dependent quark fields and meson fields, which can only be calculated with the Monte Carlo technique in lattice QCD. The philosophy of the soliton model is to calculate the correlation function with static meson fields at first and then take into account the modifications of the quantum fluctuation on the static ones.

With a static meson field, we can get the total mass of the soliton

$$M_{cl} = N_c\varepsilon_{\rm val} + E_m, \quad (8)$$

where $\varepsilon_{\rm val}$ is the lowest pole of the quark propagator and E_m is the contribution from the meson fields. One can find the detail of the formulas in [7,8]. The subscript cl means the mass is the classical mass without considering the quantum fluctuation of the fields. The mass is just a mass of the soliton and one cannot identify this with the nucleon in the experiment because there are no quantum fluctuations. One should quantize the fluctuation above this classical soliton to get the quantum state corresponding to the nucleon. However, the quantization is very complicated and hard to deal with. We are working on this direction and will show the recent results of the quantization in the other places. We take the saddle point condition $\delta M_{cl}/\delta B^\theta = 0$ to determine the meson fields. The equation of the saddle point condition became a group of coupled Dirac-like equation and Klein-Gordon–like equations, which can be found in [7,8].

III. NUMERICAL RESULTS AND DISCUSSION

To solve the soliton equations, one needs at first the effective gluon propagator to get the scalar functions $A(p^2)$ and $B(p^2)$ in the dressed quark propagator. In the present work, we take the Maris-Tandy model [32] for the effective gluon propagator, which reads

$$g^2D(q) = \frac{4\pi^2 d}{\omega^6} q^2 e^{-q^2/\omega^2} + \frac{8\pi^2\gamma_m \pi}{\ln\left[\tau + \left(1 + \frac{q^2}{\Lambda_{\mathrm{QCD}}^2}\right)^2\right]} \times \frac{1 - \exp\left(-\frac{q^2}{4m_t^2}\right)}{q^2}, \tag{9}$$

where the parameters γ_m, τ, Λ_{QCD}, and m_t are usually taken as constants [32], which read $\gamma_m = 12/25$, $\tau = e^2 - 1$, $\Lambda_{\mathrm{QCD}} = 0.234$ GeV, $m_t = 0.5$ GeV. The other parameters ω and d can be fixed by the best fitting of the properties of pion [32]. It is obvious that the second term in Eq. (9) serves the one-loop asymptotic behavior of the gluon propagator. The first term displays the infrared enhancement and vanishes at zero exchange momentum. It is reported that such an effective gluon propagator coincides with the behavior of gluon propagator given in lattice QCD simulations (see for instance Ref. [28]) qualitatively well. Then this effective gluon propagator is believed to be quite sophisticated and close to the realistic one.

For the Dyson-Schwinger equation (3), we first solve it in the Euclidean space. In order to continue our soliton calculation, we do the analytical continuation of the propagator from the space-like real axis to the time-like region, by taking the same method as that in Ref. [29] and smoothly changing the external momentum to virtual space while keeping the loop momentum space-like.

For the Dirac-like quark equation, we solve it in momentum space. For the soliton of a nucleon, we consider only the orbital angular momentum $L = 0$ states,

$$u_j(\vec{p}) = \begin{pmatrix} f_j(p) \\ i\vec{\sigma}\cdot\hat{p}g_j(p) \end{pmatrix}, \tag{10}$$

and take the Hedgehog [38] form solution of the meson equations

$$\sigma(\vec{r}) = \sigma(r), \tag{11}$$

$$\pi_i(\vec{r}) = \hat{r}_i\pi(r)(i = 1, 2, 3), \tag{12}$$

with $\hat{r}_i$ being the polar direction in the coordinate space.

In the process of solving the equations, we first take trial meson fields to solve the quark equation and then input the quark wave function to solve the meson equations to get new meson fields. Continuing the iteration to a desired precision, we obtain the final solutions.

After solving the coupled equations, we obtain the eigenenergy $\varepsilon_{\mathrm{val}}$ and the wave function $u_j(\vec{p})$ of the quarks and the meson fields, and, in turn, their potential energy $E_p = \int d^3x U(\chi^2)$ and kinetic energy $E_k = \int d^3x[\frac{1}{2}(\nabla\sigma)^2 + \frac{1}{2}(\nabla\vec{\pi})^2]$. The energy of the soliton can then be determined by $E_{\mathrm{tot}} = 3\varepsilon_{\mathrm{val}} + E_p + E_k$. Furthermore, We obtain the mass of the soliton with two methods. One is the naive center mass reduction, $M_{\mathrm{ncr}} = \sqrt{E^2 - \sum_j \langle p_j^2\rangle}$, where $\langle p_j^2\rangle$ is the expectation value of the square of the quark momentum, the sum is over all the valence quark levels. We also take the scheme of recoil correction [39] to deal with the center of mass motion. The formulae to calculate M_{rec} and R_{rec} can be found as Eqs. (29), (34), and (35a) in the paper [39], which we do not give here for concise description.

TABLE I. Calculated properties of the GCM soliton with the Maris-Tandy model of the effective gluon propagator.

	2T	2F	3T	3F	4T	4F
ω (GeV)	0.401		0.450		0.472	
d (GeV^2)	0.930		0.830		0.790	
$\varepsilon_{\mathrm{val}}$ (MeV)	107	162	189	237	270	300
E_p(MeV)	109	123	126	106	93	15
E_k(MeV)	766	608	527	355	224	18
E_{tot} (MeV)	1196	1217	1220	1172	1127	933
M_{ncr}(MeV)	1060	1094	1084	1044	996	890
M_{rec}(MeV)	956	909	892	814	761	618
R_{ncr}(fm)	0.67	0.71	0.79	0.93	1.15	2.74
R_{rec}(fm)	0.60	0.61	0.68	0.77	0.95	2.22

In practical calculations, we take four sets of parameters, 1 to 4, along the fitted condition $\omega d \approx (0.72\,\mathrm{GeV})^3$ [18], for the effective gluon propagator in Eq. (9) which is marked with T as the last letter. It is then in fact a one parameter model of effective gluon propagator. We also take the effective gluon propagator without the ultraviolet term in Eq. (9) and name the parameter set with F as the last letter. The calculations show that the parameter set 1 ($\omega = 0.30$ GeV, $d = 1.25\,\mathrm{GeV}^2$) cannot give a static soliton solution. We list the results of some properties of the soliton obtained with parameter sets 2, 3, and 4 in Table I and illustrate the obtained quark field and chiral meson fields in the soliton in Fig. 1.

One can recognize easily from Table I that the ultraviolet term of the effective gluon propagator increases the mass (except for that with parameter set 2 and naive center of mass reduction) and shortens the radius of the soliton in the GCM soliton model. In more detail, the ultraviolet term decreases the single particle energy of the quark and enhances the kinetic energy of the chiral fields in the soliton. However, its contribution to the potential energy of the chiral fields is not monotonic with respect to the change of the chiral fields (which

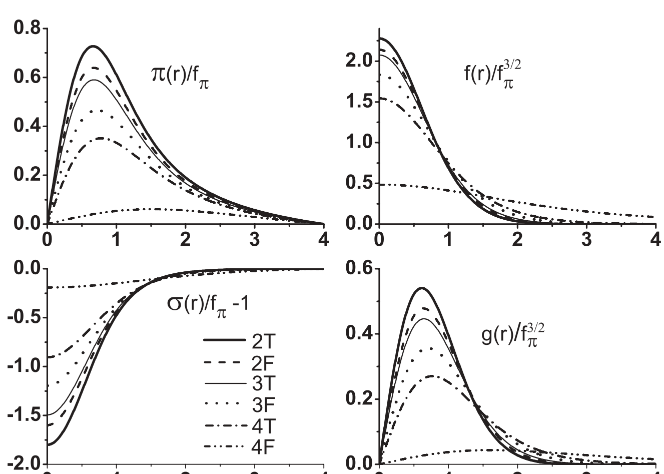

FIG. 1. Calculated distributions of the quark field and the chiral meson fields in the soliton in the Maris-Tandy model of effective gluon propagator with parameter sets 2, 3, and 4.

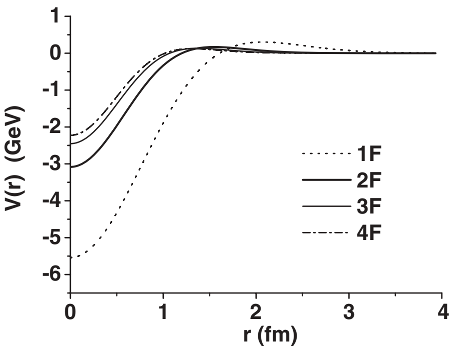

FIG. 2. Classical potential corresponding to the first term of effective gluon propagator in Maris-Tandy model with parameter sets 1, 2, 3, and 4.

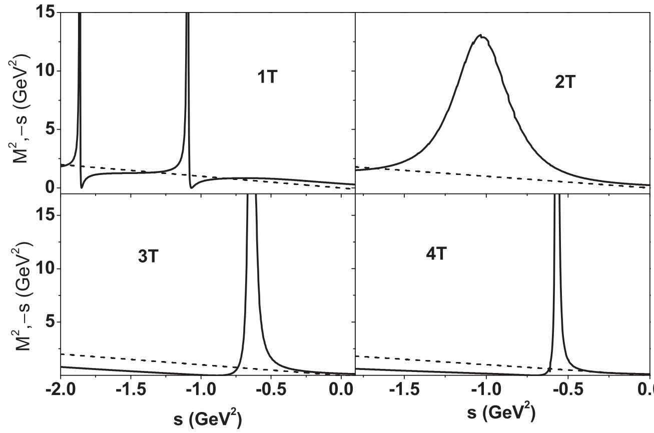

FIG. 3. Mass functions with the four sets of parameters. The solid curve shows the square of mass function and the dashed line represents the function $-s$.

can be seen directly from the formula of E_m in Refs. [7,8]). And for parameter set 4, the ultraviolet term plays an essential role in forming the soliton (the radius of the soliton without the ultraviolet term is unreasonably large). Meanwhile increasing the parameter d (decreasing the parameter ω simultaneously) influences the properties of the soliton in the same manner as that including the ultraviolet term. Such effects of including the ultraviolet term and increasing the parameter d can also be seen from the variation of the distributions of the quark and the chiral fields in the soliton illustrated in Fig. 1. Besides, parameter set 2 reproduces the properties of a soliton better.

To explore the effects of the effective gluon propagator on the properties of the soliton in the GCM soliton model, we illustrate the classical potential corresponding to the first term with the four sets of parameters of Eq. (9) in Fig. 2. From Eq. (9) and Fig. 2, one can recognize easily that the parameter d determines the depth of the classical potential and the parameter ω plays a role of the screening parameter. It is apparent that increasing the parameter d and decreasing the parameter ω increases the depth of the classical potential and enlarges the interaction radius. More explicitly, the above-mentioned variation of the parameter strengthens the interaction between quarks. As a consequence, the eigenenergy and the distribution range of the quark in the soliton decrease and the kinetic energy of the chiral fields increases. If one goes further along such a way naively, one may expect that parameter set 1T would produce a more tightly bound soliton. However, we have not obtained a soliton solution with parameter set 1.

To understand why we have not obtained a soliton solution with parameter set 1 in the Maris-Tandy model of the effective gluon propagator, we analyze the behavior of the quark propagator in the time-like space, since the calculation of the soliton needs the information of the quark propagator in the time-like region. With the measure we described above we obtain the quark propagator and, in turn, the mass function of the quark. The obtained mass functions of the quark with the four sets of parameters are illustrated in Fig. 3. The intersections of the dashed line and the solid curve in the figure fulfill the condition $p^2A^2(p^2) + B^2(p^2) = 0$, which are just the poles of the quark propagator in the time-like space. Since the functions A and B in the quark equation depend on $p^2 = |\vec{p}|^2 - \varepsilon_{\rm val}^2$. All the information one needs in the quark equation is the functions $A(s)$ and $B(s)$ with s larger than $-\varepsilon_{\rm val}^2$. One can see evidently from Fig. 3 that the quark propagator corresponding to parameter set 1T involves many poles, which locate quite near from each other and just in the time-like region where we would consider in the quark equation. These poles may then contribute to the solution of the quark equation. As a consequence, the mass functional in Eq. (8) may be so complicated that there is not a far isolated unique minimum. It makes the iteration unstable and not convergent in calculation. Because of the appearance of various nearly located minima in the mass functional and the quantum fluctuation, one may then not be able to mimic the baryons as classical solitons. However, surmising as much needs further investigation.

On the other hand, in the naive point of view, the emergence of time-like poles means that there are free quarks. However, people have not yet had clear cognition about the relations between the poles and the quark confinement [40–42] up to now. Moreover, Alkofer *et al.* [33] have shown that these poles do not have great influence on the light mesons in the framework of the Bethe-Salpeter equation. In the present case, if the poles which we referred to above locate in the region smaller than and not very close to $s = -\varepsilon_{\rm val}^2$, they will not contribute to the calculation of the quark equation. Along such a line, one can infer that the poles that lie far from the low energy region have little effect on the property of the soliton we discuss here. Therefore the effective gluon propagators with parameter sets 2 and 3 produce quite good soliton solutions.

To solidify the above argument, we should also make sure that the quarks are definitely confined in the soliton. To such an end, we refer to the axiom of reflection positivity of quantum field theory [41]. It has been shown that one can calculate the Schwinger function [19,22,41] from the quark propagator with

$$\Delta(t) = \int d^3x \int \frac{d^4p}{(2\pi)^4} e^{i(p_4 t+\vec{p}\cdot\vec{x})}\sigma(p^2), \qquad (13)$$

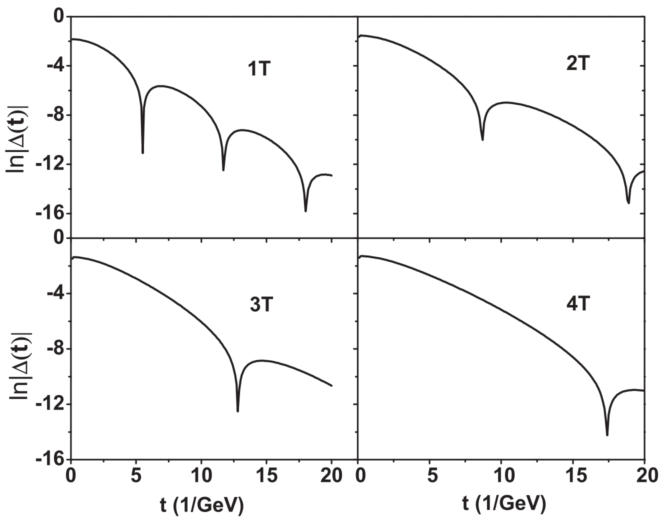

FIG. 4. Calculated Schwinger functions corresponding to the quark propagators determined with the four sets of parameters.

where the function $\sigma(p^2)$ is a scalar function extracted from the quark propagator, such as $\frac{A(p^2)}{p^2A^2(p)+B^2(p)}$ and $\frac{B(p^2)}{p^2A^2(p)+B^2(p)}$. The axiom of reflection positivity constrains the Schwinger functions as positive definite functions. If one obtains a propagator whose Schwinger function is not positive definite, the corresponding degree of freedom violates the positivity of quantum mechanics and then the propagator has no Lehmann spectral representation. Therefore, the asymptotic state cannot be the physical one and cannot be observed in experiment. On occasion, the quark propagator can show the violation of reflection positivity and then there is the quark confinement. One can see from Fig. 3 that there are time-like poles in the quark propagator and may infer that there are free quarks. However, the corresponding Schwinger functions of the presently obtained quark propagator are not positive definite, as shown in Fig. 4, where the peaks in the curves correspond to the intersection of the positive region and the negative region of the Schwinger function. Thus, all the quark propagators in this paper do have the property of quark confinement.

Recalling the above discussions, one can recognize that, even though the classical potential of the effective gluon propagator corresponding to parameter set 1T involves a very strong attraction, it cannot generate a solition because the corresponding dressed quark propagator holds many closely located poles in the time-like region. On the other hand, although all the dressed quark propagators corresponding to parameter sets 2, 3, and 4 possess less poles in the time-like space, since the poles appears at low momentum for the ones with parameter sets 3 and 4, these two sets of parameters cannot produce the soliton with an appropriate property for the nucleon. These results show evidently that the interaction among quarks in a soliton is in fact self-adjusting and then a soliton loading the properties of a nucleon can be reproduced well with the fine-tuned parameter (for example, set 2).

IV. SUMMARY AND REMARKS

In summary, we have calculated some properties of the soliton in the global color model of QCD with a sophisticated effective gluon propagator (in the Maris-Tandy model) and discussed the effect of the parameter in the infrared term as well as that of the ultraviolet term of the effective gluon propagator. We found that increasing the strength parameter d (decreasing the screening parameter ω) and including the ultraviolet term can decrease the eigenenergy of the quark as well as shorten and strengthen the distributions of the quark field and of both the pseudoscalar and the scalar chiral fields. As a consequence, the kinetic energy of the chiral fields and the interaction strength between quarks can be enhanced and the distribution range of the fields can be shortened and, in turn, increase the mass and shrink the soliton. As a result, with an appropriate choice of the parameter we produce some of the properties of the soliton in the GCM soliton model satisfactorily. Besides, we show that the interaction among quarks in a soliton is in fact self-adjusting and only the especially fine-tuned interaction (or parameters in the effective gluon propagator) can generate a soliton loading the properties of a nucleon appropriately. However we cannot identify the soliton state with a nucleon in the laboratory definitely until we quantize the angular momentum and the isospin of the Hedgehog soliton states and include the quantum excitation of the meson and quark fields. It is fortunate that the under estimate of the mass of the soliton just leaves room for including the quantum effects. The related investigations are under progress.

ACKNOWLEDGMENTS

This work was supported by the National Natural Science Foundation of China under contract Nos. 10425521, 10575004, 10575005, and 10675007, the Major State Basic Research Development Program under contract No. 2007CB815000, the Key Grant Project of Chinese Ministry of Education (CMOE) under contact No. 305001, and the Research Fund for the Doctoral Program of Higher Education of China under grant No. 20040001010. One of the authors (Y.X.L.) would also like to acknowledge the support of the Foundation for University Key Teacher by the CMOE.

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