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Deconfinement phase transition and disappearance of the soliton solution in a nontopological soliton bag model

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Based on the nontopological soliton model, a new picture for the deconfinement phase transition is proposed. The relation between the existence of the soliton solution and the nonlinearity of the potential function is analyzed. The effective potential and the equation satisfied by its extrema are given at finite temperature. It turns out that, in this model, at a critical temperature T_c, the physical vacuum is transformed into a perturbative vacuum, the soliton solution disappears, and the deconfinement transition occurs.

I. INTRODUCTION

Since the exact QCD theory in the large-scale nonperturbative region cannot be evaluated analytically, some phenomenological models, possessing the essential properties of QCD theory, have been developed, one of which is the nontopological soliton model[1] initiated by Friedberg and Lee. This model has been successfully applied to the description of the static properties of isolated hadrons by Goldflam and Wilets[2] and many other authors.[3]

The nontopological soliton model has two main features. One is the phenomenological scalar σ field which characterizes the nonperturbative property of the physical vacuum. Its interaction potential has three extrema (cf. Fig. 1 in Sec. II)—two minima, a relative one and an absolute one, corresponding to the perturbative and physical vacuum states, respectively, and a local maximum, locating between the two minima. The difference of height of the two minima measures the pressure of the physical vacuum on the perturbative vacuum, i.e., the bag constant B. The necessary condition for the existence of the soliton solution (confinement state of quarks) is that the nonlinearity of the potential function guarantees the existence of three extrema, located at the above-mentioned relative positions. Another feature of the nontopological soliton model is the color-dielectric constant κ, which is a function of σ. The form of this function gives $\kappa=1$ at the perturbative vacuum state inside the hadron and $\kappa=0$ at the physical vacuum state outside. The perfect color-dielectric property $\kappa=0$ of the physical vacuum leads to the confinement of quark and gluon fields.

Whether the confinement of quarks and gluons can be removed at finite temperature is an important question. In order to answer this question, some authors[4–6] emphasized the importance of the variation of bag constant B with temperature. They proposed the following picture for phase transition: As $T\rightarrow T_c$, $B\rightarrow 0$ and the bag radius expands without limit with the increasing of temperature, leading to deconfinement of the quark and gluon. For example, the result obtained by Takagi[5] is that the relation between bag radius and temperature satisfies $R(T)\propto(T_c^4-T^4)^{-1/4}$, where T_c is the transition temperature. When $T\rightarrow T_c$, $R(T_c)\rightarrow\infty$ and deconfinement transition occurs. Starting from a different point of view, Pisarski[6] obtained analogous results as well.

It should, however, be noticed that, in the nontopological soliton model, it is insufficient only to consider the variation of B with temperature. The existence of the confinement state (soliton solution) depends not only on the bag constant B, but also on the nonlinearity of the potential, represented by the three extrema mentioned above. In this paper, we will carefully investigate the change of the shape of the potential function with temperature. Starting from the Lagrangian of the model, we will deduce the effective potential at finite temperature and the equation satisfied by its extrema, making use of the technique proposed by Linde,[7] and will try to answer the following questions: whether the condition of the existence of the soliton solution is destroyed at finite temperature; whether the structure of the vacuum is changed in a way leading to deconfinement transition. The answer to these questions turns out to give another picture of the phase transition, different from the above-mentioned one. In this picture, as the temperature increases, the bag radius has only a minor increase. The main effect of finite temperature is that the intermediate region between the physical vacuum state outside the soliton bag and the perturbative vacuum state inside becomes wider and wider and the transition between these two states becomes smoother and smoother. When $T=T_c$, the vacuum structure with both physical and perturbative vacua becomes unstable, the physical vacuum is transformed into a perturbative one, the soliton bag disappears, and the deconfinement transition takes place.

This paper is organized as follows. In Sec. II we discuss the relation between the existence of the soliton solution and the number of extrema of the potential function at zero temperature. The effective potential at finite temperature and the equation satisfied by its extrema are given in Sec. III. In Sec. IV the variation of the effective potential with temperature is discussed in some detail. It will be shown that, as the temperature increases, the number of extrema of the effective potential reduces, the soliton solution disappears, and confinement is removed. In Sec. V the color-dielectric function at finite temperature is given and the change of the vacuum structure at critical temperature is analyzed. The main results of this paper are summarized in Sec. VI.

II. RELATION BETWEEN THE EXISTENCE OF THE SOLITON SOLUTION AND THE NUMBER OF EXTREMA OF POTENTIAL FUNCTION

The Lagrangian of the nontopological soliton model has the form[1]

$$\mathcal{L}=\overline{\Psi}(i\gamma_\mu\partial^\mu-g\sigma)\Psi+\tfrac{1}{2}\partial_\mu\sigma\partial^\mu\sigma-U(\sigma)\ , \tag{2.1}$$

$$U(\sigma)=\frac{a}{2!}\sigma^2+\frac{b}{3!}\sigma^3+\frac{c}{4!}\sigma^4+B\ , \tag{2.2}$$

where Ψ is the quark field, σ is the phenomenological scalar field, B is the bag constant, and g,a,b,c are four adjustable parameters, the values of which are restricted by the requirement of fitting the static properties of hadrons. This fitting procedure is not unique. There are different possibilities for the choice of g,a,b,c. They can be grouped into families, characterized by a parameter $f=b^2/ac$ (Ref. 2), called the "family parameter." Essentially, f characterizes the shape of $U(\sigma)$. $f=3$ and $f=\infty$ correspond to $B=0^+$ and $a=0$, respectively, as shown in Fig. 1.

From Eq. (2.1) the coupled equations of motion can be obtained:

$$\partial^2\sigma+U'(\sigma)+g\overline{\Psi}\Psi=0\ , \tag{2.3a}$$

$$(i\gamma_\mu\partial^\mu-g\sigma)\Psi=0\ , \tag{2.3b}$$

where a prime denotes a derivative with respect to σ. In the mean-field approximation, the σ field has a baglike soliton solution.[2] The existence of this kind of solution is closely related to the nonlinearity of $U(\sigma)$, represented by the number of extrema. In general, the $U(\sigma)$, leading to the soliton solution, has three extrema: one local minimum corresponding to a perturbative vacuum state located at $\sigma_{\rm vac}^{\rm per}=0$; one absolute minimum corresponding to a physical vacuum at

$$\sigma_{\rm vac}^{\rm phy}=\frac{3|b|}{2c}\left[1+\left[1-\frac{8ac}{b^2}\right]^{1/2}\right] \tag{2.4}$$

and a local maximum U_0 at

$$\sigma_m=\frac{3|b|}{2c}\left[1-\left[1-\frac{8ac}{b^2}\right]^{1/2}\right]\ , \tag{2.5}$$

lying between $\sigma_{\rm vac}^{\rm per}$ and $\sigma_{\rm vac}^{\rm phy}$, as illustrated by the short-dashed line in Fig. 1. The soliton solution has a spherical cavitylike structure: $\sigma\to\sigma_{\rm vac}^{\rm per}=0$ inside the cavity and $\sigma\to\sigma_{\rm vac}^{\rm phy}$ outside, with a smooth transition inbetween.

$f=3$ and $f=\infty$ are two limiting cases in which the soliton solution still exists. $f<3$ corresponds to the bag constant B being negative, which is prohibited physically. $f=\infty$ corresponds to the case that the local minimum at $\sigma_{\rm vac}^{\rm per}$ and the maximum at σ_m overlap, changing to an inflection point, as illustrated by the solid line in Fig. 1.

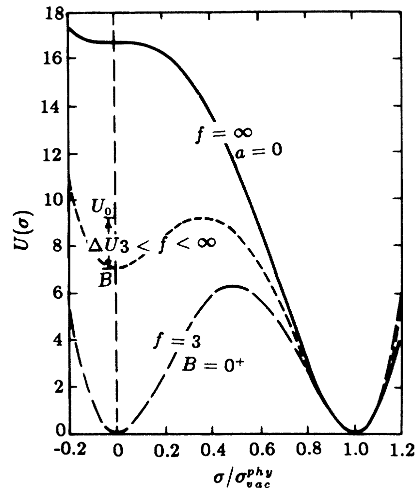

FIG. 1. Three forms of the potential function $U(\sigma)$ at zero temperature. $f=3$ and $f=\infty$ are the two limiting cases in which the soliton bag still exists. ΔU is the height of the potential barrier.

At this limiting case, the two vacuum states become metastable, very little disturbance will lead to the change of the vacuum structure. In order to ensure the stability of the two vacuum states and to guarantee the existence of the soliton solution, it is indispensable that $U(\sigma)$ has three distinct extrema. We will use the height of potential barrier

$$\Delta U=U_0-B \tag{2.6}$$

as a characteristic of the stability of this kind of vacuum structure, cf. Fig. 1. Obviously, the larger ΔU is, the stabler are the two vacuum states, and vice versa.

In this paper, we choose the color-dielectric function

$$\kappa(\sigma)=\left|1-\left[\frac{\sigma}{\sigma_{\rm vac}^{\rm phy}}\right]^2\right|^2 \tag{2.7}$$

given in Ref. 2. The form of $\kappa(\sigma)$ is restricted by the physical condition of the existence of the soliton solution, i.e., $\kappa(0)=1$ at the perturbative vacuum state and $\kappa(\sigma_{\rm vac}^{\rm phy})=0$ at the physical vacuum state. It can be seen that Eq. (2.7) satisfies these conditions.

III. EFFECTIVE POTENTIAL AT FINITE TEMPERATURE AND THE EQUATION OF ITS EXTREMA

Consider a thermodynamically equilibrium system with Lagrangian (2.1) and zero chemical potential at finite temperature. All the physical observables of this system can be described by the Gibbs average

$$\langle\ \cdots\ \rangle_\beta=\frac{\mathrm{Tr}[\exp(-\beta H)\cdots]}{\mathrm{Tr}[\exp(-\beta H)]}\ , \tag{3.1}$$

where $\beta=1/T$, and H is the Hamiltonian of the system.

Following the method proposed by Linde,[7] we perform a shift of the σ field with respect to the minimum $v(T)$ of the effective potential, i.e.,

$$\sigma=\sigma'+v(T)\ , \tag{3.2}$$

where $v(T)=\langle\sigma\rangle_\beta$. The shifted field σ' describes thermal excitation and satisfies the relation

$$\langle\sigma'\rangle_\beta=0\ . \tag{3.3}$$

Making use of the Hartree mean-field approximation[8]

$$\sigma'^3=3\langle\sigma'^2\rangle_\beta\sigma'\ , \tag{3.4}$$

we obtain from Eq. (2.3) the shifted field equation with the mass of the σ' field given by

$$m_{\sigma'}^2=a+bv(T)+\frac{c}{2}v^2(T)+\frac{c}{2}\langle\sigma'^2\rangle_\beta\ . \tag{3.5}$$

Taking the Gibbs average of the shifted field equation and noticing Eq. (3.3), we get

$$\partial^2 v(T)+\frac{c}{6}v^3(T)+\frac{b}{2}v^2(T)+\left[a+\frac{c}{2}\langle\sigma'^2\rangle_\beta\right]v(T)+\frac{b}{2}\langle\sigma'^2\rangle_\beta+g\langle\bar\Psi\Psi\rangle_\beta=0\ , \tag{3.6a}$$

$$[i\gamma_\mu\partial^\mu-gv(T)]\Psi=0\ . \tag{3.6b}$$

Equation (3.6b) gives the mass m_Ψ of the quark field as

$$m_\Psi=gv(T)\ . \tag{3.7}$$

Taking into account that $v(T)$ is a constant depending only on temperature, from Eq. (3.6a) we obtain the equation satisfied by the extrema of the effective potential $V(\sigma,T)$:

$$V'(\sigma,T)=\frac{c}{6}\sigma^3(T)+\frac{b}{2}\sigma^2(T)+\left[a+\frac{c}{2}\langle\sigma'^2\rangle_\beta\right]\sigma(T)+\frac{b}{2}\langle\sigma'^2\rangle_\beta+g\langle\bar\Psi\Psi\rangle_\beta=0\ , \tag{3.8}$$

where a prime denotes a derivative with respect to σ. In the above formula, we have substituted $\sigma(T)$ for $v(T)$. For $T=0$, $\langle\sigma'^2\rangle_\beta=\langle\bar\Psi\Psi\rangle_\beta=0$, the above formula is just the equation satisfied by $U(\sigma)$ at zero temperature, i.e., $U'(\sigma)=0$.

From Eq. (3.8), the effective potential at finite temperature can be obtained:

$$V(\sigma,T)=\frac{a}{2!}\sigma^2+\frac{b}{3!}\sigma^3+\frac{c}{4!}\sigma^4+\int d\sigma[\tfrac{1}{2}\langle\sigma'^2\rangle_\beta(b+c\sigma)+g\langle\bar\Psi\Psi\rangle_\beta]\ . \tag{3.9}$$

Obviously, the first three terms on the right-hand side give the $U(\sigma)$ at zero temperature; the last term represents the modification arising from the thermal excitation of the σ' field and the quark field.

Expanding the σ' field in terms of creation and annihilation operators, and noticing that the particle density in momentum space is

$$\langle a_{\mathbf p}^\dagger a_{\mathbf p}\rangle_\beta=1/[\exp(\sqrt{\mathbf p^2+m_{\sigma'}^2}\,/T)-1]\ , \tag{3.10}$$

where $\mathbf p$ is the momentum of the particle and $m_{\sigma'}$ is the mass of the σ' field given by Eq. (3.5), we obtain

$$\langle\sigma'^2\rangle_\beta=T^2f(m_{\sigma'}^2/T^2),\quad f(m_{\sigma'}^2/T^2)=\frac{1}{2\pi^2}\int_0^\infty\frac{x^2dx}{\sqrt{x^2+m_{\sigma'}^2/T^2}[\exp(\sqrt{x^2+m_{\sigma'}^2/T^2})-1]}\ . \tag{3.11}$$

When the chemical potential $\mu=0$, it can be shown[9] that

$$\langle\bar\Psi\Psi\rangle_\beta=-g\sigma T^2h(m_\Psi^2/T^2),\quad h(m_\Psi^2/T^2)=\frac{\gamma}{\pi^2}\int_0^\infty\frac{x^2dx}{\sqrt{x^2+m_\Psi^2/T^2}[\exp(\sqrt{x^2+m_\Psi^2/T^2})+1]}\ , \tag{3.12}$$

where γ is the total degeneracy factor,

$$\gamma=2(\text{spin})\times 2(\text{flavor})\times 3(\text{color})\ ;$$

m_Ψ is the mass of quark field given in Eq. (3.7).

IV. DISAPPEARANCE OF THE SOLITON SOLUTION AND DECONFINEMENT PHASE TRANSITION

In Sec. II we have pointed out that the existence of the soliton solution makes demands on $U(\sigma)$ having three extrema. At finite temperature, we know from the above section that the shape of the effective potential changes with the variation of temperature and the positions of its extrema shift correspondingly. It means that the extrema $\sigma_{\rm vac}^{\rm per}$, σ_m (2.5), and $\sigma_{\rm vac}^{\rm phy}$ given by Eq. (2.4) in Sec. II become temperature dependent and should be written as $\sigma_{\rm vac}^{\rm per}(T)$, $\sigma_m(T)$, and $\sigma_{\rm vac}^{\rm phy}(T)$, respectively. In order to investigate whether the soliton solution disappears at finite temperature, it is necessary only to check whether the number of extrema of the effective potential is reduced. In the following, we will do this in two different ways: first using a simple approximation and then doing a numerical calculation without approximation.

A. Approximate treatment

In order to show clearly how the shape of the effective potential and the position of its extrema change with temperature, we first make a high-temperature approximation for $\langle\sigma'^2\rangle_\beta$. As $T\gg m_{\sigma'}$, Eq. (3.11) becomes

$$\langle\sigma'^2\rangle_\beta=T^2/12\ . \tag{4.1}$$

The contribution from $\langle\bar\Psi\Psi\rangle_\beta$ has been neglected temporarily.

Under this approximation, the equation of extrema of the effective potential Eq. (3.8) becomes

$$\frac{c}{6}\sigma^3+\frac{b}{2}\sigma^2+\left[a+\frac{cT^2}{24}\right]\sigma+\frac{bT^2}{24}=0 \tag{4.2}$$

and the effective potential itself Eq. (3.9) becomes

$$V(\sigma,T)=\frac{bT^2}{24}\sigma+\frac{1}{2!}\left[a+\frac{cT^2}{24}\right]\sigma^2+\frac{b}{3!}\sigma^3+\frac{c}{4!}\sigma^4+C(T)\ , \tag{4.3}$$

where $C(T)$ is a constant depending only on temperature and has no influence on the shape of the potential function. We choose $C(T)$ to satisfy $V(\sigma_{\rm vac}^{\rm phy}(T),T)=0$.

Setting

$$\sigma=y-\frac{b}{c}\ , \tag{4.4}$$

Eq. (4.2) is reduced to a cubic equation in standard form

$$y^3+py+q=0\ , \tag{4.5}$$

where

$$p=\frac{24ac+c^2T^2-12b^2}{4c^2}\ ,$$

$$q=\frac{2b^3-6abc}{c^3}\ .$$

From the discriminant of the cubic equation

$$\Delta(T)=\left[\frac{q}{2}\right]^2+\left[\frac{p}{3}\right]^3\ , \tag{4.6}$$

we know that

$$\Delta(T)\begin{cases}<0 & \text{three real and unequal roots ,}\\ =0 & \text{three real roots of which at least two are equal ,}\\ >0 & \text{one real root and two conjugate imaginary roots .}\end{cases}$$

Figure 2 shows the result of the investigation of the relation between $\Delta(T)$ and temperature, using the values of the set of parameters a,b,c with $3<f<\infty$, which can give the static properties of the hadrons. In Figs. 3(a) and 3(b) are drawn the shape of the effective potential Eq. (4.3) at different temperatures and the temperature dependence of the position of the extrema satisfying Eq. (4.2), respectively.

From these figures we clearly see that, when $T=0$, $\Delta<0$ and the effective potential has three extrema corresponding to $\sigma_{\rm vac}^{\rm per}$, σ_m, and $\sigma_{\rm vac}^{\rm phy}$, respectively. As $T=T_c^-$, two of these three extrema coincide: The minimum $\sigma_{\rm vac}^{\rm per}(T_c^-)$ overlaps with the maximum $\sigma_m(T_c^-)$. At this moment, the shape of the effective potential is similar to that of $U(\sigma)$ with $f=\infty$ at zero temperature, as illustrated in Fig. 3(a). This is a limiting case in which the soliton solution still exists. When $T\geq T_c^+$, the effective potential has a unique minimum only, the soliton solution disappears, and the deconfinement transition occurs. It follows that T_c is the transition temperature and the discriminant (4.6) may be chosen as an "order parame-

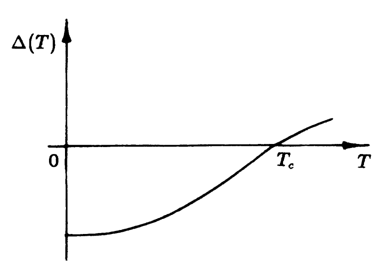

FIG. 2. The relation between order parameter $\Delta(T)$ and temperature.

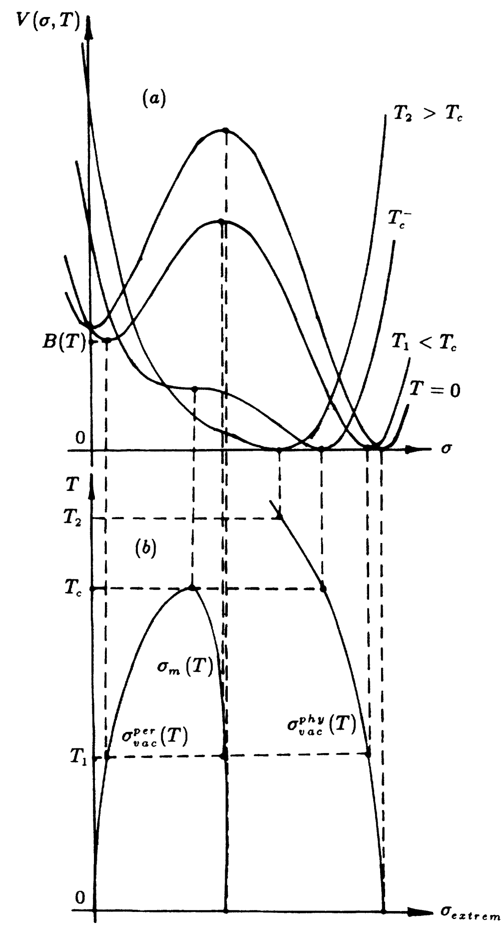

FIG. 3. (a) The effective potential at different temperatures in a high-temperature approximation with the contribution of fermions omitted. (b) The temperature dependence of the positions of the extrema of the effective potential in the same approximations.

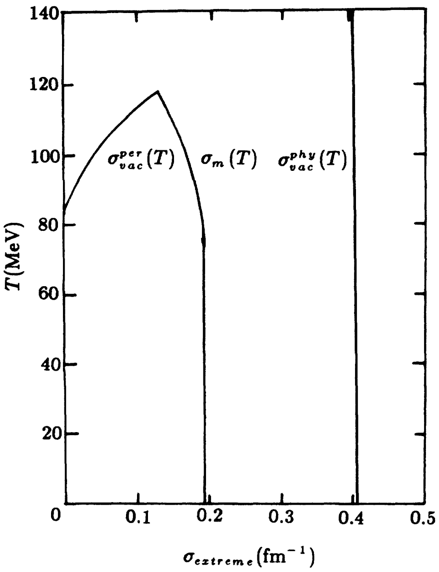

FIG. 4. The temperature dependence of the positions of the extrema of the effective potential without making the approximations in Fig. 3. The model parameters are chosen to be $a=51.6\ \mathrm{fm}^{-2}$, $b=-799.9\ \mathrm{fm}^{-1}$, $c=4000$, and $g=14.8$.

ter" to describe the deconfinement transition.

It can also be seen from Fig. 3(a) that the pressure of the physical vacuum on a bag,

$$B(T)=V(\sigma_{\mathrm{vac}}^{\mathrm{per}}(T),T)\ , \tag{4.7}$$

decreases slightly with the increasing of temperature, so that the radius of the soliton bag at thermal equilibrium increases slightly. However, when $T=T_c^-$, $B(T_c^-)$ still remains finite, and as $T=T_c^+$, the structure of the vacuum changes and the bag constant B has no definition at all. In this respect, the present deconfinement picture is different from the previous ones based on the unlimited expansion of the bag.

B. Results of numerical calculation

In view of the difficulty in analytical calculation without approximation, we have performed a numerical calculation, taking the thermal excitation Eq. (3.11) of the σ' field and Eq. (3.12) of the quark field simultaneously into account. The result tells us that the conclusions obtained in Sec. IV A do not change qualitatively. The transition temperature T_c turns out to be strongly dependent on the model parameters g, a, b, and c. Up to now, a good set of parameters for fitting all the static properties of the hadrons has not been found for the nontopological soliton model. We choose the following set of parameters:[2] $a=51.6\ \mathrm{fm}^{-2}$, $b=-799.9\ \mathrm{fm}^{-1}$, $c=4000$, and $g=14.8$, which is satisfactory for fitting a part of the static properties of the hadrons. The transition temperature T_c corresponding to this set of parameters is 120 MeV. The temperature dependence of the positions of the extrema of the effective potential is illustrated in Fig. 4.

In Sec. II we have pointed out that, at zero temperature, the height of the potential barrier ΔU describes the stability of the vacuum structure, possessing both physical and perturbative vacuum states. The larger ΔU is the stabler are the two vacuum states and so the higher will be the transition temperature. This conclusion is proved by numerical results. Table I lists the values of T_c and the corresponding ΔU for six sets of parameters taken from Ref. 2.

V. COLOR-DIELECTRIC FUNCTION AND CHANGE OF VACUUM STRUCTURE

At zero temperature, Lee's analysis[1] shows that the width of the transition region of the soliton solution is $\delta=U''(0)^{-1/2}$. We generalize this result to the case of $0<T<T_c$, for which the effective potential has three extrema and the soliton solution still exists. The width of transition region in this case can be written as

$$\delta(T)=[V''(\sigma_{\mathrm{vac}}^{\mathrm{per}}(T),T)]^{-1/2}\ . \tag{5.1}$$

The result of the calculation shows that $\delta(T)$ increases with an increasing temperature, as shown in Fig. 5. In view of this, we can give a rough sketch of the space variation of $\sigma(T)$ as shown in Fig. 6(a), where r is the radius vector and R is the radius of the bag. Note that, when $T>0$, $\sigma(T)|_{r=0}=\sigma_{\mathrm{vac}}^{\mathrm{per}}>0$, as can be seen in Fig. 3.

The existence of the soliton solution at $0<T<T_c$ indicates that the confinement of the quark and gluon is still present in this case, and the color-dielectric function remains to be $\kappa=1$ at the perturbative vacuum state and $\kappa=0$ at the physical vacuum state. On the other hand, when $T>T_c$, the confinement of quarks is removed, color antiscreening disappears, and $\kappa=1$ everywhere. In order to satisfy this requirement, we generalize Eq. (2.7) and define the color-dielectric function at finite temperature as

TABLE I. Values of T_c and corresponding ΔU for six sets of parameters taken from Ref. 2.

a (fm^{-2})	$-b$ (fm^{-1})	c	g	ΔU (MeV/fm^3)	T_c (MeV)
7.671	107.27	500	13.09	8.70	90
12.85	196.34	1000	13.06	12.22	92
40.88	783.80	5000	14.01	24.73	104
66.42	1411.60	10 000	14.83	32.65	108
107.32	2537.00	20 000	15.94	42.65	114
321.75	9824.80	100 000	19.77	76.62	118

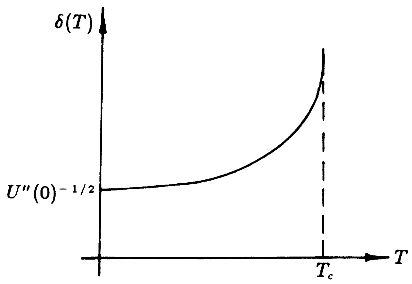

FIG. 5. The relation between the width of the transition region of the soliton and temperature.

$$\kappa(\sigma,T)=\begin{cases} \left|1-\left[\dfrac{\sigma-\sigma_{\rm vac}^{\rm per}(T)}{\sigma_{\rm vac}^{\rm phy}(T)-\sigma_{\rm vac}^{\rm per}(T)}\right]^2\right|^2 & \text{if } T\le T_c\ , \\ 1 \quad \text{if } T>T_c\ . \end{cases} \tag{5.2}$$

The relation between the color-dielectric function and the radius vector r at different temperatures is shown in Fig. 6(b).

Because the transition region connecting the physical vacuum state and the perturbative vacuum state widens with the increasing of temperature and $0<\kappa(\sigma,T)<1$ in this region, the transition of the two vacuum states becomes more and more smooth (cf. Figs. 6), leading to the instability of the vacuum structure at critical temperature T_c. At $T=T_c^+$, the structure of the vacuum changes, the physical vacuum is transformed into a perturbative vacuum, the soliton solution disappears, and the deconfinement transition occurs. The change of the vacuum structure is an abrupt process, indicating that the deconfinement transition is a first-order phase transition.

VI. CONCLUSION

In this paper, we have investigated the deconfinement transition of the nontopological soliton model at finite temperature. We get a new picture for describing the deconfinement transition, which is different from the common one with unlimitedly expanding bags.

The confinement of the quark and gluon requires the existence of the soliton solution, and the latter depends on the number of extrema of the effective potential. In view of this, we have given the effective potential at finite temperature and the relation between the positions of its extrema and the temperature. It turns out that, two of these extrema of the potential function at zero temperature overlap at a certain temperature T_c. When $T\ge T_c^+$, the effective potential has a unique minimum only. This causes the physical vacuum to transform into a perturbative vacuum, the soliton bag disappears, and the deconfinement transition takes place.

Numerical results show that the transition temperature T_c depends strongly on the model parameters. For a reasonable set of values of the parameters, the transition temperature at zero chemical potential is about 120 MeV. This result is slightly lower than the prediction of lattice gauge theory.

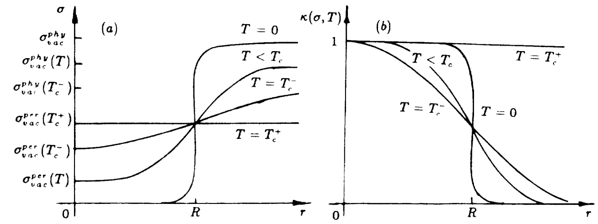

FIG. 6. (a) The soliton solutions at different temperatures. (b) Color-dielectric function at different temperatures.

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