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One π-meson and One σ-meson Exchange Potentials at Finite Temperatures

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On the basis of the field theory at finite temperature, the dependence of π-meson and σ-meson effective masses upon temperatures at nucleon level and one π and one σ exchange potentials are discussed. The dependence of the π and σ effective masses upon temperatures is found to be in good agreement with the results given by the Nambu-Jone-Lasinia model at quark level.

1. INTRODUCTION

Many authors have recently made in-depth studies on what will happen when a hadronic matter is heated. Some [1-3] have found that the confinement of quarks does not work, and the chiral symmetry of weak interactions, which is spontaneously broken at zero temperature, may be restored at high temperatures and/or at high densities. Others [4-6] have shown a liquid-gas phase transition and a superconducting phase transition in nuclear matter at high temperatures. Experimentally, the critical temperature of the liquid-gas phase transition [7--8] being in good agreement with the theoretical value has been measured and the other phase transition is hoped to be realized in relativistic heavy-ion collisions.

π-meson and σ-meson play an important role in the dynamics of hadrons at low energies. Therefore, the dynami properties of π and σ at finite temperatures have also been investigated. Bernard, Meissner, Zahed [9] and other authors [10] used the finite temperature Nambu-Jone-Lasinia model to treat it. They found that π-meson becomes heavy and tends to be decoupled with

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quarks and that σ-meson becomes light as temperature is below 215 MeV. According to field theory at a finite temperature, Chen and Su [11] have discussed one-pion-exchange potential (OPEP) at nucleon level. However, in Ref. [11], the isospin effect of π is not taken into account.

In this paper we wish to derive one π and one σ exchange potentials at finite temperatures by taking the isospin of π into account. As is well known, nucleon-nucleon (NN) interaction arises from one-π-exchange at large distances, from two-π-exchange or one-σ-exchange at intermediate distances and three-π-exchange or one-vector meson-exchange (for example, ω-meson) at short distances [12]. The one-σ-exchange potential needs to be extended to finite temperatures; therefore the thermal properties of the attractive part of nuclear forces (namely the sum of the forces arising at intermediate and large distances) may be discussed. At the same time, we attempt to study M_σ, the effective mass of σ-meson and its temperature dependence. Bernard *et al.* [9,10] have studied this problem at quark level. Can we obtain at nucleon level a result agreeable with Bernard's? This is the second problem with which we are concerned.

The organization of this paper is as follows. In Sec.2 we shall attempt to work out the effective masses of π and σ by using the imaginary time Green's function method in finite temperature field theory. In Sec.3 we shall derive one π and σ exchange potential. Our results are presented and discussed in Sec.4.

2. EFFECTIVE MASSES OF π AND σ

In this section we proceed to discuss $M_\pi(\beta)$, the effective mass of π (its isospin $T = 1$) and $M_\sigma(\beta)$, the one of σ at finite temperatures.

2.1 The Effective Mass of π Involving the Isospin Effect

Considering that the effective mass of π has been derived in Ref. [11] after neglecting its isospin effect, here we only need a short extension. Some essential results are presented below.

1). Interaction lagrangian of π-nucleon coupling reads

$$\mathcal{L}_I^{\pi\psi} = i g_p \bar{\psi}(x)\gamma_5 \boldsymbol{\tau}\psi(x)\cdot\boldsymbol{\phi}(x), \tag{2.1}$$

where ψ and $\boldsymbol{\phi}$ denote nucleon and pion fields respectively, $\boldsymbol{\tau}$ is the isospin Pauli matrix and g_p their coupling constant.

2). The OPEP at finite temperature obtained by summing up the bubble diagrams in the centre of the mass coordinates of the two-nucleon system is

$$V^\beta_{\pi\psi}(p, p') = -g_p^2 \bar{u}(p')\gamma_5\tau_i u(p)\,\frac{1}{q^2 - m_\pi^2 - i\pi^\beta_{ij}(q^2)}\cdot \bar{u}(-p')\gamma_5\tau_i u(-p), \tag{2.2}$$

where m_π stands for the rest mass of π, P and P' are the relative four-momentum in initial and final states respectively, $q = p - p'$. $\pi^\beta_{ij}(q^2)$ are the pion self-energies (see Fig.1(a)), i and j are the pion isospin indices, $i, j = 1, 2, 3$:

$$\begin{aligned}\pi^\beta_{ij}(q^2) &= -\frac{i}{\beta}g_p^2\sum_n\int\frac{d^3k}{(2\pi)^3}\mathrm{Tr}\left[\gamma_5\frac{i}{\not{k}-m}\gamma_5\frac{i}{(\not{k}-\not{q})-m}\right]\cdot\mathrm{Tr}(\tau_i\tau_j)\\ &= -i\delta_{ij}\frac{4g_p^2}{\pi^2\beta^2}\left[I_1 + q^2\beta^2\left(\frac{1}{4}I_2+\frac{1}{6}I_3\right)\right],\end{aligned} \tag{2.3}$$

where

$$\begin{cases} I_1 = \int_0^\infty dx \dfrac{x^2}{w(e^{w} + 1)} \\ I_2 = \int_0^\infty dx \dfrac{x^2}{w^3(e^{w} + 1)} \\ I_3 = \int_0^\infty dx \dfrac{x^2 e^{w}}{w^2(e^{w} + 1)^2}, \end{cases} \tag{2.4}$$

and $w = \sqrt{x^2 + a^2}$, $a = m\beta = m/k_B T$, m is the mass of nucleon, k_B and T denote the Boltzmann constant and temperature respectively.

3). The effective mass of π is

$$M_\pi(\beta) = \sqrt{\frac{m_\pi^2 + 4g_p^2 I_1/(\pi\beta)^2}{1 - 4g_p^2\left(\frac{1}{4} I_2 + \frac{1}{6} I_3\right)/\pi^2}}. \tag{2.5}$$

2.2 The Effective Mass of σ in the One-σ-exchange Potential

σ is a scalar meson, the interaction lagrangian of σ-nucleon coupling is

$$\mathscr{L}_I^{\sigma\psi} = g_s \bar{\psi}(x)\psi(x)\varphi(x), \tag{2.6}$$

where ϕ denotes the σ field and g_s denotes the coupling constant. The proper propagator of σ and the one-σ-exchange potential at finite temperature obtained by summing up the bubble diagrams are respectively

$$\begin{aligned} i\Delta'_\sigma(q) &= i\Delta_\sigma(q) + i\Delta_\sigma(q)\pi(q^2)i\Delta_\sigma(q) \\ &\quad + i\Delta_\sigma(q)\pi(q^2)i\Delta_\sigma(q)\pi(q^2)i\Delta_\sigma(q) + \cdots \\ &= \frac{i}{q^2 - m_\sigma^2 - i\pi(q^2)}, \end{aligned} \tag{2.7}$$

$$V_{\sigma\psi}^{\beta}(p, p') = g_s^2 \bar{u}(p')u(p) \frac{1}{q^2 - m_\sigma^2 - i\pi(q^2)} \bar{u}(-p')u(-p). \tag{2.8}$$

where $i\Delta_\sigma(q) = i/(q^2 - m_\sigma^2)$ is the free propagator of σ. $\pi(q^2)$ stands for the second-order σ self-energy (see Fig.1(b)). It may be expressed analytically as

$$\begin{aligned} \pi(q^2) &= -g_s^2 \frac{i}{\beta} \sum_n \int \frac{d^3k}{(2\pi)^3} \mathrm{Tr}\left[\frac{1}{\not{k} - m} \cdot \frac{1}{(\not{k} - \not{q}) - m}\right] \\ &= -\frac{i4g_s^2}{\beta} \sum_n \int \frac{d^3k}{(2\pi)^3} \frac{m^2 + k^2 - k \cdot q}{(k^2 - m^2)[(k - q)^2 - m^2]}. \end{aligned} \tag{2.9}$$

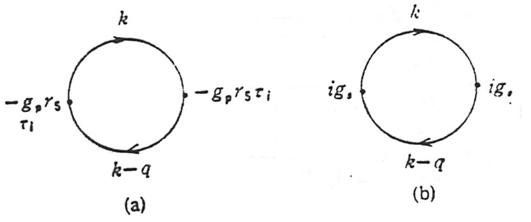

FIG.1
Two-order π and σ self-energies.

With the help of the Feynman identity

$$\frac{1}{ab} = \int_0^1 d\lambda \frac{1}{[a\lambda + b(1-\lambda)]^2} \tag{2.10}$$

Eq.(2.9) reduces to

$$\pi(q^2) = -\frac{i4g_i^2}{\beta}\int\frac{d^3k}{(2\pi)^3}\int_0^1 d\lambda \sum_n \frac{m^2 + k^2 - q^2(1-\lambda)\lambda + (1-2\lambda)\omega_n q^0}{[k^2 - m^2 + q^2(1-\lambda)\lambda]^2}, \tag{2.11}$$

where $k^0 = i\omega_n = i(2n+1)\pi/\beta$. In term of the notation $F = (k^2 + m^2 - q^2(1-\lambda)\lambda)^{\frac{1}{2}}$, Eq.(2.11) is given by

$$\pi(q^2) = -\frac{i4g_i^2}{\beta}\int\frac{d^3k}{(2\pi)^3}\int_0^1 d\lambda\{\Sigma_1 + [2m^2 - 2q^2(1-\lambda)\lambda]\cdot\Sigma_2 + (1-2\lambda)q^0\Sigma_3\}, \tag{2.12}$$

with

$$\begin{cases} \Sigma_1 = \sum_n \frac{1}{(i\omega_n)^2 - F^2} = \frac{1}{2F}\sum_n\left(\frac{1}{i\omega_n - F} - \frac{1}{i\omega_n + F}\right) \\ \Sigma_2 = \sum_n \frac{1}{[(i\omega_n)^2 - F^2]^2} = \sum_n\left\{\frac{1}{4F^3}\left(\frac{1}{i\omega_n + F} - \frac{1}{i\omega_n - F}\right)\right. \\ \qquad \left. + \frac{1}{4F^2}\left[\frac{1}{(i\omega_n + F)^2} + \frac{1}{(i\omega_n - F)^2}\right]\right\} \\ \Sigma_3 = \sum_n \frac{i\omega_n}{[(i\omega_n)^2 - F^2]^2} = \frac{1}{4F}\sum_n\left[\frac{1}{(i\omega_n - F^2)} - \frac{1}{(i\omega_n + F)^2}\right]. \end{cases} \tag{2.13}$$

Eq.(2.13) includes four types of summation over n, the discrete subscript, which always arises in quantum field theory at finite temperatures. Employing the contour integral in complex plane used in Ref. [13], we get

$$\begin{cases} \sum\limits_n \dfrac{1}{i\omega_n - F} = \dfrac{\beta}{e^{\beta F} + 1} \\ \sum\limits_n \dfrac{1}{i\omega_n + F} = \beta - \dfrac{\beta}{e^{\beta F} + 1} \\ \sum\limits_n \dfrac{1}{(i\omega_n - F)^2} = \sum\limits_n \dfrac{1}{(i\omega_n + F)^2} = - \dfrac{\beta^2 e^{\beta F}}{(e^{\beta F} + 1)^2}. \end{cases} \tag{2.14}$$

Thus we have

$$\begin{cases} \Sigma_1 = \dfrac{\beta}{2F}\left(\dfrac{2}{e^{\beta F} + 1} - 1\right) \\ \Sigma_2 = \dfrac{\beta}{4F^3}\left(1 - \dfrac{2}{e^{\beta F} + 1}\right) - \dfrac{1}{2F^2}\dfrac{\beta^2 e^{\beta F}}{(e^{\beta F} + 1)^2} \\ \Sigma_3 = 0. \end{cases} \tag{2.15}$$

Inserting Eq.(2.15) into (2.12), we obtain

$$\pi(q^2) = \pi_0(q^2) + \pi_\beta(q^2), \tag{2.16}$$

where $\pi_0(q^2)$ is temperature independent and has the form:

$$\pi_0(q^2) = i2g_i^2 \int \frac{d^3k}{(2\pi)^3} \int_0^1 d\lambda \left\{ \frac{1}{F} - \frac{1}{F^3}[m^2 - q^2(1-\lambda)\lambda] \right\}, \tag{2.17}$$

$\pi_B(q^2)$ is temperature dependent, in the case $q^2(1-\lambda)\lambda < k^2 + m^2$, $\pi_B(q^2)$ may be written as

$$\pi_\beta(q^2) = -\frac{i2g_i^2}{\pi^2}\int_0^\infty dx\, x^2 \left\{ \left[\frac{1}{\beta^2 w} + \frac{q^2}{12w^3} - \frac{m^2 - \frac{1}{6}q^2}{w^3} - \frac{\frac{1}{2}m^2q^2\beta^2}{2w^5} \right] \cdot \frac{1}{e^w + 1} - \left[\frac{m^2 - \frac{1}{6}q^2}{w^2} + \frac{\frac{1}{6}m^2q^2\beta^2}{w^4} \right] \cdot \frac{e^w}{(e^w + 1)^2} \right\}. \tag{2.18}$$

We can prove that $\pi_B(q^2) \to 0$ in the limit $T \to 0°k$. Hence, $\pi_0(q^2)$ being a divergent integral is a contribution at zero temperature. According to the renormalization theory, it maý be absorbed by the mass term of σ. For simplicity, we shall regard m_σ in the following as a renormalized mass. At low temperatures we get

$$\pi(q^2) = -i\frac{2g_i^2}{\pi^2}(A + Bq^2), \tag{2.19}$$

with

$$\begin{cases} A = I_1 - m^2\beta^2(I_2 + I_3) \\ B = \frac{1}{4} I_2 + \frac{1}{6} I_3 - m^2\beta^2\left(\frac{1}{4} I_5 + \frac{1}{6} I_4\right), \end{cases} \tag{2.20}$$

where $I_1 - I_3$ refer to (2.4), I_4 and I_5 are defined as

$$\begin{cases} I_4 = \int_0^\infty dx \frac{x^2 e^w}{w^4(e^w + 1)^2} \\ I_5 = \int_0^\infty dx \frac{x^2}{w^5(e^w + 1)}. \end{cases} \tag{2.21}$$

Consequently, the effective mass of σ is

$$M_\sigma(\beta) = \sqrt{\frac{m_\sigma^2 + 2g_s^2 A/(\pi\beta)^2}{1 - 2g_s^2 B/\pi^2}}. \tag{2.22}$$

3. ONE π AND ONE σ EXCHANGE POTENTIALS AT FINITE TEMPERATURES

It is convenient to determine from Eq.(2.2) the OPEP in coordinate space involving the effect of π isospin. It is given by multiplying a factor $(\tau_1 \cdot \tau_2)$ to the corresponding result given in Ref. [11], namely

$$V^\beta_{\pi\psi}(r) = \frac{g_p^2}{1 - 4g_p^2\left(\frac{1}{4} I_2 + \frac{1}{6} I_3\right)\Big/\pi^2} \cdot \frac{M_\pi}{4\pi} \cdot \frac{M_\pi^2}{12m^2} \cdot [Z(x)S_{12} + Y(x)(\sigma_1 \cdot \sigma_2)](\tau_1 \cdot \tau_2), \tag{3.1}$$

where $x = M_\pi r$, $Y(x) = e^{-x}/x$, $Z(x) = (1 + 3/x + 3/x^2)Y(x)$, S_{12} denotes the tenser force as usual.

In the following, we want to derive the one-σ-exchange potential from (2.8). In Eq.(2.9) $u(p)$ takes the form

$$u(p) = \sqrt{\frac{p_0 + m}{2m}} \begin{pmatrix} I \\ \frac{\sigma \cdot p}{p_0 + m} \end{pmatrix} \chi_{\frac{1}{2}},$$

where $\chi_{\frac{1}{2}} = \begin{pmatrix} 1 \\ 0 \end{pmatrix}$ is the spinor. In elastic scattering $p'_0 = p_0$, $p^2 = p'^2$, so

$$V^{\beta}_{\sigma\psi}(p, p') = g_s^2\left(\frac{p_0+m}{2m}\right)^2 \chi^+_{\frac{1}{2}}(1)\chi^+_{\frac{1}{2}}(2)\left[1-\frac{(\sigma_1\cdot p')(\sigma_1\cdot p)}{(p_0+m)^2}\right] \cdot \frac{1}{q^2-m_\sigma^2-i\pi(q^2)}\left[1-\frac{(\sigma_2\cdot p')(\sigma_2\cdot p)}{(p_0+m)^2}\right]\chi_{\frac{1}{2}}(1)\chi_{\frac{1}{2}}(2). \tag{3.2}$$

According to the discussion exhibited in Ref. [12,15], Eq.(3.2) implies

$$\int V^{\beta}_{\sigma\psi}(r)e^{-iq\cdot r}dr = g_s^2\left(\frac{p_0+m}{2m}\right)^2\left[1-\frac{(\sigma_1\cdot p')(\sigma_1\cdot p)}{(p_0+m)^2}\right]\left[1-\frac{(\sigma_i\cdot p')(\sigma_2\cdot p)}{(p_0+m)^2}\right] \cdot \frac{1}{q^2-m_\sigma^2-i\pi(q^2)}. \tag{3.3}$$

Making a Fourier transformation, we can obtain the one-σ-exchange potential in coordinate space

$$V^{\beta}_{\sigma\psi}(r) = \int\frac{d^3q}{(2\pi)^3}\left\{g_s^2\left(\frac{p_0+m}{2m}\right)^2\left[1-\frac{p'\cdot p}{2m^2} - \frac{i}{4m^2}(\sigma_1+\sigma_2)\cdot(p'\times p)\right]\frac{1}{q^2-m_\sigma^2-i\pi(q^2)}\right\}e^{iq\cdot r} \tag{3.4}$$

In the non-relativistic limit $q \to 0$, $p_0 = p'_0 = m$, so

$$\left(\frac{p_0+m}{2m}\right)^2 \cong \left(\frac{p_0+m}{2p_0}\right)^2 = \frac{1}{4}\left[1+\frac{1}{\sqrt{p^2/m^2+1}}\right]^2 \cong \frac{1}{4}\left(1+\frac{1}{1+p^2/2m^2}\right)^2 = \left(1-\frac{p^2}{4m^2}\right)^2 \cong 1-\frac{p^2}{2m^2} = 1-\frac{p^2+p'^2}{4m^2}. \tag{3.5}$$

Since $q = p - p'$ and $k = 1/2(p+p')$, therefore

$$V^{\beta}_{\sigma\psi}(r) \cong \int\frac{d^3q}{(2\pi)^3}\left\{g_s^2\left[1-\frac{k^2}{m^2}-\frac{i}{4m^2}(\sigma_1+\sigma_2)\cdot(k\times q)\right] \cdot \frac{1}{q^2-m_\sigma^2-i\pi(q^2)}e^{iq\cdot r}\right\}. \tag{3.6}$$

Eq.(2.22) means that the σ propagator can be written as

$$\frac{1}{q^2-m_\sigma^2-i\pi(q^2)} = -\frac{1}{1-2g_s^2B/\pi^2}\cdot\frac{1}{q^2+M_\sigma^2}. \tag{3.7}$$

Then since

$$q = -i\nabla,\ p' = i\nabla,\ p = i\nabla \tag{3.8}$$

Eq.(3.6) is given by

$$V^{\beta}_{\sigma\psi}(r) = \frac{-g_s^2}{1 - 2g_s^2 B/\pi^2}\left\{1 + \frac{(\nabla + \nabla)^2}{4m^2} - i\frac{1}{8m^2}(\sigma_1 + \sigma_2)[(\nabla + \nabla) \times \nabla]\right\} \times \int \frac{d^3q}{(2\pi)^3}\frac{1}{q^2 + M_\sigma^2} e^{iq\cdot r}. \tag{3.9}$$

With straightforward integrating, we find

$$\int \frac{d^3q}{(2\pi)^3}\frac{1}{q^2 + M_\sigma^2} e^{iq\cdot r} = \frac{M_\sigma}{4\pi} Y(M_\sigma r). \tag{3.10}$$

Using the formulae [15]

$$\begin{cases} \nabla f\psi = \nabla(f\psi) \\ \nabla f\psi = f\nabla\psi, \end{cases}$$

We can prove

$$V^{\beta}_{\sigma\psi}(r) = -\frac{M_\sigma}{4\pi}\frac{g_s^2}{1 - 2g_s^2 B/\pi^2}\left\{Y(M_\sigma r) + \frac{1}{2m^2}(\nabla^2 Y + Y\nabla^2) - \frac{1}{4m^2}\left[M_\sigma^2 Y(M_\sigma r) - \frac{4\pi}{M_\sigma}\delta(r)\right] - \frac{1}{2m^2}\boldsymbol{S}\cdot\boldsymbol{L}\frac{1}{r}\frac{d}{dr}Y(M_\sigma r)\right\}, \tag{3.11}$$

where $\boldsymbol{S} = 1/2(\boldsymbol{\sigma}_1 + \boldsymbol{\sigma}_2)$ and $\boldsymbol{L}$ is the orbital angular momentum of the relative motion of the two nucleons. In the intermediate range $\delta(r) = 0$. By dropping the momentum dependent terms (except for the spin-orbit coupling term), we have

$$V^{\beta}_{\sigma\psi}(r) = -\frac{M_\sigma}{4\pi}\frac{g_s^2}{1 - 2g_s^2 B/\pi^2}\left[\left(1 - \frac{M_\sigma^2}{4m^2}\right)Y(\tilde{x}) - \frac{M_\sigma}{2m^2}\boldsymbol{S}\cdot\boldsymbol{L}\frac{1}{\tilde{x}}\frac{d}{d\tilde{x}}Y(\tilde{x})\right], \tag{3.12}$$

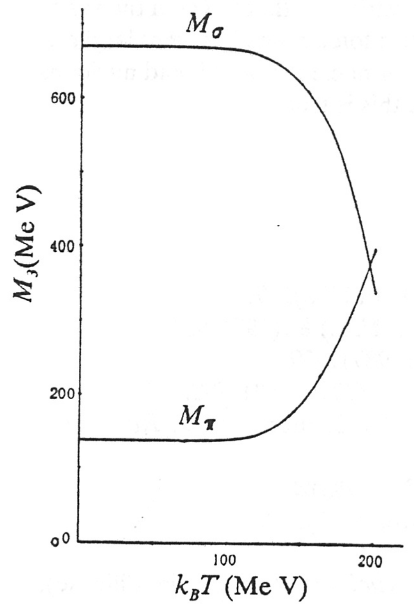

FIG.2
The temperature dependences of M_π the π effective mass and M_σ the σ effective mass.

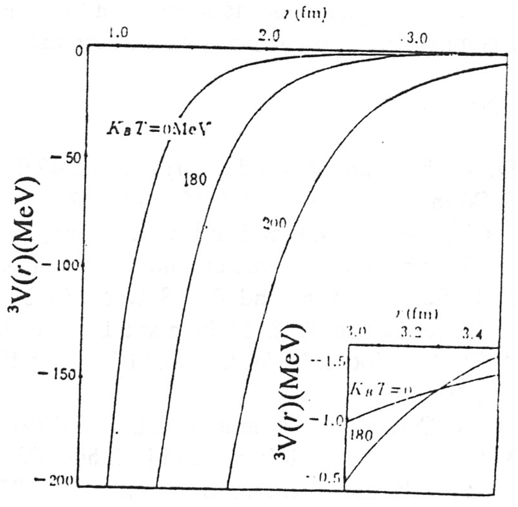

Fig.3
The attractive part of centrical NN interaction potentials for $T = 0$ and $S = 1$ NN states.

where $\tilde{x} = M_\sigma r$. In the limit $T \to 0$, all the integral $I_1 - I_5$ approach to zero, and (3.12) reduces to the one-σ-exchange potential at zero temperature as displayed in Ref. [12,15].

4. RESULTS AND DISCUSSIONS

(1) At low temperatures ($k_B T \leq 200$ MeV) the dependence of the π and σ effective masses on temperature is shown in Fig.2. We find that the M_π and M_σ curves are in good agreement with the ones in Case II given in Ref. [10]. The effective masses of π and σ become heavier and lighter respectively as temperature T increases. The approach of Ref. [10] is based on the NJL model at quark level, whereas ours is based on the interaction lagrangians of ordinary Yukawa coupling at nucleon level. Hence, the agreement of the above results are physically meaningful. It indicates that the quantum field theory at finite temperatures is suitable to deal with some correlative work at nucleon level.

(2) Attractive part of the centrical NN interaction potentials arise from one-π-exchange at large distances, and from one-σ-exchange at intermediate distances. Hence, we set

$$V^\beta(r) = V^\beta_{\pi\psi}(r) + V^\beta_{\sigma\psi}(r) \tag{4.1}$$

and compute ${}^3V(r)$, the attractive potential for iso-spin $T = 0$ and spin $S = 1$ NN states. The numerical results are displayed in Fig.3. We find the ${}^3V(r)$ curves in the range $r < 3.3$ fm drop as the temperature T increases. This indicates that the force range becomes larger when the temperature is higher. We also find that the ${}^3V(r)$ curves in the range $r > 3.3$ fm rise as T increases (see the magnified plot at the right-down corner of Fig.3). This indicates that the force range becomes shorter and the intensity weaker. Consequently, we find from the intensity variances that ${}^3V(r)$ in the range $r < 3.3$ fm becomes essential as T rises. That is to say, the attractive force range becomes larger as T rises. This result implies that the increase of the kinetic energy of nucleons would lead nucleons in a nucleus more dispersive as temperature increases. Obviously, this is true.

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