Coulomb instability of hot nuclei in a chiral symmetry model

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A chiral symmetry model for nuclear matter is extended to finite temperature and then used to study the Coulomb instability of hot nuclei. The role of the ρ meson degree of freedom is discussed. The calculated limiting temperature is compared with the results obtained from other theories and experiments.
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I. INTRODUCTION

There are various models, either relativistic or nonrelativistic, to describe the nuclear matter properties. It was already suggested in the early days of nuclear physics that the nucleon-nucleon interaction is mediated by the exchange of finite mass mesons. People thus developed the effective meson exchange models whose parameters are directly adjusted to nuclear matter properties. The Walecka model (also known as the QHD model) [1] is a typical one belonging to this kind. Despite its simplicity, it is still successful in describing the properties of both infinite nuclear matter and finite nuclei. However, it has too large a compression modulus ($\sim$540 MeV). There have been many methods to modify this model, such as the nonlinear model [2], ZM model [3], etc.

Although the effective models mentioned above are very successful, the descriptions of nuclear matter and finite nuclei are ultimately governed by the physics of low-energy quantum chromodynamics (QCD). In the absence of direct derivations from QCD, the effective hadronic models should be constrained by the underlying symmetries of QCD, such as chiral symmetry and broken scale invariance. There is a long history of attempts to generalize the linear sigma model to build models with chiral symmetry. Recently, Furnstahl, Serot, and Tang worked out a relativistic hadronic model (referred to as the FST model hereafter) [4–6], which incorporates nonlinear chiral symmetry, broken scale invariance, and the phenomenology of vector dominance. An important feature of the new model is the light scalar degree of freedom, which is given an anomalous scale dimension. The parameter sets determined by fitting to the properties of finite nuclei give also the properties of infinite nuclear matter in good agreement with the empirical values. For example, with parameter set $T1$, they obtained a saturation density of about 0.15 fm^{-3}, with a binding energy per nucleon of about 16 MeV. The nucleon effective mass M^*/M and the compression modulus at saturation density are about 0.6 and 200 MeV, respectively.

The equation of state (EOS) of nuclear matter impinges on a number of areas of physics, such as the monopole resonance, high energy nuclear collision, supernovae, and neutron stars. Therefore, the equation of state is a many dimensional function. It is both natural and interesting to extend this model to finite temperature and to study the thermodynamical properties of hadronic matter. Since the hot nuclear system formed in heavy ion collisions is of finite size, asymmetric, and charged, an important temperature, namely, the limiting temperature $T_{\rm lim}$ is introduced [7]. Below the limiting temperature $T_{\rm lim}$, the nucleus can exist in equilibrium with the surrounding vapor. But above $T_{\rm lim}$, the nucleus is unstable and will fragment. This is the so-called Coulomb instability of hot nuclei. There are some experimental results showing the existence of the limiting temperature [8,9]. It has been shown that $T_{\rm lim}$ depends on the model used in the calculation sensitively [10–13]. The limiting temperature $T_{\rm lim}$ may, therefore, serve as the additional constraint for the effective models. It is of interest to adopt the FST model to study the Coulomb instability of hot nuclei and compare the results with those obtained from other theories and experiments. Since the problem considered involves asymmetric nuclear matter, we have taken the ρ meson degree of freedom into account to make the model isospin dependent. In this way, as pointed out in Refs. [11,12], we can avoid the negative asymmetry parameter in the vapor phase in the two-phase equilibrium model used to calculate the limiting temperature. The paper is organized as follows. In Sec. II, we will derive the EOS for bulk nuclear matter at finite temperature and describe a two phase model and the resulting coexistence equations. Section III contains the results and some discussions.

II. THE FST MODEL AND TWO PHASE EQUILIBRIUM MODEL

In addition to the usual self-interaction of the scalar field in the Walecka model, the FST model also introduces the self-interaction of the vector field, and the coupling between the scalar and vector fields. The Lagarangian density for the FST model is

$$\mathcal{L}(x)=\overline{\psi}\left(i\gamma^{\mu}D_{\mu}+g_{A}\gamma^{\mu}\gamma_{5}a_{\mu}-M+g_{s}\phi-\frac{1}{2}g_{\rho}\gamma_{\mu}\tau\cdot\mathbf{b}^{\mu}+\cdots\right)\psi-\frac{1}{4}F_{\mu\nu}F^{\mu\nu}+\frac{1}{2}m_{\rho}^{2}\mathbf{b}_{\mu}\cdot\mathbf{b}^{\mu}+\frac{1}{2}\left[1+\eta\frac{\phi}{S_0}+\cdots\right]\left[\frac{1}{2}f_{\pi}^{2}\,\mathrm{tr}(\partial_{\mu}U\partial^{\mu}U^{\dagger})+m_{v}^{2}V_{\mu}V^{\mu}\right]+\frac{1}{4!}\zeta(g_{v}^{2}V_{\mu}V^{\mu})^{2}+\frac{1}{2}\partial_{\mu}\phi\partial^{\mu}\phi-H_{q}\left(\frac{S^{2}}{S_{0}^{2}}\right)^{2/d}\left(\frac{1}{2d}\ln\frac{S^{2}}{S_{0}^{2}}-\frac{1}{4}\right)+\cdots, \tag{1}$$

where $g_A=1.23$ is the axial coupling constant, $\mathcal{D}_{\mu}=\partial_{\mu}+iv_{\mu}+ig_{v}V_{\mu}$ is a chirally covariant derivative, and U, v_{μ}, and a_{μ} depend on the pion field. Here, we have introduced the ρ meson field $\mathbf{b}^{\mu}$ to describe asymmetric nuclear matter. g_s, g_v, and g_{ρ} are, respectively, the couplings of the light scalar, vector ω, and ρ meson fields to the nucleon. $F_{\mu\nu}=\partial_{\mu}V_{\nu}-\partial_{\nu}V_{\mu}$ is the ω field strength tensor. $f_{\pi}=93$ MeV is the pion-decay constant, and η and ζ are real constants. The scale dimension d of the new scalar field $S(x)$ is allowed to differ from unity. The scalar fluctuation field ϕ is related to S by $S(x)\equiv S_0-\phi(x)$. H_q is linked to the mass of the light scalar S by the relation $m_s^2=4H_q/(d^2S_0^2)$. For unpolarized nuclear matter, there are no contributions from the π field. Meanwhile, by using the mean field approximation, the Lagarangian density then takes the form

$$\mathcal{L}_{\rm MFT}=\overline{\psi}\left[i\gamma^{\mu}\partial_{\mu}-(M-g_{s}\phi_{0})-g_{v}\gamma^{0}V_{0}-\frac{1}{2}g_{\rho}\tau_{3}\gamma^{0}b_{0}\right]\psi+\frac{1}{2}m_{v}^{2}V_{0}^{2}\left(1+\eta\frac{\phi_{0}}{S_{0}}\right)+\frac{1}{4!}\zeta(g_{v}V_{0})^{4}+\frac{1}{2}m_{\rho}^{2}b_{0}^{2}-H_{q}\left(1-\frac{\phi_{0}}{S_{0}}\right)^{4/d}\left[\frac{1}{d}\ln\left(1-\frac{\phi_{0}}{S_{0}}\right)-\frac{1}{4}\right]. \tag{2}$$

In the above equation, the meson field operators are replaced by their expection values

$$\phi\rightarrow\langle\phi\rangle\equiv\phi_{0}, \tag{3}$$

$$V_{\mu}\rightarrow\langle V_{\mu}\rangle\equiv\delta_{\mu0}V_{0}, \tag{4}$$

$$\mathbf{b}_{\mu}\rightarrow\langle\mathbf{b}_{\mu3}\rangle\equiv\delta_{\mu0}b_{0}, \tag{5}$$

where ϕ_0, V_0, and b_0 are, respectively, the mean field values of the σ, ω (the time component), and ρ (the time component in the third direction of isospin) mesons. For a system at finite temperature, $\langle O\rangle$ is replaced by a thermodynamic average $\ll O\gg$ $\ll O\gg$.

The Euler-Lagarangian Equations become

$$[i\gamma^{\mu}\partial_{\mu}-g_{v}\gamma^{0}V_{0}-(M-g_{s}\phi_{0})]\psi=0, \tag{6}$$

$$\frac{1}{6}\zeta g_{v}^{4}V_{0}^{3}+\left(1+\eta\frac{\phi_{0}}{S_{0}}\right)m_{v}^{2}V_{0}-g_{v}\rho=0, \tag{7}$$

$$\frac{\eta}{2S_{0}}m_{v}^{2}V_{0}^{2}+g_{s}\rho_{s}+m_{s}^{2}S_{0}\left(1-\frac{\phi_{0}}{S_{0}}\right)^{4/d-1}\ln\left(1-\frac{\phi_{0}}{S_{0}}\right)=0, \tag{8}$$

$$\frac{1}{2}g_{\rho}\rho_{3}-m_{\rho}^{2}b_{0}=0. \tag{9}$$

In the above equations, the scalar density ρ_s, nuclear density ρ and quantity ρ_3 are, respectively, expressed by

$$\rho_{s}\equiv\langle\overline{\psi}\psi\rangle, \tag{10}$$

$$\rho\equiv\langle\psi^{\dagger}\psi\rangle, \tag{11}$$

$$\rho_{3}\equiv\langle\psi^{\dagger}\tau_{3}\psi\rangle. \tag{12}$$

By using the standard technique in field theory, it is easy to obtain the Hamiltonian $\hat{H}$ of the system as

$$\hat{H}=\sum_{\mathbf{k}\lambda}E^{*}(\mathbf{k})(A_{\mathbf{k}\lambda}^{\dagger}A_{\mathbf{k}\lambda}+B_{\mathbf{k}\lambda}^{\dagger}B_{\mathbf{k}\lambda})+g_{v}V_{0}\hat{B}+\frac{1}{2}g_{\rho}b_{0}(\hat{B}_{p}-\hat{B}_{n})+V\Bigg(H_{q}\left\{\left(1-\frac{\phi_{0}}{S_{0}}\right)^{4/d}\times\left[\frac{1}{d}\ln\left(1-\frac{\phi_{0}}{S_{0}}\right)-\frac{1}{4}\right]+\frac{1}{4}\right\}-\frac{1}{2}\left(1+\eta\frac{\phi_{0}}{S_{0}}\right)m_{v}^{2}V_{0}^{2}-\frac{1}{4!}\zeta(g_{v}V_{0})^{4}-\frac{1}{2}m_{\rho}^{2}b_{0}^{2}\Bigg), \tag{13}$$

where the baryon number operator $\hat{B}_{\tau}$ is expressed as

$$\hat{B}_{\tau}=\sum_{\mathbf{k}\sigma}(A_{\mathbf{k}\sigma\tau}^{\dagger}A_{\mathbf{k}\sigma\tau}+B_{\mathbf{k}\sigma\tau}^{\dagger}B_{\mathbf{k}\sigma\tau})\quad(\tau=p,n) \tag{14}$$

and $\hat{B}=\hat{B}_p+\hat{B}_n$. In the above equations, $A_{\mathbf{k}\lambda}^{\dagger}$, $A_{\mathbf{k}\lambda}$ ($B_{\mathbf{k}\lambda}^{\dagger}$, $B_{\mathbf{k}\lambda}$) with $\lambda\equiv\{\sigma,\tau\}$ are creation and destruction operators for positive- (negative-) energy solutions of Eq. (6). V is the volume of the box within which the system is quantized and $E^*(k)=\sqrt{M^{*2}+k^2}$, with $M^*=M-g_s\phi_0$ being the effective nucleon mass.

To derive the equation of state at finite temperature, we calculate the thermodynamic potential Ω using the standard expression from statistical mechanics.

$$\Omega(\mu,V,T;\phi_{0},V_{0},b_{0})=-k_{B}T\ln Z_{G}, \tag{15}$$

where

$$Z_{G}\equiv T_{r}\exp[-(\hat{H}-\mu\hat{B})/k_{B}T], \tag{16}$$

where k_B is the Boltzmann's constant and μ the chemical potential. Substituting Eqs. (13) and (14) into Eq. (15), we obtain

$$\Omega=V\left\langle H_q\left\{\left(1-\frac{\phi_0}{S_0}\right)^{4/d}\left[\frac{1}{d}\ln\left(1-\frac{\phi_0}{S_0}\right)-\frac{1}{4}\right]+\frac{1}{4}\right\}\right. \tag{17}$$

$$-\frac{1}{2}m_\rho^2b_0^2-\frac{1}{2}\left(1+\eta\frac{\phi_0}{S_0}\right)m_v^2V_0^2$$

$$\left.-\frac{1}{4!}\zeta(g_vV_0)^4\right\rangle \tag{18}$$

$$-2k_BT\left\{\sum_{\mathbf{k}\tau}\ln[1+e^{-\beta(E^*(k)-\nu_\tau)}]\right.$$

$$\left.+\sum_{\mathbf{k}\tau}\ln[1+e^{-\beta(E^*(k)+\nu_\tau)}]\right\}, \tag{19}$$

where $\beta=1/k_BT$ and the quantity ν_τ is related to the usual chemical potential μ_τ by the equations

$$\nu_n=\mu_n-g_vV_0+\frac{g_\rho^2\rho_3}{4m_\rho^2} \tag{20}$$

and

$$\nu_p=\mu_p-g_vV_0-\frac{g_\rho^2\rho_3}{4m_\rho^2}, \tag{21}$$

where $\rho_3=\rho_p-\rho_n$.

It is easy to prove that the thermodynamic principle conditions, $\partial\Omega/\partial\phi_0=0$ and $\partial\Omega/\partial V_0=0$ are identical with the equations of motion, Eqs. (7) and (8), respectively.

Having obtained the thermodynamic potential, one can easily calculate all other thermodynamic quantities of the system. For example, the pressure p is given by the relation $p=-\Omega/V$, the average energy density ϵ by $\epsilon V=\partial(\beta\Omega)/\partial\beta+\mu\rho V$. The resulting expressions are as follows:

$$\mathcal{E}=H_q\left\{\left(1-\frac{\phi_0}{S_0}\right)^{4/d}\left[\frac{1}{d}\ln\left(1-\frac{\phi_0}{S_0}\right)-\frac{1}{4}\right]+\frac{1}{4}\right\}+g_vV_0\rho$$

$$-\frac{1}{2}\left(1+\eta\frac{\phi_0}{S_0}\right)m_v^2V_0^2-\frac{1}{4!}\zeta g_v^4V_0^4+\frac{g_\rho^2}{8m_\rho^2}\rho_3^2$$

$$+\frac{2}{(2\pi)^3}\int d^3kE^*(k)[n_n(k)+\overline{n}_n(k)+n_p(k)+\overline{n}_p(k)] \tag{22}$$

and

$$p=-H_q\left\{\left(1-\frac{\phi_0}{S_0}\right)^{4/d}\left[\frac{1}{d}\ln\left(1-\frac{\phi_0}{S_0}\right)-\frac{1}{4}\right]+\frac{1}{4}\right\}$$

$$+\frac{1}{2}\left(1+\eta\frac{\phi_0}{S_0}\right)m_v^2V_0^2+\frac{1}{4!}\zeta g_v^4V_0^4+\frac{g_\rho^2}{8m_\rho^2}\rho_3^2$$

$$+\frac{1}{3}\frac{2}{(2\pi)^3}\int d^3k\frac{k^2}{E^*(k)}[n_n(k)+\overline{n}_n(k)+n_p(k)$$

$$+\overline{n}_p(k)]. \tag{23}$$

In the above equations, $n_\tau(k)$ and $\overline{n}_\tau(k)$ are nucleon and antinucleon distributions and expressed as

$$n_\tau(k)=\{\exp[(E^*(k)-\nu_\tau)/k_BT]+1\}^{-1} \tag{24}$$

and

$$\overline{n}_\tau(k)=\{\exp[(E^*(k)+\nu_\tau)/k_BT]+1\}^{-1}\quad(\tau=n,p). \tag{25}$$

An asymmetric nuclear matter is usually defined by the total density ρ and an asymmetry parameter α. The neutron density ρ_n and proton density ρ_p are expressed as

$$\rho_n=(1+\alpha)\rho/2,\quad \rho_p=(1-\alpha)\rho/2. \tag{26}$$

The neutron density ρ_n and the proton density ρ_p determine the chemical potentials ν_τ by the subsidiary conditions

$$\rho_\tau=\frac{2}{(2\pi)^3}\int d^3k[n_\tau(k)-\overline{n}_\tau(k)]\quad(\tau=n,p). \tag{27}$$

After a simple calculation in Eq. (10), scalar density ρ_s is expressed as

$$\rho_s=\frac{2}{(2\pi)^3}\int d^3k\frac{M^*}{E^*(k)}[n_n(k)+\overline{n}_n(k)+n_p(k)+\overline{n}_p(k)]. \tag{28}$$

With the fixed α and ρ, Eqs. (7)–(9) and (24)–(28) can be solved consistently to calculate the chemical potential μ_τ, the effective nucleon mass M^*, and the mean field values ϕ_0, V_0, b_0, and then other quantities.

To study the change in the Coulomb instability of hot nuclei with the FST model, we will adopt the same model as used in Ref. [10] . We treat the hot nucleus as a uniformly charged drop of nuclear liquid with a sharp edge at a given temperature. Its equilibrium with the surrounding vapor, both thermal mechanical and chemical, leads to a set of two-phase coexistence equations

$$p(T,\rho_L,\alpha_L)+p_{\rm Coul}(\rho_L)+p_{\rm surf}(T,\rho_L)=p(T,\rho_V,\alpha_V), \tag{29}$$

$$\mu_n(T,\rho_L,\alpha_L)=\mu_n(T,\rho_V,\alpha_V), \tag{30}$$

$$\mu_p(T,\rho_L,\alpha_L)+\mu_{\rm Coul}(\rho_L)=\mu_p(T,\rho_V,\alpha_V), \tag{31}$$

where subscripts L and V stand for liquid and vapor, respectively. In the FST model for infinite nuclear matter, the Coulomb interaction is switched off and the surface effect is not

TABLE I. Parameter sets quoted from Ref. [4]. The vector masses are $m_v=783$ MeV and $m_\rho=770$ MeV, the nucleon mass is $M=939$ MeV. Values for S_0, the scalar mass m_s are in MeV.

Set	g_s^2	m_s	g_v^2	g_ρ^2	S_0	ζ	η	d
T1	99.3	509	154.5	70.2	90.6	0.0402	−0.496	2.70
T2	96.3	529	138.0	69.6	95.6	0.0342	−0.701	2.20
T1	109.5	508	178.6	67.2	89.8	0.0346	−0.160	3.50

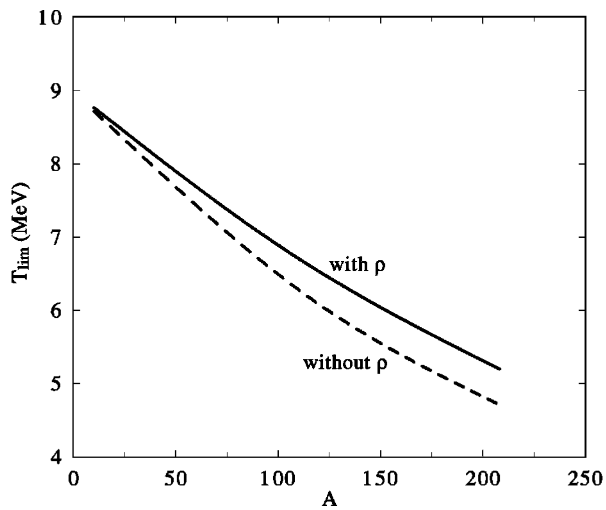

FIG. 1. The mass number dependence of limiting temperature $T_{\rm lim}$ calculated with the parameter set T1 in the FST model. The solid and dashed curves stand, respectively, for the results with and without the ρ meson degree of freedom.

considered. When the Coulomb interaction is added, the chemical potential of the proton has an additional term

$$\mu_{\rm Coul}(\rho)=\frac{6}{5}\frac{Ze^2}{R}, \tag{32}$$

where Z and R are the charge number and radius of the liquid droplet. Meanwhile, the pressure also has an extra term

$$p_{\rm Coul}(\rho)=\frac{Z^2e^2}{5AR}\rho, \tag{33}$$

where $A=N+Z$ is the number of nucleons in the liquid droplet.

For a liquid droplet with a surface, the additional pressure is given by

$$p_{\rm surf}(T,\rho)=-2\gamma(T)/R, \tag{34}$$

where R is determined by the nucleon density ρ through $A=\frac{4}{3}\pi R^3\rho$ for a given nucleon number A. $\gamma(T)$ is the surface tension. According to the suggestion of Goodman *et al.* [14], it can be written as

$$\gamma(T)=(1.14\ \ \text{MeV fm}^{-2})\left[1+\frac{3T}{2T_C}\right]\left[1-\frac{T}{T_C}\right]^{3/2}, \tag{35}$$

where T_C is the critical temperature for infinite symmetric nuclear matter.

III. RESULTS AND DISCUSSIONS

Since the parameter sets determined in Ref. [4] by fitting the properties of finite nuclei give also good descriptions for infinite nuclear matter, we will take these sets in our discussions. In Table I, we list the values of the parameter sets.

We would like to investigate the Coulomb instability of hot nuclei by calculating the limiting temperature $T_{\rm lim}$ above which the set of coexistence equations (29)–(31) have no solution. Here we consider the nuclei along the β-stability line

$$Z=0.5A-0.3\times10^{-2}A^{5/3}. \tag{36}$$

We show the mass number dependence of the limiting $T_{\rm lim}$ calculated with parameter set T1 in Fig. 1. The solid and dashed curves stand, respectively, for the results with and without the ρ meson degree of freedom. It is easily seen that the limiting temperature $T_{\rm lim}$ decreases with the increasing mass number of nuclei, and the rate of decrease is smaller for larger A. It is also seen that the limiting temperature calculated with the ρ meson degree of freedom is higher than that without the ρ meson degree of freedom. We present the solution of the coexistence equations and the equilibrium values of $\tilde{\mu}_n$, $\tilde{\mu}_p$, and p at the limiting temperature in Table II (without ρ meson) and Table III (with ρ meson). One can find that in the case of without the ρ meson (Table II), the neutron chemical potential $\tilde{\mu}_n$ is always lower than the proton chemical potential $\tilde{\mu}_p$, which results in a negative asymmetry parameter α_V of vapor. This feature is not reasonable in physics. When the ρ meson is taken into account in the model, the situation becomes very different. The asymmetry

TABLE II. Equilibrium values of densities ρ_L, ρ_V (in fm^{-3}), pressure p (in MeV fm^{-3}), chemical potentials $\tilde{\mu}_n$, $\tilde{\mu}_p$ (in MeV), and asymmetry parameter α_V for the nuclei along the β-stability line at limiting temperature $T_{\rm lim}$, calculated with parameter set T1 and without ρ meson in the FST model.

A	$T_{\rm lim}$	ρ_L	ρ_V	α_V	$\tilde{\mu}_n$	$\tilde{\mu}_p$	p
10	8.74	0.145	0.0136	−0.063	−15.90	−14.57	0.540
50	7.64	0.140	0.0116	−0.187	−14.77	−11.29	0.407
109	6.20	0.142	0.0095	−0.316	−12.69	−7.76	0.265
150	5.49	0.141	0.0077	−0.382	−11.61	−6.15	0.210
208	4.73	0.140	0.0064	−0.458	−10.39	−4.58	0.156

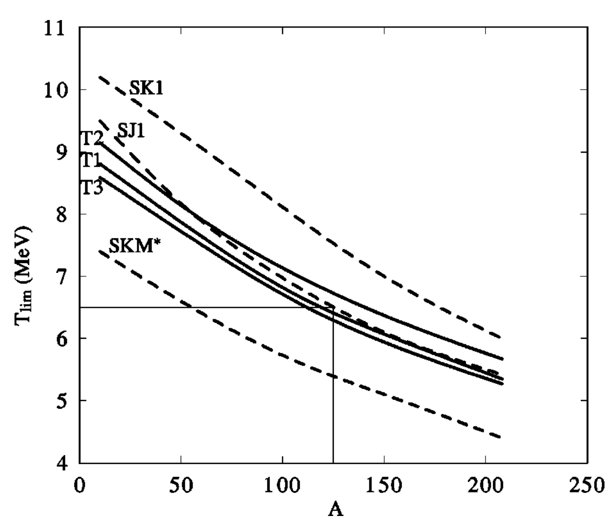

FIG. 2. The mass number dependence of limiting temperature $T_{\rm lim}$ calculated with parameter sets T1, T2, and T3 in the FST model, respectively, together with the results obtained from the nonrelativistic calculations with the Sk1, SJ1, and SkM* interactions.

TABLE III. Equilibrium values of densities ρ_L, ρ_V (in fm^{-3}), pressure p (in MeV fm^{-3}), chemical potentials $\tilde{\mu}_n$, $\tilde{\mu}_p$ (in MeV), and asymmetry parameter α_V for the nuclei along the β-stability line at limiting temperature $T_{\rm lim}$, calculated with parameter set T1 and with ρ meson in the FST model.

A	$T_{\rm lim}$	ρ_L	ρ_V	α_V	$\tilde{\mu}_n$	$\tilde{\mu}_p$	p
10	8.79	0.145	0.0136	0.0185	−15.07	−15.57	0.471
50	7.87	0.139	0.0126	0.0657	−12.54	−14.16	0.242
109	6.64	0.138	0.0101	0.167	−9.15	−12.61	0.161
150	6.03	0.136	0.0091	0.250	−7.34	−12.14	0.127
208	5.35	0.134	0.0080	0.382	−5.18	−11.86	0.121

parameter α_V in the vapor phase are all positive, in consistency with the results in either the HF calculation [15] or the nonrelativistic calculation of the limiting temperature [10].

In order to examine the difference between different parameter sets, we present the $T_{\rm lim}$-A curves in Fig. 2, calculated with parameter sets T1, T2, and T3. For a comparison, we have also shown with the dashed curves in the same figure the nonrelativistic results [10] with Sk1, SJ1, and SkM* interactions and shown with the fine lines the experimental result [9]. One can see that the difference of the results with different sets is not too large. In particular, the largest deviation from each other in $T_{\rm lim}$ is less than 1 MeV. These three curves are quite close to the SJ1 result and between the Sk1 (stiff one) and SkM* (soft one) results. The values of $T_{\rm lim}$ around $A=125$ calculated in present model are in good agreement with the experimental result.

In summary, we have extended the FST model with chiral symmetry and broken scale invariance into finite temperature and applied the extended model to study the Coulomb instability of hot nuclei. It is found that the FST model predicts reasonable limiting temperature $T_{\rm lim}$, which is quite close to experimental result. There is not large deference among the values of $T_{\rm lim}$ calculated with the different parameters. It seems that the FST model provides a good description for nuclear matter at finite temperature.

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