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EFFECTS OF SYMMETRY ENERGY ON DIRECT URCA PROCESS IN THE KAON CONDENSED PHASE

Hirotsugu Fujii, Takumi Muto, Toshihaka Tatsumi and Kyozo Tanagaki
Department of Physics, Kyoto University, Kyoto 606-01, Japan

Abstract

It is investigated whether or not the direct URCA process proceeds in neutron-star matter with kaon condensation, in connection with potential contribution of the nuclear symmetry energy. We find a region in density-symmetry-energy plane in a model independent way, where the direct URCA process occurs as a unique rapid cooling process in the kaon condensed phase. Symmetry energies with moderate density-dependence pass through this region.

A. INTRODUCTION

As one of the new hadronic phases in highly dense nuclear matter, kaon condensation has been extensively investigated from viewpoints of the quark-hadron physics [1-9] and neutron-star physics [2,8,10-13]. In the latter case, possible existence of kaon condensation is important because it leads to not only strong softening of the equation of state but also a drastic change of the chemical composition of neutron-star interior. Hence it is of interest to study cooling mechanism of neutron stars with kaon condensation, which reflects their interior structure.

In a series of our papers, we have studied onset mechanism and characteristic features of K^- -condensed phase and its implications to cooling of neutron stars on the basis of $SU(3)_L \times SU(3)_R$ current algebra and PCAC [6-8]. It has been shown that K^- -condensate appears in stable neutron-star matter through the chemical equilibrium for the weak interaction process,

$$n \rightleftharpoons p + K^-, \quad (1)$$

which results in a large proton-admixture. Because of no $K\pi$ -coupling through strong interaction, there are two independent Fermi spheres of quasi-proton ($\tilde{p}$) and quasi-neutron ($\tilde{n}$)

KUNS 1233
SW 34 10

states even in the kaon-condensed phase. (They are accompanied by the classical-kaon pair cloud non-perturbatively and coupled each other in the order of weak interaction.) Their Fermi momenta are not so different due to the large proton-admixture. Then the possibility of the so-called direct URCA process is naively expected as a unique cooling process associated with kaon condensation¹.

Following this argument, we have considered the possibility of the direct URCA process in the cold neutron-star matter with K^- -condensation,

$$\tilde{n} \rightarrow \tilde{p} + e^- + \bar{\nu}_e, \quad \tilde{p} + e^- \rightarrow \tilde{n} + \nu_e. \quad (2)$$

We have given an expression of the neutrino emissivity for the process (2) [8]:

$$\begin{aligned} E_{\text{URCA}} &= \frac{G_W^2}{16} \frac{457}{1260} \pi (1 + 3g_A^2) \cos^2 \theta_c \cos^2 \frac{\theta}{2} m_N^{-2} \mu T^6 \\ &= 8.9 \times 10^{26} \cos^2 \theta_c \cos^2 \frac{\theta}{2} \left(\frac{m_N}{m_\pi} \right)^2 \frac{\mu}{m_\pi} T^6 \quad (\text{erg} \cdot \text{cm}^{-3} \cdot \text{s}^{-1}) \end{aligned} \quad (3)$$

with the Cabibbo angle $\theta_c \sim 0.23$, chiral angle θ , charge chemical potential μ and temperature $T \equiv T/(10^9 \text{ K})$. Equation (3) holds only when the kinematical condition,

$$|p_F(\tilde{n}) - p_F(\tilde{p})| < p_F(e^-) < p_F(\tilde{n}) + p_F(\tilde{p}) \quad (4)$$

with the Fermi momenta $p_F(i)$ ($i = \tilde{p}, \tilde{n}, e^-$), is satisfied. Then the process (2) would give a larger neutrino emissivity than the other rapid cooling process under the K^- -condensation, $\tilde{n}(\tilde{p}) \rightarrow \tilde{n}(\tilde{p}) + e^- + \bar{\nu}_e$, $\tilde{n}(\tilde{p}) + e^- \rightarrow \tilde{n}(\tilde{p}) + \nu_e$ [2,11], roughly by one order of magnitude, stemming mainly from the ratio of the matrix elements between two processes, $\sim \cot^2 \theta_c$. The direct URCA process, if ever operative, would be the most efficient cooling mechanisms associated with kaon condensate.

In a simple treatment without nuclear interactions other than kaon-nucleon (KN) ones [8], however, it turns out that the kinematical condition (4) is hardly satisfied on the left inequality;²

¹In ordinary neutron-star matter, proton and electron fractions are strongly suppressed, and the direct URCA process is considered to be inhibited there because of the momentum mismatching among proton, neutron and electron near each Fermi surface.

the electron Fermi momentum decreases rapidly with the baryon number density n_B despite the large proton-admixture.

In determining the chemical composition of neutron-star matter more realistically, it is important to take into account the contribution of the isospin-dependent interactions such as nuclear symmetry energy. The symmetry energy makes matter more symmetric even in the normal phase: the larger symmetry energy gives the larger proton- and electron-admixtures in normal neutron-star matter. Hence it may modify the discussion on the kinematics of the direct URCA process mentioned above. Few works have been done on this issue so far [3, 13]. In this letter, we reexamine the kinematical condition (4) by taking into account the symmetry energy for nuclear interaction, and discuss the possibility of the direct URCA process in the K^- -condensed phase. Since the behavior of the symmetry energy at high densities has large theoretical uncertainty yet, we try to discuss the issue as independently of its details as possible.

B. FORMULATION

Our formulation is based on $SU(3)_L \times SU(3)_R$ current algebra and PCAC. This framework gives us an almost model-independent description of the kaon-nucleon dynamics as far as the s-wave condensation (kaon momentum $\mathbf{p}_K = 0$) is concerned [7, 8]. It enables us to describe in a unified way both equation of state controlled by strong interaction and cooling processes due to weak interaction within the underlying chiral symmetry [8].

The K^- -condensed state is defined by operating a unitary operator $\hat{U}_K$ on the meson vacuum $|0\rangle$,

$$|K^- \rangle \equiv \hat{U}_K(\Pi^1, \theta)|0\rangle = \hat{U}_{\Pi^1} \hat{U}_\theta |0\rangle \quad (5)$$

with $\hat{U}_{\Pi^1} = e^{i \int d^3x (\Pi_K^1 + K^+ + \Pi_K^1 - K^-)}, \quad \hat{U}_\theta = e^{i\theta \hat{Q}_K^3},$

where $\hat{Q}_a^5$ ($a=1\sim 8$) are the axial-vector charge operators which, together with the vector charge operators $\hat{Q}_a$, generate $SU(3)_L \times SU(3)_R$ algebra. Π_K^1 are the classical parts of the canonical conjugate momenta corresponding to the $K^\pm$ fields, and θ is the chiral angle. The classical kaon field is then given as

$$\langle K^- | K^\pm | K^- \rangle = \frac{f}{\sqrt{2}} \sin \theta, \quad \langle K^- | \Pi_{K^\pm} | K^- \rangle = \Pi_{K^\pm}^d \quad (6)$$

with the meson decay constant $f \sim 93$ MeV. The ground state energy is obtained as the expectation value of the effective Hamiltonian, $\mathcal{H}^{\text{eff}} = \mathcal{H} + \mu \hat{Q}$, where $\hat{Q}$ is the charge operator, and the charge chemical potential μ is introduced as a Lagrange multiplier to ensure the charge neutrality. The expectation value of $\mathcal{H}^{\text{eff}}$ for the state $|K^- \rangle$ can be evaluated as that of the modified operator, $\tilde{\mathcal{H}}^{\text{eff}}(\Pi^1, \theta, \mu) \equiv \hat{U}_K^{-1}(\Pi^1, \theta) \mathcal{H}^{\text{eff}} \hat{U}_K(\Pi^1, \theta)$, for the meson vacuum $|x\rangle$ with the baryon number density n_B and the proton mixing ratio $x \equiv n_p/n_B$.

In the original Hamiltonian $\mathcal{H}$ we assume the octet dominance of the chiral-symmetry breaking term, $\mathcal{H}_{\text{SB}} = h_0 \mathbf{1} + h_3 T_3 + h_8 T_8$, and no-derivative coupling in the Lagrangian. Thus the expectation value of $\tilde{\mathcal{H}}^{\text{eff}}$ for the state $|x\rangle$ is determined due to the symmetry property of $\mathcal{H}$, and can be written down in the following form after minimization with respect to $\Pi_{K^\pm}^d$:

$$\mathcal{E}(\theta, \mu) = \langle x | \mathcal{H}^{\text{eff}}(\mu) | x \rangle + (\cos \theta - 1) \langle x | (\hat{\Sigma}_K + \mu \hat{\rho}_V) | x \rangle - \frac{1}{2} \mu^2 f^2 \sin^2 \theta, \quad (7)$$

where $\hat{\Sigma}_K \equiv [\hat{Q}_4^5, [\hat{Q}_4^5, \mathcal{H}_{\text{SB}}]]$ is called the kaon sigma operator which reflects the explicit chiral symmetry breaking of the Hamiltonian and $\hat{\rho}_V \equiv \frac{1}{2}(V_3^{\mu=0} + \sqrt{3}V_8^{\mu=0})$ is the V-spin charge density.

Hereafter we take a definite model for the normal neutron-star matter: nucleons are regarded as the non-relativistic Fermi gas in a mean-field potential. Both relativistic free electrons and muons are taken into account in beta equilibrium. In the K^- -condensed phase, the chemical equilibrium condition for (1) should be imposed as well as the beta equilibrium one. Then we get $\mu_{K^-} = \mu_e = \mu_\mu = \mu$, where μ_i ($i=K^-, e^-, \mu^-$) are chemical potentials. Thus one gets the energy-density expression for the K^- -condensed phase,

$$\begin{aligned} \mathcal{E} = & \frac{3}{5} E_{\text{F,G}}^0 [x^{5/3} + (1-x)^{5/3}] n_B + [V_0 + V_{\text{sym}}(1-2x)^2] n_B \\ & + (\cos \theta - 1) [n_B \Sigma_{KN} + \mu n_B (x + \frac{1}{2}(1-x)) - f^2 m_K^2] \\ & - \frac{1}{2} \mu^2 f^2 \sin^2 \theta + \sum_{i=e,\mu} \Theta(|\mu| - m_i) \frac{m_i^4}{8\pi^2} \left\{ \tau_i (2\tau_i^2 + 1)(\tau_i^2 + 1)^{1/2} - \ln(\tau_i^2 + 1)^{1/2} \right\} \\ & + \mu \left[x n_B - \sum_{i=e,\mu} \Theta(|\mu| - m_i) \frac{p_{\text{F}}^3(i)}{3\pi^2} \right], \end{aligned} \quad (8)$$

where $E_{FG}^0 \equiv (3\pi^2 n_B)^{2/3}/2m_N$ and $\tau_i \equiv p_F(i)/m_i$ for $i=e^-, \mu^-$. Practically we have discarded electron mass. Both two terms in the third bracket give the s -wave K^+N attractive interactions which drive kaon condensation. For the K^+N sigma term $\Sigma_{KN} \equiv \langle N | \hat{\Sigma}_K | N \rangle$, we take the allowable value of 300~500 MeV [8]. The second term is the nuclear potential contribution, and quadratic approximation is introduced with respect to the asymmetry parameter $(1-2x)$, so that V_{sym} represents the potential contribution to the symmetry energy. This approximation has been shown to work very well even over the densities where asymmetry becomes large in normal nuclear matter [14,15]. We treat here $V_0(n_B)$ and $V_{sym}(n_B)$ unaltered in the K^- -condensed phase from ones in the normal phase; the effect of $V_0(n_B)$ does not show up explicitly in our discussion as far as we deal with the energy difference, $\Delta E^K/n_B$, between two phases.

Concerning the symmetry energy $E_{sym}(n_B)$, it consists of two kinds of contributions:

$$E_{sym}(n_B) = E(n_B; x=0) - E(n_B; x=1/2) = T_{sym}(n_B) + \bar{V}_{sym}(n_B), \quad (9)$$

where $E(n_B; x)$ is energy per baryon of asymmetric nuclear matter with proton mixing ratio x , and $T_{sym}(n_B) \equiv (1 - 2^{-2/3}) \frac{3}{8} E_{FG}^0$ is a contribution to the symmetry energy from the kinetic energy of nucleons. The empirical nuclear mass formula yields $E_{sym}(n_B=n_0)=27-36$ MeV at the saturation density, $n_0=0.17 \text{ fm}^{-3}$ [16]. There is, however, a large controversy about the density-dependence of $V_{sym}(n_B)$, depending on theoretical models (see Fig. 4).

C. MODEL-INDEPENDENT CONSIDERATION

For the moment we don't need any definite behavior of $V_{sym}(n_B)$ with density n_B . The ground state is determined under the following conditions: chemical equilibrium with respect to the process (1), $\partial \mathcal{E}/\partial x = 0$, charge neutrality, $\partial \mathcal{E}/\partial \mu = 0$, and field equation for θ , $\partial \mathcal{E}/\partial \theta = 0$. Solving these equations at various densities, we obtain a critical density for kaon condensation and an equation of state $\mathcal{E}(n_B)$ numerically [8]. We also obtain the proton-mixing ratio and charge chemical potential which are crucial to the kinematical condition (4). Then one can discuss the possibility of the direct URCA process in the kaon condensed phase.

Solving these equations consistently at every step of changing V_{sym} for a fixed density, we can get the critical value of the symmetry energy, $E_{sym, \text{crit}}^K = T_{sym} + V_{sym, \text{crit}}^K$, for each density, above

FIG. 1: The region of symmetry energy vs baryon number density for the direct URCA process. (a) Bold solid line denotes the critical value of E_{sym} , $E_{sym, \text{crit}}^K$, above which the direct URCA process proceeds in the K^- -condensed phase in case of $\Sigma_{KN}=400$ MeV. Long-dashed line is the corresponding one in the normal phase, $E_{sym, \text{crit}}^n$. At densities right to the vertical line, ng^f , the charge chemical potential is negative. Empirical value at the saturation density is also depicted for reference (thick segment). (b) The same as in (a) for $\Sigma_{KN}=300, 400, 500$ MeV. Arrow denotes ng^f for each Σ_{KN} .

which the direct URCA process proceeds under the condition (4). We depict the value $E_{sym, \text{crit}}^K$ as a function of density in Fig. 1(a) with the bold solid line for $\Sigma_{KN}=400$ MeV. For comparison, the corresponding value $E_{sym, \text{crit}}^n$ above which the direct URCA process proceeds even in the normal phase, is shown by the long-dashed line. The critical density for kaon condensation, n_B^* , at which the condition $\mu_n = \mu_p = \mu_{K^-}$ is achieved among chemical potentials in the limit $\theta \rightarrow 0$, is also shown. The behavior of the critical curves, $E_{sym, \text{crit}}^K$ and $E_{sym, \text{crit}}^n$, can be qualitatively understood as follows. Using a linear approximation around $x = 1/2$, $(1-x)^{1/3} \sim x^{1/3} \propto 1-2x$, one can find out from Eq. (4) and the chemical equilibrium condition, $\partial \mathcal{E}/\partial x = 0$, that in the normal phase $E_{sym, \text{crit}}^n$ is nearly proportional to $n_B^{1/3}$ [18], that is, monotonically increases with density. One can also see in the same approximation that once the kaon condensate appears in

the ground state, this behavior of $E_{\text{sym},\text{crit}}^K$ is modified to be

$$E_{\text{sym},\text{crit}}^K \sim 45 \frac{1 + \cos \theta}{2} (n_B/n_0)^{1/3} \text{ MeV}. \quad (10)$$

As seen in Ref. [8], $\cos \theta$ rapidly decreases with density from 1 to 0 (weakly depending on the strength of the symmetry energy itself). Therefore $E_{\text{sym},\text{crit}}^K$ decreases with density by the factor $(1 + \cos \theta)/2$. Roughly speaking, kaon condensation pulls the long dashed line down to the bold line. In other words, kaon condensation is surely favorable for the direct URCA process, and there exists the region between the curves, $E_{\text{sym},\text{crit}}^K$ and $E_{\text{sym},\text{crit}}^n$, where the direct URCA process is uniquely possible in the kaon condensed phase. In Fig. 1(b) we also show the results for $\Sigma_{KN}=300, 400, 500$ MeV. Since Σ_{KN} represents the s wave KN attractive interaction, the kaon condensate appears at the lower density and $\cos \theta$ decreases in the lower density region for the larger Σ_{KN} . Then the critical curve, $E_{\text{sym},\text{crit}}^K$, is lowered and the region where the direct URCA process occurs uniquely is enlarged for the larger Σ_{KN} .

As we have noticed in Ref. [8], $\mu=0$ and $x=1/2$ at the density $n_B^{\text{com}} = f^2 m_K^2 / \Sigma_{KN} (=4.0 n_0$ for $\Sigma_{KN}=400$ MeV) and μ becomes negative at densities above n_B^{com} (right-side of the vertical line in Fig. 1(a)). Because of no contribution of the symmetry energy at $x=1/2$, this point n_B^{com} is common for any V_{sym} , so that this behavior of μ is universal. Above n_B^{com} , μ turns to be negative and positively charged leptons, e^+ (μ^+), may appear in the system. Then the direct URCA process with positron, $\bar{p} \rightarrow \bar{n} + e^+ + \nu_e$, $\bar{n} + e^+ \rightarrow \bar{p} + \bar{\nu}_e$, may be possible in place of (2) at these high densities. In fact, the kinematical condition for this process, $|p_F(\bar{p}) - p_F(\bar{n})| < p_F(e^+) = -\mu < p_F(\bar{p}) + p_F(\bar{n})$, is satisfied in higher density region even in the case $V_{\text{sym}}=0$ (See Fig. 1). Hence we take the hatched triangle in Fig. 1(a) as a unique region for the direct URCA process in the K^- -condensed phase, distinguishing from the region where $\mu < 0$. Although the direct URCA process with e^+ (or μ^+) is certainly an interesting possibility, the results obtained in such a high density region should not be considered seriously yet, because there is still much room for elaborate study of the K^- -condensed phase far above the critical density: the energy gain, $|\Delta \mathcal{E}^K/n_B|$, is too large to be consistent with observations

of neutron stars². Hence some repulsive effects which compensate the large attraction would be desirable.

At threshold for kaon condensation n_B^* , the charge chemical potential μ exceeds the muon mass in most cases and muons exist tens of percent of the number of electrons. However their contributions to the energy gain and other physical quantities are negligible as well as electrons, though the critical density n_B^* is slightly pushed up to higher value by the inclusion of leptons. This is because lepton energies are very small as compared with those of hadrons (Eq. (8)). As K^- -condensate develops with density, muons disappear and electrons become small component of the matter because of the steep decrease of μ .

D. EXAMPLES

In the remaining discussion, we deal with some definite cases with respect to density-dependence of E_{sym} and examine the kinematical condition (4) in more detail. According to the non-relativistic potential-model calculations by the use of the $\bar{G}$ matrix and the variational methods, E_{sym} increases moderately with density. In calculations including effects of three-body forces by Wiringa, Fiks and Fabrocini [19], E_{sym} at high density has been shown to be sensitive to three-body forces; E_{sym} at $n_B \gtrsim 3n_0$ drastically decreases toward high density for TNI-V model [20], but its decrease starts at $n_B \sim 6n_0$ for UVII model [21]. On the other hand, E_{sym} obtained in relativistic Brueckner-Hartree-Fock calculations [22] continues to increase rapidly with density and has much larger values than those in the non-relativistic models. Generally E_{sym} at high density becomes lower when the isoscalar short-range interaction works more repulsively as n_B increases. We take here four typical cases in order of the magnitude of E_{sym} in our disposal: (A) as an extreme, the one obtained from the UV14 potential with a three-nucleon interaction (UV14-TNI) [19], (B) the one from the effective interaction of the asymmetric nuclear matter which is constructed by Nishizaki *et al.* [14] on

²For the values $\Sigma_{KN} \sim 500$ MeV, the equation of state with kaon condensation is too soft to construct a neutron star with an observed mass $1.44 M_\odot$ [17] in the treatment in Ref. [13].

FIG. 2. Proton mixing ratio x as a function of a baryon number density n_B for the typical cases of the symmetry energy E_{sym} : (A), (B), (C), and (D) for $\Sigma_{KN}=400$ MeV (see the text). For comparison, the case $V_{\text{sym}}(n_B) \equiv 0$ is also shown (case (O)). The density-dependence of x for the normal phase ($\theta \equiv 0$) is also shown by dashed line.

the basis of the G-matrix method with the Reid soft-core potential, (C) the one obtained by Takatsuka [23] with the G-matrix method by the use of a potential with the one-pion-exchange tail and Gaussian core (OPEG) [24], and (D) as another extreme, the one parametrized as $E_{\text{sym}}(n_B) = T_{\text{sym}}(n_0)n^{2/3} + [S_0 - T_{\text{sym}}(n_0)]F(u)$ with $u \equiv n_B/n_0$ given by Prakash et al. [15] in reference to the relativistic approach. We put here $F(u) \equiv u$ and $S_0 \approx 32$ MeV as an example of a large symmetry energy.

In Fig. 2, the density dependence of the proton-mixing ratio x is shown for the case $\Sigma_{KN} = 400$ MeV. For comparison, the value for the normal phase ($\theta \equiv 0$) corresponding to each case (A)-(D) is also shown by dashed lines. We also depict the case $V_{\text{sym}} \equiv 0$, which corresponds to the result in Ref. [8] (hereafter denoted as (O)) for reference. We can see that the value of x is considerably enhanced in the K^- -condensed phase compared with the normal phase. We note that all the lines cross the common point at the density, n_B^{gm} , where $x=1/2$ and $\mu=0$. Below the density n_B^{gm} , the larger symmetry energy leads to the larger proton mixing but

FIG. 3. Difference between proton and neutron Fermi momenta, $p_F(\bar{n}) - p_F(\bar{p})$, and the electron one $p_F(e^-) = \mu$ as functions of density for a typical case (C) of the symmetry energy E_{sym} as well as the case with $V_{\text{sym}}(n_B) \equiv 0$ (case (O)) for $\Sigma_{KN}=400$ MeV. In the case (C) bold segment on the density axis corresponds to the hatched region in Fig. 1. Those for the normal phase ($\theta \equiv 0$) are shown by dashed lines.

above it the increase of x becomes more moderate for the larger symmetry energy, because the symmetry energy term in (8) always works to make matter symmetric. ((A) is an exceptional case, where the proton-mixing continues to increase rapidly because V_{sym} for the UV14-TNI potential becomes negative around $4 n_0$ [19]. Matter favors asymmetry at high densities in this case.)

On the other hand, the relation among the Fermi momenta is given by using the resulting equations of state with K^- condensate for the four cases. The inclusion of the symmetry energy term, however, does not modify the qualitative features of the relation from the case (O): (i) the difference $p_F(\bar{n}) - p_F(\bar{p})$ decreases with density as a result of the large proton-mixing x ; (ii) $p_F(e^-) (= \mu)$ has an order of pion mass at threshold and it decreases rapidly with density. The feature (i) works favorably for the kinematical condition (4), while (ii) does against it. We give here the results for cases (C) and (O) in Fig. 3, for example. As seen in Fig. 3, the difference

FIG. 4. Region of the symmetry energy E_{sym} vs baryon number density n_B where the kinematical condition for the direct URCA process is satisfied. The density-dependence of $E_{\text{sym}}(n_B)$ for the typical cases (A), (B), (C) and (D) are shown as well as the case with $V_{\text{sym}}(n_B)=0$ (case (O)). The cross point of each symmetry energy curve with dash-dotted line means the critical density of K^- condensation, n_B^c . Bold, long-dashed and dash-dotted lines are the same as Fig. 1.

$p_\pi(\bar{n}) - p_F(\bar{p})$ becomes further smaller than that of the case (O) near threshold ($\mu > 0$) owing to the resulting larger proton-mixing (see Fig. 2) by the inclusion of the symmetry energy term. On the other hand, $p_F(e^-)$ does not change from the case (O) very much at each density. As a result, the condition is somewhat relaxed by the symmetry energy term. Nevertheless, the kinematical condition (4) is not always satisfied near threshold yet.

In Fig. 4, we overwrite $E_{\text{sym}}(n_B)$ for the typical cases (A)~(D) on Fig. 1(a), all of which pass through the empirical value 27–36 MeV at the saturation density. The cross point of each symmetry-energy curve with dash-dotted line indicates the critical density of the K^- condensation, n_B^c . One can see in Fig. 4 that the larger is the symmetry energy at the relevant densities, the more easily can the direct URCA process proceed uniquely in the K^- -condensed phase at the lower densities. In the case (A) based on the non-relativistic potential, UV14-TNI, which gives almost the same equation of state given by Friedman and Paudharipande

for nuclear and neutron matter [25], however, the direct URCA process is inhibited for both the normal and the K^- -condensed phases over the relevant densities. It should also be noted that in the cases such as (B) and (C), the direct URCA process becomes possible only in the well-developed K^- -condensed phase rather than near threshold. Almost the same results can be expected for a non-relativistic potential model with three-body forces such as AV14-UVII [19], which is an improved version of UV14-TNI. With rather extreme density-dependence of the symmetry energy which simulates the relativistic mean-field calculations (case (D)), the direct URCA process is allowed near threshold and even in the normal phase before the onset of the kaon condensation [18]. Thus, the difference of the density-dependence of the symmetry energy between the relativistic models and the standard non-relativistic many-body theories brings about the crucial difference for the possibility of the direct URCA process near threshold.

Once the direct URCA process (2) occurs at the onset density, the neutrino emissivity (3) discontinuously jumps to a finite value due to the factor $\cos^2(\theta/2)$. This is in contrast to the other rapid cooling process with the kaon condensate whose emissivity is proportional to $\sin^2\theta$ [2,8,11]. After taking the largest value at the onset density, the neutrino emissivity decreases to zero with density controlled by the falling behavior of the charge chemical potential μ in the kaon condensed phase. Above n_B^{com} , the emissivity turns up to increase due to the process with positron instead of electron.

Throughout this work, we have not discussed modification of the symmetry energy $V_{\text{sym}}(n_B)$ as well as $V_0(n_B)$ in the K^- -condensed phase. However, the isospin-dependent interaction such as the symmetry energy might be modified because the ground state with classical K^- condensate violates isospin symmetry. It is interesting to clarify whether the θ -dependence of $V_{\text{sym}}(n_B)$ is substantial or not, though various possible behaviors of $V_{\text{sym}}(n_B)$ are summarized in Fig. 4.

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