Bondi accretion of dark matter by neutron stars

Xi Huang,1,3,∗ Jian-Feng Liu,2 Wei-Hua Wang,2 Quan Cheng,1 and Xiao-Ping Zheng2,†
1School of Electronic and Electrical Engineering, Wuhan Textile University, Wuhan 430073, China
2Institute of Astrophysics, Central China Normal University, Wuhan 430079, China
3Key Laboratory of Quark and Lepton Physics (Ministry of Education), Central China Normal University, Wuhan 430079, China

In this paper, we have compared two different accretion mechanisms of dark matter particles by a canonical neutron star with \( M = 1.4 \, M_\odot \) and \( R = 10 \, \text{km} \), and shown the effects of dark matter heating on the surface temperature of star. We should take into account the Bondi accretion of dark matter by neutron stars rather than the accretion mechanism of Kouvaris (2008) [29], once the dark matter density is higher than \( \sim 3.81 \, \text{GeV/cm}^3 \). Based on the Bondi accretion mechanism and heating of dark matter annihilation, the surface temperature platform of star can appear at \( \sim 10^5 \, \text{K} \) year and arrive \( \sim 1.12 \times 10^8 \, \text{K} \) for the dark matter density of \( 3.81 \, \text{GeV/cm}^3 \), which is one order of magnitude higher than the case of Kouvaris (2008) with dark matter density of \( 30 \, \text{GeV/cm}^3 \).

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I. INTRODUCTION

According to Planck measurements of the cosmic microwave background (CMB) temperature and lensing-potential power spectra [11], we have known that the total energy of the Universe contains about 27% dark matter (DM) and about 5% baryon matter as well as about 68% dark energy. It is generally believed that DM particles only participate in gravitational interaction, but not in electromagnetic and strong interaction. Therefore, it is very difficult to detect DM particles. In recent years, there are mainly three kinds of methods to detect DM particles, namely, the direct detections, the indirect detections, and the collider detections [2,3]. Direct search experiments aim to detect the optical, thermal, and ionized signals due to the scattering between DM particles and target nuclei of detectors [4], such as CRESST-II [5], Super-CDMS [6], DarkSide [7], LUX [8], XENON100 [9], PandaX-II [10] experiments and so on. Indirect detection experiments mainly probe the signals that DM particles can decay or annihilate into gamma-rays, neutrinos or charged anti-particles such as positrons and antiprotons [11], including AMS-02, PAMELA, Fermi-LAT, IACTs, IceCube, ANTARES, Super-K experiments [12,13], etc. Finally, the collider detection is to produce DM particles via the high energy collider, such as the Large Hadron Collider (LHC) at CERN [14,15]. From a theoretical point of view, DM is stable in the cosmic age scale, lightless, and massive particle beyond the Standard Model of particle physics. To date, several candidates of DM particles are supposed [16], for example, weakly interacting massive particle (WIMP) such as Kluza-Klein particle [17] predicted by the theories of the additional spatial dimension and neutralino [18] proposed in the supersymmetric theories, axion [19,20], axino [21], gravitino [22], millicharged (MC) particle [23], self-interacting DM (SIDM) [24] and so on. At first, in order to solve the neutrino problem in the Sun and the missing mass problem in the Galaxy, Press and Spergel (1985) [25] studied that WIMPs are captured by the Sun via the captured mechanism of scattering off individual nucleons in the Sun and into bound orbits. Subsequently, Gould (1987,1988) [26,27] computed the direct and indirect capture of WIMPs by the Earth. It is interesting that Goldman and Nussinov (1989) [28] researched that WIMPs are trapped in neutron stars and concentrated towards the star center. This results in the formation of a mini black hole that consumes the neutron star (NS). Recently, basing on the work of Press and Spergel (1985), Kouvaris (2008) [29] discussed in detail that WIMPs are accreted and trapped eventually by the NS, taking into account the general relativity corrections of the star. Meanwhile, the WIMP’s accretion rate onto the NS was shown, and the effect of the WIMP annihilation on the cooling curves of a canonical NS made of ordinary nuclear matter was calculated. The main conclusions of this paper [29] are that the released energy due to WIMP annihilation inside the NS might affect the temperature of the star older than \( 10^7 \) years, appearing the plateau of surface temperature at \( \sim 10^4 \, \text{K} \) for a typical NS. By the way, besides the above-mentioned heating effects owing to the annihilation of trapped DM particles onto the NS [29,34], the rotation effect on the NS was also studied in our previous work [35], based on an enhanced slow-down of neutron stars due to an extra current yield by the accretions of millicharged dark matter particles (MCDM) [36]. We constrained the charge-mass phase space of MC particles through the simulation of the rotational evolution of neutron stars [35].

∗huangxi@mails.ccnu.edu.cn
†zhxp@phy.ccnu.edu.cn
In this paper, we revisit the accretion mechanism of DM particles by neutron stars basing on the work of Kouvaris (2008) \cite{29}. We should consider the Bondi accretion \cite{37–39} of the SIDM rather than the accretion mechanism of Kouvaris (2008), if the candidate of DM is in the form of the SIDM \cite{24, 40–43}, and the DM density is about one order of magnitude higher than the standard DM density 0.3 GeV/cm$^3$ around the Earth. As discussed in a large amount of literatures, the SIDM can affect the mass profile and shape of DM halos, reducing the central densities of dwarfs and low surface brightness galaxies in agreement with observations \cite{41}. It is very likely that the SIDM is made up of composite states of baryons or pions in the dark sector. The most attractive characteristic of SIDM is large scattering cross section with other DM particles. The commonly accepted view is that the cross section per unit mass of SIDM must be of order

$$\sigma_{X}/m_{X} \sim 1 \text{ cm}^2/\text{g} \approx 2 \times 10^{-24} \text{ cm}^2/\text{GeV} \quad (1)$$

or larger, where $X$ is the DM particle. Obviously, this value is many orders of magnitude larger than that of the WIMP. The typical cross section of WIMP is $\sigma_{X} \sim 10^{-36} \text{ cm}^2/\text{GeV}$, if the mass is $m_{X} \sim 100$ GeV, thus the cross section per unit mass is $\sigma_{X}/m_{X} \sim 10^{-38} \text{ cm}^2/\text{GeV}$. For the SIDM, we can easily deduce the DM density $\rho = \frac{1}{2\sigma_{X}/m_{X}} \approx 3.81 \text{ GeV/cm}^3$, if the mean free path takes the diameter of Milky Way galaxy $\lambda \sim 30$ kpc. We consider there is no doubt that the Bondi accretion of DM by neutron stars is dominant as long as the DM density is higher than the above value.

This paper is organized as follows. First, the thermal evolution equations of a canonical NS ($M = 1.4 \, M_{\odot}, R = 10$ km) and the accretion rate due to the Bondi accretion of DM particles, are shown in Sec. II. Then we numerically calculate the effect of SIDM annihilation on the surface temperature of the NS compared to the work of Kouvaris (2008) in Sec. III. Finally, we present the conclusions in Sec. IV.

II. THE MODELS

A. The thermal evolution equations of the NS

In general, for a newborn NS, the internal temperature can reach up to $10^{10}$ K, it cools via neutrino emission during the first million years (neutrino cooling era). Incidentally, the contribution to neutrino emission would mainly come from the interior of the NS, without considering it from the crust. The emission of photons from the surface of the star dominates the cooling (photo cooling era), as soon as the internal temperature drops below $10^{8}$ K. Besides the NS cools via the internal neutrino emission and surface photon emission, on the other hand, various heating mechanisms can be present during the late times of the thermal evolution \cite{44, 45}. These include the frictional interaction between the faster rotating superfluid core and the slower rotating outer solid crust \cite{46, 47}, crust cracking \cite{48}, magnetic field decay \cite{49, 52}, rotochemical heating \cite{53, 55}, deconfinement heating \cite{56, 57}, and dark matter heating \cite{29, 32, 34}.

Direct Urca (DU) processes and modified Urca (MU) processes are dominant for the neutrino emission. DU processes $n \rightarrow p + e + \bar{\nu}_{e}, p + e \rightarrow n + \nu_{e}$ are allowed in sufficiently dense nuclear matter. nuclear matter with pion condensation, kaon condensation, or nonzero hyperon density, and in all phases of quark matter except the Color Flavor Locked phase \cite{58}.

The temperature of the star drops very fast via the DU processes, in which the emissivity scales as $\epsilon_{\nu} \sim T^9$ ($T$ is the interior temperature of the star). However, DU processes are kinematically forbidden if the star is not in sufficiently high density. In this case, MU processes $b + n \rightarrow b + p + e + \bar{\nu}_{e}, b + p + e \rightarrow b + n + e_{\nu}$ are switched on in NS cores, which $b$ denotes the bystander (neutron or proton). The neutrino emissivity of MU processes can be written as \cite{59}

$$\epsilon_{\nu} = 1.2 \times 10^4 \left( \frac{m_{n}}{m_{0}} \right)^{2/3} \left( \frac{T}{10^8 \text{ K}} \right)^8 \text{ ergs cm}^{-3} \text{ s}^{-1}, \quad (2)$$

where $m_{n}$ is the average baryon number density, $m_{0} = 0.17 \text{ fm}^{-3}$ is the nuclear saturation density.

After the first million years, the dominant mechanism of cooling is the photon emission from the surface of the star. The surface photon luminosity of the star is $L_{\gamma} = 4\pi R^2 \sigma_{T_s}$, where $\sigma$ is the Stefan-Boltzmann constant and $T_s$ is the surface temperature of the star. The relationship between the interior temperature $T$ and the surface temperature $T_s$ is well approximated by \cite{60}

$$T_s = (0.87 \times 10^6 \text{ K}) \left( \frac{g_s}{10^{14} \text{ cm/s}^2} \right)^{1/4} \left( \frac{T}{10^8 \text{ K}} \right)^{0.55}, \quad (3)$$

where $g_s = GM/R^2$ is the surface gravity of the star. Thus, the surface photon luminosity $L_{\gamma}$ can be expressed in terms of the interior temperature as

$$L_{\gamma} = 4\pi R^2 \sigma (0.87 \times 10^6 \text{ K})^4 \left( \frac{g_s}{10^{14} \text{ cm/s}^2} \right) \left( \frac{T}{10^8 \text{ K}} \right)^{2.2}. \quad (4)$$

The emissivity from the surface photon emission for the NS of $M = 1.4 \, M_{\odot}$ and $R = 10$ km is given by

$$\epsilon_{\gamma} = \frac{L_{\gamma}}{(4/3)\pi R^3} = 1.8 \times 10^{14} \left( \frac{T}{10^8 \text{ K}} \right)^{2.2} \text{ ergs cm}^{-3} \text{ s}^{-1}. \quad (5)$$

It is well known that the NS is a compact object, thus DM around the star can be accreted. As a result of scattering with nuclei, DM can lose enough energy and get deposited in the core of the star fast. It releases a huge amount of energy to heat up the star due to the annihilation between DM particles. Because of the competition between the accretion and annihilation of DM, the population of DM inside the star is governed by

$$\frac{dN}{dt} = \mathcal{F} - \Gamma_{\text{annih}}, \quad (6)$$

where $\mathcal{F}$ (number of particles per time) represents the accretion rate of DM captured by neutron stars and $\Gamma_{\text{annih}} = \langle \sigma_{\text{annih}} v \rangle \int n_{\text{d}}^2 dV = C_{\Lambda} N(t)^2$ represents the annihilation rate.
In the above equation, \((\sigma_{\text{annih}} v)\) is the thermally averaged annihilation cross section times the DM velocity, \(n_X\) is the number density of DM inside the star, and \(C_A = (\sigma_{\text{annih}} v) / V\) is the constant where \(V\) is the volume of the star.

Obviously, the solution of Eq. (6) is

\[
N(t) = \frac{\mathcal{F}}{C_A} \tanh(t/\tau),
\]

where \(\tau = 1/\sqrt{C_A}\) is the time scale, depending on the DM density and temperature of the star [29]. Therefore, the luminosity due to the annihilation of DM particles can be expressed as

\[
L_X = \Gamma_{\text{annih}} n_X = C_A N(t)^2 n_X = \mathcal{F} \tanh^2(t/\tau) n_X.
\]

As it can be seen from above Eq. (8), the annihilation rate saturates to the accretion rate when \(t\) is larger than roughly \(3 \tau\). Hence, the emissivity of DM can be written as

\[
\epsilon_X = \frac{L_X}{4\pi R^2 \sqrt{3}},
\]

where \(R\) is the radius of the star [59].

In order to obtain the curves of thermal evolution for a NS, we also need to know the specific heat of the star. For convenience, the core of NS is made of non-superfluid neutrons, protons, and few electrons in our model. Meanwhile, we suppose that the density of the star is uniform. The specific heat of the star mainly contributed by the fermions from the core is given by [59]

\[
c_v = \frac{k_B T}{3 h c} \sum_i p_i^v \left[ m_e^2 c^2 + (p_i^v)^2 \right],
\]

where \(k_B\), \(h\), and \(c\) are the Boltzmann constant, reduced Planck constant, and speed of light, respectively, the sum runs over the different particles \(n\), \(p\), \(e\) and their corresponding Fermi momenta are

\[
p_i^v = (340 \text{ MeV}) \left( \frac{n_i}{n_0} \right)^{1/3},
\]

\[
p_i^p = p_i^e = (60 \text{ MeV}) \left( \frac{n_i}{n_0} \right)^{2/3}.
\]

However, we should emphasize that the contribution to the specific heat from DM particles can be neglected, because they consist of a tiny fraction of the whole mass of the star.

By far, we can gain the thermal evolution equation of the NS

\[
\frac{dT}{dt} = -\epsilon_s - \epsilon_r + \epsilon_X / c_v.
\]

**B. The accretion rate of Bondi accretion**

In this section we calculate the Bondi accretion rate of DM particles along the lines of [61]. As we know the continuity equation of magneto-fluid is

\[
\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{u}) = 0,
\]

where \(\rho\) and \(\mathbf{u}\) are the density and velocity of fluid, respectively. In the polar coordinate of spherical symmetry, we have

\[
\nabla \cdot (\rho \mathbf{u}) = \frac{1}{r^2} \frac{\partial (r^2 \rho \mathbf{u})}{\partial r}.
\]

Thus Eq. (14) can be written as

\[
\frac{\partial \rho}{\partial t} + \frac{1}{r^2} \frac{\partial (r^2 \rho \mathbf{u})}{\partial r} = 0.
\]

We can derive \(\frac{\partial \mathbf{u}}{\partial t} = 0\), \(\frac{\partial \rho}{\partial t} = 0\), and \(r^2 \rho \mathbf{u} = \text{const}\) easily, if the fluid flows smoothly. In general, this const is defined as \(\frac{M}{a^2}\), where \(M\) is the accretion rate (accreted mass per time). In addition, the momentum equation of magneto-fluid is

\[
\frac{\partial \mathbf{u}}{\partial t} + (\mathbf{u} \cdot \nabla) \mathbf{u} = -\frac{1}{\rho} \nabla p + \nabla \psi,
\]

where \((\mathbf{u} \cdot \nabla) \mathbf{u} = u_i \frac{\partial \mathbf{u}}{\partial x_i}\) in the polar coordinate of spherical symmetry, \(p\) is the pressure of fluid, and \(\nabla \psi = -\frac{GM}{r^2}\). Therefore, Eq. (17) can be simplified as

\[
\frac{\partial u_r}{\partial t} + u_r \frac{\partial u_r}{\partial r} = -\frac{1}{\rho} \frac{\partial p}{\partial r} - \frac{GM}{r^2}.
\]

For conveniences, the equation of state of fluid is supposed in the form of polytropic process, namely, \(p \sim \rho^\gamma\). Thus we can define the acoustic velocity \(a_s^2 = \frac{dp}{\rho} = \gamma \bar{p}\). For Eq. (18), we can deduce

\[
\frac{1}{\rho} \frac{\partial \rho}{\partial r} = -\frac{1}{\rho} \frac{\partial}{\partial r} (u_r r^2).
\]

According to continuity equation (16), we can give \(\frac{\partial}{\partial r} (r^2 \rho \mathbf{u}) = 0\), namely,

\[
\frac{1}{\rho} \frac{\partial \rho}{\partial r} = -\frac{1}{\rho} \frac{\partial}{\partial r} (u_r r^2).
\]

Combining Eqs. (19) and (20), Eq. (18) can be written as

\[
\frac{1}{2} \left( 1 - \frac{a_s^2}{u_r^2} \right) \frac{\partial (u_r^2)}{\partial r} = -\frac{GM}{r^2} \left( 1 - \frac{2a_s^2}{GM} \right).
\]

As the fluid is close to the central compact object, namely, \(r\) decreases continually, \(1 - \frac{2a_s^2}{GM}\) will increase gradually. When this value reaches zero, we can obtain

\[
r_s = \frac{GM}{2a_s^2(r_s)}.
\]

where \(r_s\) is called the critical radius. If \(r = r_s\), we can see the left hand side of Eq. (21) must be zero, requiring

\[
u_r^2 = a_s^2
\]

or

\[
\frac{d}{dr} (u_r^2) = 0.
\]
On the basis of laws of thermodynamics, there exists $dh = \frac{dp}{p}$ for the isentropic process, where $h$ is the enthalpy per unit mass of the fluid. For the polytropic process, the equation of state is $p = p_0 \rho^\gamma$ where $p_0$ is constant. Hence, we can derive

$$\frac{1}{\rho} \frac{dp}{d\rho} = \gamma \rho_0 \rho^{-\gamma - 2}. \quad (25)$$

Defining $h = h_0 \rho^\beta$ where $h_0$ is constant, thus we have

$$\frac{dh}{d\rho} = \beta h_0 \rho^{\beta - 1}. \quad (26)$$

Defining $\beta = \gamma - 1$ and combining Eqs. (25), (26), we can deduce

$$h = \left( \frac{\gamma}{\gamma - 1} \right) \frac{p}{\rho} = \frac{a_s^2}{\gamma - 1}. \quad (27)$$

In addition, since the fluid flows steadily, Eq. (18) can be written as

$$\nu_r \frac{\partial}{\partial r} \nu_r = - \frac{1}{\rho} \frac{\partial p}{\partial r} + \beta \frac{\partial h}{\partial r} = - \frac{\partial h}{\partial r} + \frac{\partial \psi}{\partial r}, \quad (28)$$

namely,

$$\frac{\nu_r^2}{2} + h - \psi = B,$$ \quad (29)

where $B$ is constant.

In the following section, we will compute the integration constant $B$ using the boundary conditions. When $r \to r_{\infty}$, the radial velocity of fluid and gravitational potential also approach to zero $\nu_r \to 0$, $\psi \to 0$. In terms of Eq. (29), we can get $h_{\infty} = B$. Basing on Eq. (27), thus we have

$$h_{\infty} = \frac{a_{s\infty}^2}{\gamma - 1} = B. \quad (30)$$

Combining Eqs. (22), (25), and (27), Eq. (29) can be written as

$$\frac{a_s^2(r_r)}{2} + \frac{a_s^2(r_r)}{\gamma - 1} - 2a_s^2(r_r) = B,$$ \quad (31)

at the critical radius $r = r_r$. In the above equation, the integration constant $B$ is given by

$$\frac{a_s^2(r_r)(5 - 3\gamma)}{2(\gamma - 1)} = B. \quad (32)$$

Comparing Eq. (30) to Eq. (32), we can easily derive

$$a_s^2(r_r) = \frac{2a_{s\infty}^2}{5 - 3\gamma}. \quad (33)$$

Hence, the critical radius can be expressed as

$$r_s = \frac{GM}{2a_s^2(r_r)} = \frac{(5 - 3\gamma)GM}{4a_{s\infty}^2}. \quad (34)$$

According to the definition of acoustic velocity $a_s^2 \sim \rho^{\gamma - 1}$, we have

$$\frac{a_s(r)}{a_{s\infty}} = \left( \frac{\rho(r)}{\rho(\infty)} \right)^{\gamma - 1}. \quad (35)$$

In terms of Eq. (33), Eq. (35) can further be written as

$$\left[ \frac{a_s(r_r)}{a_{s\infty}} \right]^2 = \left( \frac{2}{5 - 3\gamma} \right)^{\gamma - 1} = \frac{\rho(r_r)}{\rho(\infty)}^{\gamma - 1} \quad (36)$$

at the critical radius $r = r_r$. Therefore, the density of fluid at the critical radius can be expressed as

$$\rho(r_r) = \rho(\infty) \left( \frac{2}{5 - 3\gamma} \right)^{\gamma - 1}. \quad (37)$$

The accretion rate at the critical radius can represent those with different radii, because it is constant due to flowing steadily. Finally, combining Eqs. (23), (33), (34), and (37), the accretion rate (accreted mass per time) is given by

$$M = 4\pi \rho(r_r) \nu_r(r_r) r_r^2 \quad (38)$$

where $\Gamma = 4\pi \times 2^{\gamma - 1} (5 - 3\gamma)^{\gamma - 1}$. For the adiabatic process, we know $\gamma = \frac{5}{3}$, thus $\Gamma = \frac{\pi}{2}$. According to the definition of acoustic velocity, we can further calculate the acoustic velocity at infinity $a_{s\infty}^2 = \frac{\gamma}{2} \rho(\infty)^{\gamma - 1}$. Hence, Eq. (38) can further be written as

$$M = \frac{\pi}{24} \rho(r_r) G^2 M^2 \nu_{s\infty}^2. \quad (39)$$

### III. NUMERICAL RESULTS AND DISCUSSIONS

For convenience, we assume that the NS is made up of non-superfluid neutrons, protons and electrons, and suppose that the density of the star is uniform. In our specific calculations, we consider a canonical NS of $M = 1.4 M_\odot$ and $R = 10$ km, and impose the initial interior temperature $T_0 = 10^{10}$ K at very early time for the star. However, it is emphasized that the thermal evolution of the star is very insensitive to the initial condition, which affects the temperature only during the first years. Meanwhile, we safely take the average velocity at infinity $v_{s\infty} = 220$ km/s of DM particles. It is worth stressing that the DM density at infinity $\rho(\infty) = 3.81$ GeV/cm$^3$ is adopted. It is the critical DM density whether considering the Bondi accretion of DM by neutron stars or not.

The evolution of the surface temperature for a typical NS with $M = 1.4 M_\odot$ and $R = 10$ km as a function of time is shown in Fig. 1. For the “standard cooling” scenario, the cooling of a NS will undergo the neutrino emission and...
 photon emission. After $\sim 10^6$ years, the surface temperature of the star drops fast during the photo cooling era. Figure 1 shows the cases where the effect of DM annihilation is considered for three different DM densities, based on the accretion mechanism of DM particles discussed in the work of Kouvaris (2008) \[29\]. It can be seen from the figure that DM heating can affect the surface temperature of the star significantly after $\sim 10^7$ year as the temperature of the star drops. The plateau of surface temperature will appear, once the equilibrium between the heating of DM annihilation and the cooling of surface photon emission has been reached. Furthermore, the temperature platform depends on the DM density, the mass, and the radius of the star. Obviously, the temperature platform is higher and appears earlier for the case with higher DM density. However, it is not optimistic to detect the surface temperature as low as $\sim 10^4$ K for the case with the DM density of $30$ GeV/cm$^3$, which would possibly be a signature of DM annihilation. Encouragingly, the surface temperature platform can reach up to $\sim 1.12 \times 10^5$ K, if Bondi accretion will play the leading role in the accretions of DM particles by neutron stars. The temperature platform is about one order of magnitude higher and appears earlier ($\sim 10^{6.5}$ year) than that of Kouvaris (2008), because the Bondi accretion rate of DM particles is well above that of Kouvaris (2008). As it can be deduced from Eq. (39), the Bondi accretion rate of DM particles can reach up to $1.31 \times 10^8$ g/s, which is about six orders of magnitude higher than that of Kouvaris (2008) (see Eq. (18) in Ref. \[29\]). It is possible to explain the high temperature behaviors of old neutron stars.

IV. CONCLUSIONS

We have shown the effects of DM heating on the surface temperature of a canonical NS ($M = 1.4$ $M_\odot$, $R = 10$ km) for two different accretion mechanisms of DM particles. DM heating will play a leading role at the later time of NS, as the temperature of the star drops. Once the heating of DM annihilation equilibrates the cooling of surface photon emission, as a result the temperature of the star will remain flat as a function of time. The temperature plateau depends largely on the DM density, the mass, and the radius of NS. The Bondi accretion of DM by neutron stars is dominant rather than the accretion mechanism of Kouvaris (2008) \[29\], if the DM density is higher than $\sim 3.81$ GeV/cm$^3$. It is obvious that the surface temperature platform appears at $\sim 10^{6.5}$ year, which is earlier than that of Kouvaris (2008) with three different DM densities. In particular, the surface temperature platform can arrive $\sim 1.12 \times 10^5$ K for the DM density of $3.81$ GeV/cm$^3$, which is about one order of magnitude higher than the case of Kouvaris (2008) with DM density of $30$ GeV/cm$^3$. More importantly, it maybe explain the high temperature behaviors of old neutron stars.

The model we adopted is based on uniform stellar configuration. However, it is commonly considered that the neutron stars can be approximated as the uniform case, because the results will be unchanged in the order of magnitude when considering the realistic equations of states. In addition, the effect of stellar rotation is disregard, which can lead to rotochemical heating due to the derivative from beta equilibrium with the spin-down of the star. This mechanism could possibly be present for old neutron stars.

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