A Powerful $e^\pm$ Outflow Driven by a Proto-strange Quark Star

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Abstract

An electron– positron layer can cover the surface of a bare strange star (SS), the electric field in which can excite the vacuum and drive a pair wind by taking away the heat of the star. In order to investigate the pair-emission ability of a proto-SS, we establish a toy model to describe its early thermal evolution, where the initial trapping of neutrinos is specially taken into account. It is found that the early cooling of the SS is dominated by the neutrino diffusion rather than the conventional Urca processes, which leads to the appearance of an initial temperature plateau. During this plateau phase, the surface $e^\pm$ pair emission can maintain a constant luminosity of $10^{48} – 10^{49}$ erg s$^{-1}$ for about a few to a few tens of seconds, which is dependent on the value of the initial temperature. The total energy released through this $e^\pm$ wind can reach as high as $\sim 10^{51}$ erg. In principle, this pair wind could be responsible for the prompt emission or extended emission of some gamma-ray bursts.

Unified Astronomy Thesaurus concepts: Neutron stars (1108); Gamma-ray bursts (629); Stellar winds (1636)

1. Introduction

It was hypothesized that strange quark matter (SQM), consisting of roughly equal numbers of up, down, and strange quarks, is the ground state of the strong interaction (Witten 1984). If this hypothesis is true, then the nature of the compact stars known as pulsars should in fact be strange quark stars (SSs; Alcock et al. 1986; Haensel et al. 1986) rather than conventional neutron stars (NSs), at least in part. Nevertheless, the existence of NSs can still be allowed as a metastable state. It is undoubtedly of fundamental physical significance to test the SS hypothesis by searching for possible candidates from observations. According to the differences between SSs and NSs in their mass–radius relations, cooling histories, and rotational properties, several observational signatures have been proposed for identifying SS candidates (e.g., Pizzochero 1991; Page 1992; Schaab et al. 1996; Madsen 2000; Glendenning & Weber 2001; Xu et al. 2002; Ray et al. 2004; Bagchi et al. 2006; Yu & Zheng 2006; Zheng et al. 2006). Unfortunately, current observations of Galactic pulsars usually found them falling into an ambiguous parameter space that covers both predictions of the SS and NS models (Lorimer & Kramer 2012; Lyne & Graham-Smith 2012).

During the past two decades, it has been increasingly suggested that a rapidly rotating and highly magnetized pulsar can be formed from some extragalactic luminous transient phenomena such as gamma-ray bursts (GRBs; Dai & Lu 1998a, 1998b; Dai et al. 2006; Yu et al. 2010) and superluminous supernovae (SLSNe; Kasen & Bildsten 2010; Woosley 2010; Dexter & Kasen 2013). The substantial influence of the newborn pulsars on the transient emission could provide a potentially effective way to probe into the properties and nature of these compact stars (see Yu et al. 2019 for a brief review). For example, SSs are expected to be able to keep near-Keplerian rotation much longer than NSs due to the high bulk viscosity of SQM. Then, Yu et al. (2009) suggested that SS candidates could be picked up from the afterglow observations of GRBs by searching for ultra-long internal plateau afterglows, since this plateau emission is very likely to be powered by the stellar rotation. Following this consideration, Dai et al. (2016) further suggested that the central engine of the most luminous supernova ASASSN-15lh could just be an extremely rotating SS.

More excitingly, mergers of double pulsars, which can result in a short GRB, could give birth to a new massive pulsar, which is very helpful for understanding the afterglow features of short GRBs (Dai et al. 2006; Rowlinson et al. 2010, 2013; Bucciantini 2012; Gompertz et al. 2013, 2015; Zhang 2013; Lü et al. 2015; Gao et al. 2016). Some new clues to such a post-merger pulsar could have also appeared in the multi-messenger observations of GW events (e.g., GW170817), such as kilonova (Li et al. 2018; Ai et al. 2018; Radice et al. 2018; Yu et al. 2018; Margutti et al. 2019; Piro et al. 2019; Ren et al. 2020; Troja et al. 2020). Then, in comparison to the pulsars originating from core collapse of massive stars, we are more likely to expect that the nature of post-merger pulsars are inclined to be SSs. Even if the pre-merger pulsars are NSs, a phase transition from metastable NSs before merging to an SS after merging could still happen, in view of the high mass $\gtrsim 2.0 M_\odot$ of the merger product (e.g., Alford et al. 2019; Parisi et al. 2020). Such an expectation can indeed be supported by the fittings of the afterglows of some short GRBs (Li et al. 2016; Hou et al. 2018). In principle, the most direct signature of a post-merger SS can be dug from the waveform of the GW signal (Oechslin et al. 2004; Bauswein et al. 2010), which is however far beyond the ability of current GW detectors (Bauswein et al. 2019; Most et al. 2020). Therefore, at present, it is still necessary to investigate the other astrophysical manifestations of newborn SSs, which actually has long been a concern of the community (Dai et al. 1995, 2016; Cheng & Dai 1998; Ouyed et al. 2002, 2009, 2015; Leahy & Ouyed 2008; Cheng et al. 2009a, 2009b; Cheng & Harko 2010; Pagliara et al. 2013).

Due to the extremely high temperature at birth, a newborn pulsar cannot have a solid crust, which enables SSs to show a special feature. To be specific, unlike NSs, the surface of bare SSs should be covered by an electron layer of a thickness of a
where \( \zeta = 2(\alpha/\pi)^{1/2}k_{B}T_{e}/(k_{B}T_{e}) \), \( \alpha \) is the fine-structure constant, \( \varepsilon_{e,1} \) is the Fermi energy of leptons, and

\[
J(\zeta) = \frac{1}{3} \frac{\zeta^{3} \ln(1 + 2\zeta^{-1})}{(1 + 0.074\zeta)^{3}} + \frac{\pi^{5}}{6} \frac{\zeta^{4}}{(13.9 + \zeta)^{4}}.
\]

(4)

Hereafter the conventional notation \( Q_{\nu} = Q/10^{5} \) is used in the ergs units. Meanwhile, one may write the luminosity of the surface photon emission by using the Stephan–Boltzmann formula. However, as pointed out by Usov (2001), the blackbody emission can actually be significantly suppressed for \( T_{s} > 8 \times 10^{8} \text{K} \), since the diffusion of the photons of an electromagnetic frequency lower than the plasma frequency of the quark matter can be blocked. Therefore, in our calculations we can simply neglect \( L_{\nu} \).

For the neutrinos, if they can escape freely as considered for a normal SS, then the luminosity can usually be calculated by timing the stellar volume to the emissivity \( \epsilon_{\nu} \) (Iwamoto 1980)

\[
L_{\nu,\text{free}} = \frac{4}{3} \pi R^{3} \epsilon_{\nu}
\]

\[
\approx 3.7 \times 10^{57} R_{6}^{3} \alpha_{c} \left( \frac{n_{b}}{n_{0}} \right) Y_{e}^{1/3} T_{6}^{9} \text{erg s}^{-1},
\]

(5)

which is obtained from the direct Urca process \( d \rightarrow u e^{-}\bar{\nu}_{e} \) and \( u e^{-} \rightarrow d \nu_{\bar{e}} \), where \( \alpha_{c} = 0.1 \) is the strong coupling coefficient, \( n_{b} \) is the baryon number density, \( n_{0} = 0.16 \text{fm}^{-3} \) is the nuclear saturation density, \( Y_{e} \) is the electron fraction in the interior of the SS, and \( T \) is the inner temperature. Nevertheless, it should be pointed out that Equation (5) is actually inapplicable for a proto-SS, which has an extremely high temperature. In this case, the mean free path of neutrinos due to the absorption and scattering by the quark matter can be much shorter than the stellar radius, which reads (Iwamoto 1982; Berdermann et al. 2004)

\[
I_{\nu} \approx 1.4 \times 10^{3} \left( \frac{n_{b}}{n_{0}} \right)^{1/3} Y_{e}^{-1/3} \left( \frac{\mu_{e}}{\mu_{u}} \right)^{-2} \times \left[ 1 + \frac{1}{2} \left( \frac{\mu_{e}}{\mu_{u}} \right) + \frac{1}{10} \left( \frac{\mu_{e}}{\mu_{u}} \right)^{2} \right]^{-1}
\]

\[ \times \left[ \frac{\epsilon - \mu_{e}}{10 \text{MeV}} \right]^{2} + \left( \frac{\pi kT}{10 \text{MeV}} \right)^{2} \right]^{-1} \text{cm}. \]

(6)

For a rough estimation, the neutrino energy can be adopted as \( (\epsilon - \mu_{e}) \sim kT \) and the chemical potentials satisfy \( \mu_{e} \sim \mu_{u} < \mu_{\nu} \). This indicates the neutrinos can be seriously trapped in the proto-SS until the stellar temperature decreases to be lower than \( 10^{9} \text{K} \).

Therefore, for the initial cooling of the proto-SS, the neutrino luminosity should be alternatively recalculated by considering the neutrino diffusion in the stellar interior. For an approximate calculation of the neutrino emission, it could be convenient to separate the proto-SS into two parts by according to a neutrino sphere at the radius of

\[
R_{\nu} = R - I_{\nu},
\]

(7)

as illustrated in Figure 1. Furthermore, the temperatures in the inner and outer parts are both considered to be uniform, and referred to the effective internal and surface temperature,
respectively. The actual temperature gradient in the stellar interior is simplified by a temperature jump occurring at the neutrino sphere. Then, the neutrino luminosity at $R_v$ is determined by the neutrino diffusion from the stellar center to the neutrino sphere and, for an order-of-magnitude estimation, the luminosity can be given by Kasen & Bildsten (2010); Yu et al. (2015)

$$L_{v,\text{in}} = \frac{4\pi R_v^2 c}{3\kappa_v \rho} \left( \frac{\partial u_v}{\partial r} \right) \sim \frac{4\pi R_v^2 c (u_{v,\text{in}} - u_{v,\text{out}})}{3 \max(\tau_v, 1)},$$

where $u_v$ is the energy density of thermal neutrinos, $\kappa_v$ is the opacity, $\rho$ is the mass density, and $\tau_v = \kappa_v \rho R_v \approx R_v / \nu_R$ is the neutrino optical depth of the inner part. Meanwhile, the neutrino emission of the outer part of the star is contributed by two components, including the escaping of the thermal neutrinos and the neutrino production via the Urca processes. The luminosity of the outer part can be written as

$$L_{v,\text{out}} = \pi R^2 \nu_c u_{v,\text{out}} + \frac{4}{3} \pi (R^3 - R_v^3) \epsilon_v.$$  

where $u_{v,\text{in}}$ is the energy density of thermal neutrinos, $\kappa_v$ is the opacity, $\rho$ is the mass density, and $\tau_v = \kappa_v \rho R_v \approx R_v / \nu_R$ is the neutrino optical depth of the inner part and, for an order-of-magnitude estimation, the luminosity can be given by Kasen & Bildsten (2010); Yu et al. (2015)

$$L_{v,\text{in}} = \frac{4\pi R_v^2 c}{3\kappa_v \rho} \left( \frac{\partial u_v}{\partial r} \right) \sim \frac{4\pi R_v^2 c (u_{v,\text{in}} - u_{v,\text{out}})}{3 \max(\tau_v, 1)},$$

Here the appearance of the thermal emission component is due to the absorption of the neutrinos from the inner part by the outer part.

Following the above consideration, we further denote the internal energy of the inner and outer part of the proto-SS by

$$E_{\text{in}} = \frac{4}{3} \pi R_v^3 u_{v,\text{in}},$$

and

$$E_{\text{out}} = \frac{4}{3} \pi (R^3 - R_v^3) u_{v,\text{out}},$$

respectively, where the energy density are given by summarizing all particles as $u = \sum_{e=q,d,s} u_e$. Then, Equation (1) can be replaced by the following two equations:

$$\frac{dE_{\text{in}}}{dt} = -L_{\nu,\text{in}} - 4\pi R_v^2 \nu_c u_{\text{in}}$$

and

$$\frac{dE_{\text{out}}}{dt} = -L_{\nu,\text{in}} + 4\pi R_v^2 \nu_c u_{\text{in}} - L_{\nu,\text{out}} - L_{\Delta L},$$

where $\zeta = 1 - e^{-1} = 0.63$ represents the absorption by the neutrino sphere, the term $4\pi R_v^2 \nu_c u_{\text{in}}$ represents the energy transfer from the inner to the outer part due to the inward movement of the neutrino sphere, and $v_c = -dR_c/dt = dl_c/dt$ is the moving velocity of the neutrino sphere. Finally, the total luminosity of the neutrino emission is given by

$$L_{\nu} = (1 - \zeta) L_{\nu,\text{in}} + L_{\nu,\text{out}}.$$  

In our model, we do not consider the possible change of the stellar radius due to the change of the equation of state during the phase transition, which could only lead to a $\sim 1\%$ difference (Dexheimer et al. 2013).

2.2. Expressions of Internal Energy

The temperature dependence of the internal energy of the SQM is determined by the $\beta$ equilibrium in presence of neutrinos: $u + e^- \leftrightarrow d + \nu$ and $u + e^- \leftrightarrow s + \nu$. In this neutrino trapping case, the chemical potentials of the particles are suggested to approximately satisfy (Iwamoto 1982)

$$\mu_u \simeq \mu_d \simeq \mu_s$$

and

$$\mu_e \simeq \mu_\nu.$$  

Then, for relativistic quarks, their internal energy can be calculated as (Iwamoto 1982; Shapiro & Teukolsky 1983)

$$u_q = \sum_{i=u,d,s} \frac{\pi^2}{2} \frac{n_i}{\rho} \bar{e}_{F,i} \left( \frac{k_T}{\bar{e}_{F,i}} \right)^2$$

$$\approx 1.7 \times 10^{33} \left( \frac{n_b}{n_0} \right)^{2/3} \left( \frac{k_T}{10 \text{MeV}} \right)^2 \text{erg cm}^{-3},$$

where the Fermi energy of the quarks reads

$$\bar{e}_{F,q} = \left( \frac{n_b h^3 c}{8\pi} \right)^{1/3} \approx 235 \left( \frac{n_b}{n_0} \right)^{1/3} \text{MeV}$$

for $n_u = n_d = n_s = n_b$. Here, because the baryon number is conserved, the degeneracy energy of the quarks cannot be released during the stellar cooling. Thus, only the thermal component is necessary and exhibited in Equation (17). Different from quarks, neutrinos can eventually escape from the proto-SS. As a result, the number density and chemical potential of electrons and neutrinos can decrease with time. Therefore, when we describe the internal energy of the leptons, both the thermal and degeneracy components should be taken into account as

$$u_l = \sum_{i=e,\nu} \left[ \frac{3}{4} n_i \bar{e}_{F,i} + \frac{\pi^2}{2} n_i \bar{e}_{F,i} \left( \frac{k_T}{\bar{e}_{F,i}} \right)^2 \right]$$

$$\simeq \frac{4\pi}{c^3 h^3} \mu_1 \left[ 1 + \frac{2\pi^2}{3} \left( \frac{k_T}{\mu_1} \right)^2 \right],$$

where $\mu_1 = \bar{e}_{F,1}$ is used. The above expression is valid for $kT \ll \bar{e}_{F,1}$, where the Fermi energy of the leptons can be...
expressed by introducing an electron fraction $Y_e$ as

$$
\varepsilon_{F,1} = (3Y_e)^{1/3} \varepsilon_{F,q} 
\approx 200 \left( \frac{Y_e}{0.2} \right)^{1/3} \left( \frac{n_b}{n_0} \right)^{1/3} \text{MeV.} \quad (20)
$$

When most of the neutrinos escape, the chemical equilibrium in the outer part finally becomes

$$
\mu_u + \mu_e \simeq \mu_d \simeq \mu_n.
$$

Here the neutrino chemical potential $\mu_\nu$ is believed to be zero, although it is not easy to explicitly describe the decreasing behavior of $\mu_\nu$, which is non-equilibrium. Then, for an effective description of the energy density of neutrinos in the outer part, it is assumed that the neutrino chemical potential can simply co-evolve with the temperature as

$$
\mu_\nu \propto kT e^{-l_\nu/R},
$$

which is basically in agreement with the considerations in previous works (e.g., Pons et al. 1999, 2001; Benvenuto & Horvath 2013; Pagliara et al. 2013). Here the term $e^{-l_\nu/R}$ is introduced by considering that the neutrino chemical potential should quickly approach zero after the SS becomes totally transparent for $l_\nu > R$. In addition, a constant value of $10^{-3}$ will be assigned to the electron faction in the outer part, after the electrons are finally decoupled with the neutrinos.

In Figure 2, we present a comparison between the temperature dependence of $u_q$ and $u_\nu$. For a relatively high temperature, the internal energy is primarily stored in the neutrinos and thus the cooling of the proto-SS should be dominated by the release of the thermal neutrinos. On the contrary, for a relatively low temperature, the neutrino trapping is relaxed and the quark energy becomes dominate. In this case, the Urca processes are the primary channel of the stellar cooling.

3. Results and Discussion

The internal energy of a proto-SS is initially converted from the gravitational energy released during the formation of the SS and from the phase transition from nucleon to quarks. Then the initial temperature can be estimated to be on the order of magnitude of $kT_i \approx 10$ MeV by according to

$$
\frac{4}{3} \pi R^3 [u_q(T_i) + u_1(T_i)] \approx \left( \frac{3GM^2}{10R} \right) + \frac{M}{m_p} q,
$$

where $q \approx 20$ MeV is the latent heat per baryon. For such an extremely high temperature, the neutrino free path is much smaller than the stellar radius and thus we take $R_e \approx R$ in the above estimation. Then, for a reference value of $kT_i = 10$ MeV, we solve Equations (12) and (13) and obtain the cooling curves of the SS as presented in Figure 3, where both the internal and surface temperatures are shown. Obviously, the cooling curves can be separated into a plateau phase and a decline phase. Just as expected, the plateau is determined by the neutrino diffusion, while the slow decay is gradually dominated by the Urca processes.

In Figure 4, we further present the evolution of the luminosity of the neutrino and pair emissions, for four different initial temperatures. On the one hand, the obtained neutrino luminosity is basically consistent with the previous works (e.g., Pons et al. 2001; Pagliara et al. 2013), which can be estimated by

$$
L_{\nu,1} \approx \frac{4}{3} \pi R^3 \frac{u_1}{t_{d,\nu}} \approx 10^{55} \text{erg s}^{-1},
$$

with a neutrino diffusion timescale of

$$
t_{d,\nu} = \left( \frac{R}{l_\nu} \right)^2 \left( \frac{l_\nu}{c} \right) \propto l_\nu^{-1}.
$$

This diffusion timescale determines the duration of the plateau. According to Equation (6), it can be approximately obtained that $t_{d,\nu} \propto T_i^2$. So, as shown, the higher the initial temperature, the longer the plateau duration. On the other hand, after the neutrino trapping, as a result of the Urca process, the cooling for different initial temperatures would gradually approach to be a common
The neutrino luminosity evolves with time and is dependent on the initial temperature. For a sufficiently high initial temperature, the neutrino emission can take away the majority of the internal energy, leading to a temperature plateau at \( T \approx 10^9 \text{K} \), if a sufficient amount of material can be accreted onto the SS. This leads to a relatively high neutrino luminosity at the beginning. Because of the neutrino trapping, the early cooling of the SS is dominated by the neutrino diffusion rather than the Urca processes. This leads to a temperature plateau at the initial time, the duration of which is determined by the timescale of the neutrino diffusion. As a result, although the neutrino emission can take away the majority of the internal energy of the SS, the surface \( e^\pm \) pair emission due to the Uskov mechanism can still keep a constant luminosity of \( 10^{38} - 10^{39} \text{erg s}^{-1} \) for about a few to a few tens of seconds, which is dependent on the value of the initial temperature. For a relatively high temperature, the total energy carried by the \( e^\pm \) wind can reach to be as high as \( \sim 10^{51} \text{erg} \), which seems high enough to cause some observable consequences.

For example, it was suggested that such an \( e^\pm \) wind could give rise to a GRB emission, if this wind can further be accelerated to be relativistic by absorbing energy from the subsequent neutrino flux. As mentioned in the introduction, the afterglow emission of GRBs robustly suggest that some post-burst remnants could be a pulsar rather than a black hole. However, some people are concerned about how such a pulsar can produce a GRB outflow. It is even argued that only black holes can generate a GRB. Thus, the result obtained here demonstrates that the \( e^\pm \) wind from a bare proto-SS could be, at least partially, responsible for the GRB trigger. To say the least, if the GRB prompt emission indeed requires a hyper-accretion onto the SS as suggested by Zhang & Dai (2008a, 2008b), the \( e^\pm \) emission of a constant luminosity for a few tens of seconds could still provide a natural explanation for an extended emission after the prompt emission, which is usually found in short GRBs (Zhang 2013; Lü & Zhang 2014; Lü et al. 2015; Sun et al. 2017). This may indicate that these short GRBs originate from a merger of double pulsars and the merger product is a massive SS.

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4. Summary and Conclusion

In this paper, we establish a toy model to calculate the thermal evolution of a proto-SS, where the neutrino diffusion in the star is described effectively by calculating the inward movement of the neutrino sphere. This neutrino trapping effect has not been treated seriously in previous works on the same topic (Page & Uskov 2002; Cheng et al. 2009b; Cheng & Harko 2010), which leads to a relatively high neutrino luminosity at the beginning. Because of the neutrino trapping, the early cooling of the SS is dominated by the neutrino diffusion rather than the Urca processes. This leads to a temperature plateau at the initial time, the duration of which is determined by the timescale of the neutrino diffusion. As a result, although the neutrino emission can take away the majority of the internal energy of the SS, the surface \( e^\pm \) pair emission due to the Uskov mechanism can still keep a constant luminosity of \( 10^{38} - 10^{39} \text{erg s}^{-1} \) for about a few to a few tens of seconds, which is dependent on the value of the initial temperature. For a relatively high temperature, the total energy carried by the \( e^\pm \) wind can reach to be as high as \( \sim 10^{51} \text{erg} \), which seems high enough to cause some observable consequences.

Figure 4. The luminosities of the \( e^\pm \) pair emission (black) in comparison with the neutrino luminosity (gray) from proto-SSs of different initial temperatures.
