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What if pulsars are born as strange stars?

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Abstract

The possibility and the implications of the idea, that pulsars are born as strange stars, are explored. Strange stars are very likely to have atmospheres with typical mass of $\sim 5 \times 10^{-15} M_{\odot}$ but bare polar caps almost throughout their lifetimes, if they are produced during supernova explosions. A direct consequence of the bare polar cap is that the binding energies of both positively and negatively charged particles at the bare quark surface are nearly infinity, so that the vacuum polar gap sparking scenario as proposed by Ruderman and Sutherland should operate above the cap, regardless of the sense of the magnetic pole with respect to the rotational pole. Heat cannot accumulate on the polar cap region due to the large thermal conductivity on the bare quark surface. We test this “bare polar cap strange star” (BPCSS) idea with the present broad band emission data of pulsars, and propose several possible criteria to distinguish BPCSSs from neutron stars. © 2001 Elsevier Science B.V. All rights reserved.

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1. Introduction

Soon after the discovery of neutrons, the idea of neutron star was proposed [1,2]. The discoveries of radio pulsars and some other accretion-powered X-ray sources indicate that there do exist in nature certain objects with mass $M \sim M_{\odot}$, and radius $R \sim 10$ km. They were commonly regarded as neutron stars until the idea of strange star was proposed [3–5]. The basic idea of the strange star conjecture is that, strange quark matter (SQM), which is simply composed of an approximately

equal portion of up, down, and strange quarks and a few electrons to balance the non-neutral charges, might be more stable than the normal nuclear matters. According to the lattice quantum chromodynamics (QCD), a new state of strong interaction matter, the so-called quark–gluon plasma (QGP), will appear when the temperature of the matter achieves as high as ~ 150 – 200 MeV or the density of the matter achieves several times of ρ_0 , where $\rho_0 \sim 3 \times 10^{14}$ g cm⁻³ is the nuclear density. SQM is just one kind of such QGPs. Though we cannot tell whether SQM is the lowest state of hadronic matter, it is found that the energy per baryon of SQM is lower than that of the normal nuclear matter for a rather wide range of QCD parameters [3,6,7]. Assuming SQM is absolutely stable at zero pressure, strange stars, which almost

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completely consist of SQM, could exist in nature [4,5]. The assumption is *strong*, but not impossible. Therefore many objects which were previously believed to be neutron stars, including radio pulsars, might actually be strange stars. In fact, more and more strange star candidates have been proposed recently in the literatures [8–12], including both the accretion-powered and the rotation-powered compact objects.

Therefore, we have to think over a question again, which appeared to have been answered, “What is the nature of pulsars?” In their pioneering work, Alcock et al. [4] have suggested that radio pulsars might be strange stars. However, they did not fully discuss the possible observational consequences of such a strange star idea. Nor was this done by the later authors. What if pulsars are born as strange stars? How well can a strange star model interpret the pulsar observational data as compared with a neutron star model? Is there any criterion based on the radiation properties that can distinguish strange stars from neutron stars? These are some interesting issues worth exploring, and are the main topics of this paper. Suppose that strange stars are the ground state of neutron stars, the conversion from neutron stars to strange stars can occur in different stages of the neutron stars’ lifetimes. There are two main regimes discussed in the literatures. One group of models suggests that such a phase conversion occurs in the very late stage of a neutron star’s lifetime, after it accretes sufficient amount of material to make the core density exceed the critical density for phase transition [13,14]. Another group of models, on which we mainly focus in this paper, suggests that the conversion occurs in the very early stage of a neutron star’s life time, i.e., during or shortly after a supernova explosion [15–18], so that the observed radio pulsars might be strange stars. In the conventional strange star models, a crust composed of normal matter is usually assumed above a strange star’s bare quark surface [4,19], which tends to smear out the possible differences between the emission of strange stars and of neutron stars. An important issue raised in this paper which differs from the conventional models is that strange stars are almost bare, especially in the polar cap region, if they are

formed directly from supernova explosions. Such bare polar cap strange stars (BPCSSs), can also well act as radio pulsars [12,20]. In Section 2, we will discuss that phase transition from nuclear matter to SQM may be an important mechanism to retain a successful supernova explosion, and the strange stars, if born as the products of the explosions, are very likely to have *very* thin crusts ($\sim 10^{-15}M_{\odot}$) and bare polar caps almost throughout their lifetimes. The electrical and thermal properties of BPCSSs are discussed in Section 3, focusing on some distinguishing properties as compared with neutron stars. In Section 4, we test this BPCSS idea with the present available radio and high energy emission data of pulsars, and point out the strong and weak aspects of the idea. Finally we will propose some possible approaches to distinguish BPCSSs from neutron stars in Section 5, and summarize the conclusions in Section 6.

2. Newborn strange stars

2.1. *Supernova explosions: neutron stars or strange stars?*

The theory of stellar structure and evolution has been proved a great success in astrophysics. However, some challenges still exist in understanding a star’s life, such as the core-collapse supernova paradigm, which begins with the collapse of the iron core of an evolved massive star in the end of its thermonuclear evolution (for a comprehensive review, see e.g. Ref. [21]). Whether a model yields a successful explosion still challenges theorists and their numerical simulations. It is currently believed that the prompt shock, which results from the inner core’s rebound after an implosion has compressed the inner core to supranuclear density, cannot propagate directly outward and expel the entire envelope, but may stall and turn into an accretion shock at a radius of 100–200 km from the explosion center due to nuclear dissociation and neutrino cooling. It is a consensus now among different groups of supernova studies, that a successful explosion model requires a so-called “shock reheating mechanism”

or “delayed mechanism” [22,23], in which the neutrinos produced from the core are absorbed and/or scattered by the materials in the stalled shock and the shock could be revived to expel the envelope and give rise to a successful explosion. Many current researches on core-collapse supernova mechanisms are focused on the role of convection in the unstable regions below or above the neutrinosphere, and it seems that the supernova simulations without incorporating fluid instabilities may fail to explode (see, e.g., the report by Wilson and Mayle [24]). Generally, the simulations based on this mechanism give too low an energy, e.g. 0.3–0.4 foe (1 foe = 10^{51} ergs), to meet the observed energy of SN 1987A (at least 1.0 foe). Recent two-dimensional simulations by Mezzacapa et al. [25,26] showed that the neutrino-driven convection is not adequate to give an “optimistic” $15M_{\odot}$ supernova model, and the simulated time-scale for explosion is longer than what is observed. Furthermore, these authors pointed out that more realistic three-dimensional simulation may be more, not less, difficult to obtain a successful explosion.

Therefore, the key criterion for a successful explosion *and* its large enough energy should be the sufficient neutrino energy deposition behind the stalled shock. One possible way to enhance neutrino luminosity is through phase transitions, such as the transition from nuclear matter to two-flavor quark matter and from two-flavor quark matter to SQM [15–18]. After the core bounce, the central temperature and density of a proto-neutron star with a typical radius ~ 50 km [27] might be high enough to induce a QCD phase transition. A strange star will be finally formed if SQM is absolutely stable. An important issue is that whether the central density could achieve the value for phase transition. Unfortunately, this question is not easy to answer due to a large uncertainty of the MIT bag constant B . Schertler et al. [28] had included medium effects (by effective quark masses), the influence of which cannot be simulated by taking into account the coupling constant (α_s) or B or strange quark mass (m_s), in the study of SQM properties and found that medium effects reduce significantly the SQM binding energy (SQM could not be absolutely stable when $\alpha_s > 1$),

but can hardly change the mass and radius of strange stars if they exist. However, a recent investigation [29] shows that the phase transition from nuclear matter to quark matter can occur in proto-neutron stars as long as the bag constant $B \leq 126$ MeV fm $^{-3}$. Various estimates of the bag constant [30–32] indicate that the preferred value of B lies in the range of $60 \leq B \leq 110$ MeV fm $^{-3}$ (see, e.g., a review by Drago [33]), which means that a phase transition is likely to happen. Furthermore, Drago and Tambini [34] show that phase transition can occur at densities slightly larger than ρ_0 in pre-supernova matter.

Suppose that the phase transition condition is met, the energy budget problem for supernova explosions can be then naturally resolved [35]. Because each nucleon contributes about 30 MeV energy during phase conversion [11], the total released phase transition energy E_{pt} is then approximately

$$E_{\text{pt}} = Mc^2 \times \frac{30}{931} \sim 5.8 \times 10^{52} \frac{M}{M_{\odot}} \text{ ergs}, \quad (1)$$

where M is the mass of the inner core. The time-scale to burn a proto-neutron star to a proto-strange star, $\sim 10^{-4}$ s, is usually much smaller than that of neutrino diffusion and thermal evolution, ~ 0.5 s [21], since the actual combustion could be in the detonation mode [36–38] and the propagation velocity could be very probably close to the speed of light. Thus, the neutrino luminosity L caused by this conversion could be estimated as about $E_{\text{pt}}/0.5 \sim 10^{53}$ ergs $^{-1}$. In Wilson’s computations [22,23], the typical value of neutrino luminosity is 5×10^{52} erg s $^{-1}$, thus the total neutrino luminosity with the inclusion of SQM phase transition should be 1.5×10^{53} erg s $^{-1}$ (*three* times that of Wilson value). The simulated explosion energy should also be raised by a factor of 3, which is 0.9–1.2 foe, adequate to explain the observed value from SN 1987A.

Another possible way to enhance neutrino luminosity is through rapid rotation. A nascent neutron star can be rotating rapidly shortly after the supernova explosion [39], though we have little knowledge about how fast it rotates observationally and theoretically. The inner core of a rapid

rotating proto-neutron star is usually more *massive* than that of a non-rotating one in order to compress the inner core to about the same supra-nuclear density at which the prompt shock wave occurs. When the proto-neutron star de-leptonizes, more neutrinos are emitted from a larger core, and the neutrino luminosity behind the stalled shock could be enhanced to explode the supernova successfully.² Due to a larger centrifugal force from rapid rotation, the central density of the core could be much lower than the QCD phase transition density, so that a neutron star could exist after the explosion. However, such rapidly rotating neutron stars cannot keep being there for a long time because of the fine-tuning problem, proposed by Glendenning [41]. Only those neutron stars very close to the critical mass limit can rotate rapidly because they have the least radius and the greatest mass. These stars have much higher central densities than the lighter neutron stars. As these massive neutron stars spin down through dipole electromagnetic and quadruple gravitational radiations, the centrifugal force gets smaller and the central density could be increased high enough for a QCD phase transition, and the neutron stars will also inevitably be converted to strange stars.³

To sum up, the energy budget problem of the core-collapse supernova paradigm seems to require extra neutrino sources, and the phase transition from neutron matter to SQM is a natural mechanism to cure the imperfection of the paradigm. For the progenitors with negligible effect of rotation, additional neutrino emission of phase transition could sufficiently enhance the power of neutrino energy deposition behind the stalled shock to retain a successful explosion and its enough energy, with a strange star directly formed after the explosion. For the progenitors with rap-

idly rotating inner core, neutron stars might be formed as semi-products after the explosions, but such neutron stars will be finally phase-converted to strange stars when they spin down enough. Thus, it is plausible that at least some known radio pulsars may be in fact strange stars rather than neutron stars. Firm conclusions could not be drawn at present stage due to the large uncertainties involved in this problem. For example, most recently Mezzacappa et al. [42] claimed that they can successfully simulate a $13M_{\odot}$ supernova explosion with the exact Boltzmann neutrino transport, without invoking convection. Thus the question whether strange stars are “obliged” to be formed during the supernova explosions remains open. The aim of this paper is to explore the possible consequences of the assumption that strange stars are the final products, and such consequences could be tested by the observations and in turn, shed light on the plausibility of the scenario itself.

2.2. Newborn strange stars: crusts or not?

There are two kinds of strange stars discussed in the literatures, i.e., the strange stars with crusts composed of normal matter and the bare strange stars (strange stars with bare quark surfaces). Normal matter crusts above the bare quark surfaces of the strange stars were raised by Alcock et al. [4] in their pioneering paper. However, they did not discuss the formation of such crusts, but simply addressed that the universe is a “dirty” environment. Glendenning and Weber [43] discussed the maximum crust mass, $\Delta M \sim 10^{-5}M_{\odot}$ (later corrected to $\Delta M \sim 3 \times 10^{-6}M_{\odot}$ by Huang and Lu [44]), a strange star can sustain, mainly motivated by trying to reproduce pulsar glitches from strange stars. But they did not discuss the formation of such crusts, either. Usov [45] considered a strange star with a crust (or an atmosphere in his term) the mass of which is many orders of magnitude lower than the maximum value (only $\Delta M \sim (10^{-20}-10^{-19})M_{\odot}$), and studied the X-ray emission properties from such strange stars. He supposed that a strange star is more likely of this type shortly after its birth before the crust mass reaches the maximum. In principle, a

² We note that a proto-neutron star with short rotation period may generate strong magnetic field [40], which could play an important role in energizing the supernova shock if the magnetic reconnection or magnetohydrodynamic wave can significantly heat the neutrinosphere.

³ It is interesting to further investigate the corresponding astrophysical appearance of such conversion (e.g., may such kind of conversion act as the inner engine of the cosmic classical γ -ray bursts?).

pure bare strange star is unlikely to exist in the universe due to various accretion processes. The question is how thick the crust (or atmosphere) could be in reality. In this paper, we will show through some dimensional estimate that newborn strange stars are likely to have much thinner crusts than the maximum values and that they will keep like this through their lifetimes, unless they encounter some special environments that make sufficient accretion possible. The main reasons for such a conclusion are the rapid rotation and the strong mass ejection as discussed below. We note, however, that thick crusts (close to maximum) should be formed in the strange stars in the accreting systems (such as the compact objects in X-ray binaries) or the systems that clearly have strong accretion history (such as the “recycled” millisecond pulsars).

2.2.1. Rapid rotation

A rapidly rotating newborn strange star can prevent accretion from happening. Only when a strange star rotates slowly enough, could the accretion onto the surface be possible. There are three characteristic scale lengths in describing the accretion scenarios of magnetized compact stars. The magnetospheric radius (or Alfvén radius), defined by equating the kinematic energy density of free-fall particles with the magnetic energy density, is $r_m = ((R^6 B^2)/(\dot{M}\sqrt{2GM}))^{2/7} \sim 3.2 \times 10^{10} B_{12}^{4/7} R_6^{12/7} M_1^{-1/7} \dot{M}_{10}^{-2/7}$ cm, where $B = 10^{12} B_{12}$ G is the surface magnetic field of the star, $R = 10^6 R_6$ cm is the stellar radius, $\dot{M} = 10^{10} \dot{M}_{10}$ (g s^{-1}) is the accretion rate, and $M_1 = M/M_\odot$ is the stellar mass in unit of solar mass. The corotating radius, defined by the balance of gravitational force and the centrifugal force, is $r_c = (GM/4\pi^2)^{1/3} P^{2/3} = 1.5 \times 10^8 M_1^{1/3} P^{2/3}$ cm, where the rotation period P is in unit of second. The third scale length is the radius of light cylinder, which reads $r_l = cP/2\pi = 4.8 \times 10^9 P$ cm. Possible accretion of matter into a pulsar’s magnetosphere requires $r_m < r_c < r_l$ (see, e.g., Refs. [46,47]). The constraint of $r_c < r_l$ gives

$$P > 3.1 \times 10^{-5} M_1 \text{ s}, \quad (2)$$

which could be satisfied for all known pulsars. The requirement that the Alfvén radius is smaller than

the corotating radius ($r_m < r_c$), however, is very tight, which reads

$$P > 3.2 \times 10^3 B_{12}^{6/7} R_6^{18/7} M_1^{-5/7} \dot{M}_{10}^{-3/7} \text{ s}, \quad (3)$$

since the Alfvén radius could not get too low unless the accretion rate is very high, e.g., $\dot{M}_{10} \gg 1$. Note that in Ref. [47] the maximum accretion rate (Eddington luminosity) is adopted.

Let’s give a rough estimate on the accretion rate of a solitary strange star. Consider a strange star which is moving with a velocity V through a gas with a density ρ . The critical radius within which the dynamics of gas is dominated by the stellar gravity is $r_g \sim GM/V^2 \sim 1.3 \times 10^{12} M_1 V_7^{-2}$ cm, where $V_7 = V/(10^7 \text{ cm s}^{-1})$. Thus according to Bondi and Hoyle [48], the accretion rate is roughly

$$\begin{aligned} \dot{M} &\sim 4\pi r_g^2 \rho V \sim 2.2 \times 10^8 M_1^2 V_7^{-3} \rho_{24} \text{ g s}^{-1} \\ &\sim 3.5 \times 10^{-18} M_1^2 V_7^{-3} \rho_{24} M_\odot \text{ year}^{-1}, \end{aligned} \quad (4)$$

where $\rho_{24} = \rho/(10^{-24} \text{ g cm}^{-3})$. By comparing a sample of pulsar proper motion data with their Monte Carlo simulations, Hansen and Phinney [49] found that the mean speed of pulsars at birth is $V \sim (2.5\text{--}3.0) \times 10^7 \text{ cm s}^{-1}$. Diamond et al. [50] investigated the local interstellar medium (ISM) distribution with ROSAT all-sky EUV survey, and found that the ISM mean density is $\sim 10^{-25}\text{--}10^{-24} \text{ g cm}^{-3}$. Thus the minimum rotation period for accretion of a typical new born strange star in a typical ISM is $\sim 10^4$ s. In supernova remnants, the medium density may be denser so that $\rho \sim 10^{-22} \text{ g cm}^{-3}$. In this case, the critical period for possible accretion is still $\sim 10^3$ s. Since the longest period of the present known pulsars is only 8.5 s [51], we conclude that accretion is almost impossible for most isolated radio pulsars in their lifetimes, unless they are born with much less-than-average proper motion speed or they are in very dense medium. Another reason against the possible accretion is that when a pulsar is active, relativistic particles are believed to flow out from the open field line region. Accretion onto the polar cap is thus impossible with such an outflow.

2.2.2. Mass ejection

We have shown that accretion is almost impossible for a pulsar after its birth. The question

now is whether strong accretion is possible during the supernova explosion. It is known that a new born neutron star is very hot (the core temperature $T_c \sim 10^{11}$ K), and the mass ejected from such a hot neutron star during the first 10 s is about $\sim 10^{-3} - 10^{-2} M_\odot$ [52,53]. It is expected that a hot new born strange star should also have such a high mass ejection rate, so that its surface is almost bare, since the ejected mass is considerably larger than the maximum crust mass of the strange star [45,54]. As discussed in Section 2.1, the rate of energy release when forming a strange star is $\sim 10^{53}$ erg s $^{-1}$. Most of this energy is carried away by neutrinos, but part of it will be converted to electromagnetic radiation. A nuclei above the quark surface will feel strong outward pressure by photons and neutrinos (note that the Eddington luminosity is only $\sim 10^{38}$ erg s $^{-1}$), therefore normal nuclei can hardly remain being bound by the gravitational force above a strange quark surface and will be pushed away. For a neutron star, the requirement of the hydrodynamical equilibrium will make the core materials to self-adjust themselves to form a stable structure, including a very heavy crust above the superfluid neutron layer, and the backfalling supernova ejecta has little contribution to such a crust. For a strange star, things are quite different. Almost the whole iron core will be completely converted to SQM, since SQM has the property to swallow all the normal matter that directly contacts with it if Witten's [3] hypothesis is right. Also the detonation flame may expel the outer part of a proto-neutron star to form a strange star with a bare quark surface [36–38] at the very beginning. Thus a crust above the quark surface of a strange star then solely depends on the fallback of the supernova ejecta.

There are a lot of uncertainties on the possible fallback of the supernova ejecta. An important issue is whether a reverse shock can be formed. It was argued that one needs to include a reverse shock to fit the light curve of SN 1987A, and such a reverse shock might be due to a sufficiently massive hydrogen envelope around the compact star [55]. If the formation of such a reverse shock is possible, Chevalier and other authors [56–59] investigated the possibility of the accretion with hypercritical (or “super Eddington”) accretion

rate, and found that the total accreted matter could be as high as $10^{-3} M_\odot$ for a normal type II supernova, which is much higher than the maximum crust mass a strange star can sustain. In such cases, the strange star formed from the explosion should have a thick crust with maximum mass. However, the Chevalier's scenario does not include the influences of rotation and magnetic fields. When these effects are taken into account, more likely, the backfalling materials will form a disk [60] rather than just directly falling back onto the surface of the compact star. According to Michel [61], such a disk is essential for pulsar magnetospheric electrodynamics. Some models also invoked such a disk to interpret the emission behavior of the anomalous X-ray pulsars [62]. Thus it is very likely that the mass which falls back to the pulsar surface should be tiny compared to the maximum, as Usov [45,54] suggested.

Let us present a rough estimate of the mass and the thickness of the crust formed due to the direct fallback. In principle, the materials that can fall back onto the surface should be trapped by the magnetosphere and be below the corotating radius r_c . When no more outflow pressure exists, all the materials within the corotating ball will drop onto the surface if they are trapped by the magnetosphere. The condition of magnetic trapping requires the local magnetic energy density to be comparable of the kinetic energy density of the ejecta, so that the trapping radius is $r_t = RB^{1/3}v^{-1/3}(4\pi\rho)^{-1/6}$, where v is the typical velocity of the ejected materials, and ρ is the mean density of the materials. We set $r_c = r_t$ to roughly set the fallback condition, and get a rough estimate of $\rho \sim 7.0 \times 10^{-9} B_{12}^2 P^{-4} v_9^{-2} R_6^6 M_1^{-2}$ g cm $^{-3}$ (about one day after the supernova explosion for $P = 10$ ms, see Fig. 28 in Ref. [55]). The total mass that deposits onto the surface (i.e. the mass of the crust) is then $\Delta M = (4/3)\pi r_c^3 \rho$, which reads

$$\Delta M \sim 1.0 \times 10^{17} B_{12}^2 P^{-2} v_9^{-2} R_6^6 M_1^{-1} \text{ g}, \quad (5)$$

or $\Delta M \sim 5 \times 10^{-15} M_\odot$ for typical values of $P = 10$ ms, $B_{12} = R_6 = M_1 = 1$. Note that the typical velocity of the ejecta is adopted as $v = 10^9 v_9$ cm s $^{-1}$, which is the typical velocity of the revived shock [63]. The adoption of a smaller velocity will not

influence our followup conclusion qualitatively. Notice that ΔM obtained here is 10 orders of magnitude smaller than the maximum ΔM , but 5 orders of magnitude larger than the “atmosphere” mass discussed by Usov [45]. We will call the structure formed by the fallback materials a *massive atmosphere* rather than a thin crust (following Usov [45,54]).

We now estimate the depth of this atmosphere. The column density of the atmosphere is $\sigma_a \sim \Delta M / (4\pi R^2) \sim 8.0 \times 10^3 B_{12}^2 P^{-2} v_9^{-2} R_6^4 M_1^{-1} \text{ g cm}^{-2}$. The bottom pressure of the atmosphere is thus $p_b = (GM\sigma_a/R^2) \sim 1.0 \times 10^{18} B_{12}^2 v_9^{-2} R_6^2 P^{-2} \text{ dyn cm}^{-2}$. According to the equation of state [64]

$$p = 1.8 \times 10^{12} \rho^{5/3} \text{ dyn cm}^{-2}, \quad (6)$$

for a non-relativistic, completely degenerate gas (since the density is much smaller than 10^6 g cm^{-3} , see below), the bottom density is about $\rho_b = 5.8 \times 10^5 \text{ g cm}^{-3}$ for $P = 10 \text{ ms}$. Therefore the scale height of a strange star atmosphere is $h_0 \sim \sigma_a / \rho_b \sim 1.4 \times 10^2 \text{ cm}$ for a 10 ms pulsar (thinner crust for longer period pulsars). We note that such an atmosphere scale height is much larger than the characteristic thickness of neutron star atmosphere ($\sim 0.01\text{--}1 \text{ cm}$, see, e.g. Ref. [65]). This is an essential feature. It is known that a completely bare strange star has very low emissivity in X-rays [4]. We emphasize here that the existence of the massive atmosphere described above makes a strange star to radiate thermal emission in a similar way as does a neutron star (except the polar caps, see below). Another comment is that the “atmosphere” discussed here is conceptually different from the atmosphere of a neutron star, which is defined by the equilibrium of the gravitational energy and the thermal kinetic energy of the particles. Neutron star atmospheres are naturally formed even without accretion. The strange star “atmospheres” discussed in this paper, however, are formed due to the materials’ fallback, and hence, are not subject to the definition above.

Shortly after a bare strange star is born, it is believed that a typical pulsar magnetosphere with Goldreich–Julian [66] density will be formed slightly above the bare quark surface (typical height $z_c \sim 10^{-8} \text{ cm}$, see Section 3.1), mainly due

to the pair multiplication via γ - B , γ - γ [20], or γ - E [54] processes. When such a magnetosphere is formed, a space-charge-limited free flow [67] is then inevitable in the open field line regions, so that the atmosphere materials above the polar caps, which are composed of normal hadrons, will be pulled out. In neutron stars, such a flow could persist throughout the pulsars’ whole lifetimes since the thick crusts of the neutron stars can supply copious ions and electrons to be pulled out. In the case discussed here (a strange star with an atmosphere), the particles available for this extracting is limited, which is only the part of atmosphere right above the polar caps. The timescale for pulling out this atmosphere is short. With a Goldreich–Julian flow, the extracting column number density flux is $F = (\Omega B / 2\pi e) \sim 2.1 \times 10^{21} B_{12} P^{-1} \text{ cm}^{-2} \text{ s}^{-1}$. Given the atomic mass unit $u = 1.66 \times 10^{-24} \text{ g}$, the extracting column mass density flux is then $F_m \sim 2uF$, therefore the extracting timescale is typically

$$\begin{aligned} \tau &= \sigma_a / F_m \\ &\sim 3.6 \times 10^{-2} B_{12} P^{-1} v_9^{-2} R_6^4 M_1^{-1} \text{ year}, \end{aligned} \quad (7)$$

which is much shorter than a pulsar’s lifetime. Furthermore, if the timescale for the ejecta to fallback is longer than the timescale of forming a Goldreich–Julian magnetosphere, the space-charge-limited flow from the polar cap region will effectively block the deposit of the falling materials onto the polar cap area. We thus address that except for very young pulsars, the polar caps of the strange stars with massive atmospheres are likely to be completely bare when they act as pulsars.

Note that in the above estimate (Eq. (7)), we have assumed that only the matter in the polar cap regions is stripped out. This can be justified as follows. The magnetic field energy density at the pole is $\epsilon_B = B^2 / 8\pi \sim 4.0 \times 10^{22} B_{12}^2 \text{ ergs cm}^{-3}$. The energy density of the matter, which is essentially the electron degenerate energy density, could be estimated as $\epsilon_e = \epsilon_F n_e$, where ϵ_F is the electron Fermi energy, $\epsilon_F = 2\pi^4 \hbar^4 c^2 n_e^2 / (e^2 B^2 m) \sim 1.03 \times 10^{-63} B_{12}^{-2} n_e^2$ [69], and m and n_e are the mass and number density of the electrons, respectively. For n_e , we adopt the electron number density at the “effective” pulsar surface of a BPCSS (Eq. (15) for

$z = z_c$), i.e., $n_e \sim 10^{27} \text{ cm}^{-3}$. This gives $\epsilon_e \sim 10^{18} \text{ ergs cm}^{-3}$, which is much smaller than ϵ_B . This prohibits the rapid penetration of the electrons across the magnetic field lines. However, the charged particles can also diffuse across the magnetic field lines due to the collisions between the charged particles that gyrate around the lines, since a collision will alter a particle's velocity and make it to gyrate around another field line. The Fermi energy of the electrons in the case we are discussing is much higher than the energy interval between adjacent Landau levels of the electrons, we thus adopt the classical descriptions to estimate the diffusion rate. Following [68], the diffusion coefficient D_c is approximately given by the square of the mean distance travelled (which is of the order of the Larmor radius $\rho_L = mvc/eB$) divided by the mean free flight time τ_F , which is of the order of $\tau_F \sim m^2 v^3 / \pi e^4 n$ (derived from Eq. (2.55) or (2.65) of Ref. [68], assuming that light elements dominate in the atmosphere). Thus the diffusion rate is $D_c \sim \rho_L^2 / \tau_F \sim \pi c^2 e^2 n / B^2 v$. Let us estimate the electron diffusion rate first. For the bottom density $\rho_b \sim 5.8 \times 10^5 \text{ g cm}^{-3}$, the electron number density is $n_b \sim \rho_b / (2u) \sim 10^{29} \text{ cm}^{-3}$, the Fermi energy is $\sim 10^{-5} B_{12}^{-2} n_{b,29}^2 \text{ ergs}$, and the electron velocity $v \sim c$, thus one has

$$D_c \sim \frac{\pi c e^2 n}{B^2} \sim 2.2 \times 10^{-3} B_{12}^{-2} n_{b,29} \text{ cm}^2 \text{ s}^{-1}. \quad (8)$$

As particles diffuse, the scale height h of the massive atmosphere should be a function of the polar angle θ , and also a function of time t before a steady state is achieved. After a certain period of time, the diffusion flow will be steady. Let's consider the following equilibrium diffusion situation: at the polar cap boundary ($\theta = \theta_p = 1.45 \times 10^{-2} B_{12} P^{-2/3}$), $h = 0$; at a much larger polar angle (e.g. the equator $\theta = \pi/2$), $h = h_0$. Again using the equation-of-state Eq. (6), one gets

$$h(\theta) \sim \frac{\sigma_a(\theta)}{\rho_b(\theta)} \sim 3.0 \times 10^{-18} n_b(\theta)^{2/3} \text{ cm}. \quad (9)$$

For a simple estimation of the upper limit, we consider a layer with height $h(\theta)$ and with a uniform density $n(\theta) \sim n_b(\theta)$ in the following discussions. In the steady diffusion regime, the diffuse rate

$$I_{\text{dr}} = D_c \frac{dn(\theta)}{R d\theta} 2\pi R \sin \theta h(\theta) \quad (10)$$

is a constant. Solving the above equation, one obtains

$$3.1 \times 10^{21} h(\theta)^4 = I_{\text{dr}} \ln \left[\frac{\tan(\theta/2)}{\tan(\theta_p/2)} \right], \quad (11)$$

or $2.5 \times 10^{-49} n(\theta)^{8/3} = I_{\text{dr}} \ln[\tan(\theta/2)/\tan(\theta_p/2)]$ for $n(\theta) - \theta$ relation. Setting $n \sim 10^{29} \text{ cm}^{-3}$ for $\theta = \pi/2$, we get an upper limit diffusion rate $I_{\text{dr}} \sim 10^{28} \text{ s}^{-1}$. For smaller n_b , one gets smaller I_{dr} . On the other hand, the free flow rate with the Goldreich–Julian flux is $I_{\text{SCLF}} \sim \pi(R\theta_p)^2 F \sim 10^{34} \text{ s}^{-1}$ for $P = 10 \text{ ms}$. This indicates that the diffusion is unimportant. As for the collision diffusion of the ions, the rate should be even smaller as their moving velocity is much smaller due to their larger mass. Furthermore, the total number of the nuclei in the atmosphere is $\Delta M/u \sim 10^{45}$, so that the timescale for the atmosphere height to change significantly is 10^{17} s ($> 10^9$ years). All these indicate that the atmospheres can exist almost through the whole life of the pulsars, and that only the polar caps are likely bare.

Besides rapid rotation and high mass ejection discussed above, there are some more reasons which support the picture we proposed. For example, newborn strange stars usually have very high temperatures, which can significantly reduce the Coulomb barrier, and can thus result in fusion of nuclei and SQM by tunneling effect [19,45,54]. This is also in favor of the formation of strange stars without thick crusts.

In conclusion, if a strange star is born as the product of a supernova explosion, a thick crust is unlikely to be formed via the direct fallback of the supernova ejecta. Only a very thin crust, or a massive atmosphere, can exist above the bare quark surface of the strange star, and the polar caps are completely bare due to the space-charge-limited flows. Such a situation usually remains unchanged through out the pulsar's lifetime, since accretion is almost impossible due to rapid spin of the pulsar. Hereafter we will define such strange stars as BPCSSs.

3. Electric and thermal properties of the bare polar cap strange stars

In this section, we will summarize the electric and thermal properties of the BPCSSs discussed in the Section 2. We will focus on some special properties of the BPCSSs as compared with neutron stars.

3.1. Electric properties

Since a BPCSS could have a bare quark surface in the polar cap region, it is important to investigate the electric characters near the quark surface. As the strange quarks are more massive than the up and down quarks, some electrons are required to keep the chemical equilibrium of a strange star. This brings some interesting properties near the bare quark surface of a BPCSS. Since quark matter is bound through strong interaction, the density changes abruptly from $\sim 4 \times 10^{14} \text{ g cm}^{-3}$ to nearly zero in 1 fm at the surface, which is the typical length scale of the strong interaction. The electrons, which are bound by the electromagnetic interaction and are $\sim 10^3$ – 10^4 times less denser than the quark materials, can spread out the quark surface and be distributed in such a way that a strong outward static electric field is formed. Adopting a simple Thomas–Fermi model, one gets the Poisson’s equation [4]

$$\frac{d^2 V}{dz^2} = \begin{cases} \frac{4\alpha}{3\pi} (V^3 - V_q^3) & z \leq 0, \\ \frac{4\alpha}{3\pi} V^3 & z > 0, \end{cases} \quad (12)$$

where z is the height above the quark surface, α is the fine-structure constant, and $V_q^3/(3\pi^2\hbar^3 c^3)$ is the quark charge density inside the quark surface. A straightforward integration gives [70]

$$\frac{dV}{dz} = \begin{cases} -\sqrt{\frac{2\alpha}{3\pi}} \sqrt{V^4 - 4V_q^3 V + 3V_q^4} & (z < 0), \\ -\sqrt{\frac{2\alpha}{3\pi}} V^2 & (z > 0), \end{cases} \quad (13)$$

where the physical boundary conditions $\{z \rightarrow -\infty: V \rightarrow V_q, dV/dz \rightarrow 0; z \rightarrow +\infty: V \rightarrow 0, dV/dz \rightarrow 0\}$ have been adopted. The continuity of V at $z = 0$ requires $V(z = 0) = 3V_q/4$, thus the solution for $z > 0$ finally leads to

$$V = \frac{3V_q}{\sqrt{\frac{6\alpha}{\pi} V_q z + 4}} \quad (\text{for } z > 0). \quad (14)$$

The electron charge density can be calculated as $V^3/(3\pi^2\hbar^3 c^3)$, therefore the number density of the electrons is

$$n_e = \frac{9V_q^3}{\pi^2 \left(\sqrt{\frac{6\alpha}{\pi} V_q z + 4} \right)^3} \sim \frac{9.5 \times 10^{35}}{(1.2z_{11} + 4)^3} \text{ cm}^{-3}, \quad (15)$$

and the electric field above the quark surface is finally

$$E = \sqrt{\frac{2\alpha}{3\pi}} \frac{9V_q^2}{\left(\sqrt{\frac{6\alpha}{\pi} V_q z + 4} \right)^2} \sim \frac{7.2 \times 10^{18}}{(1.2z_{11} + 4)^2} \text{ V cm}^{-1}, \quad (16)$$

which is directed outward. In the above estimate, $V_q \sim 20 \text{ MeV}$ has been adopted, and $z_{11} = z/(10^{-11} \text{ cm})$.

It was believed that the strong outward electric field near the quark surface has some implications on the properties of strange stars. One implication is that a possible normal-matter crust could be formed, which in this paper we claim to be thin. Another issue is that, according to Ref. [4], “a rotating magnetized star with an exposed quark surface will not supply the charged particles necessary to create a corotating magnetosphere”, since “the electric field induced by the rotating magnetized star is small compared to the electric field at the surface”. However, here we claim that this argument is incomplete. A handy estimate from Eq. (16) shows that, although the electric field near the surface is about $5 \times 10^{17} \text{ V cm}^{-1}$, the outward electric field decreases very rapidly above the quark surface, and at $z \sim 10^{-8} \text{ cm}$, the field gets down to $\sim 10^{11} \text{ V cm}^{-1}$, which is of the order

of the rotation-induced electric field for a typical Goldreich–Julian [66] magnetosphere. Here we define the critical height, z_c , at which the strength of the intrinsic electric field is equal to that of the rotation-induced field as the effective pulsar surface [12], thus a typical pulsar magnetosphere could be naturally formed above this surface [20].

3.2. Thermal properties

Investigations of the strange star cooling behavior have been performed by many authors. It has been argued that the cooling behavior of the young pulsars can act as a definite criterion to distinguish strange stars from neutron stars, since strange stars may cool much faster than neutron stars (e.g. Ref. [71]). However, recent more complete analyses [72,73] on this issue indicate that, direct Urca process could be also forbidden in strange stars if the electron fraction of the SQM is relatively low. Furthermore, quarks may eventually form Cooper pairs, and such possible superfluid behavior of the SQM can also substantially suppress the neutrino emissivities of various processes. As a result, the surface temperature of the strange stars with crusts (typically with maximum mass) should be more or less similar to the neutron star surface temperature. Neutron stars and strange stars with crusts are hence indistinguishable in their cooling behaviors except for the first ~ 30 year after their births [72,73]. Present X-ray data from about 30 rotation-powered pulsars are consistent with the standard neutron star cooling scenario [74], and thus do not contradict the idea of strange stars with thick crusts (close to the maximum). Nonetheless, notwithstanding extensive efforts, the study on this *global* cooling behavior of strange stars does not reach an agreement among researchers, especially about the effect of the possible color superconductivity in SQM [75–77]. In this paper, however, we suggest (see Section 3.2.2) to focus on the *local* thermal properties of the polar caps, which may provide some distinguishing criteria for BPCSSs and neutron stars, while to regard the global cooling behavior as an open issue to explore.

3.2.1. Cooling of bare polar cap strange stars

For the BPCSS picture we are discussing in this paper, since the crust (atmosphere) is much thinner than the maximum, more explicit work need to be done to explore the detailed cooling behavior as compared with that of the strange stars with thick crusts. This issue is beyond the scope of this paper. Nevertheless, we expect that the cooling curves of the BPCSSs with atmospheres may not differ too much from those of the strange stars with thick crusts. The main reason is that the BPCSS discussed in our case is not completely bare. A thick atmosphere ($\Delta M \sim 5 \times 10^{-15} M_\odot$) above the surface of a bare strange star makes that the strange star can dissipate the heat from the quark core through thermal emission at the surface in a similar way that does a strange star with a thick crust (although the response function could be quite different [71]), since the thickness of the layer fulfills the optically thick condition. We note that the atmosphere described by Usov [45,54] is much thinner ($\Delta M \sim (10^{-19} - 10^{-20}) M_\odot$) than the atmosphere discussed in this paper. In his case, the atmosphere is optically thin, and the strange star has much higher hard X-ray emissivity than a neutron star or a strange star with a thick crust.

The crust or atmosphere acts as a thermal insulator between the hot quark core and the surface [71], thus the thickness of the crust determines the time delay when the surface and the core have a same temperature. A thinner crust will drive the temperature dropping point on the cooling curve to an even earlier epoch. For the case of pulsar, the youngest pulsars with measured surface temperatures (or upper limits) have the ages of several 10^3 year, much longer than the dropping point on the cooling curves of the strange stars with thick crusts (~ 30 year) [72,73], thus thinner crusts on strange stars may not bring inconsistency between the cooling theories and the observational data.

3.2.2. Polar cap heating

Pulsar polar caps are believed to be hotter than the rest of the surface due to the reheating by the bombardment of the downward-flowing particles and their radiation. The downward-flowing particles could be produced from the inner accelerator (either of vacuum type or of space-charge-limited

flow [hereafter SCLF] type) or the outer gap, and the degree of polar cap heating is hence model-dependent. The vacuum gap model proposed by Ruderman and Sutherland [78] predicts substantial polar cap heating, since the numbers of the downward and upward particles are similar when the gap breaks down due to the pair production avalanche (essentially with the Goldreich–Julian density), so that the luminosity deposited onto the polar cap is just the polar cap luminosity brought by the primary particles, which reads

$$L_{\text{pc,v}} = \gamma mc^2 \dot{N}_p \simeq 1.1 \times 10^{31} \gamma_7 B_{12} P^{-2} \text{ ergs s}^{-1}, \quad (17)$$

where $\gamma = 10^7 \gamma_7$ is the typical Lorentz factor of the primary particles, and $\dot{N}_p = cn_{\text{GJ}} \pi r_p^2 = 1.4 \times 10^{30} R_6^3 B_{12} P^{-2}$ is the particle flow rate with Goldreich–Julian density n_{GJ} . The SCLF models [67, 79,80], however, predict much less polar cap heating [81,82], since only a small fraction, $f = 10^{-4} f_{-4}$, of the primary particles (see the definition of the parameter f in Ref. [82], their Eq. (63)) could be reversed in the space-charge-limited electric field and be accelerated back to the surface. The energy deposited onto the surface in this model is then

$$L_{\text{pc,sclf}} = f \gamma mc^2 \dot{N}_p \sim 1.1 \times 10^{27} f_{-4} \gamma_7 B_{12} P^{-2} \text{ ergs s}^{-1}. \quad (18)$$

If outer gaps exist in some pulsars, the downward-accelerated particles from these outer gaps may also hit the polar cap eventually after losing substantial amount of their initial energies, and deposit $\sim 5.9 P^{1/3}$ ergs per particle when they strike the surface [83]. With Goldreich–Julian density, the luminosity deposited onto the surface is approximately

$$L_{\text{pc,og}} \simeq 5.9 P^{1/3} \dot{N}_p \simeq 8.2 \times 10^{30} B_{12} P^{-5/3} \text{ ergs s}^{-1}. \quad (19)$$

Note that both the vacuum gap model and the outer gap model predict very strong polar cap heating, the luminosities of which are comparable to the total spin-down luminosity of the pulsars, which reads

$$L_{\text{sd}} \simeq 9.7 \times 10^{30} B_{12}^2 P^{-4} I_{45} \text{ ergs s}^{-1}. \quad (20)$$

The polar cap heating temperature can be estimated by $T_{\text{pc}} = [L_{\text{pc}}/(\sigma \pi r_p^2)]^{1/4}$, where $r_p = 1.45 \times 10^4 P^{-1/2}$ cm is the polar cap radius, and $\sigma = 5.67 \times 10^{-5}$ ergs cm⁻² K⁻⁴ s⁻¹ is the Stefan’s constant. In the curvature radiation controlled vacuum gap model (see Eqs. (17) and (25)), the polar cap temperature is $T_{\text{pc,v}} = 5.9 \times 10^6 B_{12}^{3/14} P^{-3/14}$ K. The SCLF model gives a lower temperature $T_{\text{pc,sclf}} = 2.2 \times 10^6 B_{12}^{1/14} P^{-1/14}$ K (Eq. (70) in Ref. [82]). The outer gap model predicts a medium temperature $T_{\text{pc,og}} = 3.9 \times 10^6 B_{12}^{1/4} P^{-1/6}$ K.

As the polar cap is hotter than the other part of a pulsar’s surface, heat may flow from the cap to the surrounding area. If pulsars are neutron stars, such heat flow is negligible, and the kinetic energy of the backflowing particles and the energy of the electromagnetic shower produced by these particles can be almost completely converted into thermal energy and be re-radiated back to the magnetosphere. Let us give a rough estimate. The coefficient of thermal conductivity for electron transport in the neutron star surface [84]

$$\kappa^{\text{NS}} = 3.8 \times 10^{14} \rho_5^{4/3} \text{ ergs s}^{-1} \text{ cm}^{-1} \text{ K}^{-1}, \quad (21)$$

where ρ_5 is the density in unit of 10^5 g cm⁻³, is almost independent on the details of the lattice. If we assume that the horizontal temperature gradient in the crust of a neutron star is roughly $\nabla T^{\text{NS}} \sim T_{\text{pc}}/r_p$, and that the area of the heat flow is of the order of r_p^2 , then the heat flow rate is roughly

$$\begin{aligned} H^{\text{NS}} &\sim \kappa^{\text{NS}} \nabla T^{\text{NS}} r_p^2 \sim \kappa^{\text{NS}} T_{\text{pc}} r_p \\ &\sim 5.5 \times 10^{24} \rho_5^{4/3} T_6 P^{-1/2} \text{ ergs s}^{-1}, \end{aligned} \quad (22)$$

which is much smaller than L_{pc} . Therefore, the thermal conduction from the polar cap to the surrounding area is unimportant [84] for pulsars being neutron stars.

However, if pulsars are BPCSSs, the electron number density in the bare quark surface $n^{\text{BPCSS}} = 1.5 \times 10^{34}$ cm⁻³, (Eq. (15) with $z_{11} = 0$, see also Ref. [70]), is much larger than that of a neutron star, $n^{\text{NS}} = 2.8 \times 10^{28} \rho_5$ cm⁻³, thus the thermal diffusion could be much effective. The transport coefficients for degenerate quark matter

due to quark scattering had been calculated by Heiselberg and Pethick [85], from Eq. (61) of which the thermal conductivity of strange star can be obtained,

$$\kappa^{\text{BPCSS}} = 1.41 \times 10^{22} \left(\frac{\alpha_s}{0.1} \right)^{-1} \rho_{15}^{2/3} \text{ ergs s}^{-1} \text{ cm}^{-1} \text{ K}^{-1}, \quad (23)$$

where α_s is the coupling constant of strong interaction, ρ_{15} is the SQM density in $10^{15} \text{ g cm}^{-3}$. Thus we have $\kappa^{\text{BPCSS}} \sim 10^{22} \text{ ergs s}^{-1} \text{ cm}^{-1} \text{ K}^{-1}$ for $\alpha_s \sim 0.1$, $\rho_{15} \sim 1$. The corresponding horizontal heat flow rate in the bare quark surface is therefore of the order of

$$H^{\text{BPCSS}} \sim 10^{32} \text{ ergs s}^{-1}, \quad (24)$$

which is even larger than L_{pc} . This means that very likely, the heat deposited onto the polar caps by the backflowing particles will be soon dissipated to the other part of the BPCSS surface, so that no hot polar cap could be sustained.

It is worth noting that if the heat deposited onto the BPCSS surface could be dissipated to the other regions at the surface, this will add an additional full surface thermal component besides the cooling component. The temperature of this component can be estimated as $T_2 = [L_{\text{pc}}/(\sigma 4\pi R^2)]^{1/4}$. For the vacuum gap case, this is $T_2 \simeq 5.0 \times 10^5 B_{12}^{3/14} P^{-13/28}$ K. The inclusion of this component does not contradict the cooling observational data for the normal pulsars [74]. For the millisecond pulsars, since these pulsars no longer appear as BPCSSs, such a relationship no longer holds. We note that the extra full surface thermal component is also expected in the internal heating model [74] and the outer gap model [83,86].

4. Testing bare polar cap strange star idea with pulsar data

Both neutron stars and strange stars have been proposed as the nature of pulsars. However, almost all the previous researches have invoked thick crusts above the bare quark surfaces, which tend to smear out the information from the SQM itself. As a result, distinguishing strange stars from

neutron stars is a difficult task. One has to appeal to some other criteria to do the discriminations. These include equation-of-state [8,9], cooling behavior [71–73], dynamically damping effect [87], minimum rotation period (e.g. whether there exist sub-millisecond pulsars) [88], the vibratory modes [89], and so on. Some of these criteria, e.g., cooling behavior, suffer large uncertainties. We have shown that if pulsars are born as strange stars, they will very likely appear as BPCSSs. The exposure of the bare quark surface at the polar caps makes it possible that the information from the surface can be directly transmitted out, which may influence pulsar emission behavior to some extent. This opens a new window to distinguish strange stars from neutron stars according to their different emission behaviors. In this section we will discuss the consequences of the bare polar caps and test the BPCSS idea with the present available pulsar data.

4.1. Consequences of the bare polar caps

Two direct consequences follow naturally from the BPCSS picture.

The first consequence is that the inner accelerator in the polar cap region is vacuum-like, as proposed by Ruderman and Sutherland [78], since the binding energy at the surface is almost infinity for both positive and negative charges. There are two sub-types of pulsar inner gap models, according to the boundary condition at the surface, i.e., the vacuum gap model [78] and the SCLF model [67,79,80]. Though both models share some common features, they are different in some other aspects (for a comparison between the two models, see Refs. [90,91]). If pulsars are neutron stars or strange stars with thick crusts, binding energy calculations favor the SCLF scheme, both for the case of $\boldsymbol{\Omega} \cdot \mathbf{B} < 0$ and $\boldsymbol{\Omega} \cdot \mathbf{B} > 0$ (see a brief review in [12,92]). However, if pulsars are BPCSSs, vacuum gaps are preferred. This is obvious for the case of $\boldsymbol{\Omega} \cdot \mathbf{B} < 0$, i.e., anti-parallel rotators, because in this case positive charged u quarks, which are expected to flow out, are definitely impossible to be pulled out since they are bound by strong interaction. For the case of parallel rotators (i.e. $\boldsymbol{\Omega} \cdot \mathbf{B} > 0$), electrons, which are bound electro-

magnetically, might be pulled out, contingent upon the competition between the intrinsic electric field at the BPCSS surface (Eq. (16)) and the rotation-induced electric field. With Eq. (16), we see that only those electrons above the height z_c could be pulled out. Thus a vacuum gap could be formed above the “effective” BPCSS surface defined by z_c . As we have discussed in Section 2.2.2, electrons are forbidden to cross field lines due to the strong magnetic field energy density, thus only electrons above z_c in the polar cap region could be stripped out. The time scale for stripping these electrons is about 10^{-5} s [70], which is quite small. A steady SCLF accelerator is therefore not possible. Although in the very beginning of the growth of gap it is not completely vacuum, the gap will be eventually evacuated when the flow ceases. Thus the accelerators for the parallel rotators are also vacuum-like.

The basic picture of the vacuum gaps formed above the polar caps of pulsars have been delineated explicitly by Ruderman and Sutherland [78]. In this picture, a vacuum gap is formed right above the surface of the star. Primary particles are accelerated to extremely relativistic energies and emit γ -rays via curvature radiation [78] or inverse Compton scattering with the thermal photons near the surface [93–95]. These γ -rays are materialized in the strong magnetic fields via the γ - B process, and the secondary pairs screen out the parallel electric fields so that the vacuum gap is limited at a certain height (the gap height). A main feature of the vacuum gap which distinguishes from a SCLF gap is the periodic breakdown of the gaps, which produces sparks and secondary plasma clumps ejected into the outer magnetosphere. Since the curvature radius of magnetic field lines is smaller near the polar cap edge, the sparks tend to take place in a ring-like region near and within the polar cap edge. The secondary plasma produced in the sparks hence form some plasma columns or mini-tubes. Each spark forms a clump, and these plasma clouds are ejected sporadically from the gap and are separated from each other spatially. Due to the $\mathbf{E} \times \mathbf{B}$ drifting, the sparks are expected to rotate around the magnetic pole with a certain speed, and recent long term observations have revealed that the drifting sub-pulses observed in

some pulsars clearly match the Ruderman and Sutherland’s [78] prediction [96,97]. We hope to emphasize that, although the vacuum gap model was first proposed to operate on neutron stars, later binding energy calculations indicate that such gaps can be hardly rooted to neutron stars [12,92]). The inner gap rooted to a neutron star or a strange star with thick crust has to be modified into other forms, e.g., the completely free flow without any binding [67] or the flow with partial binding [92,98]. Since none of these models can reproduce the exact drifting features predicted by the vacuum model, the recent observational results [96,97] could be regarded as an important support to the BPCSS scenario [12].

The second direct consequence of the BPCSS scenario is that hot polar caps cannot be formed due to the large thermal conductivity as discussed in Section 3.2.2. We will discuss more about the implications of this feature for pulsar X-ray emission theories in Section 4.4.

More than 1000 radio pulsars have been detected so far. Among them, 35 are detected as X-ray sources, 11 are X-ray pulsars, and 8 are γ -ray pulsars [99–101]. A wealth of broad band emission data have been accumulated. In the following, we will test the BPCSS idea with the pulsar broad band emission data, focusing on the possible difference between the neutron star scenario and the BPCSS scenario.

4.2. Radio emission

Pulsar radio emission data are abundant compared with most of the other astrophysical objects. However, theories lag observations a big phase. More than 10 models have appeared in the literatures [102], and they differ from each other in many aspects. Perhaps the only consensus among these models is that pair production is the essential condition. Since both vacuum gaps and SCLF gaps can produce secondary pairs, pulsars will emit radio emission as long as one of these inner accelerators operate [90]. To distinguish BPCSSs from neutron stars, one needs to seek observational properties which characterize vacuum gaps rather than SCLF gaps, or vice versa.

As discussed above, the clearly drifting sub-pulse patterns discovered recently [96,97] seem to be a support to the BPCSS scenario rather than the neutron star scenario. Definite conclusion cannot be drawn until it is proved that (1) SCLF models or any other modified vacuum gap models definitely cannot reproduce the right drifting rate as observed, and that (2) there is no way to solve the binding energy problem within the neutron star scenario. At present, it seems that BPCSS scenario is the only way to revive the vacuum gap model, and to interpret the sub-pulse drifting rate.

Some observational features also directly infer the “sparkling” behavior from the pulsar inner gaps. Individual pulses are often composed of one or more sub-pulses, and some of these sub-pulses drift regularly. Pulsar micro-structures have fine structures of the order of 10^{-5} – 10^{-6} s, which is the breakdown timescale of a typical Ruderman–Sutherland gap.

Besides these direct effects, it seems that vacuum gaps have advantages to interpret some other radio emission data than SCLF gaps. Generally speaking, self-absorption of the radio emission limits the incoherent brightness temperature to a level much lower than what is observed. To interpret the extremely high brightness temperature observed from pulsars, certain coherent mechanisms are believed to play the role. Various methods have constrained the height of radio emission to be well within the light cylinder [103,104]. However, the strongest electromagnetic instabilities, i.e., the maser instabilities (which do not depend on the type of the inner gaps), occur at altitudes close to the light cylinder [105], which is not favored by the empirical laws of pulsar radio emission uncovered by Rankin [106,107]. Most low-altitude radio emission theories, on the other hand, require that inner accelerators should display certain “oscillation” behaviors which resemble the quasi-periodical breakdown of the vacuum gaps.

For example, the inverse Compton scattering model proposed by Qiao and Lin [108] attributes pulsar radio emission to the coherent inverse Compton emission of the secondary particles. This model can reproduce radio pulsar phenomenology well, including one core and two conal emission

components found by Rankin [106]; the linear and circular polarization features [109,110] and the frequency evolution of the pulse profiles [111]. The basic ingredient of this model is a vacuum gap, the periodic breakdown of which can naturally excite the low-frequency electromagnetic wave, which is the target of the inverse Compton scattering of the particles. Furthermore, since the pair plasma ejected from the gap is confined to plasma columns in strong fields, the outflowing plasma should be highly inhomogeneous in space, and the density between the miniflux tubes could be sufficiently low, which allows the low-frequency radio wave to propagate as if in vacuum.

Another example is the spark-associated soliton model recently proposed by Gil et al. [112,113]. According to the authors, the plasma instability invoked to interpret pulsar radio emission is the only known low-altitude instability. Such overtaking two stream instability requires pair plasma to be ejected in clumps, and the sparks produced from a vacuum gap can naturally provide such spatially separated plasma clumps.

We are not saying that SCLF model cannot interpret pulsar radio emission, though. In fact, in the neutron star scenario, a SCLF gap is preferred from the binding energy calculations, and some pulsars, e.g., the newly discovered 8.5 s pulsar PSR J2144-3933 [51], are preferably interpreted by the SCLF model [90]. Even in the strange star scenario, millisecond pulsars, which show similar emission behavior as normal pulsars, should also have SCLF accelerators since we expect thick crusts above the quark surfaces formed from their “recycling” history. What we hope to address is that, certain pulsar emission features (e.g. drifting sub-pulses) do not favor the SCLF model, and these features are not observed from the millisecond pulsars and the 8.5 s pulsar. Thus in order to interpret all the pulsar radio emission features (especially the drifting subpulse phenomenon which is popular among conal emission pulsars) within the neutron star scenario, either the SCLF model should be such modified to include certain periodic oscillation behaviors, or some fundamental progress is made to solve the binding energy problem faced by the vacuum gap model. At present stage, neither of these two possibilities are

promising, therefore at least for some pulsars, BPCSS scenario is favored.

4.3. Gamma-ray emission

In contrast to the radio emission theories, pulsar γ -ray emission theories are usually grouped into only two types, the polar cap cascade models [82,114,115] and the outer gap models [86,116,117] (for a comparison between the two models, see e.g. Ref. [82]). As outer gaps are much far away from the surface (above the null charge surface), different polar cap properties will not bring differences in the outer gaps, thus γ -ray data cannot be used to differentiate the BPCSS scenario and the neutron star scenario if γ -rays are of outer gap origin.

If pulsar γ -rays are of polar cap origin, however, different types of the inner accelerators will bring differences in the γ -ray emission properties. Contrary to radio emission, γ -ray emission data seem to favor the SCLF picture rather than the vacuum gap picture. But as we will show below, the BPCSS picture, i.e. the vacuum gap picture, cannot be completely ruled out.

The basic γ -ray emission properties of the known γ -ray pulsars include a luminosity law $L_\gamma \propto (L_{\text{sd}})^{1/2}$ [101], and wide separations between the double peaks observed in several pulsars. Within the curvature radiation controlled SCLF model [80], the luminosity law is well obeyed since the typical Lorentz factor of the primary particles does not sensitively depend on pulsar parameters (P, B_p) [82]. The wide separations of the double peaks could be interpreted by the extended polar cap scenario [114], since the SCLF accelerator may be lifted to a higher altitude due to the anisotropy of the inverse Compton scattering of the upwards versus downward primaries [80].

If the inner accelerator is of vacuum-type, the gap should be formed right above the surface. This does not favor the interpretation of the γ -ray emission. First, it is possible that there exist some strong multipole magnetic field components near the surface (these components have been long assumed [78]), which tend to lower the gap height and limit the achievable γ -ray luminosity (essentially the polar cap luminosity). However, if we assume that the near surface magnetic field con-

figuration is dominated by the dipolar component for the known γ -ray pulsars, the $L_\gamma \propto (L_{\text{sd}})^{1/2}$ luminosity law can be retained. For a curvature radiation controlled vacuum gap, with pure dipolar field, the total voltage across the gap is $\Delta V = 2.1 \times 10^{13} P^{1/7} B_{12}^{-1/7}$ V, and the typical Lorentz factor of the particles is $\gamma = 4.0 \times 10^7 P^{1/7} B_{12}^{-1/7}$. With Eq. (17), we get

$$L_\gamma \simeq L_{\text{pc,v}} = 4.4 \times 10^{31} B_{12}^{6/7} P^{-13/7} \text{ ergs s}^{-1}, \quad (25)$$

which is similar to the result of the SCLF model with no electric field saturation (Eq. (59) in Ref. [82]). This is understandable, since when the SCLF accelerator is not saturated, the $\Delta V - h$ law is also approximately quadratic ($\Delta V \propto h^2$), just as the vacuum gap model [90]. Eq. (25) will also give a similar diagram as Fig. 3 in [82], which approximately reproduces the $L_\gamma \propto (L_{\text{sd}})^{1/2}$ feature as reported by the observations [101].

Another drawback of the BPCSS scenario is that, since vacuum gaps are formed on the surfaces, very small inclination angles between the magnetic axis and the rotational axis are required to account for the observed widely separated double γ -ray peaks. The detectability of the γ -ray pulsars is also lowered. Some statistic studies show that such small inclination angles are not contradictory to the number of the γ -ray pulsars presently detected [118], but one needs to answer such questions like ‘‘Why young pulsars are born with very small inclination angles?’’. Nonetheless, keeping in mind that outer gaps are another alternative of γ -ray emission site, the BPCSS scenario does not strongly contradict the pulsar γ -ray emission data.

4.4. X-ray emission

In the neutron star scenario, X-ray emission of the spin-powered pulsars could in principle have three components, i.e., a non-thermal component with typical power-law spectrum, a thermal component from the full neutron star surface, and a hotter thermal component at the polar cap. The non-thermal component has been observed in many pulsars; the full-surface thermal component has been observed in four pulsars [119–121]; and

the evidence for hot polar cap emission is only strong for the millisecond pulsar PSR J0437-4715 [122] (for a brief review, see e.g. Ref. [82]). A luminosity law, $L_x \sim 10^{-3}L_{sd}$, was found [100].

These data again could be interpreted either by the outer gap model [86] or the full polar cap cascade model [82]. The outer gap model, which attributes the non-thermal X-ray emission to the synchrotron radiation of downward cascade particles [86], will make no difference between the BPCSS scenario and the neutron star scenario. The polar cap model, which attributes the non-thermal X-ray emission to the inverse Compton scattering of the upward cascade particles [82], may give different predictions between the two different scenarios. However, if we again assume dipolar-dominated configuration for the young pulsars, the non-thermal X-ray luminosities will not differ too much from the one predicted in the SCLF model [82], since both models have similar polar cap luminosities (Eq. (25)).

Thermal X-ray emission, especially that from the polar caps, may be a tool to differentiate between the two scenarios. In the neutron star scenario, only the SCLF model predicts a cool polar cap; both the vacuum gap model and the outer gap model predict a much hotter polar cap than observed (see Section 3.2.2). To avoid such an inconsistency, the outer gap models have invoked an assumed pair blanket slightly above the surface, which can reflect the hard thermal emission from the polar cap back to the neutron star surface [83,86]. An important consequence of the BPCSS scenario is that no hot polar cap could be formed (Section 3.2.2), which can naturally amend the hot polar cap problem encountered by both the outer gap model and the vacuum gap model. Evidence of a hot polar cap from the millisecond pulsar PSR J0437-4715 was noticed (e.g. Ref. [122]), but this poses no objection to the BPCSS scenario since millisecond pulsars are believed to have thick crusts.

4.5. Other issues

Besides the broad band emission data, some other pulsar phenomenology may also shed light on the discrimination between the two scenarios.

Pulsar “glitches” have been observed in some young pulsars. The giant glitches observed from the Vela pulsar is regarded as a strong evidence against the strange star model, even for the strange stars with the thickest crusts [123]. The BPCSS scenario is even disfavored. We may think that glitching pulsars might be neutron stars, but firm conclusion still cannot be drawn, since some efforts have been made to produce strong glitches in strange stars. Benvenuto and Horvath [124,125] presented a calculation through the introduction of the effects of the hypothetical few-quark bound state (quark alpha). An important feature emerging from their calculation is that strange stars may become shell-structured when quark-alpha particles are introduced. Glitches may arise from these shells with a similar behavior in the shells of the neutron stars. They further suggest that the pulsars’ post-glitch behavior observed should be attributed mainly to the de-coupling and re-coupling of the fluid-quark-alpha layer. However, they note that some important problems remain unsolved. More studies are necessary before firm conclusions can be drawn.

Recently, Madsen [126] argued that bare strange stars can be ruled out as the candidate of the very fast millisecond pulsars, according to the dependence of the rotational mode instabilities on the thermal behavior of the strange stars. This is an interesting test, but is unfortunately not applicable to the case discussed here, since in our BPCSS scenario, the millisecond pulsars have thick crusts due to their accretion history.

5. Distinguishing bare polar cap strange stars from neutron stars

In this paper, we are trying to open a new window to distinguish strange stars from neutron stars using the exterior properties of pulsars which are of *magnetospheric origin*. As we have shown, the idea of BPCSS raises some interesting distinguishing properties between the two scenarios. Although present data are not sufficient to draw definite conclusions, the discrimination between the two scenarios is expected as data accumulate.

Here we summarize some distinguishing criteria which may act as a guide for future observations.

1. Clearly drifting patterns from the “antipulsars”. As discussed in Ref. [12] and above, drifting pulsars give a strong support to the BPCSS idea. The support is more prevailing if the pulsar is a parallel rotator ($\boldsymbol{\Omega} \cdot \mathbf{B} > 0$), since in this case, vacuum gaps are even less impossible to be formed above a neutron star’s surface and it is referred to “antipulsar” in Ruderman and Sutherland’s [78] term. Unfortunately, the sense of the magnetic pole is indistinguishable from the polarization data. Seeking other methods to tell the sense of the magnetic pole is essential. The premise of this criterion is that the SCLF model cannot be developed to manipulate the “sparking” behavior and to predict correct drifting rate.

2. Polar cap temperature. Discussions in Section 3.2.2 suggest that most probably, BPCSSs might not have hot polar caps. This raises the possibility to distinguish strange stars from neutron stars using the *local* thermal behavior rather than the *global* thermal behavior (which is reflected in their cooling curves and is subject to many microphysics processes poorly known). A very hot polar cap on a normal isolated pulsar (no accretion history) will be a strong evidence against the BPCSS idea. Notice that the lack of polar cap temperature as high as several 10^6 K cannot rule out the neutron star model since the SCLF models predict much lower polar cap temperatures (see e.g. Eqs. (70) and (73) in Ref. [64]). However, if there is completely no evidence for hot polar cap emission, then the BPCSS idea is justified.

3. Line features. The different composition of the surfaces of neutron stars and BPCSSs may provide another interesting feature to differentiate between the two scenario. Neutron star surfaces are copious of positive ions. Within the SCLF picture, these ions should be pulled out from the surface and fill the magnetosphere in the open field line region for those pulsars with $\boldsymbol{\Omega} \cdot \mathbf{B} < 0$. For BPCSSs, however, nothing can be pulled out from the quark surface, and the magnetosphere is just filled with the electron–positron pairs. The different compositions in the open field line region of pulsar magnetospheres may result in quite different

emission spectra of the magnetospheric X-ray radiation, and some line features may be expected for the NS scenario. These line features could be above the observational level of the modern or future X-ray missions (e.g. Chandra, XMM). Unfortunately, no detailed theoretical treatment of such radiative transfer process in the magnetospheres has been published so far. Such a detailed study is desirable. An important bearing is that, if the theoretical results show that the intensity of the line features (e.g. K_{α} line) is above the detection level of some future missions, an either positive or negative detection of these lines can provide strong evidence against or for the BPCSS hypothesis. At present, no line feature has been reported from the known spin-powered X-ray pulsars.

Besides the three criteria listed above, any other distinguishing emission properties between the vacuum gap and the SCLF gap may act as a criterion for differentiating BPCSSs from neutron stars, as long as the binding energy problem cannot be solved within the neutron star scenario.

6. Conclusions and discussion

The idea that pulsars might be strange stars rather than neutron stars has long been proposed, but no detailed study on the strange star model of pulsars has been presented previously. In this paper, we have explored the possibility and the corresponding implications of the idea that pulsars are born as strange stars rather than neutron stars. Our main conclusions can be summarized as follows.

1. The lack of a mechanism producing a successful core-collapse supernova explosion has been a long-standing problem for supernova explosion studies. The phase transition from nuclear matter to SQM presents a natural way to solve this problem.
2. If pulsars are born as strange stars, very probably a thick crust cannot be formed after the explosion, nor can it be formed during the long history of its lifetime, unless the star experiences some special events or happens to be located in a superdense environment.

3. The strange stars, when they act as pulsars, may have an atmosphere with mass $\Delta M \sim 5 \times 10^{-15} M_{\odot}$, and two completely bare polar caps. We refer to such stars as BPCSSs. Pulsars in accreting binaries and the millisecond pulsars should have thick crusts.
4. A direct consequence of the bare polar caps is that BPCSSs should have vacuum-type inner gaps above their polar caps, regardless of the orientation of the magnetic axis with respect to the rotation axis.
5. The thermal conductivity in the bare quark surface is much larger than that in the neutron star surface. As a result, BPCSSs may not have hot polar caps.

We have put this BPCSS picture into test with the current available pulsar broad band observational data. We hope to emphasize that, with present data, the idea that pulsars are born as strange stars has no strong conflict with the observations, and it even has an advantage to revive the vacuum gap model, which is almost impossible in neutron stars. The vacuum gap model has been useful in various aspects of radio emission theories, and is supported by some recent observations. Detections of hot polar caps from normal isolated pulsars and line features from the X-ray spectra of some rotation-powered pulsars will rule out the BPCSS idea, but negative detections with sufficient sensitivity can be a support to the BPCSS idea.

One interesting question is that “can some pulsars be strange stars while some others be neutron stars?”. People tend to accept either “all pulsars are neutron stars” or “all pulsars are strange stars” for the sake of beauty. Caldwell and Friedman [127] proposed that if some compact objects are strange stars, essentially all “neutron stars” in the disk of the galaxy should be strange stars due to the strangeness contamination. They thus object the strange star idea using the glitching phenomenon observed from some pulsars. However, with some pulsars showing glitches, there is strong evidence that some other compact objects (e.g. the millisecond X-ray pulsar SAX J1808.4-3658) are strange stars following the equation-of-state arguments [9]. We thus suggest that the actual distribution of the known pulsars might be

bimodal. In fact, from the discussions above, we cannot draw the conclusion that all the pulsars are strange stars, nor can we say that they are all definitely neutron stars. The main motivation of this paper is to pose the possibility of regarding strange stars, especially BPCSSs, as another possible candidate for pulsars, and our idea is subject to the tests of future observations, with some possible criteria listed in Section 5.

All our discussions in this paper are based on the assumption that SQM is an absolutely stable hadronic state, since we have assumed the existence of the strange stars. Unfortunately, no hitherto-known theory can validate or reject this assumption. In terrestrial physics, searching for the new state of strong interaction matter, QGP, is the primary goal of the relativistic heavy-ion laboratories. Many proposed QGP signatures have been put forward theoretically and many experimental data have been analyzed. However, the conclusion about the discovery of QGP is still ambiguous [128]. The confirmation of the existence of strange stars in the universe, the significance of which might be comparable to that of confirming black holes, may bring profound implications to the fundamental physics, and thus will remain a hot topic in the new century.

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