Fast Radio Bursts from the Collapse of Strange Star Crusts

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Abstract

Fast radio bursts (FRBs) are transient radio sources at cosmological distances. No counterparts in other bands have been observed for non-repeating FRBs. Here we suggest the collapse of strange star (SS) crusts as a possible origin for FRBs. SSs, which are composed of almost equal numbers of u, d, and s quarks, may be encapsulated by a thin crust of normal hadronic matter. When a SS accretes matter from its environment, the crust becomes heavier and heavier. It may finally collapse, leading to the release of a large amount of magnetic energy and plenty of electron/positron pairs on a very short timescale. Electron/positron pairs in the polar cap region of the SS can be accelerated to relativistic velocities, streaming along the magnetic field lines to form a thin shell. FRBs are produced by coherent emission from these electrons when the shell is expanding. Basic characteristics of observed FRBs can be explained in our model.

Key words: radio continuum: general – radiation mechanisms: non-thermal – stars: neutron

1. Introduction

Fast radio bursts (FRBs) are a new kind of phenomenon that has been discovered in the past decade (Lorimer et al. 2007; Keane et al. 2012; Thornton et al. 2013). These transient bursts have flux densities of $S_f \sim$ a few Jy at frequencies of $f_{\text{FRB}} \sim 1$ GHz, with the waveband width of $\Delta f_{\text{FRB}} \sim$ several hundred MHz. Their durations, $\delta t$, are typically a few ms, indicating a rather compact region of emission. The observed high dispersion measurements (DMs) of $\sim$500–1000 pc cm$^{-3}$ are well above the contribution from our Galaxy for several FRBs detected at high-galactic-latitude ($\geq$40°; Cordes & Lazio 2002; Luan & Goldreich 2014), suggesting that the sources are at cosmological distances of $d \sim$ Gpc with redshifts of $z \sim$ 0.5–1. Hence the isotropic luminosities in radio waves ($L_{\text{FRB}}$) are estimated as $\sim$$10^{42}$–$10^{43}$ erg s$^{-1}$, with the total isotropic energy released in a typical burst being $E_{\text{FRB}} \sim 10^{39}$–$10^{40}$ erg. The event rate is estimated to be $\sim$2 $\times$ 10$^3$ sky$^{-1}$ day$^{-1}$ (Bhandari et al. 2018). The brightness temperatures of FRB sources can be as high as $T_B \geq 10^{36}$ $\Gamma^2$ K (Katz 2014; Luan & Goldreich 2014), where $\Gamma$ is the Lorentz factor of the emitting material. Such an extremely high temperature is far above the Compton limit for incoherent synchrotron radiation, thus a coherent origin shall be considered (Romero et al. 2016). On the other hand, no counterparts in other wavebands have been detected to be associated with non-repeating FRBs hitherto.

The origin of FRBs still remains unclear, but a number of models trying to interpret these enigmatic phenomena have been proposed, e.g., magnetar giant flares (Kulkarni et al. 2014), the collapses of magnetized supramassive rotating neutron stars (NSs; Falcke & Rezzolla 2014; Zhang 2014), binary NS mergers (Totani 2013), binary white dwarf mergers (Kashiyama et al. 2013), collisions between NSs and asteroids/comets (Geng & Huang 2015), collisions between NSs and white dwarfs (Liu 2017), and evaporation of primordial black holes (Barrau et al. 2014). Some of these models are catastrophic and the original central engines are destroyed completely by the bursts. The discovery of the repeating FRB source, i.e., FRB 121102 (Scholz et al. 2016; Spitler et al. 2016), presents new interesting clues to FRBs. Several elaborate models have been put forward to explain its repeating behaviors, e.g., highly magnetized pulsars traveling through asteroid belts (Dai et al. 2016), NS-white dwarf binary mass transfer (Gu et al. 2016), and star quakes of pulsars (Wang et al. 2018). However, it is wondered whether FRB 121102 is representative of FRBs since its features are different from others (Michilli et al. 2018; Palaniswamy et al. 2018). Naturally, people speculate that repeating FRBs such as FRB 121102 may be a separate kind of FRB source.

Here we propose a new model for FRBs. We argue that the collapse of SS crusts can also explain the main features of FRBs. The structure of this paper is as follows. In Section 2, the process of SS crust collapse is illustrated. The emission mechanism that leads to observable FRBs is described in Section 3. Possible counterparts in other wavebands are discussed in Section 4. Finally, Section 5 is a brief summary and discussion.

2. Collapse Process of SS Crust

It has been conjectured that strange quark matter (SQM), a kind of dense material composed of approximately equal numbers of up, down, and strange quarks, may have a lower energy per baryon than ordinary nuclear matter (such as $^{56}$ Fe) so that it may be the true ground state of hadronic matter (Abe et al. 1984; Witten 1984). If this hypothesis is correct, then NSs may actually be “strange stars” (Alcock et al. 1986). The bulk properties of SSs and NSs are rather similar in the typical mass range of $1 \leq M/M_\odot \leq 1.8$, and it is very difficult to discriminate between them (Haensel et al. 1986). Although several methods have been proposed to distinguish SSs from NSs (Geng et al. 2015; Lai et al. 2018), no definitive conclusions have been drawn yet.

At the SS SQM surface, the density reduces from $\sim$5 $\times$ 10$^{14}$ g cm$^{-3}$ to zero abruptly. The thickness of the SQM surface is of order 1 fm due to strong interaction between quarks, while electrons can stretch up to several hundred fm beyond the surface since they are bounded electromagnetically. An extremely intense electric field ($\sim$5 $\times$ 10$^{17}$ V cm$^{-1}$)
induced by charge separation near the SQM surface (Alcock et al. 1986). The outward-directed electric field can polarize a layer of nearby normal matter and provide a force overwhelming the gravity (Huang & Lu 1997; Steiner & Madsen 2005). As a result, a thin crust composed of normal hadronic matter may exist and obscure the whole surface of the SQM core. It has been shown by Huang & Lu (1997) that the maximum density at the bottom of the crust should be significantly less than the so-called neutron drip density. For a typical SS with a radius of \( r \sim 10 \) km, a mass of \( M \sim 1.4 M_{\odot} \), and a surface temperature of \( T_s \sim 3 \times 10^7 \) K, the crust mass is usually in the range of \( M_c \sim 10^{-7} M_{\odot} \sim 10^{-5} M_{\odot} \) and its thickness is about \( l \sim 2 \times 10^6 \) cm.

The distance between the bottom of the crust and the surface of the SQM core shall be at least \( \sim 200 \) fm so that the rate at which ions penetrate the gap through the tunneling effect is low enough to ensure the stabilization of the crust (Alcock et al. 1986). If the mass of the crust increases continuously via some accreting process, then the gap between the crust bottom and the SQM surface will become narrower and narrower to counterpoise the gravity of the crust and hence the gap width (Kettner et al. 1995). Consequently, a faster tunneling penetration is stimulated by the decreased gap width. This is a positive feedback and the SS crust will finally collapse completely on a free-fall timescale of \( \sim 0.1 \) ms.

Although the detailed mechanism for maintaining a strong magnetic field in various compact stars is still largely uncertain, it is believed that SSs can also have a strong magnetic field. When the crust of an SS breaks and falls into the SQM core, the magnetic field lines in the crust will be dragged into the core due to Alfvén’s frozen effect. A fraction of the magnetic energy originally embedded in the crust will be transferred to radiation. In fact, after the collapse of the SS crust, the magnetic field lines near the polar cap region will be disturbed and twisted because of differential rotation and/or magnetic instabilities (Kashiyama et al. 2013). Hence transient dissipation processes such as magnetic reconnection may be triggered (Thompson & Duncan 1995), and the magnetic energy will be released on a short timescale.

Theoretically, the limiting interior magnetic field is of the order of \( B_{\text{max}} \approx 10^{18} M_1 r_6^{-2} \) G, where the SS is assumed to have a mass of \( M = M_1 M_{\odot} \) and a radius of \( r = r_6 10^6 \) cm (Lai & Shapiro 1991). The convention \( Q = Q_1/10^5 \) in cgs units is adopted hereafter. Pulsars with surface magnetic fields up to \( \sim 10^{18} \) G have been reported, and there is no “smoking-gun” evidence to identify them as NSs or SSs (Kouveliotou et al. 1998). It is reasonable to postulate that some SSs could have a surface magnetic field as strong as \( B_S \sim 10^{15} \) G in the polar cap region, where the field should be the strongest for a dipole configuration. A dipole field approximation of \( B(R) \approx B_S \times (R/r)^{-3} \) is applied in our paper, where \( R \) is the distance from the SS center. The total magnetic energy stored in the crust can be expressed as \( E_B \approx 4 \pi \gamma J_1 \times B_S^2 / 8 \pi \tau \sim 5 \times 10^{43} B_S^2 L_0^{1/2} r_6^{3/2} \) erg.

In our framework, we believe that an FRB is produced mainly from the polar cap region, so not all the \( E_B \) energy is available for generating the FRB. The angular size of the polar cap region is approximately \( \theta_{\text{cap}} \sim 1.45 \times 10^{-2} P_0^{-1/2} r_6^{1/2} \), where \( P \) is the rotation period of the SS. The magnetic energy included in the polar region can then be estimated as \( E_{B,\text{cap}} \sim E_B \times \pi \theta_{\text{cap}}^2 / 4 \pi \sim 2.6 \times 10^{49} P_0^{-1/2} B_S^{1/4} L_0^{1/4} r_6^{1/2} \) erg. We suppose that the radio emission is restricted in an area with the solid angle of \( 4 \pi f \), where \( f \) is a parameter characterizing the beaming fraction. The energy needed to produce an FRB is then \( f E_{\text{FRB}} \), where \( E_{\text{FRB}} \) is the isotropic FRB energy. Assuming \( \eta \) to be the fraction of \( E_B, \text{cap} \) that is emitted, it can be calculated as \( \eta \approx f E_{\text{FRB}} / E_{B,\text{cap}} \sim 0.4 \times P_0 B_S^{2} L_0^{-1/4} r_6^{-3/4} \). If we take \( f \sim 0.1 \) and \( P \sim 1 \) s (Kondratyuk et al. 1990), \( \eta \) can be reckoned as \( \sim 0.04 \).

In short, as long as a small fraction of the magnetic energy conserved in SS polar cap regions is transferred into radiation during the collapse process, the energy is adequate for an FRB. FRBs should be connected with a coherent emission mechanism for their extremely high brightness temperature (Katz 2014). We will discuss the detailed radiation process below.

3. Emission Mechanism

After the crust collapse, the SS becomes hot and bare. It turns into a powerful source of electrons and positrons \((e^+e^-) \) pairs. These \( e^+e^- \) pairs are created in an extremely strong electric field (Usov 1998a). Near the polar cap region where the magnetic field energy is released, electrons and positrons will be accelerated to ultra-relativistic speeds (Ruderman & Sutherland 1975; Benford & Buschauer 1977). These particles coast along the magnetic field lines and form a shell with a thickness of \( \delta r_{\text{emi}} \approx c \beta t \). This process is illustrated in Figure 1.

Curvature radiation will be produced when electrons in the shell are streaming along parallel magnetic field lines. For simplicity, we postulate that all electrons are moving with almost exactly the same velocity (i.e., with the corresponding Lorentz factor of \( \gamma \)), and \( \delta r_{\text{emi}} \) remains roughly constant in the observer frame as long as \( \delta r_{\text{emi}} > \gamma^{-2} r_{\text{emi}} \), like the fireballs in
the cases of gamma-ray bursts (Rees & Mészáros 1992). Let $r_{emi}$ be the emission radius of the shell and $r_c$ be the curvature radius of the magnetic field lines, then the shell volume is $V_{emi} \approx 4\pi r_{emi}^2 \delta r_{emi}$. The patch in which electrons radiate coherently has a characteristic radial size of $\lambda \approx c/r_c$ and a solid angle of $4/\gamma^2$ according to beaming effects. Such a patch has a volume of $V_{coh} \approx (c/\nu_c) \times (4/\gamma^2) r_{emi}^2$ (Kashiyama et al. 2013; Kumar et al. 2017). In the emission region, there are $N_{pat} \approx V_{emi}/V_{coh}$ coherent patches totally, and each patch has $N_{coh} \approx n_e \times V_{coh}$ electrons in it, where $n_e$ is the electron number density. The frequency of curvature emission is $\nu_c \approx \gamma^3 (3c/4\pi r_c) \approx \nu_{FRB}$. The total coherent curvature emission luminosity from these electrons can be expressed as $L_{total} \approx (P\tau^2_{coh}) \times N_{pat}$, where $P_c = 2\gamma^2 e^2 c/3r_c^2$ is the emission power of a single electron (Kashiyama et al. 2013).

According to Benford & Buschauer (1977), the coherent radiation peaks at the places where the relativistic plasma energy density just exceeds the dipolar field energy density. With the plasma pressure of $\rho_p(R) \approx n_e^2 m_e c^2$, the magnetic energy density of $\rho_B(R) \approx B^2(R)/8\pi$, and assuming $r_{emi} \sim r_c$, the emission radius can be estimated as

$$r_{emi} \sim 0.6 \times 10^{10} (f_3^{-1} \frac{\delta L_{FRB}}{P_{FRB,9} B_{5,14}},6)^{-1} 25 \ cm.$$  
(1)

On the other hand, the electron Lorentz factor $\gamma$ can be derived as

$$\gamma \sim 1120 \nu_{FRB,9}^2 c_{10,13}^{-3/2}.$$  
(2)

The total coherent curvature emission luminosity from these electrons can be expressed as (Kashiyama et al. 2013)

$$L_{total} \approx (P\tau^2_{coh}) \times N_{pat} \sim 2.6 \times 10^{42} \ erg \ s^{-1} \times f_{nu}^2 \frac{\delta L_{FRB}}{P_{FRB,9} B_{5,14}},6 r_{emi,10}^{-1/3}.$$  
(3)

Then the observed isotropic FRB luminosity $L_{FRB} \approx f^{-1} L_{total} \sim 10^{42-43} \ erg \ s^{-1}$ is consistent with the observations.

The typical emission frequency should be lower than the plasma characteristic frequency so that the radio waves can propagate without being absorbed (Benford & Buschauer 1977; Kashiyama et al. 2013), i.e.,

$$\nu_{FRB} \gtrsim f_{pc} \approx \left( \frac{n_e e^2}{\pi m_e} \right)^{1/2}.$$  
(4)

It requires that

$$n_e \lesssim 1.1 \times 10^7 \nu_{FRB,9}^5 c_{10,13}^{-1/3} \ cm^{-3}.$$  
(5)

Another requirement is that the induced Compton scattering should not be too strong to hinder the propagation of the radio wave in the plasma (Melrose 1971). The Lorentz factor of the relativistic plasma is therefore limited by (see Equation (53) of Lyubarsky 2008)

$$\gamma \gtrsim 730 \gamma_T^{1/4} \xi^{-1/4} \nu_{FRB,9}^{1/4} \times \left( \frac{F}{1 \text{ Jy}} \right)^{1/4} \left( \frac{\delta t}{5 \text{ ms}} \right)^{-3/8} \left( \frac{D}{100 \text{ Mpc}} \right)^{1/2}.$$  
(6)

Here $\gamma_T \sim 1$ is the thermal Lorentz factor of the electrons/positrons in the plasma’s co-moving frame, $\xi$ is the fraction of the plasma energy radiated in radio, $F$ is the radio flux density, and $D$ is the distance to the source. Comparing Equation (6) with Equation (2), we find that this requirement can be reasonably satisfied.

The coherent energy loss rate per electron is $P_c \gamma^3 \sim 7.7 \times 10^6 n_e \gamma^3 r_{emi,10}^{-1} \nu_{FRB,9}^2 10^{-3} \ erg \ s^{-1}$. Meanwhile, electrons will be accelerated by the rotating SS at the rate of $2\pi B(r_{emi}) \gamma r_{emi} P_0 \sim 3 \times 10^3 r_{emi,10}^2 B_{11,14} P_0^{-1} \ erg \ s^{-1}$ (Kashiyama et al. 2013), where $B(r_{emi})$ is the magnetic field strength at $r_{emi}$. The maximum Lorentz factor during coherent emission thus can be obtained, $\gamma_{max} \sim 2.6 \times 10^6 n_e \gamma_{emi,10} r_{emi,10}^{-1} B_{5,14} P_0$, which is large enough for FRBs.

4. Counterparts in Other Wavebands

No counterparts of FRBs have been discovered in other wavebands yet, except for the repeating FRB, FRB 121102 (Michilli et al. 2018). The collapse of a SS crust might result in electromagnetic radiation besides radio (Petroff et al. 2015). It’s intriguing to check whether the emissions in other wavebands is strong or not in our model.

There are mainly two types of emission from a bare SS surface, thermal radiation and $e^+e^-$ pair emission (Usov 2001a). The plasma frequency for an SQM object can be written as

$$\omega_p = \left( \frac{8 \pi e^2 n_b}{\alpha} \right)^{1/2},$$  
(7)

where $n_b$ is the baryon number density of SQM and $\mu \approx \hbar c (\pi^2 n_b)^{1/3}$ (Alcock et al. 1986). According to the plasma dispersion relationship, propagating modes for electromagnetic waves with $\omega < \omega_p$ do not exist. For typical SQM $n_b \simeq (1.5 \sim 2) n_i$, where $n_i \approx 1.7 \times 10^{38} \ cm^{-3}$ is normal nuclear matter density, we expect $\hbar \omega_p \simeq 20 \sim 25 \ MeV$. Thus SSs with a surface temperature $T_S \lesssim 2 \times 10^{10} \ K$ are very poor radiators for blackbody radiation. However, the thermal radiation luminosity increases sharply and becomes the chief emission form if $T_S \gtrsim 5 \times 10^{10} \ K$ (Haensel & Zdunik 1991; Alcock et al. 1986).

The energy flux per unit surface in thermal photons is $F_{eq} = \frac{h}{\pi} \int_0^{\infty} \frac{\omega (\omega^2 - \omega_p^2)(g(\omega))}{\exp(h\omega/k_{B}T_S) - 1} D(\omega, \theta) \sin \theta \cos \theta d\theta$, where $g(\omega) = \frac{1}{2\pi^2} \Gamma(\gamma/2) D(\omega, \theta) \sin \theta \cos \theta d\theta$, $\nu_B$ is the Boltzmann constant, $D(\omega, \theta)$ is the coefficient of radiation transmission from SQM to the vacuum, $D = 1 - (R_+ + R_-)$, and $R_\perp = \sin^2 (\theta - \theta_0)/\sin^2 (\theta + \theta_0)$, $R_\parallel = \tan^2 (\theta - \theta_0)/\tan^2 (\theta + \theta_0)$, $\theta_0 = \arcsin(\sin \frac{1}{\sqrt{1 - (\omega_p/\omega)^2}})$ (Usov 2001a).

As pointed out by Usov (1999b), the Coulomb barrier outside the SS surface is a very powerful source of $e^+e^-$ pairs, where the electronic field, $\sim 5 \times 10^{17} \ V \ cm^{-1}$, is tens of times higher than the critical vacuum polarization field $E_{cr} = mc^2/e\hbar \approx 1.3 \times 10^{17} \ V \ cm^{-1}$. $e^+e^-$ pairs are created with the mean particle energy of $\varepsilon_{\pm} \approx m_e c^2 + KT_S$. The flux of pairs is

$$f_{\pm} \approx 10^{39.2} \left( \frac{T_S}{10^9 \ K} \right) \exp \left( -\frac{11.9 \times 10^9 \ K}{T_S} \right) J(\xi) \ \text{s}^{-1},$$  
(8)

where

$$J(\xi) = \frac{\xi^3 \ln(1 + 2\xi^{-1}) + \pi^5}{5} \frac{\xi^4}{(1 + 0.0745)^3} + \frac{\pi^5}{6} (13.9 + \xi)^3$$  
(9)

and $\xi \simeq (2 \times 10^{10} \ K)/T_S$ (Usov 2001a). The energy flux per unit surface in $e^+e^-$ pairs is $F_{\pm} \approx \varepsilon_{\pm} f_{\pm}$. 
Since the radiation features of a bare SS are determined by $T_S$, it is crucial to study the distribution of temperature and how it evolves. The heat transfer equation for SSs under the plane-parallel approximation is (Iwamoto 1982; Heiselberg & Pethick 1993; Usov 2001b)

$$\frac{C_q}{\partial T}{\partial t} = \frac{\partial}{\partial x}\left(K_c \frac{\partial T}{\partial x}\right) - \varepsilon, \quad (10)$$

where $C_q \approx 2.5 \times 10^{20} (n_b/n_0)^{2/3} T_9$ erg cm$^{-3}$ K$^{-1}$ is the specific heat for SQM per unit volume, $K_c \approx 6 \times 10^{30} \alpha_e^{-1} (n_b/n_0)^{2/3}$ erg cm$^{-3}$ K$^{-1}$ is the thermal conductivity, $\varepsilon \approx 2.2 \times 10^{26} \alpha_n^2 T_9^4$ erg cm$^{-3}$ s$^{-1}$ is the neutrino emissivity, $n_0 \approx 2n_0$ is the SQM baryon number density, $\alpha_e = g^2/4\pi \approx 0.1$ is the QCD fine structure constant with $g$ being the quark-gluon coupling constant, and $Y_e = n_e/n_B \approx 10^{-4}$ is the number ratio between electrons and baryons. Due to thermal conductivity, the heat flux is

$$q = -K_c dT/dx. \quad (11)$$

The boundary condition is $q \approx -F_\pm - F_{eq}$.

We now proceed to calculate the evolution of the radiation luminosity after the collapse of an SS crust. Our calculations are done for an SS with an initial surface and crust temperature of $T_{S0} \approx 3 \times 10^7$ K (Pizzochero 1991; Usov 1997), and a crust mass of $M_c \approx 3 \times 10^{-6} M_\odot$. There are two main kinds of energy released during the collapse: (1) gravitational energy of the crust ($\sim 0.002 M_c c^2$) due to its rapid movement to the SQM surface, and (2) deconfinement energy ($\sim 0.01$–$0.03 M_c c^2$) due to the conversion of the crust material to SQM (Cheng & Dai 1996). As a sum, we can take the typical total energy from these two reservoirs as roughly $Q \sim 0.02 M_c c^2 \approx 1 \times 10^{48}$ erg for $M_c \approx 3 \times 10^{-6} M_\odot$. After the collapse, the crust is transferred to SQM with a density of $\sim 5 \times 10^{14} g$ cm$^{-3}$ (Alcock et al. 1986) and the interaction is restricted in a thickness of $t_h \approx 1$ cm. Assuming that the actual combustion mode is detonation (Olinto 1987; Horvath & Benvenuto 1988; Heiselberg et al. 1991), the timescale of conversion is rather small and the interior SS temperature can be treated as uninfluenced while the surface layer is hot and isothermal. With $C_q \approx 4 \times 10^{11}$ erg cm$^{-3}$ K$^{-1}$, the temperature of the heated layer can be estimated as $T_S^* \approx 2 \times 10^{13}$ K. The initial temperature distribution can be expressed as

$$T(t = 0, x) \approx \begin{cases} T_{S\star}, & 0 < x < t_h \\ T_{S0}, & x \geq t_h \end{cases} \quad (12)$$

where $t$ is the time after the collapse and $x$ is a space coordinate measuring the depth below the SS surface.

Combining these postulations and approximations, we have performed numerical calculations to solve the cooling process. Figure 2 shows the evolution of the surface temperature as a function of time. The total luminosity including both photons and electrons/positrons, $L = L_{eq} + L_{\pm} = 4\pi r^2 (F_{eq} + F_{\pm})$, has also been calculated and the result is shown in Figure 3. We should note that when $L_\pm$ is very high, most of the pairs will annihilate into photons near the SS surface due to the high pair density. Hence the emerging emission consists mostly of photons. The photon spectrum is roughly a blackbody with a high energy ($\gtrsim 100$ keV) tail (Aksenov et al. 2003) since the outgoing pairs and photons are very likely in thermal equilibrium (Iwamoto & Takahara 2002). Note that bare SSs are bounded by strong interactions rather than the gravity, so the luminosity can safely exceed the Eddington limit (Alcock et al. 1986).

The luminosity distance is $d_L \approx$ a few hundred Mpc at redshift of $z \approx 0.5$–1. The hot bare SS radiates at an extremely high luminosity just after the collapse and it cools down rapidly. The surface layer will become cold soon, making the radiation power decrease quickly. Since $T_S \sim$ a few $10^9$ K during the emission, the typical photon energy can be estimated as $\sim 100$ keV for a blackbody spectrum. According to our calculations, the emitted energies are roughly $9.6 \times 10^{44}$ erg, $1.6 \times 10^{45}$ erg, and $1.7 \times 10^{45}$ erg in the first 10 ms, 100 ms, and 1000 ms after the collapse, respectively. If averaged over 100 ms, the typical luminosity is $L \sim 2 \times 10^{46}$ erg s$^{-1}$. Noting that the expected radiation flux is then $10^{-8} L_{46} (d_L/100$ Mpc)$^{-2}$ erg cm$^{-2}$ s$^{-1}$, this luminosity will be too low to trigger the Swift Burst Alert Telescope whose threshold is $\sim 10^{-7}$ erg cm$^{-2}$ s$^{-1}$ in 15–150 keV (Lorén-Aguilar et al. 2009). Also, the derived peak photon flux of $\sim 0.1 (d_L/100$ Mpc)$^{-2}$ photons cm$^{-2}$ s$^{-1}$ is lower than the

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Figure 2. Evolution of the surface temperature of a bare strange star, $T_S$, after the crust collapses.

Figure 3. Total luminosity, $L$, as a function of time after the crust collapses. In this figure, $L = L_{eq} + L_{\pm}$, where $L_{eq}$ and $L_{\pm}$ are the luminosities in thermal photons and $e^+e^-$ pairs, respectively.
trigger threshold of the Fermi Gamma-ray Burst Monitor 0.74 photons cm$^{-2}$ s$^{-1}$ (Meegan et al. 2009). Additionally, the extremely small radiation timescale of about 0.1 s makes the observation even more difficult. Although convection process may reheat the SS surface and increase its emission slightly (Usov 1998a), the light curve within the initial 0.1 s will not be influenced significantly and no afterglow could be detected in X-rays or gamma-rays by current high energy detectors.

5. Summary and Discussion

In summary, we propose that FRBs may be generated from the collapses of SS crusts. During the collapse, a fraction of magnetic energy is transferred to accelerate electrons and positions in the polar cap region to relativistic velocities. The accelerated electrons expand along magnetic field lines to form a shell. At the radius of $r_{\text{emi}} \sim 10^{10}$ cm, coherent emission in radio bands will be produced, giving birth to the observed FRB. At the same time, the emission in X-ray and $\gamma$-ray bands is too faint to be detected by current detectors.

It is argued that the magnetic field of an SS is influenced by spatial temperature variations (Xu & Busse 2001). The SQM surface will be heated up to a high temperature of $\sim 10^{10}$ K and cool down drastically via the production of electron/positron pairs and neutrinos (Usov 1998b, 2001b). A thin layer surface will be even colder than inside during the cooling process. A cold dense layer forms and then the temperature distribution seems unstable with respect to convective disturbances (Usov 1998a). The small-scale buoyant convection induced by the temperature gradient may amplify the magnetic field due to the interaction with differential rotation through fast dynamos processes (Xu & Busse 2001). The amplified magnetic field lines may emerge from the stellar interior, where the fields can be several orders of magnitude larger (Shibata & Magara 2011). A fraction of heat may be transferred to magnetic field energy and then also contribute to the emission. In other words, the SS surface magnetic field strength may increase significantly after the collapse, thus SSs originally having a weaker magnetic field may also produce FRBs. The effect of convection needs further investigation in the future.

It is an interesting question whether FRBs generated from SS crust collapses can repeat or not. The crux is to determine whether a bare SS can reconstruct its crust through accretion or other ways. In fact, most SS formation models are explosive (Haensel et al. 1991; Xu et al. 2001; Pagliara et al. 2013), thus a newly born SS is bare and how a normal matter crust forms still needs to be investigated. The free-fall kinetic energy for a proton onto the SS surface is $E_p \approx \frac{GM_m}{R_{\text{argon}}^{1.2}} \frac{GM}{R_c^{2}} \approx 138 \, M_\odot R_\odot^{-1}$ (1 + 0.2 $M_\odot R_\odot^{-1}$) MeV (Olinto 1987), where $M_m$ is the proton mass. If the SS is non-rotational and not magnetized, and the accretion is isotropic, then the free-fall energy will be high enough to overcome the electrostatic potential barrier of $eV \approx 20$ MeV outside the SQM surface (Alcock et al. 1986). So, SS crusts cannot be built in such scenarios. However, the falling material will have a non-zero angular momentum when approaching the surface of rotational magnetized SSs due to the magnetic freezing. These matter will finally hit the SS surface obliquely at a typical incidence angle with cotangent $\sim 0.05$ (Kluzniak & Wagoner 1985; Olinto 1987). Hence the accreted matter has a longer interaction with the electric field and the radial velocity will be reduced by friction and radiation. Considering this effect, an envelope built from the accreted material will cover the whole SQM core gradually. If the accretion continues, the crust will finally reach the critical mass via accretion and then collapse.

In this case, the number density of particles near the SS may be rather high due to the existence of the accretion flow. Thus the interaction between the expanding shell and the accretion flow needs to be considered. Assuming that the accretion is isotropic, the number density of the accretion flow at the emission radius is $n_a \approx 2M/\{4\pi r_{\text{emi}} m_p(2GM/r_{\text{emi}})^{1/2}\} \approx 1.6 \times 10^{3} (M/10^{-15} M_\odot \text{ yr}^{-1}) (M/1.4 M_\odot)^{2/3} \, \text{cm}^{-3}$, where $M$ is the accretion rate. If the accretion rate is high, the outgoing shell may be disrupted by the falling gas and then no FRB could be generated. However, as long as $M$ is less than a critical value of $M_c \approx 10^{-15} M_\odot \text{ yr}^{-1}$, we notice that $n_a$ will be much smaller than the number density of electrons in the expanding shell ($n_e$). Then the influence of the accretion flow can be negligible.

For SSs formed in explosive events (Haensel et al. 1991; Xu et al. 2001; Pagliara et al. 2013), the compact stars may receive a “kick” if these drastic events are asymmetric (Nordhaus et al. 2012; Bray & Eldridge 2016). The typical kick velocity, $v$, ranges from 200 to 400 km s$^{-1}$, with the highest value even in excess of 1000 km s$^{-1}$ (Lyne & Lorimer 1994; Faucher-Giguère & Kaspi 2006). SSs with a kick velocity would accelerate ambient matter at a rate of $\dot{M} \approx 2\pi \alpha (GM)^{2} v^{2} \rho \sim 2 \times 10^{-15} M_\odot^{2} v^{2} \rho \sim 3 \dot{M}_0 \text{ yr}^{-1}$, where $\alpha \sim 1.25$ is a numerical constant and $\rho$ is the ambient matter density (Bondi 1952). In this case, the typical $\dot{M}$ would not exceed $M_c$ significantly. Therefore, the expanding electron/positron shell can expand without being destroyed by the accretion flow in our scenario. For this kind of accretion, the reconstruction timescale of the crust can be roughly estimated as $\tau_{\text{rec}} \approx M_\odot/\dot{M}_0 \sim 10^{5}$ year.

Owing to this long reconstruction timescale, multiple FRB events from the same source seem not likely to happen in our scenario. Our model thus is more suitable for explaining the non-repeating FRBs, and the repeating FRB 121102 may be produced via other mechanisms (Palaniswamy et al. 2018). However, we should also note that during the collapse process, if only a small portion (in the polar cap region) of the crust falls onto the SQM core while the other portion of the crust remains stable, then the rebuilt timescale for the crust can be markedly reduced and repeating FRBs would still be possible. Further detailed studies on the crust collapse thus still need to be conducted.

The event rate of FRBs is as high as $2 \times 10^{3}$ sky$^{-1}$ day$^{-1}$ (Li et al. 2017; Bhandari et al. 2018). There are roughly $10^{9}$ galaxies within the redshift range of $z \leq 1$ and the total number of pulsars per galaxy is $\sim 10^{6}$ on average (Timmes et al. 1996). The event rate for SS crust collapse can be estimated as $\sim 2 \times 10^{3} (\tau_{\text{rec}}/10^{7} \text{ yr})^{-1} f_{\text{s}}$, where $f_{\text{s}}$ is the beaming factor of radio emission and $f_{\text{s}}$ is the number ratio of SSs with compatible conditions to generate FRBs over pulsars. Still, we would like to remind that FRBs may be of multiple origins and our model may only contribute a portion of them. Further observations and larger samples in the future would help to solve the enigma finally.

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References
Abe, K., Bacon, T. C., Ballam, J., et al. 1984, PhRvD, 30, 1
Akseenov, A. G., Milgrom, M., & Usov, V. V. 2003, MNRAS, 343, L69
Alcock, C., Farhi, E., & Olinto, A. 1986, ApJ, 310, 261
Barrau, A., Rovelli, C., & Vidotto, F. 2014, PhRvD, 90, 127503
Benford, G., & Buschauer, R. 1977, MNRAS, 179, 189
Bhandari, S., Keane, E. F., Barr, E. D., et al. 2018, MNRAS, 475, 1427
Bray, J. C., & Eldridge, J. J. 2016, MNRAS, 461, 3747
Cheng, K. S., & Dai, Z. G. 1996, PhRvL, 77, 1210
Chmaj, T., Haensel, P., & Slomiński, W. 1991, NuPhS, 24, 40
Cordes, J. M., & Lazio, T. J. W. 2002, arXiv:astro-ph/0207156
Dai, Z. G., Wang, J. S., Wu, X. F., & Huang, Y. F. 2016, ApJ, 829, 27
Falcke, H., & Rezzolla, L. 2014, A&A, 562, A137
Farhi, E., & Jaffe, R. L. 1984, PhRvD, 30, 2379
Faucher-Giguère, C.-A., & Kaspi, V. M. 2006, ApJ, 643, 332
Geng, J. J., & Huang, Y. F. 2015, ApJ, 809, 24
Geng, J. J., Huang, Y. F., & Lu, T. 2015, ApJ, 804, 21
Gu, W.-M., Dong, Y.-Z., Liu, T., Ma, R., & Wang, J. 2016, ApJL, 823, L28
Haensel, P. 1991, NuPhS, 24, 23
Haensel, P., Paczynski, B., & Amsterdamski, P. 1991, ApJ, 375, 209
Heiselberg, H., Baym, G., & Pethick, C. J. 1991, NuPhS, 24, 144
Heiselberg, H., & Pethick, C. J. 1993, PhRvD, 48, 2916
Horvath, J. E., & Benvenuto, O. G. 1988, PhLB, 213, 516
Huang, Y. F., & Lu, T. 1997, A&A, 325, 189
Iwamoto, N. 2018, ArPhy, 141, 1
Iwamoto, S., & Takahara, F. 2002, ApJ, 565, 163
Kashiyama, K., Ioka, K., & Mészáros, P. 2013, ApJL, 776, L39
Katz, J. I. 2014, PhRvD, 89, 103009
Keane, E. F., Stappers, B. W., Kramer, M., & Lyne, A. G. 2012, MNRAS, 425, L71
Kettner, C., Weber, F., Weigel, M. K., & Glendenning, N. K. 1995, PhRvD, 51, 1440
Kluzniak, W., & Wagoner, R. V. 1985, ApJ, 297, 548
Kondratyuk, L. A., Krivoruchenko, M. I., & Martemyanov, B. V. 1990, SvAL, 16, 410
Kouveliotou, C., Dieters, S., Strohmayer, T., et al. 1998, Natur, 393, 235
Kulkarni, S. R., Ofek, E. O., Neill, J. D., Zheng, Z., & Juric, M. 2014, ApJ, 797, 70
Kumar, P., Lu, W., & Bhattacharya, M. 2017, MNRAS, 468, 2726
Lai, D., & Shapiro, S. L. 1991, ApJ, 383, 745
Lai, Y.-Y., Yu, Y.-W., Zhou, E.-P., Li, Y.-Y., & Xu, R.-X. 2018, RAA, 18, 024
Lattimer, J. M., & Prakash, M. 2007, PhR, 442, 109
Li, L.-B., Huang, Y.-F., Zhang, Z.-B., Li, D., & Li, B. 2017, RAA, 17, 6
Liu, X. 2017, arXiv:1712.03509
Lorén-Aguilar, P., Isern, J., & García-Berro, E. 2009, A&A, 500, 1193
Lorimer, D. R., Bailes, M., McLaughlin, M. A., Narkevic, D. J., & Crawford, F. 2007, Sci, 318, 777
Luan, J., & Goldreich, P. 2014, ApJL, 785, L26
Lyne, A. G., & Lorimer, D. R. 1994, Natur, 369, 127
Lyubarsky, Y. 2008, ApJ, 682, 1443
Meegan, C., Chincarini, G., Bhat, P. N., et al. 2009, ApJ, 702, 791
Mészáros, P. 1992, MNRAS, 258, 41P
Meisel, D. B. 1971, PhRvD, 13, 56
Michilli, D., Seymour, A., Hessels, J. W. T., et al. 2018, Natur, 553, 182
Miralda-Escude, J., Paczynski, B., & Haensel, P. 1990, ApJ, 362, 572
Nordhaus, J., Brandt, T. D., Burrows, A., & Almgren, A. 2012, MNRAS, 423, 1805
Olinto, A. V. 1987, PhLB, 192, 71
Pagliara, G., Herzog, M., & Röpke, F. K. 2013, PhRvD, 87, 103007
Palaniswamy, D., Li, Y., & Zhang, B. 2018, ApJL, 854, L12
Petroff, E., Bailes, M., Barr, E. D., et al. 2015, MNRAS, 447, 246
Pizzochero, P. M. 1991, PhRvL, 66, 2425
Rees, M. J., & Mészáros, P. 1992, MNRAS, 258, 41P
Romero, G. E., de Valle, M. V., & Vieyro, F. L. 2016, PhRvD, 93, 023001
Ruderman, M. A., & Sutherland, P. G. 1975, ApJ, 196, 51
Scholz, P., Spitler, L. G., Hessels, J. W. T., et al. 2016, ApJ, 833, 177
Shibata, K., & Magara, T. 2011, LRSP, 8, 6
Stejner, M., & Madsen, J. 2005, PhRvD, 72, 123005
Thompson, C., & Duncan, R. C. 1995, MNRAS, 275, 255
Thorne, D., Stappers, B., & Rezzolla, L. 2013, Sci, 341, 53
Timmes, F. X., Woosley, S. E., & Weaver, T. A. 1996, ApJ, 457, 834
Totani, T. 2013, PASJ, 65, L12
Usov, V. V. 1998, PhRvL, 90, 103009
Usov, V. V. 1998b, PhRvL, 80, 230
Usov, V. V. 2001a, ApJ, 550, L179
Usov, V. V. 2001b, PhRvL, 87, 021101
Wu, R. X., & Busse, F. H. 2001, A&A, 371, 963
Witten, E. 1984, PhRvD, 30, 272
Xu, R. X., & Busse, F. H. 2001, A&A, 371, 963
Xu, R. X., Zhang, B., & Qiao, G. J. 2001, ApJ, 15, 101
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