从奇异星到致密星爆发现象

——科研经验分享

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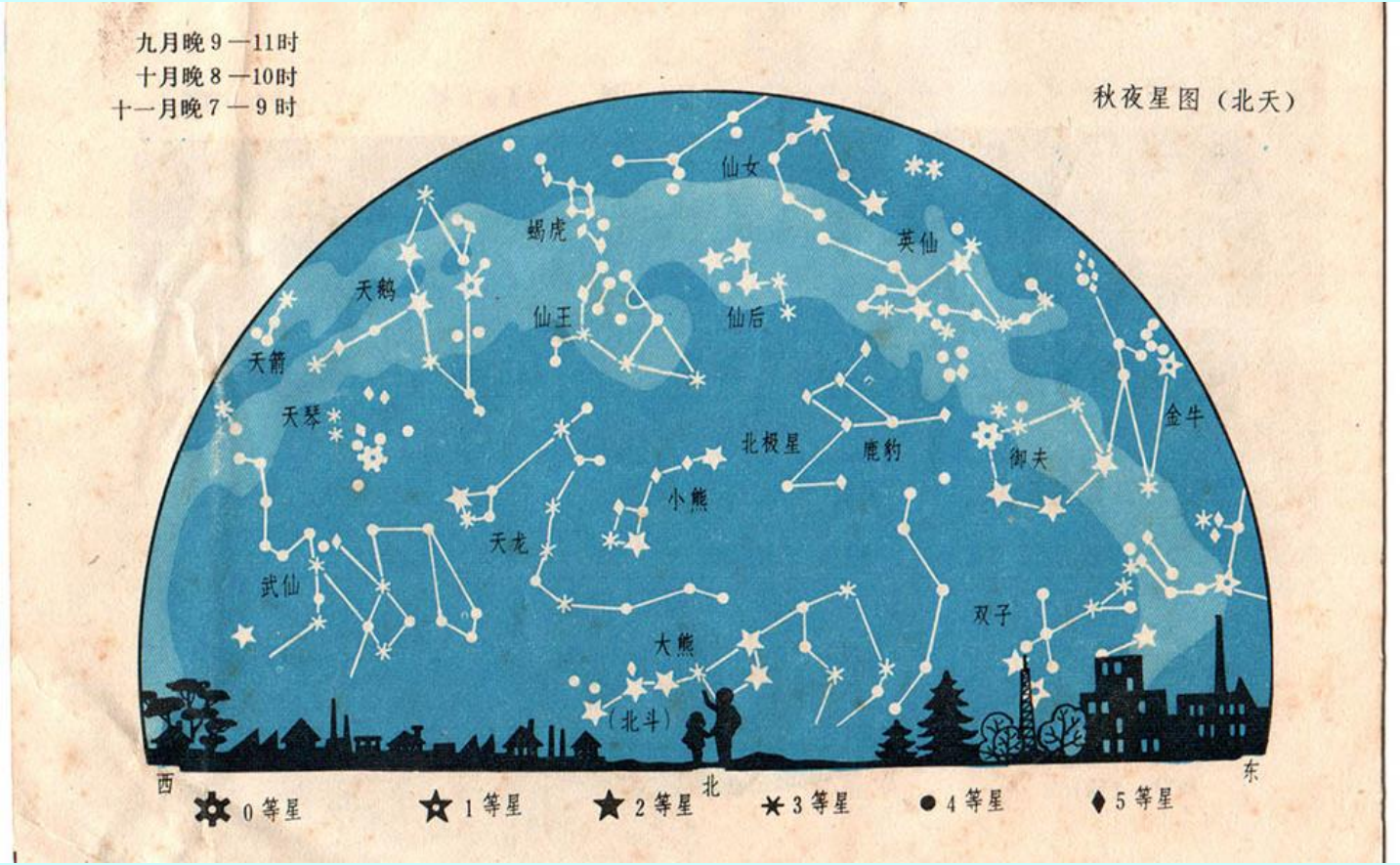
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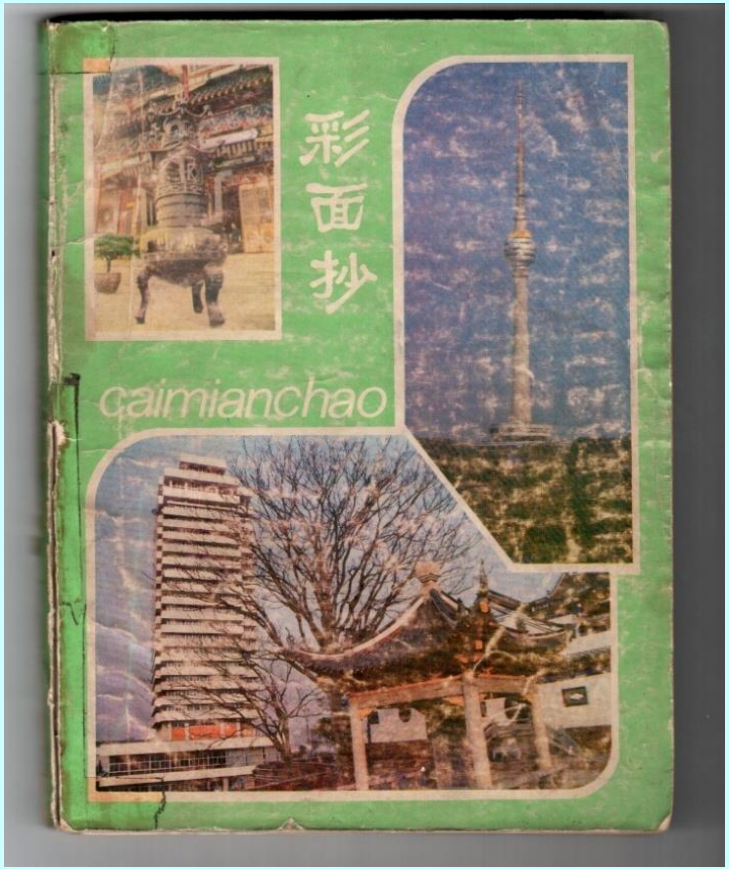
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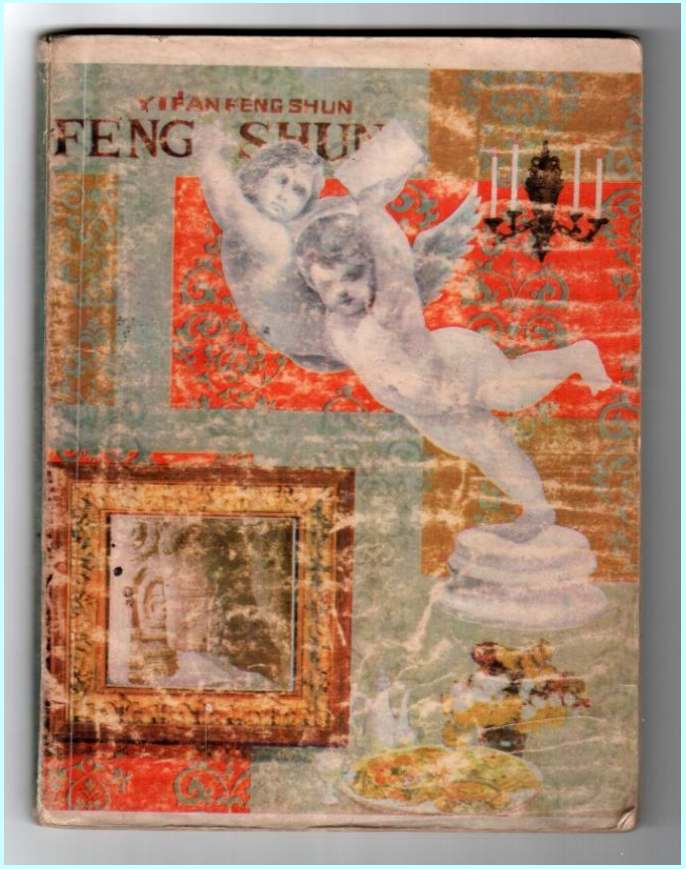
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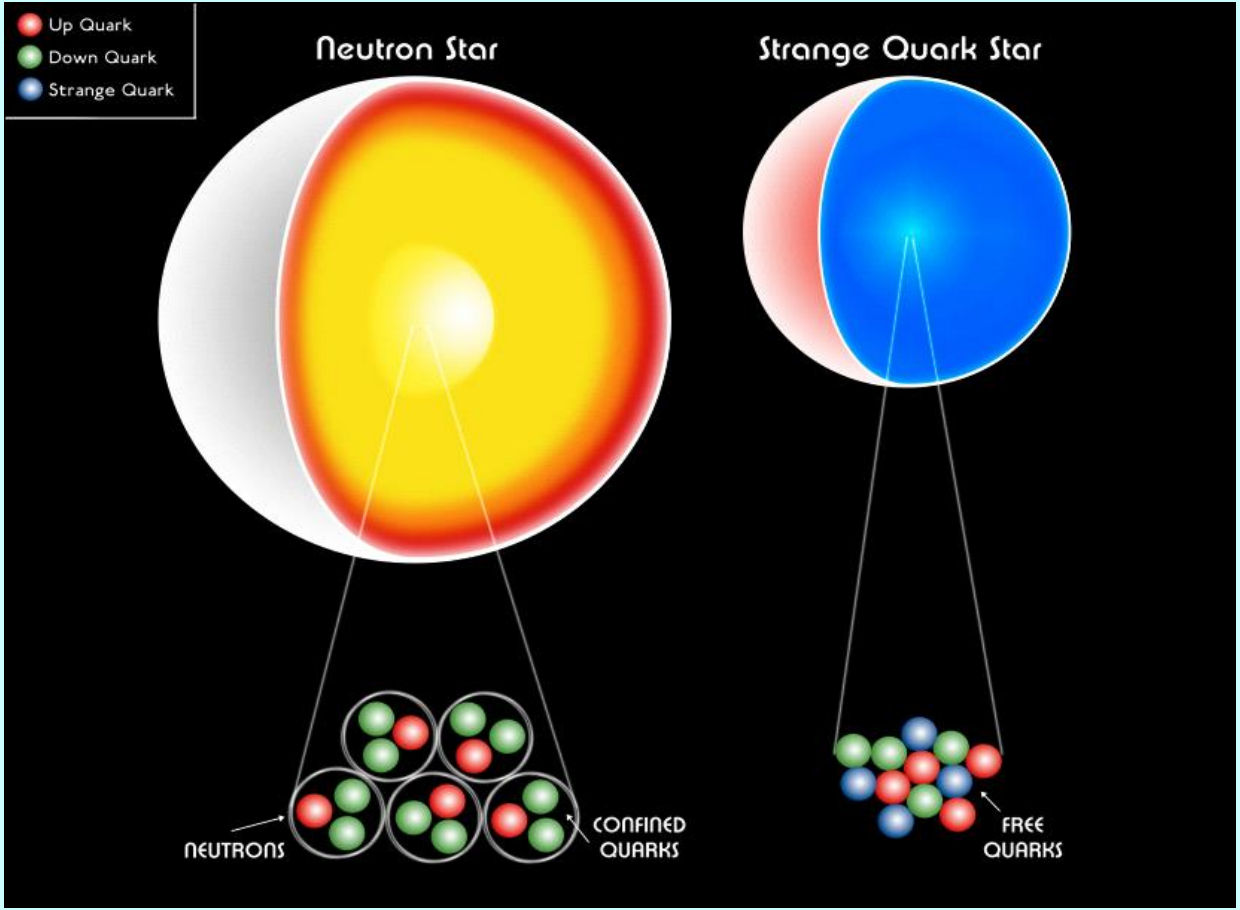

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中子星 vs. 奇异夸克星



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STRANGE STARS[1]

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ABSTRACT

Strange matter, a form of quark matter that is postulated to be absolutely stable, may be the true ground state of the hadrons. If this hypothesis is correct, neutron stars may convert to "strange stars." The mass-radius relation for strange stars is very different from that of neutron stars; there is no minimum mass, and for mass $\lesssim 1\ M_\odot$, $M \propto R^3$. For masses between $1\ M_\odot$ and $2\ M_\odot$, the radii of strange stars are ~ 10 km, as for neutron stars. Strange stars may have an exposed quark surface, which is capable of radiating at rates greatly exceeding the Eddington limit, but has a low emissivity for X-ray photons. The stars may have a thin crust with the same composition as the pre–neutron drip outer layer of a conventional neutron star crust. Strange stars cool efficiently via neutrino emission. It is not clear whether or not all neutron stars must convert into strange stars, but this seems the most likely conclusion. There may be no neutron stars, only strange stars.

Subject headings: elementary particles — neutrinos — stars: interiors — stars: neutron

I. INTRODUCTION

The true ground state of the hadrons may be "strange matter," not ^{56}Fe (Witten 1984). Strange matter is a bulk quark matter phase consisting of roughly equal numbers of up, down, and strange quarks plus a smaller number of electrons (to guarantee charge neutrality) which is *conjectured* to have a lower energy per baryon than ordinary nuclei. Strange matter is, by hypothesis, absolutely stable; has a density comparable to that of atomic nuclei; and can exist in lumps ranging in size from a few fermis up to "strange stars" of radius ~ 10 km. We discuss the properties of strange stars in this paper.

evaporates completely as the universe cools to ~ 10 MeV (Alcock and Farhi 1985).

Second, Witten pointed out that neutron stars will probably convert to strange matter. The reasons for this are discussed below, but its importance for astrophysics is clear: if the strange matter hypothesis is correct, there may be no *neutron* stars; instead, there may be *strange* stars. Accordingly, in the discussion which follows, we explicitly draw attention to the differences and similarities between neutron stars and strange stars. The possibility that neutron star interiors could be quark matter has also been discussed, and such objects are known as

内部密度轮廓

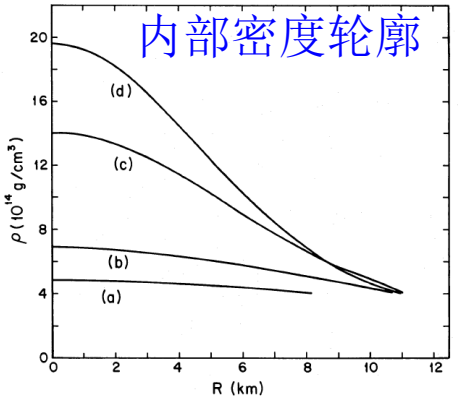


FIG. 1.—Density (ρ) vs. radius (r) for strange stars of mass (a) 0.53 $M_\odot$, (b) 1.4 $M_\odot$, (c) 1.95 $M_\odot$, and (d) 1.99 $M_\odot$.

M-ρ_c

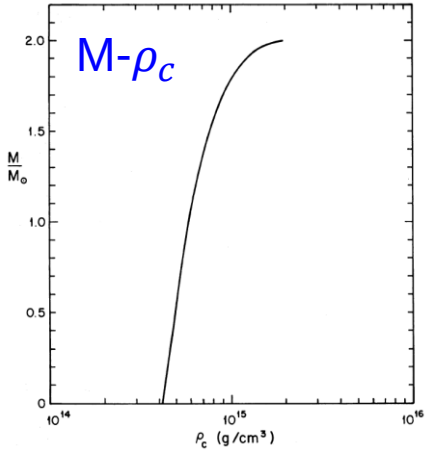


FIG. 2.—Total mass (M) vs. central density (ρ_c) for stable strange stars

M - R

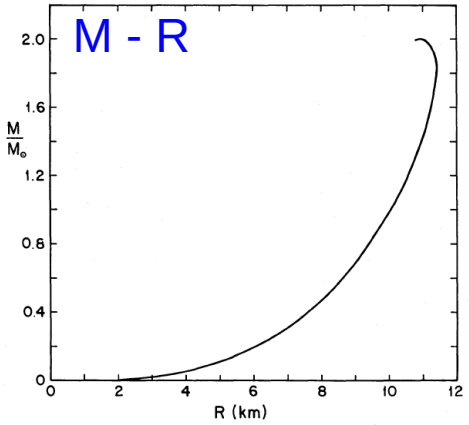


FIG. 3.—Total mass (M) vs. radius (R) for stable strange stars

I - M

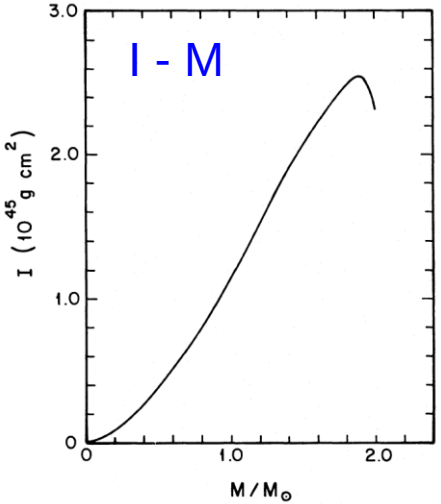


FIG. 5.—Moment of inertia (I) vs. total mass (M) for strange stars

奇异核表面电场

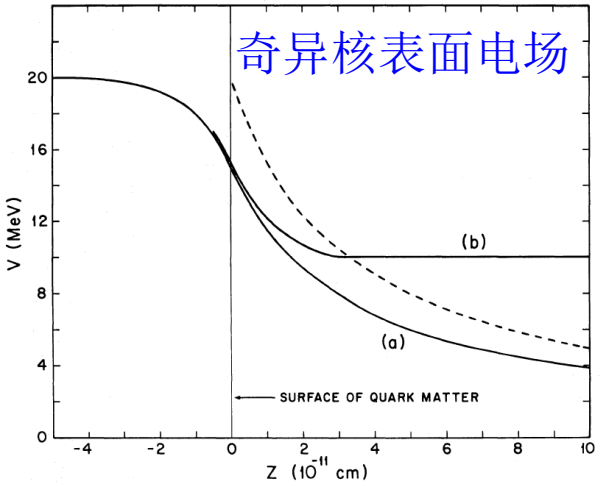
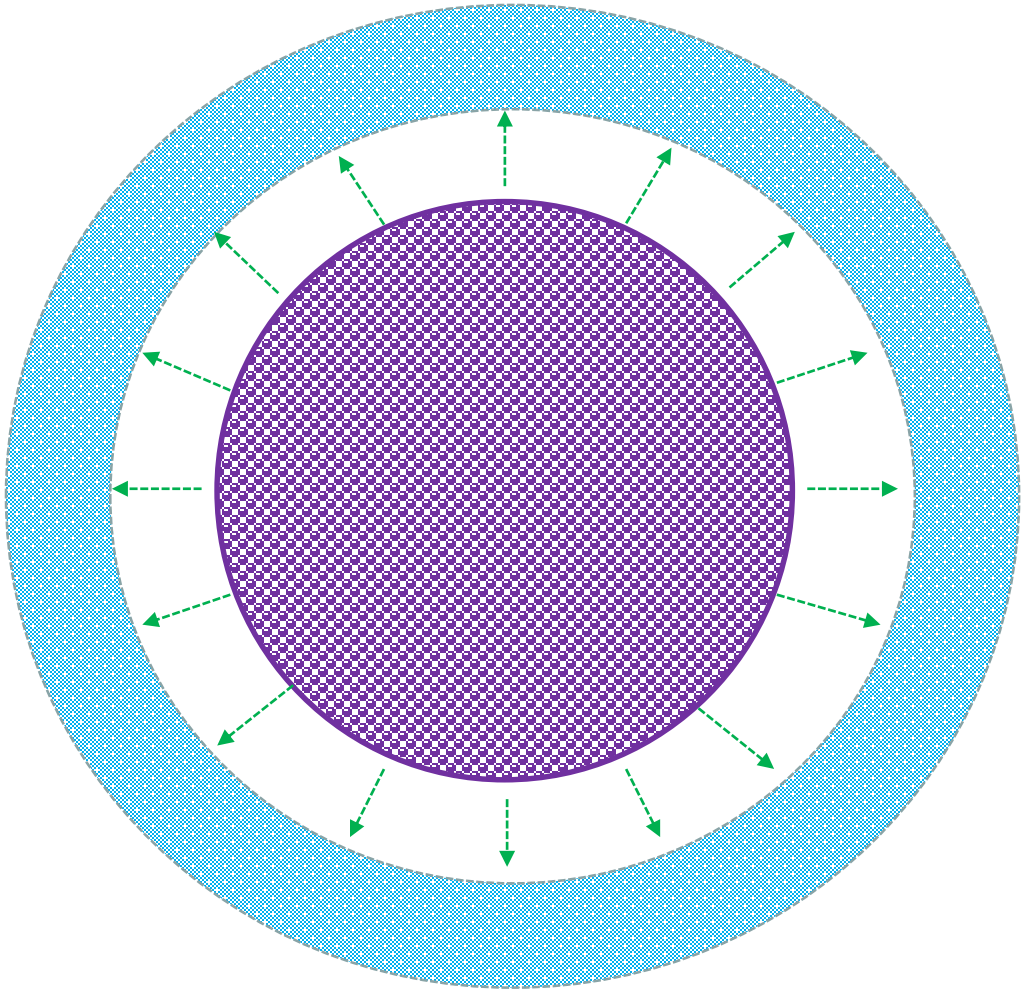
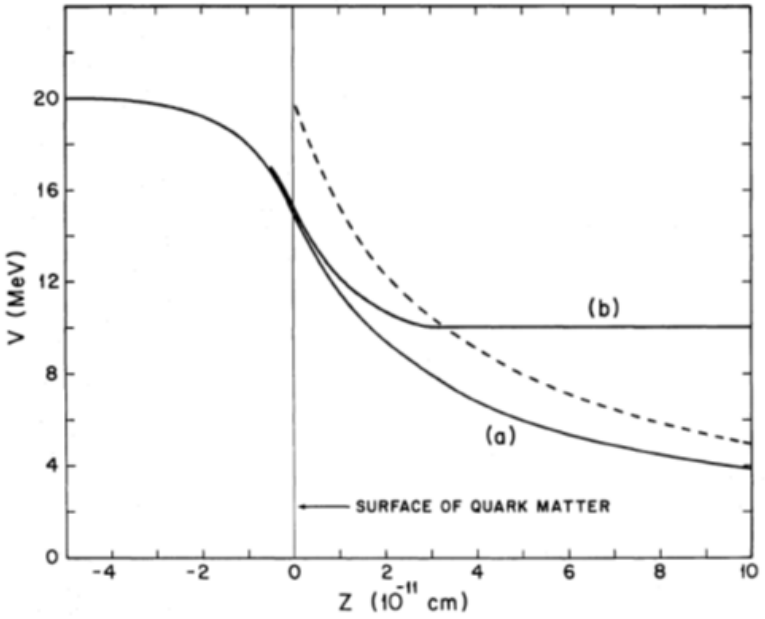


FIG. 6.—Electrostatic potential (V) vs. height (z) near the surface of the strange star for two cases with $V_q = 20$ MeV: (a) $V_c = 0$; (b) $V_c = 10$ MeV. The dashed line shows V_c vs. gap width (z_G). The vertical line at $z = 0$ represents the surface of the quark matter.




V (MeV)
20
16
12
8
4
0
-4
-2
0
2
4
6
8
10
Z (10^-11 cm)
(b)
(a)
SURFACE OF QUARK MATTER

CHIN.PHYS.LETT. Vol.14, No. 4 (1997) 314

On the Crust of a Strange Star *

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The strong electric field near the surface of a strange star is discussed, and its supporting interaction to the normal nuclear crust is calculated by using a self-consistent method. It is proved that the crust would begin not to be supported by this electric field at a quite low density at the crust base, much lower than the previously assumed neutron drip density.

PACS: 97.20.Rp, 97.60.Jd, 97.60.Sm

Though some ideas of quark stars were proposed as early as in 1969, no great achievements had been made until 1984. Witten[1] and Farhi *et al.*[2] first predicted the existence of strange stars by proving their energy to be lower than that of neutron stars for quantum chromodynamical parameters within rather wide range, hence, the stable ones should be strange stars rather than neutron stars. At the same time, Wang *et al.*[3] first studied the dynamical behavior of a strange star and found its most important property: the very high bulk viscosity in strange matter. Two years later, Alcock *et al.*[4] even said that there may be no neutron stars, only strange stars.

In the same paper, Alcock *et al.*[4] first studied the structure of strange stars with crusts. As strange matter is bound mainly by quark confinement, not by gravitation, a strange star may have a surface where the density drops from bulk quark matter density abruptly to zero. The thickness of the quark matter surface will be order of magnitude about 1 fm, determined by the typical strong interaction length scale. As the s(strange) quark is more massive than d(down) and u(up) quarks, the s quark matter in equilibrium is not charge neutral, there must be a small part of electrons in order to keep the strange matter wholly neutral. However, the electrons are held to the quark matter electromagnetically not strongly, they will extend farther than the quark matter surface, then a strong outward-directed electric field ($\sim 10^{17}\,\mathrm{Vcm^{-1}}$) will exist beyond the quark surface, within the extension region of electrons. This electric field could effectively prevent normal nuclear matter from falling into the strange quark matter. Thus, a strange star will usually be composed of a strange core, a crust of normal nuclear matter and a gap of strong electric field between them. The gap is usually only thick of several hundred fm.

The crust could not be very thick. For a very thick crust, the density at its base will be very high. As that density reaches the neutron drip point ($\rho_{\mathrm{drip}} = 4.3\times10^{11}\,\mathrm{g\,cm^{-3}}$), neutrons will begin to drip out from nuclei. Because neutrons are neutral, they will not be prevented by strong electric field and will fall freely into the strange core and then be deconfined to be strange matter. Alcock *et al.*[4] argued that the force of strong electric field on a single nucleus overwhelms the force of gravity which could be neglected. This leads them to the conclusion that the only limit on the crust mass is determined by the neutron drip. Based on this result, in recent years, Glendenning *et al.* studied the crust structure of strange stars,[5–7] leading to a very interesting picture of the series of strange stars-strange dwarfs.

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However, the crust is only weakly charged (nearly neutral), the electric force on the crust is rather weak, mechanical balance should be held between electric and gravitational forces on the crust, not on a single nucleus. In real calculation, we consider the balance between the electric force and the pressure which is determined by gravitation. Based on this analysis, we recalculate the question in a self-consistent way, and prove that the gap height will begin to be zero (the crust could not be supported) for a density ρ_{bottom} at the base of the crust far lower than the neutron drip density, quite different from results previously known.

In order to calculate the electron distribution over the background of "sandwich" structure (quark matter-gap-nuclei), we use a simple Thomas-Fermi model.[4] The number density n_{e} of electrons is given locally by the electron Fermi momentum p_{f} as follows:

$$n_{\mathrm{e}} = p_{\mathrm{f}}^3/3\pi^2\,, \tag{1}$$

where units of $c = \hbar = 1$ are used. Alcock *et al.*[4] proved, in equilibrium, $p_{\mathrm{f}} = V$, where V/e is the electrostatic potential. The potential energy V obeys the Poisson equation:

$$\frac{\mathrm{d}^2V}{\mathrm{d}z^2} = \begin{cases} 4\alpha(V^3 - V_{\mathrm{q}}^3)/(3\pi)\,, & z \le 0\,, \\ 4\alpha V^3/(3\pi)\,, & 0 \le z \le z_{\mathrm{c}}\,, \\ 4\alpha(V^3 - V_{\mathrm{e}}^3)/(3\pi)\,, & z_{\mathrm{c}} \le z\,, \end{cases} \tag{2}$$

where V is the potential energy in units of MeV, z the height from the surface of the strange star, $V_{\mathrm{q}}^3/(3\pi^2)$ the charge density of the quark matter, $V_{\mathrm{e}}^3/(3\pi^2)$ the charge density of the crust nuclei. Of course, $V_{\mathrm{c}} < V_{\mathrm{q}}$. Denote z_{c} the gap height, namely, the distance between the base of the crust and the surface of the strange core, and α the fine-structure constant.

Related to nuclear matter (typical: $A = 118, Z = 36$) of neutron drip density ρ_{drip}, Alcock *et al.* took $V_{\mathrm{q}} = 20\,\mathrm{MeV}$, $V_{\mathrm{c}} = 10\,\mathrm{MeV}$. They integrated Eq. (2) from a deep place of $z < 0$ outward until $V = V_{\mathrm{c}}$, they denoted this place as the base of the crust, and the corresponding z as z_{c}. As they perfectly neglected the gravitation, they regarded the crust completely neutral, V would be constant within the crust, the 3rd equation of Eq. (2) be in fact unused in their calculating procedure. So, the upper limit of the density at the base of the crust, they found, was just the neutron drip density.

As we have argued above, the calculation of Alcock *et al.* is inconsistent. In our opinion, the crust, especially near its base, is not neutral, V will not be constant there, V_{c} does not denote the potential energy at the base of the crust, but denotes the potential energy at some outer place (denoted as z_{p}, it is actually very small, only about some hundreds of fm) within the crust. Thus, at the base of the crust, $V(z_{\mathrm{c}}) \neq V_{\mathrm{c}}$. Only from z_{p} outward, one could have $V(z \ge z_{\mathrm{p}}) = V_{\mathrm{c}}$. We calculate this question (for given V_{q} and V_{c}) in a self-consistent way as follows: For a given density of $\rho = \rho_{\mathrm{c}}$ near the base of the crust, the corresponding pressure $P = P_{\mathrm{c}}$ can be evaluated according to the equation of state, for which we use that of Baym *et al.*[8] For an arbitrarily chosen z_{c} we integrate Eq. (2) from a deep place of $z < 0$ outward until $V = V_{\mathrm{c}}$. The z corresponding to $V = V_{\mathrm{c}}$ will be denoted as z_{p}. The region from z_{c} to z_{p} is not neutral. The force on this region due to electric field can be calculated, and the equivalent pressure P_{E} resulting from this force can also be calculated. Usually, $P_{\mathrm{E}} \neq P_{\mathrm{c}}$. We adjust z_{c} until $P_{\mathrm{E}} = P_{\mathrm{c}}$, then we obtain the final solution of z_{c}, from which the final z_{p} can be calculated. For different ρ_{c}, we will have different z_{c} and z_{p}, and will have also different crust

mass M_{crust}. An evident difference between our calculating procedure and that of Alcock *et al.* is that the 3rd equation of Eq. (2) is really used by us but not by them.

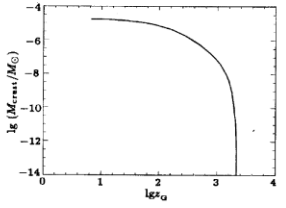


Fig. 1. $M_{\mathrm{crust}} - z_{\mathrm{c}}$ relation. Here unit of z_{c} is fm.

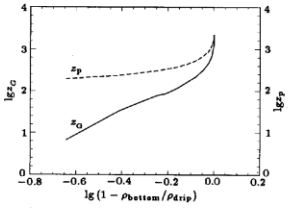


Fig. 2. The z_{c} and z_{p} as functions of ρ_{bottom}. Here units of z_{c} and z_{p} are fm.

The following results are obtained for the typical strange star of mass $M = 1.4\,M_\odot$ ($M_\odot$ is the mass of the sun), and radius $R = 1.097\times10^6$ cm [the bag constant of strange matter is taken to be $B = 57\,\mathrm{MeV\,fm^{-3}}$ (or $B^{1/4} = 145\,\mathrm{MeV}$)]:

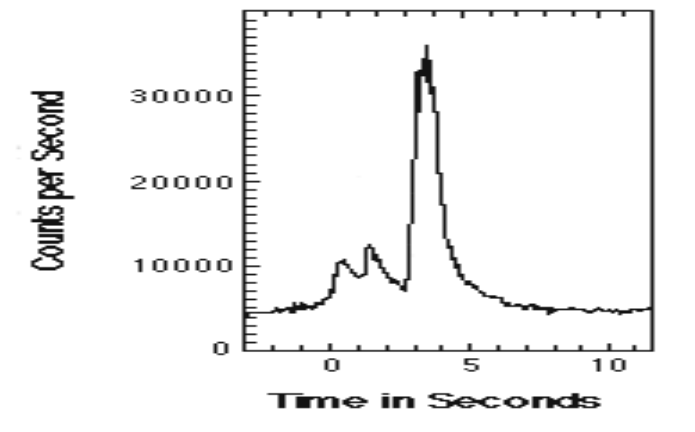
1. As the total crust mass gets bigger and bigger, the gap width z_{c} gets smaller and smaller. For $M_{\mathrm{crust}} = 2\times10^{-8}\,M_\odot$, we have $z_{\mathrm{c}} = 1260$ fm; for $M_{\mathrm{crust}} = 1.0\times10^{-5}\,M_\odot$, we have $z_{\mathrm{c}} = 62$ fm; for $M_{\mathrm{crust}} = 1.8\times10^{-5}\,M_\odot$, we will have a very small width $z_{\mathrm{c}} = 6.6$ fm. The relation between M_{crust} and z_{c} is shown in Fig. 1.

2. In order to prevent the strong interaction between the crust and the strange core, z_{c} should at least be about 200 fm.[4–7] Thus we could conclude that the maximum crust mass should be $3.4\times10^{-6}\,M_\odot$, the corresponding density near the base of the crust should be $8.33\times10^{10}\,\mathrm{g\,cm^{-3}}$, much lower than ρ_{drip}. The situation that the density at the base of the crust is still much lower than ρ_{drip} when z_{c} decreases toward 200 fm or even toward zero is evidently shown in Fig. 2. A crust density limit of less than ρ_{drip} is still obtained when V_{q} is evaluated as low as 15 MeV or V_{c} as high as 15 MeV.

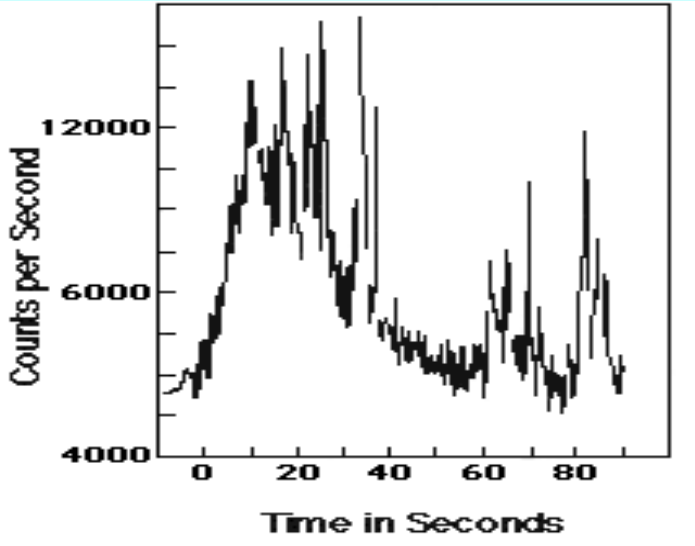
3. These new results will also influence the series of strange stars-strange dwarfs, which will be published later.

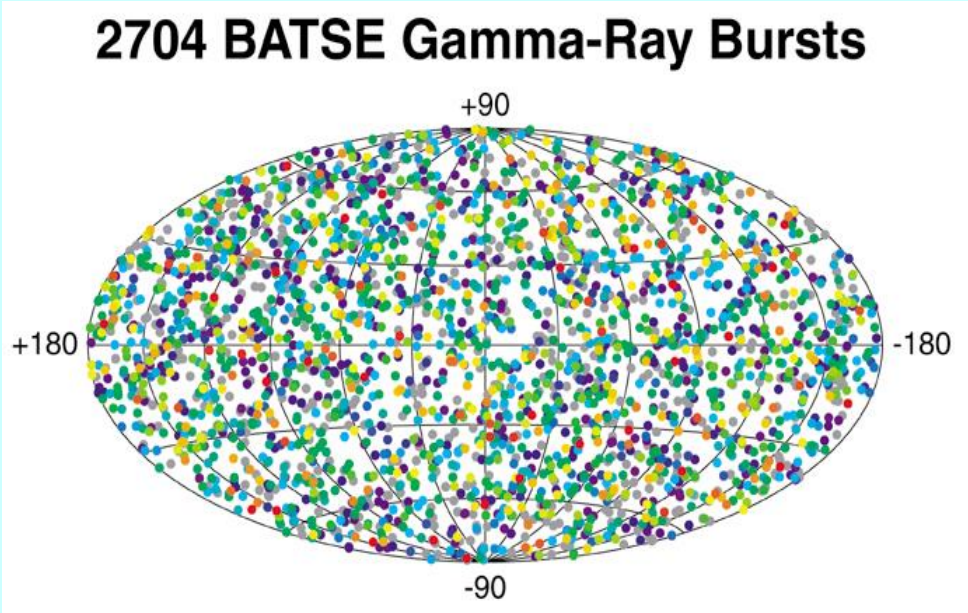
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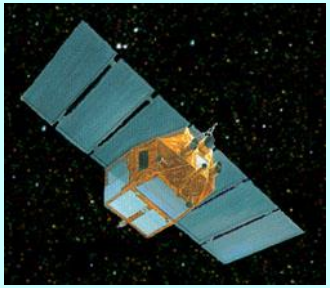


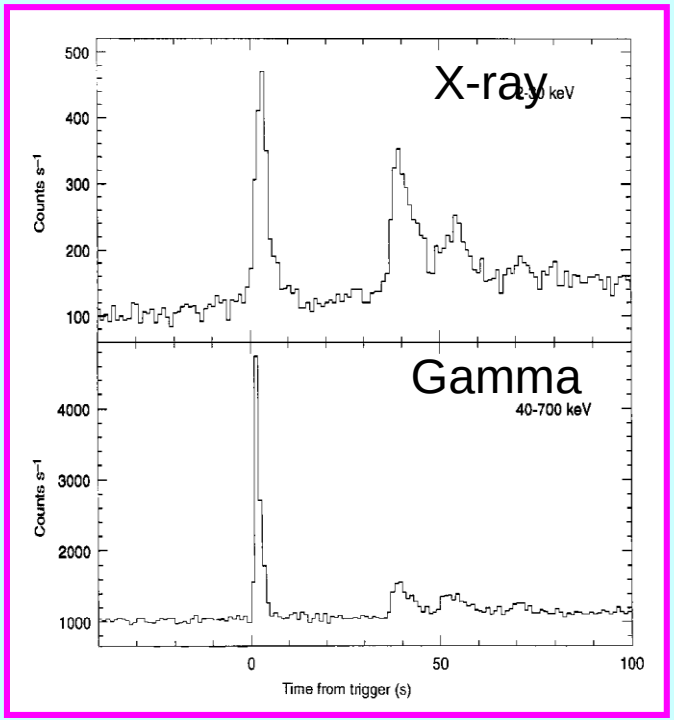
1997年：伽马射线暴研究进入新阶段



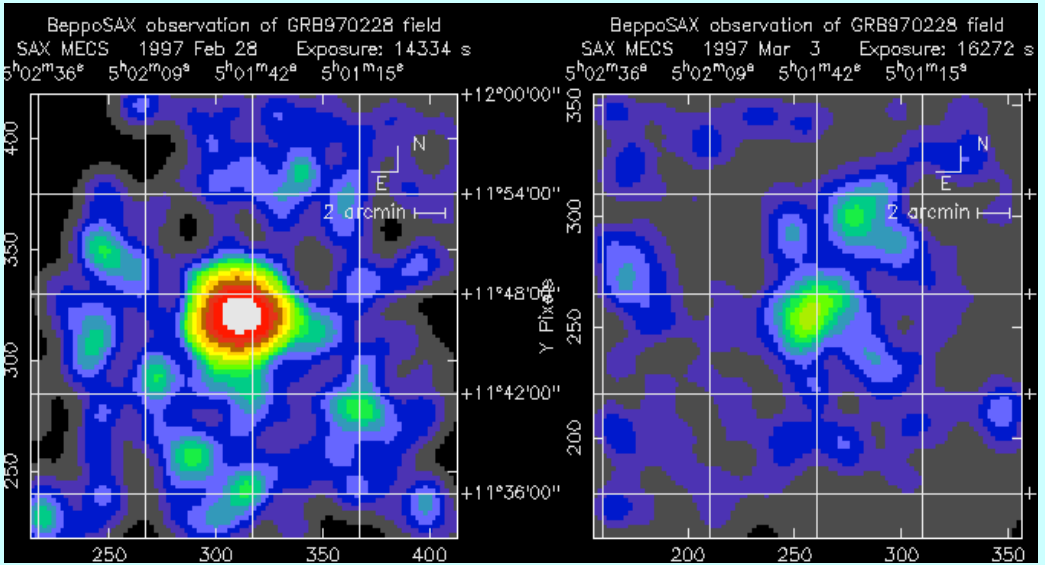


GRB 970228: the milestone



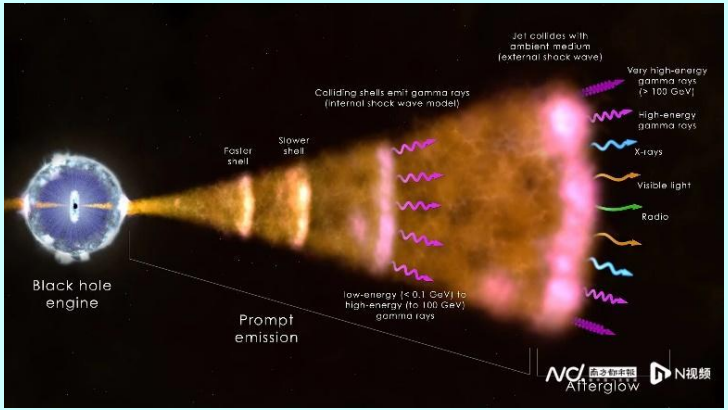


Fading X-ray afterglows

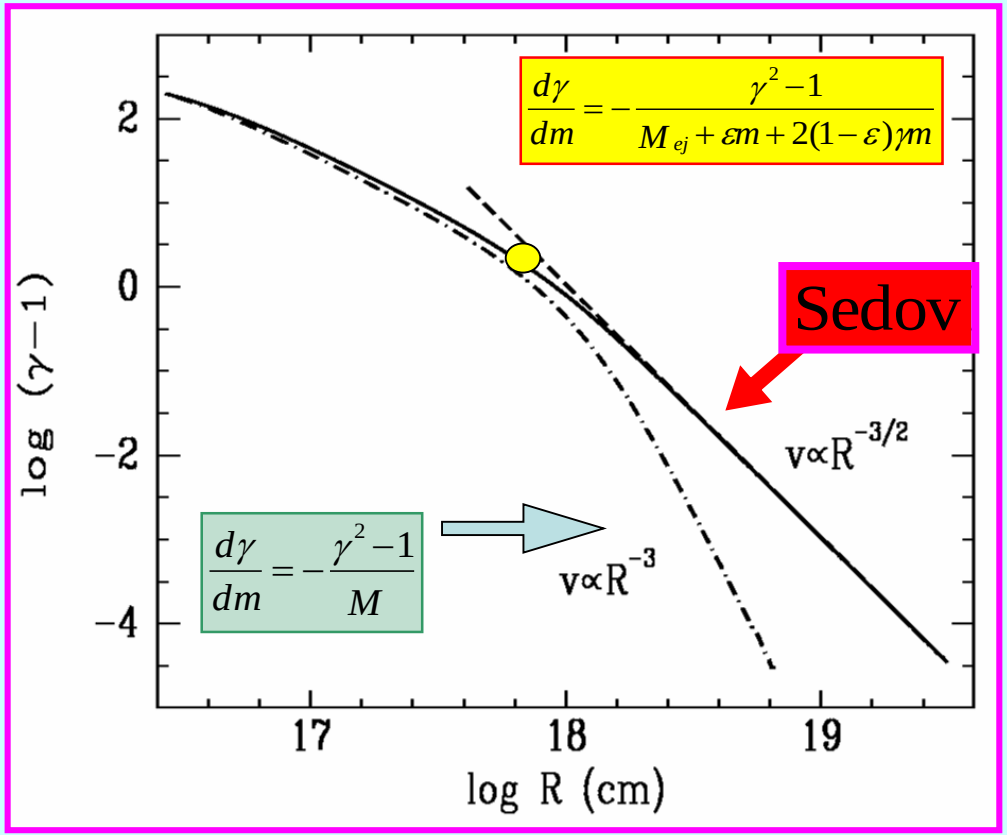


On Feb 28

3 days later



我们给出了正确的统一动力学公式：



Huang, Dai & Lu 1999, MNRAS, 309, 513

文献中的统一动力学公式：

$$\frac{d\gamma}{dm} = -\frac{\gamma^2-1}{M}$$

$$(dM = [(1-\varepsilon)\gamma+\varepsilon]dm)$$

但我们发现，该公式给出：$v\propto R^{-3}$;
而Sedov近似却要求：$v\propto R^{-3/2}$

Mon. Not. R. Astron. Soc. **309**, 513–516 (1999)

A generic dynamical model of gamma-ray burst remnants

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ABSTRACT

The conventional generic model is considered to explain the dynamics of gamma-ray burst remnants very well, no matter whether they are adiabatic or highly radiative. However, we find that, for adiabatic expansion, the model cannot reproduce the Sedov solution in the non-relativistic phase, and thus it needs to be revised. In this paper a new differential equation is derived. The generic model based on this equation is shown to be correct for both radiative and adiabatic fireballs, and in both ultrarelativistic and non-relativistic phases.

Key words: hydrodynamics – relativity – shock waves – gamma-rays: bursts.

1 INTRODUCTION

Since the *Beppo*SAX detection of GRB 970228, X-ray afterglows have been observed from about 15 gamma-ray bursts (GRBs), of which 10 have been detected optically and five also at radio wavelengths (Costa et al. 1997; Kulkarni et al. 1998; Bloom et al. 1998; Piran 1999, and references therein). The cosmological origin of at least some GRBs is thus firmly established. The so-called fireball model (Goodman 1986; Paczyński 1986; Rees & Mészáros 1992, 1994; Mészáros & Rees 1992; Katz 1994; Sari, Narayan & Piran 1996) is strongly favoured, and is successful at explaining the major features of the low-energy light curves (Mészáros & Rees 1997; Vietri 1997; Tavani 1997; Waxman 1997; Wijers, Rees & Mészáros 1997; Sari 1997; Huang et al. 1998b; Dai & Lu 1998a; Dai, Huang & Lu 1999). A variant of this model, in which central engines (e.g. strongly magnetized millisecond pulsars) supply energy to post-burst fireballs through magnetic dipole radiation, has been proposed to account for the special features of the optical afterglows from GRB 970228 and 970508 (Dai & Lu 1998b,c).

Since the expansion of a fireball may be either adiabatic or highly radiative, extensive attempts have been made to find a common model applicable to both cases (Blandford & McKee 1976; Chiang & Dermer 1999; Piran 1999). As a result, a conventional model has been suggested by various authors (e.g. Chiang & Dermer 1999; Piran 1999). A dynamical model should be correct not only in the initial ultrarelativistic phase, which is well described by those simple scaling laws (Mészáros & Rees 1997; Vietri 1997; Waxman 1997), but also in the consequent non-relativistic phase, which is correctly analysed by using the Sedov solution (Sedov 1969; Wijers et al. 1997). Although the conventional model is correct for the ultrarelativistic phase, we find that it cannot match the Sedov solution in the non-relativistic limit. In this paper we will therefore construct a dynamical model that is really capable of describing generic fireballs, no matter whether they are radiative or adiabatic, and no matter whether they are ultrarelativistic or non-relativistic.

2 CONVENTIONAL DYNAMICAL MODEL

A differential equation has been proposed to depict the expansion of GRB remnants (Chiang & Dermer 1999; Piran 1999):

$$\frac{\mathrm{d}\gamma}{\mathrm{d}m} = -\frac{\gamma^2 - 1}{M}, \tag{1}$$

where m is the rest mass of the swept-up medium, γ is the bulk Lorentz factor and M is the total mass in the comoving frame, including the internal energy U. Since the thermal energy produced during the collisions is $\mathrm{d}E = (\gamma - 1)\,\mathrm{d}m\,c^2$, we usually assume that $\mathrm{d}M = (1 - \epsilon)\,\mathrm{d}E/c^2 + \mathrm{d}m = [(1 - \epsilon)\gamma + \epsilon]\,\mathrm{d}m$, where ϵ is defined as the fraction of the shock-generated thermal energy (in the comoving frame) that is radiated (Piran 1999). It is supposed that equation (1) is correct in both the ultrarelativistic and non-relativistic phases, for both radiative and adiabatic fireballs. However, after careful inspection, we find that, during the non-relativistic phase of an adiabatic expansion, equation (1) cannot give a solution consistent with the Sedov results (Sedov 1969).

2.1 Radiative case

In the highly radiative case, $\epsilon = 1$, $\mathrm{d}M = \mathrm{d}m$, equation (1) reduces to

$$\frac{\mathrm{d}\gamma}{\mathrm{d}m} = -\frac{\gamma^2 - 1}{M_{\mathrm{ej}} + m}, \tag{2}$$

where M_{ej} is the mass ejected from the GRB central engine. Then an analytic solution is available (Blandford & McKee 1976; highly radiative GRB remnants, but it has difficulty in reproducing the Sedov solution for adiabatic fireballs. This has gone completely unnoticed in the literature. We have constructed a new generic model to overcome this shortcoming. Numerical evaluation has proved that our model is highly credible. We hope that this work will remind researchers of the importance of the transition from the ultrarelativistic to the non-relativistic phase, which might occur as early as 10^6–10^7 s after the initial burst (Huang, Dai & Lu 1998a).

In the above analysis, for simplicity, we have assumed that ϵ is a constant. In realistic fireballs, however, ϵ is expected to evolve from 1 to 0 owing to the changes in the relative importance of synchrotron-induced and expansion-induced loss of energy (Dai et al. 1999). Assuming that electrons in the comoving frame carry a fraction $\xi_e = 1$ of the total thermal energy and that the magnetic energy density is a fraction $\xi_B' = 0.01$ of it, we re-evaluate equation (6) numerically. The results are plotted in Figs 1–4 with dash–dotted lines. We see from Fig. 4 that the evolution of ϵ changes $E_{\mathrm{K}}(t)$ dramatically.

It is worth mentioning that SNRs evolve from the non-radiative to the radiative stage, but GRB remnants are just the opposite.

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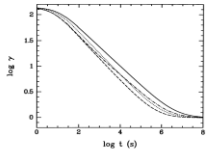


Figure 1. Evolution of the bulk Lorentz factor γ. We take $E_0 = 10^{52}$ erg, $n = 1\,\mathrm{cm}^{-3}$ and $M_{\mathrm{ej}} = 2 \times 10^{-5}\,\mathrm{M}_{\odot}$. The full, dotted and dashed lines correspond to $\epsilon = 0$ (adiabatic), 0.5 (partially radiative) and 1 (highly radiative) respectively. The dash–dotted line is plotted by allowing ϵ to evolve with time (see Section 5).

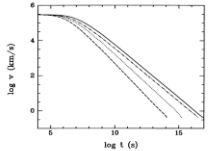


Figure 2. Evolution of the bulk velocity v. Parameters and line styles are the same as in Fig. 1.

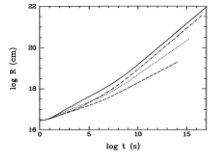


Figure 3. Evolution of the shock radius R. Parameters and line styles are the same as in Fig. 1.

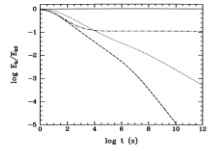


Figure 4. Evolution of the total kinetic energy E_{K}. Parameters and line styles are the same as in Fig. 1.

This is not surprising, because GRB remnants radiate mainly through synchrotron radiation while SNRs lose energy as a result of excited ions. It is reasonable to deduce that at very late stages, when the cooling resulting from ions becomes important, GRB remnants may become highly radiative again, in the same way as SNRs do. The transition may occur when the temperature drops to below $\sim 10^8$ K and the velocity is just several tens of kilometres per second. This needs to be addressed in more detail.

Another interesting problem is the possibility that HI super-shells might be highly evolved GRB remnants (Loeb & Perna 1998; Efremov, Elmegreen & Hodge 1998). Our Figs 3 and 4 have shown that typical adiabatic GRB fireballs can evolve to $R \sim 1$ kpc at $t \sim 10^6$–10^7 yr, with $v \sim 10\,\mathrm{km\,s^{-1}}$, but highly radiative fireballs are obviously not powerful enough. To discuss this in detail, we should pay attention to the possible adiabatic-to-radiative transition mentioned above.

ACKNOWLEDGMENTS

We are very grateful to R. Wijers for his valuable comments and suggestions. This work was partly supported by the National Natural Science Foundation of China, grants 19773007 and

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Piran 1999):

$$\frac{(\gamma - 1)(\gamma_0 + 1)}{(\gamma + 1)(\gamma_0 - 1)} = \left(\frac{m_0 + M_{\mathrm{ej}}}{m + M_{\mathrm{ej}}}\right)^2, \tag{3}$$

where γ_0 and m_0 are initial values of γ and m respectively. Usually we assume $\gamma_0 - \eta/2$, $m_0 - M_{\mathrm{ej}}/\eta$, where $\eta \approx E_0/(M_{\mathrm{ej}}c^2)$ and E_0 is the total energy in the initial fireball (Waxman 1997; Piran 1999).

During the ultrarelativistic phase, $\gamma \gg 1$, $M_{\mathrm{ej}} \gg m$, equation (3) gives $(\gamma + 1)m = M_{\mathrm{ej}}$, or equivalently the familiar power law $\gamma \propto R^{-3}$, where R is the radius of the blast wave. In the later non-relativistic phase, $\gamma - 1$, $m \gg M_{\mathrm{ej}}$, we have $m\beta^2 = 4M_{\mathrm{ej}}^2$, or $\beta \propto R^{-3}$, where $\beta = v/c$ and v is the bulk velocity of the material. This is consistent with the late isothermal phase of the expansion of supernova remnants (SNRs) (Spitzer 1968). From these approximations, we believe that equation (2) is really correct for highly radiative fireballs.

2.2 Adiabatic case

In the adiabatic case, $\epsilon = 0$, $\mathrm{d}M = \gamma\,\mathrm{d}m$, equation (1) also has an analytic solution (Chiang & Dermer 1999):

$$M = (M_{\mathrm{ej}}^2 + 2\gamma_0 M_{\mathrm{ej}} m + m^2)^{1/2}, \tag{4}$$

$$\gamma = \frac{m + \gamma_0 M_{\mathrm{ej}}}{M}. \tag{5}$$

During the ultrarelativistic phase, $\gamma_0 M_{\mathrm{ej}} \gg m \gg M_{\mathrm{ej}}/\gamma_0$, $\gamma \gg 1$, this solution can produce the familiar power law $\gamma \propto R^{-3/2}$, which is often quoted for an adiabatic blastwave decelerating in a uniform medium. In the non-relativistic limit ($\gamma - 1$, $m \gg \gamma_0 M_{\mathrm{ej}}$), Chiang & Dermer (1999) have derived $\gamma \propto 1 + \gamma_0 M_{\mathrm{ej}}/m$, so that they believe it also agrees with the Sedov solution (Lozinskaya 1992). However, we find that their approximation is not accurate enough, because they have omitted some first-order infinitesimals of $\gamma_0 M_{\mathrm{ej}}/m$. The correct approximation can be obtained only by retaining all the first- and second-order infinitesimals, which in fact gives $\gamma = 1 + (\gamma_0 M_{\mathrm{ej}}/m)^2/2$; then we have $\beta \propto R^{-2}$. This is not consistent with the Sedov solution! We have also evaluated equation (1) numerically; the result is consistent with equations (4) and (5), all pointing to $\beta \propto R^{-2}$, not the relation $\beta \propto R^{-3/2}$ as often quoted in the literature (Chiang & Dermer 1999; Piran 1999).

This discrepancy is serious. First, it means that equation (1) is not a dependable model for non-radiative fireballs, although it can reproduce the major features in the ultrarelativistic phase. Secondly, the expansion of a realistic fireball is widely believed to be highly radiative at first, but after only a few days the expansion will become non-radiative (Sari, Piran & Narayan 1998; Dai et al. 1999). In the non-relativistic phase, therefore, the fireball is likely to be adiabatic rather than highly radiative. However, it is under just this condition that the conventional model fails. So any calculation made according to equation (1) will lead to serious deviations in the light curves in the non-relativistic phase.

3 OUR GENERIC MODEL

Equation (1) is not consistent with the Sedov solution, and we need to revise it. In the fixed frame, since the total kinetic energy of the fireball is $E_{\mathrm{K}} = (\gamma - 1)(M_{\mathrm{ej}} + m)c^2 + (1 - \epsilon)\gamma U$ (Panaitescu, Mészáros & Rees 1998), and the radiated thermal energy is $\epsilon\gamma(\gamma - 1)\,\mathrm{d}mc^2$ (Blandford & McKee 1976), we have

$$\mathrm{d}[(\gamma - 1)(M_{\mathrm{ej}} + m)c^2 + (1 - \epsilon)\gamma U] = -\epsilon\gamma(\gamma - 1)\,\mathrm{d}mc^2. \tag{6}$$

For the variable U, it is usually assumed that $\mathrm{d}U = (\gamma - 1)\,\mathrm{d}mc^2$ (Panaitescu et al. 1998). Equation (1) has been derived just in this way. However, the jump conditions (Blandford & McKee 1976) at the forward shock imply that $U = (\gamma - 1)mc^2$, so we suggest that the correct expression for $\mathrm{d}U$ should be $\mathrm{d}U = \mathrm{d}[(\gamma - 1)mc^2] = (\gamma - 1)\,\mathrm{d}mc^2 + mc^2\,\mathrm{d}\gamma$. Here we simply use $U = (\gamma - 1)mc^2$ and substitute it into equation (6); then it is easy to obtain

$$\frac{\mathrm{d}\gamma}{\mathrm{d}m} = -\frac{\gamma^2 - 1}{M_{\mathrm{ej}} + \epsilon m + 2(1 - \epsilon)\gamma m}. \tag{7}$$

We expect this equation to describe a generic fireball correctly.

Indeed, in the highly radiative case ($\epsilon = 1$) equation (7) reduces to equation (2) exactly, while in the adiabatic case ($\epsilon = 0$) equation (7) reduces to

$$\frac{\mathrm{d}\gamma}{\mathrm{d}m} = -\frac{\gamma^2 - 1}{M_{\mathrm{ej}} + 2\gamma m}. \tag{8}$$

This equation has an analytic solution:

$$(\gamma - 1)M_{\mathrm{ej}}c^2 + (\gamma^2 - 1)mc^2 = E_{\mathrm{K}0}, \tag{9}$$

where $E_{\mathrm{K}0}$ is the initial value of E_{K}. In the ultrarelativistic phase ($\gamma_0 M_{\mathrm{ej}} \gg m \gg M_{\mathrm{ej}}/\gamma$) we obtain the familiar relation $\gamma \propto R^{-3/2}$, and in the non-relativistic phase ($m \gg M_{\mathrm{ej}}$) we obtain $\beta \propto R^{-3/2}$ as required by the Sedov solution.

For any other ϵ-value between 0 and 1, equation (7) describes the evolution of a partially radiative fireball. Unfortunately, we now cannot find an exact analytic solution for equation (7). In the non-relativistic phase, however, by assuming that $m \gg M_{\mathrm{ej}}$, we can still get $m\gamma - 1)^{2-\alpha} \propto$ constant, that is,

$$\beta \propto R^{-3/2\cdot\alpha}. \tag{10}$$

4 NUMERICAL RESULTS

We have evaluated equation (7) numerically, bearing in mind that (Huang et al. 1998b)

$$\mathrm{d}m = 4\pi R^2 n m_{\mathrm{p}}\,\mathrm{d}R, \tag{11}$$

$$\mathrm{d}R = \beta c\gamma\left(\gamma + \sqrt{\gamma^2 - 1}\right)\mathrm{d}t, \tag{12}$$

where n is the number density of the interstellar medium, m_{p} is the mass of a proton, and t is the time measured by an observer. We take $E_0 = 10^{52}$ erg, $n = 1\,\mathrm{cm}^{-3}$ and $M_{\mathrm{ej}} = 2 \times 10^{-5}\,\mathrm{M}_{\odot}$. Figs 1–4 illustrate the evolution of γ, v, R and E_{K} respectively. In these figures, we have set $\epsilon = 0$ (full lines), 0.5 (dotted lines) and 1 (dashed lines). It is clearly shown that our generic model overcomes the shortcomings of equation (1).

For example, for highly radiative expansion, the dashed lines in these figures approximately satisfy $\gamma \propto t^{-3/7}$, $R \propto t^{1/7}$, $\gamma \propto R^{-3}$, $E_{\mathrm{K}} \propto t^{-3/7}$ when $\gamma \gg 1$, and $\gamma \propto t^{-3/4}$, $R \propto t^{1/4}$, $v \propto R^{-3}$, $E_{\mathrm{K}} \propto t^{-3/4}$ when $\gamma - 1$. For adiabatic expansion, the full lines satisfy $\gamma \propto t^{-3/8}$, $R \propto t^{1/4}$, $\gamma \propto R^{-3/2}$ when $\gamma \gg 1$, and $v \propto t^{-3/5}$, $R \propto t^{2/5}$, $v \propto R^{-7/2}$ when $\gamma - 1$.

5 DISCUSSION AND CONCLUSION

The conventional dynamical model is successful at describing

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19825109, and the National Climbing Project on Fundamental Researches.

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A generic dynamical model of GRB remnants

Huang, Dai, Lu, 1999, MNRAS, 309, 513

MNRAS审稿意见

This paper sets straight an obvious problem in the equations for the evolution of fireballs, and as is usual with things that are right, you almost wonder how we could have got it wrong in the first place. It is a must-publish, since it prevents the further propagation of the error addressed by the authors.

van Paradijs et al. (2000, *ARA&A*, 38, 379) 的引用

$E \simeq E_0 \propto \gamma^2 m = const.$, so γ slowly decreases as $\gamma \propto m^{-\frac{1}{2}}$. For the simplest case of a uniform ambient medium, we then have the standard result for a self-similar relativistic blast wave: $\gamma \propto r^{-3/2}$.

To get the corresponding observer times, we need to introduce the kinematic relation between radius and observer time for a relativistically approaching shell (Rees 1964):

$$dt = dr/2\gamma^2 c. \tag{2}$$

Before serious deceleration, we may set $\gamma = \gamma_{\rm fm}$ and omit the differentials. For a pure hydrogen medium with number density $n = \rho/m_{\rm p}$ we then get the radius and observer time of the onset of deceleration (Rees & Mészáros 1992):

$$r_{\rm nm} = [3E_{\rm iso}/4\pi\gamma_{\rm fm}^2 nm_{\rm p}c^2]^{1/3} = 3.8 \times 10^{16}\,(E_{\rm iso}/n)^{1/3}\,\gamma_{\rm fm}^{-2/3}\ {\rm cm} \tag{3}$$

$$t_{\rm nm} = r_{\rm nm}/2\gamma_{\rm fm}^2 c = 63\,(E_{\rm iso}/n)^{1/3}\,\gamma_{\rm fm}^{-8/3}\ {\rm s}. \tag{4}$$

Compared with the supernova case, deceleration takes place at somewhat smaller radii, but most starkly within seconds rather than centuries; this is caused by the higher speed of the ejecta, but even more by the two Lorentz factors in Equation 2. The relativistic phase ends and turns into the usual supernova behavior (Sedov 1969 Ch. IV, Taylor 1950) when the energy per particle approaches $m_{\rm p}c^2$. At that time the Lorentz factor corrections disappear, and we may estimate the size simply as ct, hence

$$r_{\rm snr} = (3E_{\rm iso}/4\pi nm_{\rm p}c^2)^{1/3} = 1.2\,(E_{\rm iso}/n)^{1/3}{\rm yr}. \tag{5}$$

As we see from the preceding qualitative discussion, the relativistic self-similar phase can span many orders of magnitude in time. We now derive the evolution for the blast wave more precisely and generally, following Huang et al (1999; see also Katz & Piran 1997). We assume that material passing through the shock quickly radiates a fraction ϵ of the post-shock thermal energy, and retains the rest as thermal energy. The total (kinetic plus thermal) energy of the burst is then $E = (\gamma - 1)(M_0 + m)c^2 + (1 - \epsilon)\gamma U$, and the radiated energy is $dE_{\rm rad} = \epsilon\,\gamma(\gamma-1)c^2\,dm$ (Panaitescu & Mészáros 1998a, Blandford & McKee 1976). From the shock jump conditions, one finds $U = (\gamma - 1)mc^2$, which in combination with $dE = dE_{\rm rad}$ gives

$$\frac{d\gamma}{dm} = -\frac{\gamma^2 - 1}{M_0 + \epsilon m + 2(1-\epsilon)\gamma m}. \tag{6}$$

For fully radiative ($\epsilon = 1$) and fully adiabatic ($\epsilon = 0$) blast waves, one can write analytic solutions for the expansion. Equation 6 differs somewhat from the sometimes quoted

$$\frac{d\gamma}{dm} = -\frac{\gamma^2 - 1}{M}, \tag{7}$$

in which M is the total rest-frame mass (inclusive of mass equivalent of internal energy; Chiang & Dermer 1999, Katz 1994b). For $\epsilon = 1$, the forms are equivalent, but the adiabatic form of Equation 7 is only valid in the ultrarelativistic case: its solution near $\gamma = 1$ should yield the classical Sedov-Taylor solution but does not. Equation 6 does yield the Sedov-Taylor solution in the non-relativistic limit.

3.1.1 Fully Radiative Case, $\epsilon = 1$ The equation of motion is easily integrated to give (Blandford & McKee 1976, Huang et al 1999, Katz & Piran 1997)

$$\left(\frac{\gamma - 1}{\gamma + 1}\right)\left(\frac{\gamma_0 + 1}{\gamma_0 - 1}\right) = \left(\frac{M_0 + m_0}{M_0 + m}\right)^2, \tag{8}$$

where $\gamma_0 \sim \eta/2$, $m_0 \sim M_0/\eta$ are the initial conditions of the deceleration. The ultrarelativistic limit is obtained for $M_0 \gg m$, since the deceleration starts in that limit, and the large radiative losses mean that the shock comes to a halt after having swept up only a few times m_0. Expanding Equation 8, we obtain $\gamma m = M_0$, i.e. $\gamma \propto m^{-1}$. For a uniform ambient medium, this gives the well-known $\gamma \propto r^{-3}$. Inserting this into the kinematic relation between radius and observer time (Equation 2), we find that $\gamma \propto t^{-3/8}$ and $r \propto t^{1/8}$; the Lorentz factor decreases substantially at very slowly changing radius. In the non-relativistic limit, the solution becomes $\beta \propto m^{-2}$, which for a uniform medium is the snowplow (radiative) phase of supernova remnant expansion.

3.1.2 Adiabatic Case, $\epsilon = 0$ The equation of motion for this case is easily integrated to give $\gamma M_0 + (\gamma^2 - 1)m = const.$ To interpret this result, we subtract M_0 and multiply by c^2, and note that the result is the sum of thermal and bulk kinetic energies of the blast wave, hence

$$(\gamma - 1)M_0c^2 + (\gamma^2 - 1)mc^2 = E_{\rm tot}, \tag{9}$$

where $E_{\rm tot}$ is a constant. In the ultrarelativistic limit ($m \gg M_0/\gamma$ and $\gamma \gg 1$), this implies $\gamma^2 m = const.$ For a uniform ambient medium, this then gives $\gamma \propto r^{-3/2}$ as derived qualitatively above. With the kinematic r–t relation, this then gives $\gamma \propto t^{-3/8}$ and $r \propto t^{1/8}$. This demonstrates that although the self-similar relativistic case can cover 6–7 orders of magnitude in time, the range in radii it covers is not nearly as large. In the non-relativistic limit ($\gamma \sim 1, m \gg M_0$) we get $\beta \propto r^{-3/2}$, which is the proper Sedov-Taylor solution.

The above relations define the limiting cases of the dynamical behavior of the blast wave. Unfortunately, the great distances make the blast waves unresolved on the sky, so we cannot directly test the relations between ϵ, t, and β, γ. [The scintillation size determination of GRB 970508 (Frail et al 1997c) is the one exception to this.] Therefore, we will have to describe the flux and spectrum of radiation from the expanding shock before we can find direct tests of our model from data.

博士期间完成了10篇论文

GAMMA-RAY BURSTS
AND STRANGE STARS

Huang Yong-feng

Department of Astronomy
Nanjing University

April, 1999

APPLYING FOR Ph. D.

ADVISOR: Professor Lu Tan

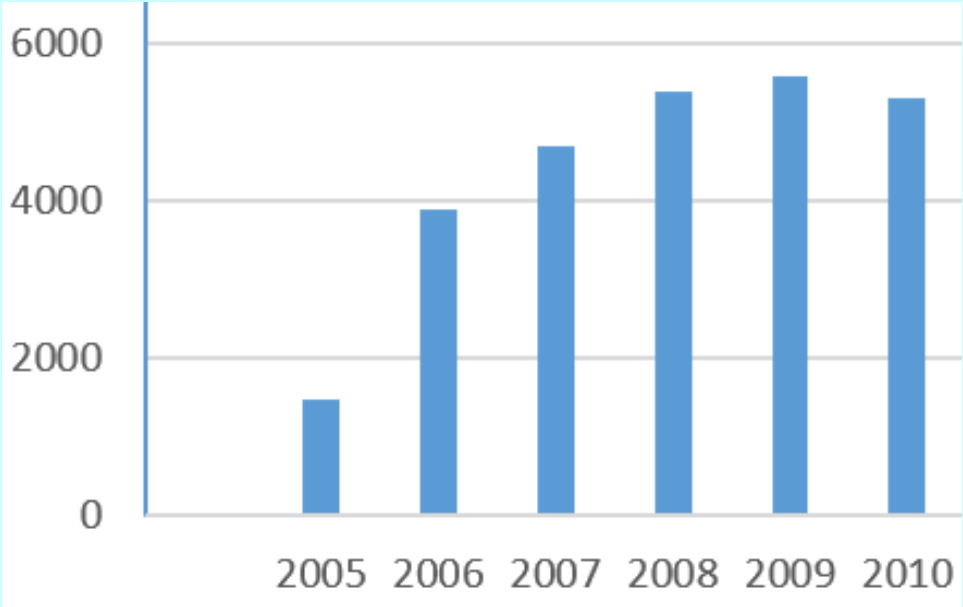
我们每个人都会有迷茫、彷徨和困难的时候。

关键的是，越是感到迷茫，我们越是要努力，
才能改变困难的局面。

如果沉沦下去，可能就永远无法摆脱困境。

随时把自己的科研想法记下来

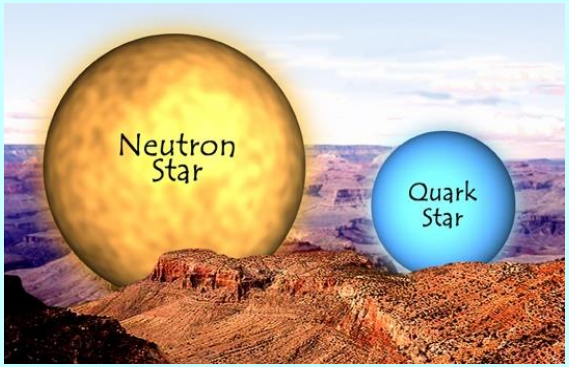
1999：博士毕业，留校工作



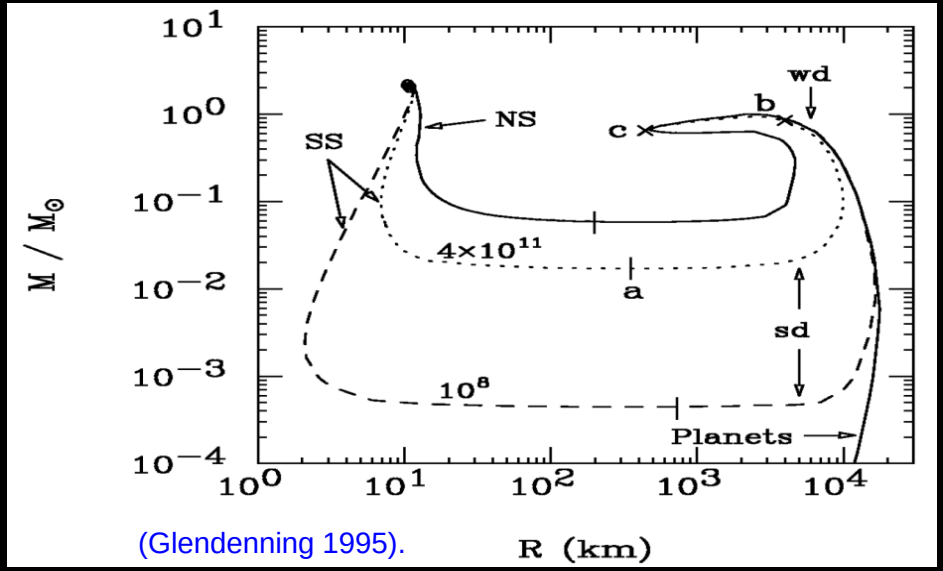
三、科研之路，创新相伴

如何观测区分中子星和奇异星?

- Different radius at M ~ 1.4 M_{sun}?
- Different M~R relation?
- Different M_{max} ?
- Different cooling rate?
- Can spin more quickly (P < 1ms)?
- Gravitational Waves?
- ……

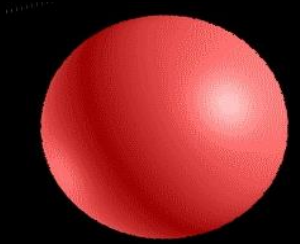


- We do not have an unambiguous criterion!



(Glendenning 1995).

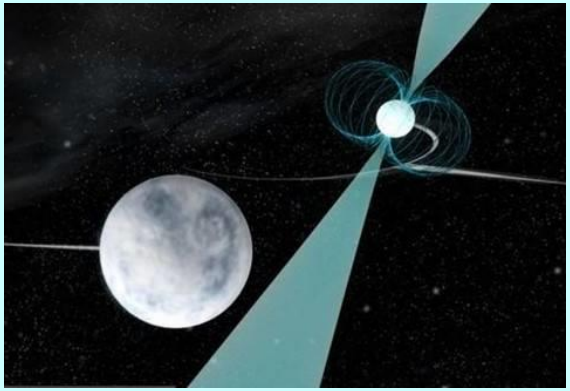
（1）寻找奇异夸克行星？

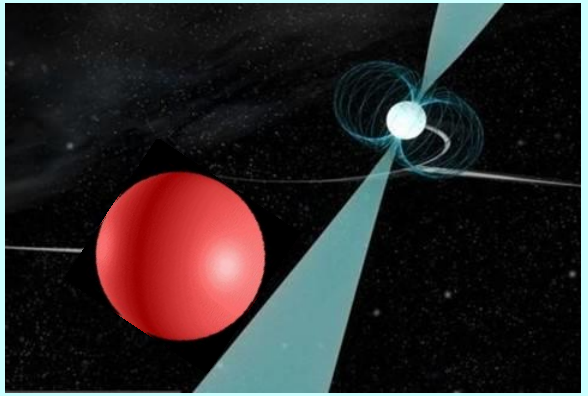


$(\rho \sim 10^{15}\, g/cm^3)$

Parameters of bare strange planets ：

M	$10^{-9}M_{sun}$	$10^{-6}M_{sun}$	$10^{-5}M_{sun}$	$10^{-4}M_{sun}$	$10^{-3}M_{sun}$
R	8m	80 m	170 m	370 m	800m



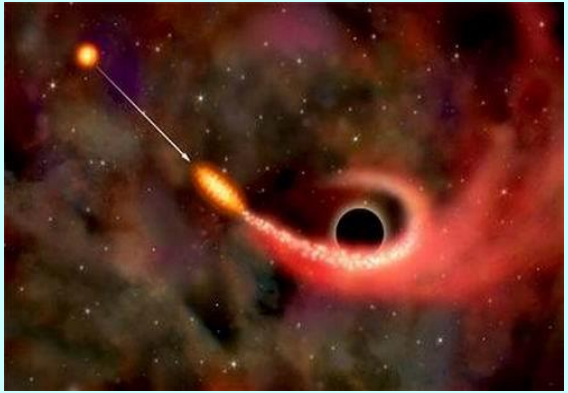


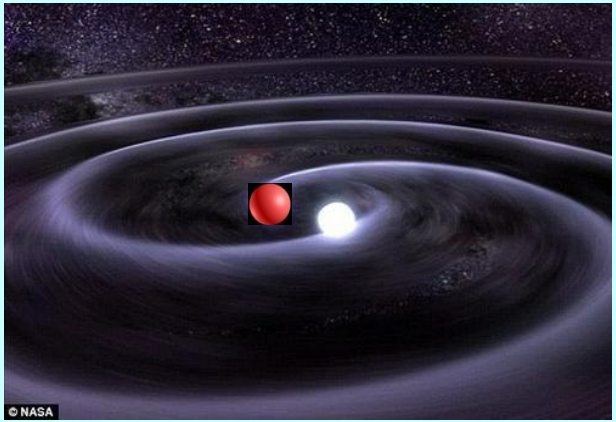
Tidal disruption radius of normal planets：

$$r_{\rm td} \approx 5.1 \times 10^{10} \left(\frac{M}{1.4\,M_\odot} \right)^{1/3} \left(\frac{\rho_0}{10\ {\rm g\ cm^{-3}}} \right)^{-1/3} {\rm cm}.$$

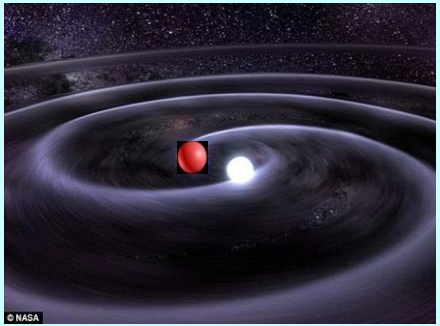
For a strange planet:

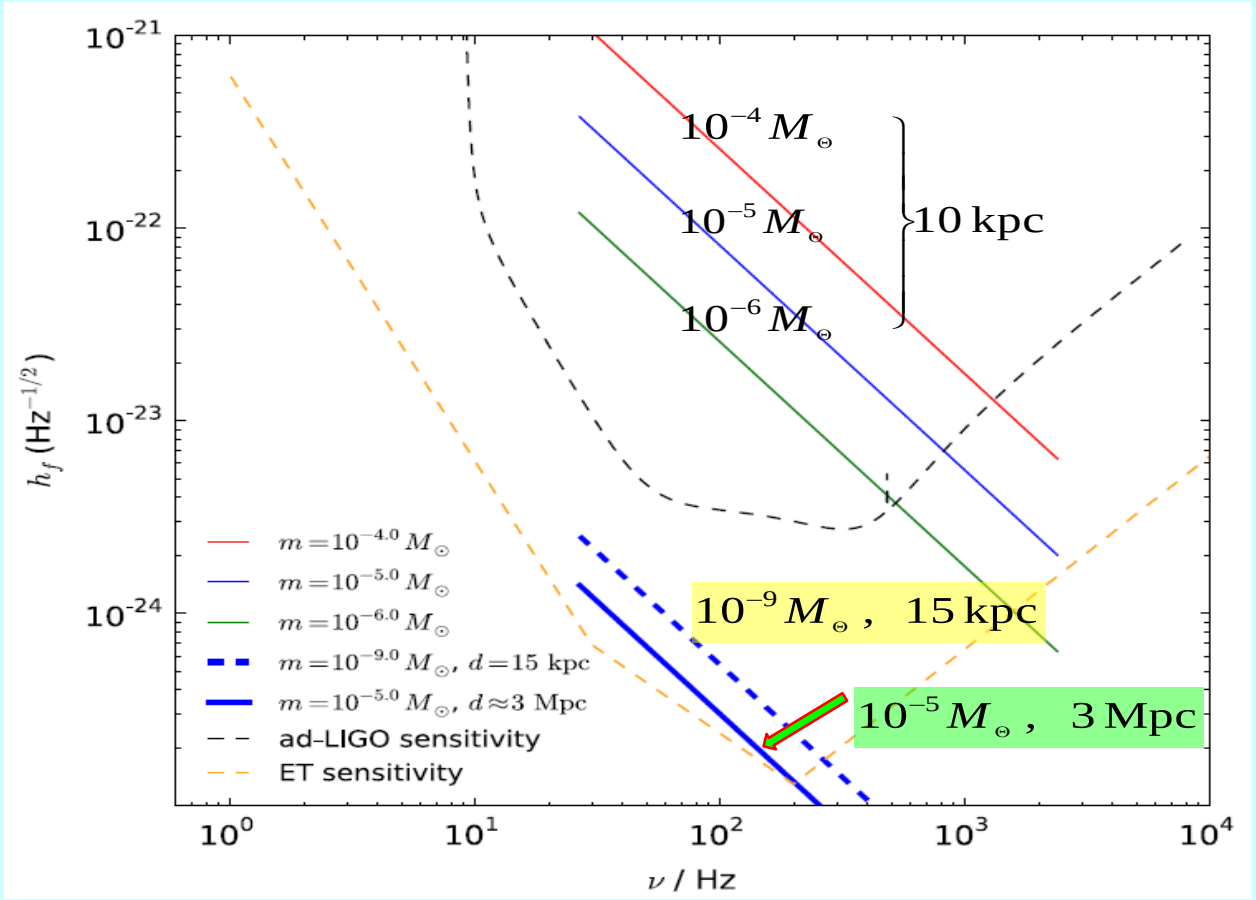
$$r_{td} = 1.5 \times 10^{6}\ \text{cm}$$



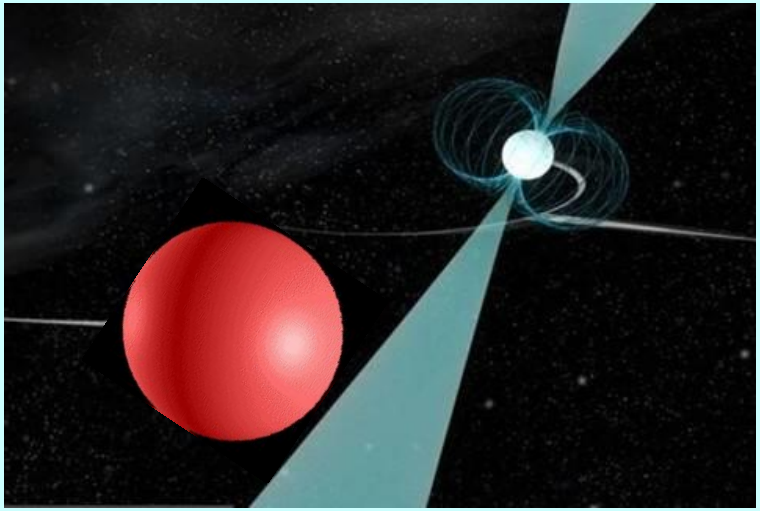


The GW bursts are detectable for adv-LIGO and ET.





Geng, Huang, Lu, 2015, ApJ, 804, 21



- If people detected a “strange” GW burst event from:

  a “neutron star” + a “planet”

- Then, it must be a SQM planetary system.

（2）以原初黑洞为探针，深入脉冲星内部?

The Astrophysical Journal Letters, 928:L13 (8pp), 2022 April 1 https://doi.org/10.3847/2041-8213/ac5ea6
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Gravitational-wave Emission from a Primordial Black Hole Inspiraling inside a Compact Star: A Novel Probe for Dense Matter Equation of State

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Abstract

Primordial black holes of planetary masses captured by compact stars are widely studied to constrain their composition fraction of dark matter. Such a capture may lead to an inspiral process and be detected through gravitational-wave signals. In this Letter, we study the postcapture inspiral process by considering two different kinds of compact stars, i.e., strange stars and neutron stars. The dynamical equations are numerically solved, and the gravitational-wave emission is calculated. It is found that the Advanced LIGO can detect the inspiraling of a $10^{-5}\ M_{\odot}$ primordial black hole at a distance of 10 kpc, while a Jovian-mass case can even be detected at megaparsecs. Promisingly, the next generation of gravitational-wave detectors can detect cases of $10^{-5}\ M_{\odot}$ primordial black holes up to ~1Mpc and Jovian-mass cases at several hundred megaparsecs. Moreover, the kilohertz gravitational-wave signal shows significant differences for strange stars and neutron stars, potentially making it a novel probe to the dense matter equation of state.

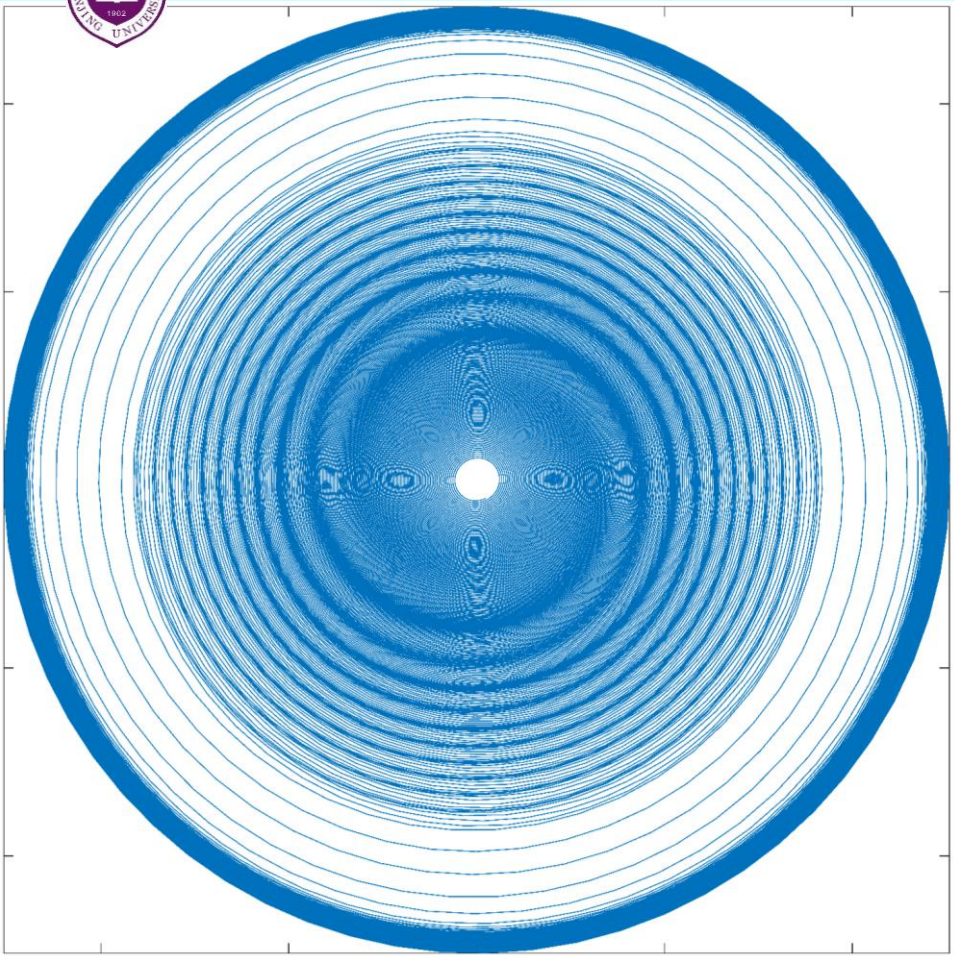
Unified Astronomy Thesaurus concepts: Gravitational waves (678); Neutron stars (1108); Nuclear astrophysics (1129); Primordial black holes (1292)

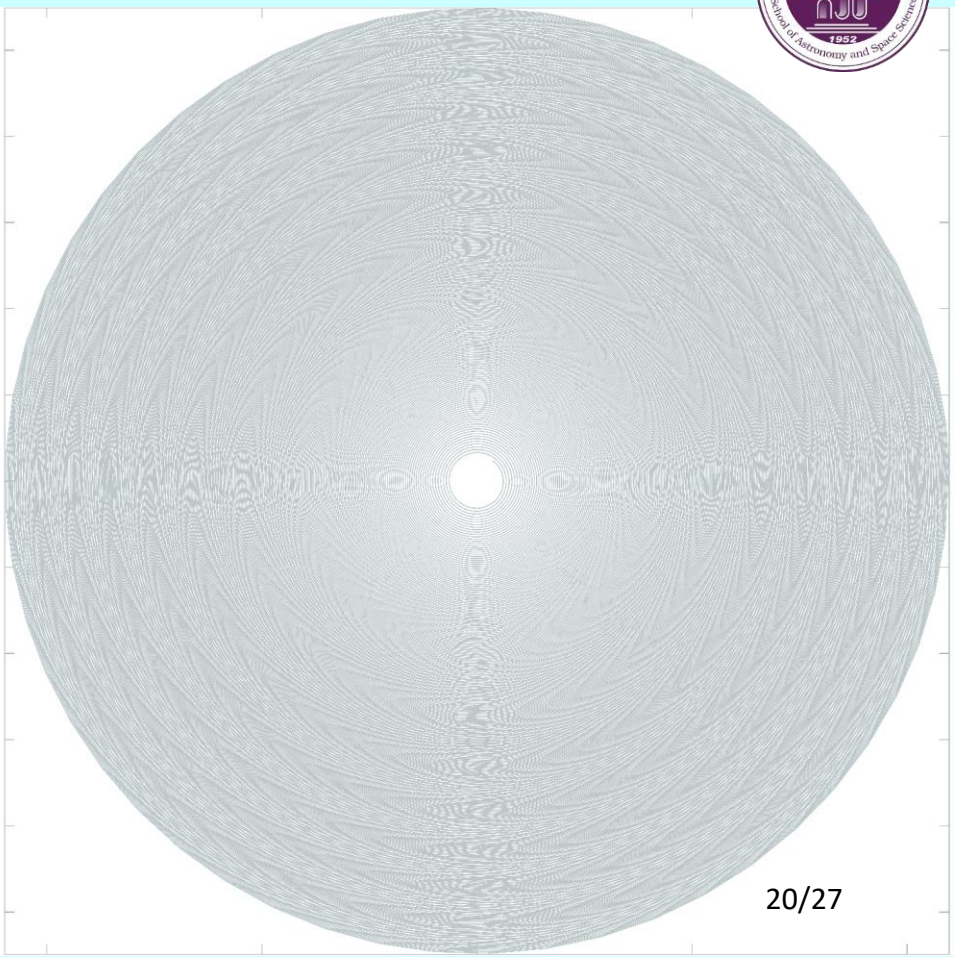


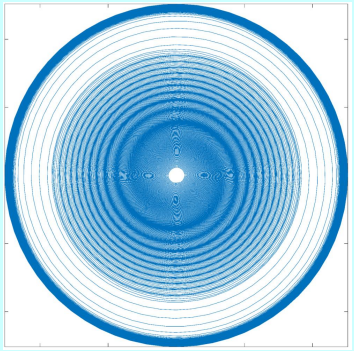
图片来自网络

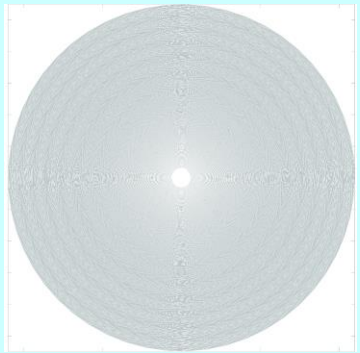
Zou & Huang, 2022, *ApJL*, 928, L13

Orbit: Neutron star vs. Strange star









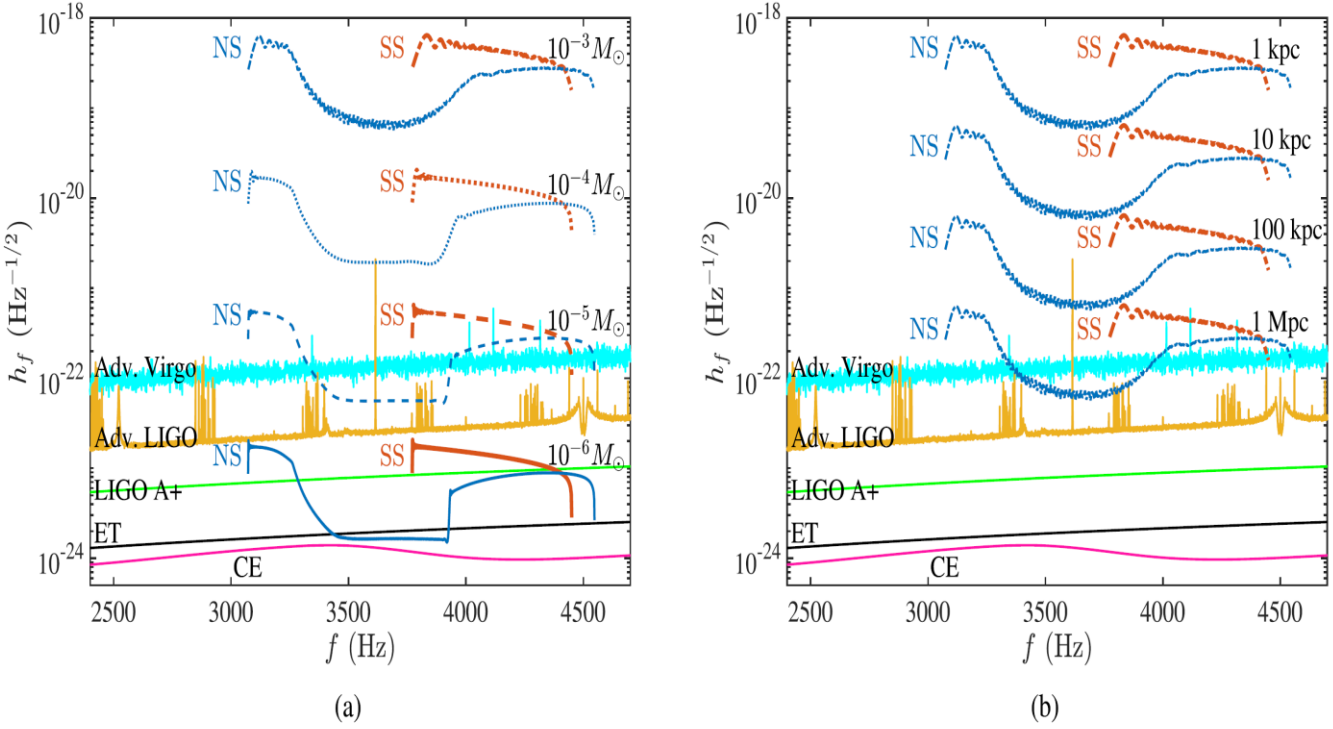
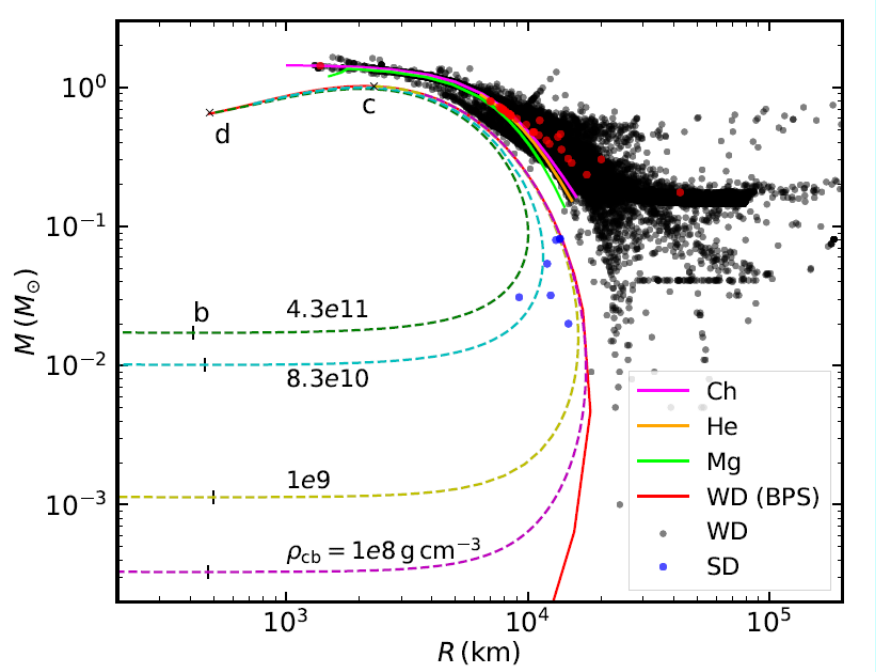


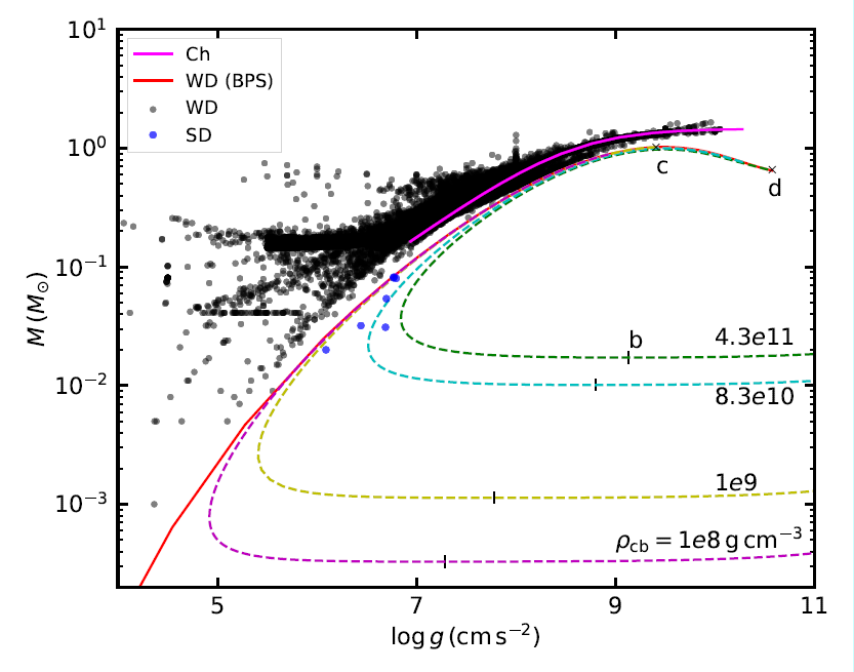
Figure 7. Strain spectral amplitude of GWs against frequency for a PBH inspiraling inside a compact star. In panel (a), the system is fixed at 1 kpc away, but different initial masses are assumed for the PBH and marked near the corresponding curves. In panel (b), the initial mass of the PBH is fixed as $10^{-3}\,M_\odot$, but the systems are

Zou & Huang, 2022, *ApJL*, 928, L13

（3）Strange quark dwarfs: do they exist?



M vs. R for observed “white dwarfs”



M vs. g for observed “white dwarfs”

Kuerban, Huang, Geng, Zong, 2022, Physics Letters B, 832, 137204 (arXiv:2012.05748)

（4）Explain the anti-glitch of 1E2259+586 in April 2012

LETTER

doi:10.1038/nature12159

An anti-glitch in a magnetar

R. F. Archibald[1], V. M. Kaspi[1], C. -Y. Ng[1,2], K. N. Gourgouliatos[1], D. Tsang[1], P. Scholz[1], A. P. Beardmore[3], N. Gehrels[4] & J. A. Kennea[5]

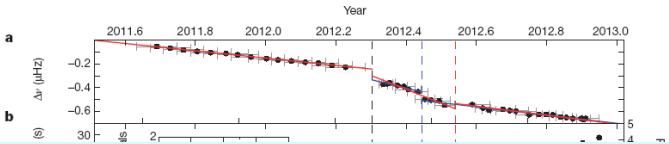
Magnetars are neutron stars with X-ray and soft γ-ray outbursts thought to be powered by intense internal magnetic fields[1]. Like conventional neutron stars in the form of radio pulsars, magnetars exhibit 'glitches' during which angular momentum is believed to be transferred between the solid outer crust and the superfluid component of the inner crust[2–4]. The several hundred observed glitches in radio pulsars[5,6] and magnetars[7] have involved a sudden spin-up (increase in the angular velocity) of the star, presumably because the interior superfluid was rotating faster than the crust. Here we report X-ray timing observations of the magnetar 1E 2259+586 (ref. 8), which exhibited a clear 'anti-glitch'—a sudden spin-down. We show that this event, like some previous magnetar spin-up glitches[9], was accompanied by multiple X-ray radiative changes and a significant spin-down rate change. Such behaviour is not predicted by models of neutron star spin-down and, if of internal origin, is suggestive of differential rotation in the magnetar, supporting the need for a rethinking of glitch theory for all neutron stars[10,11].

1E 2259+586 is a magnetar with a rotation period of about 7 s, with a spin-inferred surface dipolar magnetic field strength of 5.9×10^{13} G. Over 16 years of monitoring with the Rossi X-ray Timing Explorer, 1E 2259+586 has shown a very stable spin-down rate and pulsed flux, with the exception of two spin-up glitches in 2002 (ref. 9) and 2007 (ref. 12), and a small timing event in 2009 (ref. 12).The 2002 glitch was also accompanied by an increase in X-ray luminosity by a factor of 20 (ref. 9) and X-ray bursts[13], neither of which was seen in the 2007 glitch.

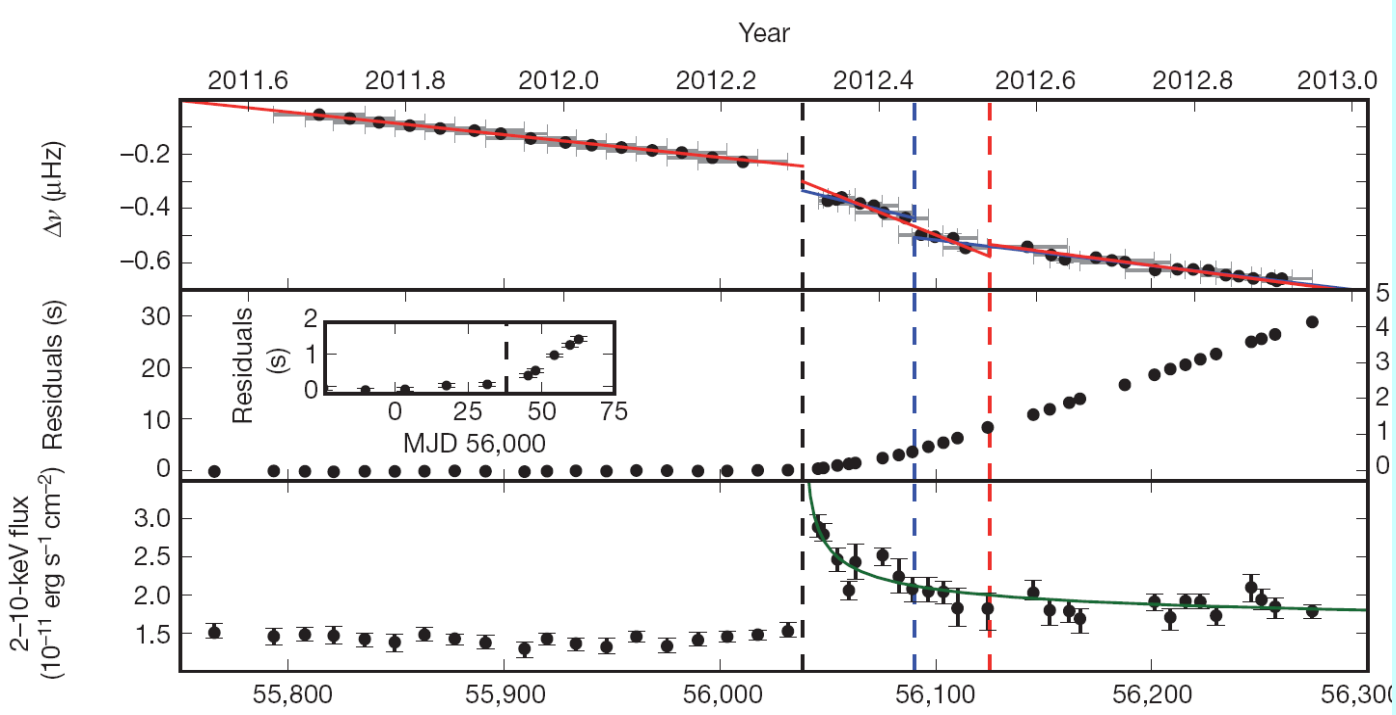
We began monitoring 1E 2259+586 with NASA's Swift X-ray Telescope[14] in July 2011. Observations were made every 2–3 weeks, with typical exposure times of 4 ks. From each observation, we obtained a pulse time-of-arrival (TOA) by folding the X-ray time series at the current pulse period and aligning this folded light curve with a high signal-to-noise template.

We fitted the pulse TOAs to a long-term timing model that keeps track of every rotation of the neutron star. This model predicts when the pulses should arrive at Earth, taking into account the pulsar rotation as well as astrometric terms. We compared the observed TOAs with the model predictions, and obtained best-fit parameters by χ^2 minimization, using the TEMPO2[15] software package. Until the observation on 14 April 2012 (modified Julian date, MJD 56,031.18), these TOAs were well fitted using only a frequency and first frequency derivative, as shown in Fig. 1.

The subsequent data, however, clearly were not predicted by this simple model. TOAs starting on 28 April 2012 (MJD 56,045.01) showed an apparently instantaneous change of the frequency—which we dub an 'anti-glitch'. On 21 April 2012 (MJD 56,038), consistent with the epoch of this sudden spin-down, a 36-ms hard X-ray burst was detected by the Fermi Gamma-ray Burst Monitor[16], with a fluence of about 6×10^{-8} erg cm^{-2} in the 10–1,000 keV range. No untriggered bursts from the Gamma-ray Burst Monitor were seen within three days of the observed burst[16]. Also, on 28 April 2012 (MJD 56,045.01), coincident with the nearest post-anti-glitch observation, we detected an increase in the 2–10-keV flux by a factor of 2.00 ± 0.09



Archibald et al. 2013, Nat.

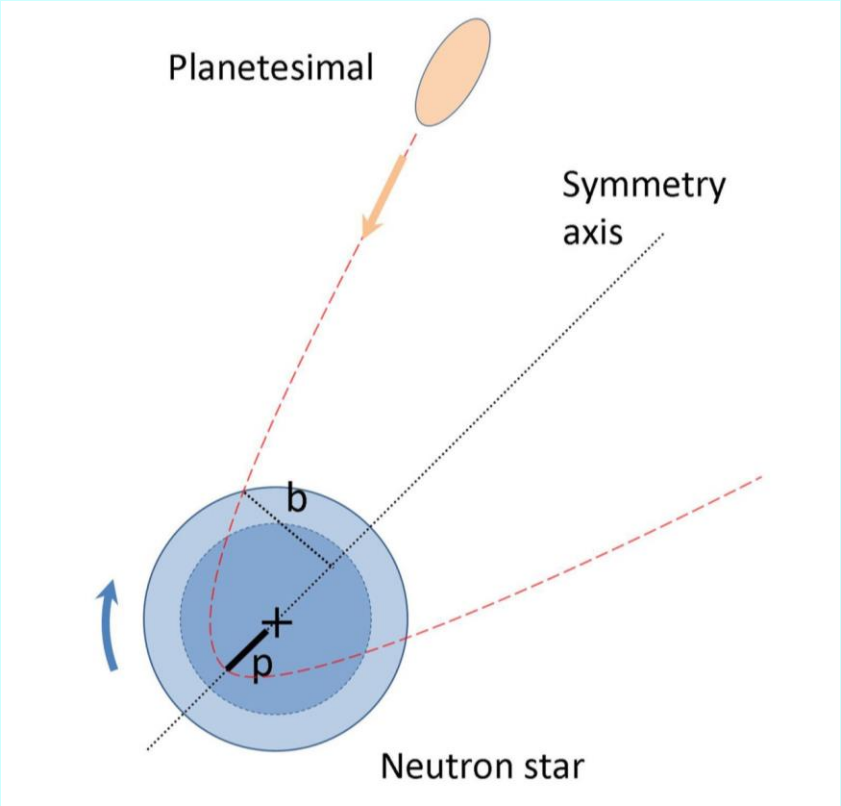


Model 1: $\begin{cases} \Delta\nu = -4.5(6)\times10^{-8}\,Hz \\ \Delta\dot{\nu} = -2.7(2)\times10^{-14}\,Hz/s \end{cases} + \begin{cases} \Delta\nu = 3.6(7)\times10^{-8}\,Hz \\ \Delta\dot{\nu} = -2.6(2)\times10^{-14}\,Hz/s \end{cases}$ x^2=45.4/44

Model 2: $\begin{cases} \Delta\nu = -9(1)\times10^{-8}\,Hz \\ \Delta\dot{\nu} = -1.3(4)\times10^{-14}\,Hz/s \end{cases} + \begin{cases} \Delta\nu = -6.8(8)\times10^{-8}\,Hz \\ \Delta\dot{\nu} = 1.1(4)\times10^{-14}\,Hz/s \end{cases}$ x^2=38.1/44

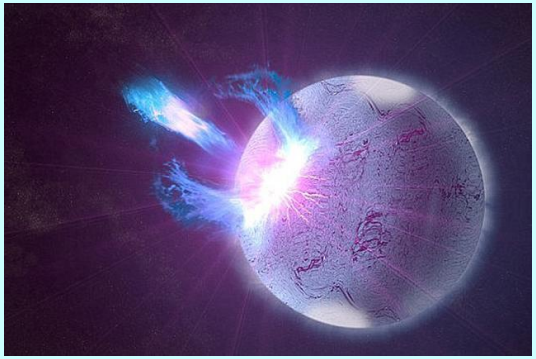
Archibald et al. 2013, Nat.

Our model: collision of a solid body with the magnetar

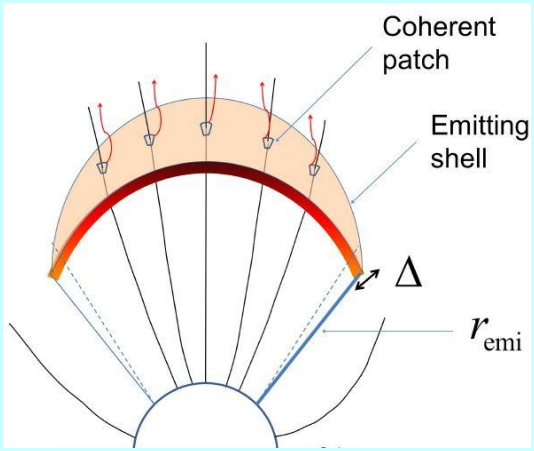


Huang & Geng, ApJ, 2014, 782, L20

（5）中子星-小行星碰撞产生FRB



http://tech.qq.com/a/20150111/007468.htm



Duration $\delta t = \Delta / c$

Emission volume $V_{\rm emi} \approx 4\pi f \Delta r_{\rm emi}^2$

Characteristic frequency $\nu_c = \gamma^3 \dfrac{3c}{4\pi r_{\rm emi}}$

Luminosity (Kashiyama et al. 2013)

$$L_{\rm tot} \approx \left(P_{\rm e} N_{\rm coh}^2\right) \times N_{\rm pat}$$

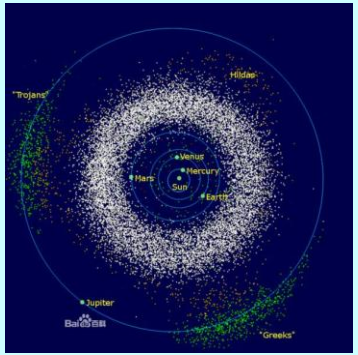
$$P_{\rm e} = 2\gamma^4 e^2 c / 3r_{\rm emi}^2,\ N_{\rm coh} \approx n_{\rm e} \times V_{\rm coh}$$

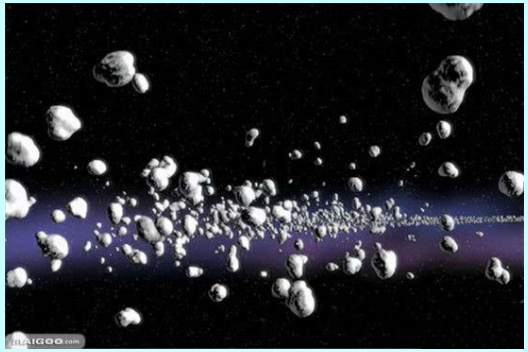
$$N_{\rm pat} \approx V_{\rm emi}/V_{\rm coh}$$

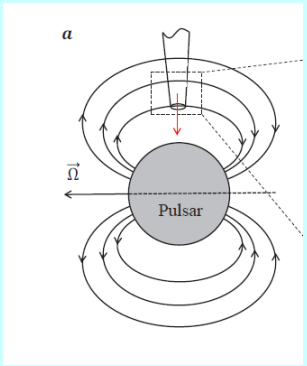
$$V_{\rm coh} = (4/\gamma^2)\, r_{\rm emi}^2 \times (c/\nu_{\rm c})$$

Geng & Huang, 2015, ***ApJ,*** 809:24 (arXiv:1502.05171)

Dai et al.：重复FRB 121102来自于脉冲星穿过小行星带





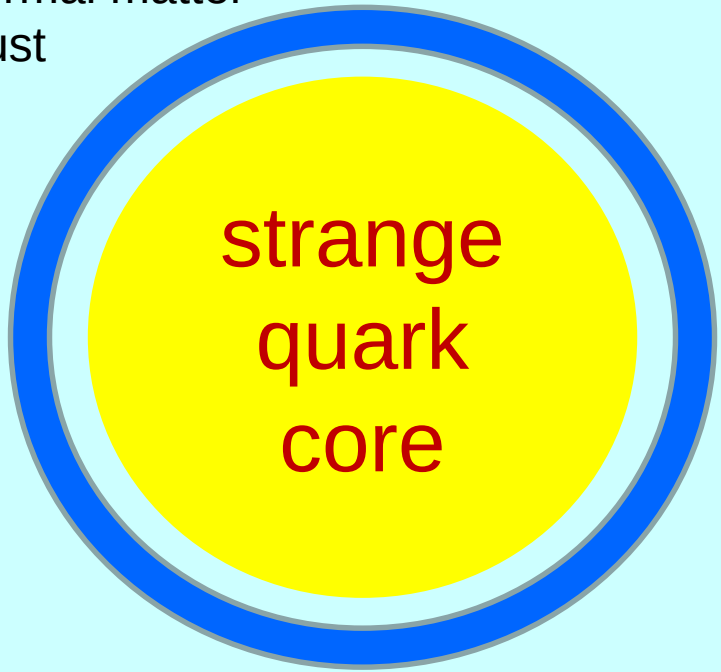


Description	Symbol	Dependence
Duration of an FRB♯	Δt	$1.58 m_{18}^{4/9}$ ms
Luminosity of an FRB♯	$L_{\rm tot}$	$\sim 2.63\times10^{40} m_{18}^{8/9}$ erg s^{-1}
Total energy release rate♯	$\dot{E}_{\rm G}$	$\sim 1.18\times10^{41} m_{18}^{5/9}$ erg s^{-1}
Duration distribution¶	$N(>T)$	$\propto T^{-\alpha_T}$, where $\alpha_T = 3(\beta-1)/4$
Luminosity function¶	$N(>L)$	$\propto L^{-\alpha_L}$, where $\alpha_L = 3(\beta-1)/8$
Repetitive rate‖	$\mathcal{R}_a$	$\sim 3.12(v_{*,7}/2)^{-1}(N_{a,10}/5)$ h^{-1}
Occurrence rate‖	$\mathcal{R}_b$	$\sim 2.0\times10^{-6}(v_{*,7}/2)(N_{\rm b,11}/2)(N_{*,8}/2.5)$ gal^{-1} yr^{-1}

Dai ZG, Wang JS, Wu XF, Huang YF, 2016, ApJ, 829, 27

（6）奇异夸克星壳层塌缩产生FRB

Normal matter crust



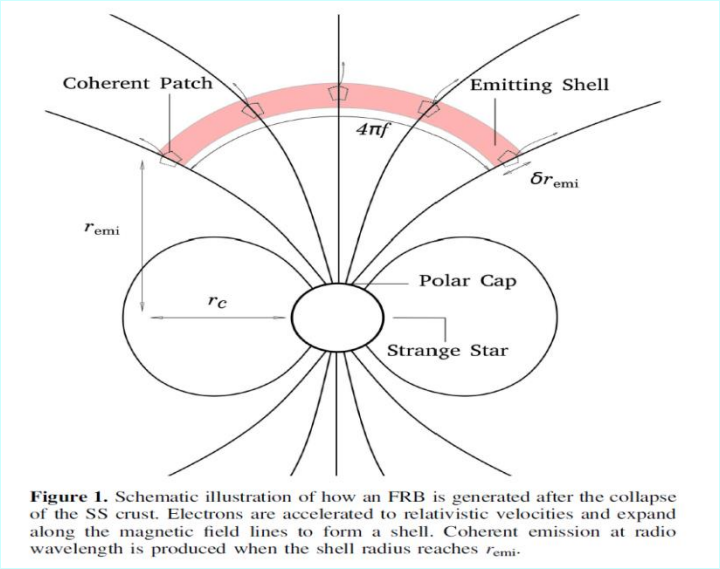
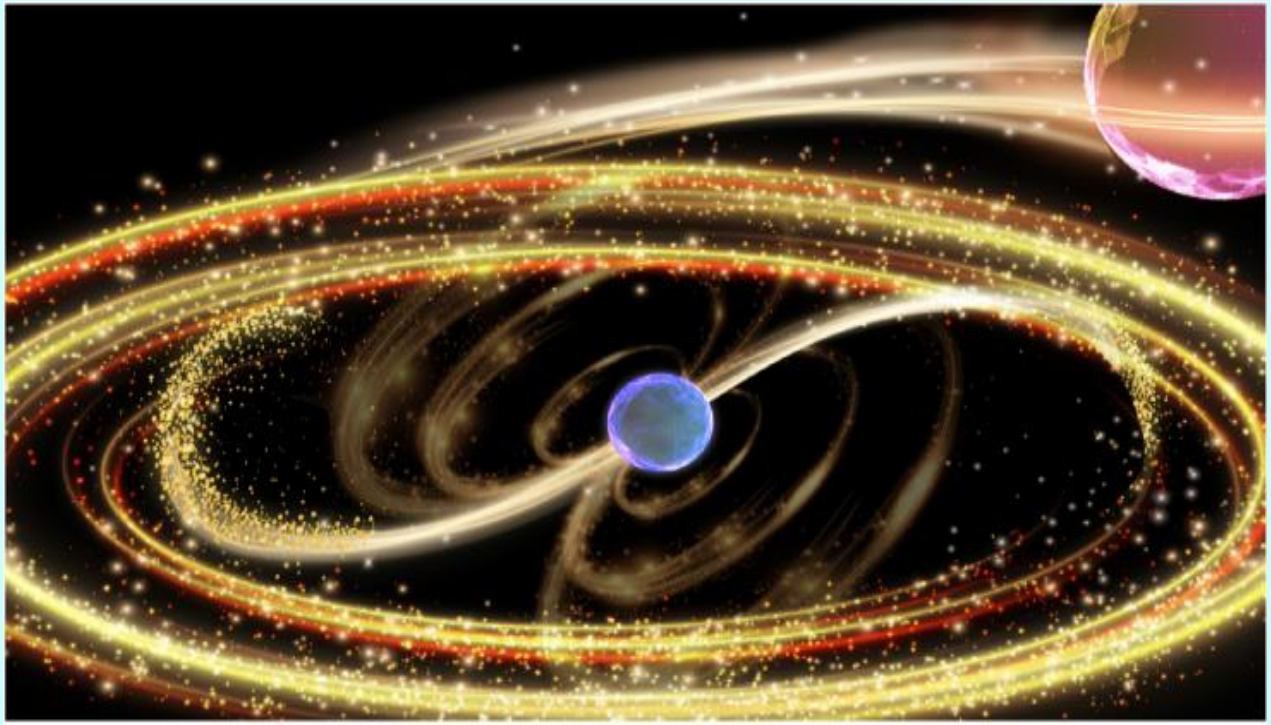


Figure 1. Schematic illustration of how an FRB is generated after the collapse of the SS crust. Electrons are accelerated to relativistic velocities and expand along the magnetic field lines to form a shell. Coherent emission at radio wavelength is produced when the shell radius reaches $r_{\rm emi}$.

$$L_{\rm total} \approx (P_{\rm e} N_{\rm coh}^2) \times N_{\rm pat}$$
$$\sim 2.6 \times 10^{42}\ {\rm erg\ s^{-1}}$$
$$\times f n_{\rm e,7}^2\, \delta t_{-3} \nu_{\rm FRB,9}^{-1/3} r_{\rm emi,10}^4 r_{\rm c,10}^{-4/3}.$$

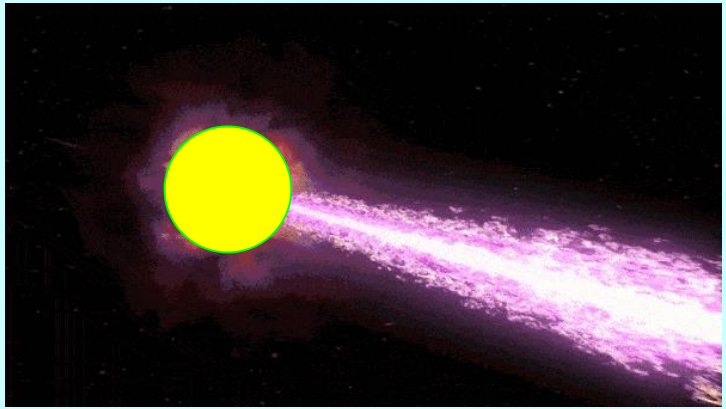
Zhang Y., Geng J.J., Huang Y.F., 2018, ***ApJ***, 858, 88 (6pp):
“Fast radio bursts from the collapse of strange star crusts”

奇异夸克星通过吸积还可产生周期性重复FRB



Geng Jinjun, Li Bing, Huang Yongfeng, 2021, The Innovation, 2, 100152

(7) GRBs from NS kicks



图片来自网络

Assume: NS kick is associated with the launch of a relativistic jet.

Momentum conservation: $E_{\rm flow}/c = M_{\rm NS}V_{\rm NS}$

Jet energy:

$$E_{\rm flow} = M_{\rm NS}V_{\rm NS}c = 8.3\times10^{51}\ {\rm ergs} \times \left(\frac{M_{\rm NS}}{1.4\ M_\odot}\right)\left(\frac{V_{\rm NS}}{1000\ {\rm km\ s^{-1}}}\right)$$

Radiation energy:

$$E_{\rm iso} = \frac{2\epsilon E_{\rm flow}}{1-\cos\theta} \approx 4\epsilon M_{\rm NS}V_{\rm NS}c\theta^{-2} = 3.3\times10^{53}\ {\rm ergs} \times \left(\frac{\epsilon}{0.1}\right)\left(\frac{\theta}{0.1}\right)^{-2}\left(\frac{M_{\rm NS}}{1.4M_\odot}\right)\left(\frac{V_{\rm NS}}{1000\ {\rm km\ s^{-1}}}\right)$$

$$\varepsilon = 0.3, \theta = 0.05 \Rightarrow E_{\rm iso} \sim 4\times10^{54}\,{\rm ergs}$$

Huang,Dai, Lu, Cheng, Wu, 2003, ApJ, 594, 919

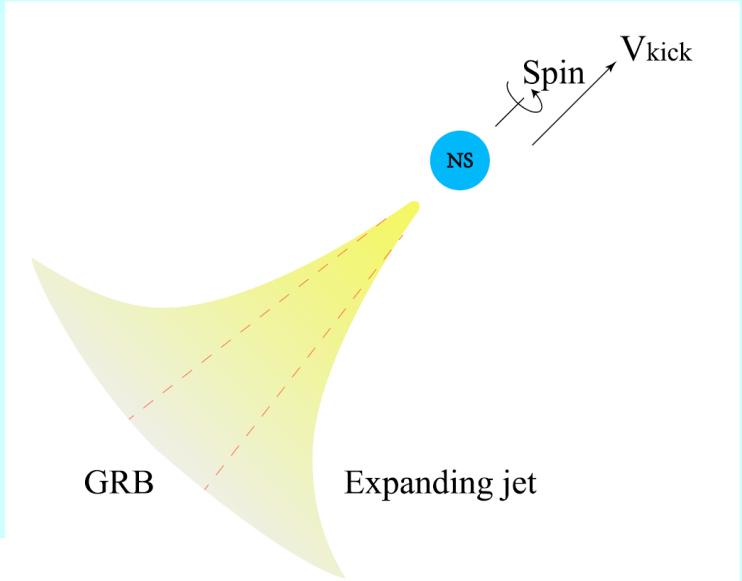
Monthly Notices of the ROYAL ASTRONOMICAL SOCIETY

MNRAS **509,** 4916–4922 (2022) https://doi.org/10.1093/mnras/stab3342
Advance Access publication 2021 November 23

Is the birth of PSR J0538+2817 accompanied by a gamma-ray burst?

Fan Xu [1], Jin-Jun Geng [1,2], Xu Wang,[1] Liang Li[3] and Yong-Feng Huang[1,4]★

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[2] *Purple Mountain Observatory, Chinese Academy of Sciences, Nanjing 210023, People's Republic of China*
[3] *ICRANet, Piazza della Repubblica 10, I-65122 Pescara, Italy*
[4] *Key Laboratory of Modern Astronomy and Astrophysics, Nanjing University, Nanjing 210023, People's Republic of China*



Yao et al. (2021, NA) reported the evidence for 3D spin-velocity alignment of PSR J0538 + 2817.

Table 1. Parameters of PSR J0538+2817.

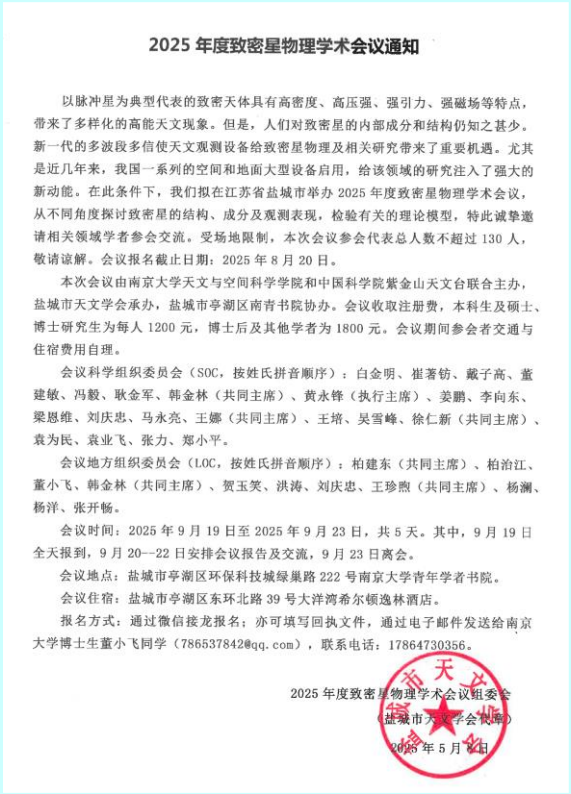
Observed parameters	Value	Ref.[a]	Derived parameters	Value	Ref.
RA (J2000)	$05^{h}\ 38^{m}\ 25\overset{s}{.}0623$	2	Characteristic age (kyr)	600	4
Dec. (J2000)	$28^{\circ}\ 17'\ 09''.1$	2	Kinematic age (kyr)	34.8 ± 0.4	4
Period, P (ms)	143.157 776 645(2)	1	DM distance, $D_{\rm DM}$ (kpc)	1.2	3
First derivative, $\dot{P}$ ($\times 10^{-15}$)	3.6681(1)	1	Parallax distance, D_{π} (kpc)	$1.30^{+0.22}_{-0.16}$	3
Dispersion measure (pc cm^{-3})	39.57	3	Transverse velocity, $V_{\perp}$ (km s^{-1})	357^{+59}_{-43}	3
μ_{α} (mas yr^{-1})	$-23.57^{+0.10}_{-0.10}$	3	3D velocity, $V_{\rm 3D}$ (km s^{-1})	407^{+79}_{-57}	4
μ_{δ} (mas yr^{-1})	$52.87^{+0.09}_{-0.10}$	3	Magnetic field (G)	7×10^{11}	1
π (mas)	$0.72^{+0.12}_{-0.09}$	3	Spin-down luminosity (erg s^{-1})	5×10^{34}	1

Note. [a] List of references: 1: Anderson et al. (1996); 2: Kramer et al. (2003); 3: Chatterjee et al. (2009); 4: Yao et al. (2021).

总结与建议

1. 必要的专业兴趣
2. 积极+主动+专注
3. 勇于创新
4. 随时记录想法，争取1/10的完成率

欢迎参加9月下旬的盐城致密星物理会议



2025 年度致密星物理学术会议通知

以脉冲星为典型代表的致密天体具有高密度、高压强、强引力、强磁场等特点，带来了多样化的高能天文现象。但是，人们对致密星的内部成分和结构仍知之甚少。新一代的多波段多信使天文观测设备给致密星物理及相关研究带来了重要机遇。尤其是近几年来，我国一系列的空间和地面大型设备启用，给该领域的研究注入了强大的新动能。在此条件下，我们拟在江苏省盐城市举办 2025 年度致密星物理学术会议，从不同角度探讨致密星的结构、成分及观测表现，检验有关的理论模型，特此诚挚邀请相关领域学者参会交流。受场地限制，本次会议参会代表总人数不超过 130 人，敬请谅解。会议报名截止日期：2025 年 8 月 20 日。

本次会议由南京大学天文与空间科学学院和中国科学院紫金山天文台联合主办，盐城市天文学会承办，盐城市亭湖区南青书院协办。会议收取注册费，本科生及硕士、博士研究生为每人 1200 元，博士后及其他学者为 1800 元。会议期间参会者交通与住宿费用自理。

会议科学组织委员会（SOC，按姓氏拼音顺序）：白金明、崔著钫、戴子高、董建敏、冯毅、耿金军、韩金林（共同主席）、黄永锋（执行主席）、姜鹏、李向东、梁恩维、刘庆忠、马永亮、王娜（共同主席）、王培、吴雪峰、徐仁新（共同主席）、袁为民、袁业飞、张力、郑小平。

会议地方组织委员会（LOC，按姓氏拼音顺序）：柏建东（共同主席）、柏泊江、董小飞、韩金林（共同主席）、贺玉笑、洪涛、刘庆忠、王珍煦（共同主席）、杨澜、杨洋、张开畅。

会议时间：2025 年 9 月 19 日至 2025 年 9 月 23 日，共 5 天。其中，9 月 19 日全天报到，9 月 20—22 日安排会议报告及交流，9 月 23 日离会。

会议地点：盐城市亭湖区环保科技城绿巢路 222 号南京大学青年学者书院。

会议住宿：盐城市亭湖区东环北路 39 号大洋湾希尔顿逸林酒店。

报名方式：通过微信接龙报名；亦可填写回执文件，通过电子邮件发送给南京大学博士生董小飞同学（786537842@qq.com），联系电话：17864730356。

2025 年度致密星物理学术会议组委会
（盐城市天文学会代章）
2025 年 5 月 1 日

时间： 9月19 – 23 日

报名： 微信接龙（8月20日截止）

联系： 董小飞
786537842@qq.com
17864730356

Thank You!

举办： 南京大学、紫台、
盐城天文学会、盐城南青书院

SOC： 白金明、崔著钫、戴子高、董建敏、冯毅、耿金军、韩金林（共同主席）、黄永锋（执行主席）、姜鹏、李向东、梁恩维、刘庆忠、马永亮、王娜（共同主席）、王培、吴雪峰、徐仁新（共同主席）、袁为民、袁业飞、张力、郑小平