An Accreting Stellar Binary Model for Active Periodic Fast Radio Bursts

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Abstract

In this work, we propose an accreting stellar binary model for understanding the active periodic fast radio bursts (FRBs). The system consists of a stellar compact object (CO) and a donor star (DS) companion in an eccentric orbit, where the DS fills its own Roche lobe near the periastron. The CO accretes the material from the DS and then drives relativistic magnetic blobs. The interaction between the magnetic blobs and the stellar wind of the DS produces a pair of shocks. We find that both the reverse shock and the forward shock are likely to produce FRBs via the synchrotron maser mechanism. We show that this system can in principle sufficiently produce highly active FRBs with a long lifetime, and also can naturally explain the periodicity and the duty cycle of the activity that appeared in FRBs 180916 and 121102. The radio nebula excited by the long-term injection of magnetic blobs into the surrounding environment may account for the associated persistent radio source. In addition, we discuss the possible multiwavelength counterparts of FRB 180916 in the context of this model. Finally, we encourage the search for FRBs in ultraluminous X-ray sources.

Unified Astronomy Thesaurus concepts: Radio transient sources (2008)

1. Introduction

Fast radio bursts (FRBs) are intense radio transients with an extremely short duration, and their physical origin is a mystery (Cordes & Chatterjee 2019; Petroff et al. 2019; Zhang 2020b; Xiao et al. 2021). Observationally, some FRBs show repeat bursts, but most bursts do not (Spitler et al. 2016; Fonseca et al. 2020; James et al. 2020). It is also a mystery whether all FRBs have the same origin (Caleb et al. 2018, 2019; Palaniswamy et al. 2018).

The CHIME/FRB Collaboration et al. (2020) found a 16.35 day periodicity in FRB 180916 with a 4.0 day phase window, implying a duty cycle $D \approx 0.24$ of activity. In addition, Rajwade et al. (2020) reported a possible 157 day periodicity in the FRB 121102 with a duty cycle $D \approx 0.56$. These are very important clues for studying the physical origin of FRBs. Such a periodicity may be explained either by the orbital period of the binary system (Dai & Zhong 2020; Gu et al. 2020; Ioka & Zhang 2020; Lyutikov et al. 2020; Mottez et al. 2020; Zhang 2020a) or the precession of the emitter (Katz 2020; Levin et al. 2020; Tong et al. 2020; Yang & Zou 2020; Zanazzi & Lai 2020).

Recently, the CHIME/FRB Collaboration et al. (2020) and Bochenek et al. (2020) reported a single FRB (FRB 200428) with two pulses in association with the active Galactic magnetar SGR 1935+2154, which is located at a distance of $\sim 9$ kpc (Zhong et al. 2020). Moreover, it was found that an X-ray burst was almost simultaneous with the FRB (Li et al. 2020; Mereghetti et al. 2020; Ridnaia et al. 2020; Tavani et al. 2020). But strangely, no more FRBs were observed during the active phase of this magnetar (Lin et al. 2020). In any case, this discovery shows that at least some of the FRBs are produced by magnetars (Lyubarsky 2014; Beloborodov 2017; Mezger et al. 2019). However, an important lesson learned from FRB 200428 is that though extragalactic analogs of Galactic magnetars could explain some of the FRB population, much more active sources are still required to explain the highly active periodic repeaters like FRBs 121102 and 180916 (CHIME/FRB Collaboration et al. 2020). If the active repeaters are also powered by magnetars, they must be produced by a type of rare and active magnetars not seen in the Milky Way (Lu et al. 2020; Mereghetti et al. 2020).

In addition, the observations of the environment of FRB 180916 found that it is 250 pc away from the nearest young stellar clump and suggested that its age may be $\sim 1$ Myr, which is in line with the hypothesis that high-mass X-ray binaries or gamma-ray binaries are the progenitors rather than young magnetars (Tendulkar et al. 2020). Moreover, Pastor-Marazuela et al. (2020) rule out the scenario in which companion winds cause FRB periodicity by using simultaneous Apertif and LOFAR data. Motivated by the observational results mentioned above, we propose an alternative model for understanding the highly active repeating FRBs with periodicity in this work. The model consists of a stellar compact object (CO) and a donor star (DS) with its filled Roche lobe, in which the CO can be a neutron star (NS) or a black hole (BH). We will show that this model can explain the actively repeating FRBs themselves and their periodicity behaviors, as for the cases FRBs 180916 and 121102.

Sridhar et al. (2021) also link periodic FRBs to accreting binaries. In Sridhar et al. (2021), the FRBs are produced by the interaction of intermittent jets (of high luminosity) with the quiescent jet (of low luminosity). However, in this work, we discuss the process by which magnetic blobs interact with the donor star’s wind to produce FRBs. This is the key difference between this work and their’s.

2. The Model

In this section, we propose a binary model for active repeating FRBs. The binary system consists of a stellar CO and a DS. The CO may be a BH or an NS. As illustrated in Figure 1, when the DS fills its Roche lobe, significant mass transfer will occur from the DS to the CO. Then an accretion disk will form around the CO. As we know, jets are widely present in accretion systems.
However, in addition to the usual continuous jets, the accretion process may also produce episodic jets. The energetic, collimating episodic jets have been observed in active galactic nuclei, stellar binaries, and protostars. For instance, in X-ray binaries, episodic jets are usually observed during their X-ray outbursts and intense radio flares (Fender et al. 2004; Zhang & Yu 2015). Practical models and numerical simulations suggested that the episodic jets may be driven by a magnetic instability in the accretion disk (Yuan et al. 2009; Yuan & Zhang 2012; Zhao et al. 2020). Because of shear and turbulent motion of the accretion flow, a flux rope system is expected to form near the disk. The energy is accumulated and stored in the system until a threshold is reached, then the system loses its equilibrium and the energy is released in a catastrophic way, i.e., by ejecting episodic magnetic blobs. By assuming the accretion flow is advection dominated, the available isotropic free magnetic energy of one blob is (Yuan & Zhang 2012)

$$E_f \sim 10^{42} f_{b, -3}^{-1} \alpha^{-1} \beta^{-1} \frac{M_u}{10^{-7} M_\odot \text{yr}^{-1}} \left(\frac{M}{3 M_\odot}\right) \beta^2 f_w^2 \text{erg}. \quad (1)$$

Here, we consider the magnetic blobs is ejected in a collimated angle, and $f_b = (1 - \cos \theta)/2$ is the beaming factor, where $\theta$ is the jet opening angle. $\alpha$ is the viscous parameter, and $\beta$ is the ratio of the magnetic pressure over the total pressure in the accretion disk. We adopt $\alpha = 0.01$ and $\beta = 0.1$ as typical values. $M_u$ is the mass accretion rate, $M$ is the mass of the CO, and $\hat{r}$ is the radius at which the magnetic blobs are formed in units of $2GM/c^2$. Here and hereafter, we employ the short-hand notation $q_s = q/10^6$ in cgs units.

The binary system is surrounded by the stellar wind from the DS. Therefore, we suggest that the interaction between the magnetic blobs and the stellar wind could lead to shock formation, and then this shock process would produce powerful coherent radiation (FRBs) through the synchrotron maser mechanism analogous to that in flaring magnetar models (Lyubarsky 2014; Beloborodov 2017; Metzger et al. 2019).

Taking FRB 180916 as a template, the typical energy of the bursts is $\sim 10^{37} - 10^{38}$ erg (CHIME/FRB Collaboration et al. 2020). In fact, it turns out that the typical energies of FRB 121102 and of FRB 180916 are similar, as found by recent observations (Pastor-Marazuela et al. 2020; Li et al. 2021). One will also see from Equation (17), $E_f = 10^{42} \text{erg}$ is reasonable to power FRB 180916 and 121102. According to Equation (1), a mass accretion rate of $\sim 10^{-5} M_\odot \text{yr}^{-1}$ is required. Note that this high accretion rate $M_u \sim 10^{-3} M_\odot \text{yr}^{-1}$ can occur in high-mass Roche lobe filling binaries, such as SS433, which is accreting at a rate of $\sim 10^{-4} M_\odot \text{yr}^{-1}$ (Fabrika 2004). Moreover, Wiktowicz et al. (2015) showed that the Roche lobe overflow rate can be up to $10^{-3} M_\odot \text{yr}^{-1}$ if the DS is evolving in the Hertzprung gap.

The wind emerged from the DS is filled around the CO, and its density distribution $n_w$ can be estimated as

$$n_w = \frac{\dot{M}_w}{4\pi m_w R_w^2 v_w}, \quad R \simeq (r^2 + a^2)^{1/2}$$

where $\dot{M}_w$ is the wind mass-loss rate of the DS, and $v_w$ is the speed of the wind. Variable $a$ is the semimajor axis of the binary, and $r$ is the distance from the CO. For $r \ll a$, we have

$$n_w \approx 1.0 \times 10^3 \dot{m}_w (M_\odot \text{yr}^{-1}), \quad \beta_w = v_w/c, \quad \text{and } c \text{ is the speed of light.}$$

We adopt $\beta_w = 0.01$ as a typical value for massive stars, and we adopt $\dot{m}_w = 10^{-11}$ for self-consistency in the working model. The choice of $a = 10^3$ cm is motivated by the analysis for FRB 180916 in Section 3.

Therefore, according to Sari & Piran (1995), the Lorentz factor of the blast wave during the early reverse shock crossing phase is $\Gamma = (n_\infty / n_g)^{1/4}$, where $n_g = E_f / (4\pi m_e c^2 a^2)$ is the co-moving density of the ejecta at a radius $r$, $\Gamma_{ej}$ is the initial Lorentz factor of the ejecta, and $\delta t$ is the duration of the central energy activity. Then we have

$$\Gamma (r < r_{dec}) = \left( \frac{E_f \beta_w}{4 M_u c^2 \delta t} \right)^{1/4} \left( \frac{r}{a} \right)^{-1/2}$$

Figure 1. Schematic illustration of the model in this work. (a) The DS fills its Roche lobe when orbiting near the periastron, and mass transfer to the CO occurs. (b) The transfer of mass leads to the formation of an accretion disk around the CO. Magnetic blobs are ejected from the accretion disk due to magnetic instability. The magnetic pressure over the total pressure in the accretion disk is $\sim 10^{-1}$, i.e., $\sigma_0$ is the initial magnetization parameter of the blobs. Then the blobs interact with the stellar wind, which is driven by the DS and immerses itself around the binary system, to induce shocks powering FRBs by the synchrotron maser mechanism. (c) The interactions of the magnetic blobs and the stellar wind produce FRBs. Long-term energy injection into the surrounding medium produces a nebula that powers the persistent radio source associated with the FRB.
where
\[
r_{\text{decl.}} \simeq 1.1 \times 10^{12} E_{f,4}^{1/4} m_{w,-11}^{1/4} \beta_{w,-2}^{3/4} \delta_{3}^{1/4} a_{13}^{1/2} \text{ cm} \tag{5}
\]
is the deceleration radius. One sees that it satisfies \( r_{\text{decl.}} \ll a \). Therefore, the Lorentz factor at the deceleration radius is
\[
\Gamma(r_{\text{decl.}}) \simeq 136 E_{f,4}^{1/8} m_{w,-11}^{1/8} \beta_{w,-2}^{3/8} \delta_{3}^{3/8} a_{13}^{1/4} . \tag{6}
\]
That said, we get the co-moving density in the ejecta at \( r_{\text{decl.}} \)
\[
n_{ej} \simeq 1.4 \times 10^{8} \Gamma_{ej}^{-2} E_{f,4}^{1/2} m_{w,-11}^{1/2} \beta_{w,-2}^{-1/2} \delta_{3}^{-3/2} a_{13}^{-1} \text{ cm}^{-3}. \tag{7}
\]
Combining Equation (3) with Equation (7), one gets that the reverse shock will become relativistic, which is the case considered in this work, if
\[
\Gamma_{ej} \gtrsim 194 E_{f,4}^{1/8} m_{w,-11}^{1/8} \beta_{w,-2}^{-3/8} \delta_{3}^{-3} a_{13}^{1/4}, \tag{8}
\]
according to the condition \( \Gamma_{ej} > f = n_{ej}/n_{w} \) (Sari & Piran 1995).
In the context of this model, the reverse shock is ultrarelativistic only for \( \Gamma_{ej} > 10^{2} \), and it is mildly relativistic for \( \Gamma_{ej} \sim 200 \), which is the easier case to achieve. Can accreting stellar-mass COs produce ejecta with \( \Gamma_{ej} > 100 \)? As argued in Sridhar et al. (2021), it cannot be excluded that highly super-Eddington systems are capable for brief periods of generating outflows with such a large \( \Gamma_{ej} \). Especially for ejecta with a large initial magnetization \( \sigma_{0} \gg 1 \), which is the case in this work: it can be accelerated to a high Lorentz factor \( \Gamma_{ej} \sim \sigma_{0}^{2} \) (Drenkhahn 2002).
Therefore, assuming the reverse shock is mildly relativistic, then Equation (7) is reduced to
\[
n_{ej} \simeq 3.7 \times 10^{7} E_{f,4}^{1/4} m_{w,-11}^{3/4} \beta_{w,-2}^{-3/4} \delta_{3}^{-3/4} a_{13}^{3/2} \text{ cm}^{-3}. \tag{9}
\]
However, if the reverse shock is ultrarelativistic, this is the upper limit for \( n_{ej} \).
In previous studies that discussed the production of FRBs by magnetized shocks, both forward and reverse shocks have been involved (Lyubarsky 2014; Beloborodov 2017; Waxman 2017; Metzger et al. 2019; Beloborodov 2020). In this work, we consider the synchrotron maser emission produced by both the reverse shock and forward shock. The equipartition of energy in the reverse and forward shock is assumed in this work for simplicity. The theory of the synchrotron maser instability has been developed both for isotropic (Sagiv & Waxman 2002; Gruzinov & Waxman 2019) and ring-like (Lyubarsky 2006) particle distributions in momentum space; also see a review Lyubarsky (2021). In our scenario, we consider weakly magnetized ejecta at the deceleration radius and a stellar wind that is itself also weakly magnetized (Weber & Davis 1967; Sakurai 1985; Ignace et al. 1998; ud-Doula & Owocki 2002; Puls et al. 2008; Harvey-Smith et al. 2010). For the weakly magnetized plasma, i.e., \( \sigma \ll 1 \), the characteristic frequency measured in the rest frame of the synchrotron maser may be given by Sagiv & Waxman (2002), Waxman (2017), Gruzinov & Waxman (2019), and Lyubarsky (2021):
\[
\nu_{c} \approx \sigma^{-1/4} \nu_{p}, \tag{10}
\]
where \( \sigma \) is the magnetization parameter of the plasma with magnetic field \( B \) and number density \( n \),
\[
\sigma = \frac{B^2}{4\pi n m_p c^2}. \tag{11}
\]
The magnetization \( \sigma \) is a very uncertain and difficult parameter to know, so it can be regarded as a free parameter in principle in this work. As an example, in the case of gamma-ray bursts (GRBs), after the highly magnetized jets are accelerated through the magnetic dissipation process, the remaining magnetization \( \sigma_{ej} \) may be small, \( \sim 0.1 \), which can naturally account for the observed high radiative efficiency of most GRBs (Zhang & Yan 2011). The magnetization \( \sigma_{w} \) of the DS’s wind is more difficult to determine, but perhaps we can make a very rough estimation. Suppose the DS has a surface magnetic field strength \( B \sim 10^{5} \text{ G} \), a radius \( R \sim 10^{2} \text{ cm} \), and a mass-loss rate \( \dot{m}_{w} \sim 10^{-10} \) (the range of \( \dot{m}_{w} \) covered in this work is \( 10^{-9} \sim 10^{-11} \)). One gets the magnetization of the wind, at \( r_{\text{decl.}} \), \( \sigma_{w} \sim 10^{-3} \) (Lamers & Cassinelli 1999; Harvey-Smith et al. 2010). In addition, it can be seen from Equation (10) that the result is only slightly dependent on \( \sigma \), so it is not unreasonable to directly use a rough value of \( \sigma \) for orders of estimation within the range of parameters considered in this model.
Therefore, we adopt \( \sigma_{ej} = 0.1 \) for the ejector at \( r_{\text{decl.}} \) and \( \sigma_{w} = 10^{-3} \) for the DS’s wind, respectively. In the context of this work, the material shocked by the reverse shock and forward shock are both baryon-dominated plasma. Therefore, the plasma frequency \( \nu_{p} \) is determined by Sagiv & Waxman (2002):
\[
\nu_{p} = \left( \frac{n_{e} e^{2}}{\pi \gamma_{e} m_{e}} + \frac{n_{p} e^{2}}{\pi \gamma_{p} m_{p}} \right)^{1/2}, \tag{12}
\]
where \( \gamma_{e} \) and \( \gamma_{p} \) are the Lorentz factors of the electrons and the protons, respectively. Assuming that the electrons are near equilibrium with the protons with \( \gamma_{e} m_{e} \approx \gamma_{p} m_{p} \) in the downstream, we get the characteristic frequency of the synchrotron maser produced by the reverse shock and the forward shock, respectively, measured in the observed frame,
\[
\nu_{pk} \approx \left\{ \begin{align*}
\frac{8 n_{ej} e^{2}}{7 \pi m_{p}}^{1/2} & \quad \Gamma_{ej}^{-1/4} \nu_{pk} \\
0.88 \sigma_{ej}^{-1/4} E_{f,4}^{1/4} m_{w,-11}^{1/4} \beta_{w,-2}^{-1/4} \delta_{3}^{1/4} a_{13}^{-1/2} \text{ GHz}, & \text{ for reverse shock} \\
\frac{8 n_{w} e^{2}}{7 \pi m_{p}}^{1/2} & \quad \Gamma_{ej}^{-1/4} \nu_{pk} \\
0.01 \sigma_{w}^{-1/4} E_{f,4}^{1/4} m_{w,-11}^{3/8} \beta_{w,-2}^{-3/8} \delta_{3}^{-3/4} a_{13}^{-3/4} \text{ GHz}, & \text{ for forward shock}.
\end{align*} \right. \tag{13}
\]
One sees that the maser emission is determined by protons, in contrast to the case for the pair plasma.
The optical depth due to free—free absorption is estimated as
\[
\tau_{\text{ff}}(\nu) \sim \frac{4 e^{6}}{3 k_{B} m_{e} c} \left( \frac{2 \pi}{3 k_{B} m_{e}} \right)^{1/2} \left( \frac{\nu_{\text{decl.}}}{\nu_{p}} \right)^{-3/2} \nu^{-2} \sim 10^{-8} \frac{\nu_{c} E_{f,4}^{1/4} m_{w,-11}^{3/8} \beta_{w,-2}^{-3/8} \delta_{3}^{-3/4} a_{13}^{-3/4}}{7 \pi^{2} T_{w,d}^{3/2} \nu_{0}^{2}}, \tag{14}
\]
where \( k_B \) is the Boltzmann constant, \( g_e \) is the mean Gaunt factor, and \( T_{\text{obs}} = 10^4 \text{ K} \) is the temperature of the wind of the DS. For \( \nu / k_B T_{\text{obs}} \ll 1 \), the optical depth to Thomson scattering, \( \tau \approx (\sqrt{3} / \pi) \ln(2.2 k_B T_{\text{obs}} / \nu) \), and one has \( g_e (\nu, T_{\text{obs}}) \approx 7 \).

The optical depth to Thomson scattering, \( \tau \approx \sigma T_{\text{obs}} n_e (\nu_{\text{det}}) \sim 10^{-3} \), is very small. However, due to the extremely high brightness temperature of FRBs, the induced scattering process becomes important (Lyubarsky 2008; Lyubarsky & Ostrovska 2016). Because the electrons downstream of the shocks are ultrarelativistic, the induced Compton scattering caused by them can be negligible (Wilson & Rees 1978). Therefore, the optical depth due to induced Compton scattering is mainly contributed by the DS’s wind, which is given by

\[
\tau_{\text{IC}} \sim \frac{1}{10} \frac{3 \sigma T_{\text{obs}} n_e (\nu_{\text{det}}) E_{\text{FRB}} (\nu_{\text{det}})}{32 \pi m_e r_{\text{dec}}^2 \nu_{\text{pk}}^3} \\
\approx \begin{cases} 
77 \sigma_T \gamma_{\text{ic}} - \frac{3}{4} \nu_{\text{pk}}^4 m_{\text{w}, -11}^{3/4} \beta_{\text{s}, -2}^{3/4} \delta_{\text{t}}, a_{\gamma, 3}^{-3/2} & \text{for reverse shock} \\
10^7 \sigma_T \gamma_{\text{ic}} \nu_{\text{pk}}^4 m_{\text{w}, -11}^{3/4} \beta_{\text{s}, -2}^{3/4} \delta_{\text{t}}, a_{\gamma, 3}^{-3/2} & \text{for forward shock.}
\end{cases}
\]

(15)

It can be seen that the emission at \( \nu_{\text{pk}} \) cannot be transmitted freely due to the induced Compton scattering. Here, \( E_{\text{FRB}} \sim 10^{50} \text{ ergs} \) is applied, and \( \epsilon \) is the efficiency of the synchrotron maser around \( \nu = \nu_{\text{pk}} \). The radiation efficiency \( \epsilon \) of the synchrotron maser is highly uncertain. PIC simulations show that the efficiency \( \epsilon \) depends on the magnetization and temperature for pair plasma. Plotnikov & Sironi (2019) found that \( \epsilon \approx 10^{-2} \) for an upstream magnetization \( \sigma \sim 0.1 \), and it decreases for \( \sigma > 1 \). When the upstream plasma is nonrelativistic, \( \epsilon \) would be independent of the temperature of the plasma for \( \sigma > 1 \) (Babul & Sironi 2020), but its behavior remains unclear for \( \sigma \ll 1 \). However, for a proton plasma, which is the case considered in this work, the efficiency \( \epsilon \) is even less known and has yet to be investigated in detail (Lyubarsky 2021). For the model to be consistent with observations, the efficiency should not be too low to account for the FRBs’ energetics, so we assume the efficiency to be \( \epsilon \approx 10^{-3} \) in this work.

According to Equation (15), we have \( \tau_{\text{IC}} (\nu) = \tau_{\text{IC}} (\nu_{\text{det}}) (\nu / \nu_{\text{det}})^{(3 + \sigma)} \) if we assume the spectrum has the form \( E_{\text{FRB}} \propto \nu^{-3} \). As a result, the observed peak frequency of the radiated spectrum moves to \( \nu_{\text{m}} \) where \( \tau_{\text{IC}} (\nu_{\text{m}}) = 3 \).

\[
\nu_{\text{m}} \sim
\begin{cases}
0.88 \times 10^{10} \nu_{\text{pk}}^{+3/2} \nu_{\text{det}}^{-1/3} & \text{for reverse shock} \\
0.01 \times 10^{10} \nu_{\text{pk}}^{+3/2} \nu_{\text{det}}^{-1/3} & \text{for forward shock}
\end{cases}
\]

One sees that \( \nu_{\text{m}} \) can account for the emission of FRBs around a GHz. In addition, during the reverse shock crossing phase \( t < \delta t \), this peak frequency will evolve with the shock, decelerating as \( \nu_{\text{m}} \propto t^{-1/3} \) for the reverse shock, and as \( \nu_{\text{m}} \propto t^{-4/3} \) for the forward shock. This temporally decreasing peak frequency may explain the observed downward drifting frequency structure in the subpulses of some repeating FRBs (CHIME/FRB Collaboration et al. 2019; Hessels et al. 2019). The drift rate depends on the specific value of \( s \).

According to the latest observations for FRB 121102, its typical energy is \( \sim 10^{27} - 10^{29} \text{ erg} \) (Li et al. 2021), and \( a \sim 10^{3.5} \) cm (see Section 3). If \( \dot{m}_{\text{w}} = 10^{-9} \), based on Equation (17), we have \( \nu_{\text{m}} \approx 3.83 \nu_{\text{det}}^{15/36} \nu_{\text{det}}^{-5/6} \) GHz for the reverse shock and \( \nu_{\text{m}} \approx 1.03 \nu_{\text{det}}^{11/16} \nu_{\text{det}}^{-11/16} \) GHz for the forward shock, adopting \( s = 1.5 \) (Macquart et al. 2019). This is consistent with the fact that FRB 121102 has few detections in the sub-GHz bands (Houben et al. 2019; Josephy et al. 2019), which may be because the DS in the case of FRB 121102 has a stronger stellar wind \( \dot{m}_{\text{w}} \approx 10^{-9} \) than in the case of FRB 180916, \( \dot{m}_{\text{w}} \approx 10^{-11} \).

By definition, the radiation energy around \( \nu_{\text{pk}} \) is \( E_{\text{FRB}} (\nu_{\text{pk}}) = \epsilon E_{\gamma} \); then, the radiation energy around \( \nu_{\text{m}} \) is given by \( E_{\text{FRB}} (\nu_{\text{m}}) = (\nu_{\text{m}} / \nu_{\text{pk}})^{-3} E_{\text{FRB}} (\nu_{\text{pk}}) \), namely,

\[
E_{\text{FRB}, 39} (\nu_{\text{m}}) =
\begin{cases}
10^{43/3} \nu_{\text{m}, -11}^{-2} \nu_{\text{det}, 3}^{2} & \text{for reverse shock} \\
10^{44/3} \nu_{\text{m}, -11}^{-2} \nu_{\text{det}, 3}^{2} & \text{for forward shock}
\end{cases}
\]

Since the maser energy is mainly concentrated around \( \nu_{\text{pk}} \), in fact, the efficiency near the observed frequency may be much lower than \( 10^{-3} \), which can be clearly seen from the above equation. That is, the efficiency near the observed frequency is equal to \( \epsilon \) multiplied by a factor of less than unity, such as \( 5 \times 10^{-3} \) and \( 10^{-7/3} \) for the reverse shock and forward shock, respectively. In this way, the low radiation efficiency \( \sim 10^{-5} \) of the radio bursts at the observed frequency in the observations of the Galactic burst FRB 200428 (Margalit et al. 2020) can be understood. Moreover, one sees that the observed energy of the synchrotron maser from the reverse shock is significantly greater than that from the forward shock. It indicates that the observed energy may exhibit a bimodal distribution in a single repeating FRB. The energetic bursts come from the reverse shock, while the less energetic bursts come from the forward shock. This may provide an explanation for the bimodal burst energy distribution in FRB 121102 found by FAST recently (Li et al. 2021). A more detailed analysis will be presented elsewhere.

In this model, according to Equation (17), the burst energy may be adjusted mainly by the model parameters such as \( E_{\gamma} \), \( \dot{m}_{\text{w}} \), and \( a \). For different FRB sources, there may be different model parameters that result in different observed energies. Moreover, FRBs may be generated by both the reverse and forward shocks and observed by telescopes with different frequencies and thresholds, so it is very easy to generate diverse distributions of energy. In order to compare directly with observations, one needs to do detailed modeling, but it is beyond the scope of this work. One thing needs special attention, however: Equation (16) and Equation (17) are only valid when \( \tau_{\text{IC}} > 3 \). If \( \tau_{\text{IC}} \ll 3 \), Equation (16) and Equation (17) will degenerate to \( \nu_{\text{m}} = \nu_{\text{pk}} \) and \( E_{\text{FRB}, 39} = 1 \).

3. Constraints from the Periodicity and Duty Cycle

Periodicities have been observed in some X-ray sources, and mechanisms have been proposed to explain this behavior, mainly including orbital period modulation (Strohmayer 2009).
and disk/jet precession (Begelman et al. 2006; Foster et al. 2010; also see a review in Kaaret et al. 2017).

For periodic FRBs, under the framework of our model, the periodicity might be explained by the precession of the jets (Katz 2020, 2021; Sridhar et al. 2021). However, FRB 180916 shows that the signals are concentrated in a narrow active window with a duty cycle $D \approx 0.24$, which is very different from the cases of X-ray binaries (Strohmayer 2009; Foster et al. 2010). As discussed in the previous section, FRB emission is produced when the DS fills its Roche lobe. Therefore, we introduce an eccentric orbit modulation mechanism to explain the narrow duty cycle. The periodicity may be explained by periodic orbital motion of the binary, and the duty cycle may be explained by the eccentricity of the orbit. It can be naturally understood by this picture: the DS fills its Roche lobe when it is near the periastron, where the FRB emission is produced, if the orbit of the binary is eccentric. After the DS moves away from the periastron, the process of producing FRB emission stops due to a significant decrease in the accretion rate.

Define the mass ratio $q = M/M_\odot$, where $M_\odot$ is the mass of the CO, and $M$ is the mass of the DS. Then the orbital period of the binary is

$$ T = 2\pi q^{1/2}(1 + q)^{-1/2}a^{3/2}(GM)^{-1/2} \quad (18) $$

where $G$ is the gravitational constant, and $a$ is the semimajor axis. The effective radius $R_{L,0}$ of the Roche lobe of the DS at the periastron can be estimated as (Eggleton 1983)

$$ R_{L,0} \approx \frac{0.49q^{2/3}}{a(1 - e)} \left( \frac{0.6q^{2/3} + \ln(1 + q^{1/3})}{3} \right) \approx \chi, \quad (19) $$

where $e$ is the orbital eccentricity. With the assumption that the DS has filled its Roche lobe, its average density is

$$ \bar{\rho} = \frac{3M}{4\pi f_{RL}^3 R_{L,0}^3}, \quad (20) $$

where $f_{RL} = R/R_{L,0} \geq 1$ is the Roche lobe filling factor, and $R$ is the radius of the DS. We expect that $f_{RL}$ is just slightly greater than 1, i.e., $f_{RL} - 1 \ll 1$. Combining Equations (18) and (20), one gets

$$ T \approx (3\pi)^{1/2}f_{\rho}^{1/2}(1 - e)^{-3/2}(G\bar{\rho})^{-1/2}, \quad (21) $$

where $f_{\rho} = q^{1/2}(1 + q)^{-1/2}\chi^{-3/2}$. Note that $f_{\rho} \in (1.4, 1.8)$ for $q \geq 0.1$. For simplicity, we take $f_{\rho} = 1.5$ because the parameter range of interest in this work is $q \gtrsim 0.1$. Thus, one can simply deduce that the average density of the DS is

$$ \bar{\rho} \approx (1 - e)^{-3}(\frac{T}{0.21 \text{ day}})^2 \rho_{0,\odot}. \quad (22) $$

where $\rho_{0,\odot} \approx 1.4 \text{ g cm}^{-3}$ is the current average density of the Sun. It is worth noting that Equation (22) is a generalization of Equation (4.10) in Frank et al. (2002), $\rho \propto T^{-2}$; namely, it is generalized from the case of a circular orbit to the case of an eccentric orbit. Next, we discuss the constraint on $e$ from the observed duty cycle $D$.

Without loss of generality, taking the CO as the reference, the DS moves on an elliptic orbit with respect to the CO. The distance from the CO to the DS is $r(\theta) = a(1 - e^2)/(1 + e \cos \theta)$. The DS is at periastron when $\theta = 0$, where $\theta$ is the angle between the vector diameter from the CO to the DS and the polar axis in the polar coordinate system. For the DS at different positions, its Roche lobe radius is denoted as $R_{L,0} = \chi r(\theta)$. Assume that when $\theta = \pm \alpha (\alpha < \pi)$, the DS can just fill its Roche lobe, i.e., $R_{L,0} = R$. Then, we have

$$ \frac{1 + e}{1 + e \cos \alpha} = f_{R \text{ e f l}}, \quad (23) $$

Therefore, according to Kepler’s second law, the duty cycle of activity can be calculated as $D = \Delta S/\Delta t$, where $\Delta S$ is the area swept out by $r(\theta)$ from $\alpha$ to $\alpha’$, and $\Delta t$ is the total area enclosed by the DS’s elliptical orbit. That is,

$$ D = \frac{1}{2\pi} (1 - e^2)^{3/2} \int_{-\alpha}^{\alpha} \frac{d\theta}{(1 + e \cos \theta)^2}, \quad (24) $$

where, by letting $\lambda = 1 - f_{RL}^{-1}$,

$$ \alpha = \begin{cases} \pi, & \lambda(1 + e)/(2e) \geq 1 \\ 2 \arcsin \sqrt{\lambda(1 + e)/(2e)}, & \lambda(1 + e)/(2e) < 1 \end{cases} \quad (25) $$

For $\lambda(1 + e)/(2e) \geq 1$, it means that the DS fills its Roche lobe throughout the cycle, thus $D = 1$. Therefore, based on the observed duty cycle, we can discuss the constraint to $e$. Figure 2 shows the duty cycle as a function of $e$ for $\lambda = 0.24$ (FRB 180916) and $D = 0.56$ (FRB 121102), respectively. Note that $\lambda = 1 - f_{RL}^{-1} = (R - R_{L,0})/R$ describes the degree of the Roche lobe overfilling of the DS at periastron. One naturally expects $\lambda \ll 1$, otherwise the Roche lobe overflow would be violent. By adopting $\lambda \ll 0.1$, we have $e \leq 0.22$ for FRB 180916 and $e \leq 0.08$ for FRB 121102, which are also shown in Figure 2. It can be seen that the required orbital eccentricity is not large.

Now we can discuss what kind of DS is needed to explain the periodicity of FRBs 180916 and 121102. For FRB 180916, $T = 16.35$ days and $e \leq 0.22$, and one gets $\bar{\rho} \sim 10^{-3.5}\rho_{0,\odot}$. According to Equation (22). Similarly, for FRB 121102, $T = 157$ days and $e \leq 0.08$, and one gets $\bar{\rho} \sim 10^{-6}\rho_{0,\odot}$. This indicates that the DSs might be supergiant stars. For red supergiants, their average density can be as low as $10^{-8}\rho_{0,\odot}$. Therefore, we expect that this model can explain a period of up to $T \sim 10^9$ days when the companion is a red supergiant. Note that, however, there is some uncertainty on the phase.
window of FRB 180916. It is pointed out in the CHIME/FRB Collaboration et al. (2020) that 50% of the CHIME bursts are detected in a 0.6 day phase window, with the event rate dropping rapidly toward the edges of the active phase, and the duty cycle would be $D = 0.04$ if 0.6 day is the width of the active phase. Moreover, Pastor-Marazuela et al. (2020) found that its activity window is narrower at higher frequencies, namely, the FWHM of Apertif bursts is 1.1 day compared to CHIME bursts’ 2.7 day. If one adopts $D = 0.04$, $e \leq 0.65$ and the density of the DS for FRB 180916 would be $\rho \sim 10^{-2.4} \rho_0$, which is in line with that of massive OB stars.

Moreover, it should be noted that the density obtained from Equation (22) is the average density, which means that it is not necessarily the true density of the DS. For example, Be stars, although they themselves have radii of only $\sim 10R_\odot$, their accretion disk radius can be as large as $\sim 100R_\odot$ (Rivinius et al. 2013). The CO can accrete material from the disk of the star, although it does not accrete material from the star itself directly (Karino 2021). In this case, for example, for a Be star with a mass of $10M_\odot$, and an accretion disk of radius $\sim 10–200R_\odot$, the effective average density would be $10^{-7}\rho_0$ to $10^{-5}\rho_0$. It can be seen that a Be companion is also consistent with the density requirements in this model.

That said, according to the results in Section 2, the model requires a relatively weak stellar wind environment, specifically, $M_w \sim 10^{-11} M_\odot$ yr$^{-1}$ for FRB 180916 and $M_w \sim 10^{-9} M_\odot$ yr$^{-1}$ for FRB 121102 may be appropriate. It is also strongly constrained by the small DM variations observed in FRB 180916 at low frequency bands (Pleunis et al. 2021). Therefore, it is unlikely that the DSs are supergiants, as they tend to have a much stronger wind, unless they happen to be in periods of cooling with a weak wind or if the wind is highly inhomogeneous/clumpy (Puls et al. 2008; Stairs et al. 2001). As mentioned above, a cold OB star companion with $M_w \sim 10^{-11} M_\odot$ yr$^{-1}$ may also be reasonable for FRB 180916. However, the DSs are more likely to be Be stars because their polar wind may be relatively weak (Kervella & Domiciano de Souza 2006; Kanaan et al. 2008), which could provide a unified picture for the cases of FRB 180916 and FRB 121102. But then again, current observations are not enough to tell us exactly what kind of stars the DSs are, and future observations are needed to provide more clues. In any case, this model’s requirement for massive stars to be DSs is consistent with the fact that FRB 121102 and FRB 180916 are associated with star-forming regions (Chatterjee et al. 2017; Marcote et al. 2020; Tendulkar et al. 2020).

Given the presence of the DS’s wind, one needs to consider its contribution to the dispersion measure (DM). According to Equation (3) and Equation (5), one can roughly estimate the DM associated with the local wind, $DM_{loc} \sim n_w \langle v_{dec} \rangle$, $v_{dec} \sim 0.03 \rho_2^{1/4} M_{wind}^{3/4} \sim 0.03 \rho_2^{1/4} \rho_t^{-1/2} \rho_s^{-3/2} \rho_{wind}^{-3/2}$ pc cm$^{-3}$ which is small enough to be negligible compared to the total DM of FRBs, even if a relatively large wind rate $n_w \sim 10^{-9}$ is adopted. Therefore, we do not expect that there is an obvious periodic modulation in the observed DM for FRB 180916B (Pastor-Marazuela et al. 2020), and of course the same is true for FRB 121102.

Interestingly, the DM of FRB 121102 seems to have a slow growth trend with a rate of $\sim 0.85$ pc cm$^{-3}$ yr$^{-1}$ (Li et al. 2021). Based on the above analysis, it is clear that the evolution of the DS’s wind is not sufficient to lead to such an outcome. The DM variation may depend on the environment in which FRB sources are located. For example, an FRB source in an expanding supernova remnant around a nearly neutral ambient medium during the deceleration phases or in a growing HII region can increase the DM (Yang & Zhang 2017). We are going to discuss this issue, in the context of this model, in detail elsewhere.

4. Summary and Discussions

In this work, we propose a model for understanding the highly active periodic FRBs. The system consists of a stellar CO and a DS, in which the DS fills with its own Roche lobe. The CO accretes material from the DS and ejects relativistic magnetic blobs. The interaction between the magnetic blobs and the stellar wind of the DS produces a pair of shocks, the reverse shock traveling through the ejecta and the forward shock traveling to the wind of the DS. We find that both of these shocks are likely to produce FRBs. The energy of the FRBs from the reverse shock is greater than that from the forward shock. It indicates that the observed energy in a single repeating FRB may exhibit a bimodal distribution. This may provide an explanation for the bimodal burst energy distribution in FRB 121102 found by FAST recently (Li et al. 2021).

Moreover, such a Roche lobe filling accretion system can in principle sufficiently power the highly active periodic FRBs over a long lifetime. The orbital motion of the binary can explain the periodicity of the FRBs such as FRBs 180916 and 121102, if the DSs are giants/supergiants or Be stars. To produce the narrow duty cycle of the activity, such as for FRB 180916, the orbit of the binary needs to be moderately eccentric, with the DS filling its Roche lobe only near the periastron. It should be noted that for our model to work, it requires (1) a sufficiently large accretion rate and (2) a weak stellar wind from the DS. If not, if the accretion is too weak, the FRB energy will be too low to be observed. If the stellar wind is too strong, GHz radiation cannot pass through freely. We realize that this is reasonable because FRBs would have been observed in a large number of binary systems in the Milky Way if it were not the case that the right conditions to produce FRBs are required in binary systems. It is these particular low wind binary systems that produce such rare sources like FRB 180916 and FRB 121102. Therefore, if this model is correct, it gives us a great opportunity to study these special binary systems.

Recent observations of FRB 180916 revealed that the burst activity is frequency dependent; namely its activity window is both narrower and earlier at higher frequencies (Pastor-Marazuela et al. 2020; Pleunis et al. 2021). The causes of this observed phenomenon may be complicated. It may be due to a combination of the absorption of FRBs by the surrounding environment and instrument selection effects, as shown in Li et al. (2021). It may also arise from the structured and beaming effects of the jet (Sridhar et al. 2021). But we will not discuss this issue in this work because it requires detailed modeling in the context of this model, and we plan to study it in detail elsewhere.

The more recent observation found subsecond periodicity in FRB 20191221A, which may indicate that the central engine of this FRB is an NS, and the period corresponds to the rotation period of the NS (the CHIME/FRB Collaboration et al. 2021). However, the duration of this FRB is actually $\sim 3$ s, which is different from any other known FRB, and it has not yet been discovered whether it repeats, so it may be an entirely new class of FRBs (the CHIME/FRB Collaboration et al. 2021). Nevertheless, a similar subsecond periodicity is known to exist in some X-ray binaries, which may result from the rotation of the accreting NSs (Kaaret et al. 2017; Patruno & Watts 2021).
A similar periodicity is also expected in GRBs, although no unambiguous periodicity has been found in GRB pulses (Tarnopolski & Marchenko 2021 and references therein). In this model, the rotation of the accreting NS (or the fluctuations in the accretion disk) may also modulate the accretion process and thus the generation of the jets, and whether it ultimately results in the production of a periodicity in FRB pulses deserves further study.

It is expected that, in the context of this model, the energy injection by the ejection of long-term blobs from the system into the surrounding environment may excite a radio nebula that can explain the persistent radio source associated with FRB 121102, inspired by the fact that the Galactic X-ray binary SS433 powers a similar radio nebula (Fabrika 2004). If this model is correct, the luminosity of the persistent radio source may be estimated as $L_R \propto E_{\text{FRB}}^3$, where $E_{\text{FRB}} \equiv E_{\text{FRB}} \Omega$, $E_{\text{FRB}}$ is the typical energy of the FRB, and $\Omega$ is the repetition rate. Here, we adopt $\gamma = 1$ for a rough estimation, although the exact value is expected to be slightly greater than 1 in the context of synchrotron radiation (Dai et al. 2017). Taking the typical isotropic energy $\sim 10^{37}$ erg with a repetition rate $\Omega \sim 10^{-1}$ h$^{-1}$ for FRB 180916 (Pastor-Marazuela et al. 2020), and a similar isotropic energy but with a much higher repetition rate $\Omega \sim 10^3$ h$^{-1}$ (Li et al. 2021) for FRB 121102, we have $E_{\text{FRB}} \sim 10^{36}$ erg for both FRB 180916 and FRB 121102. Therefore, we predict that the persistent radio emission associated with FRB 180916 would be $\approx 10^{35}$ erg s$^{-1}$, which is in line with the observational limit (Marcote et al. 2020).

In addition, it is expected that the accreting CO will also have persistent X-ray emission. Assuming solar abundances, the X-ray luminosity of the accretor is (Shakura & Sunyaev 1973; Poutanen et al. 2007)

$$ L_X \approx 1.3 \times 10^{38} \text{erg s}^{-1} \left[ \frac{\dot{m} M_{\odot}}{M_{\text{Edd}}} \right], \quad \frac{\dot{m}}{1} \leq 1 $$

$$ \left(1 + \ln \dot{m} \right) \left[ \frac{\dot{m} M_{\odot}}{M_{\text{Edd}}} \right], \quad 1 \leq \dot{m} \leq 100 $$

(26)

where $\dot{m} = \dot{M}/M_{\text{Edd}}$, with the super-Eddington accretion rate $M_{\text{Edd}} \approx 2.5 \times 10^{31} (M_{\odot}/M) M_{\odot}$ yr$^{-1}$. Based on the discussions in Section 2, we have $L_X \approx 10^{38}$ erg s$^{-1}$ for an accretion rate $\dot{M} = 10^{-5} M_{\odot}$ yr$^{-1}$, which is below the detection limit for both FRB 180916 and FRB 121102 (Tavani et al. 2020; Scholz et al. 2020). However, we expect to detect the X-ray emission from the accretor for sources at close distances, i.e., within a few tens of Mpc. In addition, we have confirmed that the prompt gamma-ray emission radiated by the thermalized electrons behind the (reverse and forward) shocks associated with FRB 121102 and FRB 180916 are also expected to be too dim to be detected by the current gamma-ray detectors.

Finally, we anticipate that if those active periodic FRBs are observed at close distances in the future, multiband observations will verify or falsify our model. Also, we encourage the search for FRBs in ultraluminous X-ray sources.

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