On the nature of Fast Blue Optical Transients

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20 June 2022

ABSTRACT

Short rise times of Fast Blue Optical Transients (FBOTs) require very light ejected envelopes, $M_{ej} \leq 10^{-1} M_\odot$, much smaller than of a typical supernova. Short peak times also mean that FBOTs should be hydrodynamically, not radioactively powered. The detection by Chandra of X-ray emission in AT2020mrf of $L_X \sim 10^{42}$ erg s$^{-1}$ after 328 days implies total, overall dominant, X-ray energetics at the Gamma Ray Bursts (GRBs) level of $\sim 6 \times 10^{49}$ erg. FBOTs show no evidence of relativistic motion, hence no beaming: the observed X-ray luminosity is similar to the true isotropic luminosity.

We further develop a model of Lyutikov & Toonen (2019), whereby FBOTs are the results of a late accretion induced collapse (AIC) of the product of super-Chandrasekhar double white dwarf (WD) merger between ONeMg WD and another WD. Small ejecta mass, and the rarity of FBOTs, result from the competition between mass loss from the merger product to the wind, and ashes added to the core, on time scale of $\sim 10^3 - 10^4$ years. FBOTs occur only when the envelope mass before AIC is $\leq 10^{-1} M_\odot$. FBOTs proper come from central engine-powered radiation-dominated forward shock as it propagates through ejecta. FBOTs’ duration is determined by the diffusion time of photons produced by the NS-driven forward shock within the expanding ejecta. All the photons produced by the central source deep inside the ejecta escape almost simultaneously, producing a short bright event, violating the “Arnett’s law”. The high energy emission is generated at the highly relativistic and highly magnetized termination shock, qualitatively similar to Pulsar Wind Nebulae. The X-ray bump observed in AT2020mrf by SRG/eROSITA, predicted by Lyutikov & Toonen (2019), is coming from the break-out of the engine-powered shock from the ejecta into the preceding wind. The model requires total energetics of just few $\times 10^{50}$ ergs, slightly above the observed X-rays. We predict that the system is hydrogen poor.

Keywords: white dwarfs; supernovae: general; neutron stars: neutron stars; white dwarfs: supernovae: general; neutron stars: neutron stars; white dwarfs: supernovae: general

1 INTRODUCTION

Fast-rising blue optical transients (FBOTs, Drout et al. 2014) is a class of bright, short supernova explosions. AT2018cow (Prentice et al. 2018; Ho et al. 2019; Perley et al. 2019; Margutti et al. 2019) and AT2020xnd Ho et al. (2021a) had similar properties. Recent observations of AT2020mrf (Yao et al. 2021) further constraint the properties of the FBOTs’ progenitors. Similarity between AT2018cow, AT2018lqh, AT2020xnd and AT2020mrf imply similar physical system, with some variations of the parameters (see Ho et al. 2021b, for a classification of rapid optical transients - most are classified as core-collapse events).

The new constraint provided by the AT2020mrf is the relatively bright X-ray emission detected nearly a year after the original explosion. Let us give estimates of the different observed channels in FBOTs, taking AT2020mrf as an example (Yao et al. 2021) ($E_{v,r,mm,X}$ below is the total energetics in optical, radio, millimeter and X-rays)

- Optical: $M_v = -20$ for 3.7days at $3 \times 10^{43}$ erg/sec, $E_v = 10^{49}$ erg.
- Radio emission: $\nu F_\nu = 1.2 \times 10^{49}$ erg/s at 261 days: $E_r = 2.7 \times 10^{46}$ erg.
- Millimeter (Table 3 of Yao et al. 2021): 50μJy at 16 GHz at 417.5 days: $E_{mm} \sim 1.5 \times 10^{47}$ erg.
- X-rays: (i) $2 \times 10^{43}$ erg s$^{-1}$ at 20 days, $E_X^{\text{early}} = 3.5 \times 10^{49}$ erg; (ii) $10^{42}$ erg s$^{-1}$ at 328 days, $E_X^{\text{late}} = 2.8 \times 10^{49}$ erg. (The early X-ray luminosity of AT2018cow ($10^{43}$ erg s$^{-1}$ at 20days, $E_X^{\text{early}} = 1.7 \times 10^{49}$ erg. Prentice et al. 2018).

Thus, the most energetically constraining observation of FBOTs is the high energy X-ray emission. Especially surprising is the detection by Chandra of GRB-like emission of $L_X \sim 10^{42}$ erg s$^{-1}$ after 328 days (AT2020xnd had a similar flux at $\sim 50$ days Perley et al. 2021). Optical emission of FBOT proper comes close. (The nature of the hot and luminous source detected by Sun et al. 2022, in the direction of AT2018cow is still uncertain.) The implied total X-ray energetics for AT2020mrf is $\sim 6 \times 10^{49}$ erg s$^{-1}$. Thus, energies of AT2020mrf matches those of...
GRBs; but unlike GRBs they do not show relativistic velocities, hence their luminosities/energetics are of the order of the true luminosities, while in GRBs true luminosities are smaller by the beaming factor, $\sim 10^{-2}$.

Explaining the power and the photon energies a year after an explosion is the most challenging. In what follows we demonstrate that the model of Lyutikov & Toonen (2019) both predicted the X-ray bump observed by SRG /eROSITA, can account for the total energetics, and generally explains all the observed phenomena of FBOTs.

We also mention models of FBOTs by Leung et al. (2020); Gottlieb et al. (2022); Soker (2022). All these models rely on massive hydrogen-rich stars. To keep the energy budgets under control the models require highly jetted, GRB-like outflows. Explaining late X-ray emission, after nearly a year, is most challenging within these models. The main observations distinction is that the present model in contrast advocates light hydrogen-poor ejecta.

2 HYDRODYNAMIC AND RADIOACTIVE CONTRIBUTIONS TO SN LIGHT-CURVES

Supernova light curves is a complicated combination of hydrodynamic/internal heat dissipation (Grassberg et al. 1971) and radioactive decay (Colgate & McKee 1969), see also Filippenko (1997); Woosley et al. (2002). The SN-Ia light curves are powered mostly by $^{56}$Ni beta-decay of $\sim 0.5 M_\odot$ (Arnett 1982; Pinto & Eastman 2000). SN-Ib/c and SN-II are powered by a combination of shock heating, recombination of hydrogen (Grassberg & Nadyozhin 1976), and later by the Co56 $\rightarrow$ Fe56 decay (see reviews by Nadyozhin & Imshennik 2005; Smartt 2009; Livio & Mazzali 2018). “FBOTs proper” ¹, with the short rise time, cannot be powered by the radioactive decay (Lyutikov & Toonen 2019; Pasham et al. 2021, yet the long term properties, on the scale of $\sim$ months, can be/are affected). FBOTs proper must be hydrodynamically-powered.

Let us take an extreme position, and neglect the energy contribution from the radioactive decay. It is overall mildly significant, but comes at a later time, §5.

Hydrodynamically powered light curves, with short rise time, require small mass of the ejecta, §3.1, otherwise most of the internal heat or internal shock power is lost to the adiabatic expansion. The required ejecta mass is $M_{ej} \lesssim 0.1 M_\odot$. This is an order of magnitude smaller than a typical SN-Ia/b/c or SN-II ejecta. Alternative possibility - massive and very fast ejecta moving nearly with the speed of light - requires enormous energy budget, on par with GRBs.

Small ejecta mass is the key ingredient of a model by Lyutikov & Toonen (2019), whereby FBOTs result from an electron-capture collapse to a neutron star of a merger product a massive ONeMg white dwarf (WD) with another WD, Fig. 1. Two distinct evolutionary channels lead to the disruption of the less massive WD during the merger and the formation of a shell burning non-degenerate star incorporating the ONeMg core. After the transients settle down, the result is a special type shell-burning star with a size few times $10^7$ cm, fast rotating (at the surface), with luminosity $L \sim 10^4 L_\odot$, producing nearly hydrogen-clear winds. The star lives for $\sim 10^4$ years, while the envelope mass is both lost to the wind and added to the core as nuclear ashes. If the mass of the core exceeds the Chandrasekhar mass, the electron-capture collapse follows after $\sim 10^4 - 10^5$ years.

The collapse produces various observed phenomena that depend on the particular properties/parameters of the merging system. In particular, the observed properties of the collapse depend on (i) duration of shell burning affecting the amount of envelope mass left at the moment of collapse; (ii) duration of shell burning and the corresponding amount of angular momentum transferred to the core; (iii) the viewing angle with respect to the axis. Eventually, the amounts of mass left in the shell and the core’s angular momentum at the moment of collapse depend on the masses of the merging WDs and the orbital separation of the Main Sequence stars. This scenario explains a small envelope mass of FBOTs: as little as $\sim 10^{-2} M_\odot$ of the material is ejected with the total energy $\sim$ few $10^{50}$ ergs. This ejecta becomes optically thin on a time scale of days.

During the collapse, the neutron star is spun up and magnetic field is amplified (Obergaulinger et al. 2009; Barkov & Komissarov 2011; Mösta et al. 2015). The ensuing fast magnetically-dominated relativistic wind from the newly formed neutron star shocks against the ejecta, and later against the pre-collapse wind. The radiation-dominated forward shock produces the long-lasting optical afterglow, while the termination shock of the relativistic wind produces the high energy emission in a manner similar to Pulsar Wind Nebulae.

3 FBOT PROPER

In the case of AT2020mrf most constraints come from the late X-ray detection. In this section we address the nature of FBOT proper (bright optical transients lasting a few days), but we use numerical parameters demanded by the X-ray, see §4.

¹ We call “FBOTs proper” the short, few days, bright optical transients

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Figure 1. Outline of the model of Lyutikov & Toonen (2019). A binary system, via two separate evolutionary channels, leads to the merger of a heavy ONeMg WD with another WD. The merger produce consists of a non-degenerate envelope around the ONeMg core. Shell material is both lost to strong winds with \( \dot{M} \sim 10^{-5} - 10^{-3} M_\odot/\text{yr} \), and ashes added to the core (IRAS 00500+6713 Oskinova et al. 2020, is currently at this stage, see §6). The core experiences AIC (Nomoto & Kondo 1991), producing fast rotating neutron star and ejecting a light remaining shell of \( M_{\text{ej}} \sim 10^{-2} \) few \( 0.1 M_\odot \). The secondary wind from the neutron star propagates first through the ejecta (the radiation-dominated forward shock produces FBOTs at this stage) and breaks out into the wind (this produces SRG /eROSITA X-ray bump predicted by Lyutikov & Toonen 2019). At later stages the high energy emission is generated in a PWN-manner, at the termination shock of highly relativistic and highly magnetized pulsar wind.

3.1 Small ejecta mass

An AIC of a WD produces a central engine, a fast rotating neutron star, while ejecting some mass \( M_{\text{ej}} \). For homologous expansion of the ejecta with \( v \propto r \) (see Sakurai 1956; Matzner & McKee 1999; Ro & Matzner 2013, for a more detailed modeling), the energy in the ejected part is

\[
E_{\text{ej}} = \frac{3}{10} M_{\text{ej}} V_{\text{ej}}^2,
\]

while the density evolves according to

\[
\rho_{\text{ej}} = \frac{3}{4\pi} \frac{M_{\text{ej}}}{(V_{\text{ej}1})^3}
\]

where \( V_{\text{ej}} \) is the maximum velocity of the ejecta.

Using scattering cross-section \( \kappa \approx 0.1 \text{ cm}^2 \text{ g}^{-1} \) (Arnett 1982), total optical depth through ejecta

\[
\tau_{\text{tot}} = \frac{3}{4\pi} \frac{M_{\text{ej}} \kappa}{V_{\text{ej}}^2} = 640 m_{\text{ej},-1} t_d^{-2} V_{\text{ej},4}^{-2}
\]

where \( m_{\text{ej},-1} = M_{\text{ej}}/(0.1 M_\odot) \), \( V_{\text{ej},4} = V_{\text{ej}}/(10^4 \text{ km s}^{-1}) \) and time \( t_d \) is measured in days.

The diffusion time is then

\[
t_{\text{FBOT}} = \left( \frac{3}{4\pi} \frac{M_{\text{ej}} \kappa}{c V_{\text{ej}}} \right)^{1/2} = 4.6 m_{\text{ej},-1}^{1/2} V_{\text{ej},4}^{-1/2} \text{ days}
\]

We identify the diffusion time \( t_{\text{FBOT}} \) with the observed peak of the light curve.

Several arguments can be used to derive (4). In §3.3 we derive photon’s Green function in the expanding ejecta, showing the scaling for the diffusion radius \( r_d \propto t^2 \), Eq. (20). Second, the time for a shock traveling with velocity \( V_s \) through ejecta of thickness \( V_{\text{ej}} t \) should be set equal to the photon diffusion time \( (V_s t)^2/\langle \rho \rangle \) where \( \langle \rho \rangle = 1/(c \rho) \) is mean free path (Ohyama 1963; Castor 1972); equating the two gives (4). Equivalently, for radiation-dominated shocks, the photons diffusing ahead of
Figure 2. Velocity structure: Left panel before NS-drive shock breaks out into the wind (top right panel in Fig. 1); Right panel after break-out (bottom right panel in Fig. 1). At large radii this is a pre-explosion wind with \( v_w \sim 10^4 \) km s\(^{-1}\). Ejecta, with a linear velocity profile, propagates into the wind, launching a forward shock (FS-1) and reverse shock (RS-1). At smaller radii the wind from the neutron star generate forward shock (FS-2) in the ejecta and reverse shock (RS-2) in the neutron star wind. X-rays are produced at RS-2, FBOT proper at FS-2. Photons (red dashed lines) diffuse ahead of the radiation-dominated shock. At later times the neutron star wind may broke out into the pre-explosion wind. Long term X-rays are generated at the RS, mm-radio emission - at the FS.

Figure 3. Diffusion of NS-shock produced photons in optically thick ejecta. The diffusion time corresponds approximately to the rise time of FBOTs, Eqns. (4) and (5)

The shock escape when optical depth to the emitting surface is \( \tau \sim c/V_s \). Estimating shock velocity \( V_s \sim V_{ej} \) this occurs at (4) (see also Arnett 1982).

If scattering is purely Thomson, then one needs to change \( \kappa \to \sigma_T/m_p = 0.4 \) cm\(^2\) g\(^{-1}\),

\[
\tau_{tot} = \frac{3}{4\pi} \frac{M_{ej}\sigma_T}{m_p V_{ej}^2} = 2.5 \times 10^3 \times m_{ej,-1} \tau_{e}^{-2} V_{ej,4}^{-2}
\]

\[
t_{FBOT} = \left( \frac{3}{4\pi} \frac{M_{ej}\sigma_T}{cm_p V_{ej}} \right)^{1/2} = 9.2 m_{ej,-1} V_{ej,4}^{-1/2} \text{ days},
\]

a mild correction. The gray scattering cross-section \( \kappa \sim Y_e x_e \) cm\(^2\) s\(^{-1}\), \( Y_e \) is the electron fraction, \( x_e \) is the ionization degree.
For AT2020mrf the peak is at 3.8 days, hence

\[ M_{ej} = 6.7 \times 10^{-2} M_\odot V_{ej,A}^{1/2} \]

\[ E_{ej} = \frac{3}{10} M_{ej} V_{ej}^2 = 4 \times 10^{49} V_{ej,A}^2 \text{erg} \]

\[ R_{ej} = V_{ej} t_{FBOT} = 3 \times 10^{14} \text{cm} \tag{6} \]

(At times much shorter than \(7\) the photospheric radius is just a bit smaller, by \(\sim 1 - (2/3)V_{ej}/c\).)

Estimate \(6\) implies that ejecta mass must be small, \(\leq 10^{-1} M_\odot\), much smaller of a typical ejecta of any conventional supernova. If \(\sim 1 M_\odot\) is ejected, short diffusion time \(4\) would require \(V_{ej} \sim c\), with the total energetics at \(\sim 10^{54} \text{ergs}\). Estimate \(6\) on the mass is the upper limit: (i) Wheeler et al. (2015) estimated \(5\) as the rise time (not peak time), which is shorter; (ii) we omitted a factor 1/3 in the diffusion coefficient.

The ejecta itself becomes fully transparent at somewhat longer times

\[ t_r \approx t_{FBOT} \sqrt{\frac{c}{V_{ej}}} = 25 m_{ej,-1}^{1/2} V_{ej,-1}^{1/2} \text{days} \rightarrow 21 V_{ej,-1}^{-1} \text{days} \tag{7} \]

(the last relation uses estimate of the ejecta mass \(6\)). Thus the FBOT proper is contributed both by the energy input from the NS wind and somewhat later by cooling of the ejecta.

The effective emission radius starts to decrease at approximately half of the time \(7\). Qualitatively, the radius \(R_{ph}\) where optical depth to the surface equals unity evolves according to

\[ R_{ph} = V_{ej} t \left(1 - \frac{4\pi (V_{ej} t)^2}{3 M_{ej} c^2}\right) \tag{8} \]

It decreases after half of the time \(7\), until the ejecta becomes fully transparent at time \(7\).

This is consistent with observation of (Perley et al. 2019, their Fig. 8) which shows long-term decreasing effective radius (ejecta velocity somewhat smaller that \(10^4 \text{km s}^{-1}\) is required to extend this to \(\sim 1 \text{ month}\)). We also note that similar to early stages of terrestrial nuclear explosions, the radiative cooling of the ejecta may be faster via the formation of the cooling wave (Zeldovich & Raizer 2003; Grassberg & Nadyozhin 1976).

Small mass of the ejecta is a challenge to the conventional models: typical SN ejected mass is \(\geq 1 M_\odot\). To produce mildly relativistic motion of large mass, even for jetted outflows, requires GRB-type energies, well in excess of \(10^{51}\) of kinetic energy (and correspondingly large total rates.) The model of Lyutikov & Toonen (2019) offers a natural explanation.

### 3.2 Blue color: radiation-dominated shock within the ejecta

Consider a two-stage explosion: the initial neutrino-driven ejection of a shell with mass \(M_{ej}\) and maximal velocity \(V_{ej}\), followed by the central source (neutron star-) driven (second) wind (the “first wind” is the pre-explosion wind from the progenitor star.) Thus, the initial “heavy lifting” is done by SN shock. The energy of the central engine is then mostly spent on producing the non-thermal emission, as opposed to generating heavy slow outflows (see discussion by Lyutikov 2011a, in applications to GRBs). The required delay time is only few seconds - time for the SN shock to cross the collapsing envelope. In the magnetar model of relativistic explosion (Usov 1992; Komissarov & Barkov 2007; Metzger et al. 2011) it is expected that in few seconds the neutron star cools sufficiently so that it’s wind becomes clean, pulsar-like.

After the SN shock propagated through the ejecta-to-be, the NS-driven wind launches another shock in the expanding envelope. We consider its dynamics next.

Qualitatively two regimes for the shock propagation through ejecta may be realized: energy conserving Sedov-type (Sedov 1959), and momentum driven (thin shell) Kompaneets-type (Kompaneets 1960). Sedov-type flow would occur if the termination shock in the wind is close to the source (e.g., as in the case of Crab PWNe). The wind-blown bubble is then in approximate causal contact. The thin shell case is realized if the termination shock is close to the contact discontinuity - it is much more energetically demanding, by a factor \(c/V_s\), and is less likely to be realized. Below we use the Sedov scaling (see also Chevalier 1982)

At sufficiently short times, less than a month, we can assume a constant spin-down luminosity \(L_{sd}\) (see Eq. (35) for justification). For a contact discontinuity at radius \(R\) the swept-up energy and mass are

\[ E_s = \frac{3}{10} M_{ej} R^5 \]

\[ M_s = \frac{M_{ej} R^3}{V_{ej}^3} \tag{9} \]

(since \(R \propto t^{6/7}\), Eq. (11), both \(E_s\) and \(M_s\) increase with time).

The energy balance and the equation of motion then read

\[ L_{sd} t + E_s = \frac{M_s (\dot{h} R)^2}{2} \]
The diffusion equation reduces to the familiar form with solution

\[ \partial_t R = \sqrt{\frac{2 L_{sd} V_{ej}^3 r^4}{M_{ej} R^2} + \frac{3 R^2}{5 t^2}} \]  

(10)

Eq. (10) has a solution

\[ R = \left( \frac{50 L_{sd} V_{ej}^3 r^4}{21 M_{ej}} \right)^{1/5} \]

\[ \frac{R}{V_{ej} t} = \left( \frac{t}{t_{br}} \right)^{1/5} \]  

(11)

Velocity of the shock with respect to the ejecta

\[ V_s = \partial_t R - \frac{R}{t} = \left( \frac{t}{t_{br}} \right)^{1/5} \frac{V_{ej}}{5} \]  

(12)

where \( t_{br} \) is given by (25).

The radiation-dominated shock jump condition (it may be verified that early on the shock is radiation-dominated)

\[ \frac{4}{c} \sigma_{SB} T^4 \sim \rho_{ej} V_s^2 \]  

(13)

gives

\[ T = 1.5 \times 10^3 L^{1/10}_X \frac{M_{ej}^{3/20}}{V_{ej}^{13/20} L^{1/10}_X} = 4 \times 10^4 L^{1/10}_X \frac{e_{X,-1} V_{ej}^{-9/20}}{c V_{ej}^{1/10} \rho_{ej}} K \]  

(14)

at time \( t = 3.8 \) days. This is an upper limit on temperature as it assumes that all the enthalpy is provided by the radiation; matter contribution would reduce the estimate of temperature. A blackbody fit to the spectrum suggests a temperature of \( T = 2 \times 10^4 \) K and a radius of \( R = 7.9 \times 10^{14} \) cm (Yao et al. 2021). Our estimates (6) and (14) are consistent.

Predicted evolution of temperature \( T(t) \propto t^{-13/20} \) (14), corresponding to the constant density of the ejecta, can be tested against observations. Our fit to the evolution of temperature in FBOT AT2018cow (Table 4 of Perley et al. 2019) gives \( T \propto t^{-0.23} \) (limited to 27 days; after this time there is a clear break in the evolution of temperature; in our model the break is due to the transition of the NS-powered shock into the preceding wind Lyutikov & Toonen 2019). Freely expanding ejecta models predict much steeper decay \( T \propto t^{-1} \) (in the adiabatic limit). Thus the present model fairs better. Also, variations of the density of the ejecta give more freedom.

Additional complication comes from the fact that in the highly radiation-dominated regime the shock, defined as a hydrodynamic discontinuity, may disappear, so that fluid properties evolve smoothly (Zeldovich & Raizer 2003; Weaver 1976). Alternatively, an isothermal jump may form (Landau & Lifshitz 1959). Qualitatively, the above estimates remain valid.

The energetics of FBOTs come both from the internal energy of the hot ejecta and the luminosity of the central source. Both sources of energy also experience adiabatic losses. For the energy contributed by the central source all the photons emitted by the central source at times \( t \leq t_{FBOT} \) come out in a narrow window \( \sim 1/3 \) near \( t_{FBOT} \), see Fig. 6). In addition, a photon that is injected at day 1 and finally diffuses out on day 3 will have dropped in energy by a factor of \( \sim 1/3 \). We can estimate the peak luminosity then as \( L_{FBOT} \sim L_{sd} \sim 10^{44} \) erg s\(^{-1}\).

3.3 Short duration of FBOTs: photon diffusion from the central source through expanding ejecta

3.3.1 Photon’s Green’s function in expanding ejecta

Consider early times when the NS-powered shock is at small radii within the ejecta. Approximate the shock as a source of photons located at \( r = 0 \) (according to (11) the shock remains deep with the the ejecta for a long time). In an expanding medium the diffusion coefficient \( \eta_d \) varies as

\[ \eta_d = \frac{\kappa_0 (t_+ + t_0)^3}{3 M_{ej} c} \]

(15)

where \( t_0 \) is the delay time before the beginning of the expansion and the injection of photons and \( t_+ \) is time since injection. Neglecting advective transport, making a change in time variable,

\[ \tilde{t} = \frac{1}{4} t \left( t_+^3 + 4 t_0 t_+^2 + 6 t_0^2 t_+ + 4 t_0^3 \right) \]

\[ \tilde{t} \approx t^4/4 \text{ for } t \gg t_0, \]  

(16)

the diffusion equation reduces to the familiar form with solution

\[ G(r, t) = \frac{1}{\sqrt{8 \pi \eta_0 t}} e^{-r^2/(4 \eta_0 t)} \]  

(17)
Figure 4. Photons’ Green function for diffusion in expanding ejecta with time-dependent diffusion coefficient (18) (solid lines) if compared with Green’s function for constant diffusion coefficient (dashed lines). Two sets of plotted curves correspond to $t = 0.1t_0$ and $t = 10t_0$. Notice that at $t = 10t_0$ the Green’s function extends to $\sim t^2$.

see Fig. 4. $G(r, t)$ is normalized so that $4\pi\int_0^\infty r^2 G(r, t) dr = 1$. Eq. (18) gives Green’s function for $\delta(t)$ injection at $r = 0$ that occurred at $t_0$ after explosion, as measure at time $t + t_0$ after the injection.

Shifting time $t + \rightarrow t_0$ ($t$ is time measured from the initial explosion),

$$G(r, t) = \frac{1}{8(\pi \kappa_0)^{3/2}} e^{-r^2/(\kappa_0(t^2 - t_0^2))}$$  

(18)

This is Green’s function for injection at time $t_0$ after the explosion, as measured at time $t$ after the explosion; $t > t_0$.

For zero delay, $t_0 = 0$,

$$G(r, t) = \frac{1}{8(\pi \kappa_0)^{3/2}} e^{-r^2/(\kappa_0 t^2)}$$  

(19)

Thus, the diffusion radius scales as

$$r_d \approx \sqrt{\kappa_0 t^2}$$  

(20)

The diffusion radius (20) then becomes $\sim V_{ej} t$ at time (4)-(5).

3.3.2 Short duration: “diffusive caustic”

During the FBOT proper the wind shock remains deep inside the ejecta. We can then use the Green’s function (18) with a given wind luminosity to find distribution of photons inside the ejecta. (One also has to take account of the fact that photons escape from the edge.) Let us give semi-qualitative estimates of the expected light curve.

Using (18) the typical photon trajectory is

$$r_d = \sqrt{t^4 - t_0^4 \sqrt{\kappa_0}}$$  

(21)

(Again, here $t_0$ is injection time, $t$ is time since the explosions). Equating $r_d$ with the location of the surface of the ejecta, $r_d = V_{ej} t$, we find the escape time

$$t_{esc} = \sqrt{1 + \sqrt{4t_0^4 + 1}}$$

$$t_{esc} = \frac{t_{esc}}{t_{FBOT}}$$

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Figure 5. Photon trajectories (21) as a function of the injection time $\tilde{t}_0$. Dashed line with inclination of 1 is the surface of the ejecta. In this approximation nearly all the photons produced before $t_{FBOT}$ escape nearly simultaneously.

$$\tilde{t}_0 = \frac{t_0}{t_{FBOT}}$$

Thus, all the photons produced by the central source deep inside the ejecta, $\tilde{t}_0 \leq 1$ escape almost simultaneously, Fig. 5. This reminds of a caustic, hence the name.

To get an analytical estimate of the flux, we can integrate (18) times the luminosity over the injection time $t_0$. For constant photon production rate, using steepest decent method we find the photon density at $r = V_e t$

$$n_{ph}(r = V_e t) \propto \frac{1}{t_{FBOT}} \frac{1}{t^2} \frac{1}{(1 - (2/3)(t_{FBOT}/t)^2)^{3/4}}$$

Divergence at $t = \sqrt{2/3} t_{FBOT}$ is an artifact of the analytical approximation. Yet it demonstrates qualitatively that short, bright events can be produced.

Relation (23) also explains why short $t_{FBOT}$ are needed to produce bright FBOTs: the flux at $t = t_{FBOT}$ is (neglecting the divergent component)

$$F \propto n_{ph} \big|_{r = V_e t} \times r^2 \propto t_{FBOT}^{-3}$$

3.3.3 Monte Carlo simulations of photon escape

We performed simple 1D Monte Carlo simulations of photon propagation within the ejecta. Photons are injected at $r = 0$ with variable rate (impulsive $\delta(t)$, constant and declining rates). The diffusion coefficient scales as $\propto t^3$. Coordinate of each photon $r$ is traced until it reaches $V_e t$. It is then assumed that the photon scapes. In Fig. 6 we plot the rate of photon escape. At early times all the photons are trapped. Near $t_{FBOT}$ all the trapped photons escape with a short time $\sim t_{FBOT}/3$. Thus, the peak flux is $\sim$ three times above the average source luminosity. These are clear counter example to the “Arnett’s law” (Arnett 1982, that the luminosity at the peak is equal to the instantaneous luminosity at that time). For example, for impulsive injection the instantaneous luminosity at the peak escaping luminosity is zero. At later times the flux is approximately the injected flux (zero for impulsive, constant and/or declining).
Note that at long times the optical light curve reflects the (decreasing) rate of the production of the optical photons in the forward shock, not the power of the central engine, which may remain constant.

Our results for light curves are different from Kasen & Bildsten (2010). We treat the input from a magnetar as perturbation, not as a dominant energy source in the ejecta.

3.4 SRG /eROSITA X-ray bump: break-out of the NS-driven shock from ejecta into the wind

The NS-driven shock breaks from the ejecta at

\[ t_{br} = \frac{21}{50} \frac{M_{ej} V_{ej}^2}{L_{sd}} \approx \frac{E_{ej}}{L_{sd}} \]

approximately when the NS-injected energy becomes comparable to the ejecta energy. (At times (26) the spin-down power can be assumed constant, see Eq. (35)).

Using (32) this occurs at

\[ t_{br} = \epsilon_e \frac{E_{ej}}{L_X} = 46 \epsilon_e^{-1} V_{ej,A}^2 L_X^{-1} \text{ days} \]

(the ejecta mass is incorporated into ejecta energy, see Eq. (25)).

We associate this time with the SRG /eROSITA X-ray bump. (Production of UV and soft X-rays during shock breakout has been discussed by Imshennik & Nadezhin 1988; Ensmann & Burrows 1992; Blinickov et al. 2000; Calzavara & Matzner 2004). The X-ray feature during the break-out was predicted by Lyutikov & Toonen (2019).
Figure 7. Cartoon of late X-ray and IR emission. The central neutron star produces a relativistic wind with Lorentz factor $\Gamma_w$ and power $L_{sd}$. The wind shocks at the reverse shock, where particles are accelerated to the Lorentz factor $\gamma \sim \Gamma_w$, and produce X-ray synchrotron emission in the wind’s magnetic field. The IR emission originates from the synchrotron emission in the turbulent, amplified magnetic field behind the forward shock.

4 LONG TERM X-RAY EMISSION FROM RELATIVISTIC TERMINATION SHOCK

4.1 Dynamics of wind-wind interaction

The long term X-ray emission provides the most energetic constraints on the model. At this stage (last panel in Fig. 1), the central engine - newly formed neutron star - produces relativistic, highly magnetized wind that produce two shocks, the forward shock (FS) in the preceding wind, and the reverse shock (RS) in the relativistic wind. Thus, it is relativistic wind-wind interaction (see Lyutikov 2017; Lyutikov & Camilo Jaramillo 2017; ?; Barkov et al. 2021, for dynamics of such double explosions). Next we discuss the model of Lyutikov & Toonen (2019) in the light of new observations of AT2020mrf by Yao et al. (2021), see Fig. 7 for a qualitative description.

At late times, months, the forward shock propagates through a powerful pre-explosion wind with mass loss rate

$$\dot{M} = 4\pi R^2 \rho_w v_w$$

At times of few months - years the evolution of the spin-down luminosity $L_{sd}$ of the central neutron star should be taken into account,

$$L_{sd} = \frac{L_0}{(1 + t/t_\Omega)^2}$$

where $L_0$ is the initial spindown power, $t_\Omega$ is the spindown time.

Neglecting the swept-up momentum and energy of the preceding wind, in the Kompaneets approximation (Kompaneets 1960; Bisnovatyi-Kogan & Silich 1995) the non-relativistic source-powered forward shock propagates according to

$$\frac{L_{sd}}{4\pi R_s^2 c} = \rho_w (\partial_t R_s - v_w)^2$$

where $R_s$ is the location of the shock (more precisely, this is the location of the structure of the reverse shock-contact discontinuity - forward shock in the thin shell approximation; this overestimates the location of the reverse shock). Formally this is a condition at the contact discontinuity, which is located close to the FS. (In passing we note that conclusion of Bietenholz et al. 2020, that absence of overall relativistic expansion excludes long-lived relativistic outflow is incorrect: overall expansion traces the post reverse shock flow, not the fast upstream wind, Fig. 2. Like in PWNe, the wind is highly relativistic, while the overall expansion is non-relativistic.)
In the thin shell approximation the reverse shock-contact discontinuity-forward shock system move according to
\[ R_s = v_w t + \left( \frac{L_0 v_w}{c M} \right)^{1/2} t_\Omega \ln(1 + t/t_\Omega) \rightarrow \left( 1 + \frac{L_0}{c v_w M} \right)^{1/2} \] \[ v_w t \approx 2 \times 10^{37} \text{ cm} \left( \frac{0.43}{L} \right)^{1/2} \] \[ L_{sd} \geq c v_w M = 2 \times 10^{41} \dot{M}_{-4} v_{w,4} \text{ erg s}^{-1} \] (31)
where \( t_{0,43} = L_0/(10^{43} \text{ erg s}^{-1}) \) (this is fixed by the late time X-ray power, see below) and numerical estimates are at \( t = 328 \) days. Thus the shock is mildly relativistic (we also assumed \( t \leq t_\Omega \), see below). This is consistent with the limit derived by Bietenholz et al. (2020).

Strong shock condition requires
\[ L_{sd} \approx c v_w M = 2 \times 10^{41} \dot{M}_{-4} v_{w,4} \text{ erg s}^{-1} \]
This is satisfied by our fiducial parameters.

### 4.2 Properties of the central neutron star

In a PWN paradigm (e.g., Kennel & Coroniti 1984; Luo et al. 2020), Fig. 7, the X-ray emission occurs in fast cooling regime: hence all the energy put into particles is radiated. The observed X-ray luminosity \( L_X \) is then a fraction \( \epsilon_e \) of the spin-down luminosity \( L_{sd} \)
\[ L_X = \epsilon_e L_{sd} \]
\[ L_{sd} \approx B_{NS}^2 R_{NS}^2 c \left( \frac{\Omega_{NS} R_{NS}}{c} \right)^4 \] (32)
Where \( B_{NS} \) is the surface magnetic field, \( R_{NS} \) is neutron star radius, and \( \Omega_{NS} \) is the spin of the central neutron star.

Scaling the surface magnetic field with the quantum field,
\[ B_{NS} = b_q B_Q \]
\[ B_Q = \frac{c^3 m_e^2}{\hbar} = 4 \times 10^{13} \text{ Gauss}, \] (33)
and given the observed X-ray luminosity \( L_X \), Eqns (32) give the required period of a neutron star:
\[ P_{NS} = 2\pi b_q^{1/2} \epsilon_e^{1/4} B_Q^{1/2} R_{NS}^{1/2} c^{3/4} L_X^{-1/4} = 10 b_q^{1/2} \epsilon_e^{1/4} L_{X,42}^{-1/4} \text{ msec} \] (34)
The corresponding spin-down time
\[ t_\Omega = b_q^{-1/2} \epsilon_e^{1/2} \frac{c^3 I_{NS}}{2 B_Q L_X^{1/4} R_{NS}^{1/4}} = 215 \epsilon_e^{-1/2} b_q^{-1} L_{X,42}^{-1/2} \text{ days} \] (35)
A surface field just below the quantum field, \( b_q \sim 0.5 \) is required to have spindown time longer than the time of Chandra observations at 328 days.

The initial rotational energy of the NS is high, but not extreme:
\[ E_{NS} = L_{sd} t_\Omega = \frac{1}{2} \epsilon_e b_q \frac{c^3 I_{NS} L_X^{1/4}}{B_Q R_{NS}^{1/4}} = 1.8 \times 10^{39} b_q^{-1} L_{X,42}^{1/2} \epsilon_e^{-1/2} \text{ erg} \] (36)
As a check, rotational energy of the neutron star is somewhat larger than the expected ejecta energy, \( E_{NS} \geq E_{ej} \), Eq. (6).

Thus, the requirement on the properties of the central source from X-ray observations are fairly mild: central neutron star with surface fields below the quantum field, spinning at \( \sim 10 \) milli-seconds (not a millisecond magnetar).

### 4.3 X-rays: emission from the relativistic termination shock: power, frequency and variability

As we discussed in the Introduction, it is the late X-ray emission that is the most energetically demanding, both in terms of the photon energy and the emitted power. Explaining late X-ray emission is most challenging; even the most energetically powerful, collimated and relativistically beamed GRBs typically do not produce X-ray emission after a year.

Following Lyutikov & Toonen (2019), we outline a PWNe-like picture for the high energy emission, Fig. 7, that the X-ray emission is generated at the termination shock of a long lasting relativistic wind (see also Lyutikov 2011b; Lyutikov & Camilo Jaramillo 2017; Barkov et al. 2021, with applications to GRBs proper).

Unlike the case of emission from the FS, where magnetic field needs to be amplified, at the termination shock the magnetic
field is supplied by the central source. At the location of the reverse shock the magnetic field in the plasma associated with the second, NS-generated, wind is

\[ L_{sd} = B_{RS}^2 R_{RS}^2 c \]

\[ B_{RS} = \frac{\sqrt{L_{sd}/c}}{R_{RS}} = \sqrt{\frac{M}{\sqrt{\beta_w} t}} \]

(37)

where we used \( R_{RS} \sim (L_0v_w/(\epsilon M))^1/2 \) t, Eq. (30), for the location of the shocks and approximated \( L_{sd} \sim L_0 \). At time of 328 days the shocks are located at \( \sim 10^{17} \) cm, see (30), with magnetic field \( \sim 0.1 \) Gauss. (This parameterization formally assumes mild magnetization \( \sigma \), but it is in fact applicable to high-sigma flows, as demonstrated by the resolution of the \( \sigma \)-paradox by Porth et al. 2013).

The estimate of the properties of the RS (37) formally matches the estimate at the forward shock (42) for \( \epsilon_B \sim 1 \) (\( \epsilon_B \) is a common parameter in the GRB/SNR theory, Piran 2004, the ratio of the post-shock amplified magnetic field energy density to the shocked plasma energy density) The main difference in the forward shock and RS emission is the typical Lorentz factor of the accelerated particles: in the PWN paradigm it is determined by the Lorentz factor of the wind, not the energy flux. If relativistic second wind has Lorentz factor \( \Gamma_w \), estimating the particles’ post-RS random Lorentz factor as \( \gamma \sim \Gamma_w \), the peak frequency is

\[ \epsilon_X \approx h \gamma^2 \omega_B \approx h \Gamma_w^2 \left( \frac{\epsilon B_{RS}}{e_m c} \right) = \Gamma_w^2 \frac{e h \sqrt{M}}{c m_e \sqrt{\beta_w} t} = 1 \text{keV} \left( \frac{\Gamma_w}{6} \right)^2 \left( \frac{v_{w,4}}{v_{w,4}} \right)^{-1/2} \left( \frac{t}{328 \text{days}} \right)^{-1} \]

(38)

(Like in PWN, the post-RS flow is non-relativistic in our frame, hence no bulk Lorentz factor is involved.)

The required Lorentz factor \( \Gamma_w \sim 10^5 \) is somewhat higher than is usually assumed for pulsars. For example in Crab the termination shock is located at similar distance of \( \sim 10^{17} \) cm, while magnetic field is three orders of magnitude smaller; with the wind’s Lorentz factor of \( \Gamma_w \sim 10^3 \) the peak emission falls into near IR (e.g., Lyutikov et al. 2019). But recall that the thin shell approximation overestimates the location of the reverse shock (hence underestimates the value of magnetic field). In addition, high-sigma flows will have reverse shock at smaller radii. All these effects will lead to less strict requirements on the Lorentz factor of the wind. In addition, acceleration at the termination shock may be more efficient, e.g., to additional effects of reconnection (e.g., Sironi & Spitkovsky 2011). These points are also important for the variability, see below.

Since the synchrotron cooling time at the RS is short,

\[ \tau_{cool} = \frac{m_e^2 c^5 v_{w,4}^2}{\Gamma_w^4 \epsilon W} = 4 \times 10^3 \text{sec} \left( \frac{\epsilon}{v_{w,4}} \right) \left( \frac{t}{328 \text{days}} \right)^2 \]

(39)

All the energy of the wind that is injected into accelerating particles is emitted, Eq. (32).

Another important observational fact is fast, erratic intra-day variability of the X-ray emission (Ho et al. 2019; Margutti et al. 2019). It is hard to reproduce fast variability within the forward shock scenario since the forward shock emission properties depend on the integrated quantities - central engine total energy and total matter swept. But fast variability can be reproduced within the internal shock paradigm, especially in the highly magnetized winds (Lyutikov & Toonen 2019). First, at the termination shock particles emit in the fast cooling regime, (39), hence any variation of the plasma parameters are reflected in the emission (as advocated by Lyutikov & Camilo Jaramillo 2017, for flares and sudden changes observed in GRB afterglows).

Second, variations of the plasma properties need not be global, but can be local. Examples include: (i) variations of Crab wisps (Hester 2008) due to the non-stationarity of the termination shock; (ii) Crab flare-like reconnection processes in the shocked pulsar wind (Clausen-Brown & Lyutikov 2012; Lyutikov et al. 2018); (iii) anisotropic local emission (“jet-in-jet” model of Lyutikov 2006a; Giannios et al. 2010, the latter works especially well in the exhaust jets of highly magnetized reconnection regions).

### 4.4 Radio-mm emission from the forward shock in the wind

Self-consistent models of radio emission of astrophysical sources is a notoriously difficult problem, partly because radio emission is energetically subdominant (e.g., the problem of radio emission of Crab PWN, Kennel & Coroniti 1984). Microphysics of electron injection at shock is most complicated (Riquelme & Spitkovsky 2011). In the case of highly magnetized winds, a combination of Fermi and reconnection processes are likely at play (Sironi & Spitkovsky 2014; Lyutikov et al. 2019).

A simple approach is to use the classic GRB prescription (Sari & Piran 1995; Sari 1997), see also Chevalier & Fransson (2017). The peak Lorentz factor of particles accelerated at the forward shock propagating into unmagnetized medium with velocity \( \beta_{FS} \) is expected to be

\[ \gamma_{FS} \sim \frac{m_e}{m_p} \beta_{FS}^2 = \frac{m_e}{m_e} \frac{L_{sd} v_w}{M c^3} \]

(40)

with \( \beta_{FS} \) given by (30).
In the preceding wind the magnetic field is weak, and needs to be amplified to produce synchrotron emission. The post-forward shock energy density

\[ u_{FS} \sim \rho_w \beta_{FS}^2 c^2 = \frac{L_{wd}}{4\pi R^2 c} = \frac{\dot{M}}{4\pi \dot{\nu}_w t^2} \]

(41)
gives the magnetic field

\[ B_{FS}^2 = \frac{\dot{\Sigma}_B u_{FS}}{8\pi} = \frac{\epsilon_B u_{FS}}{\dot{M}} = \frac{\sqrt{L_c}}{\sqrt{\epsilon_c}} = \frac{\sqrt{2\epsilon_B}}{\sqrt{\epsilon_c}} \frac{\dot{M}}{\dot{\nu}_w t} \]

(42)

The magnetic field (42) and post-forward shock Lorentz factor (40) give the emission frequency

\[ \frac{\omega}{2\pi} \sim \gamma_{FS}^2 \frac{eB_{FS}}{2\pi m_e c} = \frac{\sqrt{\epsilon_B}}{2\pi} \frac{e^2 m_p^2}{c^2 m_e^2} \frac{L_{wd}^{3/2}}{M^{3/2} t} = 1.43^{1/2} \frac{i_{\nu,4} v_{\nu,4}^{3/4} M_{-4}}{t_{261 \text{days}}} \] GHz

(43)
at 261 days. This matches the value of the peak emission at 261 days, and decrease with time (Fig. 7 of Yao et al. 2021). The peak emission frequency (43) is a sensitive function of the central source’s luminosity, pre-explosion wind parameters and time of observations.

The above estimates also exclude forward shock as the origin of the X-ray emission. Though the injection problem at shocks is notoriously difficult, the estimates are off by \( \sim 10 \) orders of magnitude in emitted frequency, \( \sim 5 \) orders of magnitude in particle energy. Recall, that the late X-ray emission is the dominant energy channel in AT2020mrf. To bring emission energy to the keV range, the Lorentz factor of the X-ray emitting particles should exceed the peak Lorentz factor (40) by factor \( \sim 10^5 \). In that case the total energetics is higher by a similar factor (for distribution function \( f \propto \gamma^{-2} \)).

5 CONTRIBUTION FROM RADIOACTIVITY

The envelope ejected during AIC will contain some \( ^{56}\text{Ni} \); up to \( \sim \) few times \( 10^{-2} \) or \( ^{56}\text{Ni} \) are expected (Fryer et al. 1999). This will not affect the FBOT light curve (half-life time of \( ^{56}\text{Ni} \) is 6.1 days), but this does affect the longer light curves. For example, a decay of \( 10^{-2} M_\odot \) of \( ^{56}\text{Ni} \) ejected would produce (at 3.6 MeV emitted as radiation per decay), total electromagnetic energy \( 1.3 \times 10^{48} \) erg, comparable to the overall energetics (if reprocessed to optical).

An important issue here is that for small ejecta masses the radiative energy is more efficiently converted into radiation, as it is not degraded by adiabatic losses in the typical months-long optically thick stage of expansion. The ejecta becomes fully transparent at time (7), just somewhat longer than the \( ^{56}\text{Ni} \) decay half-life. Hence, unlike the conventional SNIa where most of the \( ^{56}\text{Ni} \) energy is “wasted” into the kinetic energy of expansion, a larger fraction is radiated now.

6 TRANSIENTS FOLLOWING WDs MERGERS

Mergers of WDs is one of the most frequent catastrophic events (Shen et al. 2012; Schwab et al. 2016). Supernova of Ia type is the most frequently discussed channel of WDs mergers (Maoz et al. 2014; Soker 2019). The rates of SNIa are much smaller than of the WD mergers (Yungelson & Kuranov 2017). Mergers of CO-CO WDs with a combined mass above Chandrasekhar, is the most frequently discussed channel of WDs mergers (Maoz et al. 2014; Soker 2019). The rates of SNIa are much smaller.
found that during the merger itself very little mass, \( \sim 10^{-3} M_\odot \), is ejected hydrodynamically. If AIC occurs \( \geq 100 \) yrs after the merger, that primary ejected shell is reached by the AIC ejecta/shock on time scale of a year. Dall’Osso et al. (2014) argued that the pulsar J0737-3039B was a product of an AIC of ONeMg WD with small shell mass; in our model this would require that the initial system was a triple.

Finally, there are possibly related Type Iax class of supernovae (subluminous Ia-s, Foley et al. 2013) and Type Ibn (Hosseinzadeh et al. 2017). Type Iax may be the results of AICs with larger ejecta mass \( M_{ej} \geq 0.1 \); Type Ibn may be the result of merger with a He WD. (Somewhat similarly Nicholl et al. 2015, argued that the ejecta mass is the dominant factor affecting the light curves ) We live investigation of these possibilities to a future research.

7 CONCLUSION AND PREDICTION

The present model of broad band emission from FBOTs is highly constrained - it reproduces in a self-consistent manner the optical FBOTs, late radio-mm emission and the dominant long term X-ray emission. One of the key advantages of the model is the low requirement on energetics, Eq. (36), only few \( \times 10^{50} \) ergs, just a bit higher than the observed X-ray energy. This is a consequence of small ejecta mass. This low energy budget compares favorably with GRB-like energy budgets of \( \geq 10^{52} \) ergs in models that advocate large ejecta mass (e.g., Gottlieb et al. 2022; Chen & Shen 2022). Our inferred rotational period of few milliseconds is consistent with possible detection of periodicity by Pasham et al. (2021).

There are several principal points we make. First, the present model naturally explains small ejecta mass. This is due to the competition of envelope mass lost to the wind, and ashes added to the core. FBOTs results from AIC when the shell mass is \( \leq 0.1 M_\odot \). If AIC occurs with larger shell mass then the shell remains optically thick longer and most of the energy dissipated at the neutron star-driven shock is not radiated, but is spent on \( pdV \) work. FBOTs thus correspond to a narrow range of initial parameters of the binary system - this also explains their rarity. FBOTs’ emission is generated as the radiation-dominated forward shock propagates through ejecta.

The second major point is that the overall energetics is dominated in the case of AT2020mrf by long-term X-ray emission. Both the production of X-rays at times \( \sim \) one year, and the total energetics are the most demanding. Even the most energetic GRBs do not produce X-ray that late. We outlined a PWN-type model of high energy emission, where X-rays are generated at the relativistic termination shock (not the forward shock). The demands on the central neutron star are not extreme - a combination of high magnetic field and fast spins (e.g., 10 milliseconds) is needed, but not a millisecond magnetar with quantum field. Both the early and late X-rays are produced at the NS-driven termination shock: first inside the ejecta, and later in the preceding wind. Since X-ray emission is in the fast cooling regime (high magnetic fields in the wind) a large fraction of the wind energy is radiated.

Third, the duration of the FBOTs is determined by the diffusion of photons produced deep in the ejecta, in the regions of high optical depth, Eqns. (4-5). This is somewhat new regime, as typically supernova light curves were determined mostly by distributed source of photons (bremsstrahlung, recombination and radioactive decay). In our case a powerful, highly non-equilibrium source of photons is located deep in the ejecta.

In the case of centrally produced photons, within a high optical depth ejecta, nearly all the photons escape at the same time (assuming spherical symmetry of the ejecta - a far from certain assumption). Thus, the peak flux is an integrated property: all the photons produced by the NS shock before \( t_{FBOT} \) escape within a narrow time near \( t_{FBOT} \).

Another constraint comes from the requirement that the NS-generated shock should break through ejecta to produce the SRG/eROSITA X-ray bump. For this the energy deposited by the NS should be larger than the energy of the ejecta. Sufficiently light ejecta, and sufficiently powerful NS are needed. SRG /eROSITA X-ray bump corresponds to the break-out of the shock from the ejecta into the preceding wind; this was predicted by Lyutikov & Toonen (2019).

The present model explains

- short rise times of FBOTs (light ejecta)
- blue color (radiation-dominated shock)
- decreasing photospheric radius (Perley et al. 2019)
- change of properties at \( \sim 1 \) month (shock exiting the ejecta into the wind)
- predicted the X-ray bump seen by eROSITA in AT2020mrf (due to the shock breakout)
- long late X-ray emission (from the termination shock)
- long radio-mm (from the forward shock)
- presence of dense circumburst medium (inferred by Bright et al. 2022; Nayana & Chandra 2021) - this is the post-merger/pre-collapse wind
- Host galaxies: as discussed by Lyutikov & Toonen (2019), their Fig. 6, the merger rates of the CO-ONeMg WDs peaks at short delay times of about \( \sim 50-100 \)Myr, with a long tail to long delay times. Thus the typical delay time of the CO-ONeMg mergers is closer to that of core-collapse supernovae. Lyutikov & Toonen (2019) expected the host galaxies of CO-ONeMg mergers to be more similar to those of core-collapse supernovae instead of Type Ia supernovae. This explains detection of

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FBOTs preferentially in star-forming galaxies (Prentice et al. 2018; Lyman et al. 2020). We expect further detections in the early-type galaxies though.

The model has a number of predictions:

- The system should be hydrogen poor, but not necessarily hydrogen free. If the secondary WD was of the DA type, the most common, the ejecta and the wind will likely have $\sim 10^{-4}M_{\odot}$ of hydrogen. (Fox & Smith 2019; Pellegrino et al. 2021, noted that H-poor spectra of SN-Ibn are similar to that of AT 2018cow at later times). An advanced radiation transfer modeling is required.
- One expects evidence (e.g., spectral) of shock interaction with fast and dense pre-explosion wind
- Anisotropy is expected: both the ejection of the envelope during AIC, as well as neutron star-driven winds are expected to be anisotropic.
- Pre-FBOT archival data should show a bright persistent hydrogen-poor source, possibly surrounded by a nebula (similar to IRAS 00500+6713, Gvaramadze et al. 2019; Oskinova et al. 2020)
- In the case when AIC does not occur, a hydrogen poor envelope around the central (massive) WD is expected. Its observability depends on how long ago the last part of the envelope was lost. In the case of massive WD observed by Caiazzo et al. (2021), the age is few million years, so the envelope is dissipated by now.

ACKNOWLEDGMENTS
I would like to thank Igor Andreoni, Ilaria Caiazzo, Paul Duffel, Ori Fox, Dimitrios Giannios, Daniel Kasen, Danny Milisavljevic, Lida Oskinova, John Raymonds, Noam Soker, Bhagy Subrayan, Silvia Toonen, Alexander Tutukov, Beatriz Villarroel for discussions. Comments by late Vasilii Gvaramadze are acknowledged. This work had been supported by NASA grants 80NSSC17K0757 and 80NSSC20K0910, NSF grants 1903332 and 1908590.

8 DATA AVAILABILITY
The data underlying this article will be shared on reasonable request to the corresponding author.

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