Supercollapsars and their X-ray bursts

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ABSTRACT

The very first stars in the Universe can be very massive, up to $10^3 \, M_\odot$. If born in large numbers, such massive stars can have a strong impact on the subsequent star formation, producing strong ionizing radiation and contaminating the primordial gas with heavy elements. They would leave behind massive black holes that could act as seeds for growing supermassive black holes of active galactic nuclei. Given the anticipated fast rotation, such stars would end their life as supermassive collapsars and drive powerful magnetically dominated jets. In this Letter, we investigate the possibility of observing the bursts of high-energy emission similar to the long gamma-ray bursts associated with normal collapsars. We show that during the collapse of supercollapsars, the Blandford–Znajek mechanism can produce jets as powerful as few $\times 10^{52} \, \mathrm{erg \, s^{-1}}$ and release up to 10^{54} erg of the black hole rotational energy. Due to the higher intrinsic time-scale and higher redshift, the initial bright phase of the burst can last for about 10^4 s, whereas the central engine would remain active for about 1 d. Due to the high redshift the burst spectrum is expected to be soft, with the spectral energy distribution peaking at around 20–30 keV. The peak total flux density is relatively low, $10^{-7} \, \mathrm{erg \, cm^{-2} \, s^{-1}}$, but not prohibitive. If one supercollapsar is produced per every minihalo of dark matter arising from the 3σ cosmological fluctuations, then the whole sky frequency of such bursts could reach several tens per year.

Key words: black hole physics – magnetic fields – relativity – early Universe – gamma-rays: bursts – X-rays: bursts.

1 INTRODUCTION

According to the modern hierarchical clustering theories of galaxy formation, the first stars are born within collapsed haloes of dark matter of $\simeq 10^6 \, M_\odot$ at $z \simeq 20$. The primordial gas falls into the potential well of these haloes and fragments into clamps of $\simeq 10^3 \, M_\odot$ via gravitational instability (Bromm, Coppi & Larson 2002). Because this gas is metal-free, its cooling is rather slow and further fragmentation into smaller clamps seems to be avoided (cf. Stacy, Greif & Bromm 2009; Turk, Abel & O'Shea 2009). Instead, the clumps contract in a quasi-static fashion as a whole, suggesting that the first stars can be very massive indeed, $M > 100 \, M_\odot$. The actual initial mass function (IMF) of first metal-free stars (Population III stars), however, is not known yet since too many factors come into play, making the problem intractable analytically and rather challenging numerically. In particular, the initial mass of protostars can be very small, down to $10^{-3} \, M_\odot$, and the eventual accumulation of mass proceeds via accretion of the surrounding gas. A very high accretion rate, $\dot{M} > \dot{M}_c \simeq 4 \times 10^{-3} \, M_\odot \, \mathrm{yr}^{-1}$, may limit the final mass to few hundreds of solar masses as the protostellar luminosity reaches the Eddington limit (Omukai & Palla 2003). For a lower accretion rate, the accretion may proceed even after the onset of nuclear burning in the stellar core and result in the final mass $M \simeq 10^3 \, M_\odot$. Numerical studies of cosmological gravitational instability suggest that, although in principle the accretion rate can be as high as few $\times 10^{-2} \, M_\odot \, \mathrm{yr}^{-1}$, in reality the rotational support against gravity often becomes important and reduces the rate below $\dot{M}_c$ (Gao et al. 2007). Ohkubo et al. (2009) studied the evolution of accreting Population III stars from the pre-main-sequence evolution to the core-collapse and confirmed that the final mass can be as large as $10^3 \, M_\odot$. Very massive first stars are also predicted in theories involving dark matter annihilation (e.g. Natarajan, Tan & O'Shea 2009).

Population III stars with masses $140 \, M_\odot \leq M \leq 260 \, M_\odot$ most likely end their life as pair-instability supernovae which leave no compact remnant behind (Fryer, Woosley & Heger 2001). If such stars were the main outcome of initial star formation, they would overproduce heavy elements in the early Universe, in conflict with

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the observations of extremely metalpure stars in Galactic bulges and the observed abundances of intergalactic and intercluster media (Umeda & Nomoto 2002; Heger & Woosley 2002; Chieffi & Limongi 2002).

More massive stars, which will be referred to as Very Massive Stars (VMSs), are expected to collapse into black holes with very little mass loss (Fryer et al. 2001). They would leave behind massive black holes (MBHs), which could play the role of seeds for the supermassive black holes (SMBHs) of active galactic nuclei (AGNs). Assuming that MBHs are formed at the rate of one per minihalo developed from a 3σ fluctuation, Madau & Rees (2001) estimated their density to be around 5000 per galaxy like the Milky Way making the total mass of MBHs comparable to the current total mass of SMBHs. This suggests that SMBHs could form via mergers of MBHs, the idea that has been actively developed in recent years. Even more massive VMSs, with $M \simeq 10^6 \, M_\odot$, could be formed in more massive dark matter haloes, with total mass $M \simeq 10^8 \, M_\odot$, collapsed at $z \simeq 10$ (Bromm & Loeb 2003; Begelman, Volonteri & Rees 2006). Although much more rare, such events can provide an alternative way of producing SMBHs.

From the observational perspective, it is difficult to distinguish between a VMS and a cluster of less massive Population III stars. This suggests us to investigate the potential observational consequences of VMS collapse, which could be quite spectacular because of the very high mass involved. Given their expected fast rotation, it seems likely that supercollapsars (classified as type-III collapsars in Heger et al. 2003) develop accretion discs, drive relativistic jets and produce bursts of high-energy emission in the fashion similar to their less massive relatives (MacFadyen & Woosley 1999; Barkov & Komissarov 2008). If detected, they would become the most distant strong sources of light and provide us with a new way of probing the physical conditions in the epoch close to the 'Dark ages'. Even non-detection could be useful, allowing us to set constraints on models of star formation in the early Universe and the origin of SMBHs. There has been a number of papers looking into the cosmological evolution of gamma-ray bursts (GRBs), including those from the Population III stars (e.g. Bromm & Loeb 2002; Le & Dermer 2007; Naoz & Bromberg 2007). They assumed that the Population III GRBs are similar to those from the lower redshift Population II stars whose mass is significantly below $100 \, M_\odot$. However, because of their very high mass and redshift, the GRB-like bursts of supercollapsars can be rather special.

There are two crucial differences between a normal collapsar and a supercollapsar. One is that instead of a protoneutron star of solar mass, the supercollapsars develop protoblack holes of tens of solar masses, within which the neutrinos from electron capture are trapped (Fryer et al. 2001; Suwa et al. 2007). The other is that the accretion discs of supercollapsars are far too large and cool for the neutrino annihilation mechanism. This has already been seen in the numerical simulations of a supercollapsar with mass $M = 300 \, M_\odot$ (Fryer et al. 2001). Utilizing the study of hyper-accreting discs by Beloborodov (2008), we find that at best the rate of heating due to this mechanism is

$$\dot{E} \simeq 2 \times 10^{48} \dot{M}_0^{9/4} M_{h,3}^{-3/2} \, \text{erg s}^{-1}, \qquad (1)$$

where $\dot{M}$ is the accretion rate and M_h is the black hole mass. (Here and in other numerical estimates below we use the following notation: $\dot{M}_k$ is the mass accretion rate measured in the units of $10^k \, M_\odot \, \text{s}^{-1}$ and M_k is the mass measured in the units of $10^k \, M_\odot$.) Such low values have led Fryer et al. (2001) to conclude that the magnetic mechanism is the only candidate for producing GRB jets from supercollapsars. In the following, we analyse one particular

version of the mechanism, namely the one where the jets are powered by the rotational energy of the black hole via the Blandford–Znajek (BZ) process (Blandford & Znajek 1977; Barkov & Komissarov 2008).

2 BLANDFORD–ZNAJEK JETS FROM SUPERCOLLAPSARS

VMSs are expected to rotate rapidly, close to the break-up speed and hence produce rapidly rotating MBHs. Moreover, in the absence of a strong magnetic field in the pristine primordial gas, VMSs will be weakly magnetized and as a result could develop rapidly rotating cores (Woosley & Heger 2006). This suggests that the spin parameter of MBHs can be very high, $a \simeq 1$, yielding the enormous rotational energy $E_{\text{rot}} \simeq 5 \times 10^{56} M_{h,3}$ erg. In order to estimate the BZ luminosity, we need to know the strength of the magnetic field accumulated by the hole. The usual approach is to relate it to the gas pressure in the disc and for this we need to know the parameters of the accretion disc itself. Accurate determination of these parameters and their time evolution requires us to know the structure of VMS prior to the collapse and the physics of the accretion disc. Unfortunately, this information is lacking at the moment. In particular, although the structure of supermassive stars $M_s \gg 10^3 \, M_\odot$ is very well described by a polytropic model with $n = 3$ (Zeldovich & Novikov 1971), we are more interested in stars with $M_s \leq 10^3 \, M_\odot$. Fryer et al. (2001) studied the structure and evolution of a $300 \, M_\odot$ star. Prior to the collapse, the star entered the red giant phase and expanded from $R_s = 4 \times 10^{12}$ to 1.5×10^{14} cm. However, the initial rotation rate of this star was small compared to the one required in the single-star model of GRB progenitors (about $\simeq 50$ per cent of the break-up speed). In this model, the progenitors remain chemically homogeneous and compact all the way up to the collapse (Yoon & Langer 2005; Woosley & Heger 2006). Given this lack of information about the progenitor structure, we will follow Bethe (1990) and assume the $\rho \propto R^{-3}$ ($R_c \leq R \leq R_s$) distribution of mass density in the star prior to collapse. In fact, this distribution agrees reasonably well with the numerical models of rapidly rotating low metallicity stars considered as likely progenitors of a normal GRB (see fig. 2 in Kumar, Narayan & Johnson 2008). As to the stellar rotation, we will assume that it is uniform ($\Omega = \text{constant}$) in the stellar envelope, with 50 per cent of break-up speed at the stellar surface.

Due to the slow neutrino cooling, the accretion discs of supercollapsars are expected to be radiatively inefficient, with possible exception only for the very inner region. This suggests us to use the Advection Dominated Accretion Flow (ADAF) model (Narayan & Yi 1994) to describe these discs. Since the radiation pressure dominates, we can use the ratio of specific heats $\gamma = 4/3$ which gives us

$$v_{\text{in}} \simeq \frac{3\alpha}{7} v_k, \qquad c_s^2 \simeq \frac{2}{7} v_k^2, \qquad H \simeq R c_s / v_k, \qquad (2)$$

where v_{in} is the accretion speed, c_s is the sound speed, $v_k = \sqrt{GM/R}$ is the Keplerian speed, H is the vertical disc scale and α is the effective viscous stress parameter of the α-disc model (Shakura & Sunyaev 1973). The disc density and pressure can be estimated combining the above equations with the expression for the mass accretion rate, $\dot{M} \simeq 2\pi R H \rho v_{\text{in}}$. Straightforward calculations yield

$$P \simeq \frac{\sqrt{14}}{12\pi\alpha} \frac{\dot{M}(GM)^{1/2}}{R^{5/2}}. \qquad (3)$$

The poloidal magnetic field should scale with the thermodynamic pressure, so we write $B^2 = 8\pi P/\beta$, where β is the magnetization parameter. Applying this equation to the radius of the marginally

bound orbit $R_{mb} = f_1(a)R_g$, where $f_1(a) = 2 - a + 2(1 - a)^{1/2}$ and $R_g = GM_h/c^2$ is the BH's gravitational radius, we find

$$B_{mb} \simeq 3 \times 10^{13} f_1^{-5/4} \beta_1^{1/2} \dot{M}_0^{1/2} \text{ G}, \tag{4}$$

where $\beta_1 = \beta/10.$[1] If the magnetic field is generated in the disc, then it is likely to change polarity on the scale $\simeq H$. This may lead to significant variation in the strength and polarity of the magnetic field accumulated by the black hole and reduce the BZ luminosity (e.g. Barkov & Baushev 2009). We will assume that this effect is accounted for in the value of β.

The power of the jet energized via the BZ mechanism can be estimated using the monopole solution for magnetospheres of rotating black holes (Blandford & Znajek 1977), which gives

$$L_{BZ} = \frac{1}{3c} \left(\frac{\Psi_h \Omega_h}{4\pi} \right)^2, \tag{5}$$

where $\Omega_h = f_2(a)c^3/GM_h$ is the angular velocity of a BH, $f_2(a) = a/2(1 + \sqrt{1 - a^2})$ and Ψ_h is the magnetic flux threading one of the BH's hemispheres. Inside the marginally bound orbit, the disc plasma quickly dives into the BH and the magnetic flux can be roughly estimated as $\Psi = 2\pi R_{mb}^2 B_{mb}$ (Reynolds, Garofalo & Begelman 2006). Combining this result with equations (4) and (5), we find

$$L_{BZ} \simeq \frac{\sqrt{14}}{9} \frac{f_1^{3/2} f_2^2}{\alpha\beta} \dot{M}c^2 \simeq \frac{0.05}{\alpha_{-1}\beta_1} \dot{M}c^2, \tag{6}$$

where $\alpha_{-1} = \alpha/0.1$ (for $0.5 < a < 1$, the combination $f_1^{3/2} f_2^2$ depends weakly on a and is approximately $1/4$).

The mass accretion rate can be estimated following the procedure described in Barkov & Komissarov (2009). The total accretion time includes the travel time of the rarefaction wave sent into the stellar envelope by the core collapse, the time of the envelope collapse and the disc accretion time, which gives the largest contribution. Accounting only for the disc contribution, the accretion time-scale for the stellar matter located in the progenitor at radius R is $t \propto l^3/M^2$, where $l = \Omega R^2$ and $M(R)$ is the stellar mass enclosed within radius R. Then $\dot{M} = dM/dt \simeq (dM/dR)/(dt/dR)$, where for Bethe's model we have $dM/dR \simeq M_s/(R \ln(R_s/R_c))$ and $dt/dR \simeq 6t/R$, where R_c is the stellar core radius. Collecting the results, we obtain

$$\dot{M} \simeq \frac{1}{6} \frac{M_s}{\ln(R_s/R_c)} \frac{1}{t} \simeq 36 \frac{M_{s,3}}{t} \text{ M}_\odot \text{ s}^{-1}, \tag{7}$$

where t is measured in seconds and we used $R_s/R_c = 100.$[2] Here, we assumed that the whole of the disc is accreted by a BH, following the original ADAF model. However, it has been argued that this model has to be modified via including disc wind (Advection Dominated Inflow Outflow Solution (ADIOS); Blandford & Begelman 1999], which implies a mass loss from the disc and a smaller accretion rate compared to equation (7). While the arguments for disc wind are very convincing, the actual value of mass loss is not well constrained and can be rather low. Given equation (7), the power of the BZ jet is

$$L_{BZ} \simeq \frac{3.2 \times 10^{52} \epsilon_m}{\alpha_{-1}\beta_1} \frac{M_{s,3}}{t_2} \text{ erg s}^{-1}, \tag{8}$$

where $t_2 = t/100$ and $\epsilon_m < 1$ is the fraction of the disc mass reaching the BH. Given the jet propagation speed inside the star, $v_j \simeq 0.2c$, deduced from axisymmetric numerical simulations (Barkov & Komissarov 2008), the jet breakout time is expected to be of around a few hundred seconds and, thus, the numerical factor in equation (8) gives us the optimistic jet power at the time when it becomes observable. In fact, the initial influx of mass through the polar column is very large and activation of the BZ mechanism can be delayed (Komissarov & Barkov 2009). The very latest time for the activation is given by the free-fall time of the whole star:

$$t_{ff} \simeq 1000 R_{s,12}^{3/2} M_{s,3}^{-1/2} \text{ s}, \tag{9}$$

as by this time the polar column becomes completely empty.

The total duration of the jet production phase has to be similar to the disc lifetime. If VMS is rotating at half of the break-up speed, then the initial outer edge of the disc is at $R_d \simeq R_s/4$. Ignoring the edge expansion due to accumulation of angular momentum, the disc lifetime is given by its 'viscous' time-scale

$$t_{ce} \simeq \frac{2R_s}{3v_{in}(R_s)} \simeq 5000\alpha_{-1}^{-1} R_{s,12}^{3/2} M_{s,3}^{-1/2} \text{ s}, \tag{10}$$

where $R_{s,12}$ is the stellar radius measured in 10^{12} cm. By this time, the BZ power will be significantly reduced but could still play a role in shaping the light curve of afterglow emission (Barkov & Komissarov 2009).

Using the mass accretion rate given by equation (7), we can check if the neutrino cooling needs to be included in the model. Under the conditions of the supercollapsar's disc, its cooling is dominated by pairs. Using the well-known equation for this cooling rate (e.g. Yakovlev et al. 2001), we can compare the cooling time with the accretion time at a given disc radius. The result is

$$\frac{t_d}{t_{cool}} \simeq 0.3\alpha_{-1}^{-9/4} (R/R_g)^{-13/8} \dot{M}_{-1}^{5/4} M_{s,3}^{-3/2}. \tag{11}$$

Thus, except for the very inner part of the disc, the neutrino cooling is indeed inefficient.

The high BZ power given by equation (8) suggests that the GRB-like burst emission from such jets could be seen even from $z \simeq 20$ and in the next section we discuss the properties of such bursts in more details.

3 OBSERVATIONAL SIGNATURES

Assuming that the radiation mechanism of the supercollapsar jets is similar to that of normal GRB jets, we expect the peak in the spectral energy distribution of the prompt emission in the source frame to be around 0.5 MeV. However, the cosmological redshift effect reduces the peak down to

$$E_{max} \simeq 25 \text{ keV} \left(\frac{1 + z}{20} \right)^{-1}, \tag{12}$$

which is still inside the energy window of *Swift*'s Burst Alert Telescope (BAT). For the same reason, the observed total duration of the burst increases up to

$$t_b \simeq 1 \left(\frac{1 + z}{20} \right) \alpha_{-1}^{-1} R_{s,12}^{3/2} M_{s,3}^{-1/2} \text{ d}. \tag{13}$$

The characteristic source frame time-scale for the decay of BZ luminosity in the model presented above is given by the time since the onset of the collapse. Thus, the initial time-scale for the burst decay will be of the order of the jet breakout time, few $\times 10^2$ s, or a bit longer if the activation of the BZ mechanism is significantly delayed. In the observers frame, this translates into few

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[1] The inner edge of the disc varies between the marginally bound and the marginally stable orbits, depending on its thickness. In our calculations, this does not make much difference.

[2] For a 300 M$_\odot$ star at $t \simeq 8$ s, this gives $\dot{M} \simeq 1.3$ M$_\odot$ s^{-1}, which agrees reasonably well with the numbers given in Fryer et al. (2001).

$\times 10^3$–10^4 s. Thus, these bursts would appear not only unusually soft but also unusually long-lasting.

The total flux density of the burst emission received on Earth and the isotropic luminosity are related via $F = L/(4\pi r_L^2)$, where r_L is the luminosity distance to the source (e.g. Peebles 1993). However, the emission from GRB jets is highly anisotropic due to the relativistic beaming. Moreover, not all of the BZ power is converted into the radiation within the energy window of the receiver. This leads to

$$F = \epsilon_c \frac{L_{BZ}}{4\pi r_L^2 \mathcal{A}}, \tag{14}$$

where $\mathcal{A} \ll 1$ is the solid angle of the radiation beam and $\epsilon_c < 1$ is the conversion efficiency. In flat Universe,

$$r_L = \frac{c}{H_0}(1+z) \int_0^z (\Omega_m(1+z)^3 + \Omega_\Lambda)^{-1/2} dz. \tag{15}$$

For $z = 20$ and the density parameters $\Omega_\Lambda = 0.72$, $\Omega_m = 0.28$ (Komatsu et al. 2009) this gives us

$$F \simeq 2 \times 10^{-7} \frac{\epsilon_{c,-1}\epsilon_m}{\alpha_{-1}\beta_1 \mathcal{A}_{-3}} M_{s,3} t_3^{-1} \text{ erg cm}^{-2} \text{ s}^{-1}, \tag{16}$$

where $\epsilon_{c,-1} = \epsilon_c/0.1$, $\mathcal{A}_{-3} = \mathcal{A}/10^{-3}$ and $t_3 = t/10^3$. One can see that for the first 10^4 s, this is above the sensitivity of BAT, 10^{-8} erg cm^{-2} s^{-1}, and thus such a burst could trigger BAT. Having said this, we keep in mind that there is a great deal of uncertainty with respect to the values of various parameters appearing in equation (16).

The time dependence in equation (16) gives the evolution of the mean bolometric flux. It is not clear if the supercollapsar bursts will also exhibit the fine substructure characteristic of normal GRBs. If the variability of normal GRBs is due to internal shocks in baryon-dominated flow, as this is proposed in the currently most popular model of prompt gamma-ray emission (Mészáros & Rees 1994), then the supercollapsar burst produced by the magnetically dominated BZ jet may well be smooth and featureless. However, there are models of normal GRBs that attribute the observed variability to unsteady magnetic dissipation (Lyutikov & Blandford 2003; Giannios et al. 2009; Kumar & Narayan 2009). If they are correct, then the supercollapsar bursts will also show fine substructure.

In order to estimate the observed rate of such a burst we assume, following Madau & Rees (2001), that the dark matter minihaloes that host supercollapsars arise from 3σ fluctuations that constitute only $\simeq 0.3$ per cent of the dark matter of the Universe and that only one supercollapsar per minihalo is produced. The total mass per Mpc^3 at $z = 20$ is $M_{Mpc} \simeq 1.5 \times 10^{15}$ M$_\odot$. The number density of 3σ minihaloes is then

$$n_{mh} \simeq 0.003 \frac{\Omega_{dm} M_{Mpc}}{10^6 M_\odot} \simeq 10^6 \text{ Mpc}^{-3}.$$

Let us assume, for the sake of simplicity, that all supercollapsars go off simultaneously at cosmological time t_e corresponding to $z = 20$ (a moderate spread around this redshift will not significantly change the result). In flat Universe, the observed time separation between events occurring simultaneously at r_0 and $r_0 + dr_0$, where r_0 is the comoving radial coordinate, is $dt_o = c\, dr_0$. The corresponding physical volume within 1 sr of the BAT's field of view is

$$dV = a^3(t_e)r_0^2 dr_0,$$

where $a(t_e) = (1+z)^{-1}$ is the scaling factor of the Universe at $t = t_e$ [in the calculations, we fix the scaling factor via the condition

$a(t_o) = 1$]. r_0 and t_e are related via $r_0 = r_L(1+z)^{-1}$. Putting all this together, we find the rate to be

$$f_c = \mathcal{A}\frac{c n_{mh} r_L^2}{(1+z)^5} \simeq 4\, \mathcal{A}_{-3}\left(\frac{n_{mh}}{10^6}\right) \text{ yr}^{-1} \text{ sr}^{-1}.$$

Recent high-resolution simulations of cosmological star formation indicate the possibility of further fragmentation of gas clumps in minihaloes, resulting in the formation of binary or even multiple protostars in some realizations (Stacy et al. 2009; Turk et al. 2009). Thus, the theoretical rate of VMS formation can be significantly smaller compared to the one used in our calculations, making the supercollapsar bursts rare events. This may explain why such bursts have not been seen so far.

4 CONCLUSIONS

In spite of the significant progress in the astrophysics of GRBs, both observational and theoretical, it may still take quite a while before we fully understand both the physics of the bursts and the nature of their progenitors. At the moment, there are several competing theories and too many unknowns. Similarly, we know very little about the star formation in the early Universe. For this reason, the analysis presented above is rather speculative and the numbers it yields are not very reliable. Further efforts are required to develop a proper theory of supercollapsars and to make firm conclusions on their observational impact. On the other hand, our estimates suggest that if we are on the right track, then the X-ray bursts of supercollapsars may already be detectable with *Swift*. The expected very long duration of bursts and their relatively low brightness imply that a dedicated search programme using the image trigger may be required. Such search would be useful even in the case of non-detection as this would put important constraints on models of star formation in the early Universe, models of the GRB progenitors and the origin of SMBHs.

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