Exploring the Properties of Choked Gamma-ray Bursts with IceCube’s High-energy Neutrinos

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Received 2017 November 7; revised 2018 January 26; accepted 2018 January 26; published 2018 March 5

Abstract

Long duration gamma-ray bursts (GRBs) have often been considered the natural evolution of some core-collapse supernova (CCSN) progenitors. However, the fraction of CCSNe linked to astrophysical jets and their properties are still poorly constrained. While any successful astrophysical jet harbored in a CCSN should produce high-energy neutrinos, photons may be able to successfully escape the stellar envelope only for a fraction of progenitors, possibly leading to the existence of high-luminosity, low-luminosity, and not-electromagnetically bright (“choked”) GRBs. By postulating a CCSN–GRB connection, we accurately model the jet physics within the internal-shock GRB model and assume scaling relations for the GRB parameters that depend on the Lorentz boost factor \( \Gamma \). The IceCube high-energy neutrino flux is then employed as an upper limit of the neutrino background from electromagnetically bright and choked GRBs to constrain the jet and the progenitor properties. The current IceCube data set is compatible with up to 1% of all CCSNe harboring astrophysical jets. Interestingly, those jets are predominantly choked. Our findings suggest that neutrinos can be powerful probes of the burst physics and can provide major insights on the CCSN–GRB connection.

Key words: gamma-ray burst: general – neutrinos – supernovae: general

1. Introduction

Gamma-ray bursts (GRBs) are among the most energetic astrophysical transients (Mészáros 2006; Kumar & Zhang 2014; Mészáros 2017). GRBs are expected to be sources of high-energy neutrinos produced through hadronic and lepto-hadronic interactions. Neutrinos from GRBs could be possibly detected by neutrino telescopes. However, targeted searches of neutrinos from high-luminosity GRBs (HL-GRBs) have reported evidence for a lack of statistically significant temporal and spatial correlation of neutrino data (Schmid & Turpin 2015; Aartsen et al. 2017b) constraining the proposed theoretical models (Baerwald et al. 2012; He et al. 2012; Bartos et al. 2013).

At the same time, a flux of astrophysical neutrinos has been detected by the IceCube Neutrino Observatory (Aartsen et al. 2013a, 2013b, 2014, 2015a, 2015b, 2015c, 2015d). While the flux is held to be predominantly extragalactic (Palladino & Vissani 2016; Aartsen et al. 2017a; Denton et al. 2017), its origin is currently unknown. In this context, “choked” GRBs (see, e.g., Mészáros & Waxman 2001; Ando & Beacom 2005) have been considered as possible sources of some of the IceCube neutrinos (Murase & Ioka 2013; Tamborra & Ando 2015, 2016; Murase et al. 2016; Senno et al. 2016, 2018).

A choked jet is one where the jet is successful in accelerating particles, but electromagnetic radiation is unsuccessful in escaping the stellar envelope. Choked GRBs may even be more abundant than the GRBs ordinarily observed in photons (Mészáros & Waxman 2001; Razzaque et al. 2003a, 2003b, 2004; Ando & Beacom 2005; Horiiuchi & Ando 2008; Murase & Ioka 2013). Neutrinos and possibly gravitational waves may be the only messengers from these sources.

There is solid evidence that GRBs and core-collapse supernovae (CCSNe) are related (Paczynski 1998; Modjaz 2011; Hjorth & Bloom 2012; Lazzati et al. 2012; Hjorth 2013; Margutti et al. 2014; Sobacchi et al. 2017). This was also foreseen in the so-called collapsar model (MacFadyen & Woosley 1999; MacFadyen et al. 2001; Woosley & Bloom 2006) and is supported by the fact that we expect a comparable amount of energy to be released in CCSNe and GRBs. Recent works, see, e.g., Sobacchi et al. (2017), suggest that possibly a large fraction of jets harbored in CCSNe might not be electromagnetically visible. In this work, we take the CCSN–GRB relationship seriously and investigate a model of astrophysical jets originating from CCSNe.

We build up a model where the same physics applies in jets that produce \( \gamma \)-rays and those that do not when the jet is trapped within the stellar envelope. According to our scenario, one could think of high (HL)- and low-luminosity (LL)- GRBs as sub-classes of one larger ensemble to which choked GRBs also belong (Bromberg et al. 2011a; Nakar 2015). Early literature (Murase et al. 2006; Gupta & Zhang 2007; Liu et al. 2011; Liu & Wang 2013; Murase & Ioka 2013; Tamborra & Ando 2015, 2016) proved that the neutrino production from LL- and choked GRBs can even be larger than the one expected for ordinary GRBs.

We explore a general model where the jet properties scale as a function of the Lorentz boost factor \( \Gamma \) and estimate the total diffuse neutrino background from both bright and choked jets by linking choked GRBs with observed GRBs under the same model. To do this, we perform detailed simulations within the internal-shock GRB model and vary the jet parameters to provide a realistic picture of the diffuse flux from the whole jet population. We include both pp and p\( \gamma \) interactions as well as all relevant cooling processes for protons and intermediate accelerated particles and adopt energy-dependent cross sections for all cooling processes relevant to our purpose.

For the sake of completeness, we distinguish among two GRB models. The first one is a “simple” GRB model that is based on the commonly used scaling law \( \theta_\gamma = 1/\Gamma \) applied to the whole jet, with \( \theta_\gamma \) being the half opening angle; in this model, \( \Gamma \) varies across the GRB population. The second one is an “advanced” GRB model that contains a \( \Gamma \)-dependent GRB
population distribution as does the simple model but also contains a distribution of $\Gamma$ within the jet. Finally, we use IceCube data to define upper limits on the jet energy and the fraction of CCSNe that form jets.

The outline of the paper is as follows. In Section 2, we review standard jet physics and explain how each jet parameter is related to the others. We then calculate the review standard jet physics and explain how each jet parameter fraction of CCSNe that form jets.

2. High-energy Neutrino Production in Astrophysical Jets

In this section, we will overview the properties of the astrophysical jets and discuss the relevant cooling processes affecting the protons and secondary particles. We will then convert those jet properties into the diffuse neutrino intensity observed at the Earth. For simplicity, we will now rely on the simple GRB model, wherein each jet is described by a single value of $\Gamma$, until otherwise specified.

2.1. Properties of the Astrophysical Jet

We parameterize the astrophysical jet by the amount of kinetic energy in the jet $E_j$, bulk Lorentz factor $\Gamma$, and electron (magnetic) energy fraction $f_e$ ($g_B$). There are various internal properties that are functions of jet parameters. We enumerate them here for reference.

First, we take the standard theoretical $\Gamma-\theta_j$ relation from special relativity, $\theta_j=1/\Gamma$ (see, e.g., Mészáros 2006). This form is often used throughout the literature with the argument that given a Lorentz boost factor $\Gamma$, the typical angular scale is $\theta_j=1/\Gamma$. We hereby adopt a modified version of this $\Gamma-\theta_j$ relation to match observed jet angles:

$$\theta_j = \begin{cases} 
\frac{1}{\Gamma} & \Gamma \leq 100 \\
30 \frac{1}{\Gamma} & \Gamma > 100.
\end{cases}$$

(1)

The above relation for a typical HL-GRB with $\Gamma = 300$ gives an opening angle of $\theta_j = 6^\circ$, consistent with observations (Goldstein et al. 2016). The break at $\Gamma = 100$ is taken from Cenko et al. (2011), Ackermann et al. (2011), Dermer et al. (2014), Tamborra & Ando (2015). Measurements of the jet opening angle for LL-GRBs are more uncertain. Nevertheless, for the range $\Gamma \in [3, 100]$ the opening angle varies in the range $\theta_j \in [0^\circ 6, 19^\circ]$ which is consistent with estimations of LL-GRBs jet opening angle reported in the literature (Daigne & Mochkovitch 2007; Liang et al. 2007; Toma et al. 2007; Bromberg et al. 2011a; Zhang et al. 2012; Nakar 2015).

The magnetic field strength is given by,

$$\frac{B'^2}{8\pi} = 4\epsilon_B \frac{E_j'}{V'},$$

(2)

where the jet volume is given by,

$$V' = \Omega_j \tilde{r}_j^2 c \tilde{t}_j \Gamma.$$
and estimate it for a typical Wolf–Rayet star with mass $M_\star = 20 M_\odot$ and radius $R_\star = R_\odot$. The isotropic luminosity is $L_{\text{iso}} = 4\pi E_j / \Omega_j$. As many of the parameters in Equation (9) scale with $\Gamma$, we note that the overall $\Gamma$ and $E_j$ dependence of the jet head radius is $r_\text{h} \propto E_j^{2/5} \Gamma^{-8/5}$. We use the standard notation, $Q_\gamma = Q/10^4$ in cgs units, unless otherwise specified. The condition for a jet to be choked in photons is when $r_\text{h} < R_\star$. If the internal shock radius is larger than the stellar radius then no cocoon can form and the jet is not collimated. In this case, Equation (9) no longer applies as there is no jet head, as the jet is visible (not choked).

2.3. Particle Acceleration and Cooling Processes

All of the charged particles in the jet, protons, pions, kaons, and muons lose energy to various processes. To determine the final spectrum of neutrinos, the cooling process of each particle needs to be determined. In different regimes of energy and $\Gamma$, as well as the other parameters, different cooling processes dominate.

2.3.1. Protons

Protons are accelerated on a timescale given by the magnetic field strength,

$$ t_{p,\text{acc}}' = \frac{E_p'}{B'c}. \quad (10) $$

Protons continue to be accelerated until they lose energy faster than their acceleration timescale. The energy loss mechanisms that they may suffer are listed in the following equations. Protons lose energy to synchrotron losses in magnetic field,

$$ t_{p,\text{sync}}' = \frac{3m_p^4c^3}{4\pi \sigma_T m_e^2 E_j B'^2}. \quad (11) $$

where $\sigma_T$ = $6.65 \times 10^{-25}$ cm$^2$ is the Thomson cross section. Protons are cooled by inverse Compton scattering, which we split into two regimes,

$$ t_{p,\text{IC}}' = \begin{cases} \frac{3m_p^4}{4\pi \sigma_T m_e^2 E_j E_{\text{syn}},} & E_p' f_p' < m_p^2 c^4, \\ \frac{3m_p^4}{4\pi \sigma_T m_e^2 c^2 E_{\gamma},} & E_p' E_\gamma > m_p^2 c^4. \end{cases} \quad (12) $$

The Bethe–Heitler process ($p\gamma \rightarrow pe^+e^-$) has a cooling time of,

$$ t_{p,\text{BH}}' = \frac{E_p' \sqrt{m_p^2 c^4 + 2E_p' E_{\gamma}}}{2n_\gamma \sigma_{\text{BH}} m_e c^3 (E_p' + E_{\gamma})}, \quad (13) $$

where $\sigma_{\text{BH}}$ is,

$$ \sigma_{\text{BH}} = \alpha r_e^2 \left[ \frac{28}{9} \ln \left( \frac{2E_p' E_{\gamma}}{m_p m_e c^4} \right) - \frac{106}{9} \right] \quad (14) $$

with $\alpha$ being the fine structure constant and $r_e$ the classical electron radius. Protons are also cooled via $p\gamma$ and $pp$ interactions. Their cooling times are,

$$ t_{p,p\gamma}' = \frac{E_p'}{\sigma_{p\gamma} m_p E_p' \Delta E_p'}, \quad (15) $$

$$ t_{p,pp}' = \frac{E_p'}{\sigma_{pp} m_p E_p' \Delta E_p'}. \quad (16) $$

with $\Delta E_p' / E_p' = 0.2, 0.8$ for the $p\gamma$, $pp$ cases respectively. Energy-dependent cross sections, $\sigma_{p\gamma}$ and $\sigma_{pp}$, are taken from (Patrignani et al. 2016). Finally, protons lose energy due to adiabatic cooling from the expansion of the jet,

$$ t_{p,\text{ac}}' = \frac{\bar{r}_j}{c\Gamma}. \quad (17) $$

Together, the inverse of the total cooling for protons is the sum of the inverses of each individual cooling time, $t_{p,\text{c}}^- = \sum_i t_i^-$. The cooled proton spectrum is the un-cooled spectrum scaled by an additional factor of $[1 - \exp(-\eta_p)]$, where the number of energy losses is, $\eta_p = t_{p,\text{c}}^- / t_{p,\text{acc}}$.

Figure 1 shows the cooling times for protons from each process for our canonical high- and low-$\Gamma$ bursts as a function of the proton energy for jet energy $E_j = 10^{31}$ erg, energy fractions $\epsilon_p = \epsilon_B = 0.1$, and redshift $z = 1$. The solid lines mark the various cooling processes, while the dashed–dotted line represents the total cooling, and the dotted line is the acceleration time.

2.3.2. Intermediate Particles

Pions and kaons created in $p\gamma$ and $pp$ interactions decay into a muon neutrino and muons. The muons subsequently decay into a muon neutrino and an electron neutrino (and an electron):

$$ \pi \rightarrow \mu + \nu_\mu, \quad K \rightarrow \mu + \nu_\mu, \quad \mu \rightarrow e + \nu_\mu + \nu_e. \quad (18) $$

It is important to include the kaon contribution. In fact, even though the branching ratio to produce kaons is $\sim 30$ times less than that for pion, as their maximum energy is often much higher, they dominate at high energies (Ando & Beacom 2005; Asano & Nagataki 2006).

Each intermediate particle also experiences cooling in a similar fashion to protons. The total cooling time for each of these is the same as for protons after changing $m_\pi \rightarrow m_i, i = \pi, K, \mu$, and there is no contribution from the Bethe–Heitler process or $p\gamma$. In addition, while muons do undergo hadronic cooling, the process is negligible (Bulmahn 2010). The fractional energy loss for hadronic interactions for pions and kaons is $\Delta E_i(\pi, K)/E_i' = 0.8$, the same as for $pp$ interactions.

The final cooled spectra for neutrinos coming from the decay of the intermediates are modified in a similar way to the proton spectrum with a factor of $\eta_\mu^{-1} + \eta_\pi^{-1} + \eta_p^{-1}$, where the number of energy losses due to the intermediates is $\eta_i = t_i'^{\text{acc}} \epsilon_i / E_i' \bar{r}_j$ and $\bar{r}_j$ is the rest frame lifetime of the particle. The proton parameters are calculated at the proton energy that corresponds to the given neutrino energy, related by $\bar{a}_i$. The

$^3$ For low energies ($\sqrt{s} < 10$ GeV, with $s$ being the Mandelstam variable), the PDG data are used, while the In’s parameterization is used for high-energy interactions.
muon term is included only for neutrinos from a muon; for neutrinos directly from the mesons, no muon term is included. The multiplicity factors for these cooling processes are given in Table 1.

Figure 2 shows the cooling times for pions, muons, and kaons for GRB models with the same input parameters as for Figure 1 as a function of the neutrino energy. For the adopted GRB parameters, the synchrotron cooling is the only process that affects the spectrum.

2.4. Input Energy Spectra for Protons and Photons

We assume that the central engine is accelerating protons to high energies (where the maximum energy is determined by the acceleration time and the cooling time, see Figure 1 and Section 2.3.1). For protons, we assume an initial Fermi shock accelerated proton spectrum, \( E_p^{-2} \). This sets the initial power law for all subsequent spectra.

2.5. Neutrino Energy Spectrum

For the neutrino spectrum, we take the proton spectrum \( \propto E_p^{-2} \) and multiply it by \( \tau_{\nu\gamma}, \tau_{pp} \) for \( p\gamma, pp \) interactions respectively. As \( \tau_{\nu\gamma} = \sigma_{\nu\gamma} \langle E_{\nu\gamma} \rangle / \Gamma \) (with \( \nu = p, \gamma \)), the effect of the photon break energy is automatically included by integrating over photon energies weighted by the photon spectrum. Then, the un-normalized un-cooled (unc) neutrino spectrum from initial \( pa \) interaction and intermediate...
where proton and neutrino energies are related by \( a_i \) (see Table 1) depending on which intermediate particle the neutrino comes from: \( E'_p = E'_\nu / a_i \). The photon spectrum is normalized such that \( \int dE'_\gamma \frac{dN_{\gamma}}{dE'_\gamma} = 1 \).
We note that in the simple case where $\sigma_{\nu p}$ is given by a step function at the $\Delta b$ baryon threshold energy, we see that the $p\gamma$ correction to the $E_{\nu}^{-2}$ part of the neutrino spectrum is $\propto E_{\nu}^{\gamma}$ before the first break and then $\propto \log E_{\nu}$ after the first break until the spectrum cools to a softer spectrum after the second break. This is different than the conventionally used correction to $E_{\nu}^{-1}$ which is $\propto E_{\nu}^{0}$, $\propto 1$, $\propto E_{\nu}^{-1}$. Our numerical results using the full $p\gamma$ cross section confirm this behavior shown in Figure 3.

The total neutrino flux, accounting for the various energy loss mechanisms, is:

$$ F_{\nu,i,(a)}(E_{\nu}) = \frac{(1+z)^{3}}{\Omega L d_{L}^{2}} N_{a} f_{p} \left[ 1 - \left(1 - \frac{\Delta E_{\nu}'}{E_{\nu}'} \right)_{\nu} \right] \frac{dN_{i}}{dE_{\nu}'}_{\nu,a}, $$

where $i = \pi, \mu, K, \mu K, a = p, \gamma$, and the neutrino spectrum is normalized to the total jet energy, $\int dE_{\nu}' E_{\nu}' d_{L} = E_{\nu}$. The $[1 - (1 - \Delta E_{\nu}' / E_{\nu}')^{\nu}a]$ term accounts for the energy loss due to multiple $p\nu$ and $pp$ interactions. The remaining term accounts for energy loss from intermediate cooling, and is

$$ f_{p} = \frac{\int dE_{\nu}' E_{\nu}' d_{L} \left[ E_{\nu}' / E_{\nu} \right]_{pp+p\gamma}}{\int dE_{\nu}' E_{\nu}' d_{L} \left[ E_{\nu}' / E_{\nu} \right]_{pp+p\gamma}}, $$

where $pc$ in the denominator refers to the fact that we include only proton cooling and not the cooling of secondaries in that integral. Figure 3 shows the neutrino fluence observed at Earth from one source for $\Gamma = 100$, jet energy $E_{j} = 10^{51}$ erg, energy fractions $\epsilon_{p} = 0.1$, and redshift $z = 1$ for each of the four intermediates. The contribution due to $pp$ interactions is mostly visible at low energies (the flat tail of the spectrum), while the $p\gamma$ component is rising at lower energies and then experiences the various cooling processes described above.

The per-flavor neutrino flux before flavor oscillations is,

$$ F_{\nu,\text{unosc}} = F_{\nu,p} + F_{\nu,\mu} + F_{\nu,K} + F_{\nu,\bar{\nu}_{e}}. $$

Neutrinos oscillate en route to Earth and the distance averaged oscillated flux is (Anchordoqui et al. 2014b):

$$ F_{\nu,\text{osc}} = \frac{1}{4} \sin^{2} 2\theta_{12} F_{\nu,\text{unosc}} $n_{i} + \frac{1}{8} (4 - \sin^{2} 2\theta_{12}) F_{\nu,\text{unosc}}, $$

$$ F_{\nu,\text{osc}} = \left(1 - \frac{1}{2} \sin^{2} 2\theta_{12}\right) F_{\nu,\text{unosc}}. $$

with $\theta_{12} = 33^\circ$ (Esteban et al. 2017). While neutrino oscillations through the stellar envelope do result in matter effects for $E_{\nu} \lesssim 10$ TeV (Mena et al. 2007; Sahu & Zhang 2010; Osorio Oliveros et al. 2013), in our analysis we will focus on neutrino energies larger than $\gtrsim 10$ TeV, as those are better constrained by the IceCube data; hence, neutrino oscillations in the source are neglected in this work.

Figure 4 shows the oscillated muon neutrino fluence from one source for jet energy $E_{j} = 10^{50}$ erg, energy fractions $\epsilon_{p} = 0.1$, and redshift $z = 1$ as a function of the neutrino energy for different values of $\Gamma$. As expected, the $pp$ contribution (the nearly flat part at low energies) is subdominant compared to the $p\gamma$ contribution except at energies where IceCube’s sensitivity to astrophysical neutrinos is low due to large atmospheric backgrounds. Moreover, we have restricted ourselves to the high-energy starting event (HESE) data set, which contains neutrinos with $E_{\nu} \gtrsim 40$ TeV and has very low backgrounds.

3. Neutrino Diffuse Emission in the Simple GRB Model

The first model we will consider is the simple GRB model where every jet has one Lorentz boost factor $\Gamma$ for the entire jet. The population of jets will be sampled from a distribution of $\Gamma$’s that describe the data well.

To calculate the diffuse neutrino intensity from GRBs, we assume that the GRB rate, $R(\Gamma)$, is separable into $R(\Gamma) = R_{0}(\Gamma) \xi(\Gamma)$ with

$$ R(\Gamma) \propto \left[ \left(1 + z \right)^{p_{1}} + \left(1 + z \right)^{p_{2}} + \left(1 + z \right)^{p_{3}} \right]^{1/k}, $$

where $k = -10$, $p_{1} = 3.4$, $p_{2} = -0.3$, $p_{3} = -3.5$ are the fit parameters to the star formation rate from (Yuksel et al. 2008). In fact, we assume that the CCSN rate (and in turn the rate of choked and bright GRBs) follows the star formation rate (Dahlen et al. 2012; Horiuchi et al. 2013). This function $R(\Gamma)$ is composed of three parts with power laws $p_{1}, p_{2},$ and $p_{3}$ with breaks at $z_{1} \approx 1$ and $z_{2} \approx 4$, and is normalized to $R(0) = 1$. We make the ansatz that $\xi(\Gamma)$ follows a power law $\xi(\Gamma) = \beta_{i} \Gamma^{\alpha_{i}}$ (Tamborra & Ando 2016). We then constrain

Note that a linear function does not in general remain non-zero when fit to the given criteria. Hence, we only use the power-law parameterization.
the power law by the measured HL-GRB rate for jets with \( \Gamma > 200 \), and the known CCSN rate for all jets:

\[
R_{\text{SN}}(0)\zeta_{\text{SN}} \left( \frac{\Omega_j}{4\pi} \right) = \int_1^{1000} d\Gamma \xi(\Gamma),
\]

\[
\rho_{0,\text{HL-GRB}} = \int_{200}^{1000} d\Gamma \xi(\Gamma),
\]

where \( R_{\text{SN}}(0) \approx 2 \times 10^5 \text{ Gpc}^{-3} \text{ yr}^{-1} \) (Dahlen et al. 2004; Strolger et al. 2015) is the local CCSN rate, \( \zeta_{\text{SN}} \in (0, 1] \) is the fraction of CCSNe that form jets and is taken to be redshift independent. The rate \( \rho_{0,\text{HL-GRB}} \approx 0.8 \text{ Gpc}^{-3} \text{ yr}^{-1} \) is an optimistic estimation for the observed local HL-GRB rate (Wanderman & Piran 2010). We use the range \( \Gamma \in [200, 1000] \) to define HL-GRB’s motivated by Fermi-LAT (Ackermann et al. 2011). The mean fraction of jets pointing toward the Earth \( \langle \Omega_j \rangle /4\pi \) is

\[
\langle \Omega_j \rangle = \int_1^{1000} d\Gamma \xi(\Gamma) \Omega_j,
\]

as \( \Omega_j \) is a function of \( \Gamma \). For example, for our canonical GRB model with \( \zeta_{\text{SN}} = 0.1 \), we get \( \alpha_\Gamma = -2.6 \) and \( \beta_\Gamma = 6.7 \times 10^5 \text{ Gpc}^{-3} \text{ yr}^{-1} \).

The resultant diffuse neutrino intensity from GRBs is

\[
I_\nu(E_\nu) = \int_1^{1000} d\Gamma \int_{z_{\text{min}}}^{z_{\text{max}}} dz \frac{c^2(1 + z)^{-3}}{H_0 E(z)} R(z, \Gamma) F_\nu(E_\nu),
\]

with \( d_L(z) \) the luminosity distance, \( E(z) = \sqrt{\Omega_M(1 + z)^3 + \Omega_\Lambda} \) (Hogg 1999) computed by taking \( \Omega_M = 0.31, \ Omega_\Lambda = 0.69, \ H_0 = 68 \text{ km/s/Mpc} \) (Ade et al. 2016), and \( [z_{\text{min}}, z_{\text{max}}] = [0, 10] \). We require \( \theta_j < \pi/2 \) in Equation (32); if the jet would be larger than that we set the flux to zero.

The resultant diffuse neutrino intensity is plotted in the top panel of Figure 5 for different values of \( \zeta_{\text{SN}} \) along with the six-year HESE data from IceCube (Aartsen et al. 2015d).

Noticeably, according to \( \tilde{E}_j \) and \( \zeta_{\text{SN}} \), only a small fraction of all CCSNe gives origin to a successful jet, as shown in Figure 6.

4. Neutrino Diffuse Emission in the Advanced GRB Model

In this section, we take a somewhat more realistic model of the physics within a jet and allow \( \Gamma \) to vary across the jet angle from some maximum value \( \Gamma_{\text{max}} \) at the center of the jet (\( \theta = 0 \)) to \( \Gamma = 1 \): the advanced GRB model. In this model, we do not rely on \( \theta_j = 1/\Gamma \) anymore. Because we anticipate that multiple shocks will be accelerating the protons, the resulting \( \Gamma \) distribution is a von-Mises–Fisher distribution,

\[
\Gamma(\theta) = \Gamma_{\text{max}} \exp \left[ \kappa (\cos \theta - 1) \right],
\]

where the concentration \( \kappa \approx 1/\sigma^2 \) for \( \sigma \) small, with \( \sigma \) the usual standard deviation. We take \( \sigma = 1/\sqrt{\Gamma_{\text{max}}} \) motivated by random walks of the accelerated particles within the jet. In this model, we define the volume of the jet (Equation (3)) with...
The component of the GRB rate $R(z, \Gamma)$ introduced in Section 3 depending on $\Gamma$ is assumed to follow a distribution similarly defined as in the previous section but this time this is a function of $\Gamma_{\text{max}}$: $\xi(\Gamma_{\text{max}}) = \beta_\Gamma \Gamma_{\text{max}}^{\alpha_\Gamma}$.

The constraint in Equation (30) to reproduce the observed HL-GRB rate becomes,

$$\rho_{\text{HL-GRB}} = \int_{200}^{1000} d\Gamma_{\text{max}} \int_{(0 < \theta < \theta_{\text{max}})} \frac{d\Omega}{4\pi} \xi(\Gamma_{\text{max}}).$$

The maximum value of $\cos \theta$ is 1 and the minimum value of $\cos \theta$ (the maximum value of $\theta$) is when $\Gamma = 200$, which is $\cos \theta|_{\text{min}} = 1 - \ln(\Gamma_{\text{max}}/200)/\kappa$.

$$\rho_{\text{HL-GRB}} = \frac{\beta_\Gamma}{\alpha_\Gamma} \left[ (200^{\alpha_\Gamma} - 1000^{\alpha_\Gamma}) + \alpha_\Gamma 1000^{\alpha_\Gamma} \ln \frac{1000}{200} \right].$$

Similarly to Equation (29), for the advanced GRB model we have

$$R_{\text{SN}}(0) \zeta_{\text{SN}} = \int_{1}^{1000} d\Gamma_{\text{max}} \xi(\Gamma_{\text{max}})$$

$$= \frac{\beta_\Gamma}{\alpha_\Gamma + 1} (1000^{\alpha_\Gamma+1} - 1),$$

where the beaming angle factor on each side has canceled.

The diffuse neutrino intensity is

$$\mathcal{L}(E_\nu) = \int_{1}^{1000} d\Gamma_{\text{max}} \int_{\cos \theta_{\text{min}}}^{1} d(\cos \theta) \int_{\theta_{\text{max}}}^{\theta_{\text{max}}} d\zeta$$

$$\times \frac{cd^2}{H_0 E(\zeta)} R(\zeta, \Gamma_{\text{max}}) F_\nu(E_\nu).$$
typical GRB jet energy ($\bar{E}_j \sim 3 \times 10^{51}$ erg), we expect that ~70% of the jets are choked.

5. IceCube Constraints on the CCSN–GRB Connection

We then construct a $\chi^2$ test with the IceCube data where we allow for neutrinos from jetted bursts to contribute a subdominant component of the observed astrophysical flux,

$$\chi^2 = \sum_i \left[ \frac{I(E_i) - \bar{I}_C(E_i)}{\bar{I}_C(E_i)} \right]^2 \Theta \left[ I(E_i) - \bar{I}_C(E_i) \right],$$

where the sum is over nine energy bins in the range [40 TeV, 20 PeV], which includes four bins with zero events.

We do not include two of the energy bins in the $\chi^2$ that are likely under-fluctuations: the “dip” and the “Glashow” bins centered at $E_\nu = 570$ TeV and 5.3 PeV respectively. These bins have been investigated in the literature as possible evidence of new physics (Anchordoqui et al. 2014a; Learned & Weiler 2014; DiFranzo & Hooper 2015; Tomar et al. 2015). Under the assumption of no new physics, these bins are almost certainly under-fluctuations and unduly push up the $\chi^2$. Moreover, IceCube has seen a through going track event with deposited energy of 2.6 PeV (Schoenenn & Raedel 2015), which corresponds to a higher neutrino energy of ~5–10 PeV. This would suggest that the zero bins in the data will in fact be filled in with future data. Finally, the “dip” bin has already begun to be filled in since the initial deficit, also suggesting that it is an under-fluctuation from the first few years of data.

We then scan jet energies ($\bar{E}_j$) and the fraction of CCSN that form jets ($c_{\text{SN}}$) and determine the significance of the resulting intensity given the data. Figure 8 shows the 90% contour levels for both the simple and advanced GRB models. We note that the simple GRB model starts becoming increasingly consistent with the IceCube data above $\bar{E}_j \sim 10^{51}$ erg due to an increasing fraction of the jets becoming unsuccessful at accelerating protons. On the other hand, the two GRB models give comparable results for $\bar{E}_j \lesssim 3 \times 10^{51}$ erg. Our findings are discussed in the next section.

6. Discussion

For a typical jet energy $\bar{E}_j \sim 3 \times 10^{51}$ erg, Figures 7 and 8 suggest that less than 1% of all CCSN can harbor jets. Most of those jets are predicted to be choked. This fraction should be compared with existing empirical constraints, based on electromagnetic observations of bright jets, suggesting that most likely a sub-sample of CCSNe could further evolve in jets (Guetta & Della Valle 2007; Grieco 2012; Margutti et al. 2014; Modjaz et al. 2014; Sobacchi et al. 2017).

It has been estimated that, under the assumption that all SN Ib/c harbor a jet, less than 10% of these manages to break out from the stellar envelope and further power a prompt emission visible in $\gamma$'s (Soderberg et al. 2004, 2006; Sobacchi et al. 2017). Similarly, Grieco (2012) finds that the ratio of GRB to type Ib/c SNe is about 0.1%–1% in the local universe where type Ib/c SNe comprise up to ~10% of all SNe. The smaller sub-class of broadlined SNe has been linked to GRBs even in the absence of observed gamma-ray emission (Podsiadlowski et al. 2004; Mazzali et al. 2005; Soderberg 2006; Soderberg et al. 2006; Milisavljevic et al. 2015; Milisavljevic & Fesen 2017). Interestingly, our findings are compatible with the above observational constraints and suggest that the large majority of jets harbored in CCSN should be not electromagnetically visible. Additionally, Fesen & Milisavljevic (2016) finds evidence that Cassiopeia A, categorized as a type IIb, may have had a weak jet, which could have either been unsuccessful at accelerating particles or choked and further suggests that all jets may be a part of a single continuous distribution.

Given the uncertainties on the CCSN explosion mechanism and concerning the conditions leading to the jet formation, we refrain from establishing a firm connection to a specific CCSN sub-class in our model and instead rely on the whole CCSN population. Nevertheless, our model, being very general, does automatically take into account any eventual relativistic supernovae (Soderberg et al. 2006, 2010; Chakraborti et al. 2015) not directly linked to an electromagnetically bright GRB as a sub-class of all stellar explosions possibly harboring jets (Lazzati et al. 2012; Margutti et al. 2014).

It is worth noticing that our estimation does not take into account any dependence of the CCSN–GRB rate as a function of the redshift and progenitor metallicity (Levesque et al. 2010; Grieco 2012; Perley et al. 2013). Moreover, there may be additional metallicity/redshift dependence for the choked jets that does not apply to the electromagnetically bright GRBs.

Our findings depend on the assumption of maximally efficient particle acceleration in the jet. The fraction of the jet energy accelerating particles and further leading to the production of neutrinos and photons is currently unconstrained; this should enter as a constant normalization factor in the estimation of the diffuse neutrino intensity. However, if particle acceleration should not be fully efficient in the jet, then one should expect a correspondingly weaker upper bound on $c_{\text{SN}}$. The scan of different possible values of $\bar{E}_j$ should anyway give an idea of the range of $c_{\text{SN}}$ compatible with the data even in the case of less efficient particle acceleration.

7. Conclusions

Jets harbored in CCSN are promising sources of high-energy neutrinos. By relying on the collapsar model, we...
assume that similar physical processes govern both electromagnetically bright and “choked” GRBs. Our calculations include neutrino production from both $p\gamma$ and $pp$ interactions and account for cooling effects of protons, pions, kaons, and muons.

Three relevant classes of jets are investigated. These classes are based on whether or not a jet successfully accelerates protons and whether or not the jet escapes the stellar envelope. If the jet is optically thick, then it is unsuccessful and no high-energy photons or neutrinos are produced. Successful jets can then be either visible, if photons escape the stellar envelope, or energy photons or neutrinos are produced. Successful jets can be either visible, if photons escape the stellar envelope, or choked if the jet does not escape. In either of these two cases, high-energy neutrinos are produced.

We calculate the neutrino diffuse intensity for two different scaling relations between the opening angle $\theta$ and the Lorentz boost factor $\Gamma$. The “simple” GRB model assumes the classical $\theta = 1/\Gamma$ relation, with $\Gamma$ considered to be the same throughout the jet. On the other hand, the “advanced” GRB model takes a $\Gamma$ distribution throughout the jet. $\Gamma$ is assumed to be highest along the jet axis and decreases down to one on the edges of the jet, similarly to what should occur in a more realistic case. In the advanced GRB case, the characteristic width of the jet is given by $1/\sqrt{\Gamma_{\text{max}}}$ where $\Gamma_{\text{max}}$ is the Lorentz boost factor at $\theta = 0$. Our model is then tuned on the observed rate of high-luminosity GRBs and the CCSN rate. We adopt the flux of high-energy neutrinos measured by IceCube as an upper limit to the possible neutrino flux coming from bright and choked GRBs.

We find that while all of the jets in the simple GRB model are electromagnetically bright, the majority of the jets in the advanced GRB model are choked for jet energies $E_j \lesssim 5 \times 10^{51}$ erg, given the differences in the scaling laws of the two GRB models. This implies that it is crucial to adopt a refined modeling of the GRB microphysics in order to constrain the jet properties. In fact, for both models, the compatibility with the IceCube data (Aartsen et al. 2017b) is similar for $E_j \lesssim 10^{51}$ erg with the advanced model being slightly more constrained. Starting above $E_j \sim 10^{51}$ erg, an increasing number of jets in the simple GRB model is unsuccessful, leading to a smaller diffuse intensity.

Noticeably, our findings suggest that, at most, 1% of all CCSNe can harbor jets. Interestingly, those jets are mostly choked. This fraction is competitive with existing empirical and observational constraints, suggesting that even a smaller fraction can further lead to electromagnetically bright GRBs.

Our study constitutes a step toward a realistic and general modeling of the neutrino production within bright and choked jets. It still relies on several simplifying assumptions; however, e.g., it does not take into account any feature due to the metallicity and progenitor dependence of the CCSN population and only considers acceleration at the internal-shock radius. As a consequence, our bounds should provide results in the correct ballpark but may still suffer changes within a more sophisticated population-dependent modeling.

This work proves that neutrinos could be powerful messengers of burst physics. In light of the increasing IceCube statistics, neutrinos could provide major insights on the CCSN–GRB connection in the near future.

We are grateful to Jochen Greiner, Jens Hjorth, and Hans-Thomas Janka for useful discussions. P.B.D. and I.T. acknowledge support from the Villum Foundation (Project No. 13164), and the Danish National Research Foundation (DNRF91). P.B.D. thanks the Danish National Research Foundation (grant No. 1041811001) for support during the final stages of this project. The work of I.T. has also been supported by the Knud Højgaard Foundation and the Deutsche Forschungsgemeinschaft through Sonderforschungsbereich SFB 1258 “Neutrinos and Dark Matter in Astro- and Particle Physics (NDM).
