Enhanced high-energy neutrino emission from choked gamma-ray bursts due to meson and muon acceleration

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It has been suggested that a potentially large fraction of supernovae could be accompanied by relativistic outflows that stall below the stellar surface. In this letter we point out that internal shocks that are believed to accelerate protons to very high energies in these flows will also accelerate secondary mesons and muons. As a result the neutrino spectrum from meson and muon decay is expected to be much harder compared to previous estimates, extending as a single power law up to $\sim 10^3$ TeV. This greatly improves the detection prospects.

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Based on the observational connection (see Ref. [1] for a review) between gamma-ray bursts (GRBs) and supernovae (SNe), it has been hypothesized that a sizable fraction of all SNe is accompanied by a relativistic outflow similar to those that are believed to lie at the base of observed GRBs [2]. Initially the flow accelerates by radiation pressure and ploughs through the pre-burst stellar material. The flow may however be 'choked' below the stellar surface when the central engine is not active for a sufficiently long time. As electromagnetic radiation that is dissipated by these flows will be absorbed by the stellar material, neutrinos are likely the only particles that could indicate the existence of this phenomenon.

High-energy ($\gtrsim 1 \text{ TeV}$) neutrinos arise predominantly in the decay of charged mesons and muons that are created in the interactions of shock-accelerated protons with target protons or photons. In particular, neutrino emission due to proton acceleration in internal shocks in the relativistic flow has been studied in detail [2, 3, 4, 5, 6]. It has been argued that the fluence of high-energy neutrinos in this scenario is suppressed because the mesons and muons lose a substantial amount of energy before decay due to synchrotron emission and hadronic interactions [4, 6]. This energy loss strongly limits the neutrino detection prospects. However, as we point out here, the mesons and muons are also deflected and scattered by the strong magnetic field on very short timescales. These particles have ample time to cross the shock several times before decay and hence they are subject to shock acceleration. In this work we investigate the observational consequences of the acceleration of neutrino parent particles. For concreteness we restrict ourselves to internal shocks in choked GRB outflows. We expect however that the acceleration mechanism is quite general in GRBs, so that our results may affect estimates of neutrino fluxes in other scenarios as well.

The model — We consider a relativistic outflow with total energy $E = 10^{52} E_{52}$, Lorentz factor $\Gamma_j = 10 \Gamma_{j,1}$

and opening angle $\theta = 0.1\theta_{-1}$ (we use the notation $Q = 10^{x}Q_{x}$ throughout this letter). Here θ parameterizes the combined effect of collimation and relativistic beaming. The isotropic-equivalent burst energy is $E_{\rm iso} = 2E/\theta^2 = 2 \times 10^{54} \text{ erg} \times \theta_{-1}^{-2} E_{52}$. Following Refs. [4, 5, 6], we assume that internal shocks occur in the flow at a radius $r_{\rm int} = 2c\Gamma_j^2\delta t = 6 \times 10^{11} \text{ cm} \times \Gamma_{j,1}^2 \delta t_{-1}$, where $\delta t = 0.1 \delta_{j,-1}$ s is the variability timescale of the central engine. The comoving proton density at the internal shock radius is $n_p' = E_{\rm iso}/4\pi r_{\rm int}^2 \Gamma_j^2 m_p c^3 t =$ $10^{19} \text{ cm}^{-3} \times \Gamma_{j,1}^{-6} \theta_{-1}^{-2} E_{52} t_1^{-1} \delta t_{-1}^{-2}$, where $t = 10t_1$ s is the burst duration (quantities in the comoving frame are denoted with a prime) and m_p is the proton mass. The large proton density gives rise to a very high Thomson optical depth $\tau \gtrsim 10^6$. This implies that synchrotron photons that are emitted by shock-accelerated electrons will thermalize. The number density of photons at the internal shock radius is $n'_{\gamma} = 0.33 (\epsilon_e U'/\hbar c)^{3/4} = 2 \times 10^{23} \text{ cm}^{-3} \times \Gamma_{j,1}^{-9/2} \theta_{-1}^{-3/2} (\epsilon_e E)_{51}^{3/4} t_1^{-3/4} \delta t_{-1}^{-3/2}$, where $U' = E_{\text{iso}}/4\pi r_{\text{int}}^2 \Gamma_j^2 ct$ is the comoving energy density and $\epsilon_e = 0.1 \epsilon_{e,-1}$ denotes the fraction of the total energy in the flow that is transferred to the thermal photon distribution. The magnetic field strength at the internal shock radius is $B' = 2 \times 10^8 \,\mathrm{G} \times \Gamma_{j,1}^{-3} \theta_{-1}^{-1} (\epsilon_B E)_{52}^{1/2} t_1^{-1/2} \delta t_{-1}^{-1}$, where $\epsilon_B = 0.1 \epsilon_{B,-1}$ denotes the ratio of electromagnetic energy to the total energy in the flow.

We adopt the SN rate parameterization presented in Ref. [7]. The SN rate within proper distance d_p can be approximated with $\dot{N}_{\rm SN} = 4 \times 10^2 \ {\rm year}^{-1} \times d_{p,2}^3$, where $d_p = 100 \ d_{p,2}$ Mpc. This estimate is within ~30% for $d_p \lesssim 500$ Mpc. Due to collimation of the relativistic outflow, the rate of observable choked GRBs associated with these SNe is $\dot{N}_{\rm CGRB} = \theta^2 \dot{N}_{\rm SN}/2 = 2 \ {\rm yr}^{-1} \times \xi_{\rm SN} \theta_{-1}^2 d_{p,2}^3$, where $\xi_{\rm SN} \leq 1$ is the fraction of SNe that is endowed with the type of outflows considered in this work.

Proton acceleration — We assume that internal shocks accelerate a fraction of the protons in the flow

to high energies (see also below). We express the energy spectrum of accelerated protons as $dN_p/dE'_p =$ $\xi_p E_{\rm iso}(p-1) \Gamma_i^{-1} E'_p^{-p} (m_p c^2)^{p-2}$, where $\xi_p = 0.01 \xi_{p,-2}$ denotes the fraction of shock-accelerated protons to all nucleons in the flow, and p is the shock-acceleration power-law index. Theoretical studies [8] indicate that $p \simeq 2.3$, but recent observations suggest that p could vary [9]. To keep the discussion general we consider here the range p = 2.0...2.6. The proton acceleration timescale is equal to $t'_{p, \text{acc}} = E'_p/qcB'\Gamma'_s = 6 \times 10^{-11} \text{ s} \times E'_{p,0} {\Gamma'}^2_{s,1} \theta_{-1} (\epsilon_B E)_{51}^{-1/2} \delta t_{-1} t_1^{1/2}$, where $E'_p = 1E'_{p,0} \text{ TeV}$ denotes the proton energy and $\Gamma'_s = 10\Gamma'_{s,1}$ is the Lorentz factor of the shock [8]. We take $\Gamma'_s = \Gamma'_j$ since the variation in Lorentz factors between two subsequent shells of material $\Delta \Gamma_j \sim \Gamma_j$. The maximum proton energy may be limited both by energy losses and by the finite shock lifetime. The lifetime is roughly equal to the dynamical timescale $t'_{\rm dyn} = r_{\rm int}/c\Gamma_j = 2\,{\rm s} \times \Gamma_{j,1}\delta t_{-1}$. The dominant proton energy-loss mechanisms are synchrotron radiation, photopion production and protonproton (pp) collisions. The synchrotron energy-loss timescale is $t'_{p,\,\rm sync} = (6\pi m_p^4 c^3) / (\sigma_{\rm T} m_e^2 {B'}^2 E'_p) = 0.1\,{\rm s} \times$ $E'_{p,0}^{-1}\Gamma_{j,1}^{6}\theta_{-1}^{2}(\epsilon_{B}E)_{51}^{-1}\delta t_{-1}^{2}t_{1}$, where we assume that the protons are relativistic. The energy-loss timescale due to pp collisions is $t'_{p,pp} = 1/(cK_{pp}\sigma_{pp}(1-\xi_p)n'_p) = 10^{-4} \text{s} \times \Gamma^6_{j,1} \theta^2_{-1} E^{-1}_{52} \delta t^2_{-1} t_1$, where we assume that $\xi_p \ll 1$, and we approximate the cross section for pp collisions with $\sigma_{pp} \simeq 5 \times 10^{-26} \text{ cm}^2$ and the fractional energy loss with $K_{pp} \simeq 0.5$. At center-of-mass energies well above the pion creation threshold we approximate the protonphoton $(p\gamma)$ cross section with $\sigma_{p\gamma} \simeq 10^{-28} \text{ cm}^2$, and the fractional energy loss to pion production with $K_{p\gamma\pi} \simeq$ 0.2. In this regime the bulk of the photons participates in photopion production so that we estimate the photopion energy-loss timescale as $t'_{p,p\gamma\pi} \simeq 1/(c\sigma_{p\gamma}K_{p\gamma\pi}n'_{\gamma}) =$ $9 \times 10^{-6} \text{ s} \times \Gamma_{j,1}^{9/2} \theta_{-1}^{3/2} (\epsilon_e E)_{51}^{-3/4} \delta t_{-1}^{3/2} t_1^{3/4}$. Equating the acceleration timescale to the dynamical timescale and the energy-loss timescales we find that, unless extreme values of the parameters are invoked, the maximum proton energy $E'_{p,\max}$ is determined by synchrotron energy loss. In this case

$$E'_{p,\max} = 5 \times 10^4 \,\text{TeV} \times \Gamma_{j,1}^2 \theta_{-1}^{1/2} (\epsilon_B E)_{51}^{-1/4} \delta t_{-1}^{1/2} t_1^{1/4} \,. \,(1)$$

For protons with sufficient energy the optical depth for photopion production $\tau_{p\gamma} \simeq 2\sigma_{p\gamma}n'_{\gamma}r_{\rm int}/\Gamma_j = 2 \times 10^6 \times \Gamma_{j,1}^{-7/2} \theta_{-1}^{-3/2} (\epsilon_e E)_{51}^{3/4} \delta t_{-1}^{-1/2} t_1^{-3/4}$ is always larger than unity (unless Γ_j is very large). On the other hand, the optical depth for pp interactions $\tau_{pp} \simeq \sigma_{pp}n'_pr_{\rm int}/\Gamma_j = 3 \times 10^4 \times \Gamma_{j,1}^{-5} \theta_{-1}^{-2} E_{52} \delta t_{-1}^{-1} t_1^{-1}$ is less than unity when $\Gamma_j \gtrsim 80$. In this case protons with energy below the photopion threshold may traverse the jet relatively unhindered and impact directly on the jet head or the stellar material, which could have interesting observational consequences. Here we assume that protons lose all their energy in the outflow and that 20% of this energy is transferred to secondary mesons.

Meson and muon acceleration — The secondary mesons and muons are created in a strongly magnetized environment. The magnetic field deflects and scatters the particles through electromagnetic interactions on timescales comparable to the gyration timescale $t'_{x, \text{gyr}} =$ $\begin{aligned} \epsilon'_x/qcB' &= 6 \times 10^{-10} \,\mathrm{s} \times \epsilon'_{x,0} \Gamma^3_{j,1} \theta_{-1} (\epsilon_B E)_{51}^{-1/2} \delta t_{-1} t_1^{1/2}, \\ \text{where } x \text{ denotes either } \mu \text{ (muon)}, \pi \text{ (pion) or } K \text{ (kaon)}, \end{aligned}$ and $\epsilon'_x = 1 \epsilon'_{x,0}$ TeV is the particle energy. The comoving decay time is given by $t'_{x, \text{dec}} = \tau_x \epsilon'_x / m_x c^2 =$ $2\times 10^{-2}\,\mathrm{s}\,(2\times 10^{-4}\,\mathrm{s},\,2\times 10^{-5}\,\mathrm{s})\times\epsilon_{x,0}'$ for muons (pions, kaons), where τ_x denotes the proper decay time and m_x denotes the mass. Since the decay time is much longer than the gyration time, the particles have ample time to be deflected and scattered by the magnetic field. This allows them to cross the shock repeatedly, thereby gaining energy in a stochastic way. This is essentially the same mechanism of shock acceleration that applies to protons. A necessary condition for this mechanism to work is that the acceleration timescale $t'_{x,\,\rm acc}=t'_{x,\,\rm gyr}/\Gamma'_{s}\ll t'_{x,\,\rm dec}.$ Notice that the ratio $t'_{x,\,\rm acc}/t'_{x,\,\rm dec}$ is independent of the particle energy, so that there is no intrinsic maximum energy to the acceleration process.

For stable particles the energy spectrum due to shock acceleration can be approximated with a power law with index $p = 1 - \ln(\mathcal{P}_{ret}) / \ln \chi$, where $\mathcal{P}_{ret} \simeq 0.5$ is the return probability (i.e. the probability that a particle completes a full cycle of two shock crossings), and $\chi \simeq 1.6 - 2.0$ is the average fractional energy gain per cycle [8]. For unstable particles the energy spectrum at decay (which determines the energy spectrum of the daughter particles) is, in principle, expected to be softer than the energy spectrum of accelerated stable particles because fewer particles complete many cycles. The effect of particle decay can be accounted for by adopting the return probability $\tilde{\mathcal{P}}_{ret} = \mathcal{P}_{ret} - \mathcal{P}_{dec}$, where \mathcal{P}_{ret} is the return probability for stable particles and $\mathcal{P}_{\rm dec} \simeq t'_{x,{\rm acc}}/t'_{x,{\rm dec}}$ is the probability that a particle decays during one cycle. The acceleration timescale of charged particles in the internal shock environment is $t'_{x, \text{acc}} = \epsilon'_x / qcB'\Gamma'_s = 6 \times 10^{-11} \text{ s} \times \epsilon'_{x,0} \Gamma^2_{j,1} \theta_{-1} (\epsilon_B E)_{51}^{-1/2} \delta t_{-1} t_1^{1/2}$. Since $t'_{\text{dec}} \gg t'_{\text{acc}}$ over a wide range of parameters, $\mathcal{P}_{\text{dec}} \ll \mathcal{P}_{\text{ret}}$ and we expect that the mesons and muons are accelerated to a power law with index $p \simeq 2.3$ that extends to the maximum energy determined by the shock lifetime or by energy losses. We find that, both for mesons and muons, the maximum energy ${\epsilon'_x}^{\max}$ is determined by synchrotron losses. Equating the acceleration timescale to the synchrotron energyloss timescale $t'_{x, \text{ sync}} = t'_{p, \text{ sync}} (m_x/m_p)^4$, we estimate:

$$\begin{aligned} \epsilon_{\mu}^{\prime\,\mathrm{max}} &= 6 \times 10^2 \,\mathrm{TeV} \times \Gamma_{j,1}^2 \theta_{-1}^{1/2} (\epsilon_B E)_{51}^{-1/4} \delta t_{-1}^{1/2} t_1^{1/4} \,; (2a) \\ \epsilon_{\pi}^{\prime\,\mathrm{max}} &= 10^3 \,\mathrm{TeV} \times \Gamma_{j,1}^2 \theta_{-1}^{1/2} (\epsilon_B E)_{51}^{-1/4} \delta t_{-1}^{1/2} t_1^{1/4} \,; (2b) \\ \epsilon_{K}^{\prime\,\mathrm{max}} &= 10^4 \,\mathrm{TeV} \times \Gamma_{j,1}^2 \theta_{-1}^{1/2} (\epsilon_B E)_{51}^{-1/4} \delta t_{-1}^{1/2} t_1^{1/4} \,. (2c) \end{aligned}$$

Neutrino fluence — We denote the average neutrino multiplicity per proton with $\mathcal{M}_{p\nu(x)}$ and the average fraction of the proton energy that is transferred to the neutrino with $\xi_{p\nu(x)}$. Here x labels the intermediate particle (muon, pion, kaon). Following Refs. [5, 6], we assume that a proton transfers 20% of its energy per collision to the secondary mesons. Furthermore we take the average pion (kaon) multiplicity per proton interaction equal to 1 (0.10). The branching ratio of pion (kaon) decay to a muon and a neutrino is virtually unity (0.63) and the neutrino receives ~0.25 (0.50) of the meson energy. A muon transfers ~0.33 of its energy to each of two daughter neutrinos. Hence $\mathcal{M}_{p\nu(\mu)} = 2$, $\mathcal{M}_{p\nu(\pi)} = 1$, $\mathcal{M}_{p\nu(K)} = 0.06$; $\xi_{p\nu(\mu)} = 0.05$, $\xi_{p\nu(\pi)} = 0.05$, $\xi_{p\nu(K)} = 0.1$.

The differential neutrino fluence in the observer frame (all flavours combined; neutrinos and antineutrinos combined) can be expressed as follows:

$$\Phi_{\nu(x)}(\epsilon_{\nu}) = \frac{1}{4\pi d_{p}^{2}} \frac{1}{\Gamma_{j}} \frac{\mathcal{M}_{p\nu(x)}}{\xi_{p\nu(x)}} \frac{dN_{p}}{dE_{p}} \qquad (3)$$

$$= \tilde{\Phi}_{\nu(x)} \Gamma_{j,1}^{p-2} \theta_{-1}^{-2} E_{52} \xi_{p,2} d_{p,2}^{-2} \left(\frac{\epsilon_{\nu}}{1 \text{ TeV}}\right)^{-p} ,$$

where we assume that the redshift $z \ll 1$, and

$$\tilde{\Phi}_{\nu(x)} = 5.2 \times 10^{-4} \,\text{TeV}^{-1} \text{cm}^{-2}$$

$$\times \mathcal{M}_{p\nu(x)} \left(\frac{\xi_{p\nu(x)}}{0.05}\right)^{p-1} (p-1)(4.7 \times 10^{-4})^{p-2} \,.$$
(4)

From eqs. (2), we find that the maximum neutrino energy in the observer frame is:

$$\begin{aligned} \epsilon_{\nu(\mu)}^{\max} &= 2 \times 10^3 \,\text{TeV} \times \Gamma_{j,1}^3 \theta_{-1}^{1/2} (\epsilon_B E_{51})^{-1/4} \delta t_{-1}^{1/2} t_1^{1/4} \,; \\ \epsilon_{\nu(\pi)}^{\max} &= 3 \times 10^3 \,\text{TeV} \times \Gamma_{j,1}^3 \theta_{-1}^{1/2} (\epsilon_B E_{51})^{-1/4} \delta t_{-1}^{1/2} t_1^{1/4} \,; \\ \epsilon_{\nu(K)}^{\max} &= 6 \times 10^4 \,\text{TeV} \times \Gamma_{j,1}^3 \theta_{-1}^{1/2} (\epsilon_B E_{51})^{-1/4} \delta t_{-1}^{1/2} t_1^{1/4} \,, \end{aligned}$$

where we take the neutrinos to be isotropic in the comoving frame.

Detection prospects — Based on preliminary results presented in Ref. [10] we conservatively approximate the effective area of IceCube for muon neutrinos with:

$$A_{\rm eff}(\epsilon_{\nu}) = \begin{cases} 5.0 \times 10^2 \,\mathrm{cm}^2 \,\times \left(\epsilon_{\nu,0}\right)^{1.7} & \left(\epsilon_{\rm th}^0 < \epsilon_{\nu} < \epsilon_{\rm br}\right) \\ 1.5 \times 10^5 \,\mathrm{cm}^2 \,\times \left(\epsilon_{\nu,0}\right)^{0.3} & \left(\epsilon_{\nu} > \epsilon_{\rm br}\right) \end{cases}$$

where $\epsilon_{\rm th} = 0.1$ TeV denotes the detector threshold energy, $\epsilon_{\rm br} = 60$ TeV is a break energy, and $\epsilon_{\nu} = 1 \epsilon_{\nu,0}$ TeV. As the neutrino flavour ratio at the source is roughly $\nu_e : \nu_\mu : \nu_\tau \simeq 1 : 2 : 0$, the expected neutrino flavour ratio at the detector is $\simeq 1 : 1 : 1$ due to neutrino oscillations over very large distances [11]. Hence we approximate the fluence of muon neutrinos at the detector with

$$\Phi_{\nu_{\mu}}(\epsilon_{\nu}) \simeq \frac{1}{3} \left(\Phi_{\nu(\mu)}(\epsilon_{\nu}) + \Phi_{\nu(\pi)}(\epsilon_{\nu}) + \Phi_{\nu(K)}(\epsilon_{\nu}) \right) .$$
 (5)

We estimate the number of muon-neutrino interactions $N_{\nu_{\mu}}$ in IceCube by multiplying the muon-neutrino fluence

with the effective area and integrating over ϵ_{ν} to find:

$$N_{\nu_{\mu}} \simeq \tilde{\Phi}_{\nu_{\mu}} \tilde{N}_{\nu_{\mu}} \Gamma_{j,1}^{p-2} \theta_{-1}^{-2} E_{52} \xi_{p,2} d_{p,2}^{-2} , \qquad (6)$$

where $\tilde{\Phi}_{\nu_{\mu}} = (\tilde{\Phi}_{\nu(\mu)} + \tilde{\Phi}_{\nu(\pi)} + \tilde{\Phi}_{\nu(K)})/3$, and

$$\tilde{N}_{\nu\mu} = \frac{9 \times 10^3 \,\text{TeV}\,\text{cm}^2}{60^{p-2}} \left(\frac{1 - x_{\text{th}}^{2.7-p}}{2.7-p} - \frac{1 - x_{\text{max}}^{1.3-p}}{1.3-p} \right) \,.$$

Here we use the shorthand notation $x_{\rm th} \equiv \epsilon_{\rm th}/\epsilon_{\rm br}$ and $x_{\rm max} \equiv \epsilon_{\nu}^{\rm max}/\epsilon_{\rm br}$, and assume that $\epsilon_{\nu}^{\rm max} > \epsilon_{\rm br}$ and that $1.3 . In the limit that <math>\epsilon_{\nu}^{\rm max} \gg \epsilon_{\rm br}$, we find that $\tilde{N}_{\nu\mu} = 2 \times 10^4 \ (8 \times 10^3, 4 \times 10^3)$ TeV cm² for p = 2.0 (2.3, 2.6). Combining this with eqs. (4) and (6), we predict $N_{\nu\mu} = 14 \ (0.63, 0.038) \Gamma_{j,1}^{p-2} \theta_{-1}^{-2} E_{52} \xi_{p,-2} d_p^{-2}$ muon-neutrino interactions in IceCube for $p = 2.0 \ (2.3, 2.6)$. Hence, for model parameters similar to those adopted in this study, a choked GRB at 100 Mpc could be observed by IceCube provided that the shock-acceleration index p is not too large. For reference values of the other parameters the detection of one neutrino requires $p \lesssim 2.3$.

The diffuse flux (per sterad) due to unresolved choked GRBs can be estimated with

$$\Phi_{\nu_{\mu}}^{\text{diff}}(\epsilon_{\nu}) = \frac{\xi_{\text{SN}}\theta^2}{8\pi} \int_0^\infty dz \left(\frac{dV}{dz}\right) \dot{n}_{\text{SN}} \Phi_{\nu_{\mu}}(\epsilon_{\nu}) \,, \qquad (7)$$

where $\xi_{\rm SN}$ denotes the fraction of SNe that is accompanied by a choked GRB, V is the comoving volume, z is the redshift, and $\dot{n}_{\rm SN}$ in the SN rate per unit volume for which we adopt the parameterization given in Ref. [7] (see also Ref. [5]). We find that:

$$\Phi_{\nu_{\mu}}^{\text{diff}}(\epsilon_{\nu}) = \tilde{\Phi}_{\nu_{\mu}}^{\text{diff}} \Gamma_{j,1}^{p-2} E_{52} \xi_{p,2} \xi_{\text{SN}} \left(\frac{\epsilon_{\nu}}{1 \text{TeV}}\right)^{-p} , \qquad (8)$$

where $\tilde{\Phi}_{\nu_{\mu}}^{\text{diff}} = 9 \times 10^{-6} \,\tilde{\Phi}_{\nu_{\mu}} \,\text{sr}^{-1}$. In figure 1 we have plotted our estimate (8) of the diffuse muon-neutrino flux for three values of p and reference values of the other parameters. Also shown are the 90% confidence level upper limits of the AMANDA-II [12] and IceCube (3 year) [13] experiments, the atmospheric neutrino background with the parameterization used in Ref. [5], and the Waxman-Bahcall bound [14]. As can be seen in the figure, the existing limit from the AMANDA-II experiment is already constraining the parameter space of choked GRBs. For p = 2.0 (2.3), we find that $E_{52}\xi_{p,2}\xi_{\rm SN} \lesssim 10^{-2}$ (1). In this regime the predicted diffuse flux is above the Waxman-Bahcall bound, which applies to optically thin sources. When p = 2.6 the expected diffuse flux is below the detector sensitivity for reference values of the other parameters. IceCube will be able to put more stringent constraints on the model parameters. A more detailed analysis of the detection prospects is beyond the scope of this work. We note however that a visible SN counterpart, which would provide evidence in favour of our model, would strongly reduce the neutrino background.

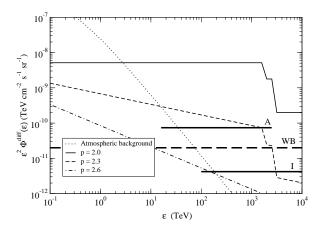


FIG. 1: Diffuse muon-neutrino flux at Earth for three values of the power-law index *p* and reference values of the other parameters. Also shown are the atmospheric neutrino background, the AMANDA-II (A) and IceCube (I) upper limits, and the Waxman-Bahcall bound (WB).

It may also be feasible to use neutrino detectors to initiate a SN search.

Conclusions — In this letter we have found that muons, pions, and kaons created by pp and $p\gamma$ interactions in the internal shocks of choked GRB outflows are also accelerated by these shocks, essentially in the same way as protons and electrons are. Although this acceleration mechanism appears to be unavoidable under the circumstances believed to be present in these flows, it was (to the best of our knowledge) not considered before. Due to meson and muon acceleration the resulting neutrino spectrum is expected to be dominated by neutrinos from muon decay and to follow a single power law with index $p \simeq 2.3$ up to a maximum energy $\sim 10^3$ TeV (and larger than 10^4 TeV for the subdominant contribution of neutrinos from kaon decay). This is in contrast to previous estimates that predict spectral breaks at energies ≤ 1 TeV for neutrinos from meson decay and virtually no high-energy neutrinos from muon decay [4, 5, 6]. The relatively hard neutrino spectrum strongly increases the detection prospects. In fact, the current AMANDA-II limit on the diffuse neutrino background is already mildly constraining the model parameters. The upcoming IceCube neutrino detector will be in a good position to test the model and, possibly, to observe neutrino emission from single choked GRBs. We have estimated that a single choked GRB at 100 Mpc aimed toward Earth will result in 14 (0.63, 0.038) muon-neutrino interactions in IceCube for p = 2.0 (2.3, 2.6) and reference values of the other parameters. The rate of choked GRBs emitting neutrinos toward Earth within 100 Mpc may be as large as a few per year.

An important caveat in our results is that the model relies on the existence of internal shocks that occur at a radius $r_{\rm int}$ due to variability in the flow. It is however not clear whether these shocks can indeed develop while the jet is traversing the star and has not yet created a low-density funnel (Thomas Janka, private communication). In contrast to this, a forward shock and a reverse shock seem unavoidable in the interaction of the relativistic outflow with the stellar environment. Also internal shocks are expected to occur behind the forward shock as it propagates through the star. Neutrino production in these shocks will be studied in a forthcoming publication. We have also assumed that a fair fraction of the secondary mesons makes its way to the shock after being produced. This has to be verified in a more detailed study. We expect that this assumption is best justified at high proton energies, where the proton mean free path is not much larger than the meson gyroradius.

Shock acceleration of mesons and muons may be a fairly general phenomenon in GRBs. The necessary condition that the acceleration timescale is smaller than the decay timescale is fulfilled when $B\Gamma_j > 5$ (6×10^2 , 5×10^3) G for muons (pions, kaons), which is easily achieved in GRBs. Hence our results may also affect the estimates for neutrino emission from successful GRBs.

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