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Charm decay in slow-jet supernovae as the origin of the IceCube ultra-high energy neutrino events

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Abstract. We investigate whether the recent ultra-high energy neutrino events detected at the IceCube neutrino observatory could come from the decay of charmed mesons produced within the mildly relativistic jets of supernova-like astrophysical sources. We demonstrate that the allowed region in the astrophysical and QCD parameter spaces permit an explanation of the 5.7σ excess of neutrinos observed by IceCube in the energy range 30 TeV–2 PeV as a diffuse flux of neutrinos produced in such slow-jet supernovae. We discuss the theoretical uncertainties inherent in the evaluation of charm production in high energy hadronic collisions, as well as some of the astrophysical uncertainties associated with slow-jet supernova sources. These sources result in a diffuse neutrino spectrum that exhibits a sharp drop at energies above a few PeV. We incorporate the effect of energy dependence in the spectrum-weighted charm production and decay cross sections and show that this has a very significant effect on the shape, magnitude and cutoff energies for the diffuse neutrino flux.

Keywords: neutrino astronomy, core-collapse supernovas, neutrino theory

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Contents

1 Introduction

Recently the IceCube (IC) neutrino telescope at the South Pole has reported the observation of 37 neutrino events in the energy range 30 TeV–2.1 PeV, accumulated over three years of runtime [1–3]. These events are 5.7σ above the atmospheric neutrino background, and present, possibly, the first observation of astrophysical neutrinos. The reconstructed flux-spectrum from these events suggest conformity with an isotropic E^{-2} spectrum up to energies of ~ 2 PeV, with the best-fit per-flavor $\nu + \bar{\nu}$ flux in this energy range being given by

$$E^2\Phi = (0.95 \pm 0.3) \times 10^{-8} \text{ GeV cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}, \tag{1.1}$$

assuming an E^{-2} flux. The overall best-fit uniform power-law flux consistent with the lack of events above 3 PeV is given by a slightly more steeply falling flux [3]:

$$E^2\Phi = 1.5 \times 10^{-8} \, (E/100 \text{ TeV})^{-0.3} \text{ GeV cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}. \tag{1.2}$$

The ANTARES neutrino telescope, meanwhile, reports an upper limit on this flux of $E^2\Phi = 4.8 \times 10^{-8}$ GeV cm^{-2} s^{-1} sr^{-1} at 90% confidence level [4]. The IC measurement is, in principle, consistent with the theoretical expectations for a diffuse neutrino flux from extragalactic sources; however, contrary to these expectations, at energies above 2 PeV, the IC event rate drops, hinting at a steep decline in the incident neutrino flux itself at these multi-PeV energies. One theoretical challenge is to explain the apparent cutoff of the neutrino spectrum. The low number of observed events makes it difficult to conclusively determine the nature of the astrophysical sources responsible for the all-sky diffuse flux of neutrinos leading to these events. Several possible origins have been suggested, both astrophysical sources [5–15] and dark matter interactions [16–19].

The role of slow-jet supernovae (SJS) as a possible source of UHE neutrino fluxes has been previously suggested by Razzaque, Meszaros and Waxman (RMW) in refs. [20, 21], and has been explored in detail, see, e.g., [22, 23]. SJS are core-collapse supernovae (SNe) that have jets, similarly to gamma-ray bursts (GRB), although the jets in SJS have much lower Lorentz factors than the jets in a GRB and do not reach the envelope of the star. The environment is optically thick to photons and charged particles; therefore, the only visible sign of the jets may be the emitted neutrinos.

The neutrino flux produced from pion and kaon decays within these sources lies below the atmospheric neutrino background at TeV energies and beyond [20, 22]. However, it has been shown [24] that the decay of charmed D-mesons (D^0, $\bar{D}^0$, $D^\pm$) produced in pp collisions within these sources may lead to considerably higher neutrino fluxes with a spectrum resembling the shape of the proton flux, including the cutoff in the PeV energy range. The proton energy cutoff is due to proton cooling processes starting to dominate over proton acceleration. Previous estimates were done assuming an energy-independent, proton-spectrum weighted charm production cross section [24, 25]. In this paper, we show that the charm production in pp collisions in the source is significantly affected by the proton spectrum and its cutoff. The charm spectrum then has implications for the spectrum of neutrinos from their decays. The effect of the proton-spectrum weighted moments of the charm production cross section and the D-meson-spectrum weighted moments of their decays modify the previously considered [24, 25] effect of the proton energy cutoff. The main qualitative feature of the PeV scale neutrino energy cutoff is retained.

We demonstrate here how the energy dependence of the spectrum weighted charm production and decay is translated to the diffuse neutrino flux from SJS sources. We can account for the observed IC excess events with SJS sources characterized by luminosity, jet bulk Lorentz factor and other parameters similar to those presented by RMW in ref. [21]. The energy dependent production and decay factors make the prediction less compelling than predictions made with approximate energy independent factors. Nevertheless, uncertainties in the inputs discussed below are large enough to accommodate the current IC results.

Our starting point for evaluating the diffuse flux from SJS sources is the RMW model [21] with astrophysical parameters that satisfy the observational constraints set by IceCube [26]. We consider the dependence of the diffuse flux on some of the uncertainties in the astrophysical parametrization of the source distribution in the universe as well as from theoretical uncertainties in the charm production cross section from pp collisions [27].

Apart from the astrophysical uncertainties, the charm production cross section has theoretical uncertainties due to the choice of the charm quark mass and QCD scales, as well as the parton distribution functions and fragmentation functions. As discussed in [27], the choice of charm mass $m_c = 1.5\,\mathrm{GeV}$ seems to underestimate the total charm production cross section seen at recent experiments, e.g., ALICE [28, 29], LHCb [30] and ATLAS [31]. We use $m_c = 1.27\,\mathrm{GeV}$ [32]. Our QCD scales are guided by a recent comparison of theoretical predictions for the $pp \to c\bar{c}X$ cross sections and experimental measurements in [33]. A similar evaluation of the scale dependent uncertainties in the context of the atmospheric neutrino flux appears in [34]. By considering the uncertainties in the relevant astrophysical parameters and in the charm cross section, we compute the plausible range of variation of the resulting diffuse neutrino flux from SJS sources. For a range of choices of the theoretical parameters, we show that the diffuse neutrino flux coming from such sources could explain the most striking features of the flux that reproduces the IC events, namely

1) the approximate $E^{-2.3}$ behavior at energies 30 TeV–2 PeV consistent with the IC best-fit shown in eq. (1.2), and

2) the drop in the flux at energies beyond 2 PeV.

The IC events further seem to derive from an isotropic flux, with no clustering in neither time nor space. The diffuse flux naturally fulfills this requirement.

2 Neutrinos from charm decay in individual slow-jet source

We consider neutrino production in the "choked jets" of mildly relativistic, massive (more than $28M_\odot$) supernovae with bulk-Lorentz factor $\Gamma_j \sim 5$ and a jet angle of $\theta_j \sim 1/\Gamma_j$ guided by the RMW model of ref. [20]. The number densities of electrons and protons in such sources are given by [21]

$$n'_e = n'_p = \frac{1}{2\pi m_p c^3} \frac{L_j}{(\theta_j r_j \Gamma_j)^2}\,, \tag{2.1}$$

where primed quantities are given in the jet comoving frame. Here L_j is the jet luminosity, $r_j = 2\Gamma_j^2 c t_v$ is the jet radius, and t_v is the jet variability time scale ~ 0.1 s. For our benchmark estimate, we use $L_j = 10^{50}$ erg s^{-1}, a jet burst duration of 10 s, typical of such sources [20, 21]. This gives a jet energy of $E_j = 10^{51}$ erg. We take $\Gamma_j = 5$, a choice that allows an observed neutrino energy cutoff in the PeV range yet is not strongly constrained by observations. For the same quantities, the average photon energy and number densities are given by

$$U'_\gamma = \frac{\epsilon_e L_j}{2\pi(\theta_j r_j \Gamma_j)^2 c}\,, \qquad n'_\gamma = \frac{2\zeta(3)}{\pi^{7/2}}\left(\frac{15 U'_\gamma}{\hbar c}\right)^{3/4} \tag{2.2}$$

respectively, where ϵ_e represents the fraction of energy transferred to photons. For the rest of the work, we assume the standard value for $\epsilon_e = 0.1$ as in [21].

Given the relatively high proton content in such slow jets, proton-proton collisions at high energies can dominate over $p\gamma$ interactions and lead to the production of D-mesons. The decay lengths of D-mesons are much shorter than the corresponding interaction lengths, so they decay almost instantly, producing neutrinos (see, e.g., [24]).

In the limit where meson decay dominates over meson cooling, the neutrino fluence (defined as the total particle flux emitted during a complete burst of duration t_j, i.e., $\mathcal{F}_\nu = \phi_\nu t_j$) expected at Earth due to the decay of meson M, from a source at a luminosity distance d_L, is given by [24]

$$\mathcal{F}_\nu(E) = \sum_{M=D^0,\bar{D}^0,D^\pm} Z_{M\nu}(E) Z_{NM}(E) \mathcal{F}_N(E), \tag{2.3}$$

where the spectrum-weighted moments Z_{ij}, discussed below, encode the production and decay of D-mesons. This expression does not hold when significant cooling of the meson occurs before it decays; for this case we use the full expressions from [24],

$$Z_{NM}(E) \to \frac{L_M^{\rm eff}}{L_M^{\rm dec}(\ell_N^{\rm had} + \ell_N^\gamma)}(Z_{NM}(E)\ell_N^\gamma + Z_{NM}^\gamma(E)\ell_N^{\rm had}) \tag{2.4}$$

where the effective length is given by [24]

$$(L_M^{\rm eff})^{-1} = (L_M^{\rm dec})^{-1} + (L_N^{\rm had})^{-1} + \frac{m_p^4}{m_M^4}[(L_N^{\rm IC})^{-1} + (L_N^B)^{-1}]. \tag{2.5}$$

This includes the decay length of the meson M (dec) and cooling lengths from hadronic interactions (had), inverse compton scattering (IC), and cooling from synchrotron radiation because of the magnetic field in the jet, as applicable. For the later two, the mass factor m_p^4/m_M^4 scales the scattering lengths of the proton. The quantities $\ell_N^{\rm had} = (\sigma_{pp} n'_p)^{-1}$ and $\ell_N^\gamma = (\sigma_{pp} n'_\gamma)^{-1}$ describe interactions of the protons with the ambient protons and photons in the (co-moving) jet. Cooling is important for the kaon contribution from SJS, where first hadronic cooling, then radiative cooling dominate. For charmed meson contributions to the neutrino flux from SJS, eq. (2.3) is applicable because decays dominate below the proton energy cutoff.

The quantity $\mathcal{F}_N$ represents the proton fluence within the source, as described in an Earth-observer frame, and $E'_{\rm min} = m_p c^2$. The shape of the proton spectrum is dependent upon the shock acceleration parameters within the source. As long as the protons take longer to cool (due to synchrotron radiation, inverse Compton scattering with thermal photons and interactions with hadrons or gammas) than to be accelerated to the particular energy, the proton spectrum is a power-law $\propto E'^{-2}$. Since acceleration times increase linearly with proton energies, as energies of the protons reach $\mathcal{O}(1)$ PeV (in the comoving frame), the cooling times fall below the acceleration time, and with the cooling processes now dominating, the corresponding proton flux falls off steeply at higher energies.

The exact energy at which the crossover between proton acceleration time and cooling times occurs depends on the specifics of the conditions inside the source. For the acceleration time, these include the magnetic field in the jet, the fraction of jet kinetic energy converted to magnetic field energy and the diffusion coefficient. The diffusion coefficient characterizes the orientation of the magnetic field relative to the shock and is parametrized by κ, a quantity which is inversely proportional to it (see e.g. [24]). Depending on the specifics of the shock acceleration in the source, κ can vary from about 10 to more than an order of magnitude lower. The proton acceleration time is roughly proportional to κ, and in view of the IceCube result, we take $\kappa = 1$ to extend the proton accelation to the $\sim$ PeV energies. Specifically, for our choice of astrophysical source parameters, the crossover energy is $E'_{\rm max} = 10.2$ PeV for $\Gamma_j = 5$. Depending on the orientation of the magnetic fields, κ could be as large as 10 [35, 36], implying that with the other parameters held fixed, the cutoff energy $E'_{\rm max}$ could be lower by about a factor of 10 from what we use here.

We take the proton fluence to be given by

$$\mathcal{F}_N(E') \propto E'^{-2} f_N(E', E'_{\rm max}), \tag{2.6a}$$

where

$$f_N(E', E'_{\rm max}) = \left[1 + \left(\frac{E'}{E'_{\rm max}}\right)\right] e^{-E'/E'_{\rm max}}\,, \tag{2.6b}$$

describes the energy cutoff behavior [36]. The proton fluence is normalized so that for a source at redshift z, if the jet were not choked,

$$\mathcal{F}_N(E) = \frac{E_j(1+z)}{2\pi\theta_j^2 d_L^2 E^2 \log(E_{\rm max}/E_{\rm min})}\, f_N(E, E_{\rm max})\,, \tag{2.6c}$$

for observed proton energy E.[1] The overall normalization of the diffuse neutrino flux is governed by the astrophysical inputs in eq. (2.6c) and the evolution of the SJS population, as discussed in the next section.

In eq. (2.3) $Z_{M\nu}$ and Z_{NM} account for the energy distribution in the decay of the meson M to neutrinos and its production from NN interactions, respectively. It is most convenient to evaluate these Z-moments in the frame co-moving with the jet. The Z_{NM} are defined as follows:

$$Z_{NM}(E') = \int_0^1 \frac{\lambda_N(E')}{\lambda_N(E'/x_E)} \frac{\mathcal{F}_N(E'/x_E)}{\mathcal{F}_N(E')} \frac{\mathrm{d}n_{N\to M}}{\mathrm{d}x_E} \frac{\mathrm{d}x_E}{x_E}\,, \tag{2.7}$$

where $x_E \equiv E'_M/E'_N$, $\lambda_N(E')$ is the hadronic cooling length for protons (see e.g. [24]), and dn/dx_E is the energy distribution of the meson M produced by $N = p$. We compute the Z_{pD}

[1] The normalization in this work is different from that in [24]. The normalization of eq. (2.6c) is consistent with [20–22].

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using the differential charm cross section calculated with the next-to-leading order K-factor following [37] updated with CTEQ6.6 parton distributions [38] and fragmentation of charm quarks into charmed mesons.[2]

We incorporate the energy dependence of the proton-fluence in Z_{pD} by including its energy cutoff factor f_N, and in addition Z_{pD} depends on energy through the energy dependence of the charm total and differential cross sections. We find that the mean x_E for $pp \to DX$ is around 0.2. Therefore, the most significant effect of incorporating the energy dependence in the Z_{pD} is that the D-meson fluence has a lower cutoff than the proton fluence.

In a similar vein, we also compute the energy-dependent decay moments $Z_{D^\pm\nu}$ and $Z_{D^0\nu}$ $(= Z_{\bar{D}^0\nu})$. This is given by

$$Z_{M\nu}(E') = \int_0^1 \frac{d_M(E')}{d_M(E'/x)} \frac{\mathcal{F}_M(E'/x)}{\mathcal{F}_M(E')} \frac{\mathrm{d}n_{M\to\nu}}{\mathrm{d}x} \frac{\mathrm{d}x}{x} \tag{2.8}$$

where, $\mathcal{F}_M(E')$ represents the meson fluence at the energy E' and $d_M(E')$ is the decay length of the meson M at the energy E'. The meson fluence can be related to the proton fluence in eq. (2.6): $\mathcal{F}_M(E') \propto Z_{pM}(E')d_M(E')\mathcal{F}_N(E')$. Thus, the decay moments can be expressed in terms of the production Z-moments and the proton fluence as

$$Z_{M\nu}(E') = \int_0^1 \frac{\mathrm{d}x}{x} \frac{Z_{pM}(E'/x)\mathcal{F}_N(E'/x)}{Z_{pM}(E')\mathcal{F}_N(E')} \frac{\mathrm{d}n_{M\to\nu}}{\mathrm{d}x}. \tag{2.9}$$

The decay distribution $\mathrm{d}n_{M\to\nu}/\mathrm{d}x$ yields roughly 1/3 of the charm energy going into each final state particle since at the parton level, we have, e.g., $c \to s\mu^+\nu_\mu$. We use the decay distribution parameterizations of [40] for charmed mesons to muons.

We present in figure 1 our results for the Z-moments as a function of energy in the frame co-moving with the jet. These are the moments for $Z_{pD^+} = Z_{pD^-}$ and $Z_{pD^0} = Z_{p\bar{D}^0}$, and the corresponding decay moments. We use $m_c = 1.27\,\mathrm{GeV}$ [32], and we show the uncertainty due to the choice of renormalization scale μ_R and factorization scale m_F. For figure 1, we take $[\mu_R, m_F]$ within the range $[1.71m_c, 4.65m_c]$ to $[1.48m_c, 1.25m_c]$. This range of scales is guided by a comparison of the next-to-leading order QCD evaluation of the $c\bar{c}$ cross section with experimental measurements, as discussed in [33].

The Z-moments in figure 1 fall off at energies lower than the proton fluence cutoff. Previous estimates of the neutrino fluences from the decay of D-mesons in slow-jet sources in the literature [24, 25] were made assuming constant Z_{pD} and $Z_{D\nu}$, which, consequently, were larger in magnitude at high energy. This led to neutrino spectra that closely mirrored the proton spectrum, which is not the case when accounting for the energy dependence in these quantities.

3 Diffuse flux and IceCube

The diffuse flux from SJS is given by (see, e.g., [20])

$$\Phi(E_\nu) = \frac{\xi_{\mathrm{sn}}}{2\Gamma_j^2} \int_0^\infty \frac{\dot{n}_{\mathrm{sn}}(z)d_L^2 c}{(1+z)^2} \mathcal{F}_\nu \left|\frac{\mathrm{d}t}{\mathrm{d}z}\right| \mathrm{d}z\,, \tag{3.1}$$

[2]In [24] we used the charm cross section from the dipole picture calculation of [39], but we have found [34] that this cross section has a slower growth with energy that falls below the recently measured charm cross section at the LHC [28–31].

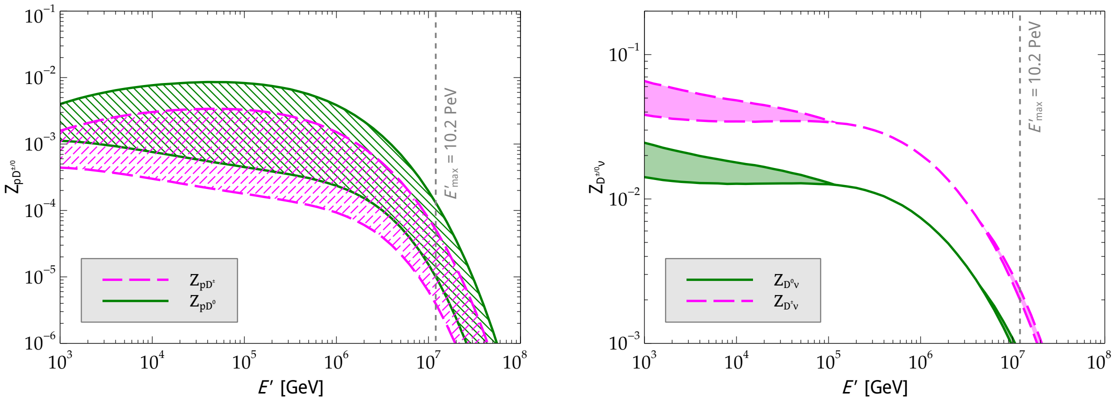

Figure 1. Z-moments for $D^0/D^{\pm}$ production (left) and decay (right) in the jet-comoving frame for $m_c = 1.27\,\mathrm{GeV}$, shown with uncertainties due to variation the renormalization (μ_R) and factorization (μ_F) scales within the range $[1.71m_c, 4.65m_c]$ and $[1.48m_c, 1.25m_c]$. $\Gamma_j = 5$ is used in each case, and the dotted vertical line indicates $E'_{\rm max}$ for the proton flux in the source.

where, $\dot{n}_{\rm sn}(z)$ represents the cosmic supernova-formation rate (SNFR) at a redshift z. At red-shift z, for jet bulk-Lorentz factor Γ_j,

$$\mathcal{F}_\nu = \frac{E_j(1+z)}{2\pi\theta_j^2 d_L^2 E_\nu^2 \log(E'_{\rm max}/E'_{\rm min})} Z_{M\nu}(E')Z_{NM}(E')f_N(E') \tag{3.2}$$

for $E' = (1+z)E_\nu/\Gamma_j$. The red-shift evolution of the universe is given by

$$\frac{\mathrm{d}z}{\mathrm{d}t} = H_0(1+z)\sqrt{\Omega_\Lambda + \Omega_{\rm M}(1+z)^3}. \tag{3.3}$$

Following standard ΛCDM cosmology, the Hubble constant is $H_0 = 68$ km s^{-1} Mpc^{-1}, and $\Omega_{\rm M} = 0.3175$, $\Omega_\Lambda = 0.6825$ [41], and $0 \leqslant \xi_{\rm sn} \leqslant 1$ represents the fraction of supernovae with slow-jets. Further, only a fraction $1/2\Gamma_j^2$ of the total SJS population are directed toward Earth. The SNFR closely follows the cosmic star-formation rate (SFR), $\dot{\rho}(z)$, and is given by $\dot{n}_{\rm sn}(z) = f_{\rm SN}\dot{\rho}(z)$, where $f_{\rm SN}$ (in $M_\odot^{-1}$) represents the fraction of stars converting into supernovae. We use the SFR modeled in [42], with the normalization $\dot{\rho}(0) = 1.3 \times 10^{-4} M_\odot$ yr^{-1} Mpc^{-3} and $f_{\rm SN} = 0.0122\ M_\odot^{-1}$ for the computation of our central flux and corresponding event-rates. The local star formation rate [$\dot{\rho}(0)$] has large uncertainties itself, varying between $(0.6\text{–}2) \times 10^{-2} M_\odot$ yr^{-1} Mpc^{-3} [43–45] depending on the model, while $f_{\rm SN}$ can vary between $(0.8\text{–}1.22) \times 10^{-2} M_\odot^{-1}$ (see, e.g., [21]). These uncertainties in the astrophysical modeling of the sources add to the already existing uncertainties in the diffuse flux which originates from the charm cross section.

To achieve a sufficiently high normalization of the diffuse neutrino flux from SJS, we take the fraction of SNe with jets $\xi_{\rm sn}$= 1. An earlier search by IC and the ROTSE Collaborations [26] for SJS put limits on the bulk-Lorentz factor Γ_j, the jet energy $E_j = L_j t_j$, and $\xi_{\rm sn}$ in the SJS model of [22]. For larger $\Gamma_j \sim 10$, the limit on $\xi_{\rm sn}$ is at the percent level for typical values of $E_j \sim 3\text{–}10 \times 10^{51}$ erg, but for $\Gamma_j \lesssim 6$ these parameters are largely unconstrained.

We show the diffuse neutrino flux produced from decay of kaons and charmed mesons in SJS sources in figure 2. The QCD uncertainties in these fluxes due to the range of renormalization and factorization scales span the region between the long-dashed and solid lines. For our benchmark parameters with $\Gamma_j = 5$, $\xi_{\rm sn}$=1 and $E'_{\rm max} = 10.2$ PeV, using the

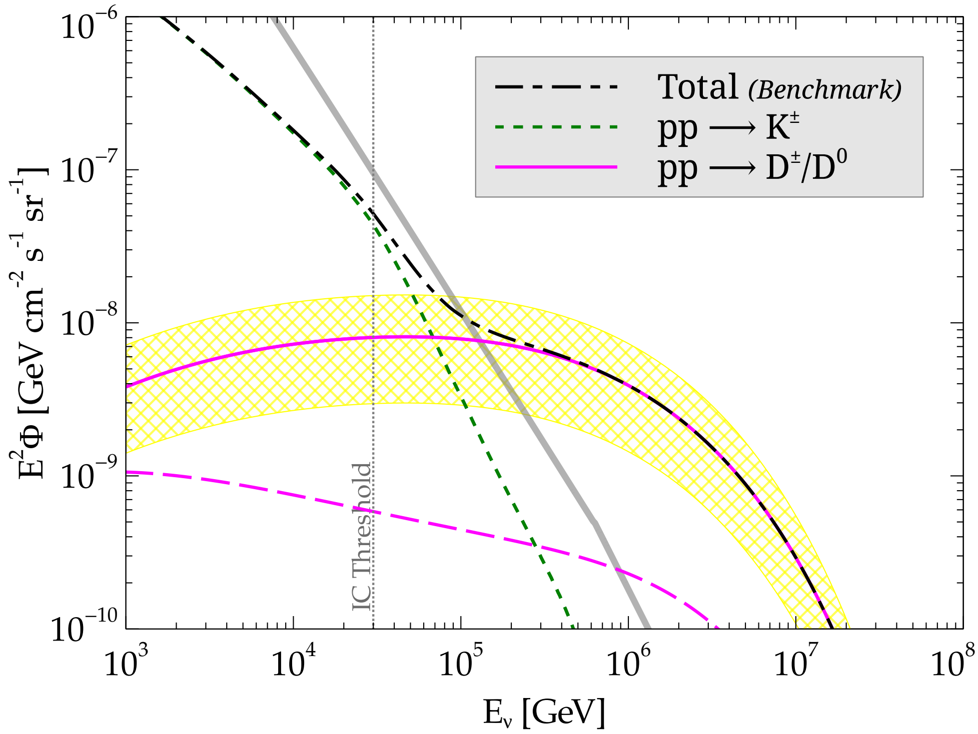

Figure 2. The (unoscillated) diffuse $\nu_\mu + \bar{\nu}_\mu (= \nu_e + \bar{\nu}_e)$ flux obtained with the jet luminosity $L_j = 10^{50}$ erg s^{-1}, $\Gamma_j = 5\,, E'_{\rm max} = 10.2$ PeV and $\xi_{\rm sn} = 1$. The upper solid line and lower long-dashed line show the range of QCD uncertainties from the scale choices in evaluating the charm production cross section. The yellow hatched region shows the variation of the QCD upper limit flux (using $[\mu_R, \mu_F] = [1.71 m_c, 4.65 m_c]$) from uncertainties in the SN formation rate. The short-dashed lines show the kaon contributions to the diffuse neutrino flux from SJS. For comparison, the gray curve shows the vertical flux of conventional atmospheric $\nu_\mu + \bar{\nu}_\mu$ (see, e.g., [46]).

upper curve associated with the QCD uncertainty due to scale dependence, we find that the predicted event-rate from SJS is consistent with the IC observation, dropping off sharply beyond ~ 2 PeV. In contrast, the flux from the decay of kaons is more steeply falling, and only contributes noticeably to the total flux at energies just above the IceCube threshold.

Using the total SJS diffuse flux (kaon + D-meson), we show the estimated 988-day event-rate at IC in figure 3, comparing it to actual observations and against the event rates predicted from the IC best-fit $E^{-2.3}$ flux in eq. (1.2). The number of events expected in IceCube from our predicted diffuse flux is computed by convoluting the diffuse flux with the effective neutrino area given in [2]. Thus, the total number of events (shower + track) expected within the energy range E_0–E_f is given by

$$N_e = 4\pi T \int_{E_0}^{E_f} {\rm d}E \; \Phi_{\nu_\alpha + \bar{\nu}_\alpha}(E) S_\alpha(E)\,, \tag{3.4}$$

where T represents the runtime of the experiment and the IC effective area at the energy E for the neutrino flavor α is given by $S_\alpha(E)$. $\Phi_{\nu_\alpha + \bar{\nu}_\alpha}$ indicates the total $\nu + \bar{\nu}$ flux at earth corresponding to the flavor α. D-mesons decay to produce a predominantly $1:1:0$ (for $\nu_e + \bar{\nu}_e : \nu_\mu + \bar{\nu}_\mu : \nu_\tau + \bar{\nu}_\tau$) flavor composition of neutrinos at source, and we account for neutrino oscillation as they propagate to the Earth while evaluating the predicted event rates at IC. The corresponding neutrino mixing parameters are set at their present best-fit values [47]. The shaded region between the solid and long-dashed histograms shows the QCD scale uncertainty in the SJS diffuse flux prediction. As the figure shows, only the upper end of the QCD prediction gives a reasonable level of high energy IC events. In both figures 2 and 3, we show the variation of the upper end of the QCD prediction due to uncertainties in the SN formation rate.

The natural drop in the event rates shown by the solid histogram beyond 2 PeV is consistent with the lack of events at $E \geqslant 2.1$ PeV at the IC, and in contrast to the ~ 4 events predicted by a uniform E^{-2} flux with normalization set by eq. (1.1). Indeed, for $\Gamma_j = 5$, the predicted event-rate between 3–10 PeV from SJS is only about 1 in 10 yrs. Although, the $E^{-2.3}$ best-fit predicts event-rates consistent with the lack of events in the 988-day IC data, these event-rates are still higher than predictions from the SJS diffuse flux with our parameters. A power-law neutrino flux should be observable over a longer runtime of the experiment. If despite a significantly longer run, the IC nevertheless fails to observe events at these multi-PeV energies, even steeper fluxes like eq. (1.2) will start being disfavored. Indeed, a recent analysis including events above 25 TeV point to an $E^{-2.46}$ power law [48]. Softer neutrino spectra and stronger cutoff features near ~ 2 PeV strengthen the proposal of SJS being important sources of the diffuse neutrino flux.

The normalization of the diffuse neutrino flux from SJS cannot be dramatically larger than what is shown in figure 2. The flux depends on L_j (and E_j) and Γ_j, but so does the cutoff energy. The SJS luminosity used here is in the lower range of such sources, but a significant increase in the jet luminosity is problematic. Photon number densities increase proportionally with luminosities, leading to more effective cooling for the protons in sources with higher L_j. Effective cooling translates to a decrease in $E'_{\rm max}$. Simply scaling up the neutrino energy by increasing Γ_j is not an option because it has also implications for cooling times. Furthermore, larger Γ_j values are more strongly constrained by observations.

Other uncertainties in some of the astrophysical parameters translate the overall magnitude of the diffuse flux, in particular $\xi_{\rm sn}$ and the star formation rate. To achieve the overall normalization shown here, we have taken $\xi_{\rm sn}$=1 and a star formation rate parameterization on the upper end of the 1σ uncertainty. Depending on what the precise values of these parameters are, the diffuse flux from the decay of charmed meson in slow-jet sources might form a smaller fraction of the total diffuse flux seen at IC. Thus, if a significantly lower event-rate than predicted by uniform power-law fluxes $\phi_\nu \propto E^{-\alpha}$ for $\alpha = 2.0$–2.3 persists at these high energies, SJS could be viable candidates for the origin of the incident neutrino flux. However, should future observations reveal significantly higher multi-PeV event-rates, it would disfavor SJS as the main contributor to the diffuse flux. Future observations with higher statistics in the high energy region of the diffuse neutrino flux are needed. Observations of individual point sources may be able to throw light on this as well.

While our focus here has been on the diffuse neutrino flux, SJS are sufficiently bright sources of neutrinos that an explosion at a single nearby source might be visible as a point source over the atmospheric neutrino background. Although the likelihood of a direct nearby SJS burst during the lifetime of IC's run is small [20, 21], a temporally short but visible spike in the muon-track event rates above the atmospheric background could point to such a source. For example, following the procedure outlined in ref. [49], we find that a 10 s burst directed to Earth from a 20 MPc distant source at 0°–30° angle of declination could produce an observable excess at IC.

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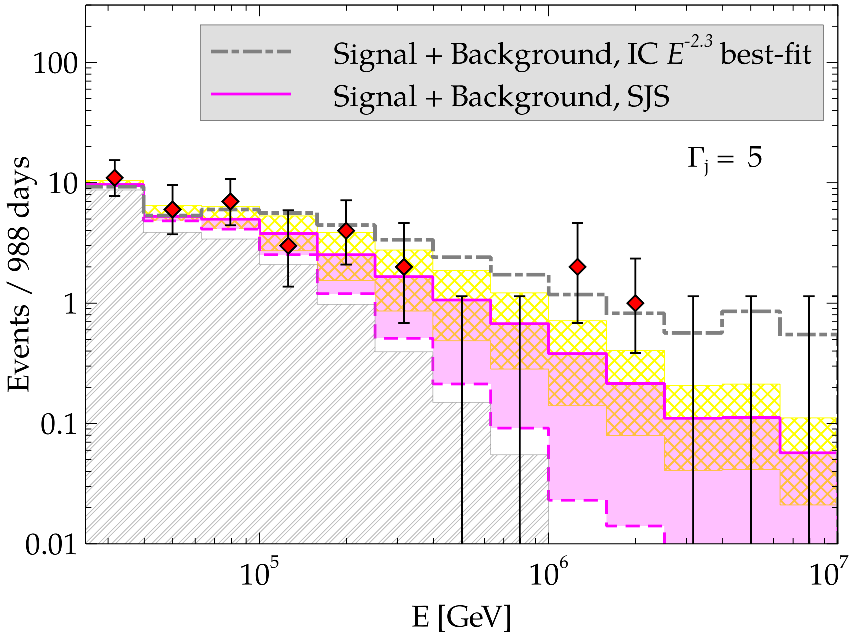

Figure 3. Predicted 988-day total (shower + track) event rates at IC from slow-jet sources for $L_j = 10^{50}$ erg s^{-1}, $\Gamma_j = 5\,, E'_{\rm max} = 10.2$ PeV and $\xi_{\rm sn} = 1$. The solid shaded histogram reflects the QCD scale uncertainties in the charm pair production cross section calculation, with the solid (dashed) histogram showing the upper (lower) range of the SJS diffuse plus atmospheric background number of events. The variation in event-rates relative to the solid histogram from uncertainties in the SN formation rate is shown as a yellow hatched area. Observed event-rates from [3] along with 1σ statistical error bars are shown (red diamonds), as is the total atmospheric neutrino + muon background estimated in the same reference (grey shaded region).

4 Conclusions

We have shown that the recent UHE events seen at IC are consistent with having their origin in a diffuse flux generated by decays of charmed mesons within the mildly relativistic jets of supernovae. To achieve a normalization of $\Phi(E_\nu)$ that approximates eq. (1.2), the upper range of the uncertainty bands must be used.

In contrast to previous work, we have now included the energy dependence in the Z-moments for both the D-meson production and decay. Successively through production and decay, the proton energy cutoff in the jet is translated to lower cutoff energies in the observed neutrino flux. The astrophysical parameters of the SJS sources determine the proton energy cutoff and magnitude of the neutrino flux. We find that for choices of the astrophysical parameters for SJS which have not yet been directly constrained [26] and QCD parameters in the production of charm, the diffuse neutrino flux at Earth from such sources could be enough to explain the observed event rates at energies of 30 TeV to 2 PeV, while also exhibiting a sharp drop in the flux at energies above 2 PeV, in conformity with the lack of events at IC at such high energies. However, QCD uncertainties in the charm production cross section are large. Nevertheless, for a range of parameters, the neutrino flux from D-meson decays within slow jet astrophysical sources could form a significant component of the total diffuse flux seen at IC. If D-meson contributions to the diffuse flux are not included, SJS are not good candidate sources for the observed flux.

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With more IC events, it should be possible to ascertain if the observed neutrinos indeed originate from charmed meson decays because of the distinctive cutoff-like spectral nature of the flux. If, in the future, IC were to find that the diffuse flux were consistent with an unbroken power-law spectrum even at energies beyond 2 PeV and extending into the tens of PeV's, the slow-jet supernovae charmed-meson-origin hypothesis of the incoming neutrino flux would be disfavored. In this latter scenario, incorporating the neutrino flux from charm decay would supplement the IC's capability of constraining the $(\Gamma_j, \xi_{\rm sn})$ parameter space further.

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