TOWARDS AN UNDERSTANDING OF THE SOURCE OF
VERY-HIGH-ENERGY ASTROPHYSICAL NEUTRINOS

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Abstract

Over the past five decades, the frontiers of neutrino astronomy have expanded from the Sun’s interior to the Large Magellanic Cloud with the observation of SN 1987A. Recently, the limits of the observable neutrino sky were expanded to include extra-galactic sources with the discovery of $\sim 10$ TeV-PeV neutrinos at the IceCube Observatory. In this dissertation, I review the unique properties of neutrinos starting from their hypothesized existence in 1930 and including the phenomena of neutrino flavor oscillations, which resulted in the Nobel Prize in Physics for 2015. I go on to present three astrophysical models—supernovae (SNe) and hypernovae in starburst galaxies, choked jet and low-luminosity Gamma-Ray Bursts, and jetted Tidal Disruption Events—in an attempt to explain the source of the observed diffuse flux of neutrinos. Multi-messenger studies that combine electromagnetic (e.g., gamma-ray) and neutrino observations have begun to rule out gamma-ray bright source models such as starburst galaxies. In chapter 5, I present the first coincidence search between a sub-set of gamma-ray dark transients (Type-Ibc core-collapse SNe) and IceCube neutrino events. Because I do not find a connection between core-collapse SNe, I can place limits on the fraction of the explosion energy that can be converted into cosmic-rays—and ultimately non-thermal very-high-energy neutrinos—as well as the fraction of such SNe that may harbor a jet pointed towards Earth. I conclude this work with a review of very-high-energy neutrino astronomy over the last decade and some future work that may help us to understand what the sources of astrophysical neutrinos are.
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1.1 Properties of Neutrinos

Neutrinos are neutral, nearly massless $m_\nu \lesssim 2 \text{ eV}$, particles (See particle data group (PDG) booklet and references within [1]). They are leptons, with three known flavors corresponding to the three charged leptons (the electron, muon, and tau), which they can convert into via charged current weak interactions. Because neutrinos have only been observed to undergo weak interactions (both charged and neutral current) they can be used to probe dense environments that would destroy electromagnetic signals. For example, a beam of neutrinos passing through a one light-year thick slab of lead will only be attenuated by a factor of $\sim 1/2$.

Neutrinos were originally postulated to explain a 1914 observation by Chadwick who showed that electrons produced during $\beta$-decay follow a continuous energy distribution, with energies below the expected value [2]. Note that for a two body interaction, (e.g., a daughter nucleus and a $\beta$ particle), the energy and momentum of both particles are uniquely specified. In 1930 Pauli invented a hypothetical neutral, massless particle to explain the “missing” energy and momentum [3]. These hypothetical particles were included in the standard atomic model in 1934 and given the name neutrino (which translates to “little neutral one” in Italian) by Enrico Fermi [4]. The first experimental detection of neutrinos was announced in 1956 by Reines and Cowan, indirectly detecting neutrinos produced from a fission reactor at the Savannah River Plant in South Carolina [5].

In the Standard Model of particle physics, there are three flavors of neutrinos each of which is massless. The acceptance of the theory of neutrino oscillations indicated that
neutrinos do have (small) mass—which resulted in the Nobel Prize in Physics for 2015 [6]—paved the way for one of the first observations of beyond-the-Standard-Model physics. Recent observations from reactor and neutrino beam experiments have produced some tension in the parameters needed to explain oscillation with only three neutrinos flavors, hinting at the possibility of “sterile” neutrinos, which do not interact weakly [7, 8] (note that cosmological observations have begun to discredit the sterile neutrino hypothesis [9]).

1.1.1 History of Neutrino Detection

Neutrinos rarely interact with conventional matter (i.e., electrons and atomic nuclei). Neutrino detection experiments typically require a massive detector, an intense neutrino source, or both. The first successful neutrino detection by Cowan & Reines utilized two target tanks of a water solution of cadmium chloride which totaled $\sim 400$ L in volume to measure a reactor (anti)neutrino flux of $\Phi_{\nu} \sim 10^{18} \text{cm}^{-2} \text{s}^{-1}$ [5]. They utilized anti $\beta$-decay interactions, where an incoming anti-neutrino interacts with a proton target (in this case from the water nuclei) to create a neutron and positron.

$$\bar{\nu} + p \rightarrow n + e^+.$$ (1.1)

The measurement was of the time delay between $\gamma$-rays produced by the positron annihilation with ambient electrons and the neutron capture on the cadmium.

The first neutrino astronomy experiment was performed by Ray Davis, who was able to detect solar neutrinos [10]. These neutrinos are produced by the fusion processes that power the sun. Davis was able to utilize the relatively large cross section of electron neutrinos interacting with chlorine nuclei to produce a radioactive isotope of argon

$$\nu_e + ^{37}\text{Cl} \rightarrow ^{37}\text{Ar} + e^-,$$ (1.2)

as well as a significantly larger detector ($\sim 4$ kL in volume). It is important to note that the Chorine experiment was only sensitive to electron neutrinos. The dominant form ($\sim 99\%$) of energy production is from the proton-proton reaction $p + p \rightarrow D + e^+ + \nu_e$, which converts two protons into deuterium, a positron, an electron neutrino, and some free energy. However, the neutrinos produced via this mechanism had a maximum energy $\varepsilon_{\nu, ^7}\text{Be} \lesssim 0.5 \text{MeV}$, and therefore were not detectable by a Chlorine experiment, which was sensitive to $\varepsilon_{\nu} \gtrsim 1 \text{MeV}$ (See Fig. 1.1 [11, 12]). Instead, he measured neutrinos which were produced by the decay of $^8\text{B}$ and and $^7\text{Be}$. The former produces a neutrino
spectral line at $\varepsilon_{\nu,7\text{Be}} \sim 0.9 \text{MeV}$ (and an undetectable line at $\sim 0.4 \text{MeV}$), while the latter produces a continuum spectrum with a maximum energy of $\varepsilon_{\nu,8\text{B}} \lesssim 15 \text{MeV}$. Their predicted flux was $\Phi_{\nu,7\text{Be}} \sim 4 \times 10^9 \text{cm}^{-2} \text{s}^{-1}$ and $\Phi_{\nu,8\text{B}} \sim 10^7 \text{cm}^{-2} \text{s}^{-1}$. Fig. 1.2 shows the number of neutrinos observed by the Davis Homestake mine experiment, compared with the predicted number [12]. A deficit of $\sim 1/3$ the expected number of (electron) neutrinos was consistently observed.

This tension between the predicted and detected number of neutrinos went on to fuel the famous solar neutrino problem, which was only resolved by the Super-Kamiokande [13–15] and Sudbury Neutrino Observatory (SNO) collaborations [1,16,17]. They were able to show that the correct total expected number of neutrinos were observed if one summed the contributions of electron, muon, and tau neutrinos. This implied that electron neutrinos in the Sun had been converted to muon and tau neutrinos by the time they had reached the detectors on Earth, in other words, that while neutrinos traveled in vacuum or matter, their flavor could oscillate.

1.1.2 Neutrino Oscillations

This section is based on the Neutrino Properties section of the Particle Data Book [1]

Neutrino oscillations can be described using simple quantum mechanics, if one assumes that there is some mixing between the neutrino flavor states $|\nu_\alpha\rangle$, where $\alpha$ represents the electron, muon, or tau, and the neutrino mass states $|\nu_i\rangle$, where by convention $i = 1, 2, 3$. Mathematically, this mixing is expressed as a mixing matrix such that

$$|\nu_\alpha\rangle = \sum_i U_{\alpha i}^* |\nu_i\rangle$$

(1.3)

$$|\nu_i\rangle = \sum_i U_{\alpha i} |\nu_\alpha\rangle.$$ By convention, the matrix $U$ is a unitary matrix, and its elements are expressed in terms of sines and cosines of mixing angles (e.g., $\theta$). In the two neutrino mixing case (i.e., $i = 1, 2$) the mixing matrix can be expressed as a two-dimentional rotation matrix,
Figure 1.1. Theoretical spectra of solar neutrinos from various nuclear interactions and decays, compared with the sensitivity of various neutrino detection experiments [11,12]. Note that the majority of the neutrino energy from the Sun comes from the $pp$ interaction, but is hard to detect due to a cutoff at $\varepsilon_\nu \sim 0.4$ MeV. The Homestake mine experiment (a Chlorine detector) was able to measure the continuum neutrino flux from $\text{B}^8$ and was marginally sensitive to the neutrino line at $\varepsilon_\nu \sim 0.9$ MeV from the decay of $\text{Be}^7$ nuclei.
Figure 1.2. Comparison between the predicted number of measured solar neutrinos based on the Standard Solar Model (see Fig. 1.1), and the actual number measured by various experiments [12]. The relevant contributions from different nuclear interaction and decay pathways are shown. Note that a deficit of $\sim 1/3$ neutrinos are measured across different types of neutrino detectors that are sensitive to electron neutrinos alone.
\[ U = \begin{pmatrix} \cos \theta & \sin \theta \\ -\sin \theta & \cos \theta \end{pmatrix}. \] (1.4)

To find the probability that a neutrino with initial flavor \( \alpha \) at \( t = 0 \) now has flavor \( \beta \) for \( t > 0 \), one can calculate
\[
P_{\alpha \to \beta} = |\langle \nu_\beta(t) | \nu_\alpha(0) \rangle|^2
\]
with the assumption that the neutrino mass eigenstates can be expressed as a sum of plane wave solutions, i.e.,
\[
|\nu_i(t)\rangle = \exp \left[ -i \left( \epsilon_i t - \vec{p}_i \cdot \vec{x} \right) \right] |\nu_i(0)\rangle
\] (1.5)
where in this case \( i = \sqrt{-1} \) is the imaginary number, and \( \epsilon_i \) and \( \vec{p}_i \) represent the energy and momentum respectively of the neutrino eigenstate \( i \) at time \( t \) located at position \( \vec{x} \). Note that for simplicity \( \hbar = c = 1 \) for this subsection. With the additional assumption that the neutrinos are ultra-relativistic (i.e., \( |\vec{p}_i| \gg m_i \)), the substitution \( t \approx L \) is appropriate, where \( L \) is the distance the neutrino has traveled—recalling that \( c = 1 \). With these simplifications Eq. 1.4 becomes
\[
|\nu_i(t)\rangle = \exp \left[ -i \frac{m_i^2 L}{2\epsilon} \right] |\nu_i(0)\rangle
\] (1.6)
where I have used the approximation \( \epsilon_i = \sqrt{p_i^2 + m_i^2} \approx p_i + \frac{m_i^2}{2p_i} \approx \epsilon + \frac{m_i^2}{2\epsilon^2} \) and \( \epsilon \) is the total energy of the neutrino (i.e., \( \epsilon \approx p_i \) for an ultra-relativistic particle). It can be shown that in the two neutrino mixing scenario, for \( \alpha \neq \beta \)
\[
P_{\alpha \to \beta} = \sin^2(2\theta) \sin^2 \left( \frac{\Delta m^2 L}{4\epsilon} \right)
\] (1.7)
where \( \Delta m^2 = m_1^2 - m_2^2 \) or mass splitting is related to the difference of the square of the two neutrino masses. The probability of finding a neutrino in a particular flavor state depends on two measured parameters (\( \theta \) and \( \Delta m \)) and two parameters that can be tuned—the distance \( L \) and the neutrino energy \( \epsilon \). For neutrino astronomy, obviously the later two are set by the distance the neutrino source is from Earth, and the spectrum of neutrinos it produces.

In general, for larger numbers of neutrino eigenstates, an equation similar to Eq. 1.7 will be found with terms with depend on mixing angles, and terms which depend on the mass splitting(s), distance traveled, and neutrino energy in some combination. For neutrinos which travel through many oscillations (e.g., \( L \gg \frac{\hbar c}{\Delta m^2 c^2} \epsilon \)), the terms with
distance and energy dependence are averaged out through decoherence, resulting in a
probability which depends only on the measured mixing angles. For three neutrino mass
eigenstates

\[ P_{\alpha \rightarrow \beta} = \sum_{j=1}^{3} |U_{\alpha j}|^2 |U_{\beta j}|^2. \]  

(1.8)

For \( \sim 100 \) TeV neutrinos – such as those observed by the IceCube Observatory [18]–
it can be shown using Eq. 1.8 and the measured mixing angles that a population of
neutrinos with flavor ratios \( \nu_e : \nu_\mu : \nu_\tau \sim 1 : 2 : 0 \) at a distance astrophysical source, will
be observed with a flavor ratio of \( \sim 1 : 1 : 1 \) at Earth.

To date, there are known to be three neutrino mass eigenstates and three flavor
eigenstates. Recent measurements have produced some tension between the measured
values of some mixing angles and mass splittings, which may indicate a fourth “sterile”
neutrino mass eigenstate [7,8]. However, sterile neutrinos may be ruled out by \textit{Planck}
measurements of the Cosmic Microwave Background-radiation (CMB) [9].

\section*{1.2 Modern Methods of Astrophysical Neutrino Detection}

As mentioned above, neutrino astronomy started at MeV energies with large tanks of
cleaning solution, and an elemental reaction caused by neutrino interactions [10]. Many
modern beam experiments such as the Deep Underground Neutrino Experiment [19] use
large tanks of similar chemicals as their detectors. However, these modern detectors
allow for real time position, direction, and energetic reconstruction. In this work, we
focus on very-high-energy (VHE), namely GeV to TeV astrophysical neutrinos. The two
modern methods of detection both involve observing secondary electromagnetic radiation
produced by the interaction of a neutrino with a large mass of instrumented “natural”
detector material (e.g., seawater, Antarctic ice, or a large mountain). Because these
detectors are naturally occurring, they can be large-scale \( \gtrsim \) km\(^3\) in size.

\subsection*{1.2.1 Cherenkov Radiation}

Cherenkov radiation is an electromagnetic analog of a sonic shock in fluid dynamics. The
speed of light in most materials is typically slower than that of light in vacuum and is
related to the index of refraction of the material [20]. For example, air typically has an
index of refraction of $n_{\text{air}} = 1.0003$ (CRC Handbook of Chemistry and Physics [21]) which
reduces the speed of light by three parts in $10^4$ (i.e., $c_{\text{air}} = c_{\text{vacuum}}/1.0003$). Note that for
water, $n_{\text{water}} = 1.33$. While such an effect is small when charged particles move through
a given material with a speed $c_{\text{material}} < v_{\text{particle}} < c_{\text{vacuum}}$, they emit a cone of optical
light – referred to as Cherenkov radiation. The best Cherenkov detectors are made of
materials that are transparent to optical light, and will not significantly attenuate or
scatter photons so that the initial Cherenkov signal can be reconstructed. The majority
of Cherenkov based neutrino detectors are made by instrumenting ice (e.g., IceCube [22])
or large bodies of water (e.g., ANTARES in the Mediterranean sea [23]).

Neutrinos can typically produce two types of Cherenkov signals: tracks and cascades.
Figure 1.3 shows an example of the former [24], while Fig. 1.4 is a famous example of
the latter (the so-called “Bert” $\sim$ PeV cascade event) [25]. Track events are produced by
muon neutrinos that undergo a charged current weak interaction with a detector nucleus.
The majority of the energy of the the incoming neutrino is transferred to a muon – with
some energy producing secondary hadronic and electromagnetic (EM) cascades – which,
for sufficiently high energies will move faster than the speed of light in the detector
material.

$$\nu_\mu + N_X \rightarrow \mu + \text{hadrons} + \text{EM cascade}$$

(1.9)

Muons have a relatively long lifetime, and therefore leave a track of Cherenkov emission
that can be meters long in some cases. Such long structures allow for accurate direction
information of the incoming neutrino since the angle between the direction of the neutrino
and muon are small. However, the energy reconstruction of such events can only serve
as a lower limit on the neutrinos total energy if the muon track leaves the detector
volume. Tau neutrinos of sufficiently high energy can also produce track events, often
in a “double-bang” shape, with the first interaction corresponding to the creation of a
charged tau lepton, and the second to the tau’s subsequent decay. No such event has
been observed to date.

Cascade events comprise the remaining types of neutrino/nuclei interactions: electron
weak charged current, low energy tau weak charged current, and all neutral current
interactions. In this scenario, a quasi-spherical cascade of light is seen as the secondary
particles quickly decay or give up their energy to the detector material. Cascade events
provide poor directional reconstruction due to the sphericity of the light emitted. However,
Figure 1.3. An example of a typical HESE neutrino track event observed by IceCube [24]. Note that the earliest detected light comes from a quasi-spherical blob, which shows the location of the initial muon neutrino interaction. Later, a track is formed as the resulting charged muon moves through the instrumented ice.
Figure 1.4. An example of a $\sim$ PeV HESE cascade event. This event – known as Bert – was one of the first astrophysical neutrinos identified by IceCube [25]. Because the high energy electron produced by the neutrino interaction with the ice decays very quickly, there is no associated track, and the resulting light is quasi-spherical.
because the cascade light is typically completely contained inside the detector volume, these events provide better neutrino energy reconstruction. Note, however, that some of the initial neutrino’s energy is dissipated through non-electromagnetic channels that do not emit Cherenkov radiation.

1.2.2 The IceCube Observatory

The IceCube neutrino observatory is located at the geographic South Pole in 2.8 km thick glacial ice [22]. It was constructed between 2005-2010 and consists of 5160 digital optical modules (DOMs) which are frozen into the ice on 86 strings (see the diagram presented in Fig. 1.5 [26]). The DOMs measure the Cherenkov light (Section 1.2.1) produced by the charged particles produced by neutrino interactions (both charged and neutral current weak interactions) with the interior ice. It is currently the largest neutrino detector, with \( \sim 1 \text{ km}^3 \) of instrumented ice.

IceCube measures both track and cascade type events, and is sensitive to neutrinos with energies \( \text{GeV} \lesssim \varepsilon_\nu \lesssim 100 \text{ PeV} \). The lower energy neutrinos are generally of atmospheric origin and are produced by CRs interacting with the Earth’s atmosphere. IceCube is sensitive to neutrinos arriving from all directions, with an energy-dependent effective area (see Fig. 1.6) [27]. Note that the Earth becomes opaque to neutrinos with energies \( \varepsilon_\nu \gtrsim 100 \text{ TeV} \).

Astrophysical neutrinos make up a small sub-set of IceCube’s total sample. There are two general sample populations used to study these astrophysical neutrinos: the high-energy starting events (HESE) for which the entire neutrino interaction is contained within the detector’s effective volume, and the through-going, track-like events. For the later, a track may start inside or outside of the detector, with the resulting muon decaying outside of the effective volume. Through-going events have good directional sensitivity (\( \Delta \theta \lesssim 1^\circ \) [27]), but can only place a lower limit on the original neutrino energy. Furthermore, through-going events are unable to discriminate between muons produced by muon neutrino charged current interactions, or a track produced by a tau particle or a muon which has decayed from a tau particle. Obviously, the later scenario would be produced by a tau neutrino interaction.
Figure 1.5. Schematic of the IceCube Observatory that shows its scale [26]. The detector is located at the geographic South Pole, and is comprised of 5160 digital optical monitors spread between 86 strings. These optical modules are frozen more than a kilometer below the ice’s surface, and make IceCube the largest neutrino detector with $\sim 1\,\text{km}^3$ of effective volume.
Figure 1.6. The energy dependent effective area of IceCube for a range of zenith angles (which are equivalent to the zenith angle above the detector) [27]. Notice that the detector is more effective at low energies for neutrinos arriving from the Northern Hemisphere, while the effective area is larger at high energies for neutrinos from the South. This is because the Earth becomes opaque to neutrinos with energies $E_\nu \gtrsim 100$ TeV.
1.3 Cosmic-Rays

Cosmic-rays (CRs) are charged particles produced in astrophysical environments (e.g., the Sun and both Galactic and extra-galactic sources) [28]. While CRs have been studied for over one hundred years, starting with their discovery by Victor Hess and his famous balloon experiment, the sources of high-energy CRs are not yet known.

CRs make up a population of non-thermal particles and follow a relatively unbroken power-law spectrum over many orders of magnitude in CR energy, (see Fig. 1.7) from semi-relativistic $\varepsilon_{\text{cr}} \sim \text{GeV}$ to the ultra-high-energies with $\varepsilon_{\text{cr}} \sim 10^{21} \text{eV}$. Such a uniform spectrum with index $\alpha_{\text{cr}} \sim 2.7$ suggests that the CRs are accelerated by the Fermi mechanism or diffusive shock acceleration [29], although CR accelerated by magnetic reconnection in relativistic jets has also been suggested [30].

It is generally accepted that CRs below the so-called knee $\varepsilon_{\text{cr}} < \sim 10^{15} \text{eV}$ are Galactic in origin and are produced by young supernova remnants (SNR) [31]. Recent observations of TeV gamma-ray emission from dense molecular clouds near such SNR supports this hypothesis [32], but important questions – such as the composition of CRs around the knee – still remained unresolved. The Hillas approximate expression [33]

$$\varepsilon_{\text{cr, max}} \lesssim q B R$$

(1.10)

where $q$ is the charge of the CR (i.e., an ion with atomic number $Z$ has a charge $q = Ze$), $B$ is the typical turbulent magnetic field strength of a region of size $R$, asserts that CRs with energies $\varepsilon_{\text{cr}} > \sim 10^{17} \text{eV}$ are almost certainly of extra-galactic origin (see Fig. 1.8). For the Milky Way galaxy (MW) with magnetic field strength $B_{\text{MW}} \sim 6 \mu \text{G}$ and scale-height $R_{\text{MW}} \sim 100 \text{pc}$ [34], this implies an absolute upper limit on the CR energy that can be contained inside the MW, $\varepsilon_{\text{cr, gal}} \lesssim Z \times 10^{17} \text{eV}$.

Because CRs are charged, they are deflected by galactic and intergalactic magnetic fields. The latter has not been directly measured but can be inferred to be $B_{\text{ig}} \lesssim n\text{G}$. The deflection of CRs can be as large as $\Delta \theta \sim \text{few} \times 1^\circ$, so that CRs measured at Earth do not point back directly to their source. The higher the CR energy, the less accurate is the direction, with $\Delta \theta \propto R_{\text{cr}}$, where $R_{\text{cr}} = \varepsilon_{\text{cr}}/Z$ is defined as the rigidity. There are indications of a “hot spot” of UHECRs measured by the Telescope Array [35], but the origin of the highest energy CRs is still an open question and is intimately tied to the origin of VHE astrophysical neutrinos [36,37].
Figure 1.7. Measured spectrum of all CRs. Note that the typical energy flux has been multiplied by an additional factor of energy $E^{0.7}$ to accentuate the features inherent in the CR spectrum. For example, the knee at $\varepsilon_{\text{cr}} \sim 10^{6.5}$ GeV and the ankle around $\varepsilon_{\text{cr}} \sim 10^{9.5}$ GeV. Even with these major features, the spectrum is relatively flat, and close to the spectral index $\gamma_{\text{cr}} \sim 2$ expected from diffusive shock acceleration [29].
Figure 1.8. An updated plot that shows the relationship between the size of a CR acceleration region, as well as the strength of the turbulent magnetic field based on the Hillas criteria (see Eq. 1.10). The typical sizes and turbulent magnetic field strength of various astrophysical environments are shown (e.g., GRBs, blazars, starburst galaxy winds) as well as the LHC terrestrial particle accelerator. The dashed red lines correspond to CR energies, such as the energy at which the CR knee is measured (see Fig. 1.7). The area above and to the left of these lines are excluded as possible sites of accretion for CRs with the corresponding energy. Note that the CRs are assumed to be protons.
CR interactions with matter (i.e., $pp$ or proton-proton) and ambient light (i.e., $p\gamma$ or proton-photon) are the only conventional means to produce very-high-energy astrophysical neutrinos. For example, through creation of a resonance, which then decays into lighter hadronic, muonic, and electromagnetic components. The VHE neutrinos are then the decay products of mesons – particularly pions and kaons – and muons [38]. Unbound neutrons can also be produced, and will inverse $\beta$-decay to release an anti-electron neutrino.

\[
\begin{align*}
\pi^+ & \rightarrow \mu^+ + \nu_\mu \\
\mu^+ & \rightarrow e^+ + \nu_e + \bar{\nu}_\mu \\
\pi^- & \rightarrow \mu^- + \bar{\nu}_\mu \\
\mu^- & \rightarrow e^- + \bar{\nu}_e + \nu_\mu \\
n & \rightarrow p + e^- + \bar{\nu}_e \\
\pi^0 & \rightarrow \gamma + \gamma
\end{align*}
\]

(1.11) (1.12) (1.13) (1.14)

Note that if a single charged pion is produced as the result of a $pp$ or $p\gamma$ interaction, the resulting neutrino flavor ratio is $\nu_e : \nu_\mu : \nu_\tau \rightarrow 1 : 2 : 0$, assuming one is unable to discriminate between $\nu$ and $\bar{\nu}$ as is the case for IceCube. If however, a single neutron is produced, the flavor ratio would be $1 : 0 : 0$, with only anti-electron neutrinos produced. Additionally, the same interactions that produce neutrinos through charged meson decay also produce VHE gamma-rays through neutral meson decay, and inverse Compton up-scattering by energetic positrons/electrons. Finally, note that in strong magnetic fields, synchrotron cooling of intermediate particles such as pions or muons may significantly affect the resulting neutrino spectrum by introducing a spectral break [39].

The dominant form of neutrino production measured on Earth comes from $pp$ interactions of CRs with atmospheric nuclei (e.g., Nitrogen). The relationship between the energy of the resulting neutrinos is tied to the energy of the parent CRs in a deterministic way. Furthermore, neutrinos are neutral, weakly-interacting particles, which point back directly to their place of origin. Therefore, the detection of astrophysical neutrinos serves as a smoking-gun signature of VHE CR production and will help to illuminate the sources of the later type of particle.
1.4 Neutrino Source Models

There are two main classes of neutrino source models. Those that are steady neutrino emitters (e.g., CR calorimeters such as starburst galaxies [40]) and transient neutrino sources (e.g., GRBs [36,39], AGN flares [41]). Furthermore, there are galactic and extragalactic sources that fit into both classes of neutrino source. Because CRs are required to produce VHE neutrinos – excluding exotic models such as dark matter annihilation – most theoretical CR sources have also been considered sites of neutrino production. However, this is not true for all models, since the CRs can be produced by one type of object, escape, and then produce neutrinos in a different location. The gamma-ray bright dense molecular clouds near SNRs are an example of such a situation.

Waxman & Bahcall were among to first to identify the connection between sources of CRs and neutrinos, which forms the basis of their famous Bound which carries their names [42,43]. If a fraction $f_{p\gamma} < 1$ of extra-galactic cosmic rays lose energy before leaving their source (or host galaxy/galaxy cluster in the case of CR calorimeters) to $p\gamma$ interactions, then the expected neutrino flux would be $\varepsilon^2_{\nu} \Phi_{\varepsilon_{\nu}} \sim 0.25 f_{p\gamma} \varepsilon^2_{cr} \Phi_{\varepsilon_{cr}}$. The upper limit on the expected neutrino flux is then set by assuming $f_{p\gamma} = 1$. In 1998 they predicted that the measured muon neutrino intensity could not exceed $\varepsilon^2_{\nu} \Phi_{\varepsilon_{\nu}} \sim 10^{-8} \xi z$ GeV cm$^{-2}$ s$^{-1}$ sr$^{-1}$ (1.15) where $\xi z \sim 3$ is a factor that accounts for the cosmological evolution of the CR sources. As we will show below, this prediction is remarkably close to the diffuse muon VHE neutrino flux $\varepsilon^2_{\nu} \Phi_{\varepsilon_{\nu}} \sim 10^{-9}$ GeV cm$^{-2}$ s$^{-1}$ sr$^{-1}$ currently measured by IceCube [18]! This large measured flux, combined with the relatively hard measured spectrum of $\gamma_{\nu} = 2.13 \pm 0.13$ (for the Northern Hemisphere), is an indication that the sources of neutrinos are efficient neutrino factories with $f_{p\gamma} \sim 0.1$ for a range of CR energies so that the resulting spectrum is close to the parent CR spectrum (which is expected to be close to $\gamma_{cr} \sim 2$ from the diffuse shock acceleration model).

The sources of neutrinos must be efficient at converting large amounts of energy into a non-thermal population of VHE CRs. For $p\gamma$ interactions, the average neutrino receives $\sim 1/20$ of the original CR particles energy [42,44]. Therefore neutrinos with energy $\varepsilon_{\nu} \sim 100$ TeV indicate that their sources must be able to accelerate CRs to energies of $\varepsilon_{cr,\text{max}} \sim 10^{16}$ eV. This condition constrains the magnetic field and size of the CR
accretion region in the source via the Hillas condition from Eq. 1.10. Furthermore, because the sources are expected to be efficient neutrino factories ($f_{\nu\gamma} \sim 0.1$ from the previous paragraph), they must also contain dense photon fields for the CRs to interact with. The $p\gamma$ cross section $\sigma_{p\gamma}$ peaks for photons which have an energy $\bar{\epsilon}_\gamma \sim 0.3$ GeV in the rest frame of the CR proton (see Fig. 1.9). In the plasma frame (e.g., of a jet for the case of a GRB), for a CR with energy $\epsilon_{\text{cr}} \sim 10^{16}$ eV, the photon density must be high for photons with energy $\epsilon_\gamma \sim 10$ keV.

In the early 2000’s, analytic models were formulated to estimate the neutrino flux from jetted sources such as GRBs. They simplified the $p\gamma$ cross section to incorporate mainly resonance production (e.g., the $\Delta$-resonance). However, as can be seen from Fig. 1.9, $\sigma_{p\gamma}$ contains significant contributions from direct pion and multi-pion production also at energies above the resonance. Due to their complex energy dependence, these more complicated features were modeled numerically with software such as SOPHIA [45]. Interestingly, it was shown that including the detailed numerical calculation to describe the $p\gamma$ interaction lowered the overall normalization of the resulting neutrino spectrum by an order of magnitude [36]. This result was counterintuitive because more neutrinos were expected from higher-order channels such as multi-pion production. Hümmer at al. 2012 [37] were able to show that many of the assumptions of the original analytic models (e.g., excluding the width of the $\Delta$-resonance or substituting a single energy for the photon spectral energy distribution) which ignored the energy dependent features the $p\gamma$ interaction rate caused errors in the final neutrino flux that were all in the same direction (i.e., it artificially increased the final result). One hopes that errors of these kinds will cancel out so that the analytic result closely approximates the real one. Hümmer et al. went on to provide an improved analytic approximation that captured most of the relevant neutrino production pathways as well as the energy dependence of $p\gamma$ interactions. Despite the revision downward of the expected neutrino flux from GRBs, recent coincident searches by IceCube have indicated that GRBs can contribute at more $\lesssim 1\%$ of the diffuse measured neutrino flux [46].

Steady neutrino source models became popular around 2006. They mixed high turbulent magnetic fields, which could be used to trap VHE CRs for long periods of time with high gas densities to serve as targets for $pp$ interactions. Starburst galaxies (SBG) were seen to be the most favorable of the steady sources. Not only did they provide large magnetic fields and high gas densities, they were also (by definition) sites of large star-forming activity. A higher rate of stellar creation implies a higher rate of SNe and
Figure 1.9. Total $p\gamma$ cross section as a function of the photon energy as seen in the proton’s rest frame. The total cross section (blue solid line) is broken into three components: baryon resonances (red dashed), direct pion production (green dotted), and multi-pion production (brown solid). The theoretical cross section is also compared with cross section data from the Particle Data Group 2008.
their higher energy cousins hypernovae, which made the model attractive since SN-type explosions are still believed to be the dominant method of CR production.

SBGs and other steady neutrino sources were also attractive because they leveraged multi-messenger neutrino astronomy. The same $pp$ and $p\gamma$ interactions that produce neutrinos produce VHE gamma-rays through $\pi^0$ decay and inverse Compton up-scattering of ambient photons from secondary electrons. The VHE gamma-rays provided a second channel to observe neutrino sources since they allowed for joint gamma-ray/neutrino sources. Measurements of the extra-galactic gamma-ray background, as well as gamma-ray observations of individual SBGs, would be able to rule out steady neutrino sources sooner than neutrino observations alone. An analysis by Bechtol et al. 2017 [47] indicated the SBGs and other steady neutrino sources were likely excluded (as the source of neutrinos with energies below 100 TeV) based on measurements make by the Fermi satellite. Such objects could still be the sources of the neutrinos with higher energies ($\varepsilon_\nu \gtrsim 100$ TeV), hinting at a possible two component interpretation of the IceCube observations.

Models of so-called hidden neutrino sources have become more popular since 2014 (e.g., [48–51]). They combine the powerful jets of GRBs, AGN, or Tidal Disruption Events (TDEs) with a cloak of dense matter or photons. Note that VHE gamma-rays can be attenuated by lower energy photons via $\gamma\gamma$ absorption. Such cloaks or envelopes absorb most of the non-thermal radiation produced by a high-energy transient event but do not significantly affect the signal from weakly interacting neutrinos. These hidden models are attractive because they allow for neutrino bright events using the standard theoretical machinery dating back to the late 1990’s (e.g., GRBs) that do not violate stringent constraints set by other messengers such as gamma-rays and CRs.

1.5 Neutrino Astronomy

The era of VHE neutrino astronomy began with the detection of three PeV cascade events in the IceCube detector [52]. Since the initial discovery, IceCube has measured dozens of similar HESE events (both cascade and track) for neutrinos with energy $1\text{ TeV} \lesssim \varepsilon_\nu \lesssim 2\text{ PeV}$ [53]. They have also measured 29 through-going muon neutrinos with energies $200\text{ TeV} \lesssim \varepsilon_{\nu_{\mu}} \lesssim 8\text{ PeV}$. These neutrinos represent an excess above the expected atmospheric neutrino flux with a significance of $> 4\sigma$ [54].

The observed neutrino properties are also consistent with extra-galactic sources. Their arrival directions are consistent with an isotropic distribution, whereas a galactic source
model would have an increased signal from the direction of the Galactic Center. The measured neutrino flavor ratio is consistent with $\nu_e : \nu_\mu : \nu_\tau \sim 1 : 1 : 1$, as is to be expected from distant sources with an initial $1 : 2 : 0$ ratio (see above).

Interestingly, there is tension between the inferred spectral index for samples that include ($\gamma_\nu = 2.49 \pm 0.08$ for $30 \text{ TeV} \lesssim \varepsilon_\nu \lesssim 2 \text{ PeV}$) or exclude ($\gamma_\nu = 2.13 \pm 0.13$ for $\varepsilon_\nu \gtrsim 200 \text{ TeV}$) lower-energy neutrinos. This difference in spectral index may indicate a spectral hardening at higher neutrino energies – which may imply a two component neutrino source model – or unaccounted for contamination from atmospheric neutrinos at low energies. A neutrino spectrum of $\gamma_\nu \sim 2$ is also suggestive of an astrophysical origin since the parent CR spectrum (in or near its acceleration region) would have the same spectral index from diffusive shock acceleration. Atmospheric neutrinos, by contrast, have a relatively soft spectrum with $\gamma_{\nu,\text{atm}} = 2.7$.

### 1.5.1 Multi-Messenger Neutrino Astronomy

As pointed out by Waxman & Bahcall [39,42], the sources of VHE neutrinos must also produce CRs. Furthermore, the same hadronic interactions that produce the charged mesons that ultimately decay into neutrinos also produce neutral pions, which decay into VHE gamma-rays, as well as high-energy electron/positrons that can up-scatter ambient photons. Gravitational waves may also be associated with VHE neutrinos. Collapsars and binary neutron stars may result in GRBs (which may produce neutrinos) and are known to be loud gravitational wave sources. While there have been some binary collaborations (e.g., between IceCube and Swift), there was no real mechanism to combine results from many astrophysical channels from major detectors and experiments. To this end, the Astrophysical Multimessenger Observatory Network (AMON) started at Penn State University has been working to leverage data sets from most of the major high-energy astrophysical observatories.

The AMON project has two main objectives, 1.) to create fast (i.e., real time) alerts between observers so that different detectors can trigger off others. For example, if IceCube observes a neutrino multiplet – two or more neutrinos that arrive from roughly the same direction during a short time interval – alerts can be sent to the gamma-ray observatories such as Fermi, Swift, and the IACTs. 2.) To leverage and combine sub-threshold triggers between detectors to increase their significance to an above-threshold level. For example, if a relatively low energy neutrino is observed by
IceCube in coincidence with a GRB observed by Swift, a special alert might be sent to other observatories.

Multi-messenger astronomy places tighter constraints on neutrino source models. For example, Murase et al. 2013 [44] were able to show that the measured diffuse extragalactic gamma-ray background constrained the spectral index ($\gamma_\nu \approx 2.18$) of the diffuse neutrino flux observed by IceCube in the $pp$ scenario. Furthermore, coincident searches between gamma-ray bright transients and IceCube neutrino events have begun to rule out large classes of models as the dominant source of astrophysical neutrinos [46,47,55]. Because of the recent, stringent constraints from multimessenger studies on neutrino source models, new models of gamma-ray dark sources have begun to proliferate. To this end, new coincident studies which include large field of view optical (ASAS-SN, Zwicky Transient Facility [56]) and X-ray surveys are becoming more important to neutrino astronomy.

1.5.2 Current State of Neutrino Source Determination

The landscape of neutrino source models has changed drastically since the first observation of VHE neutrinos. Many of the top contenders have been ruled out, or are now disfavored as the primary source of astrophysical neutrinos.

Bechtol et al. have used Fermi data to claim that the majority of VHE photons from the diffuse extragalactic gamma-ray background are from unresolved blazars, meaning CR calorimeters such as SBGs, and galaxy clusters/group must be gamma-ray dim. By extension, they cannot produce enough neutrinos to explain the diffuse flux measured by IceCube.

Coincident searches between GRBs and IceCube neutrinos concluded that $\lesssim 1\%$ of the astrophysical neutrinos come from GRBs. Most searches have found no significant association between AGN gamma-ray flares and IceCube neutrinos. Kadler et al. 2016 [57] has claimed a coincident between the famous Big Bird event and an AGN flare, but the significance of the association is relatively low. Gao et al. [58] were also able to show that the gamma-ray and neutrino data could not be consistently explained by a single AGN model. Jetted TDEs were shown by Senno et al. 2017 [51] and others [49,59,59,60] to be energetically disfavored as neutrino sources.

Newer neutrino source models have focused on gamma-ray dim sources, which are not in conflict with current data. These include choked jet or low-luminosity GRBs, as well
as Galactic sources that have been screened such as pulsars. Despite being gamma-ray
dim or dark, these models are typically transients that are bright at other wavelengths
such as the optical or soft X-ray band. Therefore, multi-messenger neutrino astronomy
must expand to include the new types of wide field of view surveys such as ASAS-SN or
the Zwicky Transient Facility.

Below, Chapters 2 - 4 are based on three of my first-authored published papers (Senno
et al. [40,48,51] respectively). Chapter 5 is based on work that has been accepted for
publications with minor revisions [61].
Chapter 2  
Extragalactic Star-Forming Galaxies with Hypernovae and Supernovae as High-Energy Neutrino and Gamma-Ray Sources

2.1 Introduction

The detection of PeV and sub-PeV astrophysical neutrinos by IceCube [24, 25, 52], which more recently has been extended down to 10 TeV energies [53], is a major development. The origin of these neutrinos is a matter of intense interest (for recent reviews, see e.g., [62, 63]). Star-forming galaxies, especially starbursts, are promising candidates [44, 64–68], in which the major contributing sources may be supernovae (SNe), as well as their hyper-energetic equivalent the so-called hypernovae (HNe) [44, 69–73] or active galactic nuclei (AGN) [66, 74]. Other possible candidates include low-luminosity classes of Gamma-Ray Bursts (GRBs) [75, 76], radio-loud AGN [74, 77–80], galaxy clusters and groups [44, 81] with accretion shocks [82, 83] that may accelerate cosmic rays (CRs) to higher energies [33, 84] or other CR sources such as galaxy mergers in clusters [85].

In this work, we will concentrate mainly on the HN/SN origin of neutrinos from pp interactions in the starburst and normal star-forming intragalactic material and the intracluster medium. In particular, we discuss their implications in light of the latest IceCube data in the 10 TeV range. Constraints from the diffuse gamma-ray background measured by Fermi are even more pronounced in this case.
2.2 Hypernova and supernova energy input rate

HNe typically belong to a sub-class of broad-line Type Ibc SNe with ejecta kinetic energies of order \( E_k = 10^{52} E_{k,52} \) erg, representing a fraction \( \xi_{hn} \approx 4 \times 10^{-2} \xi_{hn,-1.4} \) of all core-collapse SNe, with substantial uncertainties [86–88]. The rate of all core-collapse SNe is \( 1.06 \times 10^{-4} \text{Mpc}^{-3} \text{yr}^{-1} \) (e.g., [89]), which implies a local HN rate of \( \mathcal{R}_{hn} \sim 4 \times 10^{-6} \xi_{hn,-1.4} \text{Mpc}^{-3} \text{yr}^{-1} \). If the fraction of HN remnant kinetic energy transferred to CRs is \( E_{cr,\,hn} \approx 2.8 \times 10^{51} \) ergs, the CR energy density input rate in the universe is \( \dot{U}_{cr} \approx 1.2 \times 10^{46} \xi_{hn,-1.4} \text{Mpc}^{-3} \text{yr}^{-1} \). Furthermore, if the CRs are protons with a power law distribution \( N(\varepsilon_p) \propto \varepsilon_p^{-\gamma_p} \) between \( \varepsilon_{p,\,min} \sim 1 \text{GeV} \) and \( \varepsilon_{p,\,max} \sim 10^{17} \) eV, the energy density input rate per logarithmic interval of energy is \( \varepsilon_p Q_{\varepsilon_p} \approx \dot{U}_{cr}/C \text{ erg Mpc}^{-3} \text{yr}^{-1} \), where the bolometric correction is \( C = \ln (\varepsilon_{p,\,max}/\varepsilon_{p,\,min}) \) for a spectral\(^{†}\) index \( \gamma_p = 2 \). Assuming \( \gamma_p \sim 2 \) and taking \( C \sim 18 C_{18} \) the local \((z = 0)\) CR energy input per logarithmic interval in the universe due to HNe is

\[
(\varepsilon_p Q_{\varepsilon_p})_{hn} \approx 6.4 \times 10^{44} \xi_{hn,-1.4} C_{18}^{-1} E_{cr,\,hn,51.4} \text{ erg Mpc}^{-3} \text{yr}^{-1},
\]

which is larger than the typical value expected for GRBs. Conventional SNe will also contribute significantly to lower-energy CRs, having a smaller kinetic energy input \( E_{k,\,sn} \approx 10^{51} E_{k,\,sn,51} \) erg but a larger rate \( \mathcal{R}_{sn} \). The typical CR energy of SNe is uncertain and could be less (e.g. \( E_{cr,\,sn} = 4.8 \times 10^{40} \) erg). In general, the energy injection rate for SNe is given by

\[
(\varepsilon_p Q_{\varepsilon_p})_{sn} = \frac{(1 - \xi_{hn})}{\xi_{hn}} C_{\,hn} \frac{E_{cr,\,sn}}{E_{cr,\,hn}} (\varepsilon_p Q_{\varepsilon_p})_{hn}.
\]

It is believed that SNe can typically accelerate CRs to a maximum energy \( \varepsilon_{p,\,max} \sim 10^{15} \) eV resulting in \( C_{\,sn} \sim 13.8 \). The energy input due to conventional SNe would be typically larger than that of HNe at lower energies. However, with the parameters given above, Eq. (2.2) implies the energy input rate of SNe for CRs with \( \varepsilon_p \lesssim \varepsilon_{p,\,max,\,sn} \) is roughly half that of HNe. Below, we leave this ratio as a free parameter.

\(^{†}\)It would be \( C = [1 - (\varepsilon_{p,\,max}/\varepsilon_{p,\,min})^{2-\gamma_p}]\varepsilon_{p,\,min}^{2-\gamma_p} (\gamma_p - 2)^{-1} \) if \( \gamma_p \neq 2 \).
2.3 Shock acceleration

A typical Type Ibc SN has a bulk ejecta mass of \( M_{\text{ej}} \sim 3M_{\text{ej},0.5}M_{\odot} \), and a HN has an average velocity \( \beta_{\text{ej}} = (V_{\text{ej}}/c) = 6.1 \times 10^{-2}E_{k,\text{hn},52}^{-1/2}M_{\text{ej},0.5}^{-1/2} \). The postshock random magnetic field strength is expected to be amplified to a fraction \( \epsilon_B \) of the postshock thermal energy, \( B_s \sim (16\pi\epsilon_B n_g M N c^2 \beta_{\text{ej}}^2)^{1/2} \), with \( n_g \) being the interstellar particle number density. The upstream magnetic field should also be amplified by, e.g., CR-streaming instabilities, but in any case the stronger magnetic fields in starburst galaxies may also be enough [44]. Diffusive shock acceleration in the blast wave leads to a power law spectrum distribution \( N(\varepsilon) \propto \varepsilon^{-\gamma_p} \), typically with \( \gamma_p \gtrsim 2 \), up to a maximum energy \( \varepsilon_{p,\text{max}} \sim (3/20)Z\epsilon_B R_{\text{dec}}\beta_{\text{ej}}^{-1} \) [29], or

\[
\varepsilon_{p,\text{hn,\text{max}}} \sim 10^{17}Zn_{g,2.3}^{1/6}E_{k,\text{hn},52}^{-2/3}M_{\text{ej},0.5}^{2/3} \text{ eV}
\]

for CRs with charge \( Z \). However, in the following work we only consider CR protons. The above equation implies that the CRs are accelerated to \( \sim 100 \text{ PeV} \) as the shock slows and enters the so-called Sedov-Taylor phase. The maximum CR energy is expected to decrease with time during this deceleration phase [90]. Also note that while many SNe and HNe may happen in relatively low-density regions such as superbubbles, the dependence on \( n_g \) is weak. For normal SNe, using the same parameters except for \( E_k = 10^{51} \text{ erg} \), the maximum energy would be \( \varepsilon_{p,\text{sn,\text{max}}} \sim 1.1 \times 10^{16}Zn_{g,2.3}^{1/6}E_{k,51}^{-2/3}M_{\text{ej},0.5}^{2/3} \text{ eV} \). As we will show below, HNe that occur in starburst galaxies can accelerate the majority of CRs which produce detectable high-energy neutrinos. Although we use typical numbers for HNe for my estimates, in star-forming galaxies hosting an AGN the latter may be also contribute as a 10-100 PeV CR accelerator [66,74].

CRs suffer energy losses both during acceleration and after escaping their source. Synchrotron losses are negligible at the energies considered here, the dominant loss mechanism being hadronuclear \( (pp) \) collisions. The effective optical depth to \( pp \) collisions undergone while advected downstream of the blast wave is \( \tau_{pp,s} \sim t_{\text{dyn}}/t_{pp} \sim \kappa \sigma_{pp} R(c/V) \), where \( \kappa \sim 0.5 \) is the inelasticity and \( \sigma_{pp}(\varepsilon_p = 100 \text{ PeV}) \sim 10^{-25} \text{ cm}^2 \) (in the numerical calculations presented below, we use the energy dependent inelastic \( pp \) cross section presented in [91]). Thus \( \tau_{pp,s} \sim 1.3 \times 10^{-6}E_{k,\text{hn},52}^{-1/2}M_{\text{ej},0.5}^{-5/6}n_{g,2.3}^{-1/3} \), which is negligible compared to losses during the subsequent propagation. Similar considerations apply also to the SNe.
2.4 Propagation effects and $pp$ optical depth

The propagation of the CRs in the turbulent magnetic field of the host galaxy and galaxy cluster depends, in the diffusion approximation, on the strength of the magnetic field $B$, the CR Larmor radius $r_L$, and the coherence length $\ell_c$ of the magnetic field fluctuations. At the highest energies $\varepsilon_p$, where $r_L(\varepsilon_p) \gg \ell_c$, the CR diffusion coefficient is $D(\varepsilon_p) \propto r_L(\varepsilon_p)^2$. At lower energies, where $r_L(\varepsilon_p) \ll \ell_c$, the diffusion coefficient is $D(\varepsilon_p) \propto r_L(\varepsilon_p)^{\alpha}$, where $\alpha = 1/3$ (1/2) for a Komolgoroff (Kraichnan) fluctuation power spectrum (e.g., [92]). The two regimes can be interpolated as

$$D(\varepsilon_p) = D_\ast \left[ (\varepsilon_p/\varepsilon_{p,\ast})^\alpha + (\varepsilon_p/\varepsilon_{p,\ast})^2 \right]$$

where $r_L(\varepsilon_{p,\ast}) = \ell_c/5$ with $D_\ast \simeq (1/4)cr_L(\varepsilon_{p,\ast})$ [93]. Below we shall use $\alpha = 1/3$ as an example, but the discussion can be generalized to a general positive $\alpha$ value.

After leaving the source (e.g., HNe or SNe), the CRs first propagate diffusively through the host galaxy or are advected away by a strong galactic wind, with typical velocities of $V_w \sim 1500$ km s$^{-1}$ in starburst galaxies [94] and $V_w \sim 500$ km s$^{-1}$ for normal star-forming galaxies [34,95]. For a starburst galaxy (SBG) the gas scale height $H_g \sim 30 – 300$ pc may be parameterized as $H_{sbg} \sim 300 \text{pc} \simeq 10^{21} H_21$ cm. We assume a magnetic field strength of $B_g \sim 200 \times 10^{-6} B_{g,-3.7}$ G and a coherence length parameterized here as $\ell_{c,g} \sim 10^{-1} H_g \sim 30 \text{pc} \simeq 10^{20} \ell_{g,20}$ cm. Both quantities are subject to large uncertainties and variations, so that the diffusion coefficient adopted here corresponds to the optimistic case. For my fiducial starburst galaxy parameters, we obtain $\varepsilon_{p,\ast,g} \sim 1.11 \times 10^{18} \ell_{g,20} B_{g,-3.7}$ eV and $D_{s,g} \sim 1.4 \times 10^{29} \ell_{g,20}$ cm$^2$ s$^{-1}$. To ensure CR confinement, we require the coherence length to satisfy $\varepsilon_{p,\ast,g} \simeq 10 – 100$ PeV [44]. For a normal star-forming galaxy (SFG), we take the typical scale height to be $H_g \sim 1000$ pc, with $\ell_{c,g} \sim 10^{-1} H_g$, a magnetic field of $B_g \sim 6 \mu$G [96] and interstellar medium density of $n_g \sim 1$ cm$^{-3}$.

Given the above, the time for CR diffusive escape from the galaxy can be calculated, which for starburst galaxies is $t_{d,g} = H_g^2/6D_g \simeq 1.5 \times 10^{12} H_{g,21}^2 \ell_{g,20} B_{g,-3.7}^{-3} \varepsilon_{p,17.2}^{-1/3}$ s. The time for advective escape is $t_{w,g} = H_g/V_w \simeq 6.2 \times 10^{12} H_{g,21} V_{w,3.2}^{-1}$ s regardless of the CR energy. Notice that advective escape dominates diffusive escape from the galaxy for CRs
with energy \( \varepsilon_p \lesssim \varepsilon_w \) with
\[
\varepsilon_w = \frac{Z e B_1^{1-1/\alpha}}{5} \left( \frac{10 V_w H_g}{3 c} \right)^{1/\alpha},
\]
yielding \( \varepsilon_w \sim 5.1 \times 10^{15} \text{ eV} \) for the fiducial parameters used here.

The effective \( pp \) optical depth undergone during propagation in a SBG is \( \tau_{pp,g} \sim n_g \kappa \sigma_{pp} c \min[t_{d,g}, t_{w,g}] \) or
\[
\begin{align*}
\tau_{pp,g} &\sim 4.9 \times 10^{-3} n_{g,2.3} H_{g,21}^2 \ell_{g,20} B_{g,20}^2 \varepsilon_{-2}^{-2} \\
\tau_{pp,g} &\sim 0.55 n_{g,2.3} H_{g,21}^2 \ell_{g,20} B_{g,20}^2 \varepsilon_{-3}^{-1/3} \\
\tau_{pp,g} &\sim 1 n_{g,2.3} H_{g,21} V_{w,3.2}
\end{align*}
\]
in the ranges \( (\varepsilon_p > \varepsilon_{p,g}) \), \( (\varepsilon_w < \varepsilon_p < \varepsilon_{p,g}) \) and \( (\varepsilon_p < \varepsilon_w) \), respectively.

One can see that starburst galaxies are efficient neutrino factories via \( pp \) interactions due to their high interstellar gas density. As seen in Figures 2.1-2.3, normal star-forming galaxies have lower values of \( \varepsilon_{p,g}, \varepsilon_w \), and \( \tau_{pp,g} \) resulting in only a modest amount of neutrinos produced at high energies. We will show below that if the starburst fraction is high and CRs with energies up to \( \sim 10 - 100 \text{ PeV} \) are sufficiently confined, the majority of the observed high-energy diffuse neutrino flux can be explained using HNe or other sources in starburst galaxies.

For the subsequent propagation in the galaxy cluster or group, the average magnetic field and coherence length are parameterized as \( B_{cl} \sim 10^{-6} B_{cl,-6} \) and \( \ell_{c,cl} \sim 30 \text{ kpc} = 10^{23} \ell_{23} \) cm. This implies \( \varepsilon_{p,cl} \sim 5.6 \times 10^{18} Z \ell_{23} B_{cl,-6} \) eV and \( D_{*,cl} \sim 1.4 \times 10^{32} Z \ell_{23} \) cm\(^2\) s\(^{-1}\). For a cluster of \( 10^{15} M_\odot \) the virial radius is \( R_{cl} \sim 2.6 M_{15}^{1/3} \) Mpc \( \simeq 8 \times 10^{24} M_1^{1/3} \) cm and the diffusion time is \( t_{d,cl} = R_{cl}^2/6D \). At the maximum proton energy this is \( t_{d,cl}(\varepsilon_{p,max}) \sim 2.3 \times 10^{17} M_{15}^{1/3} \ell_{23}^{-2/3} B_{cl,-6}^{1/3} Z^{1/3} n_{cl,-4}^{-1/18} \varepsilon_{p,0}^{-1/6} E_{B,2}^{-1/3} M_{cl,1}^{2/9} \varepsilon_{p,16.94}^{-1/3} \) s. Similarly to the galactic component mentioned above, there is a spectral break when the diffusion time exceeds the injection time of CRs [81]. Assuming CR injection effectively occurs during the Hubble time at the corresponding redshift (i.e., \( t_{age}(z) = f_z dz \left| \frac{dz}{dt} \right| \)), for a cluster located at redshift \( z = 1 \) such a break occurs at an energy \( \varepsilon_{p,cl} \sim 4 \times 10^{17} Z \) eV. If the cluster break energy is higher than the maximum HNe energy, CR diffusion does not significantly affect the fraction of CRs that interact in the intracluster medium. The cluster \( pp \) optical depth is again \( \tau_{pp,cl} = n_{cl} \kappa \sigma_{pp} c \min[t_{d,cl}, t_{age}] \), assuming a typical intracluster gas density \( n_{cl} \simeq 10^{-4} n_{cl,-4} \text{ cm}^{-3} \), at high redshifts (e.g.,
\[ z = 1 \] \( t_{\text{age}} \ll t_{d,\text{cl}}(\varepsilon_{p,\text{max}}) \) so that

\[ \tau_{pp,\text{cl}}(\varepsilon_{p,\text{max}}) \sim 2.7 \times 10^{-2} n_{\text{cl},-4}(t_{\text{age}}/5.8 \text{ Gyr}) \] (2.7)

For more nearby clusters \( \tau_{pp,\text{cl}} \) increases as the cluster age approaches the local Hubble time, although the density is also redshift-dependent.

For a diffusion exponent \( \alpha \) different from the value 1/3 used as an example above, the values of \( \tau_{pp,g} \), \( \tau_{pp,\text{cl},H} \), etc., are calculated similarly and are somewhat different, as can be seen in the numerical results discussed in the next section.

Further \( pp \) collisions occur in the intergalactic medium after the CRs escape the cluster, but with the intergalactic target density \( n_{\text{igm}} = 2.5 \times 10^{-7}(\Omega_b h^2/0.022) \text{ cm}^{-3} \), and a total flight time limited by \( t_H \sim 10^{10}\text{yr} \), the corresponding \( \tau_{pp,\text{igm}} \) is negligible compared to the previous two contributions.

### 2.5 Diffuse neutrino flux

When high-energy CRs undergo \( pp \) interactions with the ambient intragalactic and intrachannel medium, charged and neutral pions are created which subsequently decay to neutrinos and gamma-rays respectively. On average, the resulting neutrino and parent CR energies can be related by \( \varepsilon_{\nu} \sim 0.03 - 0.05 \varepsilon_p \). As a result, the diffuse neutrino flux (per flavor per logarithmic interval of energy) can be estimated using the CR energy injection rate similarly to what is done for GRBs [39,44], as

\[ \varepsilon_{\nu}^2 \Phi_{\varepsilon_{\nu}} = \frac{c}{4\pi} \int_0^z \sum_i f_{i,pp} \frac{\varepsilon_p Q_{\varepsilon_p}}{6} \left(1 + z'\right)^4 \left| \frac{dt}{dz'} \right| dz', \] (2.8)

where the physical CR energy injection rate per energy bandwidth at a given redshift \( z \) is related to Equations (2.1, 2.2), cosmological evolution is taken into account by the scale factor \( S(z) \) so the normalized physical star formation rate is

\[ \left(\varepsilon_p Q_{\varepsilon_p}\right)_{\text{phys}}(z) = \left[\left(\varepsilon_p Q_{\varepsilon_p}\right)_{\text{ln}} + \left(\varepsilon_p Q_{\varepsilon_p}\right)_{\text{sn}}\right] (1 + z)^3 S(z) \] (2.9)

with

\[ S(z) = \left[ (1 + z)^{a_\eta} + \left(\frac{1 + z}{B}\right)^{b_\eta} + \left(\frac{1 + z}{C}\right)^{c_\eta} \right]^{1/\eta} \] (2.10)
where \( a = 3.4, b = -0.3, c = -3.5, \eta \approx -10, B \approx 5000 \), and \( C \approx 9 \) \cite{97,98}.

The sum in Eq. (2.8) is over the different galactic and cluster/group contributions. We assume a fraction \( \xi_{\text{sbg}} \) of neutrinos are produced in starburst galaxies with the rest \( \xi_{\text{sfg}} = 1 - \xi_{\text{sbg}} \) produced in normal star-forming galaxies:

\[
\begin{align*}
  f_{pp,\text{sbg}} & = \xi_{\text{sbg}} \left( 1 - e^{-\tau_{pp,g,\text{sbg}}} \right) \\
  f_{pp,\text{sfg}} & = \xi_{\text{sfg}} \left( 1 - e^{-\tau_{pp,g,\text{sfg}}} \right) \\
  f_{pp,\text{cl}} & = \left( 1 - e^{-\tau_{pp,\text{cl}}} \right) \times \left[ \xi_{\text{sbg}} e^{-\tau_{pp,g,\text{sbg}}} + \xi_{\text{sfg}} e^{-\tau_{pp,g,\text{sfg}}} \right]
\end{align*}
\]  

(2.11)

Note that in the last line of Eq. (2.11) only CRs which escape from the galaxies can contribute to the cluster component.

For the cluster/group parameters and the average galaxy parameters taken in Section 2.4, the diffuse neutrino flux per flavor for a \( D \propto \varepsilon_p^{1/3} \) diffusion coefficient is shown in Figure 2.1 (top panel). Here the contributions of the \( pp \) interactions in the galaxies are indicated both for the SN and HN components. In the same figure, the resulting diffuse gamma-ray flux is also shown, resulting from the corresponding \( \pi^0 \) decays and subsequent pair cascades in the intergalactic medium, which are discussed in Section 2.6. A similar calculation for a \( D \propto \varepsilon_p^{1/2} \) Kraichnan type diffusion coefficient is shown in the bottom panel of Figure 2.1.

The situation depends strongly on the diffusion coefficients of both galaxies and clusters, which are uncertain especially at high energies, and, due in large part to uncertainties in the magnetic coherence length. For example, the diffusion coefficient for normal galaxies used in Figure 2.1 is 10 times lower than the value obtained for our Milky Way. While this discrepancy is alleviated by inhomogeneities, the diffuse Galactic emission suggests that the CR spectral break is much lower since the observed gamma-ray spectrum is already steep at GeV energies \cite{99}. In Figure 2.2, we conservatively use the diffusion coefficient suggested for our Milky Way by \cite{100} for normal star-forming galaxies. We then use the scaling relation \( D \propto r_L(\varepsilon_p)^{\alpha} \) (see Section 2.4) to determine the diffusion coefficient in starbursts. Since the break energy is sensitive to the diffusion coefficient, one sees that the diffuse neutrino background cannot be explained by star-forming galaxies in this case, even with an optimistically high fraction of HN kinetic energy converted to CRs (i.e., \( 7.5 \times 10^{51} \) erg). At high energies the galaxy contribution may not be appreciable. At and below \( \varepsilon_{p*,g} \approx 1.11 \times 10^{18} f_{g,20} B_{g,-3.7} \) eV, however, the galactic contribution becomes considerable, \( \tau_{pp,g}(\varepsilon_{p*,g}) \approx 0.33 \), overcoming the cluster.
Figure 2.1. Diffuse flux per flavor of neutrinos (solid black) and gamma-rays (dash-dot) from HNe and SNe, for a diffusion coefficient (top): $D \propto \varepsilon_p^{1/3}$, (bottom): $D \propto \varepsilon_p^{1/2}$, in both the host galaxy and cluster. For both figures HNe and SNe release on average $2.8 \times 10^{51}$ and $4.8 \times 10^{49}$ ergs of CR energy respectively, and the proton spectral index is $\gamma_p = 2$. The black line with white circles denotes the measured flux of atmospheric neutrinos [101]. The SBG scale height, density, and magnetic field strength are $H_{sbg} = 300$ pc, $n_{sbg} = 200$ cm$^{-3}$, and $B_{sbg} = 200$ $\mu$G and are represented by red lines. For normal star-forming galaxies $H_{sfg} = 1000$ pc, $n_{sfg} = 1$ cm$^{-3}$, and $B_{sfg} = 6$ $\mu$G; they are represented by blue lines. The contribution from HNe are marked with solid lines colored while those from the SNe are dashed. The solid green line denotes the total cluster contribution (i.e., HNe and SNe from both types of galaxies). Green data points correspond to the Fermi measurements of the extragalactic diffuse gamma-ray background [102]. Black points correspond to the IceCube measurements of astrophysical neutrinos [53], note that two of the low energy data points are within the gray lines of the error bars of the atmospheric flux.
Figure 2.2. Diffuse flux per flavor of neutrinos (solid black) and gamma-rays (dash-dot) from HNe and SNe which release on average $7.5 \times 10^{51}$ and $5 \times 10^{49}$ ergs of CR energy respectively, with a phenomenologically motivated diffusion coefficient based on observations of CR diffusion in the Milky Way galaxy (see text for details). In this case, cluster contributions are dominant at high energies.
contribution at the same energy. For this combination of parameters the cluster and group contribution should be dominant, and it is possible to explain the hard spectrum of the diffuse neutrino background. Note that the parameters used for the cluster/group contribution to the diffuse neutrino flux are optimistic, and massive clusters are only a fraction of the total cluster population.

Returning to the parameters used in Figure 2.1, the flux resulting from average host galaxies with a smaller (top) and larger (bottom) fraction of CRs produced in starburst galaxies is shown in Figure 2.4. Here the fraction of HN/SN CR energy was adjusted ad hoc in order to fit the observed neutrino flux with $E_{cr,hn} = 5 \times 10^{51}$ erg and $E_{cr,sn} = 2.2 \times 10^{50}$ erg for $\xi_{sbg} = 0.01$, and $E_{cr,hn} = 10^{51}$ erg and $E_{cr,sn} = 2.5 \times 10^{49}$ erg for $\xi_{sbg} = 0.5$ respectively. The diffusion coefficient was taken to be $D \propto \varepsilon_p^{1/3}$ while leaving the remaining parameters unchanged.

Figures 2.1 and 2.4 were calculated for “typical” star-forming galaxies with parameters as given above, and for a proton injection spectrum $\gamma_p = 2$. We consider next the SFG and SBG contributions using the same parameters, but with a proton injection index $\gamma_p = 2.1$; results are shown in Figure 2.3.

The effect of $p\gamma$ interactions in the galactic and intracluster medium is sub-dominant relative to the $pp$ collisions in the relevant energy range, although it becomes dominant at very high energies [103].
Figure 2.3. Same as Figure 2.1, but for a proton index $\gamma_p = 2.1$, $E_{\text{cr,ln}} = 3.5 \times 10^{51}$ erg, and $E_{\text{cr,sn}} = 10^{50}$ erg. top: $D \propto \epsilon_p^{1/3}$, bottom: $D \propto \epsilon_p^{1/2}$. 
Figure 2.4. Same as Figure 2.1, with $D \propto \varepsilon_{p}^{1/3}$ but top: $\xi_{sbg} = 0.01$ with $E_{cr, hn} = 5 \times 10^{51}$ erg, and $E_{cr, sn} = 2.2 \times 10^{50}$ erg bottom: $\xi_{sbg} = 0.5$ with $E_{cr, hn} = 10^{51}$ erg, and $E_{cr, sn} = 2.5 \times 10^{49}$ erg
2.6 Gamma-Ray Cascades

The same $pp$ interactions which produce neutrinos also produce high-energy gamma-rays with typical energy $\epsilon_{\gamma} \sim \epsilon_p/10$. Note that because of this connection, their resulting flux can be related by $\epsilon_{\gamma}^2 \Phi_{\epsilon_{\gamma}} = 2^{\epsilon_p - 1} \epsilon_{\nu}^2 \Phi_{\epsilon_{\nu}} |_{\epsilon_{\nu}=0.5 \epsilon_{\gamma}}$. When gamma-rays with energy $\epsilon_{\gamma} \gtrsim 100$ GeV are injected into intergalactic space sufficiently far from Earth (i.e., $\sim$ few Mpc), they undergo $\gamma\gamma$ interactions with extragalactic background light (EBL) photons producing electron/positron pairs. The pairs scatter additional EBL photons via the inverse Compton mechanism generating an electromagnetic cascade. The resulting cascaded gamma-ray spectrum takes a universal form, (e.g., [104,105]):

$$\epsilon_{\gamma} dN_{\gamma}/d\epsilon_{\gamma} \propto G_{\epsilon_{\gamma}} = \left\{ \begin{array}{ll} \left( \epsilon_{\gamma}/\epsilon_{\gamma}^{br} \right)^{-1/2} & \epsilon_{\gamma} \leq \epsilon_{\gamma}^{br} \\ \left( \epsilon_{\gamma}/\epsilon_{\gamma}^{br} \right)^{1-s} & \epsilon_{\gamma}^{br} < \epsilon_{\gamma} \leq \epsilon_{\gamma}^{cut} \end{array} \right. \quad (2.12)$$

The characteristic energies $\epsilon_{\gamma}^{cut}$ and $\epsilon_{\gamma}^{br}$ given above are defined by $1 = \tau_{\gamma\gamma} \left[ \epsilon_{\gamma}^{cut}, z \right]$ and $\epsilon_{\gamma}^{br} = 0.0085 (1+z)^2 \left( \epsilon_{\gamma}^{cut}/0.1 \text{ TeV} \right)^2$ respectively and the high-energy spectral index is generally taken to be $s \sim 2$. Here $\tau_{\gamma\gamma}$ is the optical depth for a high energy photon traveling through intergalactic space, the values for which are from model C of [106].

There is also an attenuated component to the observed gamma-ray flux from photons with energy $\epsilon_{\gamma} \lesssim \epsilon_{\gamma}^{cut}$ which can be calculated similarly to Equation (2.8)

$$\epsilon_{\gamma}^2 \Phi_{\gamma}^{\text{att}} = \frac{c}{4\pi} \int dz \left| \frac{dt}{dz} \right| e^{-\tau_{\gamma\gamma}[(1+z)\epsilon_{\gamma}, z]} \frac{1}{(1+z)^4} \left[ \frac{2\epsilon_p^{-2}}{3} \sum_i f_i,pp \left( \epsilon_p Q_{\epsilon_p} \right)_{\text{phys}} \right] \left. \epsilon_{\nu,i} = 10(1+z)\epsilon_{\gamma} \right|_{\epsilon_{\nu}=0.5 \epsilon_{\gamma}} \epsilon_{\nu,i}$$

which combined with Equation (2.12) can be compared with Fermi-LAT measurements of the extragalactic diffuse gamma-ray background [102]. Figures 2.1 through 2.3 show our calculated diffuse flux of neutrinos and gamma-rays along with data from IceCube and Fermi.
2.7 Discussion and Summary

In this chapter, we discuss the starburst scenario in light of the new 10 TeV neutrino data. Although there are large systematic uncertainties involved in removing the atmospheric muon background at such low energies [53], it may be challenging to explain the diffuse neutrino flux in the whole energy range with a single power-law component with $\gamma_p \sim 2$. Adding the SNe contribution enables us to explain the low-energy data, but we find that constraints from the diffuse gamma-ray background are quite stringent [44]. If the CR energy input by SNe is a factor of two larger than that by HNe, the diffuse gamma-ray background is violated. Additional constraints could be placed on the cluster contribution by considering the concomitant radio emission, although in the strong evolution case these limits are weaker than the ones imposed by the accretion shock scenario [107]. From the gamma-ray limits, we conclude that if the diffuse neutrino background in the PeV range originates mainly from HNe (and their host galaxies), the HN contribution should be larger than or at least comparable to the SN contribution. However, interestingly, in cases where the cluster/group contribution is mainly responsible for the diffuse neutrino flux, it is still possible for the SN contribution to overwhelm the HN contribution (e.g. the top panel of Figure 2.4).

The strong case scenario, where the $\sim 10$ TeV neutrino data are explained by CR reservoirs, has an interesting feature that can be tested soon. As proposed by Murase et al. [44], CR reservoirs can give a common explanation for both the diffuse neutrino and gamma-ray backgrounds. In general, the contributions from starbursts and other sources to the neutrino flux above 100 TeV result in subdominant contributions to the diffuse gamma-ray background. However, as we show above any source that contributes significantly to the 10 TeV diffuse neutrino flux in the $pp$ scenario must also account for almost all of the diffuse gamma-ray background. It is commonly believed that the diffuse gamma-ray background is dominated by unresolved blazars [108,109], implying a comparatively smaller starburst contribution. Although there are still significant uncertainties in the modeling of both blazar (e.g., [110,111]) and starburst contributions (e.g., [66,112]), my results imply that the strong case scenario can be tested by an improved understanding or characterization of the diffuse gamma-ray background.

If, for example, it is proven that blazars are responsible for $\gtrsim 50\%$ of the observed diffuse gamma-ray background, the starburst contribution to the diffuse neutrino background at low energies should be small, especially if the CR energy input from SNe is
comparable to or larger than that from HNe. Specifically, tighter constraints on the unresolved blazar fraction of the diffuse gamma-ray background measured by \textit{Fermi} and possibly the High Altitude Water Cherenkov observatory \cite{113}, as well as the low energy ($\varepsilon_{\nu} \lesssim 100$ TeV) spectral shape of the astrophysical neutrino flux as analyzed by Aartsen et al. \cite{53}, will impose limits to the corresponding fluxes from galaxy clusters/groups. If there is little room for the CR reservoirs, other sources need to be responsible for the low-energy neutrino component. For instance, there might be a significant contribution from Galactic sources. Although the Galactic diffuse emission by CRs propagating in our Milky Way cannot provide the main contribution \cite{73,114}, some extended sources such as the Fermi Bubbles \cite{73,115} or nearby HN remnants \cite{70,73} could be viable. Alternatively, the diffuse neutrino background might be produced mainly by hidden neutrino sources via $p\gamma$ processes \cite{75,80,116}. The advantage of the strong case considered here is that it can be tested by multimessenger approaches. It has been commonly believed that Galactic CRs come from SN remnants. If the diffuse neutrino background is dominated by star-forming galaxies, our results imply that even Galactic CRs may include significant contributions from past HN remnants.

CR acceleration to energies $\gtrsim 10^{16} - 10^{17}$ eV has also been proposed in other accelerators, such as shocks in AGN jets, (e.g., \cite{33,79,84}), AGN winds \cite{66,74}, and AGN cores \cite{77,78,116}. While in such cases neutrinos can come from AGNs themselves, CRs escaping from AGNs can also produce neutrinos in intergalactic space, which may give a significant contribution to the diffuse neutrino background \cite{44}. Other possibilities are accretion shocks onto clusters of galaxies \cite{81,117}, galaxy mergers in clusters \cite{85}, and GRBs including trans-relativistic SNe or low-luminosity GRBs \cite{118,119}. In principle, the discussion of the cluster/group propagation effects discussed above also applies to any intracluster sources. The main difference between these other sources and SNe/HNe (or sources inside galaxies in general) is that the CRs accelerated in the former do not undergo $pp$ interactions in the galactic gas, but only in the intracluster gas, whereas CRs from HNe and galactic sources undergo $pp$ interactions in both the host galaxy and the cluster/group. This disparity may be relevant at sub-PeV and TeV energies, where the spectral shape of the neutrino flux can provide clues to the source. In this energy range we expect the advective escape and Hubble times to dominate the galactic and cluster diffusion times respectively, and at different critical energies (i.e., $\varepsilon_{w}$ and $\varepsilon_{p,cl}$ as in Section 2.4). Therefore, for galactic sources (that are extragalactic), a soft spectrum is typically expected at energies below the maximum acceleration energy, with $\varepsilon_{\nu}^{2} \Phi_{\nu} \propto \varepsilon_{\nu}^{-\alpha}$
(for a diffusion time $\propto \varepsilon^{-\alpha}$), with a leveling off of the slope to $\varepsilon^2 \Phi_\nu \sim$constant below about $\varepsilon_{\nu,g} \sim 130 \ (2/(1+z)) \ \text{TeV}$ (assuming, e.g., a galactic magnetic field strength of 200$\mu$G and diffusion exponent $\alpha = 1/2$). Sources which release their CRs directly into the intracluster medium on the other hand are expected to produce relatively flat neutrino spectra below a break around $\varepsilon_{\nu,\text{cl}} \sim \text{few} \ (2/(1+z)) \ \text{PeV}$, steepening above that. Such a spectral softening of the cluster contribution is not seen for the HNe model considered here (e.g., in Figure 2.2) because the break energy due to diffusion is lower than the cutoff energy imposed by maximally accelerated CRs. While both $\varepsilon_{\nu,g}$ and $\varepsilon_{\nu,\text{cl}}$ are subject to large uncertainties in parameters, the presence of their corresponding features could be suggestive of sources which are embedded in star-forming or starburst galaxies.

As shown by Murase et al. [44], in the $pp$ scenario, the neutrino spectrum cannot be softer than about $\varepsilon^2 \Phi_\nu \propto \varepsilon^{-0.2}$ at low energies for the corresponding gamma-ray component not to violate the Fermi measurements of the diffuse gamma-ray background [102]. At the same time, a flat spectrum at moderate to high energies creates tension with the non-detection of neutrinos with energies near the Glashow resonance at $\varepsilon_\nu \sim 6 \ \text{PeV}$, which necessitates a neutrino spectral shape near that energy steeper than $\varepsilon^2 \Phi_\nu \propto \varepsilon^{-0.3}$ [120]. Such a spectral break can occur if acceleration stops or there is a transition in the diffusive escape time (i.e., $D(\varepsilon_{cr}) \propto \varepsilon_{cr}^\alpha \rightarrow D(\varepsilon_{cr}) \propto \varepsilon_{cr}^2$) around $\varepsilon_{cr} \sim 240 \ ((1+z)/2) \ \text{PeV}$ (e.g., [44,68]). As can be seen in Figures 2.1 and 2.4, the model presented here can also resolve this tension. The neutrino spectrum is flat at low energies $\varepsilon_\nu \lesssim 130 \ \text{TeV}$ and softens to $\varepsilon^2 \Phi_\nu \propto \varepsilon^{-\alpha}$ slightly before the Glashow resonance, while the corresponding diffuse gamma-ray spectrum is below the Fermi measured flux.

We have shown in Section 2.5 that the high-energy diffuse neutrino flux could potentially be explained by HNe, predominantly those located in dense starburst galaxies (e.g., the red solid curve in the bottom panel of Figure 2.1) especially for $\varepsilon_\nu \gtrsim 50 \ \text{TeV}$. For neutrinos with this energy and below, the SNe in both starburst and normal star-forming galaxies contribute significantly to the diffuse flux and produce a “bump” in the spectrum. Reasonable fits by eye are obtained for the diffuse neutrino flux measured by IceCube including the latest TeV data by using reasonable parameters for the sources as well as the diffusion properties in hosting structures. Such a flux also approximates but does not violate the diffuse gamma-ray background measured by Fermi. This does not mean that SNe and HNe are necessarily the only sources contributing to the neutrino

‡As we were preparing to submit our calculation, a similar qualitative conclusion was posted by Chakraborty & Izaguirre [121], who however did not consider the diffuse gamma-ray constraints from Fermi observations that are very important.
and gamma-ray diffuse backgrounds. It supports, however, the case for CR reservoirs such as clusters and groups being promising sources which could contribute at least a significant fraction of these backgrounds, without violating both CR [122] and gamma-ray constraints.
Chapter 3  
High-energy Neutrino Flares from X-Ray Bright and Dark Tidal Disruption Events

3.1 Introduction

Supermassive black holes (SMBHs) with masses $\gtrsim 10^5 M_\odot$ can disrupt stars whose orbits pass within the tidal disruption radius (see a review by Komossa [123] and references therein). For a SMBH with $M_{\text{BH},6} = 10^6 M_\odot$ disrupting a Sun-like star, this distance is $r_T \approx 10^{13} \text{ cm } M_{\text{BH},6}^{-1/3} M_\odot^{1/3} R_\odot$. Roughly half the disrupted stellar material becomes unbound from the SMBH, while the remaining plasma circularizes to form an accretion disk [124–126]. It has been argued that material from a tidal disruption event (TDE) would accrete at a super-Eddington rate $\dot{M}_{\text{Edd}} \sim 2 \times 10^{-3} M_\odot \text{ yr}^{-1}$ (e.g., [127–130]), and potentially launch a jet (e.g., through the Blandford-Znajek mechanism [131–135]).

Relativistic jets have been promising candidates of high-energy particle accelerators, as commonly considered in the literature of Gamma-Ray Bursts (GRBs) and active galactic nuclei (AGN). The X-ray detections of jetted TDEs [136–138] have stimulated proposals of TDEs as potential ultra-high-energy cosmic-ray (UHECR) accelerators [139,140] and high-energy cosmic neutrino sources* [141–143]. In particular, the recent discovery of astrophysical neutrinos by IceCube [24,52] gives us a good motivation to revisit various astrophysical transients, including TDEs, as multi-messenger sources. The sources of

* [141,142] calculated the diffuse neutrino intensity from TDEs, based on the giant flare scenario suggested by [139].
the very-high-energy neutrinos with energies $E_\nu \approx 10$ TeV are unknown. The arrival directions and flavor composition of these neutrinos are consistent with an isotropic diffuse flux of astrophysical neutrinos of extragalactic origin [18]. Sources such as GRBs [39] and Blazars [144] do not account for a significant fraction of the observed diffuse neutrino flux [55, 145, 146]. Cosmic-ray reservoirs such as Starburst Galaxies [64, 147] and Galaxy Clusters/Groups [81, 103] are both neutrino and $\gamma$-ray bright and can significantly contribute to the diffuse flux (e.g., [40, 44, 66, 107, 121, 148–150]), but only above 0.1 PeV energies because of stringent $\gamma$-ray constraints [151].

Models of $\gamma$-ray “hidden” sources—which mask their $\gamma$-ray emission at $\gtrsim 1$ GeV energies—are not constrained by Fermi-LAT or IceCube analyses [50]. Such sources include choked GRB jets [48, 75, 152], newborn pulsars [153], white dwarf mergers [154], and high-redshift galaxies$^\dagger$ [158]. X-ray bright, jetted TDEs are an example of hidden sources, since high-energy $\gamma$-rays are significantly absorbed via two-photon annihilation [136]. Alternatively, if an unbound material or optically-thick wind forms a spherical circumnuclear envelope, it could obscure or reprocess the non-thermal emission from a relativistic jet (i.e., a choked-jet TDE), making such events potential “hidden” neutrino sources [49].

In this chapter, we revisit high-energy neutrino production in TDE jets. We calculate neutrino fluxes from both X-ray bright successful TDE jets and possible choked TDE jets that could occur if the jets are not powerful enough. For Sw J1644+57 (hereafter Sw 1644), we also use IceCube data for upgoing muon neutrino events to search for coincidences between neutrino detection and the three jetted TDE candidates seen by Swift-BAT in 2011 [159]. While the right ascension and arrival time is not given for these neutrinos, we place meaningful limits on the baryonic loading factor of jetted TDEs. Then, we evaluate diffuse neutrino intensities of jetted TDEs and discuss the present constraints.

### 3.2 X-Ray Bright TDEs with Successful Jets

We first consider jetted TDEs that have non-thermal X-ray spectra ($0.3 \lesssim \epsilon_\gamma \lesssim 150$ keV) to determine the seed photon density for $p\gamma$ interactions. Only Sw 1644 was observed early enough by Swift XRT to fit an SED [136], though two jetted TDE candidates Sw

$^\dagger$Note that the $\gamma$-ray cutoff energy due to the extragalactic background cannot be lower than $\sim 10$ GeV [155–157].
J1112-82 [138] and Sw J2058+05 [137], show similar peak X-ray luminosities. We use a log-parabolic fit to the SED of Sw 1644 [136], which is given by

$$\epsilon^2 n_\epsilon = A (\epsilon/\epsilon_{pk})^{0.5-0.25 \hat{a} \log(\epsilon/\epsilon_{pk})},$$

(3.1)

where $\epsilon_{\gamma, pk} = 200$ keV, $\hat{a}$, and $A$ are fitting parameters (see Figure 3.1). The peak energy corresponds to $\bar{\epsilon}_{\gamma, pk} = 20$ keV in the jet plasma comoving frame for Lorentz factor $\Gamma = 10$. If the non-thermal X-rays are produced by synchrotron emission from leptons in the jet, the template from Sw 1644 can be adapted to fit jetted TDEs of different luminosities. The luminosity changes with time and the maximum luminosity reaches $L_{max} \equiv \epsilon L_{\epsilon, pk}^\gamma \sim 10^{48}$ erg s$^{-1}$, which lasts for 3 days (implying a duration of $t_{dur} \sim 2 \times 10^5$ s in the cosmic rest frame). However, it decreases after the peak, and the median luminosity in the 0.3-10 keV band is $L_{[0.3,10 \text{ keV}]} = 8.5 \times 10^{46}$ erg s$^{-1}$, considering the emission with duration of $t_{dur} \sim 10^6$ s. The corresponding bolometric luminosity is $L_\gamma = 5.7 \times 10^{47}$ erg s$^{-1}$, and we use the bolometric radiation energy of $L_\gamma t_{dur} = 5.7 \times 10^{53}$ erg at $L_\gamma = 10^{47}$ erg s$^{-1}$, which is relatively conservative for Sw 1644. As we see below, the meson production efficiency ($f_{\gamma \gamma}$) is proportional to $L_\gamma$ so neutrino production is expected to be dominated by the high state that lasts during $t_{\text{high}}$. Thus, for estimates of neutrino fluences, we will use $L_{[0.3,10 \text{ keV}]}^{\gamma} = 10^{48}$ erg s$^{-1}$ corresponding to $L_{\gamma, pk} \sim 12 L_\gamma$. In a relativistic jet, the turbulent magnetic field energy is parametrized as $B^2/(8\pi) = \xi_B L_{\gamma, pk}/(4\pi r^2 \Gamma^2 c)$ during the high state, and for TDEs with different luminosities, the peak synchrotron flux goes as $\epsilon F_{\epsilon, pk}^\gamma \propto L_{\gamma, pk}$ and $\epsilon_{pk} \propto B \propto L_\gamma^{1/2} r_{\text{em}}^{-1}$, where $r_{\text{em}}$ is the internal dissipation radius.

The locations of the non-thermal emission from both Sw J1644 and 2058 are believed to be close to the jet base since both show variability with $\delta t \sim 10^2$ s. Assuming—as inferred for Sw 1644 [136]—that TDE jets are modestly relativistic ($\Gamma \sim 10$), the emission radius is estimated to be $r_{\text{em}} \sim 3 \times 10^{14}$ cm $\Gamma^2 t_{\text{dur}}^2$, which corresponds to a few hundred Schwarzschild radii from the SMBH [160]. We assume $\Gamma$ and $\delta t$—and therefore the internal dissipation radius $r_{\text{em}}$—are similar for jetted TDEs of all luminosities. This location of the emission region is consistent with the observation of fast X-ray variability, while optical emission may be produced by the jet’s forward shock at a distance of $\sim 10^{15}$ cm from the SMBH (e.g., [162]). As we will show below, X-ray bright jetted TDEs are likely to have successful jets with isotropic equivalent luminosity $L \approx 10^{44.5}$ erg s$^{-1}$.

‡ See, however, Kara et al. [161] who argued that the soft X-ray emission may come from the accretion disk and produced $\sim 10$ Schwarzschild radii from the SMBH.
Figure 3.1. Non-thermal SED of Sw 1644 as a function of the comoving energy $\tilde{\epsilon}_\gamma$ in the high (solid) and median (dash-dot) luminosity states. One sees that there is a peak around 10 keV, corresponding to $\epsilon_{\gamma, pk} = 200$ keV.
3.3 X-Ray Dark TDEs with Choked Jets?

3.3.1 Jet Propagation

It has been generally assumed that TDE jets are powered by the Blandford-Znajek process. However, not all TDEs may have visible jets. The observed rate of X-ray bright jetted TDEs is $\rho_{0,X-TDE} \sim 0.03 \text{ Gpc}^{-3} \text{ yr}^{-1}$, which is much smaller than the total TDE rate at $z = 0$, $R_{0,all-TDE} \sim 10^3 - 10^5 \text{ Gpc}^{-3} \text{ yr}^{-1}$ (see Kochanek 2016 [163] and references therein). This may simply come from the fact that only a fraction of the TDEs may have jets. However, an alternative possibility is that jetted TDEs are more common but many of them are stuck in the TDE material [49]. It has been speculated that SMBHs could have a thick circumnuclear material of size $r_{out} \sim 3 \times 10^{15} \text{ cm}$ produced by the breakup of the star [131,164]. Since the visible TDE emission radius is inside of the circumnuclear envelope, the non-thermal emission will only be visible if the jet breaks out. In this scenario, the SMBH is surrounded by an optically thick material formed from the disrupted stellar material that is not bound in an accretion disk ($\sim 0.5M_\star$).

This envelope material may have a wind density distribution with $\rho_{\text{env}} \propto r^{-2}$, in which the material is dominated by disk-driven winds. We consider a density profile proposed by Loeb & Ulmer [164] and De Colle et al. [131],

$$\rho_{\text{env}}(R) = \frac{f_{\text{TDE}}M_\star}{4\pi \ln(r_{out}/r_{in})r^3} \equiv \rho_{in}r_{in}^3r^{-3} \quad (3.2)$$

where the envelope may be defined by inner and outer radii $r_{in} \sim r_T \sim 3 \times 10^{13} \text{ cm}$ and $r_{out} \sim 3 \times 10^{15} \text{ cm}$ and $\rho_{in} = 6.4 \times 10^{-10} \text{ g cm}^{-3}$. Correspondingly, the radiation temperature is assumed to be

$$T_{\text{env}} = 10^6 \text{ K} \left(\frac{r}{r_{in}}\right)^{-1}. \quad (3.3)$$

The envelope material changes the dynamics of the jet launched by the SMBH, in a way that is similar to the star/jet interaction in a long GRB in the collapsar scenario [165,166]. Specifically, the dimensionless velocity of the jet head can be related to the isotropic equivalent jet kinetic luminosity $L$ and the density of the envelope [165], where

$$\tilde{L} \approx \frac{L}{4\pi r^2 \rho_{\text{env}}c^3}.$$

When $\tilde{L} \ll \theta_j^{-4/3}$—as it is for the case of choked jet TDEs—collimation shocks are formed.
and they change the jet’s initially conical shape to a cylindrical one. The position of the jet head approximates the forward position of the collimated jet, which is given by [166]

\[ r_h \approx 2.5 \times 10^{15} \text{ cm} \ t_{\text{eng}}^{3/2} r_{\text{in}}^{1/2} \frac{t_{\text{eng}}^{-1/2}}{\theta_{j,-1}^{-3/2}} \]

assuming that the jet is powered by the accretion disk for a time \( t_{\text{eng}} \sim 10^6 M_{\text{BH}}^{1/2} \) s, corresponding to the period of the most tightly bound material in the accretion disk [132]. Using the assumed value for \( t_{\text{eng}} \) and the relationship above for the position of the jet head, we find that only weak jets with luminosity

\[ L < 2 \times 10^{44} \text{ erg s}^{-1} \ t_{\text{eng}}^{-3} r_{\text{in}}^{-3} r_{\text{in,13.5}}^{3} r_{\text{out,15.5}}^{2} \theta_{j,-1}^{2} \]

can be choked by the envelope (i.e., \( r_h(t_{\text{eng}}) \lesssim r_{\text{out}} \)). Assuming the fraction of jet luminosity converted to photons \( \xi \simeq 0.2 \) is the same for all jets, we extrapolate the luminosity function of Sun et al. [167] to estimate the number of “hidden” jetted TDEs, which occur with jet luminosities \( L \lesssim 10^{44.5} \text{ erg s}^{-1} \). Note that the existence of jets with \( L_j \approx (\theta_{j,-1}^2/2)L \lesssim 10^{42} \text{ erg s}^{-1} \) (that is a sub-Eddington luminosity) is not guaranteed, and numerical simulations have not found such weak jets in the current setup [133,135]. If the cocoon pressure is assumed to be constant, the collimation shock radius can be as large as \( r_{\text{cs}} \sim r_h \). However, as indicated in Mizuta & Ioka [166], a pressure gradient exists in more realistic situations (especially if \( r_{\text{cs}} \ll r_h \)), and the collimation shocks may occur near the inner envelope radius \( r_{\text{cs}} \sim r_{\text{in}} \). In this chapter, in order to discuss an optimistic case, we assume that the collimation shock radius is large enough that internal dissipation occurs at \( r_{\text{em}} \sim 3 \times 10^{14} \text{ cm} \).

### 3.3.2 Photon Field Modeling

For bright jetted TDEs such as Sw 1644, the jet’s non-thermal spectrum, and therefore the target photons for \( p\gamma \) interactions are readily available from observations. Choked jets by definition have their photon fields hidden from us. Thus, we must extrapolate their SEDs under several assumptions.

In this chapter, we consider two emission regions. The first is the internal shock site, which can occur near the collimation shock. This scenario is an extrapolation of the X-ray bright TDE jet scenario. In addition to non-thermal photons produced inside the jet, we consider thermal radiation fields provided by the external material. The second is the
termination shock, which exists around the jet head. This is unique to choked jet models. However, if the shock is radiation-mediated, the diffusive shock acceleration of cosmic-rays is inefficient so that high-energy neutrino production is not expected [75]. In the TDE cases, both internal and termination shocks are collisionless and radiation-unmediated, so that one may expect efficient particle acceleration, in principle, if choked jets exist. Note that the meson production efficiency in such hidden cosmic-ray accelerators is always high so that the system is regarded as calorimetric. Cosmic-ray acceleration does not have to be located inside jets. One may consider disk-driven winds as alternative cosmic-ray acceleration sites. Also, observations suggested that outflows are likely to have a complicated structure [168].

In the internal shock scenario, the shock radius is assumed to be \( r_{\text{em}} = 3 \times 10^{14} \) cm. We have two radiation fields. One is thermal radiation from the envelope, which has temperature of \( T = 10^5 \) K at this radius. We checked that thermal radiation from the hot cocoon is sub-dominant in our setup. However, for the radiation from the optically-thick envelope, only a fraction of \( f_{\text{esc}} \approx 1/\tau_T \approx 1.2 \times 10^{-2} e_{\text{in}}^{-1} r_{\text{in},13.5}^{-3} r_{14.5}^2 \) can escape to the optically-thin jet region [75]. Taking into account that the thermal photon number density is boosted by \( \Gamma \) in the jet comoving frame, the comoving thermal photon density in the jet is

\[
n_{\gamma,\text{th}} \approx \Gamma f_{\text{esc}} 16\pi \zeta(3)(kT/hc)^3.
\]

For the non-thermal synchrotron spectrum, we again use a template synchrotron spectrum of Sw 1644 to estimate the target photon density in lower-luminosity jetted TDEs [136], assuming \( E_{pk} \propto L_1^{1/2} r_{\text{em}}^{-1} \). Note that the synchrotron photons are produced by electrons co-accelerated in the internal shocks and are not significantly modified. See §3.2.

The second scenario is the termination shock scenario, in which cosmic-rays are accelerated at the termination shock caused by the cylindrical jet with \( \Gamma_{cj} \approx \theta_j^{-1} \). Note that, in the GRB case, this scenario does not work because the termination shock is radiation-mediated and efficient particle acceleration would not occur. However, in the TDE case, the shock is collisionless since the jet density is small enough. The blackbody spectrum from the jet head is estimated using the photon energy density \( U_{\gamma h} \approx L/(4\pi \Gamma_h^2 r_h^2 c) \approx 84 L_{44.5} \Gamma_h^{-2} r_{h,15.5}^{-2} \) erg cm\(^{-3}\) and the relationship \( U_{\gamma h} = aT_h^4 \), where the jet head temperature is \( T_h \approx 1.0 \times 10^4 L_{\gamma,44.5}^{1/4} \Gamma_h^{-1/2} r_{h,15.5}^{-1/2} \) K [169, 170]. The jet head is non-relativistic and \( \Gamma_h \sim 1 (\beta_h \ll 1) \) is obtained. Assuming the thickness of the jet head is \( \Delta r_h \approx r_h/\Gamma_{cj} \sim 0.1 r_h \), the jet is found to be optically-thin with \( \tau_h \approx n_h \sigma_T \Delta r_h < 1 \), so we do not have to include the suppression factor in this estimate. The photon density is
simply estimated to be \( n_{\gamma,ts} \approx 16\pi \zeta(3)(kT_h/hc)^3 \). Note that the circumnuclear envelope is still optically thick and manages to absorb and scatter any \( \gamma \)-rays so the system is hidden in \( \gamma \)-rays.

### 3.4 Detection of neutrinos from individual TDEs

#### 3.4.1 Semi-Analytic Method of Calculation

We evaluate the fraction of cosmic-rays accelerated in TDE jets that undergo \( p\gamma \) interactions. For \( p\gamma \) interactions, we use a parameterization of the \( p\gamma \) cross section that has been used in the previous publications (e.g., [48, 75, 171, 172]), which utilizes the experimental cross section data and Geant4 simulation package to treat multi-pion production. However, contrary to the above previous works, we take a semi-analytical method rather than calculate neutrino spectra by solving full kinetic equations taking into account all relevant cooling processes for pions, muons, and kaons, as well as neutrino flavor mixing. This is because such a semi-analytical method is faster and more useful for qualitative comparisons with the real data, especially when parameter scans are necessary. Note that detailed numerical treatments of pion and muon cooling can be important if the cooling is strong, as often expected in choked jets [48, 75], but our semi-analytical method works reasonably well for TDE jets in which cooling effects are moderate.

As described above, an X-ray spectral template based on observations of Sw 1644 is used to model the non-thermal spectra from jetted TDEs [136]. The fraction of cosmic-rays which produce neutrinos is given by the ratio between the cooling and dynamic times \( f_{p\gamma} \approx t_{dy} / t_{p\gamma} \), where \( t_{dy} \approx r_{em} / (\Gamma c) \). The former depends on the photohadronic cross section \( \sigma_{p\gamma}(\bar{\epsilon}) \) and the TDE target photon spectra \( n_{\epsilon} \) (see §3.3.2), and we use

\[
t_{p\gamma}^{-1}(\epsilon_p) = \frac{1}{2} \frac{m_p^2 c^3}{\epsilon_p^2} \int \frac{d\epsilon n_{\epsilon}}{\epsilon_{th}^2} \int_{2\epsilon_p^2/\epsilon_{th}}^{2\epsilon_p^2} d\bar{\epsilon} \bar{\epsilon} k_{p}(\bar{\epsilon}) \sigma_{p\gamma}(\bar{\epsilon}),
\]

where \( \bar{\epsilon} = \epsilon_p c/m_p (1 - \cos \theta_{p\gamma}) \) is the photon energy in the CR rest frame and \( \epsilon_{th} \approx 140 \text{ MeV} \) is the threshold energy for \( p\gamma \) interactions.

Our results for X-ray bright TDEs are shown in Fig. 4.4. For X-ray bright TDEs, interactions with synchrotron photons are dominant. The physical setup is similar to X-ray flares of GRBs, as pointed out by [143]. For a photon spectrum given by \( n_{\epsilon} \propto \epsilon^{-\beta} \), the efficiency is \( f_{p\gamma} \propto \epsilon_{\gamma}^{\beta-1} \). Below the peak in \( \epsilon L_{\epsilon} \), the synchrotron spectrum has a
Figure 3.2. Cooling timescales for visible TDEs as a function of parent cosmic-ray energy in the plasma rest frame of the jet. The cooling timescales were calculated in the high state, in which an isotropic X-ray luminosity $L_{[0.3,10 \text{ keV}]} = 10^{48} \text{ erg/s}$ and $\xi_B = 1$ are assumed, with $r_{em} = 3 \times 10^{14} \text{ cm}$ and $\Gamma = 10$. The acceleration timescale is also shown. Note that $f_{p\gamma}$ is enhanced by a factor of 10 during $t_{\text{high}}$. 
spectrum of $\beta \sim 1.5$, and it has a break at $\varepsilon_b(<\varepsilon_{pk})$ (where the photon number becomes the maximum). Around the cosmic-ray energy of interest – $\varepsilon_{pk}^b \approx 1.6 \times 10^6(0.1 \text{ keV}/\varepsilon_b)$ GeV corresponding to the break photon energy of the TDE photon number density – the photomeson production efficiency which is estimated to be $[172]$

$$f_{p\gamma} \sim 1 \frac{L_{\text{max,48}}(E_b/0.01E_{pk})^{0.5}}{r_{\text{em,14.5}}^{2}}(E_b/1 \text{ keV})^{-\frac{1}{2}}(E_b/\varepsilon_{pk}^b)^{\frac{\beta-1}{2}},$$

(3.8)

where an additional factor of 2-3 enhancement due to multi-pion production is included. This is in agreement with our numerical results, taking into account the difference between a log parabolic function and power law (See Fig 4.4).

Our results for TDEs with choked jets are also shown in Figs. 3.3 and 3.4, where low-power jets with $L = 10^{44.5} \text{ erg/s}$ are assumed. The thermal radiation is relevant in both scenarios. In the internal shock scenario, we have

$$f_{p\gamma} \sim 80 \varrho_{\text{in}}^{-1}r_{\text{in,13.5}}^{-3/2} T_{\text{in,13.5}}^{3/2} T_{5}^{3} \Gamma_{1},$$

(3.9)

which is again consistent with our numerical result shown in Fig. 3.3. Note that for the luminosity regime relevant for producing a choked jet, the non-thermal spectrum peaks at frequencies $10^{14} \text{ Hz} \lesssim \varepsilon/\hbar \lesssim 3 \times 10^{15} \text{ Hz}$, leading to $p\gamma$ energy conversion at energies $\varepsilon_p \sim 10 \text{ PeV} - 300 \text{PeV}$. However, non-thermal contributions are smaller than thermal contributions in our case. Note that we do not consider $pp$ interactions. In the internal shock scenario, cosmic-rays may lose their energies via adiabatic cooling before they reach the dense material. (However, cosmic-rays accelerated at collimation shocks may experience subsequent inelastic $pp$ collisions without significant adiabatic losses $[75]$.) In the termination shock scenario, cosmic-rays may cause $pp$ interactions, but the $pp$ efficiency is typically small at $\sim r_h$.

Using our results on $f_{p\gamma}$, we can estimate neutrino fluxes from a single TDE event. For a flat energy spectrum with $\varepsilon_p L_{\varepsilon_p} \propto \text{const.}$, we have

$$\varepsilon_\nu L_{\varepsilon_\nu} \approx \frac{3}{8} \min[1, f_{p\gamma}(\varepsilon_p)] f_{\text{sup}} \varepsilon_p L_{\varepsilon_p} \varepsilon_p \left(\frac{L_{\text{cr}}}{\ln(\varepsilon_{p,\text{max}}/\varepsilon_{p,\text{min}})}\right)^{1/2},$$

(3.10)

where $f_{\text{sup}}$ is the suppression factor due to meson and muon cooling (see below). The observed neutrino energy is therefore related to the jet comoving proton energy for a TDE located at a redshift $z$ with bulk Lorentz factor $\Gamma$, $\varepsilon_p \approx 20(1+z)E_{\nu}/\Gamma$. The baryon
Figure 3.3. The same as Fig. 2, but for the choked TDE jet with $L = 10^{44.5}$ erg/s and $\xi_B = 1$ in the internal shock scenario, with $r_{em} = 3 \times 10^{14}$ cm and $\Gamma = 10$. The temperature of the external field is set to $T = 10^5$ K, and the escape fraction is included.
Figure 3.4. The same as Fig. 2, but for the choked TDE jet with $L = 10^{45.5}$ erg/s and $\xi_B = 1$ in the termination shock scenario, with $r_h = 3 \times 10^{15}$ cm and $\Gamma_h = 1$. The temperature of the jet head is $T_h = 10^4$ K.
loading parameter \( \xi_{cr} \equiv L_{cr}/L_\gamma \) is the ratio of the amount of TDE jet luminosity in protons and high-energy photons [171]. We assume that it has \( \xi_{cr} \sim \epsilon_p/\epsilon_e \sim 1 - 100 \), which is typically required for GRBs and blazars to explain UHECRs. Strictly speaking, its definition depends on the photon energy band and SED. For Sw 1644, we also define \( \tilde{\xi}_{cr} \equiv L_{cr}/L_{[0.3,10 \text{ keV}]} \).

High-energy neutrino spectra can be modified when mesons and muons cool before they decay. In this work, we semi-analytically take into account the meson suppression factor, which can be approximated to be

\[
 f_{\text{sup}}(\varepsilon_\pi) = \frac{t_{\text{dec} - \pi}^{-1}}{t_{\text{dec} - \pi}^{-1} + t_{\text{syn} - \pi}^{-1} + t_{\text{IC} - \pi}^{-1} + t_{\text{ad}}^{-1}},
\]

where \( t_{\text{dec} - \pi} \) is the pion lifetime, \( t_{\text{IC} - \pi} \) is the pion inverse-Compton cooling time, and \( t_{\text{syn}} \) is the pion synchrotron cooling timescale. For example, \( t_{\text{syn} - \pi} \) is given by

\[
 t_{\text{syn} - \pi}^{-1}(\varepsilon_\pi) = \frac{4}{3} \sigma_T \frac{m_e^2}{m_\pi^2} \frac{U_B}{c^3} \varepsilon_\pi,
\]

where \( \varepsilon_\pi \approx 0.2 \varepsilon_p \) is the pion energy in the plasma comoving frame. Synchrotron cooling is important for pions with energy \( \varepsilon_\pi \gtrsim 16 \text{ PeV} \). [136] concluded that TDE jets contain a large amount of turbulent magnetic field energy. For high magnetic fields, the charged decay products of \( p\gamma \) interactions will cool before decaying if \( t_{\text{syn}} \approx t_{\text{dec}} \), resulting in a cooling break in the final neutrino spectrum. Our results on neutrino fluences are shown in Figs. 5 and 6, and the results agree with analytical expectations based on the evaluation of \( f_{p\gamma} \).

### 3.4.2 Constraints on Neutrinos from Sw 1644 and Detection Prospects

Among the three cases, the X-ray bright TDE jet is of particular interest since we can expect a time and space-coincidence. The three jetted TDE candidates occurred during 2011 while the IC detector was fully operational with a total of 86 strings. Furthermore, two events including Sw1644 located at \( z \simeq 0.354 \) occurred above Earth’s northern hemisphere, and would be seen by IC as up going track events. Using our model of the neutrino flux from visible
Figure 3.5. Muon neutrino fluence for our canonical values for an X-ray bright TDE at $z \approx 0.3$. The cosmic-ray energy input is $\mathcal{E}_{\text{cr}} = \xi_{\text{cr}} L_{\gamma} t_{\text{dur}} = 17 \times 10^{53}$ erg with $\xi_{\text{cr}} = 1$ (thick) and $\xi_{\text{cr}} = 3$ (thin). Note that in our model Sw 1644 is more energetic by a factor of 2, and would have been marginally detectable with $\xi_{\text{cr}} \gtrsim 20$. 
Figure 3.6. Muon neutrino fluence for our canonical values for an X-ray dark TDE with choked jets at $z = 0.3$. Neutrino contributions come from the internal shock emission region (IS) and the terminal shock region (TS), where the jet head stalls. The cosmic-ray energy input is $\mathcal{E}_{\text{cr}} = \epsilon_p L_{\text{dur}} = 10^{50.5}$ erg.
TDEs combined with the effective area of IC, we discuss constraints on the detectability of the jetted TDE candidate Sw 1644.

The neutrino fluence (per flavor) for a single TDE is:

\[ F_{\nu} \sim 0.5 \times 10^{-4} \text{ erg cm}^{-2} \left( \frac{\xi_{\text{cr}}}{40} \right) \left( \frac{\min[1, f_{\gamma}]}{0.5} \right) \times \left( \frac{L_{\gamma, pk} t_{\text{hs}}}{2 \times 10^{53} \text{ erg}} \right) \left( \frac{d_L}{2 \text{ Gpc}} \right)^{-2} \]  

(3.13)

Using the IceCube effective area, we estimate that a signal neutrino could be detected from a fluence of \( F_{\nu} \sim 10^{-4} \text{ erg cm}^{-2} \). For Sw 1644, a naive estimate implies that \( L_{\gamma, pk}^{0.3,10 \text{ keV}} t_{\text{high}} = 2 \times 10^{53} \text{ erg} \) and \( \xi_{\text{cr}} \gtrsim 80 \) lead to a marginally detectable event, given that \( n_{\text{bkg}} \approx 10^{-4} \) atmospheric background neutrino events are expected from the same region of the sky with \( \bar{\varepsilon}_{\nu} \gtrsim 10^4 \) p.e.u. (Here the neutrino counts are binned in proxy energy units (p.e.u) related to the total electromagnetic energy observed by IC optical sensors.) Our results are also consistent with [143]. Comparing the calculated number with the number observed in the IC sample of up going muon neutrino events allows us to place upper limits on \( \xi_{\text{cr}} \). More quantitatively, one needs to take into account the effect of neutrino attenuation in the Earth. Also, the neutrino spectrum of Sw 1644 is harder than a simple \( E^{-2} \) spectrum so that the expected number of muon events is smaller.

From that fact that no neutrino events were observed from the relevant part of the sky, we can set a limit of \( \xi_{\text{cr}} \approx 100 - 250 \) (or \( \xi_{\text{cr}} \approx 20 - 50 \)) With this upper limit, our model predicts that Sw 1644 would have produced a neutrino detection if it was located at redshift \( z \sim 0.2 \) for \( \xi_{\text{cr}} \sim 10 \). The local rate of X-ray bright TDEs with peak luminosity \( L_{\gamma, pk} \gtrsim 10^{48} \text{ erg/s} \) and \( L_{\gamma, pk} \gtrsim 10^{47} \text{ erg/s} \) are \( \rho_0 \sim 0.03 \text{ Gpc}^{-3} \text{ yr}^{-1} \) and \( \rho_0 \sim 0.3 \text{ Gpc}^{-3} \text{ yr}^{-1} \), respectively. This implies that a TDE with that luminosity or higher occurs within \( z = 0.1 \) once every 10-100 years. Stacking searches can be more powerful. With X-ray sky monitors with ultimate sensitivities, which can detect TDEs up to \( z \sim 1 \), the detection rate by IceCube could be improved to \( \sim 0.1 - 1 \text{ yr}^{-1} \).

### 3.5 Diffuse Neutrino Intensity

We calculate the diffuse flux from jetted TDEs. The Eddington luminosity of a \( 10^6 M_\odot \) SMBH is \( L_{\text{Edd,BH}} \sim 10^{44} M_{\text{BH,6}} \text{ erg s}^{-1} \). The corresponding isotropic equivalent jet luminosity is \( L \sim 2 \times 10^{46} \theta_{j,1}^2 \text{ erg s}^{-1} \). One can calculate the diffuse neutrino flux from
“visible” (i.e., not choked) jetted TDEs (e.g., Sw 1644) by integrating luminosities for $10^{44.5} \text{ erg s}^{-1} < L_\gamma < 10^{49.5} \text{ erg s}^{-1}$. In this work, we allow $L_\gamma < 10^{46} \text{ erg s}^{-1}$ to consider a maximum contribution, but this does not change our results for the luminosity function we use. We consider choked jet contributions by extrapolating the luminosity function down to $L_\gamma < 10^{44.5} \text{ erg s}^{-1}$, which can be choked by a $0.5 M_\odot$ material with a radius of $r_{\text{out}} \sim 3 \times 10^{15} \text{ cm}$ (See Sec. 3.3.1).

One of the important predictions for TDEs is that they typically have a negative or weak evolution. We use the following TDE redshift evolution function [167],

$$f_{\text{TDE}}(z) = \left[ (1 + z)^{0.2\eta} + \left( \frac{1 + z}{1.43} \right)^{-3.2\eta} + \left( \frac{1 + z}{2.66} \right)^{-7.0\eta} \right]^{1/\eta}$$

(3.14)

with $\eta = -2$. Note that this redshift distribution peaks for $z \approx 0$, whereas the usual star-formation rate and GRB rate tend to peak around $z \sim 1 - 3$, leading to larger values of $\xi_z$.

The individual neutrino fluxes are then integrated over redshift and isotropic equivalent luminosity [74]

$$\Phi_\nu = \frac{c}{4\pi H_0} \int_{z_{\text{min}}}^{z_{\text{max}}} dz \int_{L_{\gamma_{\text{min}}}}^{L_{\gamma_{\text{max}}}} dL_\gamma \frac{d\rho_{\text{TDE}}(z, L_\gamma)/dL_\gamma}{\sqrt{\Omega_M (1 + z)^3 + \Omega_\Lambda}} \frac{L_{\nu}^* (L_\gamma)}{E_{\nu}^* t_{\text{eng}}},$$

(3.15)

where $\frac{d\rho_{\text{TDE}}}{dL_\gamma} = \rho_0 f(z) \Lambda_{\text{TDE}}(L_\gamma)$. Although observational uncertainties are large at present, both conventional (i.e. non-jetted) and Swift X-ray TDEs appear to share a luminosity function over seven orders of magnitude in observed peak luminosity with

$$\Lambda_{\text{TDE}}(L_\gamma) \propto \left( \frac{L_{\gamma,\text{pk}}}{L_{m,\text{pk}}} \right)^{-\alpha},$$

(3.16)

$L_{m,\text{pk}} = 10^{48} \text{ erg s}^{-1}$, and $\alpha = 2.0 \pm 0.05$ [167]. Eq. 3.16 is derived using the peak X-ray luminosity of an event, whereas we take $L_\gamma$ to be the average luminosity over the initial $\sim 10^6 \text{ s}$ of a jetted TDE. Again using observations of SW 1644, we assume the rough relationship $L_{\gamma,\text{pk}} \approx 12 L_\gamma$ (compare the maximum and median X-ray luminosity of SW 1644 in Supplementary Table 7 of [136]).

Using $t_{\nu\gamma}(\varepsilon_p, L_{\gamma,\text{pk}})$ (see Eq. 3.7), we integrate Eq. 3.15 and show the results of the
Figure 3.7. Contributions to the diffuse neutrino background due to $p\gamma$ interactions from X-ray bright visible jets and possible choked jets. For successful jets leading to X-ray bright TDEs, the cosmic-ray luminosity is given by $\xi_{cr} = L_{cr}/L_\gamma = 1 - 3$. For choked jets, the internal shock and termination shock scenarios are considered, and the cosmic-ray luminosity is assumed to be comparable to the total luminosity $L$. For the diffuse neutrino data, we use the muon neutrino data obtained by IceCube with the multiplication by a factor of 3 [54].
diffuse neutrino flux from visible and choked-jet TDEs in Fig. 3.7. The results can be understood by virtue of analytical estimates. The all flavor diffuse flux of neutrinos from extragalactically distributed sources can be estimated from the amount of cosmic-ray energy released per burst $\mathcal{E}_{\text{cr}} = \xi_{\text{cr}} \mathcal{E}_{\gamma}$ and the local rate of such explosions $\rho_0$ as

$$E_\nu^2 \Phi_\nu \approx \frac{3}{8} \frac{c}{4\pi} t_H \xi_z \min[1, f_{p\gamma}] f_{\text{sup}} f_{\text{cho}} \frac{\mathcal{E}_{\text{cr}}}{C} \rho_0$$

$$\sim 1 \times 10^{-9} \text{ GeV cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \min[1, f_{p\gamma}] f_{\text{sup}}$$

$$\times f_{\text{cho}} \xi_{\text{cr},0.5} \mathcal{E}_{\gamma,53.8} (\rho_0/0.1 \text{ Gpc}^{-3} \text{ yr}^{-1}) (\xi_z/0.5) \quad (3.17)$$

where $f_{\text{cho}} (\geq 1)$ is the enhancement factor (for energy budgets) due to the existence of choked jets [75]. We assume the Hubble time $t_H \sim 13.7$ Gyrs, a red-shift correction factor $\xi_z < \sim 0.6$ since most TDE are located at low redshift (See Eq. 3.14). The amount of cosmic-ray energy released from a TDE, approximated by $\varepsilon_p Q_{\varepsilon_p} \sim t_{\text{eng}} L_{\text{cr}}/C$. The factor $C = \ln(\varepsilon_{p,\text{max}}/\varepsilon_{p,\text{min}}) \sim 20$ converts the total bolometric cosmic-ray energy into the cosmic-ray energy per logarithmic energy interval which is generally what is expressed for the diffuse neutrino flux. The rate of TDEs is expressed in terms of the local TDE rate $\rho_0 = 0.03^{+0.04}_{-0.02} \text{ Gpc}^{-3} \text{ yr}^{-1}$ for events with peak luminosity greater than $L_m = 10^{48} \text{ erg/s}$, and a luminosity function $\rho_0 f(z) \Lambda_{\text{TDE}}(L_\gamma) \equiv \frac{\rho_{\text{TDE}}}{dL_\gamma} \propto L_{\gamma}^{-\alpha}$ (See Eq. 3.16). As a result, the TDE rates above $L_\gamma \sim 10^{46} \text{ erg/s}$ and $L_\gamma \sim 10^{46.5} \text{ erg/s}$ are $\rho_0 \sim 0.3 \text{ Gpc}^{-3} \text{ yr}^{-1}$ and $\rho_0 \sim 0.1 \text{ Gpc}^{-3} \text{ yr}^{-1}$, respectively.

For our luminosity function with $\alpha = 2$, the energy generation rate of visible jets and choked jets are similar. Thus, the maximum enhancement factor is $f_{\text{cho}}$ cannot be as large as $\sim 10 - 1000$. Thus, the neutrino flux from choked jets is comparable to that of visible jets and IceCube's neutrino flux cannot be explained. However, the luminosity function is uncertain. If the luminosity function is significantly steeper than the one we use and almost all TDEs have choked jets, $f_{\text{cho}}$ could be as large as $\sim 100 - 1000$. With such a bit extreme assumption, it is possible for them to significantly contribute the diffuse neutrino flux.

The calculated fluxes are significant, but below the observed diffuse flux of IC neutrinos [54], even with optimistic parameters. We have other constraints from multiplet searches. Using the absence of significant clustering in the 6 or 7 year data of IceCube, [151] obtained

$$n_0^{\text{eff}} \sim 1.1 \times 10^{-7} \text{ Mpc}^{-3} q_L^2 \left( \frac{\xi_z}{3} \right)^{-3} F_{\text{lim},-9}^{-3} \left( \frac{\Delta \Omega}{2\pi} \right)^2. \quad (3.18)$$
By replacing $n_0$ with $\rho_0 t_{\text{dur}}^3/T_{IC}^2$, the above result is readily rewritten as

$$\rho_0^{\text{eff}} \gtrsim 1.7 \times 10^3 \text{Gpc}^{-3} \text{yr}^{-1} \frac{q_L^2 (\Delta \Omega/2\pi)^2 (T_{IC}/6 \text{ yr})^2}{\xi_3^3 F_{\text{lim}, -6.9}^3 d_{\text{dur}, 6}^3}.$$ (3.19)

For TDEs with $t_{\text{eng}} \sim 10^6$ s, we have used the sensitivity $F_{\text{lim}} \sim 10^{-7}$ GeV cm$^{-2}$ s$^{-1}$. Note that the above limit is weaker by a factor for hard spectra [151], but we can conclude that it is very unlikely that X-ray bright TDEs are the sources of IceCube’s neutrinos. X-ray bright TDEs with the local rate of visible TDEs $\rho_0 \sim 0.1$ Gpc$^{-3}$ yr$^{-1}$ can give $\lesssim 5 - 10\%$ of the diffuse neutrino flux, which is consistent with our diffuse flux calculations. In other words, $\xi_{\text{cr}} \gtrsim 100$ is unlikely. However, in principle, X-ray dim TDEs with choked jets can avoid this constraint by achieving $f_{\text{cho}} \gg 1$ §. The luminosity function should be different from that obtained by [167]; that is, $\alpha \approx 3.2 - 4.0$ is required.

X-ray bright TDEs are also expected to be $\gamma$-ray dim. While GeV-TeV $\gamma$-rays are expected from $p\gamma$ interactions, they are attenuated by the X-rays that are the target photons for cosmic-ray interactions. The $\gamma$-ray optical depth in X-ray bright sources can be related to the fraction of cosmic-rays that undergo $p\gamma$ interactions, $\tau_{\gamma\gamma} \sim \sigma_{\gamma\gamma}/(\kappa_p \sigma_{\gamma p}) f_{p\gamma} \sim 1000 f_{p\gamma}$ [50]. From Fig. 4.4 we see that $f_{p\gamma} \sim 0.1$ for $\epsilon_p \gtrsim 10$ TeV. Since $\epsilon/\epsilon_p \sim m_e^2 c^2/(0.15 \text{ GeV} m_p)$, X-ray bright TDEs are opaque to $\gamma$-rays with energies $E_\gamma \gtrsim 10 - 100$ MeV, which is consistent with [136].

### 3.6 Conclusion and Discussion

We studied different possibilities of high-energy neutrino production inside TDE jets. Sw 1644 is the closest and brightest jetted TDE candidate, and the non-detection would give us a limit on the cosmic-ray loading factor, $\xi_{\text{cr}} \approx 20 - 50$. If Sw 1644 had occurred within $z \sim 0.1 - 0.2$, it would have produced a definitive neutrino detection. Since such a bright TDE only occurs once every 10-100 years within that distance, the association of X-ray bright TDEs with neutrinos will require aggregate neutrino coincidence searches such as stacking analyses. If we have X-ray all-sky monitors with ultimate sensitivities allowing us to see TDEs up to $z \sim 1$, IceCube may detect neutrino signals from TDEs, with $0.1 - 1$ yr$^{-1}$. But new X-ray satellites with a wide field of view and a sensitivity of $F_{\text{lim}} \sim 5 \times 10^{-11}$ erg cm$^{-2}$ s$^{-1} L_{\gamma, pk, 48}(d_L/6.6 \text{ Gpc})^{-2}$ would be necessary for efficient

---

§For choked ultra-long (or low-luminosity) GRB jets with $\alpha = 2.3$, $f_{\text{cho}} \gtrsim 2 - 3$ (or $f_{\text{cho}} \gtrsim 1$) can avoid constraints thanks to larger $\xi_3$ as long as cosmic-rays can be accelerated.
searches of high-energy neutrinos from jetted TDEs. With IceCube-Gen2 [173], our results imply that it is possible to detect the brightest TDE events such as Sw 1644.

The method we use to determine the limit on the neutrino brightness of Sw 1644 can be applied to other non-\(\gamma\)-ray bright sources (e.g., Type Ibc SNe), assuming that a consistent model of the broadband electromagnetic and neutrino emission is formulated. It is also important to determine the duration of the neutrino emission, and any time delays between photon and neutrino arrival times since this affects the background rate of IC neutrinos. Understanding the size of the signal time interval also affects the energy limit in the search. For example, using one year of IC data required a cut at \(\tilde{\varepsilon}_{\nu\mu} \gtrsim 10^4\) p.e.u, while a time window of \(10^6\) s around the triggering time of Sw 1644 would lower the atmospheric background rate of neutrinos and lead to a lower energy threshold of \(\tilde{\varepsilon}_{\nu\mu} \gtrsim 10^3\) p.e.u. Future searches of \(\gamma\)-ray dim transient sources (i.e., sources other than GRBs and Blazar flares) will require an optimization of the threshold energy and time interval search size.

We also found that the flux and spectral index of the diffuse neutrino flux from electromagnetically bright TDEs are typically sub-dominant as an origin of IceCube’s neutrino flux, but could be significant at very high energies (\(E_{\nu} \gtrsim 1\) PeV). As in \(\gamma\)-ray bursts [36] and X-ray flares in GRB afterglows [172], neutrinos flares associated with TDEs, which may have \(10^{-9}\) GeV cm\(^{-2}\) s\(^{-1}\) sr\(^{-1}\), are interesting targets for next-generation neutrino telescopes that are suitable for extremely high-energies, such as IceCube-Gen2 [173], Askaryan Radio Array [174], ExaVolt Antenna [175], and Giant Radio Array for Neutrino Detection [176].

We discuss possible contributions made by TDEs with choked jets. Although this scenario is still speculative, we found that the magnitude of the diffuse neutrino flux expected from choked jet TDEs is typically low to be compatible with IC measurements. However, we caution that our estimate is based on the large extrapolation of the luminosity function. The luminosity function obtained by [167], \(\Lambda_{\text{TDE}}(L_\gamma) \propto L_\gamma^{-2}\), all TDEs with different luminosities contribute approximately the same amount to the diffuse cosmic-ray flux. The neutrino flux is dominated by TDEs with \(f_{p\gamma} \approx 1\). The situation is different from other classes of choked jet objects (e.g., [48,75] for low-power GRBs) that have been found to be more efficient neutrino producers. For example, low-luminosity GRBs have \(\Lambda_{\text{LLGRB}} \propto L_\gamma^{-2.3 \pm 0.2}\) so that the lowest-luminosity events largely contribute to the diffuse neutrino flux. Because of large uncertainty, it would be possible to assume steeper values of \(\alpha\), and then one could increase the contributions from choked jets. A steep luminosity
function leading to $f_{\text{cho}} \sim 100 - 1000$ would be required to explain the IceCube data. Besides this issue, the setup expressed by Eq. (3.2) should be tested by observations. For example, the radio emission is useful to probe densities of the circumnuclear material around a SMBH [177]. Also, how sub-Eddington luminosity jets are launched by spinning BHs should be justified.
Chapter 4  
Choked Jets and Low-Luminosity Gamma-Ray Bursts as Hidden Neutrino Sources

4.1 Introduction

A diffuse flux of very high-energy (VHE) neutrinos with energies $10 \text{ TeV} \lesssim E_\nu \lesssim 2 \text{ PeV}$ has been reported from the Antarctic neutrino detector IceCube [18,24,25,52,53,159,178,179]. The sources of these neutrinos are currently unknown, but they appear to be extragalactic in origin [62,63,180]. Gamma-Ray Bursts (GRBs), energetic supernovae (SNe), active galactic nuclei (AGNs), starburst and star-forming galaxies, galaxy groups and clusters, as well as some Galactic sources have been proposed as potential candidates (see reviews [181–183]). In particular, GRBs, which are believed to be caused by ultra-relativistic jets launched by the collapse of a massive star (i.e. collapsars or long GRBs) or the merger of two compact binary objects (e.g. [184] for review) have been investigated as potential sources of ultrahigh-energy cosmic rays (CRs) and secondary VHE neutrinos for more than a decade [39]. Stacking analyses [46,185] lead to the conclusion that $\lesssim 1\%$ of the measured diffuse flux can be explained as prompt emission from observed high-luminosity (HL) GRBs [75,76,120]. However, low-power GRBs∗, including low-luminosity GRBs (LL GRBs) and ultralong GRBs (UL GRBs), may be largely missed by current GRB satellites such as Swift and Fermi. LL GRBs may be more common than classical

∗“Low-power GRBs” here are introduced based on the observed luminosity. Even if the observed luminosity of LL GRBs is $L_\gamma \sim 10^{47}$ erg s$^{-1}$, which is much lower, the intrinsic isotropic-equivalent luminosity of choked jets may be as high as that of classical GRBs.
GRBs [167, 186, 187] and UL GRBs may also be as common as high-luminosity GRBs [188]. In addition, choked jets that do not escape their progenitor star in the collapsar scenario – so called failed GRBs [169] – have been suggested as possible sources for the observed diffuse neutrino flux [50, 75, 189]. The present stacking limits on prompt neutrino emission from classical high-luminosity GRBs are not applicable to such low-power or dark GRBs, and these may give a significant contribution to the diffuse neutrino flux [190, 191]. For the first time, we study the expected neutrino contribution from jets that successfully escape from their progenitor star, but are subsequently smothered by a dense, optically thick external material resulting in a LL GRB or a failed GRB. This scenario, described by Nakar [192], allows for a unified picture for HL, LL, and dark GRBs which have similar intrinsic progenitor and jet properties, but different circumstellar environments (e.g. the presence or absence of a dense wind or outer envelope). It is conjectured that LL GRBs may occur if the relativistic jet becomes smothered by the extended wind/outer envelope, and acts as a piston driving a quasi-spherical shock into the circumstellar material. The GRB event occurs when this shock breaks out in the optically-thin region.

Another possible subclass of interest are UL GRBs, which have a much longer duration compared to classical GRBs (but see also Ref. [193]). Their long duration may suggest a long-lasting fall-back accretion from an extended progenitor onto a black hole. Blue supergiants (BSGs) are possible UL GRB progenitors, and are believed to be common at very high redshifts [194, 195]. Alternatively, such long durations may be explained by a fast-rotating pulsar, which could account for the connection between UL GRBs, super-luminous SNe and hypernovae (e.g., [196–198]). Although we do not consider potential sources of UL GRBs in this work, these low-power GRBs can also contribute to neutrino emission [75].

Predictions for high-energy neutrino emission from GRB jets of both high and low luminosity are still uncertain despite recent improvements in theoretical calculations (e.g., [36, 37, 171, 199–202] although guaranteed emission is expected in the GeV-TeV range for neutron-loaded outflows e.g., [203–206]). Irrespective of their viability as VHE neutrino factories, the mechanisms for producing, and the physical processes associated with low-power GRBs are still largely unknown and remain intriguing open questions. Nearby long GRBs have been associated with broad-line Type Ic supernovae (SNe) (e.g., GRB 980425, 060218, and 100316D), which are known to be caused by the collapse of massive stars that eject of their outer envelopes. LL GRBs have been of special interest since they show intermediate properties between GRBs and SNe, and have been associated
with transrelativistic SNe [207]. Both types of transients may be driven by jets [192, 208] and the study of LL GRBs may offer clues to the GRB-SN connection [209, 210].

In this work, based on the above motivation we consider the VHE neutrino emission from jets choked by dense external material, as well as any subsequent shocks resulting from the jet acting as a relativistic piston. In particular, we focus on scenarios which may produce LL GRBs. Under the current constraints imposed by the IceCube analyses mentioned above, such LL GRBs are attractive as the originators of the diffuse VHE neutrino flux (i) because of their high local rate relative to their high-luminosity cousins, and (ii) because their low gamma-ray flux make them difficult to detect with conventional electromagnetic detectors (e.g. Swift). Recently, Murase & Ioka [75] showed that choked jets may be more favorable as sites of efficient neutrino production. Jets which successfully penetrate both the progenitor star – and if applicable a circumstellar envelope – (i.e. emergent jets) typically have high-luminosities such that they form radiation-mediated shocks, which are unfavorable for CR acceleration and neutrino production. Taking into account the luminosity and redshift distribution of LL GRBs, we show that they and the choked jets may contribute to the diffuse neutrino flux while remaining absent from GRB joint electromagnetic-neutrino searches. We also explicitly show the conditions required to produce choked jets with radiation-unmediated shocks.

4.2 Dynamics of Relativistic Jets

4.2.1 Model Setup for Emergent Jet, Shock Breakout, and Choked Jet Scenarios

GRBs are thought to result from the intense emission from relativistic jets that successfully penetrate a progenitor star, and an understanding of jet propagation is undoubtedly relevant (e.g., [165, 166, 169]). It would be natural to expect that the radiation mechanism of LL GRB gamma-ray emission is similar to that of classical GRBs [208, 211, 212]. The simplest such model is a scaled-down version of the classical GRB, where dissipation occurs in a mildly relativistic jet which has emerged outside of the progenitor star and any circumstellar material. We call this scenario the emerging jet (EJ) model (see Fig. 4.3 right panel). For EJs, prompt neutrino emission is produced together with prompt gamma-ray emission outside the star, identical to the scenario expected from classical GRBs [190, 191, 213].
Figure 4.1. The choked jet model for jet-driven SNe. Orphan neutrinos are expected since electromagnetic emission from the jet is hidden, and such objects may be observed as hypernovae.
Figure 4.2. The shock breakout model for LL GRBs, where transrelativistic shocks are driven by choked jets. A precursor neutrino signal is expected since the gamma-ray emission from the shock breakout occurs significantly after the jet stalls (e.g., [169]).
Figure 4.3. The emerging jet model for GRBs and LL GRBs. Both neutrinos and gamma-rays are produced by the successful jet, and both messengers can be observed as prompt emission.
Another interpretation of LL GRBs which has received attention is the shock breakout emission model, where the prompt emission is attributed to dissipation caused by a transrelativistic, aspherical shock in a dense wind [214–217]. The origin of the relativistic velocity components in the ejecta is an issue. One of the promising possibilities is that the fast shock is driven by a choked jet. The jet stalls close enough to the photosphere so that a transrelativistic shock breaks out through the star and its extended material. We call this model the choked jet-shock breakout (CJ-SB) model, and the middle panel of Fig. 4.2 shows its schematic.

A luminous jet naturally leads to an easier break out from a typical compact progenitor such as a Wolf-Rayet (WR) star. One possible cause for an energetic jet to stall is through smothering by an extended envelope of wind material. Such an environment need be only marginally more massive than what is inferred, e.g. by [214,215]. For an extended material mass \( M_{\text{ext}} \sim 10^{-3} - 10^{-2} M_\odot \), a typical GRB jet penetrating the compact star is expected to be choked at a distance of \( \sim 10^{12} - 10^{13} \) cm from the central engine [192]. Although such a radius is 10 times larger than that of WR stars believed to be GRB progenitors, it is appealing that the model can also explain the mysterious UV component in GRB 060218. Envelopes of \( M_{\text{ext}} \sim 10^{-7} M_\odot \) are observed for pre-exploded WRs, but there is increasing evidence for exceptionally high mass loss in the weeks prior to collapse [88,218]. Accompanying theoretical explanations have been proposed [219–221].

Choked jets with a shock breakout component may or may not produce prompt gamma-rays (e.g., [215,222]). The jet may stall sufficiently far below the photosphere that the piston action of the jet does not lead to transrelativistic velocity components of the shock, i.e., such objects are simply observed as energetic SNe or hypernovae in the optical band with no accompanying gamma-ray emission. We call this final scenario the choked jet (CJ) model. That is, the SN has a “normal” quasispherical shock, without a jet-driven (aspherical, transrelativistic) extra shock breakout component (see Fig. 4.1 left panel where we have neglected to show the subrelativistic shock for clarity).

Thus, it is possible to have a unified picture where the GRB-SN connection is explained by the strength of the choked jets. Both of the CJ and CJ-SB models provide favorable environments for neutrino production since shock acceleration may occur inside the jet before the shocks become radiation mediated, allowing for efficient neutrino emission. On the other hand, the gamma-rays produced deep inside the choked jets will not be able to escape through the extended material to observers, and bright gamma-rays are observed from the shock breakout component only. A similar picture for neutrino production
is considered in Ref. [75] for low-power jets with \( L \lesssim 10^{47} \) erg s\(^{-1}\) and/or extended progenitors such as blue supergiants (BSGs).

### 4.2.2 Hydrodynamical Constraints on Choked Jets

We will consider VHE neutrino emission from choked jets and it is relevant to consider the condition where relativistic jets are stalled. The dynamics of a relativistic jet is determined by its interaction with the ambient medium in the progenitor star and circumstellar envelope, which can change the shape of the jet through collimation shocks [165,166]. While this work focuses on neutrino emission only from internal and termination shocks, it is important to note that collimation shocks near the base of the jet will also affect estimates of VHE neutrino production [75] but are generally not considered in most of the previous literature [152,170,223–225]. As the jet drills through the star, a contact discontinuity is formed between the shocked jet material and the shocked ambient matter. This region of shocked material is often referred to as the jet head. Balancing the jet’s internal pressure with the ram pressure of the ambient material determines the head’s dynamics (see, e.g., [165,166,169]), and the head velocity is given by

\[
\beta_h = \frac{\beta_j}{1 + \tilde{L}^{-1/2}},
\]

with dimensionless luminosity

\[
\tilde{L} \approx \frac{L_{0j}}{\pi (t_h \theta_0^2)^2 \rho_a c^3},
\]

where \( L_{0j} \) is the one-sided jet luminosity, \( \rho_a \) is the ambient density, \( r_h \) is the distance of the jet head from the central engine, and \( \theta_0 \) is the initial opening angle inside the star. Assuming the jet material is relativistic \( \beta_j \sim 1 \) before reaching the head it is obvious from Eq. (4.1) that for \( \tilde{L} \gg 1 \) the head moves relativistically as well. Then, the jet will be collimated for \( \tilde{L} \ll \theta_0^{-4/3} \) or uncollimated for \( \tilde{L} \gg \theta_0^{-4/3} \) [165].

First, let us consider a jet propagating inside its progenitor star. As shown in Refs. [165,166], such a jet is typically collimated. Let us assume that the density profile is approximated to be \( \rho_a = (3 - \alpha)M_*(r/R_*)^{-\alpha}/(4\pi R_*^3) \) (\( \alpha \sim 1.5 - 3 \)). Here \( M_* \) is the progenitor mass and \( R_* \sim 0.6 - 3R_\odot \). For WR progenitors, we may take \( \alpha = 2.5 \) [226], leading to the jet head radius \( r_h \simeq 5.4 \times 10^{10} \) cm \( t_1^{6/5}L_{0j,52}^{2/5}(\theta_0/0.2)^{-4/5}(M_*/20 M_\odot)^{-2/5} R_*^{1/5} \), where \( L_0 = 4L_{0j}/\theta_0^2 \) is the isotropic-equivalent total jet luminosity [75,165]. The
classical GRB jet is typically successful (i.e., it emerges from the progenitor), since the time required for the jet to escape the progenitor \( t_{\text{jbo}} \approx 17 \, s \, L_{0.5}^{-1/3} (\theta_0/0.2)^{2/3} (M_*/20 \, M_\odot)^{1/3} R_{*11}^{2/3} \) is shorter than the jet duration \( t_{\text{eng}} \sim 10^{1.5} \, s \). This time is in good agreement (i.e. within a factor of a few) with numerical studies of jet emergence [165,227–229]. See also Fig. 15 of Ref. [166].

Toma et al. [211] suggested that the prompt emission of GRB 060218 may come from an emerging jet with a Lorentz factor of \( \Gamma \sim 5 \), and this possibility of marginally successful jets has been further investigated by Irwin & Chevalier [208]. The jet has more difficulty in penetrating the progenitor star due to its lower luminosity, but on the other hand, its longer duration helps in achieving breakout. In this model, the prompt gamma-ray emission may come both from relatively low radii around the photosphere or large radii. Such marginally successful jets are expected for larger radius progenitors such as BSGs, and UL GRBs may correspond to the case of successful GRBs [75].

Next, we consider jets embedded in an extended, massive envelope. The jet can be choked if the mass of the extended material is sufficiently large. Motivated by the CJ-SB model for LL GRBs, we consider an extended material with mass \( M_{\text{ext}} \sim 10^{-2} \, M_\odot \) and radius \( r_{\text{ext}} \sim 3 \times 10^{13} \, \text{cm} \). WR stars have been observed with such unusually massive envelopes in the months leading up to their SN explosion [230–232]. Nakar [192] suggested similar envelope parameters for LL GRB 060218, but without strong constraints on the density profile. For simplicity, we therefore assume the same wind profile for all LL GRBs, namely

\[
\rho(r) = 5.0 \times 10^{-11} \, \text{g} \, \text{cm}^{-3} \left( \frac{M_{\text{ext}}}{0.01 \, M_\odot} \right) r_{\text{ext},13.5}^{-3} \left( \frac{r}{r_{\text{ext}}} \right)^{-2}, \tag{4.3}
\]

with the density at the outer envelope edge \( \rho_{\text{ext}} \equiv \rho(r_{\text{ext}})/(5.0 \times 10^{-11} \, \text{g} \, \text{cm}^{-3}) \). Assuming – as Nakar did – that the majority of the envelope’s mass is located near the outer radius (i.e. the quantity \( \rho(r) \, r^3 \) increases up until \( r_{\text{ext}} \)) different wind profiles do not significantly affect the dynamics of the jet head. The jet is typically uncollimated for sufficiently luminous jets and the Lorentz factor of the jet head is given by

\[
\Gamma_h \approx \frac{\tilde{L}^{1/4}}{\sqrt{2}} \simeq 3.5 \, L_{0.52}^{1/4} \rho_{\text{ext}}^{-1/4} r_{h,13.5}^{-1/2}, \tag{4.4}
\]
while the jet head radius is estimated to be

\[ r_h \approx 2 \Gamma_h^2 c t \approx 2.3 \times 10^{13} \text{ cm} \ L_{0.52}^{1/2} \rho_{\text{ext}}^{-1/2} r_{\text{ext,13.5}}^{-1} t_{1.5}. \]  

(4.5)

The condition \( r_h = r_{\text{ext}} \) gives the jet breakout time \( t_{\text{jbo,ext}} \), and the condition \( t_{\text{jbo,ext}} \lesssim t_{\text{eng}} \) gives the jet-stalling condition

\[ L_\gamma \lesssim L_\gamma^{\text{JS}} \approx 0.95 \times 10^{48} \text{ erg s}^{-1} \left( \frac{\xi_\gamma}{0.25} \right) \left( \frac{\theta_j}{0.2} \right)^2 t_{\text{eng,1.5}}^{-1} \times T_{3.5}^{-1} \rho_{\text{ext}} r_{\text{ext,13.5}}^4, \]  

(4.6)

where we have used

\[ L_\gamma \approx \xi_\gamma \frac{\theta_j^2}{2} L_\gamma t_{\text{eng}} \frac{t_{\text{eng}}}{T}. \]  

(4.7)

\( L_\gamma \) is the observed luminosity of the LL GRB, \( \xi_\gamma \) is the gamma-ray emission efficiency, \( \theta_j \) is the choked jet opening angle in the extended material, and \( T \) is the typical observed duration of LL GRBs. For choked jets, the jet head radius at \( t_{\text{eng}} \) is defined as the jet-stalling radius \( r_{\text{stall}} \). In our CJ-SB scenario, the central engine activity time \( t_{\text{eng}} \) is unrelated to the duration of the prompt emission \( T \), since the later only depends on the shock velocity and breakout radius \( T \approx r_{\text{sb}} / (\Gamma_{\text{sb}}^2 c) \), which are determined from the envelope properties. The former timescale is deduced from the lifetime of GRB jets that are seen in high-luminosity GRBs, while the latter reflects the typical observed duration of a LL GRB. Note that the hydrodynamic constraints are relevant for neutrino production. First, they restrict the emission radius, which limits the overall non-thermal particle energy density as well as the maximum neutrino energy. Since the emission region of traditional GRB jets assume a wide variety of values (e.g. \( 10^{11} \text{ cm} \lesssim r_{\text{em}} \lesssim 10^{17} \text{ cm} \) for the classic fireball model in [36]), the phenomenology of neutrinos from choked jets can be quite different. Additionally, the jet luminosity and central engine duration need to be consistently determined. If the jet is too powerful or its duration is too long, it is no longer choked and should be reduced to the classical GRB case.

If the jet is choked in the dense wind close to the edge of the star, it will launch a transrelativistic shock that becomes an aspherical shock breakout. As described in Refs. [233, 234], breakout nonthermal emission may be released when the optical depth of the shock reaches unity. The emission time of the breakout – and therefore the approximate duration of the subsequent GRB – is \( T \approx r_{\text{sb}} / (\Gamma_{\text{sb}}^2 c) \sim 10^{3.5} \text{ s} \) in agreement
with the average duration of LL GRBs, where \( r_{\text{sb}} \) is the shock breakout radius and \( \Gamma_{\text{sb}} \) is the Lorentz factor of the shock. It has been shown that shock breakouts produce smooth light curves similar to what are seen in LL GRBs.

### 4.3 Radiation Constraints on Shock Acceleration

CRs are generally assumed to be accelerated with a power-law distribution by the first-order Fermi process in the presence of shocks or turbulence. As known from the literature of nonrelativistic shocks, (e.g., [233,235]), efficient conversion of the fluid kinetic energy to a non-thermal particle population can occur if the shocks are collisionless (i.e. mediated by plasma instabilities), requiring the upstream plasma to be optically thin for relativistic shocks. CRs gain energy thanks to the shock compression. If the shock is mediated by radiation, efficient acceleration is prevented [75,236] since the shock width is larger than the CR Larmor radii and particles cannot efficiently cross between the upstream and downstream fluids. This subtle feature of CR acceleration in relativistic jets is often not considered in the literature [152,170,223–225].

The radiation constraints give us stringent restrictions on the rate of VHE neutrino production from choked jets. Two shells in the jet have a relative Lorentz factor \( \Gamma_{\text{rel}} \approx \Gamma_r/2\Gamma \), where the rapid shell is moving with velocity \( \Gamma_r \sim \text{few} \times \Gamma \). Murase & Ioka [75] derived radiation constraints. For internal shocks, efficient CR acceleration can occur if

\[
\tau_T = n_j' \sigma_T (r_{\text{is}}/\Gamma) \lesssim \min[\Gamma_{\text{rel}}^2, 0.1C^{-1}\Gamma_{\text{rel}}^3] \text{ or }
\]

\[
L_{52}r_{\text{is},10}^2\Gamma_2^{-3} \lesssim 5.7 \times 10^{-3} \min[\Gamma_{\text{rel},0.5}^2, 0.32C_1^{-1}\Gamma_{\text{rel},0.5}^3].
\]  

(4.8)

where \( \Gamma \sim 100 \) is the bulk Lorentz factor of the jet, \( L \) is the isotropic-equivalent kinetic luminosity of the jet, and the comoving density is \( n_j' \approx L/(4\pi\Gamma^2r_{\text{is}}^2m_pc^3) \). The factor \( C = 1 + 2\ln\Gamma_{\text{rel}}^2 \) accounts for possible pair production effects in the shocked material [217,237]. The jet propagating in the star is usually collimated, and the collimation shock radius \( r_{\text{cs}} \) is smaller than the jet head radius (\( r_{\text{cs}} < r_h < R_\star \)). We expect radiation-mediated shocks for the jet inside a WR star, but the shocks can be collisionless if the ambient material is more extended [75]. Setting the shock radius equal the stall radius calculated in the previous section (\( r_{\text{is}} = r_{\text{stall}} \)) and ignoring the prefactor depending on details of the shock structure, we obtain the following upper limit for the
luminosity based on radiation constraints,

\[ L \lesssim 1.7 \times 10^{54} \text{ erg s}^{-1} \Gamma_{2}, \rho_{\text{ext}}^{-1} r_{\text{ext},13.5}^{-2}. \]  

(4.9)

In the CJ-SB model, using the observed gamma-ray luminosity, the constraint becomes

\[ L_{\gamma} \lesssim L_{\gamma}^{\text{RC}} \approx 8.6 \times 10^{49} \text{ erg s}^{-1} \left( \frac{\xi_{B}}{0.25} \right) \left( \frac{\theta_{j}}{0.2} \right)^{2} \times \Gamma_{2}^{-2/3} \rho_{\text{ext}}^{-1} r_{\text{ext},13.5}^{-2}. \]  

(4.10)

As the luminosity increases or the radius decreases, the optical depth largely exceeds unity and CRs cannot be accelerated efficiently at the shock. While we assume the plasma is optically thin inside the jet, the envelope or circumstellar medium outside is largely optically thick to Thomson scattering. Therefore, photons are free to move inside the jet but cannot escape. The free streaming of photons from outside the jet core – specifically thermal emission from the jet head – allow for efficient neutrino production through \( p\gamma \) interactions.

The radiation constraints apply to shocks in the envelope material as well as those in the choked jet. Before breakout, the shock is radiation mediated. As photons diffuse out from the system, the shock becomes collisionless and CRs may be accelerated [233, 234]. The condition is given by \( \tau_{T} \lesssim \beta_{\text{SH}}^{-1} \), where \( \beta_{\text{SH}} = V_{\text{SH}}/c \) is the shock velocity.

### 4.4 Neutrino Production in Choked Jets

In the present work, we focus on the CJ and CJ-SB models, so we assume that CRs are accelerated in jets choked by the circumstellar material. For simplicity we assume that internal shocks occur near the collimation shock radius since they are expected to occur most frequently at \( r_{\text{is}} \approx 2\Gamma_{2}^{2} c \delta t \) when the jet is not stalled, where \( \delta t \) is the variability timescale. We calculate neutrino spectra only for GRB jets satisfying the following condition,

\[ L_{\gamma} \lesssim \min[L_{\gamma}^{\text{JS}}, L_{\gamma}^{\text{RC}}], \]  

(4.11)

where \( L_{\gamma}^{\text{JS}} \) comes from the jet-stalling condition (Eq. 4.6) and \( L_{\gamma}^{\text{RC}} \) comes from the radiation constraint (Eq. 4.10). The neutrino spectra are calculated numerically, taking into account various microphysical processes such as multipion production. For details
of the method, see Refs. [75, 171, 199, 238]. We explicitly calculate the CJ component of the neutrino flux from jets with luminosities $L_\gamma = 10^{45}$ erg s$^{-1}$, $L_\gamma = 10^{46}$ erg s$^{-1}$, $L_\gamma = 10^{47}$ erg s$^{-1}$, and $L_\gamma = 10^{48}$ erg s$^{-1}$ and use the results to infer the approximate contribution from CJs with other luminosities in Eq. 4.14 below. For the GRB jet parameters, assuming that the jets leaving the star are similar to those of classical GRBs, we use $t_{\text{eng}} = 10^{1.5}$ s, $\theta_j = 0.2$, and $\Gamma = 100$. In this work, we calculate the neutrino emission for different luminosities but fix the other parameters because of computational limitations. Note that our GRB jet parameters have been widely used in the GRB literature as typical values. Also, thanks to the high neutrino production efficiency, the flux level is insensitive to the GRB jet parameters [75, 199]. The resulting neutrino flux effectively scales with the CR loading parameter. Note that lower Lorentz factors further restrict the parameter space of the VHE neutrino production since the shock becomes radiation mediated. For slow jets with $\Gamma \lesssim 10$, VHE neutrino production inside a typical GRB progenitor requires $L \lesssim 10^{47}$ erg s$^{-1}$ [75]. In addition, we assume $M_{\text{ext}} = 0.01 M_\odot$ and $r_{\text{ext}} = 10^{13.5}$ cm for the extended material, which are based on the parameters suggested for the explanation of LL GRBs.

CRs accelerated at internal shocks interact with photons produced from both internal shocks and the jet head. The properties of electrons accelerated at internal shocks are uncertain so we conservatively consider only thermal photons coming from the jet head [169]. In rest frame of the head the photon temperature is $kT_h \simeq 0.37$ keV $L_h^{1/4} r_{h,13}^{-1/2} \Gamma_h^{-1/4}$. This temperature is Lorentz boosted by a factor $\Gamma_{\text{rel}} \sim \Gamma / 2 \Gamma_h$ in the jet rest frame. Likewise, the photon energy density seen in the jet is $U_{\gamma,j} \sim \Gamma_{\text{rel}}^2 U_{\gamma,h}$, which can make the thermal component from the head the most significant photon field. We also take into account the photon escape fraction $\sim (n_{\gamma,h} \sigma_T r_h / \Gamma_h)^{-1}$. The photomeson production efficiency satisfies

$$\min[1, f_{\text{p}}] \approx 1$$

(4.12)

Thus, CRs that exceed the pion production threshold are depleted by the photomeson production. Choked jets can be regarded as “calorimetric” sources in the sense that all of the available CR energy goes into making neutrinos and the observation of neutrinos allow us to directly probe the amount of accelerated CRs. Note that, although there are nonthermal populations of photons radiated by co-accelerated pairs, this point is unchanged. Additional photons enhance the efficiency of the photomeson production.

At subphotospheric radii, inelastic $pp$ interactions are shown to be relevant below 100 TeV and the photon meson efficiency is found to dip due to the Bethe-Heitler process

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When radiation constraints are satisfied, the \( pp \) optical depth during the dynamical time in the jet is limited to \([199,234]\)

\[
f_{pp} \lesssim \frac{N_{pp} \sigma_{pp}}{\sigma_T} \approx 0.04.
\]

Thus, for the CR spectrum with \( s \sim 2 \), the energy flux of the \( pp \) component is typically lower than that of the associated \( p\gamma \) component, and the main production mechanism for VHE neutrinos inside the jet is the photomeson production process. Note that, as pointed out in Ref. [75], low-energy CRs in the jet may eventually be advected along the collimated jet and all CRs can be depleted for neutrino production via subsequent \( pp \) interactions. However, for uncollimated shocks, CRs accelerated at the shocks may simply cool via adiabatic losses during the jet expansion. In this work, to be conservative, we do not consider effects of the remaining CRs.

CRs lose their energy via photomeson production (\( p\gamma \)), Bethe-Heitler pair production (BH), hadronuclear (\( pp \)), synchrotron radiation (syn) and inverse-Compton radiation (IC) processes. Since in our case the jet is expanding, adiabatic losses are also included. The numerically calculated acceleration and cooling timescales are shown in Fig. 4.4. At higher energies, photomeson production is the dominant cooling channel for CRs, and sets a maximum CR energy of \( \varepsilon_M \sim 3 \) PeV in the comoving frame. At lower energies \( \varepsilon_p \lesssim 1 \) TeV, BH cooling dominates and suppresses the neutrino spectrum around 100 GeV. The magnetic field is set by using the parameter \( \xi_B = L_B/L_0 = 0.1 \), and the acceleration time is given by \( t_{\text{acc}} = \varepsilon_p/(eBc) \). Note that, even if the initial maximum CR energy is very high, the energy of escaping CRs is low since they are depleted advecting to the shock downstream.

The neutrino products receive \( \sim 5\% \) of their parent CR energy. For the energies of interest, \( E_\nu \sim 30 \) TeV, the \( \Delta \) resonance in \( p\gamma \) interactions requires target photons with energy \( \sim 0.2 \) keV – \( 0.2 \) MeV depending on the Lorentz factor of the target photon field (i.e. between the jet interior and head) [50]. The jet head appears as a blackbody with temperature \( kT_h \sim 0.1 – 1 \) keV as seen from the jet core providing photons within the correct energy range. The fraction of CRs that interact with these photons using the box approximation of Ref. [39] is \( f_{p\gamma} \gg 1 \), which is also seen from Fig. 4.4. As expected, \( p\gamma \) interactions dominate the neutrino production.

In addition, meson cooling is also taken into account by solving the kinetic equations numerically [75,171,238].
Figure 4.4. Various energy-loss rates of CR protons in the CJ model for LL GRBs as a function of the comoving proton energy $\varepsilon$. Photomeson production ($p\gamma$), Bethe-Heitler pair production (BH), hadronuclear ($pp$), synchrotron radiation (syn), inverse-Compton radiation (IC), adiabatic cooling (dyn), and acceleration (acc) processes are considered. The case of $L_\gamma = 10^{47}$ erg s$^{-1}$ (implying $L = 2 \times 10^{51}$ erg s$^{-1}$) is shown.
considered, we find that the pion and muon components are almost always dominant and the kaon component could be relevant only above PeV energies.

4.5 Diffuse Neutrinos from Low-Luminosity GRBs and Hypernovae

Finally, we calculate the diffuse neutrino flux by convolving the neutrino spectra for different luminosities with $10^{45}$ erg s$^{-1} \lesssim L_\gamma \lesssim 10^{48}$ erg s$^{-1}$. The upper luminosity limit is found by constraining the jet to be choked with shocks that are not radiation dominated. The lower luminosity limit is chosen such that the results are not sensitive to this choice. As emphasized above, contrary to predictions for neutrino emission from optically-thin environments, we do not have much uncertainty in values of $f_{p\gamma}$, i.e., the $p\gamma$ efficiency is always close to the maximum. Thus, as long as the rate uncertainty is not too large, the only critical parameter is the total energy of CRs even though there are other subparameters such as $M_{\text{ext}}$ and $r_{\text{ext}}$. In this work, the jet kinetic energy is assumed to be similar to that of classical GRBs. Here, importantly, even if the observed GRB luminosity is low (recall that “low-luminosity GRBs” here are defined based on the observed luminosity), choked jets themselves may be as powerful as the jets of classical high-luminosity GRBs. In the CJ-SB model, the choked jet has isotropic-equivalent luminosity $L \sim 10^{51} - 10^{52}$ erg s$^{-1}$, but the observed gamma-ray luminosity is smaller by a factor of $(2/\theta_j^2)(T/t_{\text{eng}})$. (Clearly, $t_{\text{eng}}$ can also play a large role in determining whether a jet will give rise to a classical GRB or an LL GRB). For the shock breakout luminosity $L_\gamma$, the total absolute CR energy in the jet is assumed to be $\mathcal{E}_{\text{CR}} = (\xi_{\text{CR}}/\xi_\gamma)(L_\gamma T) \approx 6.3 \times 10^{50}$ erg $(\xi_{\text{CR}}/2)L_{\gamma,47}T_{3.5}$ (where $\chi_{\text{CR}} \equiv \xi_{\text{CR}}/\xi_\gamma = 2(0.25/\xi_\gamma)(\xi_{\text{CR}}/0.5)$ is the so-called CR loading factor). Note that the total absolute CR energy scales as the observed gamma-ray luminosity. Also, the CR spectrum is assumed to be $dN_p/dE'_p \propto E'^{-2}_p$.

The diffuse neutrino flux is calculated via (e.g., [238])

$$\Phi_\nu = \frac{c}{4\pi H_0} \int_{z_{\text{min}}}^{z_{\text{max}}} dz \int_{L_{\gamma,\text{min}}}^{L_{\gamma,\text{max}}} dL_\gamma \frac{dR_{\text{cho}}(z)/dL_\gamma}{\sqrt{\Omega_M (1+z)^3 + \Omega_\Lambda}} \left( \frac{dN_\nu((1+z)E_\nu)}{dE'_\nu} \right),$$  

(4.14)
where $dN_\nu/dE'_\nu$ is the neutrino spectrum per burst, $H_0$ is the Hubble constant, $\Omega_M$ and $\Omega_\Lambda$ are cosmological parameters. If LL GRB progenitors evolve as the star-formation rate (SFR), we rescale the function found by [98]

$$R_{\text{cho}}(z) = f_{\text{cho}}R_{LL} \times \left[ (1 + z)^{p_1\kappa} + \left( \frac{1 + z}{5000} \right)^{p_2\kappa} + \left( \frac{1 + z}{9} \right)^{p_3\kappa} \right]^{1/\kappa},$$

with $\kappa = -10$, $p_1 = 3.4$, $p_2 = -0.3$, $p_3 = -3.5$, $f_{\text{cho}}$ expresses the contribution of choked jets without shock breakout (i.e., orphan neutrinos), and $R_{LL} \sim 100 - 200$ Gpc$^{-3}$ yr$^{-1}$ is the local LL GRB rate at $z = 0$. Ref. [167] constructed a luminosity function (i.e., the number of bursts with an observed isotropic-equivalent luminosity within a given luminosity interval) uniquely for the LL GRB population

$$\frac{dR_{LL}}{dL_\gamma} \approx \frac{(\alpha - 1)R_{LL}}{L_m} \left( \frac{L_\gamma}{L_m} \right)^{-\alpha}.$$ 

It was found that the data was fit best with a local rate of $R_{LL} = 164^{+98}_{-65}$ Gpc$^{-3}$ yr$^{-1}$, index $\alpha = 2.3 \pm 0.2$ and characteristic luminosity $L_m = 5 \times 10^{46}$ erg s$^{-1}$.

Fig. 4.5 shows the diffuse neutrino flux from LL GRBs for different components. For our parameter set in the CJ-SB model that explains LL GRBs, we find that the diffuse neutrino flux is compatible with the measured flux for $E_\nu \sim 0.1 - 1$ PeV. There are three relevant remarks. (i) First, since the gamma-rays and the dominant component of neutrinos are produced in different regions, a prediction of the CJ-SB model is that the majority of the LL GRB neutrino signal arrives $(r_{sb} - r_{\text{stall}})/c \sim 100 - 1000$ s before the LL GRB triggers a detector. (ii) Second, the VHE neutrino emission from choked jets is highly beamed in the CJ-SB model. On the other hand, the shock breakout contribution is nearly isotropic so that the associated neutrino emission can be observed from off-axis observers [234]. (iii) Third, precursor neutrinos from choked jets will be found within a much smaller temporal window $(t_{\text{eng}} \sim 10^{1.5}$ s) compared to the electromagnetically observed LL GRBs and/or shock breakout emission.

For comparison, we also show one of the predictions of the EJ model for $\Gamma = 5$. We assume that the luminosity function is constant and the redshift dependence is taken from Ref. [239] but also follows the SFR. Although the model uncertainty is rather large, we confirm the previous results that the EJ model may also give a significant contribution to the diffuse neutrino flux [190, 191] at large observed energies (i.e. $E_{\nu,\text{obs}} \gtrsim 1$ PeV).
Figure 4.5. All-flavor diffuse VHE neutrino fluxes from LL GRBs in various models. The choked jet CJ (this work), shock breakout (CJ-SB) [234], and emergent jet (EJ) [190] components are shown. The shock breakout component has been updated to include the newest luminosity function and redshift evolution, while the EJ component is luminosity insensitive with the redshift evolution of Ref. [239] and is shown for illustrative purposes. Note that neutrinos are observed as prompt emission or precursor emission. The IceCube data based on the combined analysis [18] and up-going muon neutrino analysis [179] are overlaid.
The spectral shape of earlier results [190,191] is seen by the recent estimate of Ref. [240]. But the overall normalization is different due to different assumptions on the CR loading factor and LL GRB rate. In Fig. 4.5, we show the $p\gamma$ component for the EJ model while the $pp$ component is less important.

By definition, LL GRBs have gamma-ray counterparts, which are attributed to the shock breakout emission in the CJ-SB model or jet emission in the EJ model. Regardless of the viability of each model as an explanation for LL GRBs, the existence of choked jets is naturally expected and should be anticipated in any situation with a jet buried deep inside a star with or without extended material around it. In the CJ model, there is no obvious high-energy electromagnetic counterpart. It is known that long GRBs are associated with core collapse SNe (e.g., GRB 060218/SN 2006aj, GRB 980425/SN 1998bw, and GRB 100316D/SN 2010bh). These SNe tend to be Type Ibc meaning little to no hydrogen or helium is observed in the ejecta. There are also broad-line Type Ibc SNe which are often referred to as hypernovae. SNe associated with LL GRBs (although they are not necessarily hypernovae) are also characterized by transreletavistic ejecta. It is then reasonable to assume that a significant fraction of broad-line Type Ibc SNe or hypernovae, even those without accompanying GRBs, contain a choked jet such as in the CJ model. A joint investigation between IceCube and the ROTSE collaboration attempted to detect optical transients from Type Ibc SNe coincident with neutrino multiplets [241]. No such events were found, but an upper limit on the rate of SNe with a jet was found to be $\approx 4.2\%$ of the assumed rate of ccSNe within 10 Mpc. This study could be replicated using neutrino singlets from IceCube and the All-Sky Automated Survey for Supernovae (ASAS-SN; www.astronomy.ohio-state.edu/~assassin) network to further constraint our CJ model.

More specific results involving the CJ scenario are shown in Fig. 4.6. Although the model uncertainty is large (since $\chi_{\text{CR}}$ for GRB jets is not well-known), our results indicate that it is possible for choked jets to achieve the observed level of the diffuse neutrino flux. In principle, lower values of $\chi_{\text{CR}}$ could be compensated by larger values of $f_{\text{cho}}$. We set a rough upper limit on the choked jet contribution by using the observed hypernova rate, $R_{\text{HN}} \approx 4000 \text{Gpc}^{-3}\text{yr}^{-1}$ [88,186], which gives $f_{\text{cho}} \approx 40$ [63]. A CJ rate similar to that of HNe is also in agreement with Ref. [170] who considered neutrinos from electromagnetically dim sources. They found that transients with a rate of $\approx 10^5 \text{yr}^{-1}$ up to $z \sim 1$ could produce a detectable flux of neutrinos. A similar result is also obtained by Ref. [152].
Figure 4.6. All-flavor diffuse neutrino fluxes from choked jets. Neutrino emission from LL GRBs is shown for the CJ-SB model (this work) and EJ model [190]. In addition, orphan neutrino emission from choked jets is included (thick curves). See the text for details.
Using the assumed rate and CR energy injection per event, the all-flavor diffuse neutrino flux is analytically estimated to be

\[ E^2 \Phi_\nu \simeq 0.76 \times 10^{-7} \text{ GeV cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} f_{\text{sup}} \text{min}[1, f_{\text{p}}^4] \]

\[ \times \left( \frac{\xi_z}{3} \right) \left( \frac{f_{\text{cho}} \mathcal{E}_{\text{CR}} R_{\text{LL}}}{10^{45} \text{ erg Mpc}^{-3} \text{ yr}^{-1}} \right) R_p^{-1}, \quad (4.17) \]

where \( f_{\text{sup}} \) is the suppression factor due to meson and muon cooling, \( \xi_z \) is a factor accounting for redshift evolution of the rate [42, 43], and \( R_p = \ln(\varepsilon_p^M / \varepsilon_p^\text{min}) \sim 10 \) is the bolometric correction factor. Interestingly, even this simple-minded calculation is remarkably close to the measured all-flavor diffuse flux of neutrinos [18],

\[ E^2 \Phi_\nu^{\text{obs}} |_{30 \text{ TeV}} \sim 10^{-7} \text{ GeV cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}. \quad (4.18) \]

Murase et al. [50] showed that neutrino sources obscured in the GeV-TeV gamma-ray range are necessary to explain the IceCube data below 100 TeV with extragalactic sources, independently of the neutrino production mechanism. LL GRBs and choked jets satisfy this criterion, and their contribution to the extragalactic gamma-ray background is negligible.

As a lower limit, for given parameters we can use the rate of observed LL GRBs (i.e., excluding choked jets without prominent shock breakout emission). Noting that the emitted CR energy is roughly the same as that for hypernovae, \( R_{\text{LL}} \sim 100 \text{ Gpc}^{-3} \text{ yr}^{-1} \) results in \( E^2 \Phi_\nu \sim 10^{-8} \text{ GeV cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \), which is compatible with the IceCube data above 100 TeV. Using modest values of \( f_{\text{cho}} \sim \) a few allows us to reasonably fit the IceCube data obtained from the combined analysis.

One can set an optimistic upper limit for the contribution of orphan neutrinos from choked jets by assuming the same spectral shape as in the CJ-SB model with CR energy injection rate of \( Q_{\text{CR}} \sim 10^{46} \text{ erg Mpc}^{-3} \text{ yr}^{-1} \). This upper limit is set by the reasonable expectation that the CR injection by GRB jets does not exceed the CR injection by SN remnants (see Ref. [50]). Note that the CR injection rate inferred by observations of the Galactic CRs is \( \sim 10^{45} - 10^{46} \text{ erg Mpc}^{-3} \text{ yr}^{-1} \) [65]. Fig. 4.6 indicates the the optimistic upper limit can exceed the IceCube data in principle. The spectral shape is suggestive as it is globally soft for \( 10 \text{ TeV} \lesssim E_\nu \lesssim 5 \text{ PeV} \), but avoids the constraints set by the Fermi extragalactic gamma-ray background measurement in the sub-TeV range [40, 44, 102].

While our results show that the choked jets are energetically plausible as high-energy neutrino sources, we have not tuned parameters to fit the IceCube data quantitatively.
Because of the limited statistics of the IceCube data and a tension among the different analyses, such an attempt is beyond the scope of the present work. Also, better fits of the spectral shape and normalization of the diffuse flux would be possible by changing the parameters of the jet and/or extended material within model uncertainties.

4.6 Summary and Discussion

We have revised the VHE neutrino emission from LL GRBs, taking into account the jet-stalling condition for a dense circumstellar wind and radiation constraints. Lower-power jets and/or more extended external material are more favorable for both jet-stalling and VHE neutrino production. This implies the relevance of “orphan” neutrinos from choked jets with no prominent electromagnetic counterparts. Using properties inferred from LL GRB observations (where we did not tune the parameters to explain the IceCube data), we found that the spectral shape and flux normalization of the CJ model can be consistent with the present IceCube data. Although the rate of choked jets with dim shock breakout emission is unknown, it is plausible to use the rate of broad-line Type Ibc SNe as an upper limit. It is currently difficult to exclude parts of the CJ/CJ-SB parameter space in our model since there are many degeneracies (e.g. the amount of CR energy injected by CJs versus the rate of HNe which have an accompanying luminous, relativistic jet as seen in Eq. 4.17). However, by assuming that the seed photons for p\gamma interactions in the CJ scenario come predominately from the black body emission of the jet head, it should be possible to constrain the combined parameter space of the jet luminosity, head position, and jet Lorentz factor (see §4.4 above) by using the energy at which the observed neutrino flux is at its maximum. An important prediction of the CJ-SB model is that the majority of neutrinos will be precursors to the prompt gamma-ray emission. Therefore, for a neutrino-LL GRB coincidence search it is imperative to look for neutrinos in temporal blocks \(\sim 100 - 1000\) s before the GRB trigger. Based on the currently implied rate of LL GRBs, two such coincident detections can be expected to occur within the next five years of IceCube operation. For this purpose, we emphasize that better all-sky monitors in the x-ray and gamma-ray range, which are also suitable for detections of high-redshift GRBs, are necessary. Coincident searches can also be expanded to include hypernovae and other energetic SNe. While the rate of such events are higher, the delay between the neutrino and optical/x-ray signal is unknown but may be \(\gtrsim 1000\) s.
Detecting precursor neutrinos with short duration would support the CJ-SB model for LL GRBs. Further observations of $1 \text{ TeV} \lesssim E_\nu \lesssim 100 \text{ TeV}$ neutrinos may provide further information about the envelope mass, radius, and density profile for the extended material around WR stars. Such regions are hard to probe observationally, especially if the mass loss rate of GRB progenitors significantly increases in the months-weeks before collapse. While we considered only environments around WR stars in this work, the treatment may be equally important in BSGs, e.g., [242], which are also believed to be collapsar progenitors and tend to have stellar envelopes that extend up to $10^{13.5} \text{ cm}$. 
5.1 Introduction

The IceCube Antarctic neutrino observatory has been observing high-energy (HE) neutrinos with energies of $E_\nu \gtrsim 10$ TeV for almost half a decade [24, 25, 52-54, 159]. They observe an apparently isotropic diffuse neutrino flux that is equally distributed between the neutrino flavors (i.e., $\nu_e : \nu_\mu : \nu_\tau \approx 1 : 1 : 1$) [18]. The sources of these HE neutrinos are still unknown, despite observations of $\sim 100$ HE contained events, and tens of thousands of through-going track events to date. Gamma-ray bright sources such as gamma-ray bursts (GRBs) [39, 171] (see also [36, 243, 244] for the latest papers after the discovery of IceCube neutrinos) and blazars [41, 74, 79, 144, 245, 246] are now disfavored as the main origins of IceCube’s neutrinos [46, 55, 151].

During the last few years, $\gamma$-ray dark sources have been investigated as potential sources of the HE diffuse flux (see Ref. [50] and references therein). Such sources may be able to accelerate HE cosmic-rays (CRs) which interact with their surrounding environment to produce $\gamma$-rays and neutrinos via $pp$ or $p\gamma$ interactions. While the $\gamma$-rays are attenuated by a dense photon field (via $\gamma\gamma$ interactions), the HE neutrinos escape. Studies of $\gamma$-ray dark sources attempt to determine their potential maximum contribution to the diffuse neutrino flux, while still respecting multi-messenger observations such as the diffuse extragalactic $\gamma$-ray background [48, 49, 51, 60, 75, 116, 152]. By definition, the
amount of non-thermal energy in CRs is difficult to directly determine for γ-ray dark transients, making it harder to constrain the suggested theoretical models (although rare sources such as jetted tidal disruption events (TDEs) are already disfavored by the absence of clustering [51] and diffuse emission [51,59,60]).

It has been known since the combined observation of SN 1998bw/GRB 980425 [247,248] that so-called long gamma-ray bursts (GRBs) share a common progenitor with core-collapse supernovae (ccSNe). Prompt GRB emission is observed once a relativistic jet escapes from the stellar envelope of its progenitor during core-collapse. It is then natural to suppose that some jets do not escape the stellar envelope [166,249], instead resulting in a trans-relativistic explosion or stripped-envelope SNe [209]. While no events have been definitively identified as an example of a “choked jet” SN (cjSNe), it is believed that unusually energetic explosions or those with trans-relativistic ejecta such as a Type Ic broad-line SNe or hypernovae may be the result of such a scenario [248,250,251]. Low-luminosity GRBs or trans-relativistic SNe have been suggested as HE neutrino emitters [190,191,234].

Conventional (i.e., γ-ray bright) GRBs were thought to be prime candidates to produce HE neutrinos [39], but they can only contribute to ≲ 1% of IceCube’s neutrinos [46,46,145,185,252,253]. While it stands to reason that their choked-jet brethren may deposit a significant amount of energy into CRs and neutrinos [169,170,223,224,254,255], the non-detection of precursor neutrinos from high-luminosity GRBs already indicate that HE neutrino production should be suppressed inside a star for powerful GRBs. However, Ref. [75] showed that this is consistent with the theoretical expectation that CR acceleration is inefficient at radiation-mediated shocks, and proposed that low-power choked jets may give a significant contribution to the observed neutrino flux without contradicting the non-detection of GRB neutrinos. Alternatively, Ref. [48] considered HE neutrino production in choked jets embedded in the circumstellar material or extended envelope. It has also been suggested that choked jets may account for the observed high-energy neutrinos [75], especially in light of the medium-energy excess [48,75,152,192].

A search for HE neutrinos coincident with SN 2008D – assuming it contained a choked jet – has been performed [256]. In this current work, we investigate a population of γ-ray dark transients – stripped-envelope SNe – as the potential sources of HE neutrinos. We use observed SNe events to constrain the fraction of the population that can produce neutrinos, and the amount of energy each event can deposit into CRs (E_{cr}). The background only hypothesis cannot be ruled out using only the present neutrino
data (i.e., there is not a significant association between neutrinos and Type Ibc SNe). We are therefore able to place upper limits on $E_{\text{cr}}$ and the fraction of SNe that could have a choked jet pointed towards Earth ($f_{\text{jet}}$). Note that the latter parameter cannot discriminate between SNe that do not have jets, and those which have jets that are not pointing towards us because of beaming considerations.

The IceCube collaboration is expected to release the remainder of their upgoing track-like events (roughly 7 years in total), which will improve the limits of this analysis. Furthermore, our procedure can be applied to test neutrino coincidences with any $\gamma$-ray dark source (e.g., low-luminosity GRBs, choked-jet TDEs). With the introduction of wide field-of-view optical surveys such as the Palomar Transient Factory [257] and All-Sky Automated Survey for Supernovae (ASAS-SN)*, new rich datasets of $\gamma$-ray dark transient events will be made available to test as potential neutrino sources. Furthermore, cjSNe are important targets for the Astrophysical Messenger Observatory Network (AMON) [258]. Because current optical telescopes view the entire sky every few days, it will soon be possible to understand the initial period of a SN explosion; and, improving this analysis by determining the exact time at which a choked jet would occur.

5.2 Data

We consider one-year of IceCube data that was taken between May 2011 and May 2012 using the 86 string detector array [159]. This sample contains $\sim 70,000$ upgoing, track-like events. To eliminate contamination from atmospheric muons, we consider only upgoing track events. Each neutrino event contains information on its energy proxy – which is related to the total amount of photoelectrons observed in the detector [259] – the day on which it was detected, its arrival direction, and corresponding angular uncertainty. Track-like events typically have median angular uncertainties of $\lesssim 2^\circ$, although they depend on their energies. In particular, the low-energy events are affected by large kinematic angles. There are also misreconstructed events. We note that $< 100$ events have large angular uncertainties of $\gtrsim 50^\circ$.

There are 29 Type Ibc SNe $^\dagger$ that were detected in the Northern hemisphere during the neutrino data collection period. From the Open SNe catalog (See [260] and references

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*http://www.astronomy.ohio-state.edu/~assassin

$^\dagger$Note that an additional SN – SN2012bz [251] – is omitted from this analysis because it does not have a measured date on which its optical flux reaches a maximum, and therefore does not meet the requirements of this analysis.
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**Table 5.1.** Observations of the 29 Type Ibc SNe which were detected in the Northern Hemisphere between May 2011 and May 2012.
within), we determine the date of their maximum optical brightness, their location in the
sky, and redshift (the latter for all but one event). Table 5.1 lists the SNe observations
used in our analysis. Note that available catalogues are quite incomplete at present. Also,
with the inclusion of new optical transient “factories” such as the Palomar Transient
Factory, and ASAS-SN the number of Type Ibc SNe detected each year following 2012 is
significantly larger.

5.3 Analysis

5.3.1 Signal and Background PDFs

We use the unbinned maximum log-likelihood method first developed for neutrino
astronomy by Braun et al. [261] to search for associations of neutrino arrival directions
with source positions. This technique was later adapted to analyze the correlation
between neutrino events and the prompt emission of GRBs [46]. The same analysis
is appropriate for determining the correlation of SNe explosions with neutrino events,
although the likelihoods of the neutrino arrival time, direction, and deposited energy
must be adjusted for these new potential sources. The method we present can be easily
extended to include additional years of IceCube data, or modified for other gamma-ray
dark sources, such as choked-jet TDEs.

The neutrinos in our model are assumed to be produced by a jet that is launched
during the core-collapse of the SN progenitor, but is subsequently choked off by the
stellar or circumstellar envelope. For typical GRB jet luminosities, this process occurs
within \( \sim 10 - 100 \) s after the initial stellar explosion [166, 249]. However, almost all
SN are not detected until days after the initial explosion, and it is hard to precisely
reconstruct the stellar explosion time. By studying GRBs that are associated with SNe,
it was found that the prompt GRB emission was detected \( \sim 13 \pm 2.3 \) days before the SN
reached its maximum optical brightness (see [262] and references within).

We therefore use the maximum SN brightness time as a proxy for the stellar explosion
time, and assume that for any given Type Ibc SNe, the difference in time between the
two is given by a Poisson distribution with \( \lambda_T = 13 \) days. Note that since the arrival
times for the neutrinos are coarsely binned by the day they were detected, a discrete
probability distribution is appropriate as an approximation here. The signal probability
mass function (i.e., the analog of a probability density function but for discrete values)
for the arrival time of neutrinos is therefore

\[ S_T(T_{\text{arr}, \nu}, T_{\text{max}, \text{sn}}) = e^{-\lambda_T} \frac{\lambda_T^{T_{\text{max}, \text{sn}} - T_{\text{arr}, \nu}}}{(T_{\text{max}, \text{sn}} - T_{\text{arr}, \nu})!}, \tag{5.1} \]

where \( T_{\text{max}, \text{sn}} \) and \( T_{\text{arr}, \nu} \) are integers rounded to the nearest Modified Julian Day.

The background neutrino events for this analysis are primarily neutrinos produced by CR interactions in the Earth’s atmosphere, and are assumed to occur at a constant rate. Therefore, the background probability density function (PDF) for the arrival time of neutrinos (denoted as \( B_T \)) is constant during the observation window. This window is taken to be the central 99% confidence interval for Eq. 5.1, \( T_{\text{max}, \text{sn}} - 19 \text{ days} \leq T \leq T_{\text{max}, \text{sn}} - 4 \text{ days} \).

Unlike Ref. [261], we assume that the signal PDF for neutrino arrival directions is the von Mises-Fisher (aka the Kent [263]) distribution [253,264]

\[ S_{\text{dir}}((\alpha, \delta)_\nu, (\alpha, \delta)_{\text{sn}}) = \frac{\kappa}{4\pi \sinh \kappa} e^{\kappa \mu}, \tag{5.2} \]

where \( \kappa = 1/\sigma_{\nu}^2 \) is related to the angular error of the neutrino arrival direction, since it can be assumed that the uncertainty of the SN position is negligible compared to \( \sim 1^\circ \). The cosine of the angular separation between the neutrino and SN position is \( \mu = \cos \Delta \psi \). Note that for \( \Delta \psi \ll 1 \) and \( \kappa \gg 1 \) Eq. 5.2 reduces to a 2D Gaussian PDF \( S_{\text{dir}} = \frac{1}{2\pi \sigma_{\nu}^2} e^{-\Delta \psi^2/2\sigma_{\nu}^2} \) (compare with Eq. 9 in [261] and Eq. 3 in [46]).

The background PDF for the arrival direction of neutrinos (denoted as \( B_{\text{dir}} \)) is the product of the PDFs for the neutrino right ascension or RA \( (\alpha) \) and declination \( (\delta) \). The former is adequately approximated by a uniform distribution between 0 and \( 2\pi \), RA \( \sim U(0, 2\pi) \), while the latter is constructed using the declinations of neutrinos which are outside of the acceptance windows of any SNe (See Fig. 5.1). The acceptance window for arrival direction is determined by the 99% lower confidence interval of Eq. 5.2 (i.e., all neutrinos which satisfy \( \mu_{\nu} \geq \mu_{99\%} \)).

The energy and zenith dependent IceCube effective areas provided in Aartsen et al. (2015) [159] are convolved with an unbroken \( E_{\nu}^{-2} \) spectrum for the neutrino fluence to construct the signal PDF \( S_{\text{ene}} \) for the neutrino energy proxy, which is related to the total amount of photoelectric energy a neutrino deposits in the detector (see Fig. 5.2, dash-dotted lines). The total amount of CR energy in the jet – which can be related to the neutrino energy if one assumes that the choked jet is calorimetric – is left as a tunable parameter so we can set a limit on the CR acceleration efficiency of cjSNe.
**Figure 5.1.** Background PDF for neutrino arrival declination. Constructed using neutrino data from events not within the acceptance window (i.e., outside of the 99% central (lower) confidence level of Eq. 5.1 (5.2))
The luminosity distance is obtained from the measured redshift of each SN. For the one SN for which no redshift is determined (PTF11ilr), we randomly draw a value from the remaining SN redshifts. The background energy PDF \( B_{\text{ene}} \) is again constructed using neutrinos outside of the acceptance windows of all SNe, and is in good agreement with the proxy energy distribution of atmospheric neutrinos for proxy energies \( \gtrsim 10^3 \) (see Fig 5.2, dashed lines) [259]. Because of the geometry of the IceCube detector, and the direction and energy dependent neutrino opacity of the Earth, both the signal and background energy PDFs depend on the IceCube zenith angle (i.e., the declination of the neutrino event, see Fig 5.2). Therefore, we produce energy PDFs for each zenith bin of the IceCube effective area from Aartsen et al. 2015.

5.3.2 Construction of Test Statistic

The likelihood that a given neutrino event is associated with a particular SN is related to the product of the ratio of signal(-like) and background PDFs defined above

\[
\frac{S}{B} = \frac{S_T S_{\text{dir}} S_{\text{ene}}}{B_T B_{\text{dir}} B_{\text{ene}}}. \tag{5.3}
\]

The composite likelihood for each SN is the product of the likelihoods for each neutrino event associated with a SN. We perform a stacking analysis by forming a likelihood composed of the product of all SN likelihoods. Each SN likelihood is weighted by an appropriate Poisson factor, which accounts for the number of neutrinos associated with the \( j \)th SN as a function of its background \( b_j \) and signal \( s_j \) rates (i.e., the number of signal-like and background neutrinos found in a SN’s acceptance window)

\[
P_j(s_j, b_j) = \frac{(s_j + b_j)^{N_j}}{N_j!} e^{-(s_j + b_j)}. \tag{5.4}
\]

The background rate \( b_j \) is calculated using data from randomized synthetic experiments. The neutrino arrival data are scrambled to create synthetic data sets. The expected background rate for the \( j \)th SN is taken to be the average number of associated neutrinos from these scrambled datasets. The signal rate \( s_j \) is calculated by maximizing the log-likelihood, where the likelihood function is given by

\[
\mathcal{L}(\{s, b\}) = \prod_{j=1}^{N_s} \prod_{i=1}^{N_i} \mathcal{L}_{i,j}(\{s, b\}).
\]

The likelihood that an individual neutrino event \( i \) is associated with SN \( j \) is given by

\[
\mathcal{L}_{i,j}(s_j, b_j) = \frac{s_j S_i + b_j B_i}{s_j + b_j}. \tag{5.5}
\]
Figure 5.2. Signal and background energy PDFs for neutrino arriving in the IceCube zenith bins $-1 \leq \cos \theta_z < -0.9$ (top) and $0.1 \leq \cos \theta_z < 0.0$ (bottom). Note that the IceCube detector is most sensitive to very high-energy neutrinos coming from the horizon, since the Earth is opaque to neutrinos with energy $E_\nu \gtrsim 1$ PeV. The proxy energy for each neutrino is related to the amount of photo-electrons produced in the detector.
For simplicity, we maximize the ratio of the log-likelihoods assuming a signal and no signal hypothesis, with \( TS = \ln \left[ \frac{\mathcal{L}(s)}{\mathcal{L}(0)} \right] \). Our test statistic is then

\[
TS = \sum_{j=1}^{N_{\text{sn}}} \left[ -s_j + \sum_{i=1}^{N_j} \ln \left( \frac{s_j S_i}{b_i B_i} + 1 \right) \right].
\]  

From Eq. 5.3 \( S_i = S_{\text{dir}} S_{\text{ene}} \) and likewise for \( B_i \). The significance of our observed test statistic \( TS_{\text{obs}} \) is determined using a frequentist method. We produce synthetic background only data sets by randomizing the arrival times, directions, energy, and angular uncertainty of our original neutrino sample. We then compute a distribution of test statistics \( TS_{\text{bkg}} \) by applying Eq. 5.6 to 100,000 of these synthetic data sets. Note that in the background only hypothesis, some neutrinos in the scrambled data are identified as signal-like neutrinos (i.e., \( s_j > 0 \), see the gray band of Fig. 5.4). Therefore, we rejected the background only hypothesis if the observed test statistic \( TS_{\text{obs}} \) is > 90% of the distribution \( TS_{\text{bkg}} \).

### 5.3.3 Placing Upper Limits

The test statistic calculated using observed (i.e., unscrambled) data \( TS_{\text{obs}} \) does not deviate from the distribution of test statistics calculated using scrambled data \( TS_{\text{bkg}} \) (See Fig. 5.3 below). This latter scenario is referred to as the background only hypothesis. Therefore, we can place upper limits on the total SN energy deposited in CRs \( E_{\text{cr}} \) (assuming the cjSNe are calorimetric with regards to neutrino production), and the fraction of SN which have choked jets pointed towards earth \( f_{\text{jet}} \). To do this, we calculate the probability that the distribution of SNe with fixed \( E_{\text{cr}} \) and \( f_{\text{jet}} \) would produce a > 90% CL detection. In practice, this is given by the fractional number of \( TS_{\text{sig}} \) that are greater than 90% of \( TS_{\text{bkg}} \).

We accomplish this by injecting true signal neutrinos into our randomized synthetic data sets. For each data set, \( f_{\text{jet}} N_{\text{sn}} \) SNe are chosen from random to produce signal neutrinos. These signal neutrinos have arrival time and direction, as well as energies drawn from the signal PDFs described above. The number of true signal neutrinos injected in the acceptance window of each SN is determined using the expected fluence from that SN. The muon neutrino fluence per logarithmic energy interval from a calorimetric SN jet can be approximated as [42]
\[ \mathcal{F}_{\nu} \approx \frac{1}{8 \pi D_L^2 \mathcal{R}} \frac{\mathcal{E}_{\text{cr}}}{\text{erg cm}^{-2}}, \quad (5.7) \]

where \( D_L \) is the luminosity distance of the SN, and \( \mathcal{R}_{\text{cr}} = \ln(\varepsilon_{\text{cr,max}}/\varepsilon_{\text{cr,min}}) \approx 18 \) is a bolometric correction factor for a \( n_{\varepsilon_{\text{cr}}} \propto \varepsilon_{\text{cr}}^{-2} \) CR spectrum. The number of signal neutrinos is drawn from a Poisson distribution, with mean given by the expected number of signal neutrinos. The latter is determined by the product of the SN neutrino fluence with IceCube’s effective area. While we use the neutrino effective area given by Ref. [159], the energy and zenith-angle averaged value is approximately \( \bar{\mathcal{A}}_{\text{eff}} \sim 10^4 \text{ cm}^2 \).

### 5.4 Results

Based on the distribution of TS assuming a background only model (i.e., scrambling the existing neutrino data without injecting signals), we do not see any significant association between the HE IceCube neutrinos and optically detected Type Ibc SNe for the time period May 2011-May 2012.

Fig 5.3 gives the cumulative distribution function (CDF) of TS assuming a background only hypothesis, with \( \text{TS}_{\text{obs}} \) given by the dashed line. Each SN was associated with \( \sim 15 \) track events on average. This level of association was expected from the synthetic experiments. Large values of TS correspond to experiments with many signal-like neutrinos, while values of TS \( \approx 0 \) indicate few signal-like neutrinos.

Looking at the best fit values for the number of signal-like neutrinos in our observed data, we find that they are within the 99% confidence interval determined by the synthetic background data distribution (See Fig 5.4). This gives a further indication that we do not see a significant association between HE neutrinos and the 29 Type Ibc SNe in our sample.

We also produce distributions of TS_{\text{sig}}, where true signal neutrinos are injected into our sample. Fig. 5.5 (left) compares the CDF of the background only scenario (solid line) with different distributions of TS_{\text{sig}} with fixed \( f_{\text{jet}} = 1.0 \) and \( \mathcal{E}_{\text{cr}} = 10^{51.3} \text{ erg} \) (dotted), \( \mathcal{E}_{\text{cr}} = 10^{52} \text{ erg} \) (dash-dotted). Fig. 5.5 (right) compares the background only CDF with two different signal distributions, with fixed \( \mathcal{E}_{\text{cr}} = 10^{52} \text{ erg} \) and \( f_{\text{jet}} = 0.56 \) (dashed) and \( f_{\text{jet}} = 1 \) (dash-dotted). These distributions are compared with the observed value TS_{\text{bkg,90}} (thick dashed) to determine what is the probability of making a 90% CL observation for a particular signal hypothesis. The wavy features for the signal CDFs are a result of SNe.
Figure 5.3. Cumulative distribution function of $TS_{bkg}$ generated from synthetic experiments of randomized neutrino data with the observed test statistic $TS_{obs}$ (dashed line) and the 90% upper limit on the background distribution $TS_{bkg,90}$ used to set exclusion contours for the signal hypothesis (dash-dot, see text for details). $TS_{obs}$ is consistent with a background only hypothesis.
Figure 5.4. Best fit value for the number of signal-like neutrinos ($s_j$) for each SN in our sample (with their label number on the x-axis). The black dots correspond to the best fit values from our observed data, while the gray band is the 99% confidence interval determined from the synthetic data samples of randomized data. Note that $s_j \neq 0$ for all cases, meaning that some neutrinos from the scrambled data sets are identified as signal-like neutrinos. We do not see a significant number of signal-like neutrinos above what is expected from a background only model.
Figure 5.5. Cumulative distribution functions of TS comparing the background and signal scenarios with $TS_{\text{bkg},90}$. We show how the CDF produced by the signal hypothesis changes as $E_{\text{cr}}$ is varied with $f_{\text{jet}} = 1.0$ (left), and for different values of $f_{\text{jet}}$ with $E_{\text{cr}} = 10^{52}$ erg (right). Insets show where $TS_{\text{bkg},90}$ intersects with each CDF. The wavy features of the signal CDFs are a results of SNe from different luminosity distances being included or excluded in a synthetic experiment. When $E_{\text{cr}}$ is the same for all SNe, closer jetted SNe produce a higher fluence of signal neutrinos in the detector.
from varying luminosity distances being included or excluded in a synthetic experiment. When $\mathcal{E}_{\text{cr}}$ is the same for all SNe, closer jetted SNe result a higher fluence of signal neutrinos in the detector. Note that, as the number of signal neutrinos in our synthetic data set decrease – as either $\mathcal{E}_{\text{cr}}$ or $f_{\text{jet}}$ decrease – the signal distribution of TS begins to resemble that of the background only scenario.

We are able to place upper limits on the fraction of SNe which have choked jets pointed towards us $f_{\text{jet}}$, and the total amount of CR energy contained in such jets $\mathcal{E}_{\text{cr}}$. Fig 5.6 shows a heat map of the exclusion region in the $\mathcal{E}_{\text{cr}} - f_{\text{jet}}$ plane which is given by the probability of observing $\text{TS}_{\text{sig}} > 90\%$ of $\text{TS}_{\text{bkg}}$ (see §5.3.3). We are able to place limits on $f_{\text{jet}}$ down to $\mathcal{E}_{\text{cr}} \sim 10^{51.5}$ erg, which is comparable to a typical SN explosion energy. Note that $\mathcal{E}_{\text{cr}}$ is the isotropic equivalent energy, so the true amount of jet energy contained in CRs would be reduced by a factor $\theta^2_j/2 \sim 10^{-1}$ depending on the opening angle of the choked jet $\theta_j$.

We can compare our result with a simple analytic argument, using the typical IceCube fluence sensitivity $\phi_{\text{lim}} \sim 10^{-4}$ erg cm$^{-2}$ and summing over each SN. The limit on $\mathcal{E}_{\text{cr}}$ and $f_{\text{jet}}$ is set using the Poissonian probability of observing more neutrinos than the 90% upper limit assuming a background only hypothesis ($N_{\text{bkg},90}$). With the total neutrino background rate (i.e., $n_b = \sum_{j=1}^{N_{\text{sn}}} b_j$) the 90% upper limit on the number of observed neutrinos in the background only hypothesis is calculated solving [265]

$$
\sum_{x=0}^{N_{\text{bkg},90}} \frac{n_b^x e^{-n_b}}{x!} \leq 0.1
$$

(5.8)

The probability $P_{>90}$ of observing more neutrinos given a signal rate ($n_s$), assuming the average isotropic equivalent CR energy released per burst is $\tilde{\mathcal{E}}_{\text{cr}} = \mathcal{E}_{\text{cr}} f_{\text{jet}}$, is given by

$$
P_{>90} = \sum_{y=0}^{N_{\text{bkg},90}} \frac{(n_s + n_b)^y e^{-(n_s + n_b)}}{y!},
$$

(5.9)

where $n_s$ is estimated to be

$$
n_s = \phi_{\text{lim}}^{-1} \frac{\tilde{\mathcal{E}}_{\text{cr}}}{32\pi C} \sum_{j=1}^{N_{\text{sn}}} \frac{1}{D^2_{L,j}},
$$

(5.10)

For the 28 SNe in our sample with a measured redshift, this gives $\tilde{\mathcal{E}}_{\text{cr},90\%} \sim 10^{52}$ erg.

Fig. 5.7 compares the heat map of our numerical results (as seen in Fig. 5.6) with the analytic results produced by Eq. 5.9 (white solid line). We see that there is reasonable
Figure 5.6. Heat map that shows the exclusion limits for values of $E_{cr}$ and $f_{jet}$ based on synthetic experiments with true signal neutrinos added. Darker color means higher confidence in exclusion (e.g., the top right corner is excluded at 90% confidence).
agreement between the shape of the exclusion region from both methods, as well as the location of the 90% confidence limit at $E_{\text{cr}} \sim 10^{52}$ erg (for $f_{\text{jet}} = 1$). With an additional 6 years of IceCube data, we find using Eqs. 5.8-5.9 that the 90% confidence limit on $E_{\text{cr}}$ can be improved by a factor of $\sim 10$ (see Fig. 5.7 black dashed line, which was calculated using 131 Type Ibc SNe that were observed between May 2010-May 2017 with an extrapolation of the expected neutrino background rate for 7 years of data from 1 year of data).

5.5 Summary and Discussion

We performed an unbinned log-likelihood analysis of the $\sim 70,000$ up-going, track-like neutrino events observed by IceCube between May 2011 and May 2012 and 29 Type Ibc SNe observed during this same time period. We found no significant excess of signal neutrinos associated with these SNe (see Fig. 5.3). Based on this non-detection, we were able to place upper limits on the fraction of Type Ibc SNe which may harbor a choked jet pointed towards Earth, and the total amount of CR energy they can produce assuming such jets are efficient neutrino factories. Our upper limits will be improved by a factor of $\sim 10$ once the remaining 6 years of IceCube data are made public.

If cjSNe significantly contribute to the observed IceCube flux below 100 TeV, the required CR energy production rate is $3 \times 10^{53}$ erg Mpc$^{-3}$ yr$^{-1}$. For Type Ibc SNe with $2 \times 10^4$ Gpc$^{-3}$, the required CR energy is $E_{\text{cr}} \sim 10^{49}$ erg for $f_{\text{jet}} = 1$. In reality, $f_{\text{jet}} \ll 1$ is expected, so the number of “nearby” SN samples should be large enough ($\gg 1/f_{\text{jet}}$) to obtain meaningful constraints to test the choked jet scenario for IceCube neutrinos [75]. IceCube currently has upgoing track-like data for $\sim 7$ years (2010-2017). When this data is made available to the public, combined with the increased detection rate of ccSNe with ASAS-SN and other all-sky optical surveys – there are additional 104 Type Ibc SNe that were detected from May 2012 to May 2017 – cjSNe can be more tightly tested as potential sources of HE astrophysical neutrinos in future, especially with IceCube-Gen2.

Non-detection of neutrinos from cjSNe may not be surprising at all. Theoretically, the standard Fermi acceleration mechanism is inhibited in radiation mediated shocks, so powerful jets or compact progenitors would not be ideal as high-energy neutrino emitters via this process [75]. This also implies that one has to be careful to use the neutrino data to constrain physical quantities of choked jets, except for the CR acceleration efficiency. On the other hand, a correlation between HE neutrinos and cjSNe would indicate that
Figure 5.7. Comparison of the numerical results of our analysis (heat map, see Fig. 5.6) with the 90% upper limit determined by Eq. 5.9. The white solid line corresponds to the 90% upper limit using the 28 SNe with measured redshift, while the black dashed line corresponds to the analytic 90% confidence level using Eq. 5.9 and 131 SNe from May 2010-May 2017 (see text for details).
high-energy CRs can be accelerated in relativistic jets launched during stellar collapse, which has important implications for CR acceleration in dense environments. Sub-TeV neutrino signals may allow us to probe other alternative acceleration mechanisms, such as the neutron-proton-converter acceleration process \cite{205,206}. Also, a new window will have been opened into the interior of old, massive stars. Except for the thermal neutrinos observed from SN1987A, these high energy neutrinos would constitute the first direct observation of an exploding star beneath its photosphere\textsuperscript{1}. Such observations could allow stellar interior models to be directly tested using methods such as “neutrino tomography” \cite{170}.

Our analysis is one of the first systematic searches for neutrino sources from a population of choked jet SNe. We have assumed a $E_\nu^{-2}$ spectrum. While it is reasonable to assume the neutrino energy spectra is given approximately flat, these spectra are generally soft at high energies $E_\nu \sim 0.1 - 1$ PeV, and peak in the energy range that IceCube is most sensitive to $E_\nu \sim 100$ TeV. Future works will be able to improve the accuracy of this analysis by constructing an energy signal PDF by convolving different cjSNe neutrino models with the exact IceCube effective area.

Searches that rely on optical data only, or with some combination of soft $\gamma$-rays/X-rays are becoming increasingly important, as $\gamma$-ray bright sources such as GRBs and blazars are being ruled out as the main sources of the IceCube neutrinos. Furthermore, current and soon to be operational optical surveys will view the entire optical sky within 3 days up to a depth of 20.4 mag. By constantly monitoring the sky, we will gain a better understanding of the early explosion processes of ccSNe (i.e., $\lesssim 1$ day after the initial explosion). The upgraded Zwicky Transient Facility “will detect one SN within 24 hours of its explosion every night” \cite{263}. This benefits the analysis in two ways: 1.) It allows for a more accurate determination of the time window during which a choked jet may be formed in the ccSNe, reducing the signal time window from $\sim 15$ days, to hours. 2.) Early observations of SNe explosions will give the first direct observations of the material immediately surrounding a SN progenitor. Currently the mass loss processes which occur during the last $0.1 - 10$ year of a stripped-envelope progenitor are not well understood, and could provide insight into which types of stars will form a jet, and ultimately a typical Type Ibc SNe, a cjSNe, or a full GRB. This in turn will help guide searches for the most likely transients to produce HE neutrinos.

\textsuperscript{1}the region at which the stellar material becomes optically thick
Chapter 6
Discussion and Outlook

6.1 Summary

Since the original postulation of neutrinos by Pauli in 1930 [3], the study of these elusive particles has helped to push the boundaries of our understanding of the Universe. The confirmation of neutrino masses, and therefore neutrino oscillations were one of the first examples of particle physics beyond the Stand Model [6]. Neutrino astronomy – starting in the 1960’s with Ray Davis and theorist John Bahcall [10,12]– was sensitive enough to probe the central properties of the Sun. Recently, neutrino astronomy has reached the TeV-PeV energy range, with the hopes of elucidating properties of the high-energy sky. The first VHE neutrino events were first announced around 2012 [52]. Since then, many similar neutrino events – both HESE and through-going track – have been observed by the IceCube Antarctic Observatory [18]. The arrival directions of these neutrinos are consistent with an isotropic distribution, and their flavor ratio is consistent with a democratic ratio of $\nu_e : \nu_\mu : \nu_\tau \sim 1 : 1 : 1$. Both of which support extra-galactic neutrino sources.

Theoretical source models of VHE neutrinos date back at least to Waxman & Bahcall in the late 1990’s and their famous Bound [42], which relates the observed flux of VHE CRs and the maximum expected diffuse neutrino flux. The observed diffuse neutrino flux is remarkably close to this theoretical upper limit. Despite this, studies over the last five years have begun to rule out many of the top neutrino source model contenders such as conventional GRBs [46], Starburst and star-forming galaxies, galaxy clusters and groups [47], AGN [55,58] and TDEs [49,51,59,60]. VHE gamma-ray observations such as of the diffuse extra-galactic gamma-ray background and coincidence searches between
the prompt gamma-ray emission of GRBs or TeV AGN flares have been particularly constraining of such models.

The work presented in this dissertation tracks the evolution of our understanding of VHE neutrino source models from 2013-2017. In chapter 2 we discussed gamma-ray bright sources, HNe and SNe in Starburst and star-forming galaxies, and galaxy clusters and groups. Bechtol et al. 2017 [47] provided an analysis which indicated that $\sim 86\%$ of the diffuse gamma-ray background was produced by unresolved blazars, and therefore excluded a significant contribution from SBGs. Therefore, such gamma-ray bright CR calorimeters cannot provide a dominant contribution to the observed diffuse flux of VHE IceCube neutrinos.

In chapters 3 & 4 we presented models of gamma-ray dark/dim sources, including jetted TDEs and low-luminosity/choked jet GRBs. The former was found to be energetically disfavored, since the most optimistic diffuse neutrino flux expected from jetted TDEs was still a few orders of magnitude below the observed flux, while the latter type of events was found to be consistent with the IceCube neutrino observations.

The rich diversity of collapsar events from conventional Type-Ibc SNe, to trans-relativistic Type Ic-bl SNe and HNe, low-luminosity and conventional GRBs may be explained by the interaction of the relativistic jet with the stellar envelope or a circumstellar envelope. This latter scenario was postulated by Nakar 2015 [192] to explain gamma-ray and optical observations of llGRB 060218 [207], and was based on recent evidence the massive stars such as Wolf-Rayets shed high amounts of stellar material ($\sim 0.1 \, M_\odot$) in the last months/years before core-collapse. However, since ccSNe have not been observed in the days immediately following their explosion, the environment that makes up the circumstellar envelope is not well understood. We were able to show in chapter 4 that since VHE neutrinos are weakly interacting, they can therefore be used to probe the makeup of the circumnuclear envelope in such a scenario. Specifically, the time delay between the detection of non-thermal neutrinos and the prompt emission may be used to infer the properties of a collapsar jet and the environment surrounding its progenitor. A jet that successfully emerges would produce gamma-rays and neutrinos at the same time. A low-luminosity GRB in our scenario would produce neutrinos that are detectable $\sim 100 - 1000$ s before the prompt electromagnetic radiation. Finally, if a choked-jet collapsar results in an energetic Type Ic-bl SNe, the neutrino emission may proceed the detection of the electromagnetic signal – in this case thermal, optical photons – by from a few to tens of days.
In chapter 5, we performed a coincidence search between core-collapse SNe and up-going IceCube track events. This analysis is similar to earlier coincidence searches between IceCube neutrinos and electromagnetically bright GRBs [46] and AGN flares [55,57]. However, unlike conventional GRBs, the explosion time for ccSNe is not well determined, with a window spanning up to $\sim 15$ days [262]. This uncertainty is due to a lack of early observations of SNe light curves in the first hours to days after the explosion. However, there are indications that low-luminosity jets such as those believed to be in choked jet collapsars are favorable sites of CR acceleration and neutrino production, since their internal shocks are not radiation mediated, and can therefore accelerate CRs to high energies. Furthermore, because their electromagnetic radiation is attenuated by ambient matter and photons (through the $\gamma\gamma$ interaction), choked jet explosions also respect constraints set by gamma-ray observations. Using one year of ccSNe data and IceCube observations, we are able to place meaningful limits on the amount of energy such choked jets could deposit in VHE CRs, as well as the fraction of such events that contain a hidden jet pointed towards Earth. This analysis can be improved once the total seven years of IceCube is made available to the public. We show in Figure 5.7 that a non-detection from such an analysis could begin to rule out ccSNe as the sources of VHE IceCube neutrinos at the $\sim 90\%$ confidence level. The upgraded IceCube-Gen2 detector, as well as improved ccSNe observations from wide field of view surveys such as ASAS-SN and Zwicky Transient Facility [56] could also improve the significance of such an analysis.

6.2 Future Work

The search to elucidate the sources of VHE neutrinos will require a multi-disciplinary approach including advances in our understanding of particle interactions, astronomy (with both deep and wide field searches, especially in the radio, optical, X-ray, and gamma-ray bands), magnetohydrodynamics, low-number statistics, and cosmology. While progress in each of these fields has been made, there is still significant work to be done. The sources of neutrinos are also intimately connected to the yet to be determined sources of VHE and UHE CRs. In this last section, I present a short list of interesting future projects that will improve our understanding of the potential sources of VHE neutrinos. I will focus on three potential avenues of study below: 1.) Understanding the conditions which lead to a collapsar and the launching of a relativistic jet, including the conditions
that lead to a successful or choked jet. 2.) Improved implementations of heavy CR interactions (e.g., He, CNO, and Fe ions with matter and photons) and their effect on the resulting neutrino spectrum. 3.) Improved catalogs of potential VHE neutrino sources, especially in the radio, optical, and X-ray bands.

Relativistic astrophysical jets are ubiquitous in nature on a variety of scales, from stellar size microquasars to ~ kpc sized blazars and AGN. Although relativistic jets have been studied for decades, the properties of their launch, evolution, and emission mechanisms are still not well understood. The first successful model of non-thermal radiation from relativistic jets came from the famous “fireball” model (see [184] for a review), which considered a rapidly expanding, unmagnetized fireball of plasma and radiation launched by the magnetorotational interaction of an accretion disk with a central engine (e.g., a black hole or highly magnetized neutron star). This model could successfully explain important observational features of GRBs by assuming the prompt gamma-ray emission was produced by internal shocks generated by the collision of blobs of plasma in the jet. Furthermore, the major features of the fireball model could be calculated analytically.

Despite increasingly complex models of relativistic jets, (e.g., the ICMART model that incorporates strong magnetic pressures in the jet which affect its propagation [266]) and significant improvements in the field of computational magnetohydrodynamics (MHD) – from improved computer systems to streamlined computer MHD codes such as ATHENA [267]– the relationship between a progenitor star, the lunching of a jet, and the resulting non-thermal emission is not clear. That is to say, it is not known what stellar properties lead to an emerging jet GRB versus a simple ccSNe, though it is known that there is a connection between the two types of events.

VHE neutrino observations can both benefit from, and improve our understanding of the relationship between collapsars, ccSNe, and GRBs. Neutrinos are able to pass through the dense material in a stellar envelope to provide observations of the stellar interior and the internal structure of relativistic jets below the photosphere (i.e., the region that is opaque to photons because of scattering or absorption). Conversely, by understanding the types of explosions that produce relativistic jets, coincident searches between transient and neutrinos can be optimized. Note for example that the work presented in chapter 5 considers all ccSNe, while only a fraction $\approx 10\%$ are expected to have relativistic jets [152]. Improving simulations of jet propagation and our understanding of the structure of relativistic jets will increase the potency of neutrino astronomy.
Neutrino source models typically come in two major varieties: those that focus primarily on the particle physics of neutrino production, (e.g., Hüümer, Baerwald, and Winter), and those that focus primarily on the astrophysics of neutrino sources (e.g., Meszaros, Waxman, and Senno). The former camp has been able to successfully improve analytic or semi-analytic techniques to capture the most important particle physics features in their neutrino studies. Using semi-analytic over numerical methods is important because large, time-consuming parameter space searches must be performed to find good agreement with neutrino observations. For example, Hüümer et al. were able to increase the speed of their particle interaction code by $\sim 3000 \times$ with their Simulation-C over the same calculation using the SOPHIA particle interaction software [38]. However, to date these techniques only describe protonic CRs (i.e., $pp$ and $p\gamma$ interactions). There is considerable uncertainty in the composition of extra-galactic CRs, with some measurements favoring a mostly proton composition, while other prefer a mix of protons, He, CNO, and Fe ions. The composition of relativistic jets is not well known, but the massive stars that are the progenitors of collapsars tend to be metal rich CNO stars with their proton and Helium envelopes blown away (e.g., Wolf-Rayet stars). Therefore, future neutrino source models must take care to efficiently implement contributions from heavy CRs to the resulting neutrino flux by improving semi-analytic methods of hadronuclear and photohadronic interactions.

Many of the initially most promising astrophysical neutrino source models (e.g., GRBs, AGN, SBGs) have been ruled out as the dominant contributors to the measured diffuse neutrino flux. Because of this, new models of gamma-ray dim or hidden sources have risen to prominence lately (e.g., choked jet GRBs). These objects are by definition difficult to detect using conventional astronomy techniques. In the future, it is therefore important to increase our understanding of multi-messenger neutrino astronomy in two ways: 1.) by improving our theoretical understanding of the connection between potential neutrino sources and non-gamma-ray signals (e.g., radio, optical, and x-ray). For example as mentioned above by improving models of choked jet ccSNe to better understand the temporal relationship between neutrino signals and the optical emission from a SN explosion, as well as the types of SNe that will be neutrino bright. 2.) By increasing and improving the catalogs of gamma-ray dim sources, and strengthening the ties between neutrino astronomers and non-gamma-ray experiments, such as optical surveys.

To this end, multi-messenger collaborations such as AMON have signed Memoranda of Understanding with large field of view optical surveys such as Las Cumbres Observatory
Global Telescope, the Large Millimeter Telescope, and the MASTER Global Robotic Network. Large collaborations such as the Zwicky Transient Factory should also be considered, so that their large field of view surveys can be optimized to search for and quickly identify optical transients that have a high probability being neutrino bright (e.g., Type Ic-bl SNe). Because the catalogs of optical transients are so large already – and will increase even faster in the near future with ever larger all sky surveys – coincidence searches such as the one presented in chapter 5 of this work will quickly begin to rule out potential neutrino sources.

Since its beginning, neutrino astronomy has proven to be an indispensable window into the Universe. As is often the case, the most interesting results are those that do not agree with initial expectations such as the Homestake Mine experiment and the Solar neutrino problem, which took decades to resolve. It should not be surprising then that the sources of astrophysical neutrinos remain a mystery, despite significant theoretical work since the late 1990’s. While a resolution to this problem may be many years in the future, and present significant engineering, computational, and observational challenges, the study of astrophysical neutrinos will undoubtedly uncover surprising features of the way our Universe works, both at the largest (extra-galactic) and smallest (sub-atomic) scales.


[27]


[62] P. Meszaros, in Nuclear Physics B Proceedings Supplements, Center for Particle and Gravitational Astrophysics, Dept. of Astronomy & Astrophysics, Dept. of Physics, 525 Davey Laboratory, Pennsylvania State University, University Park, PA 16802, USA; >Based on a talk given at the Origin of Cosmic Rays: Beyond the Standard Model conference in San Vito di Cadore, Dolomites, Italy, 16-22 March 2014. This is not a comprehensive review of the topics in the title; it is weighted towards work in which I have been more personally involved (Nov. 2014), pp. 241–251.


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