THE THEORY AND PHENOMENOLOGY OF THE
HIGH-ENERGY AND TRANSIENT UNIVERSE

DISSERTATION

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ABSTRACT

The latest generation of telescopes and experiments have allowed for the study of previously-unexamined energy regimes and time domains, both in the Galaxy and beyond, with a variety of messenger particles. The research presented here has spanned from astronomy to high-energy physics, with a goal of bringing clarity to a variety of outstanding concerns in these fields, such as the origin of cosmic rays, the examination of the high-redshift Universe, and the natures of dark matter, supernovae, and neutrinos. I discuss how anomalous signals in the Galactic positron and electron spectra demand a local primary source (and why this may be the pulsar Geminga), the implications for reionization of the highest-redshift gamma-ray bursts, the impact of gamma-ray measurements on the dark matter parameter space, and the insight into the last days of massive stars given by the explosion in extragalactic transient detections.
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Physics is a terrible, terrible thing!

Walter Wada

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CHAPTER 1
INTRODUCTION

It is clear to you, I am sure, Lucilius, that no man can live a happy life, or even a supportable life, without the study of wisdom; you know also that a happy life is reached when our wisdom is brought to completion, but that life is at least endurable even when our wisdom is only begun.

Seneca, Epistle XVI

The fields of astrophysics and nuclear/particle physics have grown considerably over the past century, although the means by which they have done so have often differed greatly. While both are driven in large part by the new acquisition of data of increasing quality, in astrophysics these are typically obtained from the observation of distant physical systems of high complexity, as opposed to experiments in the laboratory in which controlled measurements can be made and systems broken down into their components at will in order to feed the fundamental needs of theorists. However, in recent years, these fields have been drawn together by a great collection of outstanding problems that require the services of both fields to gain insight into their inner workings.

The cardinal goal of this research has been to understand the content and dynamics of the high-energy and transient Universe. This has often involved making new connections between theoretical and observational advances within areas of physics.
and astronomy that would otherwise remain out of contact. These ventures have involved a number of collaborations with theorists, observers, and experimentalists in order to span a broad range of subjects. A selection of these form the central thesis of this Dissertation: the study of supernovae and dark matter, both of which trace back their origins in no small part to the pioneering work of Fritz Zwicky [1, 2, 3], and their relation to the problem of the origin of cosmic rays.

Our introduction begins with a brief review that touches upon the topics of core-collapse supernovae, gamma-ray bursts, and cosmic rays (produced both in our Galaxy and beyond) and their relation to high-energy neutrinos, gamma rays, and the search for dark matter. In the course of this overview, we outline a topic that will be mentioned a number of times throughout this Dissertation, from the high-energy emission of supernova remnants and pulsar wind nebulae to gamma-ray bursts: the production of charged particles via shock acceleration.

1.1 The Progenitors and Phenomenology of Supernovae

Over the past few thousand years, supernovae (SNe) have heralded both the death of kings and the impermanence of the heavens [4, 5, 6, 7]. We now know these appearances of “new stars” to be perhaps the most powerful transient events in the Universe, end products of the evolution of massive stars that typically release $\gtrsim 10^{51}$ erg of explosive kinetic energy and, in the case of core-collapse supernovae, $\gtrsim 10^{53}$ erg of neutrinos. In a groundbreaking series of papers [1, 2], Baade and Zwicky proposed that these phenomenal objects resulted from the “transition of a normal star” into a neutron star [1].

Supernovae are classified into two groups by their explosion mechanism (with further division into “Types” I/II, based on the absence/presence of hydrogen in their optical spectra). Supernovae of Type Ia are widely thought to originate from the
accretion of matter onto a white dwarf until a mass of $\sim 1.4 M_\odot$ is reached (here and throughout $M_\odot$ represents one solar mass). Upon reaching this mass (the Chandrasekhar limit), conditions are reached such that carbon burning ensues and leads to an uncontrollable thermonuclear explosion in which the central carbon/oxygen mixture is burnt to members of the iron group and other heavy elements.

Core-collapse supernovae (Types II, Ib, and Ic) are thought to be produced via an entirely different process, altogether, and exclusively by stars with initial masses exceeding a threshold of $M_i \gtrsim 8 M_\odot$ (masses below this experience a quieter death resulting in a white dwarf). Such stars quickly evolve through cycles of fusion to produce a massive core of material that cannot be fused for further energy gain [8]. Once a massive star produces a core that it can no longer profitably burn, the end comes very quickly. For $M_* \gtrsim 10 M_\odot$, iron-group elements (or O/Ne/Mg for $8 M_\odot \lesssim M_* \lesssim 10 M_\odot$ [8]) accumulate to a core mass of $\sim 1.5 M_\odot$, at which point gravity overwhelms support provided by radiation and electron degeneracy pressure, collapsing to a proto-neutron star with neutrinos carrying away the gravitational binding energy in a $\sim 10$ sec burst.

From the pioneering efforts of Fritz Zwicky in the early 20th century [13] to the automated surveys of today [14], observations have long indicated that many of these core-collapse events should be then “successful” in exploding as core-collapse supernovae [1] (see Fig. 1.1), with the type (Ib/c or II) largely determined by the initial stellar mass and metallicity [15]. Discernment of the explosion mechanism is one of the most important outstanding problems in nuclear astrophysics [16, 17, 18]. We will not discuss this mechanism further here (see Ref. [19] for a brief review), mentioning only that as a result of core collapse, an outgoing shock wave is believed to form that results in the ejection of the outer stellar envelope and the eventual optical supernova. The end result of this process is an ejection of $\sim 1 - 10 M_\odot$ of material by
Figure 1.1: The (likely-incomplete) sample [9] of local core-collapse SN distances. Denoted are the year that Fritz Zwicky retired and the distance to the Virgo Cluster. The shaded regions show the nominal reach (in order of decreasing darkness) of 32 kton (Super-K [10]), 1 Mton [11], and 5 Mton [12] neutrino detectors.

the supernova with $\sim 10^{51}$ erg of kinetic energy. Coincidentally, Type Ia supernova ejecta also carries $\sim 10^{51}$ erg, although with a fixed mass of $\sim 1.4M_\odot$ [20]. As we will be concerned with later, this makes the remnants of both types of supernovae prime candidates to accelerate cosmic rays, a possibility first advanced by Baade and Zwicky [2].

The work presented here concerns advances made within just the past few years. Indeed, there has been no shortage of novel extragalactic transients discovered in this period, allowing new types of study to be performed. In select cases, the opportunity has existed to examine the pre-supernova star in images taken before the explosion.
In this way, about a dozen SN progenitors have been studied (e.g., [21, 22, 23, 24]). The color and luminosity of each star can be used to infer the type and initial mass of the star, which is expected to be a “supergiant”, the phase in which the star vastly increases in size and luminosity during the remaining $\sim$ one million years of its life (see Fig. 1.2). However, a systematic survey of such supergiants in the local universe has yet to be performed.

In an attempt to uncover a completely new type of phenomenon via ground-based telescopes, we have recently proposed (and begun) the first survey for failed
supernovae [26]: bright stars that end with a whimper instead of a bang. Such an event may occur if a star with \( M_i \gtrsim 25M_\odot \) forms a black hole [15] before launching a shock wave to remove the outer envelope, which would have no recourse but to fall out of existence. This survey entails the systematic observation over ten years of nearly one million supergiants — the final stage in which massive stars spend their last \( \sim \) one million years as illustrated in Fig. 1.2 — so that many of these stars are expected to suffer a core collapse. Using deep imaging and image subtraction, it is possible to see whether these stars end as supernovae or form a black hole and simply fall out of sight. The latter is the scenario that we are most interested in either discovering or placing the first useful limits on the occurrence rate.

Another benefit of the aforementioned survey is the ability to identify the supergiant progenitors of supernovae occurring in nearby galaxies. In fact, when SN 2008S appeared in the nearby (\( \sim 5.6 \) Mpc) galaxy NGC 6946, it seemed an ideal candidate to find such a SN progenitor in deep archival images from the Large Binocular Telescope. Surprisingly, no massive star was seen (first panel of Fig. 3.2). However, we then found an interesting result: the first SN progenitor identified solely in the mid-infrared by using pre-explosion Spitzer Space Telescope images [23]. These results together implied a \( \sim 10M_\odot \) star surrounded by a shroud of dust expelled just thousands of years before exploding. A detailed census of analogously enshrouded stars in the very-near galaxy M33 recorded only a handful [27]. This, along with a similar transient/progenitor in NGC 300 (\( \sim 2 \) Mpc), led us to conclude that while the rate of explosions from such stars can be large, their progenitors are scarce, suggesting that this dusty phase lasts for only a brief period in the final evolutionary stages [27]. This has important implications for stars close to the critical white dwarf/core-collapse mass threshold.

As already mentioned, most of the energy involved with a core collapse actually
escapes in the form of neutrinos. These result from the cooling phase of the hot proto-neutron star, yielding a thermal spectrum with a temperature of a few MeV. Since photons cannot escape from the deep stellar core, the only way to observe the dynamics of a supernova explosion is to detect these neutrinos (or gravitational waves) produced during the core collapse, which travel essentially unhindered to us due to their weak interaction strength with matter. To do so requires extraordinarily large detectors; however, the payoff from such a detection would be tremendous. In Chapter 3, we will discuss the various information that would be learned from examining these neutrino signals. In particular, a water-based experiment at a scale of $\sim 5$ Megatons operating by detecting the Čerenkov light resulting from the energetic positron produced in an inverse-beta reaction $\bar{\nu}_e + p \rightarrow n + e^+$ would observe “mini-bursts” of neutrinos from supernovae in the nearby Universe about once per year [12]. The ability to identify every such core collapse within $\sim 5$ Mpc is in effect a continuous “death watch” of all massive stars in our corner of the Universe, making feasible studies that would otherwise be impossible. These include the abilities to independently detect prompt black hole formation, provide advance warning for searches for the “shock-breakout” preceding the supernova, and identification of “supernova impostors” by the absence of neutrinos.

1.2 The Sources of Cosmic Rays

The field of high-energy astrophysics began in earnest with the hot-air balloon measurements made by Victor Hess in 1912, in which he discovered that the rate of energetic charged particles passing through his hand-held detector increased with altitude up to at least five kilometers [28]. This implied an extra-terrestrial origin of this “cosmic radiation.” The search for the origin of these protons and electrons has continued ever since, and has involved considerable efforts in both observation and
theory to determine what can possibly be producing particles more energetic than even those in our largest laboratory accelerators. We can now estimate that, every second, about $10^{18}$ cosmic rays with energy greater than $10^9$ electron-Volts (eV) hit the Earth’s atmosphere (for comparison, a visual photon from the Sun has a typical energy of $\sim 1$ eV). This translates into an incredible amount of energy in the form of cosmic rays in the entire Milky Way. However, tracing them directly back to their sources is not possible since, being charged, they are easily deflected in the tangled magnetic fields of our Galaxy and arrive without memory of their starting point.

In recent years, new populations of gamma-ray sources have been discovered in our Galaxy in the TeV ($10^{12}$ eV) regime by the ground-based air Čerenkov telescopes HESS [30], MAGIC [31], and VERITAS [32], the all-sky gamma-ray telescope Milagro [33], and other experiments, with more now being found at lower energies by the Fermi Gamma-ray Space Telescope (hereafter Fermi) [34]. However, the details behind the production of these gamma rays remain largely unknown. It is likely, but not yet proven, that these sources are also the long-sought sites of Galactic cosmic-ray production [20]. To make this step requires demonstrating which particles are being accelerated.

Most of the energy in the form of cosmic rays in the Galaxy is carried by protons. Establishing that they are responsible for any of the gamma-ray sources would be a major step in the direction of locating their origins. However, this requires detecting a signature specific to their presence. Since protons generally only produce gamma rays through the decay of neutral pions produced in proton-proton scattering, the observation of high-energy neutrinos, produced in the decay of the charged pions that would also be present, is such a signal. Such an endeavor has long been known to be difficult [35]. However, among the known TeV gamma-ray sources, we established the most promising candidates for neutrino detection in the cubic-kilometer neutrino
telescopes now under construction, IceCube [36] and km3net [37], which would establish that these sources were cosmic proton accelerators [38, 39]. Detecting such neutrino sources can also allow for unique tests of particle physics models, since the energies of the particles produced greatly exceed what can be reached in laboratory experiments.

Several lines of reasoning converge on supernovae as the driver of Galactic cosmic-ray production. The strongest argument comes from a consideration of the energetics involved. The amount of energy supplied by the explosion of a few supernovae per century, each releasing $\sim 10^{51}$ erg in the form of ejected material, can account for the observed cosmic-ray energy density. Measurements of the energy spectrum of observed cosmic rays offer additional evidence. This spectrum very nearly follows a simple $E^{-2.7}$ power law below $\sim 3 \times 10^{15}$ eV, suggestive of a common origin for cosmic rays over a large range of energies (see Fig. 1.3). Also, the near-isotropy of cosmic-ray arrival directions helps to constrain the locations and injection spectra of sources that may contribute to the observed cosmic-ray flux [20]. How this all comes to pass depends ultimately on the mechanism of particle acceleration.

1.2.1 Shock Acceleration Theory

In the remnants of supernovae, fast-moving supernova ejecta ($v \sim 1-100 \times 10^3$ km s$^{-1}$) with total energy $\sim 10^{51}$ erg plows into the surrounding interstellar medium (ISM) forming a strong shock at the leading edge of the expanding material. It is at this shock that cosmic-ray acceleration is thought to occur. The most discussed acceleration mechanism is termed “diffuse shock acceleration”, which utilizes first order Fermi acceleration to repeatedly scatter particles crossing the shock front (e.g., Ref. [20] and
Energies and rates of the cosmic-ray particles

\[ E^2 \frac{dN}{dE} \ (\text{GeV} \ cm^{-2} \ s^{-1} \ sr^{-1}) \]

\[ E_{\text{kin}} \ (\text{GeV} / \text{particle}) \]

Figure 1.3: The observed energy spectrum of cosmic rays from \( 1 - 10^{12} \) GeV (from Ref. [29]).
references therein). One of the predictions of this model is a power-law particle spectrum, which is an attractive outcome in light of cosmic-ray observations. We will briefly discuss this scenario here.

The lives of supernova remnants typically follow a similar track, with variations depending upon the surrounding medium into which they expand (smooth ISM for Type Ia’s, wind-carved bubbles for massive CCSNe). Immediately following the supernova, an ejected shell of material begins to interact with its surroundings [40]. This process starts with what is called the “free-expansion phase”, in which the remnant expands with a nearly-constant velocity of

\[ v = \left( \frac{10E}{3M_e} \right)^{1/2} = 1.3 \times 10^4 \left[ \frac{E}{10^{51} \text{erg}} \right]^{1/2} \left[ \frac{M_e}{1 M_\odot} \right]^{-1/2} \text{km s}^{-1}, \tag{1.2.1} \]

where \( E \) is the kinetic energy of the ejecta, which has initial mass, \( M_e \) [41]. The remnant remains in this phase, sweeping up ISM, until it accumulates a mass roughly equal to \( M_e \). At this point it enters the “Sedov phase”, at a time [41]

\[ T_{\text{Sedov}} \sim 570 \left[ \frac{E}{10^{51} \text{erg}} \right]^{1/2} \left[ \frac{M_e}{1 M_\odot} \right]^{5/6} \left[ \frac{n_H}{0.05 \text{ cm}^{-3}} \right]^{-1/3} \text{yr}, \tag{1.2.2} \]

at which point the expansion is completely characterized by the explosion energy and the density of the ISM. Dimensional analysis arguments give the remnant radius as a function of time as [42] \( r \propto (E/\rho_0)^{1/5}t^{2/5} \), which has been demonstrated experimentally [43]. Eventually, the remnant will cool to the point that energy is efficiently lost to radiative processes, which quickly slow the expansion and lead to the remnant’s dissolution into the ISM.

Any model that attempts to account for the observed cosmic rays must: provide the observed cosmic-ray injection power; explain the cosmic-ray spectral index and composition; be able to accelerate particles up to at least \( 10^{18} \text{ eV} \); and do all of the above while preserving the observed isotropy. The above SNR cycle is quite amenable to developing such a model. We will first check the energetics, closely
following Hillas [20]. The cosmic-ray production rate per unit area in the Galactic plane can be found as

\[ q_{\text{CR}} = n_{\text{CR}} c \frac{\mu}{\langle g \rangle}, \]

(1.2.3)

where \( n_{\text{CR}} c \) is the cosmic-ray flux, \( \mu = 3.0 \times 10^{-3} \text{ g cm}^{-2} \) is the gas mass per unit area of the Galactic plane, averaged between \( r = 4-12 \text{ kpc} \), and \( \langle g \rangle \) is the grammage (average column density) traversed [20]. The grammage depends upon the energy-dependent rate at which cosmic rays escape the Galaxy, with \( \langle g \rangle = 34 R_{\text{GV}}^{-0.6} \text{ g cm}^{-2} \) from spallation studies or \( \langle g \rangle = 15 R_{\text{GV}}^{-1/3} \text{ g cm}^{-2} \) expected from a Kolmogorov spectrum of turbulence, where \( R_{\text{GV}} \) is rigidity in units of GV [20]. Substituting these terms (using \( R_{\text{GV}}^{-0.6} \)), along with the measured cosmic-ray proton energy density, \( 1.4 \times 10^{-10} (E/\text{GeV})(4\pi/c)(E/\text{GeV})^{-2.64} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \), we arrive at a total cosmic-ray power of \( 1.6 \times 10^{-13} (E/\text{GeV})^{-1.04} \text{ W cm}^{-2} \text{ GeV}^{-1} \) [20]. Note that the needed injected spectrum would then have to be of the form \( E^{-(2.64-0.6)} = E^{-2.04} \) (whereas a Kolmogorov spectrum implies \( E^{-(2.64-0.33)} = E^{-2.31} \)). Integrating over the 10 – 100 GeV decade of energy in the region between \( r = 4-12 \text{ kpc} \) (and multiplying by 1.5 to account for nuclei in this range), and assuming 2 supernovae occur in this region per century [20, 44], requires each explosion to contribute \( \sim 0.2 \times 10^{51} \text{ erg of cosmic-ray energy, or } \sim 20\% \) of the expected total kinetic energy.

With the energetics reasonably satisfied, we can now discuss an acceleration mechanism. The key feature of the SNR for cosmic-ray production is the strong shock. Because the velocity of the expanding shell is much greater than the sound speed in the ISM, a shock front is established in front of a contact discontinuity, with a density ratio of \( \sigma = \rho_s/\rho_{\text{ISM}} = (\gamma + 1)/\gamma - 1 = 4 \) (with a specific heat ratio \( \gamma = 5/3 \)) on either side of the shock (neglecting gas and cosmic-ray pressures for the moment) [42]. This rapid compression results in heating of the swept-up gas. The basic principle of particle acceleration in this environment is that magnetic fields in the shock region
can scatter energetic particles resulting in an incremental increase in energy after each scattering (the “first order Fermi process”). Thinking in the rest frame of the shock, a particle with a diffusive scattering mean free path somewhat larger than the compressed region can cross this boundary many times, each time gaining energy. In this fashion, an exponential energy gain can occur for a given particle, although, for the entire system, energy gains will be mediated by particles escaping the region after their mean free path grows too large, resulting in an overall particle spectrum.

Following Longair [42], for each scattering, let \( E = \beta E_0 \) be the resulting average particle energy, with \( P \) probability of remaining in the system. Beginning with \( N_0 \) particles, after \( k \) scatterings, \( N = N_0 P^k \) particles have energies \( E = E_0 \beta^k \). Using logs, we find \( \ln \left( \frac{N}{N_0} \right)/\ln \left( \frac{E}{E_0} \right) = \ln P/\ln \beta \), which gives \( N/N_0 = (E/E_0)^{\ln P/\ln \beta} \). The differential spectrum can be found from this integral quantity, giving

\[
\frac{dN}{dE} \propto E^{\ln P/\ln \beta - 1}, \tag{1.2.4}
\]

reducing the problem to finding \( P \) and \( \beta \). For a shock speed, \( v_s \), and \( \sigma = 4 \), swept-up gas is seen to approach the shock at \( v_1 = v_s \) and exit the system at a relative velocity of \( v_2 = v_s/4 \), as conservation of mass through the shock boundary requires \( \rho_{\text{ISM}} v_1 = \rho_s v_2 \). The gas behind the shock thus travels at velocity \( 3v_s/4 \) relative to the stationary ISM. Energetic particles ahead of the shock (assumed isotropized by random scatterings), see the gas behind the shock approaching at \( 3v_s/4 \), gain an amount of energy upon crossing the shock, and then are re-isotropicized. These particles, upon diffusing back toward the shock, will then see the ISM approaching at a relative velocity of \( 3v_s/4 \) and gain energy upon this crossing. Considering the angular distribution of particles approaching the shock, the average energy gain per crossing is \( \langle \Delta E/E \rangle = v_s/2c \), giving \( \beta = 1 + v_s/c \) for one “round trip” [42].

At any time, the particle flux crossing the shock boundary is \( nc/4 \). Meanwhile, particles are being “lost” from the shock region at a rate of \( nv_s/4 \). The overall
fractional loss rate is then \(v_s/c\) (which is relatively small), which is used to find \(P = 1 - v_s/c\). We can thus find the ratio \(\ln P / \ln \beta = \ln (1 - v_s/c) / \ln (1 + v_s/c) = -1\), giving us the famous \(dN/dE \propto E^{-2}\). This ideal equation will be modified by adding in the pressures of the accelerated particles and thermal gas in the shock environment. Due to diffusion, cosmic-ray pressure will be the same on both sides of the shock, while thermal pressure will only be on the “downstream” side. The cosmic-ray pressure leads to a subshock, which compresses material before reaching the shock. If we let \(P_{\text{therm}} = (2/3)w_{\text{therm}}\) and \(P_{CR} = (1/3)w_{CR} = (1/3)\alpha w_{\text{therm}}\), where \(w\) is the energy density and \(w_{CR} = \alpha w_{\text{therm}}\), the speed of gas entering the shock will be decreased as \(U = v_s - P_{CR} / (\rho_{\text{ISM}} v_s)\) before the shock and \(U_2 = v_s - (P_{CR} + P_{\text{therm}}) / (\rho_{\text{ISM}} v_s)\) after the shock [20]. The density contrast at the subshock is now

\[
\sigma = 1 + \frac{\Sigma - 1}{1 + \alpha/2},
\]

(1.2.5)

where \(\Sigma = \rho_s / \rho_{\text{ISM}} = v_s / U_2 = (8 + 7\alpha)/(2 + \alpha)\) is the overall increase in density [20]. Choosing \(\alpha = (1, 1.5, 2)\) corresponds to \(\Sigma = (5, 5.29, 5.5)\) and \(\sigma = (3.67, 3.45, 3.25)\). Lower energy particles, having smaller mean free paths, will only experience the lower preshock \(\sigma\), whereas the highest energy particles can sample the entire density contrast, \(\Sigma\). These values for \(\sigma\) and \(\Sigma\) yield spectral indices of \((2.125, 2.225, 2.333)\) and \((1.75, 1.7, 1.67)\), respectively [20]. Thus, in this scenario, hardening of the spectrum is expected at the highest energies.

When examining the maximum attainable particle energies, one can consider Hillas’s simple rule-of-thumb for acceleration in an object

\[
E_{\text{max}} < B R v_s \propto (E/10^{51}\text{erg})^{1/2} B M_e^{-1/6} \rho_{\text{ISM}}^{-1/3},
\]

(1.2.6)

where \(R\) is the radius and \(B\) is the magnetic field of said object (replaced by \(3B_{\text{upstream}}\) in our case) [20]. Assuming \(B_{\text{upstream}} \sim 3 \mu G\) with the other parameters similar in magnitude to those previously discussed, leads to \(E_{\text{max}} \sim 10^{14} Z\) eV, which falls short
of the knee. To get around this barrier, Bell and Lucek proposed to use Alfvén waves generated by cosmic rays diffusing ahead of the shock to amplify the external magnetic field \([45]\). Rather large, self-generated magnetic fields can be formed by this mechanism, with \(B \propto v_s \rho_{\text{ISM}}^{1/2}\) \([46]\). This leads to a modification of the maximum attainable energy such that

\[
E_{\text{max}} \propto \left(\frac{E}{10^{51}\text{erg}}\right) M_e^{-2/3} \rho_{\text{ISM}}^{-1/6},
\]

which is now independent of \(B\) and weakly dependent upon \(\rho_{\text{ISM}}\). This formulation suggests that a common position of the knee may possible for a wide range of objects.

### 1.2.2 Cosmic-ray Electrons and Positrons

Interest in the cosmic-ray electron content of the Galaxy arises from many quarters of astrophysics, from studies of Galactic magnetic field structure \([47]\), the acceleration of particles by supernova remnants and pulsars, and even dark matter annihilation. Understanding their spectrum is key towards achieving these and other goals (hereafter, I will refer to the sum of electrons and positrons as “electrons” unless noted); however, prior to this past year, knowledge of this quantity even at Earth could be called, at best, uncertain. Fortunately, needed clarity has recently been brought in the range of energies from \(\sim 20\) GeV up to \(\sim 1\) TeV by Fermi’s electron measurements \([48]\) and at higher energies still by HESS \([49, 50]\) as shown in Fig. 1.4.

The sum of the components that result in this total spectrum arriving at Earth remains an open question, and it is likely that this will vary across the Galaxy. A contemporaneous measurement of the relative fraction of cosmic-ray positrons to electrons by the PAMELA space mission yielded a striking rise with energy \([51]\) (see Fig. 5.3). This is not expected in conventional models (e.g., \([52]\)), where positrons are only produced as secondary by-products from cosmic-ray proton interactions with
the ISM. Thus, a need exists for a primary source of positrons. Many have suggested dark matter annihilation to fill this void (see below and Chapter 6).

In order to make proper judgments in this field, an understanding of both the multitude of data and the many high-energy astrophysical processes in the Galaxy is essential. This is particularly important in realizing the potential of the full set of Fermi data, both gamma rays and cosmic-ray electrons, especially in relation to the many other observations relevant to this purpose. These include diffuse observations, such as in the radio, microwaves, and soft X-rays, and of high-energy sources in TeV
gamma rays (by ground-based air Čerenkov telescopes), X-rays, and (soon) neutrinos. Additionally, efforts to push electron measurements to higher energies are advancing, although the signals expected near and beyond the HESS data (end of Fig. 1.4) have largely been a mystery.

Conventional thought concerning the cosmic-ray electron spectrum would favor a steep drop at an energy near 1 TeV. This can be understood by considering that cooling rate of high-energy electrons via synchrotron and inverse Compton scattering losses increases with energy, so that the distance “horizon” from which an electron with a given observed energy could have originated rapidly shrinks at high energies. The historical lack of known sources at distances sufficiently near to result in electrons with multi-TeV energies reaching Earth left the prevailing thinking unchallenged. However, recent observations suggest that this prejudice leads to an incomplete picture of the high-energy activity of the Galaxy.

A driving force behind this change has been the ∼ 100 GeV measurements from PAMELA, which introduced the need for a local source of primary positrons. There are few possible ways to accomplish this. Based on the observation by the Milagro experiment of diffuse, multi-TeV gamma rays surrounding the nearby (∼ 200 pc) pulsar Geminga [53], we found that this implied that this source was a powerful cosmic-ray accelerator within the past few hundred thousand years [54]. This established the nearest source of energetic cosmic rays to Earth, which produced a flux of electrons and positrons that may fully account for the positron excess. Furthermore, this suggests that the cosmic-ray electron spectrum can be dominated by a local source even at energies of a few hundred GeV (Fig. 1.4).

This emerging picture of the electron spectrum stresses the importance of incorporating transient acceleration events in the recent history of the Galaxy. Electrons are quite sensitive in this regard, in stark contrast to cosmic-ray protons, which can
spread throughout the Galaxy with minimal energy losses and smooth over their origins. A necessary step towards this goal is achieving a consistent treatment of time-dependent cosmic-ray propagation. As we discuss in Chapter 5, we have recently begun addressing several of the outstanding issues involved with encouraging results [55].

1.3 Gamma-Ray Bursts

Gamma-ray bursts (GRB) are perhaps the grandest spectacles in astrophysics [56, 57]. (Throughout this Dissertation, we refer only to “long” gamma-ray bursts, those with observed duration exceeding $\sim 2$ seconds.) Their connection with core-collapse supernovae [58, 59] indicates that their progenitors are very-massive, short-lived stars. Studies of GRB host galaxies demonstrate that these galaxies tend to have a low metallicity [60, 61, 62]. When compared to a naive expectation based on the galaxies’ luminosities alone, direct spectroscopic measurements reveal metallicities that are significantly lower than one would have expected [61]. Although the metallicity of the GRB progenitor is not directly measured via this technique, the metallicity of the star’s surroundings should be indicative of the star itself.

A natural explanation may be found by considering the properties of GRB host galaxies in the context of the single-star collapsar model [63, 64]. In this model, GRB progenitors are rapidly-rotating Wolf-Rayet stars that undergo a core-collapse event that produces a black hole (or possibly a rapidly-rotating neutron star) [15]. After collapse, this rapid rotation allows for the formation of highly relativistic jets which, when viewed on-axis, are observed as a burst [65]. Observationally, all known supernova counterparts of GRBs are Type Ic, with the implication that the dying star lacked an outer hydrogen/helium envelope [58]. The winds of Wolf-Rayet stars,
which are typically the cause of the loss of this envelope, are known to increase in strength with stellar metallicity (particularly iron) [66].

Importantly, in this wind-induced loss process, angular momentum (which is particularly important for forming jets [67]) is lost along with mass [68]. This loss of angular momentum can be avoided if the progenitor has a very low metal content. In addition to having weaker winds, a rapidly-rotating, metal-poor massive star can avoid the production of an envelope altogether by completely mixing its interior, which results in the hydrogen being circulated into the core and burned [64]. This would be impossible if the star was not metal-poor, as stellar mixing is expected to be inhibited by a high metal content [69]. If these indications are true, a number of implications follow.

1.3.1 The Cosmic GRB Rate

Just as the core-collapse supernova rate density seems to follow the cosmic star formation history (SFH) [70], the same might be expected of the cosmological GRB rate density [71]. However, there is mounting evidence that GRBs are not an unbiased tracer of the SFH [72, 73], in particular due to the tendency of their host galaxies to be subluminous [72] and metal-poor [60]. This has been demonstrated for GRB hosts both locally [61] and at cosmological distances [62], which suggests that low metallicity is a key ingredient in GRB production.

As we just discussed, a rapidly rotating star, as is typically required in the GRB models, can retain much of its original mass and angular momentum if it is metal-poor [64]. An anti-correlation with metallicity would imply that the cosmological GRB rate should peak at a higher redshift than normal supernovae [74, 75]. Simply put: the metallicity of the universe decreases with redshift, which would imply a stronger evolution of the GRB rate density than would be expected from the SFH.
alone. To test this directly, as we discuss in detail in Chapter 4, we examined the large set of GRBs with firm redshifts discovered in the last few years by the *Swift* satellite [76]. We found that GRBs actually *do not* follow this simplest assumption [77], instead finding that the GRB rate appears to be $\sim$ four times higher than expected at $z \sim 4$. This may confirm the picture of the progenitors of GRBs forming from metal-poor gas that allows them to maintain the angular momentum needed to launch jets following core collapse.

### 1.3.2 The History of Cosmic Star Formation

The history of star formation in the universe is of intense interest to many in astrophysics, and it is natural to pursue pushing the boundary of observations to as early of times as possible. Indeed, measuring the rate at which stars were created in the first few billion years of the Universe yields fundamental tests of cosmological evolution and galaxy formation. Foremost amongst these is determining when and how the Universe’s diffuse hydrogen transitioned from a neutral state — that resulted from cooling following the Big Bang — into the mostly-ionized gas observed today [78]. Young, massive stars are leading candidates for accomplishing this; however, it is easy to understand that measurements become increasingly difficult as we look back further into the past as the majority of galaxies become apparently, and likely intrinsically, fainter.

Our understanding of this history is ever increasing, with a consistent picture now emerging from direct observations up to redshift $z \gtrsim 4$, as summarized in Fig. 1.5 [79]. The cosmic star formation rate measurements from the compilation of Ref. [70] are shown, along with recent high-$z$ measurements based on observations of color-selected Lyman Break Galaxies (LBG) [80, 81] and Lyα Emitters (LAE) [82]. Much current
Figure 1.5: The cosmic star formation history. The compiled SFR data (light circles) and fit (dotted line) of Ref. [70] are shown, as well as newer high-z data (the LAE points only sample Ly$\alpha$ Emitters). The results of Ref. [79], as inferred using bright Swift gamma-ray bursts, are shown with dark diamonds. The solid line is a new high-z fit [79].

interest is on this high-z frontier, where the primeval stars that may be responsible for reionization reside. Due to the difficulties of making and interpreting these measurements, different results disagree by more than their quoted uncertainties.

Instead of inferring the formation rate of massive stars from their observed populations, one may directly measure the SFR from their death rate, since their lives are short. While it is not yet possible to detect ordinary core-collapse supernovae
at high $z$, gamma-ray bursts have been detected to $z \gtrsim 8$ [83, 84]. The brightness of GRBs across a broad range of wavelengths makes them promising probes of the SFH. In the last few years, Swift [76] has spearheaded the detection of GRBs over an unprecedentedly-wide redshift range, including many bursts at $z \gtrsim 4$. As just discussed, examination of the Swift data reveals that GRB observations are not tracing the SFH directly, instead implying some kind of additional evolution [77, 85, 86, 87].

GRBs can still reveal the overall amount of star formation, provided that we know how the GRB rate couples to the SFR. In Chapter 4, we proceed to use the portion of the SFH that is sufficiently well-determined to probe the range beyond $z \simeq 4$. We do this by calibrating the many bursts observed in $z \simeq 1 - 4$ to the corresponding SFR measurements, and by taking into account the known bias between observations of the GRB rate relative to the SFR. This eliminates the need for prior knowledge of the absolute conversion factor between the SFR and the GRB rate and allows us to properly relate the GRB counts at $z \simeq 4-8$ to the SFR in that range. Additionally, we make use of the estimated GRB luminosities to exclude faint low-$z$ GRBs that would not be visible in our high-$z$ sample, i.e., to compare “apples to apples”. The recent detections of GRBs at $z \sim 8$ allowed us to both push these estimates into the epoch of reionization [88] and examine properties of the host galaxies themselves [89]. Very-recent observations from the upgraded Hubble Space Telescope Wide Field Camera 3 appear to confirm our basic conclusions [90, 91]: stars born in rather small galaxies were sufficient to drive reionization.

1.3.3 Gamma-ray Bursts and Ultrahigh-Energy Cosmic Rays

Few classes of astrophysical objects can possibly account for the observed cosmic-ray spectrum at the ultrahigh energies of $\gtrsim 10^{20}$ eV [92]. The measured anisotropy in their arrival directions requires them to originate from beyond our Galaxy [93].
Active galactic nuclei (AGN) have long been considered as possible ultrahigh-energy cosmic-ray (UHECR) sources [94]. Relatively recently, a potential connection between gamma-ray bursts and UHECR has been explored. A number of models have been proposed to utilize the $\sim 10^{51} - 10^{52}$ erg of energy in their ultra-relativistic environment to accelerate protons to energies $\gtrsim 10^{20}$ eV [95, 96, 97].

To be considered as a viable source of UHECR, gamma-ray bursts must have the ability to both accelerate protons to energies $\gtrsim 10^{20}$ eV and generate a cosmic-ray flux adequate to explain the observed spectrum at the highest energies in Fig. 1.3. In conventional GRB models, a portion of the kinetic energy of a relativistically expanding fireball (with Lorentz factor $\Gamma \sim$ few hundred) is converted into internal energy [98]. Electrons are accelerated inside this jet by internal shocks and subsequently produce gamma-rays through synchrotron and inverse-Compton processes [99].

Protons may also be shock-accelerated in a similar fashion. In the internal shock model, the shocks that accelerate protons are expected to be only mildly relativistic in the wind rest frame, resulting in an $\sim E^{-2}$ spectrum. In order to efficiently accelerate protons to ultrahigh energies, the time scale of acceleration should be shorter than both the wind expansion time (to allow for an adequate period of confinement in shocked regions) and the proton energy loss time scale. The former sets the ratio of magnetic field and electron energy densities to order unity, which is necessary in order to account for gamma-ray emission from synchrotron emission boosted to the observer’s frame. The latter imposes an upper limit on magnetic field strength (and lower limit on Lorentz factor) [95].

While UHECR must be extragalactic in origin, the interaction of $\gtrsim 10^{19.5}$ eV cosmic rays with the cosmic microwave background (CMB) results in a dramatic increase in energy losses through photopion production ($p\gamma \rightarrow N\pi$), the so-called Greisen, Zatsepin, and Kuzmin (GZK) Effect [100], with the consequence that any
UHECR produced beyond a limited “horizon” of $\sim 50$ Mpc will ultimately deposit most of its initial energy into neutrinos and gamma rays. In examining the possibility of a GRB origin, we found that the observed rapid rise in the history of the cosmic GRB rate (discussed below) would result in a flux of diffuse “GZK neutrinos” larger than previously expected [86]. To conclude Chapter 4, we examine why our predicted flux is large enough to be detected by ongoing experiments, so that a positive signal could favor UHECR production by GRBs.

1.4 High-Energy Processes in the Search for Dark Matter

A variety of astronomical observations, from the level of galaxies and clusters [3, 101, 102] to the cosmic microwave background [103], strongly indicate that a majority of the mass in the universe is in the form of dark matter. Many theories predict dark matter candidates over a wide range of masses. However, since its discovery, all searches for dark matter that have succeeded have done so via its gravitational effects, while all inquiries into its particle nature have failed in their own way. The wide-range of particle candidates suggest numerous possible signals, from scatterings off nuclei [104, 105] to the byproducts of self-annihilations or decays [106, 107, 108].

To date, this search for dark matter has typically been a “diagnosis of exclusion”, with little direct evidence to go on save its energy density [109]. Observations of the products of dark matter annihilations or decays would provide invaluable information towards establishing the underlying particle physics and be a landmark discovery. The interest in such matters is well illustrated by the amount of theoretical activity stimulated by the excess of positrons seen by PAMELA. However, as previously discussed, understanding of high-energy astrophysical processes, both in the Galaxy and the rest of the Universe, is necessary in order to make a proper judgment of any prospective signal.
As we discuss in Chapter 6, one attractive model proposed that the difficult-to-explain isotropic MeV gamma-ray background [110] could be produced through the radiative decays of long-lived dark matter particles throughout the history of the Universe [111]. However, a distinctive signal would be expected from our own Galactic dark matter halo. Starting from data originally used to limit the presence of photon line signals in the Galactic Center region from the SPI gamma-ray spectrometer on the INTEGRAL satellite [112], along with the determination of the isotropic diffuse photon background by SPI [113], COMPTEL [110], and EGRET [114, 115] observations, we placed strong, model-independent constraints on late-decaying dark matter by showing that no more than $\sim 5\%$ of the unexplained MeV background can be produced by late dark matter decays either in the Galactic halo or cosmological sources [116].

Additionally, the multitude of models introduced subsequent to PAMELA (e.g., [117]) also have distinct observational implications if they are to account for the positron excess. These models generically must, of course, result in a large number of positrons (and electrons), which would radiate much of their energy in the form of gamma rays en route to Earth. However, even in a model that would be otherwise quite difficult to constrain — where most of the positrons arise from within self-bound “clumps” of dark matter that compose the substructure of the Milky Way’s halo — we found that the expected gamma-ray signals would already be in tension with a number of present data [118]. We address this in greater detail in Chapter 6, where we also discuss how new measurements from Fermi and other gamma-ray telescopes can soon considerably strengthen these conclusions.
CHAPTER 2
PROSPECTS FOR TEV NEUTRINO ASTRONOMY

Consider this fact also. Those who have never attained their mental independence begin, in the first place, by following the leader in cases where everyone has deserted the leader; then, in the second place, they follow him in matters where the truth is still being investigated. However, the truth will never be discovered if we rest contented with discoveries already made. Besides, he who follows another not only discovers nothing but is not even investigating. What then? Shall I not follow in the footsteps of my predecessors? I shall indeed use the old road, but if I find one that makes a shorter cut and is smoother to travel, I shall open the new road. Men who have made these discoveries before us are not our masters, but our guides.

Seneca, Epistle XXXIII

2.1 Introduction

The field of TeV gamma-ray astronomy is exploring energy regimes that have been, until recently, out of reach to astrophysicists.\footnote{The content of this Chapter is based in large part on our work in Refs. [38, 39].} Yet, even as the catalog of TeV sources continues to grow, it is still debated whether the observed gamma rays are produced...
leptonically, through the inverse Compton scattering of energetic electrons on ambient photons \((e^- \gamma \rightarrow \gamma e^-)\), or hadronically, though neutral pion decay \((\pi^0 \rightarrow \gamma \gamma)\). Air Čerenkov telescopes, such as HESS [30], can measure a source spectrum with high precision for \(E_\gamma \sim 1 - 10\) TeV. One might hope that as we probe higher energy gamma rays, indications of their true origin would be revealed. However, at energies \(\gtrsim 10\) TeV, the difficulties of gamma-ray astronomy become more pronounced, due to the low statistics of the quickly declining signal spectra.

It is well established that a distinctive feature of a pionic (hadronically-produced) gamma-ray spectrum is an accompanying flux of neutrinos [119, 120, 121]. These neutrinos originate from the decay of charged pions \((\pi^+, \pi^-)\), which are produced in approximately equal numbers with neutral pions in proton-proton scattering. Relative to gamma-ray telescopes, the new km\(^3\) neutrino telescopes [36, 37] will have several advantages that result in improved performance at the highest energies. The rapidly falling atmospheric neutrino background, rising neutrino-nucleon cross section \((\sigma_{\nu N} \sim E_\nu)\), and increasing muon range, which effectively expands the (already large) detector volume \((R_\mu \sim \ln E_\mu)\), all help to amplify the diminished flux at these energies. In fact, as the background quickly becomes negligible in the TeV range, the detection of any high energy neutrinos from a source could significantly indicate a hadronic production mechanism.

Neutrino telescopes have capabilities far beyond breaking the degeneracy between leptonic and hadronic production models. Spectral features in the highest energy regime, especially an expected cutoff (related to the maximum accelerated proton energies), shall not remain inaccessible to observation. Measurement of the energies of neutrino-induced muons and showers, which are related to the original charged pion energy, can probe the source proton spectrum in a complementary manner to gamma-ray observations, which effectively measure the neutral pion spectrum [122].
Figure 2.1: Ranges over which gamma-ray and neutrino observations can measure a source spectrum ($E_\gamma \simeq E_\pi/2$, $E_\nu \simeq E_\pi/4$). Squares are HESS measurements of Vela Jr., fit by a power law (dashed line). Neutrino telescopes can probe higher energies to distinguish between pionic spectra with cut-offs (dot-dashed lines) by measuring the $\nu_\mu$-induced muon and shower spectra. Muon rates of $N(E_\mu > 1\, \text{TeV}) \sim 3 - 6\, \text{yr}^{-1}$ correspond to the dot-dashed lines, illustrating the sensitivity of neutrino observations [38].
The sensitivities of these two independent approaches, including the regime where they coincide, are illustrated in Fig. 2.1. The ability to accurately measure neutrino-induced muon spectra greatly improves the prospects for detecting point sources, as the harder source spectra dominate the atmospheric background above $\sim 1$ TeV. While muon tracks have better angular resolution ($\lesssim 1^\circ$), neutrino showers ($\sim 10^\circ$ in water) more faithfully trace the spectrum. Shower observations, which measure the $\nu_e$ and $\nu_\tau$ fluxes, when combined with muon data, also allow for the study of the ratio of neutrino flavors arriving from a source [123].

Considering the latest observations of Galactic TeV sources, we calculate the corresponding spectra of detectable neutrino-induced muons and showers, assuming only that the observed gamma-ray spectra are pionic, for a range of possible high energy cutoffs in the spectra [38, 39, 124]. Relative to analyses in which the total numbers of signal and background events are counted (e.g., Ref. [125]), a maximum likelihood analysis would have much more power. For example, below $\sim 1$ TeV, a single event has a much greater probability of being background than at 10 TeV, where the source signal is dominant. Spectra allow for such an approach, which takes full advantage of experimental data, in studying the high energy behavior of TeV sources. Resolving hadronic activity at these extreme energies would provide clear evidence concerning the sources of Galactic cosmic rays [20].

Atmospheric cosmic-ray showers give rise to high rates of down-going muons, forcing a neutrino telescope to search for up-going muons resulting from neutrino interactions. IceCube is well-situated to utilize the high resolution of these $\nu_\mu$-induced muons in observing northern-sky sources. However, a detector is needed in the northern hemisphere to accurately locate southern-sky neutrino sources, although IceCube may also measure shower rates from particularly bright sources. Together, IceCube and a $1\text{ km}^3$ Mediterranean detector will provide full-time coverage of the entire sky,
a feature distinct to neutrino telescopes. The combined observations from these detectors can be used to study compound objects, like the Vela complex, by: (1) Discovering neutrino sources through high-resolution $\nu_\mu$-induced muons; (2) Confirming agreement with gamma-ray observations in the low energy regime; (3) Examining previously unexplored energies using muons and showers together.

2.2 Promising TeV Neutrino Sources

2.2.1 Vela Region

Of particular interest amongst prospective neutrino sources, the shell-type supernova remnant (SNR) Vela Jr. (RX J0852.0–4622) is one of the brightest objects in the southern TeV sky. The hard, intense TeV gamma-ray spectrum of Vela Jr., best explained as being pionic in nature, makes it an intriguing object to study [128]. Additionally, the morphology of the TeV emission may make this one of the most interesting astrophysical neutrino sources. HESS has also measured a TeV spectrum from Vela X, the pulsar wind nebula (PWN) associated with the larger Vela SNR, and advanced a leptonic origin [129]. However, if this spectrum is instead pionic, as proposed in Ref. [130], the accompanying neutrino flux would be easily detectable. The non-detection of neutrino events from such an intense source would allow for a significant test of leptonic production in a short period of observation.

2.2.2 Galactic Center Diffuse Emission

HESS has also recently discovered a region of diffuse TeV emission from the Galactic Center ridge [131]. This source has several interesting features which merit further investigation. The large extent of the emission, hardness of the spectrum, high gas density (which is well-correlated with the TeV emission) and strong magnetic fields in
the region leave very little doubt that this spectrum is pionic [131]. In addition, the
total flux from this region is actually about twice as intense as that of the previously
discovered source coincident with Sgr A* [132]. Measurement of the accompanying
neutrino flux would provide an independent confirmation of the means of production,
something which has never been possible.

2.2.3 Unknown Knowns and Known Unknowns

We also examine the neutrino detection prospects of other known TeV sources, in-
cluding SNR RX J1713–3946, the spectrum of which has been measured twice by
HESS and determined to likely be pionic [133, 134]. The calculated \( \nu_\mu \)-induced muon
fluxes based upon the two different sets of HESS data illustrate the effects of differ-
ent high energy spectral assumptions. Recently, HESS has reported the discovery of
four TeV sources in the Galactic Plane which have no apparent counterparts at other
wavelengths [135, 136].

2.2.4 Cygnus Region

Recent Milagro observations of the Cygnus region have revealed both diffuse TeV
gamma-ray emission and a bright and extended TeV source, MGRO J2019+37, which
seems to lack an obvious counterpart at other wavelengths. Study of this curious
object also promises to provide important clues concerning one of the Milky Way’s
most active environments. First, to gain insight into this mysterious source, we
consider its relation to known objects in both the Cygnus region and the rest of
the Galaxy. Second, we find that a simple hadronic model can easily accommodate
Milagro’s flux measurement (which is at a single energy), as well as other existing
observations spanning nearly seven orders of magnitude in gamma-ray energy. Third,
since a hadronic gamma-ray spectrum necessitates an accompanying TeV neutrino
flux, we show that IceCube observations may provide the first direct evidence of a Galactic cosmic-ray accelerator.

## 2.3 Basics of Neutrino Detection

In high energy $p$-$p$ scattering, $\pi^+$, $\pi^-$, and $\pi^0$ are produced in nearly equal numbers [137]. Gamma rays are the result of the decay $\pi^0 \to \gamma\gamma$, while neutrinos originate from the $\pi^+ \to \mu^+\nu_\mu \to e^+\bar{\nu}_e\nu_\mu$ and $\pi^- \to \mu^-\bar{\nu}_\mu \to e^-\nu_\mu\bar{\nu}_e\nu_\mu$ decay channels. The ratio of neutrinos to photons from pion decay is easily found. From charged pion decay, the resulting initial neutrino flavor ratio, $\nu_\epsilon : \nu_\mu : \nu_\tau$, is 1 : 2 : 0. During the traversal of astrophysical distances, vacuum neutrino oscillations transform this ratio to 1 : 1 : 1. In neutrino telescopes, neutrinos and antineutrinos are practically indistinguishable. We can then consider their sum, $\nu + \bar{\nu}$, and average cross section in all calculations. All further references to neutrinos will imply $\nu + \bar{\nu}$. Thus, for equal pion multiplicities, each photon from $\pi^0$ decay corresponds to one neutrino of each flavor ($N_\gamma = N_{\nu_e} = N_{\nu_\mu} = N_{\nu_\tau} = N_{\nu}$). The typical energy of the neutrinos resulting from these decays is $\sim 1/2$ of the gamma-ray energy from $\pi^0$ decay. The resulting $\nu + \bar{\nu}$ spectrum is then shifted, relative to the source gamma-ray spectrum of $d\Phi_\gamma/dE_\gamma = \phi_\gamma E_\gamma^{-\Gamma}$, as

$$
\frac{d\Phi_\nu}{dE_\nu} = \left(\frac{1}{2}\right)^{\Gamma-1} \phi_\gamma E_\nu^{-\Gamma} = \phi_\nu E_\nu^{-\Gamma},
$$

(2.3.1)

where we consider each neutrino flavor separately.

We will generally refer to the normalization of the gamma-ray spectrum as $\phi_\gamma$, the differential photon flux at 1 TeV (in TeV$^{-1}$ cm$^{-2}$ s$^{-1}$), which is commonly used in gamma-ray astronomy. Typical source gamma-ray spectra have $\phi_\gamma \sim 5 - 20 \times 10^{-12}$ TeV$^{-1}$ cm$^{-2}$ s$^{-1}$ and $\Gamma \sim 1.8 - 2.4$. For greater detail on the relationship between the initial proton spectrum and the resulting product spectra, see Ref. [138]. Pions
can also be produced, in a different ratio, in $p$-$\gamma$ scattering, but this is generally only important at much higher energies. A neutrino flux may be observed through neutrino-induced muons and electromagnetic (and hadronic) cascades, which are referred to as showers. We shall consider both methods in further detail.

### 2.3.1 Muon Detection

Our analysis of muon detection will be limited to $\nu_\mu$ charged-current (CC) events with an observable final energy of $E_\mu > 0.5$ TeV, which can be detected through their Čerenkov emission with an angular resolution of $\lesssim 1^\circ$ [139]. It is important to recognize that the relevant quantity in these events is the measured energy of the muon, which can be reconstructed from radiative losses in the detector [140]. As muons can be produced far outside of the detector, with energy loss prior to entering, results of calculations are given in terms of $E_\mu$ (at the detector) instead of $E_\nu$.

We can consider $\nu_\mu$-induced muons produced both inside and outside of the instrumented volume of the detector. *Contained* event rates can be found by combining the neutrino flux, detector mass, and $\sigma_{CC}$ for neutrino-nucleon scattering. We use the average of the $\nu$-$N$ and $\bar{\nu}$-$N$ cross sections, as computed in Ref. [141], and $\langle y(E_\nu) \rangle$ from Ref. [142]. We consider a detector that is entirely composed of, and enclosed by, water (ice), with a km$^2$ detector area, which is a reasonable approximation at these energies. Using an IceCube-like effective detector area, such as computed in Ref. [143], would lead to the lower energy muon spectrum being slightly suppressed. An actual km$^3$ detector will be situated on top of solid rock, which will enhance the upward-going neutrino-induced muon rates relative to our calculations. Also, muons produced via $\nu_\tau$ CC interactions, which contribute to the total muon flux through the $\tau^+ \rightarrow \mu^+\nu_\mu\nu_\tau$ decay channel, will not be considered [137, 144]. Our formulas follow
those of Gaisser [137], with appropriate approximations. Considering the uncertainties in both the gamma-ray spectra and neutrino telescope performance, which may not be fully understood until the actual detectors are calibrated, as well as the low statistics, our calculations are sufficient for the scope of this work.

Using these parameters, the spectrum of contained $\nu_\mu$-induced muon events is calculated as

$$
\left( \frac{dN_\mu}{dE_\mu} \right)_{\text{cont}} = \kappa V_{\text{det}} \frac{d\Phi_\nu}{dE_\nu} e^{-E_\nu/E_{\nu}^{\text{cut}}} \sigma_{\text{CC}}(E_\nu) e^{-\tau_\oplus},
$$

(2.3.2)

where $V_{\text{det}}$ is the detector volume, $E_{\nu}^{\text{cut}}$ is an assumed exponential cutoff in the neutrino spectrum, and the term $\kappa = N_A \rho T (1 - y(E_\nu))^{-1}$ takes into account observation time (T), normalization of the muon spectrum, and the molar density of water. The energy of the produced muon is related to the original neutrino energy as $E_\mu = (1 - y(E_\nu)) E_\nu$. The term $e^{-\tau_\oplus}$, with $\tau_\oplus = N_A \lambda_\oplus \sigma_{\text{tot}}(E_\nu)$, accounts for neutrino attenuation due to scattering within Earth. Here, $\sigma_{\text{tot}} = \sigma_{\text{CC}} + \sigma_{\text{NC}}$ and $\lambda_\oplus$ is the average column depth (in cm.w.e.) of Earth based upon the declination of the source [142]. This attenuation factor only becomes important at $E_\nu \gtrsim 10$ TeV, and varies with declination, as discussed in Ref. [122].

A high energy muon born outside of the detector may still be detectable when it enters the instrumented detector volume. To find the propagation range of a muon, we assume an average continuous muon energy loss of

$$
\frac{dE}{dX} = -\alpha - \beta E,
$$

(2.3.3)

where $\alpha = 2.0 \times 10^{-6}$ TeV cm$^2$ g$^{-1}$ and $\beta = 4.2 \times 10^{-6}$ cm$^2$ g$^{-1}$ [140, 144]. Integrating the average energy loss results in a range for a muon of initial energy $E_\mu$ of

$$
R_\mu(E_\mu, E_\mu^f) = \frac{1}{\beta} \ln \left[ \frac{\alpha + \beta E_\mu}{\alpha + \beta E_\mu^f} \right],
$$

(2.3.4)

where $E_\mu^f$ is the energy of the muon as it enters the detector. As the muon range (typically a few km) increases with energy, so does the effective volume of the detector.
The observed through-going spectrum, accounting for the probability of observing a muon entering the detector with energy $E_\mu$, can then be calculated as

$$\left(\frac{dN_\mu}{dE_\mu}\right)_{\text{thru}} = \frac{N_A \rho T A_{\text{det}}}{\alpha + \beta E_\mu} \int_{E_\mu}^{\infty} dE_\nu \frac{d\Phi_\nu}{dE_\nu} e^{-E_\nu/E_{\nu,\text{cut}}} \sigma_{\text{CC}}(E_\nu)e^{-\tau_\odot} \tag{2.3.5}$$

where $A_{\text{det}}$ is assumed to be 1 km$^2$. It should be noted that our use of the total $\sigma_{\text{CC}}$ and $\langle y \rangle$ may overestimate the muon signal rates by $\sim 20 - 30\%$ (perhaps more for spectra harder than $E^{-2}$). However, the atmospheric background will also be overestimated by a similar, even slightly larger amount.

### 2.3.2 Shower Detection

For bright point sources or low surface brightness extended sources, shower events caused by $\nu_e, \nu_\tau$ CC interactions can be used to effectively determine the source spectrum. To detect shower events, they must be at least partially enclosed within the instrumented volume. The shower signal is dominated by $\nu_e$ and $\nu_\tau$ CC events, which transfer $\sim 100\%$ of the original neutrino energy to the shower ($E_{\text{sh}} \approx E_\nu$). The measurable shower spectrum, given in terms of the observed shower energy, $E_{\text{sh}}$,

$$\left(\frac{dN_{\text{sh}}}{dE_{\text{sh}}}\right)_{\text{CC}} = 2N_A \rho T V_{\text{det}} \frac{d\Phi_\nu}{dE_\nu} e^{-E_\nu/E_{\nu,\text{cut}}} \sigma_{\text{CC}}(E_\nu)e^{-\tau_\odot}, \tag{2.3.6}$$

then effectively traces the original neutrino flux, as the neutral-current (NC) events from all three flavors, which only contain $E_{\text{sh}} = \langle y_{\text{NC}} \rangle E_\nu$ (in addition to the smaller $\sigma_{\text{NC}}$), only account for $< 10\%$ of the total signal.

Combined shower and muon observations would determine the $\nu_\mu/(\nu_e + \nu_\tau)$ ratio arriving from a source. This measured flavor ratio has many applications, as detailed in Ref. [123]. The atmospheric $\nu_e$ flux is only about $\sim 1/20$ of the $\nu_\mu$ background, which itself only adds to the shower background through the weaker NC channel, since $\nu_\mu$ CC interactions are identifiable by the resulting muon tracks. We will not consider the prompt background, which is the only source of atmospheric $\nu_\tau$, as it only becomes
important at very high energies [145]. These facts help offset the lower angular resolution of these events ($\lesssim 10^\circ$ is expected for a Mediterranean detector [146], $\lesssim 25^\circ$ for IceCube [139]), which sets the background event rate. The improved shower resolution of a Mediterranean detector, as compared to IceCube (which is naturally limited by light scattering in ice), increases the value of showers in determining the high energy properties of a source.

### 2.3.3 Neutrino Spectroscopy

To demonstrate the ability of a neutrino telescope to differentiate between various spectral properties, we first consider a general object with a fixed $\phi_\gamma = 20 \times 10^{-12}$ TeV$^{-1}$ cm$^{-2}$ s$^{-1}$ (at 1 TeV) and assume a source declination of $\delta = +10^\circ$. The resulting up-going muon fluxes for a range of spectral indices and high energy exponential cutoffs, for one year of IceCube observation, are shown in Fig. 2.2. In the top panel, we consider spectral indices with $\Gamma = 1.8, 2.0, 2.2, \text{ and } 2.4$, with $E_{\nu}^{\text{cut}} = 50$ TeV. Note that the differences are not entirely due to differing integrated neutrino fluxes, as the rising $\sigma_{\nu N}$ results in higher event rates for harder spectra. In Ref. [122], similar results were obtained by normalizing to the integrated flux above 1 TeV. The bottom panel illustrates the effect of 100, 50, 25, and 10 TeV neutrino spectrum exponential cutoffs for a fixed index ($\Gamma = 2.1$). IceCube is well-suited to isolate a similar point source located in the northern sky through $\nu_\mu$-induced muons within a short period of operation. Such a source would also be observable through shower events by a km$^3$ Mediterranean detector. The improved shower detection capabilities of a Mediterranean detector would allow for a more direct measurement of the arriving neutrino spectrum, adding greatly to the science that can be extracted from TeV sources.

All of our calculations are done on an empirical basis using the average measured parameters of gamma-ray source spectra. All high energy exponential cutoffs that
Figure 2.2: A demonstration of the sensitivity of neutrino telescopes to varying spectral characteristics. Shown is the (base-10 log) differential \((\nu_\mu + \bar{\nu}_\mu)\)-induced muon spectra for a pionic source spectrum with \(\phi_\gamma = 20 \times 10^{-12}\) TeV\(^{-1}\) cm\(^{-2}\) s\(^{-1}\). In the upper panel, the lines of decreasing height correspond to spectral indices of \(\Gamma = 1.8, 2.0, 2.2,\) and \(2.4\), with an exponential cutoff \(E_{\nu}^{\text{cut}} = 50\) TeV. In the lower panel, the lines of decreasing height correspond to neutrino spectrum cutoffs of 100, 50, 25, and 10 TeV, with a fixed spectral index of \(\Gamma = 2.1\). Each line corresponds to one year of observation [38].
we will consider are given in terms of \( E_{\text{cut}}^\nu = \frac{1}{2} E_{\text{cut}}^\gamma \), to allow for a range of calculations to be compared with future experimental data. Angle-averaged FLUKA computations of atmospheric neutrinos past \( E_{\nu} = 1 \) TeV are used to determine backgrounds [147]. For ease of comparing the expected signal to background and assessing total rates, we will mainly provide integrated rates above a given measured energy. Because the atmospheric spectrum is very steeply falling at these energies (\( \sim E_\nu^{-3.5} \)), it is often advantageous to use a low energy event cutoff \( \gtrsim 1 \) TeV when calculating significance. Our method can also yield the measured muon and shower spectra directly. Spectral information should be used for maximum likelihood analysis of observed neutrino events in a neutrino telescope, to make the best use of measured energies in determining detection significance. This spectral information requires consideration of the energy resolution of the detector. In IceCube, the energy resolution for muons and showers is expected to be \( \sim 20\% \) and \( \sim 10\% \) in the logarithm of the energy, respectively [139, 148], which does not significantly affect our conclusions.

Note that we present our results for an assumed muon effective area of 1 km\(^2\), and assumed muon angular resolution of \( \lesssim 1^\circ \), both independent of energy. These results are close to the results of detector simulations, especially \( \gtrsim 1 \) TeV, the primary area of interest [139]. Below 1 TeV, these assumptions are certainly overly optimistic; however, in that range, the atmospheric background is dominant, and so accuracy is of less importance. In regards to the angular resolution, our assumptions are somewhat too conservative (in particular, the improved resolution at higher energy will help reject background). Part of our intent is to show the likely signals, and to encourage the experimentalists to optimize their detector designs in order to achieve the necessary sensitivity to observe Galactic TeV sources. These calculations are not meant to replace a more detailed study of neutrino source detectability, which would
Figure 2.3: Vela Jr. (RX J0852.0–4622)—Muons: Integrated ($\nu_\mu + \bar{\nu}_\mu$)-induced muon rates above a given measured muon energy. The long dashed (blue), dot-dashed (red), and short dashed (gold) lines correspond to neutrino spectrum exponential cutoffs of 50, 25, and 10 TeV, respectively. The dotted line shows the expected atmospheric background in a 7 deg$^2$ bin, as discussed in the text. Rates are for one year of operation in a km$^3$ Mediterranean detector [38].

include time-dependent source locations, zenith-angle-dependent atmospheric backgrounds, angle-dependent detector sensitivities, event reconstruction methodologies, stochastic muon energy losses, source modeling, etc. Each prospective source should be subjected to such a full Monte Carlo simulation by the IceCube and Mediterranean collaborations.
2.4 The Vela Complex

Of the many known TeV gamma-ray sources, the shell-type SNR Vela Jr. (RX J0852.0–4622) is one of the most interesting. This southern-sky source has been observed at gamma-ray energies exceeding 10 TeV by HESS. Their analysis is suggestive of a hadronic origin for the gamma-ray spectrum [128]. This source is very bright in gamma rays ($d\Phi/dE = 21 \times 10^{-12} (E/\text{TeV})^{-2.1} \text{ TeV}^{-1} \text{ cm}^{-2} \text{ s}^{-1}$), with well-defined regions of gamma-ray emission. Shell-type SNRs are considered to be the most likely sites of Galactic cosmic-ray proton acceleration. As these source proton spectra are expected to be cut off near the knee ($\sim 3 \times 10^{15} \text{ eV}$), a cutoff should also be present for gamma-rays and neutrinos at a lower energy scale [138]. Thus, we calculate the expected neutrino-induced muon rate assuming a pionic spectrum with several neutrino spectrum exponential cutoffs (50, 25, and 10 TeV), as shown in Fig. 2.3. When considering emission from the entire source extension, we must accept atmospheric background from a $\sim 7 \text{ deg}^2$ area. Comparing the source and background muon rates, this SNR is expected to be significantly detectable with a km$^3$ Mediterranean detector. For the higher cutoffs, discovery may be possible in only a few years.

A number of features make Vela Jr. unique among prospective TeV neutrino sources. The TeV emission is observed to originate from several regions which are separated by $\sim 2^\circ$. In addition to detecting the presence of neutrinos, the intensity of the measured muon spectrum may provide adequate statistics to allow accurate location of neutrino-production sites. Thus, this unique gamma-ray morphology allows for the possible construction of a neutrino map of Vela Jr. This capability would enable a comparison of the production mechanism and high energy activity in different regions of the same source. When observing features of the neutrino source, the total atmospheric neutrino background enclosed is lower, which increases their individual
detectability. An example of such a map can be seen in Ref. [149]. Further HESS observations should provide improved resolution of the TeV emission.

Additionally, the brightness of the source may also make it possible to observe appreciable numbers of showers. As a consequence of the detected showers directly tracing the neutrino flux, spectral information may be found (even with a slightly higher background). Fig. 2.4 shows the rate of showers for one year of observation, compared to the irreducible atmospheric background. With enough observation time, it may be possible to determine the neutrino flavor ratio from this source by observing the $\nu_\mu$ spectrum with a km$^3$ Mediterranean detector and by also utilizing IceCube to effectively increase the shower volume (since shower events must be located within the detector) in measuring the $\nu_e + \nu_\tau$ flux.

Vela Jr. is coincident with a region of the larger Vela SNR. HESS has also reported the discovery of TeV gamma rays originating from the Vela PWN, referred to as Vela X (of no relation to the HMXB Vela X-1) [129]. Vela X has a very hard spectrum, which is best fit with a spectral index of $\Gamma = 1.45$ with an exponential cutoff beginning at $E_{\gamma}^{\text{cut}} = 13.8$ TeV. In their discovery paper, HESS concludes that this emission is produced by an inverse Compton mechanism. An independent analysis has considered the possibility that this spectrum is produced hadronically, with an accompanying flux of neutrinos [130]. This source is sufficiently separated from Vela Jr. to independently test for neutrino emission through muons. If we assume the entire measured flux is pionic, then the expected $\nu_\mu$-induced muon rate is $N(E_\mu > 1$ TeV) $\sim 4.5$ yr$^{-1}$, in a km$^3$ Mediterranean detector, which would be relatively easy to measure. While our calculation may somewhat overstate the muon rate for such a hard spectrum, nevertheless, the nondetection of a significant neutrino event excess would indicate leptonic production.
Figure 2.4: Vela Jr. (RX J0852.0–4622)—Showers: Integrated neutrino-induced shower rates above a given measured shower energy, with signal lines as in Fig. 2.3. Note that a different energy scale is used for showers, as measurable shower events are assumed to have a 1 TeV energy threshold. Background rates are for a circle of 10° radius, corresponding to the shower angular resolution in a km$^3$ Mediterranean detector. A significant reduction in background results from improved angular resolution. Rates are for one year of operation [38].

Considering the large apparent size of the Vela SNR shell ($\sim 8^\circ$), any TeV gamma-ray emission originating from it may be too diffuse to detect directly. Searching for it in neutrinos is a possibility. A northern-hemisphere detector would isolate any neutrino point source spectra located in the region through $\nu_\mu$-induced muons. Showers have the ability to observe this entire region simultaneously, raising the intriguing prospect of utilizing IceCube and a km$^3$ Mediterranean detector in concert to discover and study bright, but highly extended TeV neutrino sources.
2.5 Galactic Center Region

The HESS discovery [131] of the region of Galactic Center diffuse emission (GCD) is important for neutrino, gamma-ray, and cosmic-ray astrophysics. The spectrum of the GCD was measured over a very large region spanning the Galactic coordinates $|l| < 0.8^\circ$, $|b| < 0.3^\circ$ with $d\Phi/dE = 1.73 \times 10^{-8} (E/\text{TeV})^{-2.29} \text{ TeV}^{-1} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}$.

A number of large, dense ($n_H \sim 10^4 \text{ cm}^{-3}$) molecular clouds are known to fill this region [150]. Considering the vast extent of this hard emission ($\sim$ few hundred pc across), along with the high target density and magnetic fields in the region, the only reasonable production mechanism for these gamma rays is neutral pion decay. The close morphological correlation between the TeV emission and the gas distribution is particularly compelling [131]. This TeV gamma-ray spectrum implies a local cosmic-ray spectrum that is much harder and 3 – 9 times denser than that of the CR flux as measured at Earth. The likely source of these cosmic-ray protons is Sgr A, which hosts the remnant of a recent supernova [151, 152]. These pioneering conclusions were reached in Ref. [131].

In order to calculate the expected $\nu_\mu$-induced muon rate, we must first take into account the extent of the GCD source. The muon angular resolution of a neutrino telescope at these energies ($\sim 1^\circ$) covers the entirety of the GCD. In effect, this region can be treated as a neutrino point source. Integrating over the $\sim 1 \text{ deg}^2$ source region, the total photon flux is $d\Phi/dE = 5.2 \times 10^{-12} (E/\text{TeV})^{-2.29} \text{ TeV}^{-1} \text{ cm}^{-2} \text{ s}^{-1}$. We define the differential flux at 1 TeV as $\phi_{\text{GCD}} = 5.2 \times 10^{-12} \text{ TeV}^{-1} \text{ cm}^{-2} \text{ s}^{-1}$. We will not consider showers from the GCD; however, any such study should also take into account a similar diffuse emission around Sgr B [131].

We first consider the conservative assumption that the GCD spectrum exhibits an exponential neutrino spectrum cutoff $E_\nu^{\text{cut}} = 20 \text{ TeV}$. As a cosmic-ray production site, we also consider a higher cutoff of 50 TeV. Even though its angle-integrated
flux is larger, the GCD was revealed only after subtracting the previously discovered TeV emission from Sgr A* and a distant SNR [132]. The TeV gamma-ray flux from Sgr A*, which has $\phi_\gamma \sim 0.5 \phi_{GCD}$, can reasonably be assumed to also be pionic, as the spectral index ($\Gamma = 2.2$) is close to the GCD value [132, 153]. As the two sources are coincident, we combine their expected muon rates. The declination of the source, along with the latitude of the detector, determine the fraction of the time that the source is below the horizon. For the Galactic Center this is $\sim 0.7$. The signal rates, when compared to the expected number of atmospheric events, suggest that this source is near the threshold of discovery for the lifetime of a km$^3$ Mediterranean detector. Discovery may be hastened by utilizing maximum likelihood methods in comparing the expected signal with background for each muon event with a measured energy, as higher energy events have a much greater significance. Such an analysis is only possible with spectral information.

2.6 A Multitude of TeV Sources

Numerous discoveries of new gamma-ray sources should be expected from the HESS, MAGIC and VERITAS air Čerenkov telescopes and Fermi [31, 32, 34]. However, the characteristics of the source spectra above $\sim 10$ TeV are difficult to observe. RX J1713–3946 is also a shell-type SNR, first observed in the TeV regime by CANGAROO [154]. It has since been observed twice by HESS with greatly improved statistics. We consider both HESS observations to illustrate the potential of neutrino telescopes to determine the high energy behavior of TeV sources. The flux was initially reported up to 10 TeV as a power law ($\Gamma \sim 2.19$), with analysis favoring hadronic production [133]. Subsequent observations revealed a spectrum that extends to at least 40 TeV with steepening past $\sim 10$ TeV that can be fit by either an energy dependent spectral index or an exponential cutoff [134]. Previous neutrino
Table 2.1: Unidentified HESS Sources: Integrated ($\nu_\mu + \bar{\nu}_\mu$)-induced muon rates, assuming a pionic spectrum, for a 1 TeV $E_\mu$ threshold and neutrino spectrum exponential cutoff $E_\nu^{\text{cut}}$. For one year of $\text{km}^3$ Mediterranean detector operation, accounting for observable time below the horizon. The HESS differential flux at 1 TeV is given in terms of $10^{-12} \text{ TeV}^{-1} \text{ cm}^{-2} \text{ s}^{-1}$, with spectral index $\Gamma$.

<table>
<thead>
<tr>
<th>Source</th>
<th>$\phi_\gamma$</th>
<th>$\Gamma$</th>
<th>$E_\nu^{\text{cut}}$ (TeV)</th>
<th>$N_\mu(&gt;1 \text{ TeV})$</th>
</tr>
</thead>
<tbody>
<tr>
<td>HESS J1303–631</td>
<td>4.3</td>
<td>2.44</td>
<td>10</td>
<td>0.3</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>50</td>
<td>0.5</td>
</tr>
<tr>
<td>HESS J1614–518</td>
<td>7.0</td>
<td>2.46</td>
<td>10</td>
<td>0.5</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>50</td>
<td>0.8</td>
</tr>
<tr>
<td>HESS J1702–420</td>
<td>2.5</td>
<td>2.31</td>
<td>25</td>
<td>0.3</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>50</td>
<td>0.4</td>
</tr>
<tr>
<td>HESS J1708–410</td>
<td>1.5</td>
<td>2.34</td>
<td>10</td>
<td>0.1</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>50</td>
<td>0.2</td>
</tr>
</tbody>
</table>

studies of this source were inspired first by the CANGAROO detection [126] and later by utilizing the more accurate 2004 HESS spectrum along with a more detailed calculation [127]. We analyze the $\nu_\mu$-induced muon spectra from these two data sets independently, assuming $E_\nu^{\text{cut}} = 50 \text{ TeV}$ for the 2004 data. As seen in Table 2.2, the integrated rates above 1 TeV are similar. However, at higher energies calculations based upon the old power law spectrum exhibit a deviation from the new data. From the 2004 data, we find $N(E_\mu > 10 \text{ TeV}) \sim 0.3 \text{ yr}^{-1}$, while the 2005 data yields a rate only $\sim 1/10$ of this, which highlights the sensitivity of neutrino observations to the spectrum at the highest energies. Based on either spectrum, this SNR should be detectable in a $\text{km}^3$ Mediterranean detector.
Amongst the HESS catalog of TeV sources are several with no observed counterpart at any other wavelength. Explaining the origin of these sources is a theoretical challenge. It is possible that they represent a non-trivial new class of astrophysical objects. With so little being known about them at present, the observation of neutrinos from these sources would yield invaluable insight into their nature. The first of these sources is HESS J1303–631 [136]. Chandra observations revealed no likely x-ray counterpart [155]. This source is always visible to a Mediterranean detector. Three sources were also found in the HESS survey of the Galactic Plane that remain unidentified: HESS J1614–518, HESS J1702–420, and HESS J1708–410 [135]. The spectral information and resulting integrated $\nu_\mu$-induced muon rates from these sources are given in Table 2.1. It has been proposed that HESS J1303–631 is a remnant of a Galactic gamma-ray burst (GRB), with a harder ($\Gamma \sim 2.2$) and more extended emission than has been observed [156]. From the given predictions of the proposed hadronic source model, we would expect $N(E_\mu > 1 \text{ TeV}) \sim 2.1 \text{ yr}^{-1}$, which is potentially verifiable with a km$^3$ Mediterranean detector. From the measured spectra, detecting these sources individually may be difficult; however, the stacking method may prove advantageous. Stacking effectively increases the total observation rate by combining the individual event rates from identical sources. Stacking of the observed muon data may improve the prospects of determining whether neutrinos are present and help unveil the mystery of these unknown sources.

Other Galactic TeV sources with probable counterparts may also be examined as potential neutrino sources. Many new sources were found as a result of the HESS inner Galactic Plane survey, some of which may be SNRs [135]. MAGIC and VERITAS observations will reveal sources in the outer regions of the Plane, which has yet to be examined with an instrument of their sensitivity. Assuming that the gamma-ray spectra already observed by HESS are pionic, the calculated neutrino event rates for
<table>
<thead>
<tr>
<th>Source</th>
<th>$\phi_{\gamma}$</th>
<th>$\Gamma$</th>
<th>$E_{\nu}^{cut}$ (TeV)</th>
<th>$N_{\mu}(&gt;1\text{ TeV})$</th>
</tr>
</thead>
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<tr>
<td>Vela Jr.</td>
<td>21.0</td>
<td>2.1</td>
<td>10</td>
<td>3.1</td>
</tr>
<tr>
<td>(RX J0852.0–4622)</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>25</td>
<td>4.9</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>50</td>
<td>6.1</td>
</tr>
<tr>
<td>GC Diffuse</td>
<td>5.2</td>
<td>2.29</td>
<td>20</td>
<td>0.5</td>
</tr>
<tr>
<td>(+ GC Source)</td>
<td></td>
<td></td>
<td>50</td>
<td>0.7</td>
</tr>
<tr>
<td>RX J1713.7–3946</td>
<td>15.0</td>
<td>2.19</td>
<td>50</td>
<td>2.8</td>
</tr>
<tr>
<td></td>
<td>20.4</td>
<td>1.98</td>
<td>6</td>
<td>2.2</td>
</tr>
<tr>
<td>Vela X</td>
<td>9.0</td>
<td>1.45</td>
<td>7</td>
<td>4.5</td>
</tr>
<tr>
<td>Crab (IceCube)</td>
<td>33.0</td>
<td>2.57</td>
<td>50</td>
<td>2.7</td>
</tr>
<tr>
<td>HESS J1514–591</td>
<td>5.7</td>
<td>2.27</td>
<td>25</td>
<td>0.9</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>50</td>
<td>1.1</td>
</tr>
</tbody>
</table>

Table 2.2: HESS Sources With Counterparts: Integrated ($\nu_{\mu} + \bar{\nu}_{\mu}$)-induced muon rates, assuming a pionic spectrum, for a 1 TeV $E_{\mu}$ threshold and neutrino spectrum exponential cutoff $E_{\nu}^{cut}$. For one year of km$^3$ Mediterranean detector operation (unless noted), accounting for observable time below the horizon. The HESS differential flux at 1 TeV is given in terms of $10^{-12}$ TeV$^{-1}$ cm$^{-2}$ s$^{-1}$, with spectral index $\Gamma$.

a number of sources are shown in Tables 2.2 and 2.3. Several of these have sufficient rates to consider as possible neutrino sources independently. However, it would be difficult to significantly detect dimmer sources over the atmospheric background. Stacking, in combination with a maximum likelihood spectral analysis, may be used to examine the entire group, or subgroups, of the lower-flux sources.
<table>
<thead>
<tr>
<th>Source</th>
<th>$\phi_\gamma$</th>
<th>$\Gamma$</th>
<th>$E_{\nu}^{\text{cut}}$ (TeV)</th>
<th>$N_\mu(&gt; 1 \text{ TeV})$</th>
</tr>
</thead>
<tbody>
<tr>
<td>HESS J1616–508</td>
<td>6.0</td>
<td>2.35</td>
<td>10</td>
<td>0.5</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>50</td>
<td>0.9</td>
</tr>
<tr>
<td>HESS J1632–478</td>
<td>5.5</td>
<td>2.12</td>
<td>10</td>
<td>0.8</td>
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<tr>
<td></td>
<td></td>
<td></td>
<td>50</td>
<td>1.5</td>
</tr>
<tr>
<td>HESS J1634–472</td>
<td>2.0</td>
<td>2.38</td>
<td>10</td>
<td>0.2</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>50</td>
<td>0.3</td>
</tr>
<tr>
<td>HESS J1640–465</td>
<td>3.0</td>
<td>2.42</td>
<td>10</td>
<td>0.2</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>50</td>
<td>0.4</td>
</tr>
<tr>
<td>HESS J1745–303</td>
<td>2.5</td>
<td>1.8</td>
<td>10</td>
<td>0.5</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>50</td>
<td>1.2</td>
</tr>
<tr>
<td>HESS J1804–216</td>
<td>4.7</td>
<td>2.72</td>
<td>10</td>
<td>0.1</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>50</td>
<td>0.2</td>
</tr>
<tr>
<td>HESS J1813–178</td>
<td>2.7</td>
<td>2.09</td>
<td>10</td>
<td>0.3</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>50</td>
<td>0.5</td>
</tr>
<tr>
<td>HESS J1825–137</td>
<td>6.0</td>
<td>2.46</td>
<td>25</td>
<td>0.4</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>50</td>
<td>0.5</td>
</tr>
<tr>
<td>HESS J1837–069</td>
<td>5.0</td>
<td>2.27</td>
<td>10</td>
<td>0.3</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>50</td>
<td>0.6</td>
</tr>
</tbody>
</table>

Table 2.3: A continuation of Table 2.2.

2.7 Cygnus in the TeV

The Cygnus region is one of the most prominent features of nearly every Galactic skymap, across many orders of magnitude in energy. In gamma rays, sources have
been known to populate this region since the time of COS B [157]. In subsequent years, EGRET discovered both diffuse [158] and point-like MeV-GeV emission [159], followed by the HEGRA observation of an unidentified TeV gamma-ray source [160], TeV J2032+4130. Recently, at even higher gamma-ray energies, the Cygnus region has been discovered yet again by Milagro [33]. The gamma-ray flux measured by Milagro includes a large (∼ 15° × 10°) diffuse region and a new, unidentified TeV source, MGRO J2019+37 (see Fig. 2.5). To clarify the relative positioning of these objects, Fig. 2.6 shows the regions containing TeV J2032+4130 and MGRO J2019+37 (which are separated by ∼ 5°), as well as sources from the Third EGRET [159] and GeV [161] catalogs (which all remain unidentified), overlaid on the diffuse emission measured by EGRET [158, 162].

It is of advantage to understand this region in general and MGRO J2019+37 in particular. Towards this goal, in the following, we examine the vicinity of this source to note possible identified counterparts, which, among other things, give us concrete distances. Furthermore, in addition to Milagro, this region has been observed by other experiments in the past. We make use of these observations, particularly those of EGRET and CASA-MIA, to construct a gamma-ray spectrum for MGRO J2019+37. Finally, we discuss the possibility of using IceCube to determine whether the observed gamma-rays are the product of hadronic processes, which would indicate a site of cosmic-ray production.

2.7.1 A TeV Source of Unknown Nature

Most conspicuous in the Milagro view of the Cygnus region is the new TeV source MGRO J2019+37. While this source has been observed with very high significance (∼ 11σ) at very high energies (median energy of 12 TeV), no obvious multi-wavelength counterpart seems to be present. Analysis by Milagro does suggest that the TeV
emission originates from either a single extended source or a combination of several unresolved point sources, fit by a 2D Gaussian of width $\theta = 0.32^\circ$. This implies a radius of

$$r \simeq 5 \left( \frac{\theta}{0.3^\circ} \right) \left( \frac{D}{1\text{kpc}} \right) \text{pc},$$

which depends upon the unknown source distance (scaled to 1 kpc for convenience).

We present in Fig. 2.5 objects in the vicinity of MGRO J2019+37 that merit further
Figure 2.6: Gamma-ray sources and diffuse GeV emission in the Cygnus region. Shown are the sources discovered by Milagro (MGRO J2019+37) and HEGRA (TeV J2032+4130), along with their approximate (1σ) error circles. The fitted extent of the Milagro source is comparable to the circle shown. Also shown are nearby Third EGRET (3EG) (compiled from > 100 MeV gamma rays) and GeV (> 1 GeV gamma rays) catalog sources (all at 95% confidence); as well as gamma-ray source candidates (points), the Cyg OB2 core (dashed circle), and the region of Fig. 2.5 (box). EGRET 4−10 GeV (point-source subtracted) diffuse emission (smoothed and scaled linearly from $\sim 1 - 10 \times 10^{-6}$ cm$^{-2}$ s$^{-1}$ sr$^{-1}$, with white most intense) is also displayed [39].

study. These objects (typically located at distances of $\sim 1 - 4$ kpc) can also be seen in Fig. 2.6 in the context of the greater Cygnus region, which (for the coordinate range displayed) is also nearly completely covered by diffuse TeV emission [33].

Of most immediate interest are two unidentified EGRET sources, 3EG J2016+3657 and 3EG J2021+3716 (constructed from gamma rays with energy > 100 MeV), and the GeV catalog source, GeV J2020+3658 (> 1 GeV gamma rays [161]). These
sources have received a fair amount of attention over the years in searches for potential counterparts. For 3EG J2016+3657, a probable association with a blazar (of unknown redshift) was found. However, pair production on the extragalactic infrared background makes this an unlikely 12 TeV source [33]. 3EG J2021+3716 and GeV J2020+3658 have intriguing possible correlations with several Galactic objects. One of these is the pulsar wind nebula, PWN 75.2+0.1, which, considering the growing number of known TeV PWNe (e.g., Ref. [129]), is also one of the better candidates for MGRO J2019+37.

In light of the recent possible discovery of a TeV source (HESS J1023-575) coincident with the Wolf-Rayet eclipsing binary WR 20a, it is worthwhile to consider similar systems near MGRO J2019+37, some of which have been examined as possible counterparts to the EGRET sources. Shown in Fig. 2.5 are the Wolf-Rayet stars WR 137 [163], WR 138 [163], WR 141 [163], WR 142 [164], and HD 228766 [165]. Also displayed is V382 Cyg, a massive eclipsing binary with a significant rate of mass loss [166]. In fact, many of these systems are known to have large mass loss rates ($\gtrsim 10^{-5} M_\odot \text{ yr}^{-1}$ [167]), which may power shocks that accelerate cosmic rays. The reader is encouraged to see Ref. [168] (and references therein) for details of gamma-ray modeling in such environments. In a similar vein, the open star cluster Berkeley 87 has been proposed as a source of gamma rays.

With the class of TeV supernova remnants already having several prominent members, such an object in Cygnus would be a compelling candidate. However, a search of the area immediately around MGRO J2019+37 yields only SNR 074.8+00.6 (likely just an HII region [169]) and the quite distant (12 kpc [170]) SNR 074.9+01.2 (an unlikely association). The high energy source closest to the quoted Milagro position, IGR J20188+3647, is seen only in the $17-30$ keV band. Also worth noting is Cygnus
Figure 2.7: Data and possible hadronic spectra for MGRO J2019+37. Shown are the Milagro measurement at 12 TeV (diamond), the EGRET spectrum for 3EG J2021+3716 (circles), the upper limit from Whipple (0.3 TeV), and our inferred upper limit from CASA-MIA (115 TeV). Also shown are hadronic fits to the data, assuming $E_p^{-2.35}$ (upper) and $E_p^{-2}$ (lower) source proton spectra. The region above 1 TeV is of greatest interest to neutrino astronomy [39].

Rift, a molecular cloud complex a few kpc in extent which runs east–west through this region [170].

2.7.2 Gamma Rays and Cosmic Rays

Additional insight can be obtained by constructing a gamma-ray spectrum based on the Milagro measurement. Milagro reported a gamma-ray flux from MGRO J2019+37
of \( E^2d\Phi/dE = (3.49 \pm 0.47_{\text{stat}} \pm 1.05_{\text{sys}}) \times 10^{-12} \text{ TeV cm}^{-2} \text{ s}^{-1} \) at a median energy of 12 TeV (see Fig. 2.7), with an undetermined (but subdominant) contribution from the surrounding diffuse emission. This detection, at such a high energy, is quite useful in creating a spectrum, however, we also need additional data at higher and lower energies. The two nearby (unidentified) EGRET sources may provide information at the low end. Both have similar spectra in the MeV-GeV range [162]; however, as 3EG J2021+3716 is more likely to be associated with GeV J2020+3658, we will consider its possible association with the Milagro source. Near 1 TeV, Whipple observations place upper limits on emission from the region near 3EG J2021+3716 (see Fig. 2.7), with a stronger limit from the vicinity of PWN 75.2+0.1 [171] (not shown). Emission from Berkeley 87 is also constrained from HEGRA observations [172] (also not shown).

At energies higher than the Milagro measurement, observational limits are crucial in determining the spectral cutoff, which is of great importance to cosmic-ray and neutrino studies [38]. Fortunately, CASA-MIA observations can provide valuable limits. While a previous all-sky limit is not very constraining [173], later searches had greatly enhanced statistics allowing for an improvement of more than an order of magnitude [174]. Although there is yet no published limit specifically for the MGRO J2019+37 region, we can infer a limit based upon observations of the region containing Cygnus X-3 [174], as there were no significant sources in its general vicinity [174, 175]. The integral flux limit from CASA-MIA at the position of Cygnus X-3 is \( \Phi(E > 115 \text{ TeV}) < 6.3 \times 10^{-15} \text{ photons cm}^{-2} \text{ s}^{-1} \), given above the median detector energy in order to reduce dependence on the assumed spectral index [174]. We infer a similar limit for the region near MGRO J2019+37 by conservatively rounding the Cygnus X-3 result up to \( 10^{-14} \text{ cm}^{-2} \text{ s}^{-1} \) and by treating the integral as \( E d\Phi/dE \), giving an upper limit of \( E^2d\Phi/dE = 10^{-12} \text{ TeV cm}^{-2} \text{ s}^{-1} \) at 115 TeV for the MGRO J2019+37 region (see Fig. 2.7), which (even if somewhat higher) gives meaningful constraints.
As the properties of this source (e.g., magnetic fields) that would determine a leptonic spectrum are quite unknown, we will instead construct a simple, entirely hadronic spectrum (which requires fewer assumptions), not attempting to include contributions from primary or secondary electrons. We use the parametrization for gamma rays resulting from $p$-$p$ scattering of Kelner, Aharonian and Bugayov [138]. The gamma-ray spectrum resulting from a source distribution of protons of the form $d\Phi_p/dE_p = A_p E_p^{-\alpha} \exp \left[-\left(E_p/E_{\text{cut}}^p\right)\right]$, interacting with ambient protons of density $n_H$, is

$$
\frac{d\Phi_\gamma}{dE_\gamma} = c n_H \int_{E_{\gamma}}^{\infty} \sigma_{\text{inel}}(E_p) \frac{d\Phi_p}{dE_p} F_\gamma \left(\frac{E_\gamma}{E_p}, E_p\right) \frac{dE_p}{E_p},
$$

where $\sigma_{\text{inel}}(E_p)$ is the inelastic $p$-$p$ cross section, $c$ is the speed of light, and the function $F_\gamma(E_\gamma/E_p, E_p)$ (given by Eq. (58) of Ref. [138]) determines the number of photons produced per energy interval per scattering. We use Eq. (2.7.2) for $E_\gamma > 1$ TeV and the $\delta$-function approximation of Ref. [138] at lower energies to produce our gamma-ray spectra.

We first assume that there is an association between 3EG J2021+3716 and MGRO J2019+37. Using an $E_p^{-2.35}$ proton spectrum (with $E_{\text{cut}}^p = 1000$ TeV) and normalizing to the Milagro measurement at 12 TeV, we find the upper spectrum of Fig. 2.7. This spectrum gives a reasonable (considering the uncertainties involved) fit to the EGRET data and, importantly, satisfies the CASA-MIA limit. While the distance ($D$) to the source and ambient proton density ($n_H$) are both unknown, normalizing to the Milagro flux allows us to find the total necessary energy injected into cosmic-ray protons as

$$
\mathcal{E}_p \approx 5 \times 10^{50} \left(\frac{1 \text{ cm}^{-3}}{n_H}\right) \left(\frac{D}{1 \text{ kpc}}\right)^2 \text{ erg},
$$

which, for comparison, is similar to the total explosion energy of a typical core-collapse supernova [176] when $D \sim \text{ few kpc}$. 55
It is quite possible that neither of the EGRET sources are at all related to MGRO J2019+37. In any case, EGRET (as well as Whipple and HEGRA) constrains the spectrum from extrapolating to lower energies too steeply, if only for the simple reason that another source would have been seen. We also consider an $E_p^{-2}$ proton spectrum (with $E_p^{\text{cut}} = 500$ TeV), which gives the lower spectrum in Fig. 2.7. While safely below the low-energy observations, the CASA-MIA limit again forces a cutoff at high energies. This scenario requires $\sim 10$ times less input cosmic-ray energy, due to the lower required GeV gamma-ray flux.

The most economical approach towards explaining the Milagro observation would be to have the gamma-ray spectrum peak (in our $E^2$ plot) at $\sim 10$ TeV and then sharply fall off (to satisfy CASA-MIA). Such a spectrum has been observed from the PWN Vela X by HESS [129], which may have a hadronic origin [130]. If PWN 75.2+0.1 has a spectrum of this form, it would easily satisfy the stronger Whipple limit. However, at distance of $\sim 10$ kpc, it would need to be more than an order of magnitude more energetic than Vela X to account for the Milagro measurement.

### 2.7.3 Cygnus and IceCube

While a spectral analysis gives important guidance, the most direct way to discern the nature of a TeV source is through neutrino observations. Our interest here will be limited to the $\nu_\mu + \bar{\nu}_\mu$ flux, which can produce ultrarelativistic muons in charged-current interactions within the Antarctic ice cap.

Knowledge of the location of gamma-ray sources \textit{a priori} allows for neutrino searches without having to pay the trials factor associated with an undirected, all-sky search. Knowledge of source \textit{spectra} then allows for neutrino flux predictions, which can be compared with observations. For MGRO J2019+37, we can use the spectra shown in Fig. 2.7 to derive the associated neutrino flux and the expected event rate.
of TeV muons in IceCube. Following Ref. [138], we parametrize each gamma-ray spectrum with a fit of the form

\[
\frac{d\Phi_\gamma}{dE_\gamma} = A_\gamma E_\gamma^{-\beta} e^{-\left(E_\gamma/E_{\text{cut}}^\gamma\right)^{1/2}},
\]  

(2.7.4)

with \(A_\gamma\) chosen to fit the gamma-ray spectrum at 1 TeV. For the \(E_p^{-2.35}\) proton spectrum, we use \(\beta = 2.2\) and \(E_{\text{cut}}^\gamma = 45\) TeV; while we find \(\beta = 1.9\) and \(E_{\text{cut}}^\gamma = 20\) TeV well fits the gamma rays resulting from the \(E_p^{-2}\) input spectrum (above 1 TeV).

Using the methods detailed in prior sections, we then calculate the expected spectrum of \((\nu_\mu + \bar{\nu}_\mu)\)-induced muons (again, as a function of the observed muon energy entering or originating in the detector) for IceCube, as seen in Fig. 2.8. The uncertainties in these results are at or below that of the Milagro flux measurement itself. This result is compared to the background arising from the atmospheric neutrino spectrum (which falls off more steeply with energy) in a 3 deg\(^2\) bin, which roughly corresponds to the angular resolution of IceCube. Note that we only consider muons with \(E_\mu > 1\) TeV, which are more likely to have an accurately measured energy and direction than less energetic muons. This is important for effectively discriminating between the expected signal and background.

With such a similar gamma-ray spectrum above 1 TeV, the \(E_p^{-2}\) result is practically indistinguishable from the \(E_p^{-2.35}\) case displayed. In fact, since the Milagro measurement is at such a high energy, the expected muon rate is only weakly dependent on the spectral index, with the location of the cutoff being somewhat more important. This rate satisfies the constraints previously derived in Ref. [38] for the Cygnus region. We find that IceCube can detect neutrinos from the MGRO J2019+37 source, if indeed it is a cosmic-ray accelerator. It is worth emphasizing that not a single high-energy neutrino source has yet been detected. While the number of events will not be large, a signal could be separated from background within several years,
especially taking advantage of the rapid increase in signal to background probability with increasing measured muon energy.

2.8 A TeV Neutrino Beacon

It would be beneficial to have a bright source of TeV neutrinos that can be seen by both IceCube and a Mediterranean detector. The ideal source for this purpose would
have a declination in the range $+10^\circ \lesssim \delta \lesssim +30^\circ$, which also has the benefit of reduced neutrino attenuation in Earth at high energies, resulting in increased muon and shower event rates. The Crab nebula is already an accepted standard candle in TeV gamma-ray astronomy. From HESS measurements [177], if the entire TeV spectrum could be attributed to hadronic processes, our methods yield an event rate of $N(E_\mu > 1 \text{ TeV}) \sim 2.7 \text{ yr}^{-1}$ in IceCube. The Crab spectrum is relatively soft and well described by leptonic processes [178]. A pionic component may be uncovered by a significant observation of neutrino events.

The Vela complex that we have already discussed is located in the southern sky. Only after extensive observations was the existence of Vela Jr. confirmed [179]. It would not be surprising, then, for a similarly bright TeV source to be discovered in the northern sky (perhaps directly through neutrino events). The spectrum for such an object was calculated near the end of Section 2.3 for a variety of spectral indices and cutoffs. Future gamma-ray observations of this declination range in the TeV regime may reveal another important source for neutrino astronomy.

Multi-wavelength observations of such a source would add a rung to a TeV neutrino “distance ladder”. Cosmic-ray interactions in the Earth’s atmosphere generate the nearest guaranteed neutrino source [180, 181]. Similar processes in the atmosphere of the Sun, which is at a well known distance, are quite likely to produce TeV neutrinos as well [182, 183, 184]. Of the prospective sources, the Vela SNR is estimated to be only $\sim 250 – 300 \text{ pc}$ from Earth, while Vela Jr. may be even closer [185]. RX J1713–3946 has a distance estimate of $\sim 1 \text{ kpc}$ [186]. At $\sim 8 \text{ kpc}$, the GCD is the most distant guaranteed source with a confirmed distance. AGN and GRBs, if they produce measurable fluxes of neutrinos, would provide a range of very remote point sources. Measurements of neutrinos from a variety of distances (ideally with
flavor ratios) would provide important information for testing neutrino properties at increasing $L/E$ (distance/energy).

### 2.9 Discussion and Conclusions

Upcoming neutrino telescopes will deliver the first direct evidence concerning the production mechanism of TeV gamma-ray sources. Following the pioneering AMANDA [187] and Baikal [188] efforts, the next generation IceCube [36] and Mediterranean [37] km$^3$ detectors will reach the scale necessary to examine the Galactic sources discussed here. Hadronic mechanisms would be confirmed through neutrino detection, while a significant absence of neutrino events would imply leptonic processes. They will be able to probe the expected high energy spectral cutoffs that would otherwise be unobservable. The cutoff in a SNR may well evolve with age, so finding these cutoffs with neutrinos may also yield valuable new information in this regard.

Combined, IceCube and a km$^3$ Mediterranean detector will provide continuous, all-sky coverage and can be used together to detect neutrino-induced muons and showers from TeV sources. The good angular resolution for muon events can be used to precisely locate a source in neutrinos. Showers, and contained muon events, provide accurate reconstruction of the source spectrum at energies beyond the reach of gamma-ray telescopes. These measurements will provide important information concerning the origin of high energy Galactic cosmic rays, potentially directly observing the source population responsible for production up to the knee at $\sim 3 \times 10^{15}$ eV.

The shell-type SNR Vela Jr. (RX J0852.0–4622) is an intriguing prospective source of TeV neutrinos. The extent and intensity of the TeV emission make this, as well as the Vela complex as a whole, a unique target for neutrino telescopes. The confirmation of hadronic/leptonic processes through detection/non-detection of neutrino events
from Vela X and the possible detection of neutrinos through shower events from the Vela SNR shell are exciting possibilities.

While it is expected that other sources may eventually be more compelling, only the Galactic Center diffuse emission can presently be claimed to possess a guaranteed neutrino flux. The detection of this neutrino flux will give further insight into the complex processes occurring in the Galactic Center region. When we consider that, to date, no high energy astrophysical neutrinos have ever been positively detected, the importance of such a guaranteed flux of TeV neutrinos is difficult to overstate.

Included in the catalog of TeV sources are several that remain unidentified. As more observations are undertaken by the HESS, MAGIC, and VERITAS telescopes, this number is expected to increase. Discovering neutrino fluxes from these sources would provide invaluable information concerning their nature. As new TeV gamma-ray sources are discovered, the pool of potential TeV neutrino sources increases. While some of these may not be significant alone, when grouped into classes, stacking may increase the potential for neutrino studies. An even larger number of unidentified sources remains in the GeV regime from EGRET observations. Many of these sources have intense, hard spectra. If any of these sources, some of which are visible to IceCube, have spectra that extrapolate into the TeV regime and are pionic, they would have abundant fluxes of neutrinos.

In summary, the prospects for the near-term first discoveries of Galactic TeV neutrino sources are very good. Importantly, this conclusion is empirically based on the measured spectra of bright Galactic TeV gamma-ray sources. For some of these, there are very compelling independent arguments that the observed gamma rays arise from neutral pion decays, meaning that they must be accompanied by neutrinos. It is essential to test this directly for these sources, as well as for others where the possibility of neutrino emission is uncertain. We emphasize the importance
of using the measured muon energy spectra to discriminate against the quickly falling atmospheric neutrino backgrounds. For example, a single event near 10 TeV from a source direction is almost certainly signal, while an event near 1 TeV has a much higher probability of being background. This fact alone could be enough to help establish discovery. Due to the amplifying factors of neutrino cross section and muon range, neutrino detectors have better reach to the highest energies in the source spectra, as compared to gamma-ray telescopes, which can make precise measurements at lower energies. This complementarity can be exploited to help solve the long-standing puzzle of the origin of the Galactic cosmic rays.

For all of the sources discussed here, the rates in km$^3$ neutrino telescopes are relatively small, at most $\sim 1 - 10$ events/year, though we have argued that even these small rates could have a powerful impact. While it remains possible that various uncertainties and detector limitations may make these observations even more challenging, they might also be better than described here. If the source spectra extrapolate to higher energies than we have assumed, or especially if the emitted TeV gamma-ray spectra have been diminished by absorption in the sources, which is quite possible, then the neutrino detection event rates could be significantly larger than shown here. Given the potentially unique power of neutrino astronomy, we can only hope that Nature has been so kind.
CHAPTER 3

CORE-COLLAPSE SUPERNOVAE IN THE NEARBY UNIVERSE

That which has died falls not out of the universe. If it stays here, it also changes here, and is dissolved into its proper parts, which are elements of the universe and of thyself.

Marcus Aurelius, Book VIII of Meditations

3.1 Introduction

Core-collapse supernovae have long been suspected to be the solution of many long-standing puzzles, including the production of neutron stars and black holes, radioactive isotopes and heavy elements, and cosmic rays [1, 2]. Understanding these issues, and the properties of neutrinos and hypothesized new particles, requires improving our knowledge of supernovae. It is not enough to record their spectacular visual displays, as these do not reveal the dynamics of the innermost regions of the exploding stars, with their extremes of mass and energy density. Moreover, sophisticated simulations of the core collapse of massive stars do not robustly lead to supernova explosions [16, 17, 18], raising the suspicion that crucial physics is missing.

1The content of this Chapter is based in large part on our work in Ref. [12], and related to Refs. [23, 26, 27].
Neutrinos are the essential probe of these dynamics, as they are the only particle that escapes from the core to the observer (gravitational waves may be emitted, but they are energetically subdominant). There is an important corollary to this, namely

until supernovae besides SN 1987A are detected by neutrinos, our fundamental questions about supernovae will never be decisively answered. In fact, the most interesting problems—associated with the presence, nature, variety, and frequency of core collapse in massive stars—can only be solved by detecting many supernova neutrino bursts.

The challenges of supernova neutrino burst detection are that Milky Way sources are rare and that more common distant sources have little flux. The 32 kton Super-Kamiokande (SK) detector is large enough to detect with high statistics a burst from anywhere in the Milky Way or its dwarf companions, but the expected supernova rate is only 1–3 per century, and there is no remedy but patience. Proposed underground detectors [189, 190, 191, 192], like the ∼0.5-Mton Hyper-Kamiokande (HK), could detect one or two neutrinos from supernovae in some nearby galaxies [11]. As shown in Fig. 3.1, to robustly detect all neutrino bursts within several Mpc, where recent observations show the supernova rate to be at least 1 (2) per year within ∼6 (10) Mpc, requires scaling up the detector mass of SK by about two orders of magnitude, to at least ∼5 Mton [12].

A recent proposal for the Deep-TITAND detector shows in detail how it might be feasible to build such a large detector in a cost-effective way [193, 194]. To avoid the high costs and slow pace of excavating caverns underground, this proposal conceives of a modular 5 Mton undersea detector that could be constructed quickly. Key motivations for such a detector are superior exposure for studies of proton decay, long-baseline neutrinos, and atmospheric neutrinos. To reduce costs, the detector would be built with a shallower depth and lower photomultiplier coverage than SK; these decisions would sacrifice the low-energy capabilities for all but burst detection.
Figure 3.1: Probabilities to obtain the indicated numbers of $\bar{\nu}_e$ neutrino events (with $E_{e^+} > 18$ MeV) in a 5 Mton detector as a function of the supernova distance. We assume a Fermi-Dirac $\bar{\nu}_e$ spectrum with an average energy of 15 MeV and a total energy of $5 \times 10^{52}$ erg. Optical supernovae observed in the last 10 years are noted at their distances; those in red indicate multiple supernovae in the same galaxy [12].

There is a compelling case for a 5 Mton detector based on supernova neutrino detection alone, and the science benefits that we discuss below will hold even if a Milky Way supernova is detected first. On an annual basis, one would expect a burst of $\gtrsim 3$ events, and every several years, a burst comparable to the $\sim 10$ events from SN 1987A detected by each of Kamiokande-II [195, 196] or IMB [197, 198].
Indeed, a 5 Mton supernova neutrino detector is one of the most promising prospects for developing an observatory for non-photon time-domain astrophysics. There are no serious uncertainties in the number of sources or the strengths of their signals. The minimal size of the required detector is known now, and it is not out of reach, with costs comparable to those of existing or near-term high-energy neutrino and gravitational-wave observatories.

Before elaborating on details concerning detection rates, we will begin by exploring how the data obtained from multiple neutrino bursts would transform the way that we consider questions about supernovae; these considerations are a major part of our new results. We will then examine recent developments concerning the rate and properties of supernovae observed in the nearby universe. This will lead into our discussion of the detector properties required to measure neutrino bursts from these supernovae, focusing on the Deep-TITAND proposal [193, 194], and the quantitative neutrino yields expected.

### 3.2 Discovery Prospects

Our primary interest is on the scientific impact of measuring neutrino “mini-bursts,” detectable signals of 3 or more events within 10 seconds (the observed duration of the SN 1987A neutrino burst), from many supernovae in the nearby universe. As we will show in Sections 3.7 and 3.8, the minimum detector size for achieving this purpose is about 5 Mton. We emphasize in advance that such signals can be separated from backgrounds even at shallow depth, so that the presence of a core collapse can be deduced independently of photon-based observations. Additionally, for nearby transients identified through photons, a non-detection in neutrinos means that a conventional supernova neutrino flux was not present. These facts have new and profound implications.
While our principal focus is thus on individual objects, the aggregate data would, of course, also be useful. For science goals that require a large number of accumulated events, the most certain signal is the Diffuse Supernova Neutrino Background (DSNB), which is a steady flux arising from all core-collapse supernovae in the universe (e.g., Refs. [199, 200, 201, 202] and references therein). In the proposed $\sim 0.5$ Mton HK detector, with added gadolinium to reduce backgrounds by neutron tagging [202], $\sim 50–100$ DSNB signal events with little background could be collected per year. The ratio of DSNB signal to detector background in Deep-TITAND would be the same as in the background-dominated SK search of Ref. [203], which set an upper limit. To reach the smallest plausible DSNB signals, one needs an improvement of about a factor 3 in signal sensitivity and thus a factor of about 10 in exposure. After four years, as in the SK search, the Deep-TITAND exposure would be about 100 times larger than that of Ref. [203], thus allowing a robust detection of the DSNB flux. (To measure the spectrum well, HK with gadolinium would be needed.) The fortuitous occurrence of a supernova in the Milky Way, or even Andromeda (M31) or Triangulum (M33), would also give a very large number of neutrino events (see Table 3.2). The physics prospects associated with such yields from a single supernova have been discussed for underground detectors at the 0.5 Mton scale.

3.3 Probing the Core Collapse Mechanism

The optical signals of supposed core-collapse supernovae show great diversity [13, 204], presumably reflecting the wide range of masses and other properties of the massive progenitor stars. In contrast, the neutrino signals, which depend on the formation of a $\sim 1.4M_\odot$ neutron star, are presumed to be much more uniform. However, since we have only observed neutrinos from SN 1987A, it remains to be tested whether all core-collapse supernovae do indeed have comparable neutrino emission. The total
energy emitted in neutrinos is $\simeq 3G M^2/5R$, and some variation is expected in the mass $M$ and radius $R$ of the neutron star that is formed, though proportionally much less than in the progenitor stars.

With at least $\sim 1$ nearby supernova per year, a wide variety of supernovae can be probed, including less common types. For example, the observational Types Ib and Ic are now believed to be powered by core collapse, despite their original spectroscopic classification that defined them as related to Type Ia supernovae, which are thought to be powered by a thermonuclear runaway without significant neutrino emission. While each of the Types Ib/Ic and Ia are only several times less frequent than Type II, some of each should occur nearby within a reasonable time, so that the commonality of the Type II/Ib/Ic explosion mechanism can be tested.

### 3.3.1 SN 2008S

While the nature of the explosion in the above supernova types is very likely as expected, there are other bright transients observed for which the basic mechanism is much more controversial. For these events, the detection or non-detection of neutrinos
could decisively settle debates that are hard to resolve with only optical data. One type of so-called “supernova impostor” is thought to be the outburst of a Luminous Blue Variable (LBV) [205], which seem to require a stellar mass of $M_* \gtrsim 20 M\odot$. Since this type of outburst affects only the outer layers, with the star remaining afterward, there should be no detectable neutrino emission.

There are several recent examples in nearby galaxies where neutrino observations could have been conclusive, including the likely LBV outburst SN 2002kg in NGC 2403 [206]. SN 2008S in NGC 6946 [23] and a mysterious optical transient in
NGC 300 [27] warrant further discussion for another reason. In neither case was a progenitor seen in deep, pre-explosion optical images (Fig. 3.2 [23]); however, both were revealed as relatively low-mass stars ($M_* \sim 10 M_\odot$) by mid-infrared observations made years before the explosions. This suggests that they were obscured by dust expelled from their envelopes, a possible signature of stars dying with cores composed of O-Ne-Mg instead of iron [23, 27]. As we will address in detail later, these events were sufficiently near for a 5 Mton detector to have identified them as authentic supernovae or impostors.

### 3.4 Measuring the Total Core Collapse Rate

In the previous subsection, we implicitly considered supernovae for which the optical display was seen. However, as we will calculate, the detection of $\geq 3$ neutrinos is sufficient to establish that a core collapse occurred, including those events not later visible to telescopes. This provides a means of measuring the total rate of true core collapses in the nearby universe. A successful supernova may be invisible simply if it is in a very dusty galaxy, of which there are examples quite nearby, such as NGC 253 and M82. These are supposed to have very high supernova rates, perhaps as frequent as one per decade each, as deduced from radio observations of the number of young supernova remnants [207]. However, only a very few supernovae have been seen [9].

More interestingly, it remains unknown if, as in numerical models of supernova explosions, some core collapses are simply not successful at producing optical supernovae. As illustrated in Fig. 3.3, this can occur if the outgoing shock is not sufficiently energetic to eject the envelope of the progenitor star, in which case one expects the prompt formation of a black hole with very little optical emission [15]. Indirect evidence for such events follows from a deficit of high-mass supernova progenitors
Figure 3.3: Possible outcomes of stellar collapse resulting in the formation of a black hole [26]. Scenarios that do not result in a visible explosion have been largely unexplored.

compared to expectations from theory [26, 21], as well as from the existence of black holes recently discovered to have $M_{\text{BH}} \gtrsim 15 M_\odot$ [208].

One way to probe this exotic outcome would be to simply watch the star disappear [26]. However, a detectable burst of neutrinos should be emitted before the black hole forms (and typically, if the duration of the emission is shorter, the luminosity is higher) [209, 210, 211]. Taken together, these would be a dramatic and irrefutable signal of an otherwise invisible event. While the rate of prompt black hole formation probably cannot exceed the visible supernova rate without violating constraints on
the DSNB, reasonable estimates indicate that up to $\gtrsim 20\%$ of core collapses may have this fate [26].

### 3.5 Testing the Neutrino Signal

By measuring neutrinos from many supernovae, the deduced energy spectra and time profiles could be compared to each other and to theory. In most cases, only several events would be detected, but this is enough to be useful. The highest neutrino energies range up to $\simeq 50$ MeV. The thermal nature of the neutrino spectrum makes it relatively narrow, and since it is falling exponentially at high energies, even a small number of events can help determine the temperature. Recall that for SN 1987A, the Kamiokande-II and IMB detectors collected only $\sim 10$ events each [195, 197], but that this data strongly restricts the details of the collapsed core.

The time profile is thought to rise quickly, over perhaps at most 0.1 s, and then decline over several seconds, as seen for SN 1987A. The neutrino events collected would most likely be at the early peak of the emission, and hence the most relevant for the question of whether heating by the emergent neutrino flux is adequate for shock revival [212, 213, 214] or whether $\nu$-$\nu$ many-body effects are important.

Over time, as many supernovae are detected, the average energy spectrum and time profile will be built up. (For the time profile, there will be some uncertainties in the start times.) If there are large variations from one supernova to the next, then these average quantities will ultimately provide a more useful template for comparison than the theoretical results that must be used at present. If there is no evidence for significant variations between supernovae, then the accumulated data will be equivalent to having detected one supernova with many events. It is quite likely that such a detector would observe a supernova in one of the Milky Way, M31, or M33; the high-statistics yield from these would also provide a point of comparison. Taken
together, all of these data will provide new and exacting tests of how supernovae work.

With enough accumulated events, it is expected that neutrino reactions besides the dominant inverse beta decay process will be present in the data. One oddity still remaining from SN 1987A is that the first event in Kamiokande-II seems to be due to $\nu_e + e^- \rightarrow \nu_e + e^-$ scattering and points back to the supernova [196], which is improbable based upon standard expectations [215]. This can be tested, however, and if it turns about to be ubiquitous, could be exploited in determining the directionality of the larger future bursts without optical signals, as the inverse beta decay signal is not directional [216].

Since Earth is transparent to supernova neutrinos, the whole sky can be monitored at once. For neutrinos that pass through Earth, particularly those which cross the core, matter-enhanced neutrino mixing can significantly affect the spectrum relative to those which do not. Dividing the accumulated spectra appropriately based on optical detections, this would allow a new test of neutrino mixing, sensitive to the sense of the neutrino mass hierarchy. Detecting neutrinos from distant sources would also allow tests of neutrino decay, the equivalence principle, and other exotic possibilities.

### 3.6 Revealing Other Transient Signals

Detection of a neutrino burst means detection of the instant of core collapse, with a precision of $\sim 1$ second determined by the sampling of the peak of the $\simeq 10$ second time profile. This would provide a much smaller time window in which to search for gravitational wave signals [217, 218, 219, 220, 221] from core-collapse supernovae;
otherwise, one must rely on the optical signal of the supernova, which might optimistically be determined to a day ($\sim 10^5$ seconds). This is important, since the gravitational-wave signal remains quite uncertain, making searches more difficult.

Once core collapse occurs, the outward appearance of the star initially remains unchanged. Knowing that a signal was imminent would give unprecedented advanced warning that photons should soon be on the way, allowing searches to commence for the elusive UV/X-ray signal of supernova shock breakout [222, 223] and also the early supernova light curve. Those signals are expected to emerge within hours and days, respectively. While the neutrino signal is likely not directional, the number of events detected will provide constraints for triggered searches.

Finally, it is possible that such large detector would find not only core-collapse supernovae in nearby galaxies, but also other types of transients that are presently unknown. One such possibility is the merger of two compact objects, such as two neutron stars of a neutron star and white dwarf [224]. In the Milky Way, there would be sensitivity to any transient with a supernova-like neutrino signal, as long as its overall strength is at least $\sim 10^{-6}$ as large as that for a supernova. To be detectable, the key requirement is a $\gtrsim 15$ MeV $\bar{\nu}_e$ component.

### 3.7 Nearby Supernova Rate

Over the past decade, there has been rapid growth in the level of interest among astronomers in measuring the properties of core-collapse supernovae. There is also a renewed interest in completely characterizing the galaxies in the nearby universe, within 10 Mpc. In nearby galaxies, both amateurs and automated surveys (e.g., KAIT [14]) are finding many supernovae. For these, archival searches have revealed pre-explosion images of about a dozen supernova progenitor stars, allowing a better
Figure 3.4: Estimates of the core-collapse supernova rate in the nearby universe, based on that expected from the optical luminosities of known galaxies (line) and supernovae observed within the last decade (bins). Note that SN 2002kg is a likely LBV outburst, while SN 2008S and the NGC 300 transient are of unusual origin. These estimates are all likely to be incomplete [12].

understanding of which types of massive stars lead to which kinds of core-collapse supernovae (e.g., [23, 21, 22, 225]).

Figure 3.4 shows the expected rate of core-collapse supernovae in the nearby universe (dashed line) calculated using the galaxy catalog of Ref. [226] (designed to be ∼70–80% complete up to 8 Mpc), with a conversion from $B$-band optical luminosity to supernova rate from Ref. [227]. The effects of clustering and of incompleteness at large
distances can clearly be seen, since the histogram would rise as the distance squared for a smooth universe of identical galaxies. Ultimately, a more accurate result could be obtained by combining the information from star-formation rate measurements in the ultraviolet [228], Hα [229, 230], and infrared [231], likely leading to a larger prediction for the supernova rates.

Also displayed in Fig. 3.4 is the rate deduced from supernovae discovered in this volume in the last 10 years [9], with distances primarily from Ref. [226] (when available; otherwise from [232, 233]). While the observed rate is already \( \sim 2 \) times larger than the above calculation, even this estimate is likely incomplete, as supernova surveys under-sample small galaxies and the Southern hemisphere. As previously mentioned, supernovae with little or no optical signal, e.g., due to direct black hole formation or dust obscuration, would also have been missed. This is particularly important for nearby dusty starburst galaxies with large expected, but low observed, supernova rates, like NGC 253 and M82.

Distance measurements of nearby galaxies also stand to be improved. For example, at the largest distances, SN 1999em, SN 1999ev, SN 2002bu, and SN 2007gr are probably not all truly within 10 Mpc, as some distance estimates put them outside. We emphasize that their inclusion or not does not affect our approximate supernova rates, and barely matters for the neutrino bursts of sufficient multiplicity, which are dominantly from closer supernovae. It would be very helpful to refine distance measurements, not just for star formation/supernova rate estimates, but also to determine the absolute neutrino luminosities once a supernova has been detected.

Overall, there is a strong case that the core-collapse supernova rate within \( \sim 6 \) (10) Mpc is at least 1 (2) per year. This can be compared to the estimated Milky Way rate of \( 2 \pm 1 \) per century (see Ref. [44] and references therein), with Poisson probabilities ultimately determining the odds of occurrence, as shown in Fig. 3.5.
Figure 3.5: Probabilities for one or more supernovae in the Milky Way over time spans relevant for the lifetimes of large neutrino detectors, depending on the assumed supernova rate [12].

3.8 Neutrino Burst Detection

A goal of measuring supernova neutrino “mini-bursts” from galaxies at a few Mpc necessitates a large detector, roughly $\sim 100$ times the size of SK. We focus on the Deep-TITAND proposal for a 5 Mton (fiducial volume) enclosed water-Čerenkov detector [193, 194]. The detector would be constructed in modules sized by Čerenkov light transparency and engineering requirements. We assume a photomultiplier coverage of 20%, similar to that of SK-II (half that of the original SK-I and the rebuilt SK-III). As in SK, the detection efficiency at the energies considered here would be nearly unity.

The backgrounds present in deep detectors have been well-characterized by SK and other experiments. Deep-TITAND is proposed to be at a relatively shallow depth of 1000 meters of seawater, which would increase the downgoing cosmic ray muon rate per unit area by a factor $\simeq 100$ compared to SK, which is at a depth of 2700 meters water equivalent. A nearly perfect efficiency for identifying cosmic ray
muons in the outer veto or the detector itself is required. This was achieved in SK, where the only untagged muons decaying in the detector were those produced inside by atmospheric neutrinos [203]. Simple cylinder cuts around cosmic ray muon tracks would veto all subsequent muon decays while introducing only a negligible detector deadtime fraction.

Low-energy backgrounds include natural radioactivities, solar neutrinos, photomultiplier noise, and beta decays from nuclei produced following spallation by cosmic ray muons. Of these, only the last is depth-dependent, and this would be much larger than in SK (a factor $\simeq 30$ for the higher muon rate per area but lower muon average energy, and a factor $\simeq 30$ for the larger detector area). The high muon rate means that it would not be possible to use the cylinder cuts employed in SK to reduce spallation beta decays without saturating the deadtime fraction (note that these beta decays have lifetimes more than $10^6$ times longer than the muon lifetime). At low energies, the above background rates are large, but the spectrum falls steeply with increasing energy, essentially truncating near 18 MeV [203, 10].

This allows for a significant simplification and reduction in the background rate by considering only events with a reconstructed energy greater than 18 MeV (a neutrino energy of 19.3 MeV). Which events to reconstruct would be determined by a simple cut on the number of hit photomultipliers, just as in SK, but with a higher threshold. The backgrounds above this cut are due to atmospheric neutrinos, and thus the rates scale with the detector volume but are independent of depth. The dominant background contribution is from the decays of non-relativistic muons produced by atmospheric neutrinos in the detector, i.e., the so-called invisible muons. The background rate in 18–60 MeV in SK is about 0.2 events/day, of which the energy-resolution smeared tail of the low-energy background is only a minor component [203, 10].
Scaling this rate to a 5 Mton detector mass ($\sim 5 \times 10^{-4} \text{ s}^{-1}$) and considering an analysis window of 10 sec duration (comparable to the SN 1987A neutrino signal) allows calculation of the rate of accidental coincidences [10]. For $N = 3$ events, this corresponds to about only once every five years, and when it does, examination of the energy and timing of the events will allow further discrimination between signal and background (a subsequent optical supernova would confirm a signal, of course). For $N \geq 4$, accidental coincidences are exceedingly rare ($\sim 1$ per 3000 years), therefore we require at least $N = 3$ signal events to claim detection of a supernova (a somewhat greater requirement than in Ref. [11], where a smaller detector was assumed). Since the backgrounds observed by SK in this energy range are from atmospheric neutrinos, we expect no correlated clusters of background events.

To estimate detection prospects, for the $\bar{\nu}_e$ flavor we assume a Fermi-Dirac spectrum with an average energy of 15 MeV and a total energy of $5 \times 10^{52}$ erg. The dominant interaction for the neutrino signal is inverse-beta decay, $\bar{\nu}_e + p \rightarrow n + e^+$, where $E_{e^+} \simeq E_{\bar{\nu}_e} - 1.3 \text{ MeV}$ and the positron direction is nearly isotropic [216]. Combining the emission spectrum, cross section, and number of free target protons in 5 Mton of water, we find that the average number of neutrino events (for $E_{e^+} > 18 \text{ MeV}$) from a burst at distance $D$ is

$$\mu(D; E_{e^+} > 18 \text{ MeV}) \simeq 5 \left( \frac{D}{3.9 \text{ Mpc}} \right)^{-2}. \quad (3.8.1)$$

This is the key normalization for the supernova signal. In Table 3.8, we list recent nearby supernovae within 6 Mpc, with type, host galaxy name, distance, and the expected neutrino yields $\mu$ in a 5 Mton detector. As can be seen in Fig. 3 of Ref. [11], our $E_{e^+} > 18 \text{ MeV}$ threshold still allows us to detect $\sim 70\%$ of the total supernova signal.
<table>
<thead>
<tr>
<th>SN</th>
<th>Type</th>
<th>Host</th>
<th>D [Mpc]</th>
<th>ν events</th>
</tr>
</thead>
<tbody>
<tr>
<td>2002hh</td>
<td>II-P</td>
<td>NGC 6946</td>
<td>5.6</td>
<td>2.4</td>
</tr>
<tr>
<td>2002kg</td>
<td>IIb/LBV</td>
<td>NGC 2403</td>
<td>3.3</td>
<td>6.8</td>
</tr>
<tr>
<td>2004am</td>
<td>II-P</td>
<td>NGC 3034 (M82)</td>
<td>3.53</td>
<td>5.9</td>
</tr>
<tr>
<td>2004dj</td>
<td>II-P</td>
<td>NGC 2403</td>
<td>3.3</td>
<td>6.8</td>
</tr>
<tr>
<td>2004et</td>
<td>II-P</td>
<td>NGC 6946</td>
<td>5.6</td>
<td>2.4</td>
</tr>
<tr>
<td>2005af</td>
<td>II-P</td>
<td>NGC 4945</td>
<td>3.6</td>
<td>5.7</td>
</tr>
<tr>
<td>2008S</td>
<td>II</td>
<td>NGC 6946</td>
<td>5.6</td>
<td>2.4</td>
</tr>
<tr>
<td>2008bk</td>
<td>II-P</td>
<td>NGC 7793</td>
<td>3.91</td>
<td>4.8</td>
</tr>
<tr>
<td>2008?</td>
<td>II?</td>
<td>NGC 300</td>
<td>2.15</td>
<td>16.0</td>
</tr>
</tbody>
</table>

Table 3.2: Recent core-collapse supernova candidates within 6 Mpc, with their expected neutrino event yields ($E_{e+} > 18$ MeV) in a 5 Mton detector.

The probability to detect $\geq N$ neutrino events from a given core collapse is then

$$P(\geq N; D) = \sum_{n=N}^{\infty} P_n[\mu(D)] = \sum_{n=N}^{\infty} \frac{\mu^n(D)}{n!} e^{-\mu(D)},$$

(3.8.2)

where $P_n(\mu)$ represents the Poisson probability. $P(\geq N; D)$ is shown in Fig. 3.1 as a function of $D$ for several values of $N$. From this figure, we see, for example, that from a 4 Mpc supernova, we have an excellent chance ($\gtrsim 90\%$) to get more than 3 neutrino events. For 8 Mpc, like those shown in Fig. 3.4, there is still a $\lesssim 10\%$ chance to get $\geq 3$ events.

For a particular supernova rate, $R_{SN,i}$, we can get the expected total rate of $N$-tuplet detections from distances $D_i$ as

$$R_{N,\text{burst}} = \sum_i R_{SN,i} P_N[\mu(D_i)].$$

(3.8.3)
In Fig. 3.6, we show this as an annual rate, $R_{N,\text{burst}}$, plotted versus $N$. For the supernova rate $R_{SN,i}$, we have adopted three different models: (i) all supernova candidates shown in Fig. 3.4 (20 in total); (ii) same as (i), except excluding SN 2002kg, SN 2008S, and the NGC 300 transient as exceptional events (17 in total); the catalog-based rate estimate (line in Fig. 3.4). As the detection criterion is $N \geq 3$, the annual rate of detectable mini-bursts is obtained by summing $R_{N,\text{burst}}$ for $N \geq 3$, which yields 0.8, 0.6, and 0.4 supernovae per year, for supernova rate models (i), (ii), and (iii), respectively. Adding supernovae from beyond 10 Mpc would not change the rate of $N \geq 3$ multiplets, only increasing the number of unremarkable lower-$N$ multiplets (which, as shown, are already dominated by supernovae in the 8–10 Mpc range).

The total neutrino event counts, $N_{\text{total}}$, can be obtained from $R_{N,\text{burst}}$ by

$$N_{\text{total}} = \sum_{N=3}^{\infty} N R_{N,\text{burst}}, \quad (3.8.4)$$

which are 48, 31, and 22 per decade, for rate estimates (i), (ii), and (iii), respectively.

Since each burst is triggered with $E_{e^+} > 18$ MeV events, one would also look for somewhat lower-energy events in the same time window, potentially raising the total yield by $\simeq 20\%$.

### 3.9 Discussion and Conclusions

The $\sim 10$ neutrino events associated with SN 1987A in each of the Kamiokande-II and IMB detectors [195, 197] were the first and, thus far, only detection of neutrinos from a supernova. This detection showed that we can learn a great deal even from a small number of events, and revealed that an immense amount of energy is released in the form of neutrinos ($> 10^{53}$ erg) during a core collapse. Measuring “mini-bursts”
of neutrino events from multiple supernovae would allow for the study of the core-collapse mechanism of a diverse range of stellar deaths, including optically-dark bursts that appear to be relatively common [26, 21].

This would be made possible by a $\sim 5$ Mton scale water Čerenkov detector [193, 194], which has the special advantages of being able to trigger on supernovae using neutrinos alone, and to guarantee detection if neutrinos are produced with the expected flux. Moreover, for burst detection, a relatively-high low-energy background rate can be tolerated, significantly decreasing the required detector depth, so that construction could be relatively quick and inexpensive. Our conservative estimates
shows that the occurrence rate of mini-bursts that give $\geq 3$ neutrino events is likely $\sim 1 \text{ yr}^{-1}$ or higher.

In conclusion, we wish to reiterate that, even if a supernova occurs in the Milky Way tomorrow, the important problems discussed in Section 3.2 will remain unresolved, and can only be addressed by a suitable “census” of core collapses in the nearby universe. The possibilities mentioned here almost certainly do not exhaust the scientific potential of such an instrument. As is now almost commonplace in the business of observing supernovae with photons, it would be surprising not to find new and unexpected phenomena.
CHAPTER 4
COSMIC GAMMA-RAY BURSTS

What we anticipate seldom occurs: but what we least expect generally happens.

Benjamin Disraeli

4.1 Introduction

The connection between gamma-ray bursts and core-collapse supernovae [58] tells us that, in observing a GRB, we are witnessing the death of a massive, short-lived star.\(^1\) This alone would lead one to expect the cosmic GRB history to follow the star formation history (SFH) [71]. Any deviation from this expectation would provide new information about why a star should die as a GRB, complementing microphysical investigations.

In just the past several years, Swift [76] and a worldwide network of observers have detected gamma-ray bursts from higher redshifts than was previously possible, sparking renewed interest in the GRB redshift distribution. The Swift bursts with known redshifts now provide a sufficiently large sample upon which reasonable cuts can be made, allowing simple, model-independent tests. Additionally, improved star formation measurements [70] now provide a well-defined baseline for comparison.

\(^1\)The content of this Chapter is based in large part on our work in Refs. [77, 79, 86, 88].
Figure 4.1: The luminosity-redshift distribution of 119 Swift GRBs, as we determine from the (updated) catalog of Ref. [235]. Squares represent the 63 GRBs used in Ref. [77], with 56 found subsequently: before (grey circles) and after (red circles) the start of Fermi. Three Fermi-LAT GeV bursts (triangles) are shown (but not used in our analysis). The shaded region approximates an effective threshold for detection. Demarcated are the GRB subsamples used to estimate the SFR [88].

The intense brightness of gamma-ray bursts gives hope that we can probe the history of star formation to very early times, potentially to higher redshifts than with galaxies alone. First, we must be able to observe the GRBs and obtain redshifts for a sufficient number of events. Second, we need to understand how to calibrate the GRB rate to the star formation rate (SFR). Swift has indeed pushed the former greatly ahead, and allowed studies of the latter.

Our goals are to use the large set of Swift gamma-ray bursts with known redshifts (see Fig. 4.1) to examine the above two points in greater detail. Surprisingly, we
find that gamma-ray bursts are not unbiased tracers of the SFR and comment on the suspected origins of this revelation [77, 88]. This does not, however, prevent a study of the amount of high-$z$ star formation; it in fact allows for a more proper estimation.

Several recent high-$z$ bursts, most notably GRB 080913 at $z \simeq 6.7$ [237] and GRB 090423 at $z \simeq 8.1$ [83, 84], also allow us to extend GRB-based SFR determinations to very-high redshifts. Here, direct SFR measurements are quite challenging, particularly at the faint end of the galaxy luminosity function, where GRBs may be ideal tracers. Even with only several events, we determine [79, 88] that the SFR declines only slowly from $z \sim 4$ to $z \gtrsim 8$. This may confirm that a substantial amount of star formation occurs within faint galaxies, in agreement with extrapolations of Lyman Break Galaxy (LBG) measurements, and suggests that stars may be responsible for cosmic reionization.

The origin of ultrahigh energy cosmic rays (UHECR) is one of the great remaining mysteries in astrophysics [238]. The cosmic-ray spectrum has been measured to beyond $10^{19}$ eV, with a number of events with energy exceeding $10^{20}$ eV [239, 240]; however, it is still debated how such highly energetic particles can be produced. It is now generally accepted that UHECR are of an extragalactic origin [92, 93]. However, above $\sim 10^{19.5}$ eV, the process of photopion production ($p \gamma \rightarrow N \pi$) on the cosmic microwave background (CMB) is expected to lead to a significant diminution of the cosmic-ray spectrum, the well-known GZK effect [100]. The relatively short attenuation length associated with the GZK process [241, 242] necessitates that the observed UHECR arise from local sources.

It has been noted that the source emissivity required to account for $\gtrsim 10^{19}$ eV cosmic rays is comparable to that of gamma-ray bursts [243, 95]. When GRBs are considered as the source of UHECR (with identical cosmic-ray production per burst), the change in the cosmological cosmic-ray emissivity is simply determined by the
burst rate history. As we will see, this GRB evolution is quite strong, as illustrated in Fig. 4.8, even exceeding that of models used in past cosmic-ray studies, which have traced quasar (QSO) luminosities [245] or the SFH.

The decay of charged pions produced in the GZK process results in a flux of ultra-high energy *cosmogenic* neutrinos [246, 247]. While the observed UHECR spectrum is somewhat insensitive to variations in cosmic source evolution [248], the cosmogenic neutrino flux can be greatly enhanced by strong evolution with redshift [247, 249], as neutrinos can be produced in larger quantities due to the decreased photopion threshold (since $T_{\text{CMB}} \propto 1 + z$), and can themselves traverse cosmological distances without attenuation. We examine the effect of enhanced GRB rate evolution on the expected flux of cosmogenic neutrinos, the measurement of which may provide the only way to break degeneracies between cosmic-ray models [250]. We find that this strong evolution leads to a measurable neutrino signal [86], improving the near-term prospects for assessing this scenario with upcoming ultrahigh energy neutrino detectors [251, 252, 253].

### 4.2 The Expected GRB Rate

An ever-increasing amount of data has brought about a clearer picture of the history of cosmic star formation. As shown in Fig. 4.2, after a sharp rise up to $z \approx 1$, the SFH is nearly flat until $z \approx 4$ with relatively small uncertainties. These measurements are well-fit by a simple piecewise power law parametrization [70],

\[
\dot{\rho}_{\text{SFH}}(z) \propto (1 + z)^{3.44} \quad : \quad z < 0.97
\]
\[
(1 + z)^{−0.26} \quad : \quad 0.97 < z < 4.48
\]
\[
(1 + z)^{−7.8} \quad : \quad 4.48 < z,
\]
scaled to $\dot{\rho}_{\text{SF}}(0) = 0.0197 M_\odot \text{ yr}^{-1} \text{ Mpc}^{-3}$ (a rate per *comoving* volume). We parametrize the intrinsic (comoving) GRB rate relative to the SFH as $\dot{n}_{\text{GRB}}(z) =$
$\mathcal{E}(z) \times \dot{\rho}_{\text{SF}}(z)$. The fraction of bursts that can be seen at a given $z$, $0 < F(z) < 1$, depends on the ability to detect the initial burst of gamma rays and to obtain a redshift from the optical afterglow. We cast the distribution of observable GRBs as

$$\frac{d\dot{N}}{dz} = F(z) \frac{\mathcal{E}(z) \dot{\rho}_{\text{SF}}(z) dV/dz}{\langle f_{\text{beam}} \rangle (1 + z)}.$$  \hspace{1cm} (4.2.1)
This being an observed rate, cosmological time dilation requires the \((1 + z)^{-1}\). The comoving volume element (in terms of the comoving distance, \(d_c\)), \(dV/dz = 4\pi (c/H_0) d_c^2(z)/\sqrt{(1 + z)^3 \Omega_m + \Omega_\Lambda}\), peaks at \(z \sim 2.5\). Dividing \(dV/dz\) by the \(1 + z\) term yields a *volumetric factor* that peaks at \(z \sim 1.4\) (see Fig. 4.2). Throughout, we use \(\Omega_\Lambda = 0.7\), \(\Omega_m = 0.3\), and \(H_0 = 70\) km/s/Mpc; changing these requires correcting the SFH (see [254]). For a constant \(F(z)\), relatively fewer bursts should be observed at \(z \sim 4\) than \(z \sim 1\). GRBs that are unobservable due to beaming are accounted for through \(\langle f_{\text{beam}} \rangle\).

4.3 The *Swift* Observations

*Swift* has enabled observers to extend the reach of GRB observations greatly compared to pre-*Swift* times, resulting in a rich data set. We first consider in detail a sample of bursts from the *Swift* archive [234] up to 2007 May 15 with reliable redshifts and durations exceeding \(T_{90} > 2\) sec, and later examine a larger set using the same techniques. We estimate each GRB luminosity as \(L_{\text{iso}} = E_{\text{iso}}/[T_{90}/(1 + z)]\), where \(E_{\text{iso}}\) is the isotropic equivalent (beaming-uncorrected) \(1 - 10^4\) keV rest-frame energy release and \(T_{90}\) is the interval observed to contain 90% of the prompt GRB emission. To form a uniform \(L_{\text{iso}}\) set, as displayed in Fig. 4.3, we use the \(E_{\text{iso}}\) and \(T_{90}\) values given by Ref. [235]. Note that \(L_{\text{iso}}\) for several bursts (open symbols) may be underestimated due to \(T_{90}\) values that overestimate the GRB duration.

To make the visual density of GRBs in Fig. 4.3 more meaningful, we define a linear \(x\)-coordinate that “corrects” for the effects of the volumetric factor as

\[
Q(z) = \int_0^z dz' \frac{dV/dz'}{1 + z'},
\]

which we will quote in terms of \((c/H_0)^3 \approx 79\) Gpc\(^3\). Removing the influence of the volumetric factor allows for a better “by-eye” view, since \(d\tilde{N}/dQ = (d\tilde{N}/dz)/(dQ/dz) = \)
Figure 4.3: The luminosities ($L_{\text{iso}}$) of 63 Swift gamma-ray bursts, determined from the data of [235], versus $z$ (two $z < 0.1$ GRBs are below $10^{48}$ erg s$^{-1}$). Points are plotted linearly in $Q(z)$, as given in Eq. (4.3.1), to account for the volumetric factor. Open symbols may have underestimated $L_{\text{iso}}$. The shaded region approximates the effective Swift detection threshold. Redshifts were measured as denoted, with medians of $z = 0.8$ for emission and $z = 2.9$ for absorption [77].

$$F(Q) \hat{n}_{\text{GRB}}(Q)/\langle f_{\text{beam}} \rangle,$$ so that a flat GRB rate would appear as a constant density of GRBs with similar $L_{\text{iso}}$ per linear interval in $Q$.

The Swift trigger is quite complex, working in the 15 – 150 keV band and using time-dependent background subtraction and variable time windows in order to maximize burst detection [76, 255]. While this sensitivity is “very difficult if not
Figure 4.4: The differential GRB distribution versus $Q$. Outlined bins contain the set of all 63 GRBs (median $z = 2.3$), while shaded bins contain the 44 bursts with $L_{\text{iso}} > 10^{51}$ erg s$^{-1}$ (with a median of $z = 2.9$). The rise seen suggests that GRB rate evolves more strongly than star formation [77].

impossible” [255] to parametrize exactly, an effective luminosity threshold appears to present in the data (roughly estimated as $\propto d_l^2$ in Fig. 4.3).

A representative sample of bursts in a given redshift range can be selected, while avoiding detailed assumptions concerning this threshold, by simply choosing a lower cutoff in $L_{\text{iso}}$. Considering only bursts that could have been seen from anywhere within the redshift range examined allows for the Swift contribution to the $F(z)$ term to be treated as constant in $z$. This technique effectively integrates the GRB
luminosity function, $dN/dL_{\text{iso}}$, above the chosen $L_{\text{iso}}$ cut without assuming its functional form. This reduces the problem to number counts.

### 4.4 Testing for Evolution

Because $\dot{\rho}_{\text{SF}}(z)$ is now reasonably well-measured from $z \approx 0 - 4$ (and nearly flat for $1 < z < 4$), we consider GRBs in this range for comparison. Using only bursts with $L_{\text{iso}} > 10^{51} \text{ erg s}^{-1}$ creates a set of GRBs with a minimal expected loss of GRBs up to $z \lesssim 4$. Fig. 4.4 displays the differential distribution of these GRBs versus $Q$. The bursts above this cut (shaded bins) can be compared to the set of all bursts (outlined). As can clearly be seen, removing the population of lower luminosity bursts that are only observable at low $z$ reveals a distinct rise in the observed number of “bright” bursts. The drop at $z \gtrsim 4$ is likely due to the *Swift* threshold and an overall drop in star formation.

To compare the GRB data with the SFH expectation, we make use of Eq. (4.2.1), parameterizing the “effective evolution”, for simplicity, as $F(z) \times E(z) \propto (1 + z)^\alpha$. Since the luminosity cut removes the influence of the *Swift* threshold, any $z$-dependence of $F$ would be due to other potential observational effects. We compare the predicted and observed cumulative GRB distributions. This compares the relative trends, independent of the overall normalization and the value of $\langle f_{\text{beam}} \rangle$. A Kolmogorov-Smirnov test reveals that the SFH fit alone is incompatible at around the 95% level.

Positive evolution results in clear improvement, strengthening indications based upon different analyses and smaller data sets (e.g., [85, 86, 87]). We find that the K-S statistic is minimized for $\alpha = 1.5$. To simulate dispersion in the data, we replace the $T_{90}$ values used to calculate $L_{\text{iso}}$ with those reported by *Swift* [234] and find similar conclusions. In neither case is any significant correlation of luminosity with $z$ found in this range, disfavoring just an evolving luminosity function [256]. Although we have
attempted to minimize the loss of GRBs due to the *Swift* threshold, we are more likely to be underestimating burst counts at the upper end of the z-range examined, which could imply even stronger evolution.

### 4.4.1 Is the Evolution Intrinsic?

This evolution could result from several causes. The K-S test disfavors an interpretation as a statistical anomaly. While \( \dot{\rho}_{SF}(z) \) could just be mismeasured, the relatively-small uncertainties over the range in question suggests that this is not the likely origin. We will discuss in more detail whether it is due to some selection effect (i.e., a changing \( F(z) \)), or a real physical effect related to changes in the number of progenitors (an evolving \( \mathcal{E}(z) \)).

It has been suggested that a higher percentage of GRBs may go undetected at low redshifts [257]. We perform several simple diagnostics to determine whether the sample of *Swift* GRBs without redshifts contains an overwhelming number of bright, low-z bursts. To simplify analysis, we consider a sample of bursts that meet the “detectability” criteria of [258] (in particular, low Galactic extinction and quick *Swift* localization), a set of 50 GRBs with a confirmed \( z \) and 47 without. With this set, using the procedure of Sect. 4.4, the SFH expectation alone is still incompatible at around 95\%, with \( \alpha = 1.5 \) minimizing the K-S statistic.

On average, more distant bursts are expected to have lower observed gamma-ray fluxes, \( \mathcal{F} \), estimated by dividing each \( 15 - 350 \) keV fluence and \( T_{90} \) from [235]. Indeed, GRBs at \( z > 2 \) have \( \langle \mathcal{F}_{z>2} \rangle \sim 1.4 \times 10^{-7} \) erg cm\(^{-2}\) s\(^{-1}\), compared to \( \langle \mathcal{F}_{z<2} \rangle \sim 2.6 \times 10^{-7} \) erg cm\(^{-2}\) s\(^{-1}\) at \( z < 2 \). For GRBs without a redshift, the average flux is just \( \langle \mathcal{F}_{\text{no } z} \rangle \sim 0.7 \times 10^{-7} \) erg cm\(^{-2}\) s\(^{-1}\). A two-sample K-S test between the \( z < 2 \) and \( z \)-less sets reveals that they are incompatible at \( \sim 70\% \). Limiting the \( z < 2 \) set only to GRBs with \( L_{\text{iso}} > 10^{51} \) erg s\(^{-1}\), this increases to \( \sim 99\% \). While not conclusive, these results
could be interpreted as most $z$-less bursts being either at high $z$, in which case our evolution may be underestimated, or at low $z$ with lower intrinsic luminosities, which may not survive our $L_{\text{iso}}$ cut. Additionally, we examine the fractions of GRBs detected by Swift’s UV-Optical Telescope in the Swift archive. Of bursts with a known $z$, 34 were seen by UVOT with $\langle z \rangle = 2.2$, while the 16 not seen have $\langle z \rangle = 3.0$. Bursts lacking redshifts seem more consistent with the high-$z$ set, with only 8/47 seen by UVOT.

We also consider whether the fraction of observable bursts, $F(z) = f_z f_{\text{Swift}}$, is somehow increasing with $z$. While $f_{\text{Swift}}$ is difficult to quantify, our selection criteria disfavor incompleteness of our sample at low $z$. We focus upon the probability of determining a redshift for a given GRB, which we subdivide as $f_z = f_{E/A} f_{AG} f_{\text{human}}$. Perhaps the most obvious influence on this term is the fact that at low $z$, most redshifts are determined by observing emission lines (from the host galaxy); while at higher redshifts, nearly every redshift is found through absorption lines in the afterglow spectrum (see Fig. 4.3 and the Swift archive [234]). While this emission/absorption bias, $f_{E/A}$, is not easy to quantify, it should not cause such an evolutionary effect, since most redshifts in our sample are found through absorption.

The observability of an optical afterglow, $f_{AG}$, might be expected to be steeply falling with redshift; however, for a spectrum $\propto t^{-\alpha} \nu^{-\beta} [259]$, cosmological redshifting may allow the earliest (brightest) portion of the afterglow to be more visible. If the flatness of the obscuration-corrected SFH at moderate $z$ arose from a steeply increasing uncorrected rate and a decreasing dust correction, the detectable afterglow fraction might increase with redshift; this does not appear to be the case [260]. A paucity of IR detections of hosts of “dark” GRBs also argues against a significant obscured fraction [73]. The decision of which telescopes are made available to observe GRBs, along with other such non-intrinsic properties, can be folded into the human
factor, $f_{\text{human}}$. None of these terms appear to be able to increase overall observability with $z$, disfavoring an origin of the trend in $F(z)$.

### 4.4.2 Potential Sources of Evolution

We now investigate whether an evolving $\mathcal{E}(z)$ can explain the observed evolution. While it is now generally accepted that long GRBs arise from massive stars, the special conditions that are required for such an event are still in question. In the collapsar model, the collapse of a rapidly-rotating, massive stellar core to a black hole powers a jet that is seen as a GRB. Since every observed supernova coincident with a GRB has been of Type Ic, the progenitor star should also have lost its outer envelope (without losing precious angular momentum). Rather strong observational evidence now indicates that GRB host galaxies tend to be faint and metal-poor (e.g., [62, 61, 72, 60]), increasing interest in models that use single, low-metallicity stars as a pathway to a collapsar [64]. Decreasing cosmic metallicity may cause the GRB/SN ratio to rise with $z$. While a preference for low-metallicity environments may be the simplest explanation, absorption line metallicity studies remain inconclusive. One may therefore wonder whether it is possible to concoct an evolutionary scenario without direct progenitor metallicity dependence.

If GRBs are instead produced in binary systems (e.g., [261]), some other mechanism might lead to the appearance of evolution. Since most GRBs appear to occur in star clusters [62], where the fraction of massive stars in binaries may be high [262], such channels could be important. For example, the merger of two $\gtrsim 15 M_\odot$ stars with $M_1/M_2 \gtrsim 0.95$ (i.e., “twins”, which may be common [263]) in close orbits ($r \lesssim \text{few AU}$) can lead to a more-massive, rapidly-rotating core lacking an envelope [264]. In a dense cluster, close binaries tend to end up closer due to gravitational scatterings with interloping stars (Heggie’s Law [265]), with a scattering timescale
of \( \lesssim 10 \) Myr for a stellar density of \( \rho \sim 10^6 M_\odot \text{pc}^{-3} \) [266]. An increased rate of “interloper-catalyzed” binary mergers (ICBMs) could result from a larger fraction of star formation occurring in such environments at higher \( z \), and seen as an enhancement in the GRB rate. Such a speculative scenario presents a dynamical source of evolution (instead of altering the microphysics) and, as these rates are \( \propto \langle \rho^2 \rangle \), possibly allow for examination of the “density contrast” of star formation. It is interesting that GRBs were discovered in searches for gamma rays from explosions related to ICBMs of a different sort [56].

An evolving IMF, becoming increasingly top-heavy at larger \( z \), would increase the relative number of massive stars produced. Since star formation measurements are primarily based on radiation from such stars, this alone may not lead to apparent evolution, unless the very massive end (\( \gtrsim 25 M_\odot \)) changed significantly. Any evidence of evolution in the IMF provides motivation for considering such a change. However, this need only occur in those galaxies that host gamma-ray bursts.

4.4.3 Confirmation of the Evolutionary Trend

As shown in Fig. 4.1, many more GRBs have been detected subsequent to May 2007, warranting a reexamination of our result. This sample includes the 63 GRBs used above, supplemented by 56 subsequent \textit{Swift} events with redshifts, with calculated \( L_{\text{iso}} \) from the updated catalog of [235, 236]. Using GRBs in the range \( z = 0 - 4 \) with \( L_{\text{iso}} > 10^{51} \text{ erg s}^{-1} \) leaves us now with 66. The cumulative distribution of the bursts is shown in Fig. 4.5. A Kolmogorov-Smirnov test confirms that the GRB rate is incompatible with the expectations from the SFR, now at the \( \sim 99\% \) level (possibly higher due to likely missing bursts at \( z \lesssim 4 \)) with the present greater statistics, requiring an enhanced evolution relative to the SFR. Even if we exclude the range
Figure 4.5: The cumulative distribution of the 66 Swift GRBs with $L_{\text{iso}} > 10^{51}$ erg s$^{-1}$ in $z = 0 - 4$ (solid), as compared to the expectations from the SFH of [70] alone (dashed) and additional evolution of the form $(1 + z)^{1.2}$ (dotted). Outside of the shaded region (bounded by models with $\alpha = 0.6$ and 1.8) corresponds to an exclusion of $> 84\%$ [88].

$z = 1.5 - 2$, where a larger fraction of redshifts might be missed [257], the value remains at $\sim 98\%$.

This result suggests a slightly lower value of $\alpha$. This may be due in part to the rate of GRB observations at higher redshifts decreasing noticeably in the period following the cutoff date for our initial GRB sample (for reasons unknown). Fortunately, as can be seen in Fig. 4.1, high-$z$ GRBs detections have since increased (denoted as the period after the start of Fermi operations). Irrespective of the origins of this bias, it must be accounted for in $z = 1 - 4$ to properly relate the GRB rate to the SFR. In Fig. 4.5, we show a shaded band bounded above by a model using $\alpha = 0.6$ and below
by $\alpha = 1.8$, which can be excluded at $\gtrsim 84\%$. To be conservative, we will assume this evolution continues to higher $z$, considering $\alpha = 1.2$ in our estimates of the high-$z$ SFR.

### 4.5 The GRB-SFR Technique

It is easy to understand, with the combination of uncertain extinction corrections, cosmic variance, and selection biases, why measurements of the SFR at high redshifts are difficult endeavors. Principal among these is that flux-limited surveys observe the bright end of the galaxy luminosity function (LF) and must correct for the faint end, where much of the star formation may be occurring. The use of gamma-ray bursts as a star formation measure will have its own systematic effects; however the opportunity presented to examine very-high redshifts, and possibly unseen faint galaxies, is great, with no known backgrounds for a bona fide GRB.

We here review the framework for calibrating a GRB-based estimate of $\dot{\rho}_*$. This is based on using GRB and SFR measurements in $z = 1 - 4$ as benchmarks for comparison with bursts of similar luminosity in a higher-$z$ range. Using only GRBs that could have been detected from anywhere within the volume allows for needed empirical calibration, since neither the conversion from GRB rate to SFR nor the GRB luminosity function are known a priori. Again, part of the challenge is in determining the detection threshold versus $z$, since Swift was designed to maximize GRB detection, though not necessarily in a way well-defined for our purpose [255]. We show in Fig. 4.1 an estimated threshold ($\propto d_L^2$) based on the GRB luminosities, which acts as a guide to make cuts that maximize statistics and minimize potential “missing” bursts.

The cuts and resulting subsamples used for the SFR analysis are shown in Fig. 4.1. Note that we exclude GRB 060116 (not shown), with a possible photometric redshift.
of $z = 6.6$ [235]. Fig. 4.6 shows these in comparison to the distribution of $L_{\text{iso}}$ values for bursts in $z = 1 - 4$. Bursts within each set will be compared to GRBs within the range $z = 1 - 4$ above the given luminosity cut. We emphasize in advance that, although the final bin contains only GRB 090423 (at $z \simeq 8.1$), even this single event is significant, as it would be quite unlikely if $\dot{\rho}_s$ was too low. With this GRB we are entering a regime where the age of the star is becoming non-negligible compared to the age of the Universe, so we extend this bin to $z = 8.5$ to cover a plausible range in progenitor lifetime.

The “expected” number of GRBs in $z = 1 - 4$ is

$$N_{1-4}^{\text{exp}} = \Delta t \frac{\Delta \Omega}{4\pi} \int_1^4 dz F(z) E(z) \frac{\dot{\rho}_s(z)}{\langle f_{\text{beam}} \rangle} \frac{dV/dz}{1+z},$$

in which $A = \Delta t \Delta \Omega E_0 F_0/4\pi \langle f_{\text{beam}} \rangle$ depends on the observing time ($\Delta t$), sky coverage ($\Delta \Omega$), and luminosity range of GRBs under examination. From the average SFR, $\langle \dot{\rho}_s \rangle_{z_1-z_2}$, the same can be performed for the other ranges as

$$N_{z_1-z_2}^{\text{exp}} = \langle \dot{\rho}_s \rangle_{z_1-z_2} A \int_{z_1}^{z_2} dz (1+z)^\alpha \frac{dV/dz}{1+z}. \quad (4.5.2)$$

Our interest is in finding $\langle \dot{\rho}_s \rangle_{z_1-z_2}$ by dividing out $A$ (using Eq. 4.5.2). Taking the measured GRB counts, $N_{z_1-z_2}^{\text{obs}}$, to be representative of the expectations, $N_{z_1-z_2}^{\text{exp}}$, we find

$$\langle \dot{\rho}_s \rangle_{z_1-z_2} = \frac{N_{z_1-z_2}^{\text{obs}}}{N_{1-4}^{\text{exp}}} \int_1^4 dz \frac{dV/dz}{1+z} \frac{\dot{\rho}_s(z)}{\langle f_{\text{beam}} \rangle} (1+z)^\alpha.$$  

Note that the decrease of $(dV/dz)/(1+z)$ at $z \gtrsim 1.5$ (as shown in Fig. 4.2) gives progressively more weight to each observed higher-$z$ GRB.

We show our new determinations of the high-$z$ SFR in Fig. 4.7 (assuming a Salpeter IMF [267]). Error bars correspond to 68% Poisson confidence intervals for the binned events [268]. We also show as a shaded band the values obtained for
different assumptions of $\alpha$, bounded above by $\alpha = 0.6$ and below by $\alpha = 1.8$, which yields an uncertainty smaller than the statistics in the last bins. Variations due to changing the $L_{\text{iso}}$ cutoff can be determined from Fig. 4.6, which will typically be less than the statistical uncertainties. We have been generally conservative and have also verified that using another luminosity estimator, the peak isotropic equivalent luminosity, yields similar results. Other effects, including the selection of $z$-ranges and the inclusion/exclusion of particular bursts, are discussed in Ref. [79]. We mention only that none of these affect the basic point that the SFR must be large enough to produce the observed GRB counts.

Depending upon the source of the evolution, our bias correction may be unduly
underestimating $\dot{\rho}_s$ by a factor of a few at higher $z$. The most likely astrophysical explanation is due to metallicity. GRBs are found to favor metal-poor [61], sub-$L_*$ galaxies [60, 72, 62], so having a larger fraction of the SFR within such hosts would result in a higher GRB rate. This could be the case with a steepening faint-end slope of the galaxy LF, so that more of $\dot{\rho}_s$ arises from below $L^*_z$ ($L_*$ as defined at $z$). This has been observed between $z = 0$ and $z \approx 2 - 3$ (see Fig. 7 of [269]).

While our result at $z = 4 - 5$ is in basic agreement with earlier measurements, at the highest-$z$ ranges, LBG studies probe only the brightest galaxies and must estimate the faint end of the UV LF based on limited data. Our results diverge from these if corrections for unseen galaxies are not made. For example, we focus upon the measurements in Refs. [80, 81], which are reported (lower triangles in Fig. 4.7) for an integration down to $0.2L^*_{z=3}$ (with their adopted dust corrections). Fully integrating their UV LFs (which can be regarded as giving a maximum), with faint-end slopes of -1.73, -1.66, -1.74, -1.74 for $\langle z \rangle = 3.8, 5.0, 5.9, 7.3$, respectively, yields the upper set of triangles.

Within the uncertainties, even the highest redshift fully-integrated point now agrees reasonably well with our results, and the preference of GRBs for faint galaxies (although the exact relation between GRB hosts and star forming galaxies as a whole remains to be determined). We note that the LF slope of [81] at $\langle z \rangle = 7.3$ was taken to be the same as at $\langle z \rangle = 5.9$. If the slope is actually steeper, then these measurements could be higher, although it is difficult to draw definite conclusions, due to the limited statistics and uncertainties in dust corrections.

4.6 Implications for Reionization

Transmission in the Gunn-Peterson troughs of high-redshift quasars implies that reionization must have been accomplished before $z = 6$ [271]. AGN seem to be
Figure 4.7: The cosmic star formation history. Shown are the data compiled in [70] (light circles) and contributions from Lyα Emitters (LAE) [82]. Recent LBG data is shown for two UV LF integrations: down to $0.2L^*_{z=3}$ (down triangles; as given in [81]) and complete (up triangles). Our (bias-corrected) Swift gamma-ray burst inferred rates are diamonds, with the shaded band showing the range of values resulting from varying the evolutionary parameter between $\alpha = 0.6 - 1.8$. Also shown is the critical $\dot{\rho}_*$ from [270] for $C/f_{\text{esc}} = 40, 30, 20$ (dashed lines, top to bottom) [88].

insufficient for this purpose [272], leaving stars as the leading candidate. To address the ability of an observed population to reionize the Universe, Ref. [270] provided an estimate for the required SFR to balance recombination, $\dot{\rho}_c$, which depends upon the
fraction of photons that escape their galaxy ($f_{esc}$) and the clumpiness of the IGM ($C$) as

$$\dot{\rho}_c(z) = \frac{0.027 M_\odot}{\text{yr}} \frac{C/f_{esc}}{30} \left[\frac{1 + z}{7}\right]^3 \left[\frac{\Omega_b}{0.0465}\right]^2.$$  (4.6.1)

For comparison with our empirical SFR, we show in Fig. 4.7 curves of $\dot{\rho}_c$ as a function of $z$ for $C/f_{esc} = 40, 30, \text{and 20}$. We find that our SFR estimates can exceed the $\dot{\rho}_* \text{ required to keep the Universe ionized at redshifts as high as } z \gtrsim 8$. However, this criterion refers to an instantaneous equilibrium, and so does not address the requirement that the integrated number of ionizations exceed the number of hydrogen atoms.

Ref. [273] estimated the ionizing emissivity at $z \sim 5$ from the Ly$\alpha$ forest, and looked at simple models of the reionization history under different assumptions for the evolution of the ionizing photon emissivity at $z \gtrsim 5$. While their estimate of emissivity from the measured ionization rate is sensitive to the calculation of mean-free path, Ref. [273] reached the strong conclusion that reionization must have been an extended “photon starved” process, and that an emissivity which was constant towards higher $z$ would have been insufficient to reionize the Universe by $z \sim 6$. This implies that the ionizing emissivity must have been higher prior to the end of reionization than just after its conclusion. The origin of this higher emissivity could lie in an increase in one or all of the SFR, the escape fraction, or the fraction of massive stars in the IMF. Inspection of Fig. 4.7 suggests that the SFR is as large at or could even be higher at $z \sim 8$ than at $z \sim 6$, implying that the ionizing photon emissivity may not be falling towards redshifts greater than $z \sim 6$. Since both ionizing photons and GRBs are produced by massive stars, estimates of the ionizing photon emissivity from the GRB rate should be fairly robust against uncertainties in the IMF at the high-mass end.
We are therefore motivated to ask whether we have observed enough star formation at $z > 6$ to reionize the universe. To answer this question, we make a simple estimate, calculating the number of ionizing photons produced prior to $z \sim 6$ given the observed SFR. For a Salpeter IMF [267] and a metallicity of 1/20 Solar, $\sim 4600$ ionizing photons are produced per baryon incorporated into stars [78] (further details are given in [89]). Taking this value, together with a constant SFR for a time interval $\Delta t$, we find the number of photons produced per hydrogen in the IGM as

$$N_{\gamma} \sim 4 \left( \frac{f_{\text{esc}}}{0.1} \right) \left( \frac{\rho_*}{0.1 \, M_\odot \, \text{Mpc}^{-3} \, \text{yr}^{-1}} \right) \left( \frac{\Delta t}{400 \, \text{Myr}} \right). \quad (4.6.2)$$

Given our SFR, this implies that $N_{\gamma} \sim 3^{+3}_{-1.5} (f_{\text{esc}}/0.1)$.

In order to reionize the Universe, more than one ionizing photon per baryon is required to compensate for recombinations in the ionized IGM. Ref. [274] modeled the reionization history including evolution of the clumping factor with the restriction that reionization end at $z \sim 6$. These models yielded $N_{\gamma} \sim 4$ at $z = 6$ under a range of assumptions for the redshift range and efficiency of Population-III star formation. This is within our estimated range, provided that the escape fraction of ionizing radiation is of order 10%. This value of escape fraction lies in the range found by Ref. [275], who combined semi-analytic models of reionization and the galaxy LF with simulations of the transmission in the high-$z$ Ly-$\alpha$ forest. While not the final word, our results may thus indicate that stars produced enough ionizing photons in the range $6 < z < 8$ to reionize the Universe.

### 4.7 GRB Rate Evolution and UHECR

Any treatment of cosmological cosmic-ray or neutrino production must account for source evolution. Traditionally, the most commonly used cosmic-ray evolution models have tracked quasars [245] or the SFH [70] (particularly in GRB-related studies). As
we saw above, observations indicate, however, that GRBs do not faithfully trace the SFH, with the GRB rate density actually evolving more strongly than the standard SFH. Accounting for stellar evolution effects, an expected ratio of the rate of GRBs to core-collapse supernovae as a function of redshift was calculated in Ref. [75], which we approximate as \( \dot{n}_{\text{GRB}}(z) \propto (1 + z)^{1.4} \dot{n}_{\text{SN}}(z) \). This is consistent with our estimate of this ratio and will be used to calculate the increase in the absolute GRB rate density (of importance to cosmic-ray studies). Combining the piecewise-linear Hopkins and Beacom SFH [70] with the parametrized form of the GRB/SN ratio, we find the source evolution term [86], \( W_{\text{GRB}}(z) \), to have the form

\[
W_{\text{GRB}}(z) \propto \begin{cases} 
(1 + z)^{4.8} & : \ z < 1 \\
(1 + z)^{1.4} & : \ 1 < z < 4.5 \\
(1 + z)^{-5.6} & : \ 4.5 < z,
\end{cases}
\]

(4.7.1)

with \( \dot{n}_{\text{GRB}}(z) = \dot{n}_{\text{GRB}}(0) \times W_{\text{GRB}}(z) \). In Fig. 4.8, we present this evolution model, in comparison to the SFH alone, as well as the quasar (QSO) evolution model used in Refs. [244, 249]. Note that our rate history evolves as \( W_{\text{GRB}}(z) \sim (1 + z)^{4.8} \) for \( z < 1 \), which is substantially steeper than the models commonly considered. In particular, the QSO model, which has been extensively used, only rises as \( W_{\text{QSO}}(z) \sim (1 + z)^3 \).

One consequence of the strong evolution that we present is a predicted burst rate density of \( \dot{n}_{\text{GRB}}(z = 0) \sim 4 \text{ Gpc}^{-3} \text{ yr}^{-1} \) in the local universe (when normalized to the observed redshift distribution), which quickly rises to \( \sim 20 \text{ Gpc}^{-3} \text{ yr}^{-1} \) by redshift \( z \sim 0.4 \). The average emissivity from these bursts (at \( z \lesssim 0.4 \)) is \( E_{\text{GRB}} \sim \epsilon_\gamma \times 10 \text{ Gpc}^{-3} \text{ yr}^{-1} \sim 5 \times 10^{43} \text{ erg Mpc}^{-3} \text{ yr}^{-1} \), which is comparable to the emissivity required to account for the \( \gtrsim 10^{19} \text{ eV} \) UHECR flux, \( E_{\text{CR}} \sim \text{few} \times 10^{44} \text{ erg Mpc}^{-3} \text{ yr}^{-1} \) (as found in Ref. [95]). Considering that there is no \textit{a priori} reason for these numbers to be so similar, along with the fact that the region of the cosmic-ray spectrum of greatest interest (\( \gtrsim 10^{19} \text{ eV} \)) must arise from sources at \( z \lesssim 0.4 \) (as we will discuss...
Figure 4.8: Models of cosmic-ray source evolution (i.e. yield vs. \( z \), normalized to 1 at \( z = 0 \)). From top-to-bottom, the GRB rate density, the quasar (QSO) evolution model used in Refs. [244, 249], and the fit to the cosmic star formation history (SFH) of Ref. [70]. Models similar to the latter two have been frequently used in cosmic-ray studies [86].
in Section 4.8), this result is quite interesting. GRBs may also be more luminous in cosmic rays due to a baryon loading factor ($f_{\text{CR}}$) that may be $\gtrsim 10$ [276], which could account for any difference. These considerations, along with the isotropy of the measured cosmic-ray spectrum, allow for GRBs to be further examined as a candidate to produce the observed UHECR.

4.8 UHECR: Propagation and Spectrum

Cosmic rays mainly lose energy through three processes during propagation. At very high energies ($\gtrsim 3 \times 10^{19}$ eV), UHECR energy loss is dominated by photopion production on the CMB, $p\gamma \rightarrow N\pi$, which has a maximum cross section at the $\Delta(1232)$ resonance [241, 277]. Below the photopion threshold (and with $E_p \gtrsim 10^{18}$ eV), electron-positron pair production on diffuse photons, $p\gamma \rightarrow pe^+e^-$, becomes important [278]. While the cross section for this process is high, each interaction results in only a small energy loss. Finally, the expansion of the universe results in adiabatic energy loss, which is independent of energy.

We can account for the total proton energy loss through the characteristic timescales associated with each process (as in Refs. [243, 279]):

$$
\tau_T^{-1}(E_p, z) = \tau_{\pi}^{-1}(E_p, z) + \tau_{\text{pair}}^{-1}(E_p, z) + \tau_a^{-1}(z).
$$

The rate of energy loss for propagating protons is then $d\ln E_p/dt = \tau_T^{-1}(E_p, z)$. With the redshift-dependent energy loss rate determined, the injection energies of cosmic rays, $E'_p = E'_p(E_p, z)$, can be calculated as a function of detected energy ($E_p$) and originating redshift ($z$) through the differential equation

$$
\frac{1}{E_p} \frac{dE_p}{dz} = \frac{1}{\tau_T(E_p, z)} \frac{1}{dz/dt}
$$

where $dz/dt = H_0(1+z)[\Omega_M(1+z)^3 + \Omega_\Lambda]^{1/2}$ (with $\Omega_M = 0.3$, $\Omega_\Lambda = 0.7$, and $H_0 = 70$ (km/s)/Mpc).

The top panel of Fig. 4.9 shows cosmic-ray energy loss with redshift. The lines
correspond to the injection energy needed at redshift \( z \) in order for a cosmic ray to be observed with a given energy at Earth \((z = 0)\). For example, a cosmic-ray proton with a measured energy of \( 10^{19} \) eV must have been produced at \( z < 0.4 \). This effect is analogous to the Fazio-Stecker relation for gamma rays [280]. As energy losses above \( 3 \times 10^{19} \) eV are quite severe, the observed UHECR must be produced in the local universe. This, along with the necessary conditions required to produce UHECR at all, strongly constrains the population of prospective sources.

In evaluation of the cosmic-ray spectrum, we have assumed that, at these energies, the cosmic-ray spectrum is entirely composed of protons with an injection spectrum \( \varphi(E'_p) = N E'_p^{\alpha-\gamma} \) (per unit comoving volume per unit energy, per unit time), which is cut-off at a chosen \( E_{\text{cut}} \) and normalized to account for the observed spectrum of UHECR. Cosmic rays injected at \( E'_p \), will experience energy losses and be detected at \( E_p \). Taking into account the evolution of sources, \( W(z) \), we can calculate the UHECR spectrum as,

\[
\frac{dN_p}{dE_p} = \frac{c}{4\pi} \int_{0}^{z_{\text{max}}} \varphi(E'_p) \frac{\partial E'_p}{\partial E_p} W(z) dz / dt dz. \tag{4.8.2}
\]

We calculate the partial derivative \( \partial E'_p(E, z)/\partial E_p \) numerically from Eq. 4.8.1. In the bottom panel of Fig. 4.9, we present the expected cosmic-ray spectrum for the three source evolution models presented in Fig. 4.8, with an injection spectrum of the form \( E^{-2} \) (plotted as \( E^3 \times \text{Flux} \) to emphasize spectral features [281]). We choose the normalization of \( \varphi, N \), such that \( E^3 \times dN_p/dE_p = 2 \times 10^{20} \) eV\(^2\) cm\(^{-2}\) s\(^{-1}\) sr\(^{-1}\) at \( 10^{19} \) eV for these spectra (and in our subsequent results), as the uncertainties at higher energies await the resolution that will be delivered by Auger [240]. This usually necessitates the local emissivity of sources between \( 10^{19} - 10^{21} \) eV to be on the order of \( \mathcal{E}_{\text{CR}} \sim 5 \times 10^{44} \) erg Mpc\(^{-3}\) yr\(^{-1}\) [243, 248].

Note that the data from the AGASA, HiRes and Yakutsk experiments [239] differs by almost a factor of 2 in overall normalization in our \( E^3 \) plot (light shaded region),
Figure 4.9: Top: Cosmic-ray energy loss with redshift. Lines illustrate the injection energy required at redshift $z$ in order to be detected at a given energy at $z = 0$. Shown for illustrative purposes are the realms of photopion (dark-shaded region) and pair-production (light-shaded region) losses. Bottom: Cosmic-ray spectra expected from the GRB (solid), QSO (dashed), and SFH (dotted) source evolution models, assuming $\gamma = 2$ and normalization at $E = 10^{19}$ eV. All curves are well within range of experimental data above $10^{19}$ eV (light-shaded region, dark-shaded when normalized to the spectral dip) [86].
however, the overall spectra do not significantly disagree [282] and model-independent information can be extracted from the shape alone, especially when various experimental data are normalized to the spectral dip feature at $\sim 10^{19}$ eV [279] (dark shaded region).

When considering the cosmic-ray spectrum resulting from GRB evolution, we have followed the general form of Waxman by assuming that the entire UHECR spectrum above $10^{19}$ eV can be accounted for by GRBs, with an $E^{-2}$ injection spectrum [243, 95]. However, for an injection spectrum of this form, it is challenging to account for cosmic rays with energies $\lesssim 10^{19}$ eV [283]. Strong source evolution with redshift offers a certain degree of assistance. It has been assumed previously that an extension of the Galactic cosmic-ray spectrum may account for any remaining deficiency [284]. Contributions from extragalactic sources with a lower energy reach may also be considered (e.g., AGN, cluster shocks [285]). Alternatively, one can consider an $E^{-2.2}$ injection spectrum (as might be expected from relativistic shocks), which may offer a better fit at lower energies. This would require a larger $f_{CR}$, which itself may explain GRBs with low radiative efficiencies.

It is important to note that the cosmic-ray spectrum itself depends only mildly upon evolution (since the observed UHECR with $E \gtrsim 10^{19}$ eV must be produced at $z \lesssim 0.4$) and may be explained by a combination of GRBs and a lower-energy component. However, the strong evolution implied by the metallicity-biased GRB model results in an increased flux of cosmogenic neutrinos. This neutrino flux, which must be present if GRBs are to account for the observed UHECR, provides an independent test of the GRB–UHECR model, which we will now examine.
4.9 Implications For Cosmogenic Neutrinos

The flux of cosmogenic neutrinos produced via the GZK process is quite sensitive to cosmic-ray source evolution [250]. This is due to the unique ability of neutrinos to propagate through large distances at very high energies without appreciable energy loss (unlike cosmic rays), combined with a lower threshold for photopion production at larger \( z \) (since \( T_{\text{CMB}} \propto 1+z \)). In fact, a significant portion of the expected neutrino flux originates between from \( z \sim 1 \to 4 \).

Near the photopion production threshold, about 20\% of the original proton energy is lost in each interaction. Neutrinos are subsequently produced through the decay chain, \( \pi^+ \to \mu^+ \nu_\mu \to e^+ \nu_e \nu_\mu \), with each daughter neutrino \( (\nu_e, \nu_\mu, \bar{\nu}_\mu) \) receiving \( \sim 1/20 \) of the parent proton energy [286]. At higher energies, multi-pion production dominates [277]; however, this approximation remains viable, as the inelasticity of each interaction increases (approaching 50\%).

Waxman and Bahcall (WB) have presented an upper bound on cosmogenic neutrino production (shown in Figs. 4.10 and 4.11 as a shaded band) based on the assumption of an \( E^{-2} \) injection spectrum, with normalization chosen between \( 10^{19} \) eV to \( 10^{21} \) eV to produce the observed cosmic-ray spectrum [244]. This yields an energy-dependent rate of cosmic-ray generation of \( N_{\text{WB}} \sim 10^{44} \) erg Mpc\(^{-3}\) yr\(^{-1}\) (with \( E_{\text{WB}} \sim 5 \times N_{\text{WB}} \)). The total \( \nu_\mu + \bar{\nu}_\mu \) energy flux at Earth (not corrected for oscillations) is estimated as

\[
E_\nu^2 \frac{dN_\nu}{dE_\nu} \approx \frac{c}{4\pi} N_{\text{WB}} \frac{1}{4} \xi t_H \approx \xi \times 15 \frac{\text{eV}}{\text{cm}^2 \text{s sr}},
\]

where \( t_H \approx 10^{10} \) yr is the Hubble time and the factor \( 1/4 \) arises from the assumption that only one quarter of the energy lost is carried away by muon neutrinos. Adiabatic redshift losses and the effect of source evolution are taken into account by \( \xi \) (estimated to be \( \xi \sim 0.6 \) with no source evolution, \( \sim 3 \) with QSO-like evolution) [244].

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Figure 4.10: Expected (all-flavor) cosmogenic neutrino fluxes assuming various evolution scenarios. From top-to-bottom, are the fluxes resulting from the strongly evolving (metallicity-dependent) GRB rate density, QSO-like evolution, and the SFH. Shown for comparison are the Waxman-Bahcall bound (shaded band) and the expected sensitivities for ANITA (diamonds) and ARIANNA/SalSA (triangles) [86].

The neutrino spectrum produced in the GZK process may be better approximated for a general cosmic-ray injection spectrum and source evolution through a somewhat...
more sophisticated approach. As the energy loss distance at these energies is relatively short, we assume that cosmic rays lose all of their energy rapidly. The fraction of the original proton energy that is lost to neutrinos can then be parametrized with a gradual step function, \( S(E) = 0.45/(1 + (E_t/E)^2) \), where 0.45 is the asymptotic fraction of injected cosmic-ray energy transferred to neutrinos above \( \gtrsim 10^{21} \) eV (as shown in Fig. 1 of Ref. [249]) and \( E_t \sim 2 \times 10^{20} \) eV governs the onset of photopion production above \( \sim 3 \times 10^{19} \) eV. The total neutrino flux at Earth can be cast as

\[
\frac{dN_\nu}{dE_\nu} = \frac{c}{4\pi} \int_0^{z_{\text{max}}} 20 \varphi(E'_p) S[(1 + z) \times E'_\nu] \frac{dE'_p}{dE_\nu} \frac{dW(z)}{dz} dz,
\]

where \( E'_p = 20(1+z)E'_\nu \) and the factor 20 reflects the approximation of each daughter neutrino receiving about 1/20 of the injected proton energy. The additional factor of \((1 + z)\) in \( S \) accounts for the lowering of the photopion energy threshold as the CMB temperature increases at higher redshift. This simple formulation is similar to the notation based on neutrino yield functions often used in prior studies [249, 287, 288] and provides a reasonably accurate neutrino spectrum which agrees rather well with the literature (in the energy range most interesting to UHE neutrino detectors), with a deviation from the simulated spectra not larger than the variations typically seen between such simulations. The overall normalization (set by \( \varphi \)) is again chosen such that the corresponding predicted cosmic-ray flux agrees with the measured cosmic-ray data at \( 10^{19} \) eV (where it is well-determined), as discussed in Section 4.8.

Eq. (4.9.2) can be derived using a simple, analytic method in order to compute the cosmogenic neutrino spectrum observed at Earth, which is based upon the assumption that each daughter neutrino resulting from a photopion interaction receives \( \sim 1/20 \) of the parent proton energy. That is, a neutrino detected with energy, \( E_\nu \), was produced with energy, \( E'_\nu = (1 + z)E_\nu \), at the source, from a parent proton which had energy, \( E'_p = 20 E'_\nu = 20(1 + z)E_\nu \). We predict the observed \((z = 0)\) neutrino spectrum,
\[ \frac{dN_\nu}{dE_\nu} \text{, by integrating over the contributions to the flux from redshifts up to } z_{\text{max}} \text{ as} \]
\[ \frac{dN_\nu}{dE_\nu} = \frac{c}{4\pi} \int_0^{z_{\text{max}}} dN_\nu' dE_\nu' \frac{W(z)}{dE_\nu'/dE_\nu} \frac{dz}{dt} d\Omega \quad (4.9.3) \]

where \( dE_\nu'/dE_\nu = (1 + z) \) accounts for the fact that neutrinos are observed in a narrower energy bin than they were originally produced (due to redshifting). Evolution in the density of sources is accounted for through \( W(z) \). The neutrino spectrum at production, \( dN_\nu/dE_\nu \), can be directly related to the cosmic-ray injection spectrum \( \varphi(E'_p) \) (per unit comoving volume per unit energy, per unit time), through the conservation of the transferred energy, as
\[ E'_\nu \frac{dN_\nu}{dE'_\nu} = (E'_p \varphi(E'_p) S[(1 + z) \times E'_p]) \frac{dE'_p}{dE'_\nu} \quad (4.9.4) \]
\[ \frac{dN_\nu}{dE'_\nu} \frac{dE'_p}{dE_\nu} = 20 \varphi(E'_p) S[(1 + z) \times E'_p] \frac{dE'_p}{dE'_\nu} \quad (4.9.5) \]

with \( S \) as defined in Section 4.9. Substitution of Eq. (4.9.5) into Eq. (4.9.3) yields Eq. (4.9.2).

Given an UHECR source evolution, \( W(z) \), we can calculate the expected flux of cosmogenic neutrinos produced through photopion production on the CMB. While our simple, analytic method allows for a more transparent look at the effects of source evolution, it does not fully encompass the particle physics involved, particularly the low and high energy tails of the distributions of particle decays (which affect the low- and high-energy ends of the neutrino spectrum). We utilize the publicly-available simulation package CRPropa [289], which uses the SOPHIA [290] code to handle particle processes, for this purpose. We have made use of the analytical estimate described above and an extensive comparison to previous results presented in the literature (with similar parameters), to verify the results.

At the redshifts of greatest interest, \( z \sim 0 - 4 \), only cosmic rays with \( E \gtrsim 5 \times 10^{19}/(1 + z) \) eV can ever have the ability to produce neutrinos through the GZK
process (even with a decreased photopion threshold). At lower energies, photopion production can be facilitated by the cosmic infrared background (IRB), resulting in additional neutrinos of correspondingly lower energy [291]. Additionally, extragalactic magnetic fields may increase the path length of cosmic-ray propagation [242]. However, due to the various uncertainties involved, we have chosen not include these effects.

Fig. 4.10 compares the resulting neutrino flux from the GRB model to those expected from QSO-like evolution and the SFH, assuming an $E^{-2}$ injection spectrum (with a sharp cutoff at $10^{21}$ eV). The flux shown is the sum of all neutrino flavors (both electron and muon types are produced at the source), as a detector such as ANITA has a nearly equal taste for each flavor. As can be clearly seen, the stronger evolution of the GRB model produces a larger neutrino flux than other models. Also shown are the expected sensitivities of ANITA [251] (which began taking data in late 2006) and the proposed ARIANNA/SalSA detectors [252]. The flux predicted by strong GRB evolution, with parameters given by our metallicity-biased model, extends well into the reach of ANITA. This presents an opportunity to either confirm or significantly constrain the parameters of this model in a manner that would not be possible by cosmic-ray observations alone.

We can also consider the effect of varying the assumed parameters in the injection spectrum. Since only cosmic rays with energies greater than $10^{19}$ eV will ever contribute significantly (even at high redshifts) to the neutrino flux, harder cosmic-ray spectra will result in larger fluxes. Shown in Fig. 4.11 are the neutrino fluxes arising from cosmic-ray injection spectra of $\gamma = 1.5$, 2.0, and 2.5, assuming the strong evolution. In particular, a $\gamma = 2.2$ spectrum might arise from relativistic shocks. Since $\gamma \lesssim 2$ would be quite difficult to reconcile with cosmic-ray observations, this range
Figure 4.11: Expected (all-flavor) cosmogenic neutrino fluxes resulting from various UHECR injection spectra and assuming strong (metallicity-dependent) GRB evolution. From top-to-bottom, are the fluxes with spectral indices 1.5, 2, and 2.5. Shown for comparison are the Waxman-Bahcall bound (shaded band) and the expected sensitivities for ANITA (diamonds) and ARIANNA/SalSA (triangles) [86].

can be regarded as an upper bound, which will soon be constrained by ANITA. Assuming a lower cutoff energy would be similar in effect to a softening of the spectral
index. Another interesting scenario is producing a softer spectral index \((\gamma > 2)\), even if each GRB possesses an intrinsic \(E^{-2}\) spectrum. As proposed in Ref. [292], if the distribution of the high-energy cutoffs of cosmic-ray sources follows a power law, then the overall spectrum that is observed will follow a power law with a different slope, which may be particularly applicable to GRBs.

4.10 Discussion and Conclusions

The increase in our understanding of gamma-ray bursts can be traced to the improved capabilities now available to study these phenomena. The ability to quickly and accurately localize a GRB has led to the establishment of a GRB-supernova connection and allowed for the study of the host galaxies in which these events occur. Observations indicate that these hosts tend to be underluminous, star-forming and metal-poor. A connection between a GRB and its host galaxy metallicity is not surprising in the context of the collapsar model, which requires rapidly-rotating stars that lack a H/He envelope (in order to be in accord with supernova observations). These requirements can be satisfied by a metal-poor progenitor star.

The set of Swift gamma-ray bursts now allows for model-independent tests of the connection between the GRB and star formation rates. We have presented quantitative evidence that the GRB rate does not simply track star formation over a broad range in redshift; some mechanism, of a presently-unknown nature, is leading to an enhancement in the observed rate of high-\(z\) gamma-ray bursts. The effects of stellar metallicity appear to be a compelling explanation, but cannot be conclusively proven as of yet.

The possibility of continuing evolution in the GRB rate relative to the SFR invites some astrophysical speculations. While the overall black hole production rate in core-collapse supernovae is poorly understood empirically [26], bright GRBs might be
regarded as tracing the large angular momentum end of the black hole birth distribution (e.g., for collapsars). Evolution would mean that the high-z universe was more efficient at producing such black holes than usually considered. This may have implications for the nucleosynthetic yields from these explosions. The effect of supernovae on the gas in high-z galaxies is an ongoing field of research. Following from the considerations above, it may be that small galaxies had a disproportionately large rate of GRBs relative to normal supernovae. In this case, it would be interesting to examine the fate of such galaxies after including multiple injections of highly-asymmetric, relativistic ejecta, as well as implications for enrichment of the intergalactic medium (i.e., “GRB feedback”).

We have also developed an empirical method for estimating the high-z SFR using GRB counts, improving on earlier estimates of the high-z SFR using GRB data in several ways, not least of which by using significantly updated SFR [70] and/or GRB [235] data. Taking advantage of the improved knowledge of the SFR at intermediate z, we were able to move beyond the assumption of a simple one-to-one correspondence between the GRB rate and the SFR, accounting for an increasing evolutionary trend. The higher statistics of the recent Swift GRB data allowed the use of luminosity cuts to fairly compare GRBs in the full z range, eliminating the uncertainty of the unknown GRB luminosity function. By comparing the counts of GRBs at different z ranges, normalized to SFR data at intermediate z, we based our results squarely on data, eliminating the need for knowledge of the absolute fraction of stars that produce GRBs.

With the discovery of the first astrophysical source at z > 8, Swift has enabled GRBs to realize their potential as beacons from the distant past, both into the epoch of reionization and in adequate numbers at lower redshifts to allow for sensible use of the most remarkable events. Using this wealth of data, we have estimated the star
formation rate at the earliest times yet possible, showing that the star formation rate can remain high up to at least $z \sim 8$. From this, it is plausible that the level of star formation was sufficient to reionize the Universe, with very-recent observations from the upgraded *HST* Wide Field Camera 3 [90, 91] apparently confirming this conclusion.

The agreement with direct observations, corrected for galaxies below detection thresholds, suggests that our GRB-based estimates incorporate the bulk of high-$z$ star formation down to the faint end of the LF. We also see no evidence for a strong peak in the SFR versus $z$. This assumes that a very strong rise in the efficiency of producing GRBs (beyond that already accounted for), does not hide a drop in the SFR, although this itself would be quite interesting. While we have not included them in our analysis, of the three *Fermi* GeV-detected long GRBs with redshifts (shown in Fig. 4.1), two were at $z > 3.5$ (e.g., [293]). Their brightness raises the prospect of the independent use of GeV-selected bursts.

The current picture of small, metal-poor GRB hosts observed at low $z$ agrees well with our GRB-inferred SFR being dominated by such sub-$L_*$ galaxies at high $z$. One might wonder about the whereabouts of these GRB hosts today, whether they continued to grow, merged into more massive halos, etc. Observations of the afterglow spectrum (e.g., [294]) could determine the extent that the host had experienced the effects of the reionizing UV background. Since GRBs should originate from a different range of overdensities than quasars, potential exists for another examination of the hierarchical history of our Universe.

In models that attribute ultrahigh energy cosmic rays to gamma-ray bursts, this evolution of the cosmic GRB rate density traces the history of UHECR production in the universe. If GRBs are to account for the observed cosmic-ray spectrum, they *must* generate a flux of cosmogenic neutrinos. These neutrinos present a unique tool to
examine the GRB-as-proton accelerator conjecture. While the cosmic-ray spectrum that will be measured by Auger will allow for further assessments of the viability of prospective source models, the combined measurements of neutrinos and UHECR would break degeneracies between the various models. The enhanced cosmogenic neutrino fluxes expected to result from strong GRB evolution will allow for this model to be tested in a novel fashion.

The sensitivity afforded by ANITA will allow for near-term examination which may either affirm or, if the expected flux is not found, place substantial constraints upon the model parameters. Indeed, recent results from ANITA may indicate a flux near our expected level [295]. The lower detection threshold achieved by the next generation of detectors will allow for a realistic opportunity to discriminate between evolution models. Measurements of fluxes consistent with that expected from GRB evolution would provide compelling evidence.
CHAPTER 5
THE SOURCE OF HIGH-ENERGY POSITRONS AND ELECTRONS

The greater the difficulty, the more glory in surmounting it. Skillful pilots gain their reputation from storms and tempests.

Epicurus

5.1 Introduction

Geminga holds a place of distinction among gamma-ray sources, being the first pulsar to be discovered through gamma rays, with a history of observations through a variety of techniques [296]. While one of the brightest MeV–GeV gamma-ray point sources in the sky, there was no certain evidence of high-energy activity beyond the immediate neighborhood of the pulsar or its x-ray pulsar wind nebula (PWN) until the recent detection by Milagro of gamma rays at $\sim 20$ TeV from a region of $\sim 3^\circ$ around the pulsar [53, 297]. This detection places Geminga among the growing class of TeV PWNe (e.g., [298, 129]) and is important for understanding aged pulsars and their winds. An immediate consequence is the existence of a population of high-energy particles.

\footnote{The content of this Chapter is based in large part on our work in Refs. [54, 55].}
The relative proximity of Geminga raises an interesting possibility, namely that these high-energy particles, most likely electrons and positrons, may be at the root of the explanation of the “positron excess”, the observed $[299, 300, 51]$ overabundance of multi-GeV positrons as compared to theoretical expectations $[52]$ (see Fig. 5.3). Severe energy losses of high-energy positrons require a local source of some kind $[303]$, such as Geminga $[304]$ or even dark matter through its annihilation products $[117]$.

Indeed, interest in the cosmic-ray electron spectrum at Earth is at an all-time high, arising from both astrophysical and dark matter-related concerns. In this GeV energy range, much greater clarity than what had existed in preceding decades has been brought by the PAMELA $[51]$ and Fermi $[48]$ space missions. Such measurements, which observe electrons and positrons directly, become difficult above $\sim 1$ TeV due to their fixed detector areas relative to declining particle fluxes. Using the indirect technique of observing atmospheric air showers, HESS has pushed the energy frontier up to several TeV $[49]$, although above this, no electron measurements have been reported. With the historical lack of local multi-TeV gamma-ray sources and the soft spectra of secondary $e^\pm$ resulting from $p$–$p$ scattering $[52]$, little if any signal may have been expected.

Here, we connect the Milagro TeV gamma-ray “halo” to electrons and positrons with energies up to at least 100 TeV, expected to be accelerated in PWNe (e.g., $[305, 306, 307]$; for a review see $[308, 309]$), and present several predictions. Principally, while Geminga is apparently young enough to still produce high-energy particles, it is old enough that multi-GeV electrons and positrons from its more active past can have made it to Earth. The extended gamma-ray emission is strong evidence for $e^\pm$ production, acceleration, and escape, suggesting an explanation of the positron excess. Moreover, this single nearest high-energy astrophysical source can reasonably account for the $e^- + e^+$ spectrum as measured by Fermi $[48]$ and HESS $[49, 50]$ with
an extension to energies beyond several TeV, where no signal might be expected otherwise.

Due to the similarity of the electromagnetic showers produced in the atmosphere by energetic electrons and gamma rays, they are nearly inseparable via ground-based observations [324, 325]. Here, we proceed to translate published limits on isotropic gamma-ray fluxes (e.g., [325, 326, 328, 327]) to constrain the $e^- + e^+$ spectrum to $>\text{PeV}$ energies. We further show that our derived limits are relevant in light of recent measurements by Fermi and TeV gamma-ray telescopes, particularly the discovery and detailed observations of high-energy sources both near and far that indicate $e^\pm$ production in the GeV and TeV regimes. Considering these observations, with an improved analytical treatment of $e^\pm$ propagation to calculate expected source contributions at Earth, we conclude that it would be surprising if the influence of a nearby pulsar was not present in cosmic-ray positron data.

5.2 The Gamma-ray Source Next Door

The observation of high-energy gamma rays from an astrophysical source implies the presence of higher-energy particles, typically $e^\pm$ or protons, that gave rise to them. One striking element of the observation of $\sim 20$ TeV gamma rays (with a significance of $4.9\sigma$ in the PSF-smoothed map [53], $6.3\sigma$ for an extended source [297]) from Geminga by Milagro is the extent of the emission, $\theta \sim 3^\circ$ [53], which corresponds to a physical size of $s_G \sim 10\text{ pc}(\theta_G/3^\circ)(r_G/200\text{ pc})$, where $r_G$ is the distance to Geminga. Since the angular resolution of Milagro is better than a degree and the characteristic age of the pulsar, $t_G \sim 3 \times 10^5\text{ yr}$ [310], seemingly excludes a typical TeV supernova remnant, we shall consider an extended PWN with emission from a much larger region than seen in x-rays [311, 312]. We will draw guidance from the TeV-PWN
Figure 5.1: The cosmic-ray electron spectrum at Earth. Shown are direct $e^- + e^+$ measurements from Fermi [48], measurements of $e^- + e^+$ based on showers from HESS [49, 50], and a baseline $\propto E^{-3}$. These can be compared to our limits derived from gamma-ray experiments [325, 326, 328, 327] at $> 10$ TeV [55].

HESS J1825−137 [298], which, while only a tenth the age of Geminga, would appear tens of degrees wide if placed at $r_G \sim 200$ pc.

We first examine whether the gamma rays can be explained through inverse-Compton (IC) up-scattering of cosmic microwave background (CMB) photons by $e^\pm$. Note that the pulsar age exceeds the IC cooling time on CMB photons of the $\gtrsim 100$ TeV $e^\pm$ needed to produce $\gtrsim 20$ TeV gamma rays, $\tau_{IC} \sim 10^4(100 \text{ TeV}/E_e)$ yr in
the Thomson limit. Including synchrotron losses further decreases $\tau_{\text{cool}}$, implying fresh $e^\pm$ production. To account for the Milagro measurement of $6.9 \pm 1.6 \times 10^{-15}$ TeV$^{-1}$ cm$^{-2}$ s$^{-1}$ at 20 TeV (see Fig. 5.2), we consider a generic parent $e^\pm$ spectrum of the form $dN/d\gamma \propto \gamma^{-\alpha} e^{-\gamma/\gamma_{\text{max}}}$, with $\gamma = E/(m_e c^2)$. Lacking more detailed observations, we choose $E_{\text{min}} = 1$ GeV, $E_{\text{max}} = 200$ TeV and $\alpha = 2$ (typical to shock acceleration and as inferred in the Vela X PWN [129]). The resulting IC spectrum is

$$\frac{d\Phi}{dE_{\gamma}} = \frac{c}{4\pi r_G^2} \int d\gamma \int dE_{b} \frac{dN}{d\gamma} n_b(E_{b}) \sigma_{\text{KN}}(\gamma, E_{b}, E_{\gamma}),$$

(5.2.1)

where the Klein-Nishina cross section, $\sigma_{\text{KN}}$, is given by Ref. [313]. The dominant scattering background, $n_b(E_{b})$, in the Milagro energy range is the CMB, allowing us to construct the minimal spectrum shown in Fig. 5.2 (at lower energies, other contributions become relevant). When we normalize the IC spectrum to the Milagro TeV gamma-ray luminosity of $\dot{E}_{\gamma,\text{TeV}} \sim 10^{32}$ erg s$^{-1}$ (for $r_G \sim 200$ pc), at least $E_{e^\pm} \sim 10^{45}$ erg of $e^\pm$ is required.

We can relate this phenomenological spectrum to the pulsar. The Goldreich-Julian flux [305] is $\dot{N}_{\text{GJ}} \simeq B \Omega^2 R^3/ ec$. For Geminga, $R \simeq 10$ km, $B \simeq 1.6 \times 10^{12}$ G, and $\Omega \simeq 26.5$ s$^{-1}$ [310], which yield $\dot{N}_{\text{GJ}} \simeq 10^{32}$ s$^{-1}$. For our $E_{e^\pm}^{-2}$ spectrum, $\langle E_{e^\pm} \rangle \approx 15$ GeV. The total power is then $\dot{E}_{e^\pm} \approx \dot{N}_{e^\pm} \langle E_{e^\pm} \rangle \approx \mathcal{M} \dot{N}_{\text{GJ}} \langle E_{e^\pm} \rangle \approx 2 \times 10^{30} \mathcal{M}$ erg s$^{-1}$, where $\mathcal{M} = \dot{N}_{e^\pm}/\dot{N}_{\text{GJ}}$ is the pair multiplicity. For a pure electron flow ($\mathcal{M} \simeq 1$), the power beyond $10$ TeV is only $\sim 5 \times 10^{29}$ erg s$^{-1} \ll \dot{E}_{\gamma,\text{TeV}}$, thus requiring pair production resulting in $\mathcal{M} \gtrsim 100$. Including synchrotron losses comparable to IC increases the required $\mathcal{M}$. This can only be avoided by assuming a spectrum much-harder than $E_{e^\pm}^{-2}$ or $E_{\text{min}} > 100$ GeV (both inconsistent with, e.g., Vela X [129, 314]).

While we have been trying to explain the observed signal alone, there may well be more lower surface brightness emission at larger angles, and also at lower energies. Two considerations make this likely. First, the pulsar’s $e^\pm$ output was probably much stronger in the past when its spin-down power was higher. Secondly, no evidence of
Figure 5.2: Minimal inverse-Compton gamma-ray spectrum of the extended emission from Geminga (shaded) and the Milagro measurement at 20 TeV (left axis). Also, the energy distribution (dotted line) of the associated $e^\pm$ (right axis) [54].

a large-scale radio or x-ray nebula exists, and a substantial fraction of the $e^\pm$ may be escaping, so that the above multiplicity is only a lower limit. $\mathcal{M} \sim 10^4$ (as inferred for younger TeV PWNe [314] or pulsar models [315]) does not exceed the spin-down power of $\sim 10^{34.5}$ erg s$^{-1}$ [310], so that a large pair conversion fraction is possible.

5.3 The Origin of the Positron Excess

The confirmed presence of a nearby, ancient source of high-energy electrons and positrons immediately suggests an explanation for the positron excess. If so, then we would essentially be living “within” the extended halo of the source, seeing $e^\pm$ that were accelerated long ago when the pulsar was stronger.
In spherically symmetric geometry, the diffusion equation governing the particle density at a given location/time/energy, \(n(r, t, E)\), is \[333, 317\]

\[
\frac{\partial n}{\partial t} = \frac{D(E)}{r^2} \frac{\partial}{\partial r} r^2 \frac{\partial n}{\partial r} + \frac{\partial}{\partial E} [b(E) n] + Q ,
\]

(5.3.1)

with energy losses parametrized as \(b(E) = -dE/dt\), diffusion coefficient \(D(E)\), and source term \(Q\). We first consider a single burst from a point source with a generated particle spectrum \(dN/dE \_g\), so that \(Q(r, t, E \_g) = \delta(r) \delta(t) dN/dE \_g\). Using the Syrovatskii propagator \[333\] (as in \[304, 54\]) yields

\[
n_S(r, t, E) = \frac{e^{-r^2/r \_dif^2 g}}{\pi^{3/2} r \_dif^3} dE \_g dE ,
\]

(5.3.2)

where \(dE \_g/dE\) maps the energy at generation \(E \_g(E, t)\) to the observed \(E\) after losses.

In the high-energy regime of interest here, most relevant are inverse-Compton losses on the CMB (energy density \(\sim 0.3\) eV cm\(^{-3}\)) and synchrotron losses due to a \(\sim 3\) \(\mu\)G magnetic field (\(\sim 0.2\) eV cm\(^{-3}\)), so that \(b(E) = b_0 E^2\) with \(b_0 \simeq 5 \times 10^{-16}\) s\(^{-1}\) GeV\(^{-1}\). Integration yields \(1/E = 1/E \_g + b_0 t\), so that \(dE \_g/dE = (E \_g/E)^2\).

The diffusion radius is \(r \_dif(E, t) = 2\sqrt{\lambda(E, t)}\), with

\[
\lambda(E, t) = \int_0^t dt' \int_0^{E \_g} dE' D(E') = \int_0^{E \_g} dE' \frac{D(E')}{b(E')} .
\]

(5.3.3)

We parametrize the diffusion coefficient as \(D(E) = d_0 (1 + E/E_0)^\delta\), with \(E_0 \simeq 3\) GeV, and consider values of \(d_0 = 2 \times 10^{28}, 4 \times 10^{27}\) cm\(^2\)s\(^{-1}\), \(\delta = 0.4, 0.5\) \[317, 334\] corresponding to “fast” and “slow” models. Since \(D(E) \simeq d_0 (E/E_0)^\delta\) at energies \(E > E_0\), the diffusion radius simplifies to

\[
r \_dif(E, E \_g) \simeq 2 \left( \frac{d_0 E_0^{\delta - 1} - E \_g^{\delta - 1}}{b_2 (\delta - 1) E \_0^\delta} \right)^{1/2} .
\]

(5.3.4)

For \(t_G \sim 3 \times 10^5\) yr, \(D_0 \simeq 4 \times 10^{27}\) cm\(^2\)s\(^{-1}\), and \(\delta = 0.4\), the diffusion radius is \(r_d \simeq 150, 175, 250\) pc for \(E = 2, 10, 50\) GeV particles, respectively, so multi-GeV particles are now arriving. We caution against extrapolating to small radius, since the diffusion solution may yield spurious results, as we explain later.
For a continuously emitting source such as Geminga, the injection rate can be parametrized as \( \frac{d\dot{N}}{d\gamma} \propto L_{\epsilon^\pm}(t)\gamma^{-\alpha} e^{-\gamma/\gamma_{\text{max}}} \), with \( L_{\epsilon^\pm} \) the \( \epsilon^\pm \) luminosity. The local particle density is \( n_\odot(\gamma) = \int_{r_G}^{t_G} dt \bar{n}(r_G, t, \gamma) \). Assuming braking via magnetic dipole radiation, the spin-down luminosity evolves as \( \propto (1+t/t_0)^{- (n+1)/(n-1)} \), with \( n = 3 \) and a pulsar-dependent timescale, \( t_0 \), and \( L_{\epsilon^\pm}(t) = (E_G/t_G) [1 + (t_G - t)/t_0]^{-2} / \int_{t_G}^{t} dt' [1 + (t_G - t')/t_0]^{-2} \). For \( t_0 \sim 3 \times 10^4 \) yr, the present spin-down power, \( \sim 10^{34.5} \) erg s\(^{-1}\), corresponds to an upper limit on the total \( \epsilon^\pm \) output of \( \sim 5 \times 10^{48} \) erg (larger for smaller \( t_0 \) [316]). Geminga’s transverse velocity is \( \sim 200 \text{ km s}^{-1} \) [321]. A similar radial velocity would result in a \( \sim 100 \text{ pc} \) displacement in \( t_G \).

In Fig. 1.4, we display the local flux of \( \epsilon^- + \epsilon^+ \), \( J_\odot = (c/4\pi) n_\odot \), from our benchmark model of \( \alpha = 2 \), within a reasonable range of parameters. These have distances varying (from birth \( \rightarrow \) present) as \( r_G = 150 \rightarrow 250 \text{ pc} \), \( 220 \text{ pc} \), \( 250 \rightarrow 200 \text{ pc} \); \( \epsilon^\pm \) energy budgets of \( E_G = 1, 2, 3 \times 10^{48} \) erg; and \( \delta = 0.4, 0.5, 0.6 \), respectively (lower dotted, solid, dashed lines). The energy in \( \epsilon^\pm \) estimated for several younger TeV PWN are at least as large as these (e.g., [314, 298, 129]). Since the bulk of the energy is released in this early spin-down phase, the initial location is the most important. Adding to these the primary \( \epsilon^- \) spectrum of Moskalenko and Strong [52], with the normalization decreased by 35% and an added exponential cutoff at 2 TeV (in order to not exceed HESS data), yields the total \( \epsilon^- + \epsilon^+ \) flux (upper lines).

The spectral feature at \( \sim 1 \text{ TeV} \) naturally results from a combination of energy losses and pulsar age and distance (see Fig. 4 of Ref. [304] for comparison). The multi-TeV extension (beyond the last HESS point) is due to the continuous injection of particles, as evidenced by the Milagro observations today. Combining these with the expectations for the secondary \( \epsilon^\pm \) fluxes [52], we compare our positron fraction to measurements in Fig. 5.3 (note that solar modulation may account for disagreements between data below \( \sim 10 \text{ GeV} \) [51, 302]).
Figure 5.3: The cosmic-ray positron fraction. Shown are a compilation of data, scenarios based a secondary model (shaded), and a plausible Geminga contribution (solid, dashed, and dotted lines) dependent upon distance and energetics from Ref. [54].

It is thus plausible that Geminga is the long-sought [322] local source of electrons and positrons, influencing the spectra measured by Fermi [48] (down to tens of GeV) and HESS [49, 50] in the TeV, although we emphasize that certain parameters and the underlying Galactic primary spectrum remain uncertain. PAMELA and AMS [301] can measure the $e^{-}$ and $e^{+}$ spectra separately, which can isolate this component (since the $e^{-}$ spectrum from Geminga should be identical to the $e^{+}$).
5.4 New Cosmic-Ray $e^\pm$ Limits

While the $\sim 1 \text{ m}^2$ Large Area Telescope (LAT) of Fermi [48] has brought a sharper picture of the $e^- + e^+$ spectrum up to 1 TeV, to progress further requires a much larger effective area for particle collection. This can be accomplished by, in lieu of direct particle identification, examining the showers that result when energetic particles scatter in the upper atmosphere.

To observe an electromagnetic shower (initiated by a gamma ray or electron) requires rejecting the large background due to cosmic-ray protons. Ground-based air Čerenkov telescopes (ACTs) operate by imaging the shower and observing the differences between hadronic and electromagnetic cascade development [329], which allowed HESS to measure $e^- + e^+$ up to several TeV (Fig. 5.1). Alternatively, detectors that operate by directly detecting the long-lived products of air showers as they reach the ground make use of the low muon content of electromagnetic showers relative to hadronic events [329].

Searches for isotropic fluxes of gamma rays have been conducted over a wide range of energies by examining muon-poor showers [326, 328, 327]. While these have typically only resulted in upper limits, they can be valuable. Since a gamma ray first produces an initial $e^\pm$ pair to begin a cascade, its shower will look very similar to that of an electron of equivalent energy, only located $\sim$ one radiation length deeper in the atmosphere [324]. HESS electron data can thus be regarded as upper limits on an isotropic TeV gamma-ray flux [49, 118], while air shower arrays do not have hope of exploiting this subtle difference, rendering diffuse gamma rays and electrons inseparable [325].

By using isotropic gamma-ray limits from air shower arrays, we derive new limits on the cosmic-ray electron spectrum in a regime currently lacking constraints.
These limits are often quoted as the fraction of measured gamma-ray to proton intensity, \( I_\gamma/I_{CR} \). We use the cosmic-ray nuclei spectrum [330] in Fig. 5.1. In the range \( 10^{-1} \text{ to } 10^5 \text{ TeV} \), we obtain electron limits from HEGRA [325], CASA-MIA [326], GRAPES [327], and KASCADE [328] data, leading to the limits in Fig. 5.1.

We have conservatively assumed that the electron fraction of the electromagnetic showers seen is 100%. It is likely that dedicated re-analyses of the data can strengthen these constraints, and an analysis of Milagro [297] data could probe the 10–100 TeV region. Even so, we see that our limits are already competitive when compared to previous measurements at lower energies.

5.5 Nearby \( e^\pm \) Factories

Due to their short lifetime against radiative losses, any electron flux measured at \( E > 10 \text{ TeV} \) must be produced in the very-recent history of a nearby source. For this reason, we will consider sources within \( \sim 500 \text{ pc} \) that still exhibit evidence of particle acceleration to very-high energies, using more distant sources for added guidance. Observations, principally of extended gamma-ray emission, have recently revealed good reasons to believe that \( e^\pm \) are being produced, accelerated up to multi-TeV energies, and escaping from the high-energy-density environments of pulsars:

5.5.1 Clues from HESS J1825–137

HESS has intensively observed the distant (~4 kpc) pulsar wind nebula (PWN) HESS J1825–137 [298]. This included untypically-long exposure times due to the PWN being serendipitously within the field-of-view of the variable microquasar LS 5039, for which repeated measurements of the light curve were taken. These revealed an extended wind of \( \gtrsim 10 \text{ TeV} e^\pm \) containing \( \gtrsim 10^{48} \text{ erg} \) reaching \( \gtrsim 100 \text{ pc} \) in only \( \sim 20,000 \text{ yr} \) [298]. This indication of a pulsar “TeV mode” carrying \( \gtrsim 10^{48} \text{ erg} \) of \( e^\pm \)
to great distances likely would have been missed without these deep observations, and thus may be common.

### 5.5.2 Vela X

The proximity (290 pc) and relative youth (\(\sim 11,000\) yr) of the Vela pulsar has permitted detailed observations of its PWN, Vela X. There, HESS inferred a population of \(e^\pm\) with energies reaching \(\sim 100\) TeV [129]. Modeling source \(e^\pm\) injection spectra as \(dN/dE \propto E^{-\alpha}e^{-E/E_{\text{max}}}\), the HESS measurements imply \(\alpha = 2\) and \(E_{\text{max}} = 70\) TeV for Vela X [129]. Intriguing evidence has also been obtained at lower energies, with a distinct population of multi-GeV, radio-synchrotron emitting \(e^\pm\) (with an \(E^{-1.8}\) spectrum) modeled in Ref. [331]. The \textit{Fermi} discovery of extended GeV gamma rays in Vela X confirms this and, importantly, indicates a high-energy cutoff in the spectrum at \(\sim 100\) GeV [332]. The amplitude of this signal implies that a pulsar “GeV mode” can also produce \(\gtrsim 10^{48}\) erg of \(e^\pm\). This suggests that the total composition of Vela X includes \(\sim 100\) GeV \(e^\pm\) accelerated by the pulsar itself — likely associated with pulsed GeV gamma rays — and TeV \(e^\pm\) from shock acceleration of the pulsar wind. As observed, the GeV component contains \(\sim 100\) times more energy, although multi-TeV particles may have already exited the system.

### 5.5.3 Geminga

As discussed above, the discovery of extended \(\sim 35\) TeV gamma-ray emission by Milagro [297] surrounding the nearby (\(\sim 200\) pc) pulsar Geminga indicates a close, active source of \(e^\pm\) [54]. These data imply a firm lower limit on the maximum particle energy of \(\gtrsim 100\) TeV, and may approach 1000 TeV. In all of these systems, the implied particle multiplicities needed to account for the gamma rays via inverse-Compton scattering require \(e^\pm\) pair production. Since, like Vela, Geminga possesses bright
pulsed GeV emission, it may as well have resulted in an abundance of $\sim 100$ GeV $e^\pm$.

While Vela is too young for GeV $e^\pm$ to have reached us yet (as we will soon see), the greater age of Geminga ($\sim 300,000$ yr) may allow for a direct test of the commonality of dual high/low-energy pulsar $e^\pm$ populations. To determine this, we first address the propagation of multi-TeV $e^\pm$ in the Galaxy.

## 5.6 TeV $e^\pm$ Propagation

We first use $r_{\text{dif}}(E,E_g)$ from Eq. (5.3.2) to gain insight into which sources could contribute to the multi-TeV electron spectrum. If a particle generated with even $E_g = 100$ TeV is to be detected at $E = 10$ TeV, it may propagate a maximal distance of $r_{\text{dif}} \sim 160$ (230) pc in slow (fast) diffusion models. Note that the term $e^{-(r^2/r_{\text{dif}}^2)}/(r_{\text{dif}}^3$ in Eq. (5.3.2) strongly favors particularly nearby sources.

Closer inspection of the Syrovatskii solution reveals that more care is needed at these high energies. For $E_g \sim E$, $E \simeq b_0 t$ and $r_{\text{dif}}(t) \simeq 2\sqrt{D(E)t}$, the quantity $v_{\text{dif}}(t) = r_{\text{dif}}(t)/t \simeq 2\sqrt{D(E)/t}$ can exceed the speed of light. This is the known problem of “superluminal” diffusion, which lies in the fact that the diffusion solution is not relativistic, similar to the Maxwell distribution yielding values in excess of $c$ [335, 336, 337]. A phenomenological resolution was discussed in the context of ultrahigh-energy cosmic rays in Ref. [338], which uses a propagator based on the Jüttner particle distribution [335] to explicitly limit fluxes to $v < c$, while preserving diffusive behavior at lower energies. Now,

$$n_J(r,t,E) = \frac{\theta(1-\xi) e^{-\alpha/\sqrt{1-\xi^2}}}{4\pi (ct)^{3}} \frac{\alpha}{(1-\xi^2)^2} K_1(\alpha) \frac{dN}{dE_g} dE$$  \hspace{1cm} (5.6.1)

where $\xi(r,t) = r/ct$, $\theta$ is the step function, $K_1$ is the modified Bessel function, and $\alpha(E,t) = c^2 t^2/(2\lambda(E,t))$ [338].

This formulation has the added benefit of removing spurious features appearing
<table>
<thead>
<tr>
<th>Model</th>
<th>TeV mode (erg)</th>
<th>GeV mode (erg)</th>
<th>Distance (pc)</th>
<th>Age (yr)</th>
<th>Diffusion</th>
</tr>
</thead>
<tbody>
<tr>
<td>TeV</td>
<td>$2 \times 10^{48}$</td>
<td>0</td>
<td>220</td>
<td>$3 \times 10^5$</td>
<td>slow</td>
</tr>
<tr>
<td>Dual</td>
<td>$1.3 \times 10^{48}$</td>
<td>$1.3 \times 10^{48}$</td>
<td>300--200</td>
<td>$2 \times 10^5$</td>
<td>fast</td>
</tr>
</tbody>
</table>

Table 5.1: Parameters used for the two Geminga scenarios.

in the non-relativistic solution for source spectra harder than $E^{-2}$ and can easily be generalized to a continuously-emitting source at distance $r$ and time $t$ with a time-dependent particle injection rate $d\dot{N}/dE_g$, as $n_\odot(E) = \int_0^t dt' \dot{n}(r,t',E)$. Again, for sources such as pulsars, $r$ can vary with $t$.

5.7 Beyond the Positron Excess

Using the Jüttner formalism, we can better estimate the extent to which a source may influence the spectrum at Earth. We again assume magnetic dipole braking and that the $e^\pm$ output is simply proportional, normalizing to the $e^\pm$ luminosity of the source, $\mathcal{L}_{e^\pm}(t) = \int_{E_{\text{min}}}^{E_{\text{max}}} E d\dot{N}/dE dE$ (with $E_{\text{min}} = 1$ GeV). For Geminga, $\dot{E} \sim 10^{34.5}$ erg s$^{-1}$ at present, and we here use $t_0 \sim 3 \times 10^4$ yr (corresponding to a spin period at birth of $\sim 70$ ms). We consider two distinct scenarios: a case from above [54] in which the TeV mode dominates and another with comparable energy in the dual GeV and TeV modes. For the TeV spectra, we assume $E_{\text{max}} = 200$ TeV, and $\alpha = 2$. For the GeV spectrum, we use parameters consistent with Vela X: $E_{\text{max}} = 150$ GeV, and $\alpha = 1.8$.

In Fig. 5.4, we display the local flux of $e^- + e^+$, $J_\odot = (c/4\pi) n_\odot$, for these models, as described in Table 5.7. These require $\sim 40\%$ of the spin-down power be converted to high-energy $e^\pm$ pairs, within the range inferred from Vela X [331]. We also examine
the expected ratio of $e^+/(e^+ + e^-)$. For the denominator, we directly use a fit to the measured data (neglecting the final HESS datum), as shown in Fig. 5.4. In the numerator, we include secondary fluxes [52] to match low-energy data. As we see in the top panel, the two models can easily diverge at $\sim 100$ GeV. A more GeV-dominated case would drop more dramatically.

A few remarks are in order: (1) The underlying Galactic $e^-$ component is not known, but must be cut off at some point in order to not overshoot high-energy data (as seen in Fig. 5.4). This must make up the difference between any model and the full measured $e^- + e^+$ spectrum. (2) Multiple nearby sources may contribute, although Geminga and Vela are the only bright gamma-ray pulsars within $\sim 300$ pc. (3) Low energies are more sensitive to the source’s initial position, while at high energies the present position is most relevant (which can be directly measured). (4) The emission properties (spectral index, cutoff, pair conversion efficiency, etc.) may evolve in time, while the highest-energy emission can reasonably be tied to what is seen today. The observations of Vela do give hope of examining lost history, though.

5.8 Discussion and Conclusions

The discovery of high-energy gamma rays from an extended region around Geminga by Milagro reveals the presence of $\gtrsim 100$ TeV $e^\pm$, as observed indirectly within the x-ray PWN [311, 312]. A considerable amount of data should become available as new experiments examine the surrounding area. This will help in developing more detailed models that account for both time and spatial evolution in the $e^\pm$ spectra, directly coupled to cosmic-ray propagation. One need is a better-determined distance, the most recent quoted being $r_G \sim 250^{+120}_{-62}$ pc [321]. We briefly discuss implications for several categories of experiments.

$\textit{Fermi}$: While the observed features of Geminga will depend upon details such as
Figure 5.4: Our modeled contributions to the $e^+/e^-$ spectra. **Bottom:** Shown are $e^- + e^+$ measurements from Fermi [48] and HESS [49, 50]; direct $e^-$ from AMS [301]; our fit to these; and our $> 10$ TeV limits from Fig. 5.1. For Geminga, we show a model in which the TeV mode dominates and one with comparable GeV/TeV modes. **Top:** A comparison of the expected cosmic-ray positron fraction from the two models with data up to $\sim 100$ GeV [300, 51, 301] [55].

whether the source is roughly spherical or preferentially oriented, we would generally expect the source to become “larger” with decreasing energy, reflecting the decrease in IC cooling time with energy. Our inspection of the point-source subtracted EGRET sky-map [162] indicates emission in the GeV range of a size comparable to the Milagro
source. Fermi [34] should be able to more effectively separate the bright pulsed signal to study diffuse emission.

**TeV gamma rays:** Obtaining a detailed spectrum and morphology of the source in the TeV regime will be vital for further interpretation of the nature of the particles present. Already, VERITAS [32] has placed rather-tight upper limits on a point source [323]. Further study of the expected extended source is needed to better estimate the total energetics. Returning to HESS J1825–137 as a guide, the surface brightness was seen to drop off as $\sim 1/\theta$, inconsistent with pure diffusion and suggestive of convection, and the gamma-ray spectrum was measured to soften with increasing distance from its pulsar [298]. We expect similar behavior from Geminga if the same mechanisms are at work, the latter of which would be a distinct signature of $e^\pm$ cooling [298]. Also, studying the extended TeV emission from an old, radio-quiet neutron star should have implications for some heretofore unidentified TeV sources.

**Electrons and Positrons:** Due to the spin down of the pulsar, it is possible that the Geminga source was much brighter in the past and dominated the TeV sky. It is from this time that multi-GeV $e^\pm$ may still be reaching us today. If Geminga does account for a substantial fraction of the total $e^- + e^+$ spectrum, a mild anisotropy may be present. Since the distance to Geminga does not greatly exceed the scale for field fluctuations of $\sim 100$ pc [317], with detailed multi-wavelength studies, local diffusion parameters might be determined (which may differ from global values estimated across the Galaxy). Additionally, as Geminga remains a source of $\sim 100$ TeV $e^\pm$, it may result in a $> 10$ TeV lepton flux at Earth.

The flux level reached by our new derived high-energy limits is encouraging for the prospects of upcoming generations of electron experiments, and already constrains the presence of a nearby, very-high-energy source. Meanwhile, our handling of propagation in the presence of energy losses eliminates the appearance of spurious
superluminal solutions. This serves as a step towards a more accurate description of time-dependent $e^\pm$ propagation. We have considered the nearby pulsar Geminga, detected in multi-TeV gamma rays, as the best motivated source for multi-TeV electrons reaching Earth today, although Vela X may also contribute if $e^\pm$ have escaped.

A nearby supernova remnant could result in electrons [339], although age and distance play significant roles, with the recent history of acceleration being quite important due to the short cooling time at high energy. Prominent SNRs include [340] the Vela SNR ($\sim 250$ pc), Cyngus Loop ($\sim 440$ pc), Monogem Ring ($\sim 300$ pc), and Loop I ($\sim 100$ pc). Due to their large angular extents, it is difficult to establish with ACTs whether these are active. Searches via synchrotron radiation or wide-field gamma-ray instruments are better suited. The only SNR potentially within 500 pc and detected in TeV gamma rays is Vela Junior, with an uncertain distance ranging from $\sim 200$ pc (corresponding to an age of only $\sim 500$ yr) to $\sim 1$ kpc ($\sim 5000$ yr) [341], while the “Boomerang” SNR/PWN at $\sim 800$ pc has been seen by Milagro [297].

If features in the $e^- + e^+$ spectrum, such as the change in slope at $\sim 100$ GeV seen by Fermi or the drop at $\sim 1$ TeV measured by HESS, are due to a transition between different sources, classes, or $e^\pm$ populations, these should correspond to features in the positron fraction. Measurements of the separate $e^-/e^+$ spectra will be vital in determining the Galactic component and in distinguishing SNR, dark matter, and pulsar contributions. For regions of pulsar domination, the positron fraction should saturate at $\sim 50\%$, while SNRs result entirely in primary electrons. While PAMELA can reach 300 GeV for $e^+$ [342], AMS-02 will go to $\sim 1$ TeV [343], and ACT measurements using the varying position of the moon shadow for $e^+/e^-$ may reach several TeV [344]. These combined observations give hope for connecting very-high-energy emission to lower energies to build a complete picture of the spectra of cosmic-ray electron and positrons.
CHAPTER 6
HIGH-ENERGY ASTROPHYSICAL SIGNALS FROM DARK MATTER

Alexander, king of Macedon, began to study geometry; unhappy man, because he would thereby learn how puny was that Earth of which he had seized but a fraction!

Seneca, Epistle XCI

6.1 Introduction

Dark matter continues to live up to its name [3], despite accumulated evidence of its existence from observations of large-scale structure formation, galaxy cluster mass-to-light ratios, and galactic rotation curves.\textsuperscript{1} An attractive approach towards revealing dark matter’s particle identity is to search for its signature in radiation backgrounds, either from the Milky Way or in the isotropic diffuse photon background (iDPB), which can contain both cosmological and Galactic halo contributions. Dark matter might be comprised of particles that can decay with finite lifetimes much longer than the age of the universe. In such scenarios, the resultant fluxes of decay products depend on the amount of dark matter present alone, as opposed to self-annihilation,

\textsuperscript{1}The content of this Chapter is based in large part on our work in Refs. [116, 118].

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which, being dependent on particle density squared, is very sensitive to assumptions concerning details of dark matter clustering.

A wide variety of decaying dark matter models have been examined in regards to their observable implications [106, 345]. Among late decaying dark matter models, sterile neutrinos with multi-keV masses have been extensively studied as dark matter candidates [346], with strong constraints placed on their decays, e.g., Refs. [346, 347, 348] and references therein. The decay of moduli dark matter [349] with masses of several hundred keV may contribute to the sub-MeV iDPB. The dark matter model of Ref. [350], inspired from minimal universal extra dimensions or supersymmetry [351], with a mass scale of hundreds of GeV, is advocated as the source of the iDPB in the MeV range [111], which has yet to be accounted for with conventional sources (e.g., supernovae [352] or active galactic nuclei [353]). Similarly, decaying gravitino dark matter in R-parity breaking vacua [354], with multi-GeV masses, has been suggested as an explanation of iDPB spectral features in the GeV range [355, 356].

Rather than focusing on a particular model, we first consider a generic decaying dark matter scenario in which the decay of the parent particle is dominated by a monochromatic photon emission [116]. We assume that the lifetime of the parent particle, $\tau$, is much longer than the age of the universe ($\tau_0 \approx 4.5 \times 10^{17}$ s), thus its cosmological abundance has not changed significantly since the time of dark matter decoupling. The decay under consideration is $\chi \rightarrow \chi' + \gamma$, where $\gamma$ is a monochromatic photon emitted with energy $\varepsilon$. In general, $\tau$ and $\varepsilon$ will depend on the masses of the parent and the daughter particles ($m_\chi, m_{\chi'}$) and their splitting, $\Delta m = m_\chi - m_{\chi'}$.

The flux of photons from dark matter decays is inversely proportional to both the particle lifetime (fixing the decay rate per particle as specified by a particular theoretical model) and the mass of an individual particle (yielding the total number of particles in a fixed amount of dark matter). Thus, gamma-ray observations allow
Figure 6.1: Model-independent constraints on the product of mass and lifetime, $m\chi\tau$, versus the energy carried away by the monochromatic photon emission, $\varepsilon$, for a generic late-decaying dark matter model: $\chi \rightarrow \chi' + \gamma$. Regions excluded by either the gamma-ray line emission limits from the Galactic Center region or overproduction of the isotropic diffuse photon background are shown, together with preferred ranges of parameters from three well-studied models [116].
us to place constraints only on the degenerate product $m_\chi \tau$ versus $\varepsilon$, as we display in Fig. 6.1. As we will discuss in detail, below the jagged line between 0.02–8 MeV, the gamma-ray line signal from the Galactic Center (GC) region due to dark matter decays violates the corresponding limit obtained with the SPI spectrometer on INTEGRAL satellite [112]. Additionally, the iDPB, as determined from SPI [113], COMPTEL [110] and EGRET [114, 115] data, is overproduced (assuming it is fully accounted by late dark matter decays in a given energy band) in the triangular region, even disregarding any contributions from known astrophysical sources. We also show three representative scenarios, inspired by the theories of sterile neutrinos, R-parity breaking vacua, and mUED. Since $m_\chi$, $m_\chi'$ and $\Delta m$ are not necessarily predetermined, they may be adjusted to yield the displayed curves relating $m_\chi \tau$ and $\varepsilon$.

The possibility of using dark matter to provide energetic positrons and electrons to explain the GeV positron excess seen by PAMELA [51], the “ATIC bump” [320, 319], and the less-anomalous $e^- + e^+$ spectra measured by Fermi [48] and HESS [49, 50] has sparked considerable theoretical interest. Generating the required $e^\pm$ flux through annihilations in the smooth component of the Milky Way’s dark matter halo requires both a much larger annihilation cross section than might be expected for a thermal relic [107] and a large branching ratio to charged leptons.

We examine the observational consequences of annihilations occurring within the dark matter substructure of the Milky Way [118]. Substructure differs from the smooth halo in its spatial distribution, which is less centrally-concentrated [358, 359], and in its characteristic velocity dispersions, which are colder. We focus on the case of enhanced annihilation to charged leptons in Sommerfeld models with an annihilation cross section that increases with decreasing relative particle velocity. Annihilations to charged particles necessarily produce internal bremsstrahlung (IB) gamma rays [360]. High-energy electrons and positrons will also produce gamma rays through various
energy-loss processes. Beyond a few tens of kiloparsecs from the Galactic Center, the dominant loss process for electrons is the inverse-Compton (IC) upscattering of CMB photons.

We calculate the expected high-latitude gamma-ray emission from IB and IC resulting from dark matter annihilation in substructure throughout the Milky Way. In particular, rather than considering only the serendipitous presence of a single, nearby, massive dark matter clump, we account for the collective emission from the entire subhalo population. Noting that electron- and gamma-ray-induced showers are difficult to distinguish in air Čerenkov telescopes, we discuss how constraints on TeV gamma-ray fluxes can be obtained from the cosmic-ray electron measurements made by HESS to limit annihilations to lepton pairs via their IB emission.

6.2 Milky Way Gamma-Ray Line Search

A monochromatic line will be most detectable locally, where cosmological redshifting is of no concern. Fortunately, a search for diffuse gamma-ray line emission in the energy range 0.02-8 MeV from the GC region has been conducted by Teegarden and Watanabe using the SPI spectrometer on the INTEGRAL satellite [112], which recovered the known astrophysical diffuse line fluxes, such as the 511 keV positron annihilation line [361]. The excellent energy resolution of SPI enabled them to place very strict constraints on potential unidentified emission lines, with an energy dependent $3.5 \sigma$ flux limit, $F_{\text{lim}}(E)$, from an angular region within a 13° radius of the GC (which we refer as the GC region). This limit, reproduced in the top panel of Fig. 6.2, can be compared to the expected gamma-ray flux arising from late dark matter decays in the GC region, which we calculate following the methods in Ref. [348].
Figure 6.2: **Top:** Limits on the diffuse gamma-ray line emission from the Galactic Center region (an angular region within a 13° radius) as adopted from Ref. [112]. **Bottom:** Representative measurements of the diffuse photon background from SPI [113], COMPTEL [110] and EGRET [115] in the energy range around 0.01 MeV–100 GeV. The thick solid line, summarizing the overall trend of the data, is to be compared to predictions of decaying dark matter scenarios [116].
We first define a dimensionless line-of-sight integral at an angle $\psi$ relative to the GC,

$$J(\psi) = \frac{1}{\rho_{sc} R_{sc}} \int_0^{\ell_{max}} d\ell \, \rho \left( \sqrt{R_{sc}^2 - 2 \ell R_{sc} \cos \psi} + \ell^2 \right),$$  \hspace{1cm} (6.2.1)

where $\rho$ is the density of the dark matter in the halo as a function of the distance from the GC. This is normalized to the dark matter density ($\rho_{sc} = 0.3$ GeV cm$^{-3}$) at the solar circle ($R_{sc} = 8.5$ kpc) so that $\rho_{sc} R_{sc} \simeq 8 \times 10^{21}$ GeV cm$^{-2}$. Note that this arbitrary normalization is needed to make $J$ dimensionless and will be canceled-out later. The upper limit of this integration,

$$\ell_{max} = \sqrt{(R_{MW}^2 - \sin^2 \psi R_{sc}^2) + R_{sc} \cos \psi},$$  \hspace{1cm} (6.2.2)

depends on $R_{MW}$, the assumed size of the halo. $J$ is relatively insensitive to $\ell_{max}$ as long as $R_{MW}$ is large. The intensity of photons (number flux per solid angle) from the same direction,

$$I(\psi) = \frac{\rho_{sc} R_{sc}}{4 \pi m_\chi} J(\psi),$$  \hspace{1cm} (6.2.3)

can be integrated over a circle of radius $\psi$ around the GC (covering a patch of area $\Delta \Omega = 2\pi (1 - \cos \psi)$) to obtain the corresponding total flux,

$$F = \int_{\Delta \Omega} d\Omega' \, I(\psi') = \frac{\rho_{sc} R_{sc}}{4 \pi m_\chi \tau} \int_{\Delta \Omega} d\Omega' \, J(\psi').$$  \hspace{1cm} (6.2.4)

The limit reported in Ref. [112] has been obtained by subtracting the average flux measured at regions away from the GC region ($\psi > 30^\circ$) from the average flux measured inside the GC region ($\psi < 13^\circ$) to eliminate instrumental backgrounds.

Thus, the constraining power of this limit for decaying dark matter scenarios depends on the enhancement of the expected signal towards the GC region. Both theoretical and observational studies strongly suggest that the central regions of dark matter halos are significantly denser and, moreover, the column depth is higher towards the
GC direction relative to off-axis lines-of-sight. We have reproduced the impact of this subtraction (see Ref. [348] for details) by calculating a parameter

$$\zeta_{\text{lim}} = \int_{\Delta \Omega} d\Omega' \left[ J(\psi') - \overline{J}_{>30^\circ} \right].$$

(6.2.5)

which ranges between $\sim 0.5 - 1.5$ for various dark matter halo fitting profiles commonly used in the literature [362]. Here $\overline{J}_{>30^\circ}$ is the average of $J$ away from the GC region. We also note that the results that we adopted from Ref. [112] are based on an assumption that the expected line signal has a Gaussian source profile, while a flat source profile could yield limits that are weaker by up to a factor of $\sim 2$. Moreover, one would expect to see these limits improve as the amount of available data increases in time. In the rest of our study, we choose a conservative value, $\zeta_{\text{lim}} \simeq 0.5$, which can be realized only for profiles that are rather flat inside the solar circle. While this mostly protects our conclusions from uncertainties in the halo profile, our subsequent result can be easily rescaled for a different value.

The predicted gamma-ray emission line flux due to dark matter decays at a given $\varepsilon$ must not exceed the corresponding limits from the GC region, thus

$$\rho_{sc} R_{sc} \frac{\zeta_{\text{lim}}}{4\pi m_\chi \tau} < \mathcal{F}_{\text{lim}}(E = \varepsilon).$$

(6.2.6)

Rearranging this equation yields our model independent constraint,

$$m_\chi \tau > \frac{\rho_{sc} R_{sc} \zeta_{\text{lim}}}{4\pi \mathcal{F}_{\text{lim}}(\varepsilon)} \simeq \frac{3 \times 10^{20} \text{ GeV cm}^{-2}}{\mathcal{F}_{\text{lim}}(\varepsilon)},$$

(6.2.7)

as shown in Fig. 6.1 (region below the jagged line). The expected dark matter decay flux is inversely proportional to $m_\chi \tau$, which leads to an overproduction of gamma rays for $m_\chi \tau \lesssim 10^{25}$ GeV s in the energy range 0.02-8 MeV. Thus the area below the jagged line is excluded by the the diffuse gamma-ray line emission limits from the GC region.
6.3 Isotropic Diffuse Photon Background

While stringent limits on line emission from the GC region are only available in a rather limited energy range (0.02–8 MeV), the iDPB is measured over a broad range by many instruments. In the bottom panel of Fig. 6.2, we display three recent determinations of iDPB in different ranges of energy from SPI [113], COMPTEL [110] and EGRET [115], which are consistent with others measurements (see, e.g., Ref. [357, 363]). The thick dotted line represents the global trend of the data to be used for comparison. We choose the terminology “isotropic diffuse” photon background (iDPB), as opposed to “cosmic” or “extragalactic”, since the contribution from sources in the Milky Way or its halo is not clear, and iDPB can include gamma-ray line signals that could not have been resolved by COMPTEL or EGRET. While it is generally thought that AGN are responsible for the emission in the ∼ keV [364] and ∼ GeV [365] ranges, the origin of the iDPB, especially in the MeV regime, is far from being settled, with various scenarios having been entertained [352, 353]. It is then of interest to determine just how much of the iDPB can possibly be accounted for by late decaying dark matter.

6.3.1 Dark Matter Decays in the Halo

While the photon signal from dark matter decays in the Galactic halo is enhanced towards the GC, as has been utilized for our constraints in the earlier section, it also contains an apparently isotropic contribution. The limited energy resolution of past gamma-ray detectors could not distinguish monochromatic line emission from the Galactic halo from a truly cosmological signal. The intensity of the isotropic halo contribution, \( I_{iso} \), can be estimated from a line of sight integration in the anti-GC direction, \( J_{iso} = J(180^\circ) \sim 1 \), as this is the minimum contribution from the dark matter halo of the Milky Way. Regardless of the underlying halo profile, this number
is relatively robust, being mostly dependent on the dark matter density at the solar circle. The intensity of this isotropic component is

$$I_{\text{iso}} = \frac{\rho_{sc} R_{sc}}{4\pi m_{\chi}\tau} J_{\text{iso}}.$$  

(6.3.1)

We present a representative spectrum for this isotropic signal in Fig. 6.3 (dotted line), after convolution with a Gaussian of $\sim 10\%$ width to simulate the energy resolution of a typical detector. We have chosen $\varepsilon = 1$ MeV, with $m_{\chi}\tau = 7 \times 10^{24}$ GeV s, the maximum value allowed by the the line emission bounds from the GC region (Fig. 6.1). For these parameters, the isotropic contribution of the dark matter decays in the Galactic halo alone to the iDPB is less than 2% (in a bin of logarithmic width 0.4 dex centered around $\varepsilon = 1$ MeV). Note that the average flux expected from the decays in the Galactic halo (which is more directional, peaking toward the GC region) can be at most several times larger than this isotropic component since we are dealing with decaying dark matter particles (contrary to self-annihilating dark matter, which is highly sensitive to the details of dark matter clustering).

### 6.3.2 Cosmological Dark Matter Decays

We now evaluate the contribution of truly cosmological dark matter decays to the iDPB. For late decaying particles ($\tau \gg \tau_0$), the comoving dark matter density has remained nearly constant since the early universe. The comoving decay rate is then simply proportional to the dark matter fraction ($\Omega_{\chi} \simeq 0.25$) of the critical density of the universe, $\rho_c$, and is given as $\rho_c \Omega_{\chi}/(m_{\chi}\tau)$. The diffuse gamma-ray flux (per solid angle per unit energy) arising from the decays can be calculated by considering the contributions from all redshifts (analogous to [366]),

$$\frac{d\Phi}{dE} = \frac{1}{4\pi H_0} \int \frac{dz}{h(z)} \frac{\rho_c \Omega_{\chi}}{m_{\chi}\tau} \delta(E(1+z) - \varepsilon),$$  

(6.3.2)
Figure 6.3: Photon spectrum from isotropic Galactic halo decays (dotted line) for $\varepsilon = 1$ MeV, with $m_\chi \tau \approx 7 \times 10^{24}$ GeV s chosen from Fig. 6.1 such that the line emission bounds from the Galactic Center region are saturated. Also displayed are the spectra from cosmological decays (dashed line) and the total spectrum (solid line), which falls well below the isotropic diffuse photon background (thick solid line) [116].

where $h(z) = \left[(1 + z)^3 \Omega_M + \Omega_\Lambda \right]^{1/2}$, $\Omega_M \approx 0.3$, $\Omega_\Lambda \approx 0.7$, $H_0 = 70$ km s$^{-1}$ Mpc$^{-1}$, and $c = 3 \times 10^{10}$ cm s$^{-1}$ (so that $c/H_0 \approx 1.3 \times 10^{28}$ cm and $\rho_c = 5.3 \times 10^{-6}$ GeV cm$^{-3}$).
The integration can be eliminated after using the $\delta$-function identity: $\delta(ax - b) = \delta(x - b/a)/a$, simplifying the result to

$$
\frac{d\Phi}{dE} = \frac{1}{4\pi H_0} \frac{c \rho_c \Omega_X}{m_\chi \tau} \frac{1}{E} \frac{\Theta(\varepsilon - E)}{\sqrt{\varepsilon/E}^3 \Omega_M + \Omega_a},
$$

where $\rho_c$ is substituted and $\Theta$ is a step function. We show this, using the same parameters as in the preceding subsection and again smoothing with a $\sim 10\%$ Gaussian, in Fig. 6.3 (dashed line). As seen in the figure, this cosmological flux is slightly lower than the isotropic contribution from the Galactic halo, and their sum (solid line) still falls well short of the observed signal, restricting their combined contribution to the iDPB to be less than 4%.

To quantify and generalize our observations, we calculate the expected total (cosmological plus isotropic Galactic halo) spectrum for all values of $m_\chi \tau$ and compare to the iDPB (as denoted by the thick trend curve in Fig. 6.2), integrating both in a bin of logarithmic width 0.4 dex centered around $\varepsilon$. This choice encompasses most of the expected signal where the decay spectrum peaks, and both exceeds the experimental energy resolution and the uncertainties on the determination of the iDPB. In Fig. 6.1, the region in which dark matter decays overproduce the iDPB is shown (triangular region).

Above this region, decaying dark matter alone cannot fully account for the iDPB. In fact, since there should be additional contributions from AGN at both low and high energies [364, 365], the actual bound on the parameter $m_\chi \tau$ will be even more stringent than the one presented. Combining the iDPB overproduction constraint and the gamma-ray line emission limit from the GC region model-independently excludes a sizable region in the parameter space of $m_\chi \tau$ versus $\varepsilon$, with the latter picking up when the former is exhausted at $\varepsilon \approx 8$ MeV.
6.4 Decaying Dark Matter Models

While we derive our constraints for a decay scenario that is dominated by monochromatic photon emission, there may be additional modes of decay or self-annihilations producing other signals. Our constraints on the lifetime of the dark matter candidate via monochromatic photon emission could be generalized to the total lifetime including other decay channels, as long as the latter is long enough to justify the assumption that the cosmological abundance of the parent particle has not changed significantly.

For the generic decay we are considering, the energy of the emitted photon is dictated by the splitting, $\Delta m$, as follows. When $\Delta m \ll m_{\chi'}$ (or equivalently $m_{\chi} \simeq m_{\chi'}$), the recoil of the daughter can be neglected, so that $\varepsilon \to \Delta m$. For $\Delta m \gg m_{\chi'}$ (or $m_{\chi} \gg m_{\chi'}$), two relativistic particles are produced, so that $\varepsilon \to \Delta m/2 \simeq m_{\chi}/2$. Generally, models lie in one of these two regimes. To emphasize the generality of our constraints, now we discuss particular scenarios.

For example, WIMPs with weak-scale masses and cross sections may have monochromatic decays. The decay process between the two lightest Kaluza-Klein (KK) particles, the KK hypercharge gauge boson, $B^1$ and KK graviton, $G^1$ in mUED models, and the decay between the two lightest particles in SUSY theories, the Bino-like neutralino, $\tilde{B}$ and gravitino, $\tilde{G}$ are well-studied examples. The mass scales of these candidates are $\sim 800$ GeV for the former and $\sim 80$ GeV for the latter. The decay rates in these theories [111] are highly suppressed due to the weakness of gravity and is given by

$$\tau \simeq \frac{4.7 \times 10^{22}}{b} \frac{s}{\left( \frac{\Delta m}{\text{MeV}} \right)^3}. \quad (6.4.1)$$

where the parameter $b$ is identified as $(2, 10/3, 1, 2)$ for each of the decay reactions ($G^1 \to B^1 + \gamma$, $B^1 \to G^1 + \gamma$, $\tilde{G} \to \tilde{B} + \gamma$, $\tilde{B} \to \tilde{G} + \gamma$). The lifetime requirement of
\( \tau \gg \tau_0 \) translates into \( \Delta m < 30 \text{ MeV} \). Since \( \Delta m \ll m_\chi \), the energy carried away by the emitted photon is \( \varepsilon \simeq \Delta m \). Eq. (6.4.1) can be rearranged as

\[
m_\chi \tau \simeq 4.7 \times 10^{22} \text{ s} \left( \frac{m_\chi}{b} \right) \left( \frac{\varepsilon}{\text{MeV}} \right)^{-3},
\]

which relates \( m_\chi \tau \) to \( \varepsilon \) in terms of a single parameter: \( m_\chi/b \). We plot \( m_\chi \tau \) versus \( \varepsilon \) in Fig. 6.4 for \( m_\chi/b \simeq 300 \text{ GeV} \) to represent mUED. One sees that the Milky Way constraint requires \( \varepsilon \leq 1.5 \text{ MeV} \), which is a very strict limit as the lifetime is proportional to \( \varepsilon^{-3} \), i.e., the decay rate increases by almost an order of magnitude from 1 MeV to 2 MeV. This translates to the restriction of \( \Delta m \lesssim 1.5 \text{ MeV} \), which is far stricter than the necessary condition to have a long-lived candidate, \( \Delta m < 30 \text{ MeV} \).

In the R-parity violating supersymmetric extension of the standard model, the lightest supersymmetric particle is again a gravitino that might not be stable on cosmological time scales against decay into a photon and neutrino (\( \tilde{G} \to \nu + \gamma \)) through a small photino-neutrino mixing \( |U_{\tilde{G}\nu}| \). The lifetime of the gravitino in this model [356] is

\[
\tau \simeq 3.8 \times 10^{27} \text{ s} \left( \frac{|U_{\tilde{G}\nu}|}{10^{-8}} \right)^{-2} \left( \frac{m_\chi}{10 \text{ GeV}} \right)^{-3},
\]

with the resulting photon and neutrino each carrying an energy of \( \varepsilon = m_\chi/2 \). We can rewrite this equation in terms of \( m_\chi \tau \) versus \( \varepsilon \) as

\[
m_\chi \tau \simeq 10^{14} \text{ GeV s} \left( |U_{\tilde{G}\nu}| \right)^{-2} \left( \frac{\varepsilon}{\text{GeV}} \right)^{-2}.
\]

We plot this relation for \( |U_{\tilde{G}\nu}| = 10^{-8} \) in Fig. 6.1, which shows that the contribution of this decay model to the iDPB will be significant around \( \varepsilon \sim 5 \text{ GeV} \) (corresponding to \( m_\chi \sim 10 \text{ GeV} \)) agreeing with Ref. [356]. Slightly above/below \( m_\chi \sim 10 \text{ GeV} \), either its contribution is negligible or vastly overproduces the iDPB.

Dark matter models involving keV-mass sterile neutrinos, in their simplest description, require only two parameters, the sterile neutrino’s mass and mixing with
active neutrinos. The decay chain for sterile neutrinos is $\nu_s \rightarrow \nu_{e,\mu,\tau} + \gamma$, with a radiative lifetime [367] (for Dirac neutrinos) of

$$\tau = 1.5 \times 10^{22} s \left(\sin^{-2} 2\theta\right) \left(\frac{m_s}{\text{keV}}\right)^{-5}.$$  \hfill (6.4.5)
This can similarly be rearranged, keeping in mind that the energy of the parent sterile neutrino is split equally between the photon and the daughter neutrino ($\varepsilon = m_s/2$), as

$$m_s \tau \simeq 10^{15} \text{ GeV s} \left(\sin^{-2} 2\theta\right) \left(\frac{\varepsilon}{\text{keV}}\right)^{-4}, \quad (6.4.6)$$

which has only one free parameter, $\sin^2 2\theta$. For illustration, we plot Eq. (6.4.6) in Fig. 6.1 for $\sin^2 2\theta = 10^{-18}$. As seen in the figure, and has been established in Ref. [348], the gamma-ray line emission limit from the Galactic Center region provides quite stringent restrictions on sterile neutrino dark matter, which can be several orders of magnitude stronger than constraints from overproduction of the iDPB. Interestingly, all three models we have discussed have the form $m_\chi \tau \propto \varepsilon^{-\alpha}$, where $\alpha = 3, 2, 4$ respectively, as can also be noticed from the varying slopes of the lines representing the models in Fig. 6.1.

### 6.5 Annihilation in Substructure

To describe their structural properties, we assume each individual dark matter subhalo to be parametrized by a NFW density profile [118],

$$\rho_{\text{sub}} = \frac{\rho_s}{r/r_s (1 + r/r_s)^2}, \quad (6.5.1)$$

where $r_s$ is a scale radius and $\rho_s$ a characteristic density. The differential luminosity (photons or particles per energy per time) $L$ of a subhalo, for an annihilation cross section $\langle \sigma v \rangle_0$ that is independent of radius or velocity, is

$$L = K \int dV_{\text{sub}} \rho_{\text{sub}}^2 \propto \rho_s^2 r_s^3 \propto M \frac{c^3}{f^2(c)}, \quad (6.5.2)$$

where $M$ is the subhalo mass, $c = r_{\text{vir}}/r_s$ is the concentration, $f(c) = \ln(1 + c) - c/(1 + c)$, and the particle physics-dependence of the annihilation rate is isolated in

$$K = \frac{\langle \sigma v \rangle_0 dN}{2m_{\text{DM}}^2 dE}, \quad (6.5.3)$$
with $m_{DM}$ the dark matter particle mass and $dN/dE$ the particle spectrum produced per annihilation.

Numerical simulations find a relation between concentration and mass for subhalos that varies as a function of distance from the Galactic center [358, 359]. This is a natural consequence of tidal stripping, which more effectively removes mass from the outer regions of the subhalos while leaving the core relatively unscathed. Thus, for a given subhalo mass, subhalos located nearer to the Galactic center will be more luminous than those at large radii. We adopt the modified Bullock et al. [368] relation for low-mass halos, with radial dependence, from Ref. [369]

$$c_{\text{sub}}(M, r) = 18 \left( \frac{M}{10^8 M_\odot} \right)^{-0.06} \left( \frac{r}{r_{\text{field}}} \right)^{-0.286},$$

(6.5.4)

where $r_{\text{field}} = 402$ kpc is the radius where equal mass subhalos and field halos have the same concentration. Using this relation with Eq. (6.5.2), we can approximate the differential luminosity of a subhalo of mass $M$ at a radius $r$ from the Galactic Center by

$$L(M, r) = K \mathcal{L}(M) \left( \frac{r}{r_{\text{field}}} \right)^{-0.7},$$

(6.5.5)

where we have defined

$$\mathcal{L}(M) = \int dV_{\text{sub}} \rho_{\text{sub}}^2 \simeq \mathcal{L}_0 \left( \frac{M}{M_0} \right)^{0.87}$$

(6.5.6)

to describe the dependence of the annihilation rate on the structural properties of the subhalo. We note that the dependences on $M$ and $r$ are not formally separable, but are weak enough that Eq. (6.5.5) gives a reasonable approximation. The mass of Canes Venatici I [370], assuming an NFW density profile and concentration $c = 19.5$, is used to normalize $\mathcal{L}_0$ and $M_0$.

We assume a power-law mass function for the subhalos [358, 359], $dN/dM \propto M^{-\alpha}$ with $\alpha = 1.9$ and extrapolate this relation to a minimum subhalo mass $M_{\text{min}} =$
Figure 6.5: Angular distribution of the emission from dark matter annihilation ($J(\psi)$ from Eq. (6.5.10)). **Top panel:** Galactic substructure assuming radial distributions from Aquarius (solid) and Via Lactea II (dashed), compared to the smooth halo (dotted) assuming NFW, Einasto ($\alpha = 0.17$, $r_\text{-2} = 20$ kpc), and cored isothermal ($r_{\text{core}} = 5$ kpc) density profiles (as labeled). **Bottom panel:** The substructure emission assuming the Einasto radial distribution with Aquarius parameters, broken down into contributions from radial shells surrounding the Galactic Center (with distances as labeled). Peaks are due to shell boundaries [118].

$10^{-6}M_\odot$. Noting that $(dN/d\mathcal{L}) = (dN/dM)(dM/d\mathcal{L})$, integration over the subhalo population yields

$$\mathcal{L}_{\text{subs}} = \int_{\mathcal{L}_{\text{min}}}^{\mathcal{L}_{\text{max}}} d\mathcal{L} \frac{dN}{d\mathcal{L}},$$

(6.5.7)

which contains the dependence of the annihilation rate on the structural properties and mass function of the subhalos, and is independent of position in the Galaxy.
We model the subhalo number density (i.e., number of subhalos per volume) at a radius \( r \) from the Galactic Center with an Einasto profile

\[
n_{\text{subs}}(r) \propto \exp \left\{ -\frac{2}{\alpha_{\text{subs}}} \left[ \left( \frac{r}{r_{-2}} \right)^{\alpha_{\text{subs}}} - 1 \right] \right\},
\]

(6.5.8)

with \( \alpha_{\text{subs}} = 0.68 \) and \( r_{-2} = 199 \) kpc, as found by the Aquarius Project [359]. We assume a NFW profile for the smooth halo with \( r_s = 20 \) kpc, \( r_{\text{vir}} = 255 \) kpc, and \( M_{\text{vir}} = 1.9 \times 10^{12} M_\odot \), and normalize the subhalo distribution such that a fraction \( f_{\text{subs}} = 0.15 \) of \( M_{\text{vir}} \) is bound in substructure. For comparison, we also consider a number density distribution as in Via Lactea II, \( n_{\text{VL-II}}(r) \propto (1 + r/R_s)^{-2} \) [369], with \( R_s \approx 20 \) kpc.

The intensity \( (I) \) of gamma-ray emission resulting from annihilation in substructure at an angle \( \psi \) from the Galactic Center is

\[
I(\psi) = \frac{K}{4\pi r_\odot \rho_\odot^2} \mathcal{J}(\psi),
\]

(6.5.9)

in which \( r_\odot = 8.5 \) kpc and \( \mathcal{J}(\psi) \) is given by the line-of-sight integral

\[
\mathcal{J}(\psi) = \frac{L_{\text{subs}}}{r_\odot \rho_\odot^2} \int_{\text{los}} ds \ n_{\text{subs}}(r) \left( \frac{r}{r_{\text{field}}} \right)^{-0.7},
\]

(6.5.10)

where \( r = r(s, \psi) \). This is shown in Fig. 6.5 for the Aquarius and VL-II subhalo radial distributions, along with the analogous term for three possible smooth halo density profiles. Here, we have chosen \( L_{\text{subs}} \) such that the annihilation rate per volume is matched to the smooth halo at \( r_\odot \). Both subhalo models produce a nearly-isotropic angular signal, much less strongly-peaked towards the Galactic Center than that of the smooth halo. In the following, we adopt the Aquarius distribution from Eq. (6.5.8).

### 6.6 Sommerfeld Enhancement and the Annihilation Rate

A challenge in attributing the positron excess to dark matter annihilation is the need for a much larger annihilation cross-section than expected for a thermal relic. One
way to accomplish this is by introducing a scalar or vector boson that mediates an intermediate-range force between dark matter particles and can dramatically enhance the cross section for annihilation at low relative velocities. For dark matter, this can be through Standard Model gauge bosons (e.g., [371]) or a new mediator particle $(\phi)$. Here, we consider the possibility that annihilations in substructure, rather than in the smooth halo, are the dominant contributor to the measured local lepton flux. Ref. [118] showed in detail how this can arise for a range of parameters in model based on a Sommerfeld enhancement.

The total local annihilation rate per unit volume from subhalos is

$$\Gamma_{\text{subs}} = \frac{\langle \sigma v \rangle_0}{2m_{\text{DM}}^2} L_{\text{subs}} n_{\text{subs}}(r_{\odot}) \left( \frac{r_{\odot}}{r_{\text{field}}} \right)^{-0.7}.$$  

(6.6.1)

The large local number density of subhalos $(n_{\text{subs}}(r_{\odot}) \sim 2 \times 10^8 \text{ kpc}^{-3})$ allows us to neglect the discreteness of subhalos as sources and use the annihilation rate averaged over the subhalo population via $L_{\text{subs}}$. Including subhalos of all masses in the calculation of the local flux results in an overestimate of $L_{\text{subs}}$ of not more than 25% relative to assuming a maximum mass of $10^7 M_{\odot}$, a minor effect that we neglect. This can be understood from Eq. (6.5.7): in absence of a Sommerfeld enhancement, the total luminosity roughly scales $\propto \int dM \frac{dN}{dM} \propto M_{\text{min}}^{-0.03}$, so that each decade of mass contributes nearly equally. (This has a much weaker dependence on the choice of lower cutoff than the $L_{\text{tot}} \propto M_{\text{min}}^{-0.226}$ of Ref. [372].)

The choices we have made in defining our substructure model are conservative: using a higher normalization of the $c(M)$ relation (as in Aquarius), a steeper mass function (e.g., $\alpha = 2$), a larger mass fraction in substructure, or considering substructures-within-substructure [358, 369, 372] would all increase the subhalo annihilation rate relative to the smooth halo. Although these mechanisms can not increase the annihilation rate in substructure sufficiently to explain the electron/positron data with a
standard thermal cross section, they could significantly decrease the required saturation cross section. For the purposes of the following, we assume that annihilation to charged leptons in substructure dominates over the smooth halo and accounts for 100% of the anomalous fluxes.

6.7 Internal Bremsstrahlung

For pure leptonic final states, the only gamma-ray emission directly resulting from annihilations in substructure is internal bremsstrahlung (IB): $\chi\chi \rightarrow \ell^+\ell^-\gamma$. We consider a few representative annihilation channels: direct annihilation into $2\mu$ and $2\tau$, and annihilation to $4\tau$ through a new particle $\phi$ ($\chi\chi \rightarrow \phi\phi$ and each $\phi \rightarrow 2\tau$), with cross sections required to explain the combined PAMELA/Fermi data [375, 374].

We calculate the IB spectra for the two lepton cases as in Ref. [373],

$$\frac{dN_{IB}}{dE} = \frac{1}{E} \frac{\alpha_{em}}{\pi} \left(1 + x^2\right) \ln \left(\frac{4m_{DM}^2 x}{m_{\mu,\tau}^2}\right), \quad (6.7.1)$$

where $x = 1 - E/m_{DM}$ and $\alpha_{em} \simeq 1/137$. For $m_{DM} \gg m_{\mu,\tau}, E$, this has the behavior $dN_{IB}/dE \propto E^{-1}$ [360]. We similarly calculate the four lepton case, as detailed in Ref. [376]. For the $2\mu$ channel, we omit the gamma-ray contribution from muon decay, $\mu \rightarrow e\nu_e\nu_\mu\gamma$, which is negligible for the scenario considered here.

The IB gamma-ray intensity from annihilation in substructure at $\psi = 180^\circ$ (the minimum of the dark matter signal) is shown in Fig. 6.6 for the above cases, along with gamma rays resulting directly from pionic tau decays in the $2\tau$ scenario (using DarkSUSY [377]). Considering a smooth halo model would result in signals smaller by a factor of $\sim 2 - 3$ for all profiles at large angles. Although we do not otherwise consider them, models based on decays in the smooth halo would give signals comparable to those shown in Fig. 6.6.
Directly measuring such a diffuse gamma-ray flux at TeV energies is presently challenging, in part due to the effective area of Fermi saturating with energy [378]. While ground-based Čerenkov telescopes do not have this problem, as discussed previously, the electromagnetic showers that they observe are quite similar for TeV electrons and gamma rays, making them difficult to separate. Based on observations of fields far from the Galactic plane, HESS has reported measurements of the $e^- + e^+$ spectrum into the TeV regime [49, 50]. In principle, a nearly-isotropic ∼ TeV gamma-ray flux from dark matter annihilation could result in an apparent feature in this spectrum. As noted in Refs. [49, 50], there should be little contribution from extragalactic TeV gamma rays.

In Fig. 6.6, we show the HESS $e^- + e^+$ spectrum, which can be regarded as a conservative upper limit on isotropic TeV gamma rays. The maximum gamma-ray fraction of this measurement is likely $\lesssim 10\%$ (although systematic uncertainties could result in as much as $\sim 50\%$) [49]. It is likely that a dedicated analysis that accounts for the fields of view observed by HESS and determines a limit on the photon fraction in a given energy interval can strengthen these constraints. Improved understanding of the underlying astrophysical electron spectrum would also allow for tighter constraints from high-latitude emission, while avoiding uncertainties associated with TeV Galactic Center emission, which is highly profile-dependent and would not apply here if substructure is depleted near the center of the Galaxy.

6.8 Inverse-Compton Gamma Rays

Absent a means of containing them, high-energy electrons resulting from annihilations will escape subhalos without difficulty. However, it has longer been understood that “in this world nothing can be said to be certain, except death and taxes!” [379]. In the high-energy Universe, taxes take the form of energy losses through a variety of
agents [42, 380], which will ultimately drive these particles to their demise. Far from the Galactic disk, the most important loss channel is inverse-Compton scattering on the CMB (we neglect the cosmic IR background, which has energy density a few percent that of the CMB [381]), since the magnetic fields there should be small [382, 383] and result in negligible synchrotron losses. To calculate the gamma-ray flux, we

Figure 6.6: Isotropic gamma-ray signals resulting from dark matter annihilations in substructure (assuming an Aquarius number density profile). Left side: Inverse-Compton gamma-ray emission of the final state electron/positron population from annihilations at distances > 20 kpc from the Galactic Center (solid line) and including all radii (dotted). This can be compared to COMPTEL [110], EGRET [114], and Fermi [389] diffuse gamma-ray data. Right side: Internal bremsstrahlung associated with the birth of charged leptons is shown for annihilation to two muons ($m_{\text{DM}} = 1.6 \text{ TeV}$; dot-dashed), two taus ($m_{\text{DM}} = 4 \text{ TeV}$; dark dashed), and two $\tau^\pm$ pairs ($m_{\text{DM}} = 8 \text{ TeV}$; double-dot dashed). For the two tau case, we show the effect of including tau decays (light dashed). Cosmic-ray $e^- + e^+$ measurements from HESS (triangles) [49, 50] act as upper limits on an isotropic gamma-ray flux (see text) [118].
must first find the equilibrium $e^- + e^+$ spectrum. We start from the diffusion-loss equation for a spectrum of relativistic electrons, $n_e(E)$ [42, 317]

$$\frac{dn_e}{dt} = D(E) \nabla^2 n_e(E) + \frac{d}{dE} [b(E)n_e(E)] + Q(E), \quad (6.8.1)$$

where the diffusion coefficient, $D$, is assumed to be isotropic, $Q(E)$ is the source term, and $b(E) = b_0 E^2$ is the radiative loss term, with $b_0 \simeq 0.3 \times 10^{-16} \text{GeV}^{-1} \text{s}^{-1}$ for the CMB (in the Thomson limit). For dark matter, equilibrium can be assumed.

In an isotropic system, there is no dependence upon $D$, since particle losses are compensated for by gains. At $\sim 1 \text{ TeV}$, the electron cooling time is $\sim 10^6 \text{ yr}$, so that even if electrons propagate rectilinearly, they would only travel a distance of order the virial radius of the Milky Way. This is likely an overestimate, since their propagation should be affected by the halo magnetic field, although its structure and strength is uncertain. Considering the length scales relevant for electrons injected by annihilation in substructure, we make the simplifying assumption that this halo magnetic field results in the IC losses occurring near the injection point (more care is needed for smooth halo signals due to the steeper gradient in particle injection with radius [384]). This reduces the problem to a continuity equation [313]

$$- \frac{d}{dE} [b_0 E^2 n_e(E)] = Q(E), \quad (6.8.2)$$

which can be readily solved for a given injection spectrum. While IB signals may vary greatly between annihilation channels, essentially all models that remain viable post-\textit{Fermi} lead to nearly identical equilibrium electron spectra (up to uncertainties in the astrophysical spectrum and propagation models) [375, 374]. With generality, we consider dark matter with $m_{DM} = 2.35 \text{ TeV}$ annihilating into two $\mu^\pm$ pairs (as in [375]). The calculation proceeds with

$$\frac{d\Phi_{IC}}{dE} = \frac{1}{4\pi} \mathcal{J}(\psi) r_\odot \Gamma_\odot \frac{dN_{IC}}{dE}, \quad (6.8.3)$$
where $\Gamma_\odot$ is the local annihilation rate per volume (matched to that required to agree with PAMELA/Fermi) and the IC gamma-ray spectrum per annihilation, $dN_{IC}/dE$, is calculated using the methods of Ref. [313]. In the inner Galaxy, synchrotron and IC losses on optical/IR photon backgrounds would result in a broad range of secondary photons [385]. We thus consider the signal resulting from annihilations occurring beyond 20 kpc from the Galactic Center, where IC on the CMB can be safely assumed to be the dominant energy loss mechanism (based on modelling of the Galactic optical/IR photon field [386]). Using $J(180^\circ)$ for $r > 20$ kpc yields the solid line in Fig. 6.6, which can be seen from the bottom panel of Fig. 6.5 to be nearly isotropic. Naively including radii interior to 20 kpc would result in the dotted line.

For the scenario considered here, the IC spectrum happens to peak at a similar energy to the pionic spectrum from cosmic-ray interactions [387]. We note that this IC signal retains less angular information concerning substructure [388] than direct gamma rays (such as IB). In comparing to isotropic gamma-ray data [110, 114, 389], we have made no attempt to account for other astrophysical contributions.

The velocity-dependence of the annihilation cross-section in Sommerfeld models makes calculating the cosmic signal in this scenario more complicated than in the standard picture [390]. This requires moving beyond the assumption of a constant boost due to a dependence of the velocity dispersion and hence the cross-section on halo mass. Also, the effects of baryons on dark halos vary with mass, since low-mass halos likely were never able to retain gas to form stars, and even in halos containing dwarf galaxies dark matter governs dynamics in the inner regions. Although a detailed treatment of these matters is beyond our scope, a simple estimate based on the total annihilation rate within the Milky Way halo (and scaling the amount of mass within substructure with host halo mass) suggests that this could be a factor of a few larger than the above IC flux with a similar spectral shape.
6.9 Discussion and Conclusions

We have shown that the gamma-ray line emission limits from the Galactic Center region, along the isotropic diffuse photon background, allow for stringent constraints to be placed on late decaying dark matter scenarios that produce monoenergetic photons. We emphasize that the Galactic and cosmic constraints are not independent of each other, with the GC region providing stronger limits in its range of applicability due to new spectroscopic data. Rather than attempting to explain various gamma-ray phenomena with a specific model, we report model-independent bounds on the decaying dark matter parameter space (as defined by $m_\chi \tau$ versus $\varepsilon$). Our general constraints are applicable to a number of models, and can be used as a guide for future model building. Upcoming gamma-ray telescopes with improved energy resolution, such as *Fermi* or ACT [391], can improve upon these bounds, particularly by making use of the unique spectral shape and directionality of decays from the Galactic halo.

One interesting application of our study is to assess the recent suggestion that cosmological late dark matter decays can explain the isotropic diffuse photon background in the 1-5 MeV range, whose origin remains a mystery [111]. We plot $m_\chi \tau$ versus $\varepsilon$ in Fig. 6.4 for $m_\chi/b \simeq 300$ GeV and $m_\chi/b \simeq 50$ GeV to represent the aforementioned mUED and SUSY models of Ref. [111], respectively. We also show the range of parameters, $m_\chi \tau$ versus $\varepsilon$, that can lead to a substantial ($> 10\%$) contribution to the iDPB (shaded diagonal band) or overproduce them (triangular region) through the sum of the local decays (Galactic halo) or decays from truly cosmological sources (all distant dark matter halos). The region excluded by the gamma-ray line emission limits from the GC region is below the jagged line. As seen in the figure, even the combined emission from the Galactic halo and cosmological sources due to either the mUED or SUSY inspired decaying dark matter models of Ref. [111] cannot make a significant contribution to the iDPB while abiding by the gamma-ray line
emission limits from the GC region. The mUED model can contribute to the iDPB only for $\Delta m \lesssim 1.5$ MeV with a contribution of $\lesssim 5\%$, while the SUSY model is even more severely constrained. Even relaxing our assumptions on the distribution of dark matter in the halo does not increase these fractions dramatically, thus, dark matter cannot decay in the late universe at a high enough rate to make a prominent contribution to the iDPB in the MeV range.

If one takes the position that astrophysical resolutions of the positron excess are untenable, then to have a viable dark matter scenario requires invoking a relatively-large annihilation cross section. One possibility for achieving this goal is a velocity-dependent cross section due to the presence of a new medium-range force resulting in a Sommerfeld enhancement. Here, we have examined the observational consequences of annihilation to charged particles in the context of a halo populated with dark matter substructure.

In determining the dark matter annihilation signatures of Galactic substructure, it must be kept in mind that the microphysics (annihilation) occurs within kinematically distinct subhalos so that their macroscopic distribution sets many aspects of the problem. Models in which annihilations in substructure, rather than in the smooth halo, are the dominant source of the locally-measured lepton flux imply associated IB and IC gamma-ray emission at high latitudes at a level accessible to (and possibly already in tension with) current observations. We have also argued that HESS TeV electron measurements can be regarded as limits on the isotropic TeV gamma rays arising from IB. Importantly, these new prospective signals can be tested with upcoming gamma-ray observations by Fermi and air Čerenkov telescopes.
CHAPTER 7
CONCLUSIONS

The following also was nobly spoken by someone or other, for it is doubtful who the author was; they asked him what was the object of all this study applied to an art that would reach but a very few. He replied: “I am content with few, content with one, content with none at all.”

Seneca, Epistle VII

With a number of long-planned experiments and telescopes having recently come to fruition, and with more on the way, the present era is an exciting time in the field of particle astrophysics. From ground-based telescope arrays to see showers produced by gamma rays and cosmic rays, to satellites that directly detect both gamma rays and charged particles, to the vast detectors needed to study neutrinos and ultrahigh-energy cosmic rays, to recent systematic surveys to find supernovae and their progenitors, these have begun to provide the data necessary for answering some of the major outstanding questions in physics and astronomy. To take full advantage of the opportunities afforded requires spanning a broad range of fields and interactions between traditionally-separated theorists, experimentalists, and astronomers. The research presented here has attempted in some small way to follow the path laid out over the course of the preceding half century [394, 395, 396, 397] to accomplish these and other goals.
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