Research in Theoretical Physics

Final Report¹

DOE/ER/40211--9

G. Domokos and S. Kovesi-Domokos
The Henry A. Rowland Department of Physics and Astronomy
The Johns Hopkins University
Baltimore, MD 21218

June 1998

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¹Report Number: DE-FG-85ER402211-9

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Abstract This report summarizes the research carried out under Grant # DE-FG-85ER402211. The main topics covered are: astroparticle physics at very high and ultrahigh energies; search for new physics by means of detectors of ultrahigh energy particles of extraterrestrial origin. Methods for searching in heavy quark decays for signatures of physics beyond the standard model are developed.

1 Introduction

Astroparticle physics gained renewed prominence in connection with searches for physics beyond the the standard model of elementary particle interactions. While the standard model appears to be in a very good agreement with current experimental results, it is clearly incomplete as a theory. Despite its predictive power, it contains a large number of input parameters (nineteen). In addition, the theory of electroweak symmetry breaking is, at best, incomplete; at worst, it is internally inconsistent. Hence, from the point of view of further developmernt of particle physics, it is vital to develop methods for discovering physics beyond the standard model.

2 Astro Particle Physics

There are, in essence, two reasons for the revival of the interest of high energy physicists in some branches of cosmic ray physics. Historically, the first one came with the publication of the paper by the Kiel group [1] in which they describe their observation of CYG X-3 at ultrahigh energies (UHE); they found an anomalously large number of muons in the extensive air showers (EAS) associated with that point source. At the time, it was generally assumed that the only particles which initiate the air showers associable with a point source are photons. According to the Standard Model, photons generate muon poor showers; hence too many muons mean either "*new physics*" in photoproduction or that the primaries are particles other than photons. The Kiel results have been confirmed by other observations; for a review, see G. Yodh's talk at the Seventh International Symposium on Very High Energy Cosmic Ray Interactions [2]. From the theoretical point of view, the puzzle remains, in essence, unresolved.

The second reason arose from the realization that experimentation with high energy lepton and photon beams is severely limited in energy if one relies entirely on accelerator based physics. At present, HERA operates at $\sqrt{s} \approx 30$ GeV. Photon energies extracted from an electron machine are less than that. Likewise, if one imagines that at one time a neutrino beam will be constructed at the future SSC by extracting and dumping one of the proton beams, the achievable energy

is $\sqrt{s} \lesssim 300 \text{ GeV}$ or so. Both energies are considerably lower than the expected characteristic energy at which “*new physics*” is expected to show up¹. By contrast, it is not unreasonable to expect neutrino energies from X-ray binaries of $\sqrt{s} \approx 10 \text{ TeV}$, see *e.g.* [3, 4] for reviews. Of course, the expected fluxes are much lower than those familiar from accelerator based experiments. Nevertheless, high energy, low flux experiments are likely to supplement the ones performed at lower energies at the various accelerators.

Due to the fact that particle beams from extraterrestrial sources contain less information about the beam itself than one is familiar with in the case of accelerator based particle physics, a part of the theoretical activity in this area is necessarily connected with the investigation of the properties of the primary beam and its origins. Activities of this type are necessarily interdisciplinary in their nature, requiring some knowledge both in particle physics and in astrophysics.

3 Methods in Data Analysis

All our efforts in this area were devoted to a better understanding of signals emitted by pulsed UHE sources. From the point of view of particle physics, these are the best sources of UHE photons and neutrinos, since the production mechanism is likely to be a hadronic one, see *e.g.* [4] and, due to the pulsed nature of the signal, the separation of the signal from the background can be reasonably efficiently accomplished. Even though only a few such sources have been so far observed, one hopes that as soon as a new generation of detectors will start observations, there will be a dramatic increase in the data sample. Our work anticipates this increase in the available data sample.

One of the most important questions connected with the observation of UHE signals from X-ray binaries is the nature of the primary particle.

There is a consensus emerging that, contrary to original expectations, the primaries giving rise to muon-rich EAS cannot be photons. For this reason, it is of importance to gather as much information about their nature as possible. As early as 1985, Halzen [5] described a simple formula to estimate the mass of the primary. It is based on a very simple argument: if a particle has a non-vanishing rest mass, its speed depends on its energy. Hence, if the source emits a continuum of energies, the effective time resolution of a detector far away from the source is deteriorated if the carriers of information (the primaries of EAS, for instance) are massive. The formula reads as follows.

$$\delta t = \frac{dm^2}{2} \left(\frac{1}{E_0^2} - \frac{1}{E_M^2} \right). \quad (1)$$

Here d stands for the distance between the source and the detector; E_0 and E_M are the minimal and maximal energies observed by the detector and m is the rest

¹Typically, one estimates that characteristic energy to be of the order of a few TeV.

mass of the carrier particle. Finally, δt is the shortest characteristic time observed by the detector (*e.g.* a pulsar period). In all practical cases, $E_M \gg E_0$; hence, the contribution of the second term on the right hand side of eq. (1) is negligible. For the same reason, we neglect terms of the order of E_0^2/E_M^2 .)

The validity of eq. (1) has been questioned by a number of authors, including Halzen and his collaborators, [6]. Due to the potential importance of the physics involved (are we seeing a breakdown of the Standard Model ? — see *e.g.* [3]), the question of placing reliable limits on the mass of the carrier particle is worth re-examining. Furthermore, some of the present, and certainly the majority of the future generation of detectors will be sufficiently sensitive to collect data at energies of 10^{15} TeV or more. At such energies, neutrons can reach terrestrial detectors from some of the nearby point sources, such as HER-X1 or CYG-X3. Consequently, one needs a method which is sufficiently sensitive even at these energies in order to distinguish between a hadronic and some other primary.

In essence, we reformulated the problem in analogy with the standard treatment of the propagation of a periodic signal in a dispersive transmission line. This provides us with a method which, unlike eq. (1), utilizes the complete information available from the observation of a periodic source. The method requires the separation of muon rich and muon poor showers in the data sample to be analyzed. Currently available data samples are too small to allow for such a separation in a meaningful way. However, should larger arrays, currently under construction or in the planning stage, observe point sources, the separation will be possible.

We consider a periodic source a distance d away, of fundamental period T and of an arbitrary pulse shape. Depending on the physical situation, the quantity T may be a pulsar period, an orbital period, *etc.* Implicitly, it is assumed that any change in the period takes place on a time scale much longer than the period itself, *i.e.* $dT/dt \ll 1$.

Under these assumptions, the emitted intensity at the source can be expanded into a Fourier series:

$$I(E, t_s) = \sum_{n=-\infty}^{\infty} I_n(E) \exp in\omega t_s, \quad (2)$$

where $\omega = 2\pi/T$ and t_s stands for the time measured at the source. Of course, the intensity is real, hence, $I_{-n} = I_n^*$. In most cases, the "DC component" of the intensity is either vanishing or it can be subtracted whenever necessary, hence, one can assume $I_0 = 0$ if needed.

At the detector, the observed intensity as a function of the time, t , is given by the expression:

$$I(E, t) = \sum_{n=-\infty}^{\infty} I_n(E) \exp in\omega \left(t - \frac{d}{v(E)} \right) \quad (3)$$

In eq. (3), the quantity in brackets is the retarded time, with $v(E)$ being the group velocity of the signal. For energies much larger than the rest mass, m , of

the carrier particle, we have:

$$\frac{1}{v(E)} \approx 1 + \frac{m^2}{2E^2}. \quad (4)$$

If the rest mass is different from zero, the group velocity is energy dependent; the various components suffer a *phase shift* which depends, in an obvious way, both on the energy of the carrier particle and on the order of the Fourier coefficient.

Knowing the intensity at the source and the time at which each particle of a given energy started, allows one to determine the energy dependence of the group velocity (i.e. the rest mass of the carrier particle) by measuring the relative phase shifts between the various Fourier components. This is in complete analogy with the way one can determine the dispersion of a transmission line. In an astrophysical situation, however, one cannot directly measure the signal at the source. Moreover, even with the arrival of a new generation of detectors, it is unrealistic to expect that the data sample will be sufficiently large in order to determine a significant number of Fourier coefficients as functions of the primary energy. What one can hope to do is to determine the time averaged energy spectrum, i.e. the energy dependence of the Fourier coefficient I_0 and the Fourier coefficients averaged over energy. (One notes that I_0 remains undistorted, no matter how big is the rest mass of the primary particle.)

We now argue that at ultrahigh energies, at least the first few Fourier coefficients at the source depend on the energy in the same way: hence, their energy dependence can be read off from that of I_0 . In fact, UHE particles must come typically from the first interaction of the high energy protons with the target. The reason is the following. Assuming the most favorable acceleration mechanism, the upper limit on the proton energy hitting a target is of the order of 10^{18} eV [7], corresponding to a C.M.S. energy of about 50 TeV in a nucleon-nucleon collision. An extrapolation of currently available accelerator based data [8] predicts an average hadron multiplicity, $\langle N_h \rangle \approx 100$ in this energy range. Hence, in a typical collision, the average energy per hadron decreases by two orders of magnitude. (Obviously, this also gives the rate at which the average γ and lepton (in particular, ν), energies degenerate.

As a consequence, if one is to observe events with primary energies of the order of a few hundred or a few thousand TeV, one cannot allow for more than about two collisions in the target star in a binary system (typically, a main sequence star). It follows that the pulse at these energies closely reflects the properties of the primary beam: in particular, its shape changes little with energy. Hence, the Fourier coefficients exhibit roughly the same energy dependence, viz. that of the primary beam. For the same reason, one also expects the pulse to be considerably narrower than at lower energies: only extreme glancing collisions with the target can produce UHE photons or neutrinos. This is apparently borne out by the data, see the next subsection in this report, strengthening one's belief in the basic correctness of the picture.

A consequence of the above reasoning is that for an UHE signal, the Fourier coefficients in eq. (31) approximately factorize:

$$I_n(E) \approx S(E) I_n, \quad (5)$$

where now I_n no longer depends on the energy. In order to simplify matters, we perform a time translation: we replace $t - d$ by t in eq. (3). If the source is, as assumed, stationary during the observation period, this is permissible. We assume as usual that the spectral intensity, $S(E)$ is given by a power law; correspondingly, the normalized distribution for a detector with threshold energy E_0 , is given by the expression:

$$S(E) = (\alpha - 1) \frac{1}{E_0} \left(\frac{E}{E_0} \right)^{-\alpha}. \quad (6)$$

Using eq. (5) and the shape of the spectrum, eq. (6), the Fourier coefficients at the detector, averaged over the spectrum, are of the form:

$$I_n^d := I_n D_n(E_0), \quad (7)$$

where the *distortion factor*, $D_n(E_0)$ is given by the expression:

$$D_n(E_0) = \int_{E_0}^{\infty} \frac{dE}{E_0} (\alpha - 1) \left(\frac{E}{E_0} \right)^{-\alpha} \exp \left(-inz \left(\frac{E}{E_0} \right)^2 \right). \quad (8)$$

The variable z is defined by:

$$z = \omega d \frac{m^2}{2E_0^2} \quad (9)$$

The Fourier coefficients at the detector are distorted by a factor due to the velocity dispersion. As a consequence of the approximate factorization of the coefficients, eq. (5), each Fourier coefficient is multiplied by its own distortion factor.

The integral in eq. (8) can be expressed in terms of an incomplete Γ function of a pure imaginary argument. For our purposes, it is sufficient to give its expression for $z \ll 1$; this will be justified *a posteriori*. The distortion factor is a complex quantity; its phase and absolute value can be expanded in a Taylor series in powers of z . We give here the first few terms of the series expansion:

$$\begin{aligned} |D_n| &= 1 - \frac{n^2 z^2}{2} K(\alpha) + \dots; \\ \arg D_n &= nz \frac{\alpha-1}{\alpha+1} + \dots \end{aligned} \quad (10)$$

Here $K(\alpha)$ stands for the quantity

$$K(\alpha) = \frac{4(\alpha-1)}{(\alpha+3)(\alpha+1)^2}$$

This series is generated by means of expanding the exponential in equation (8) into a power series and integrating termwise. Unless the mass is assumed to

be unreasonably large, the variable z is quite small. To quote an example, let us assume that one observes a pulsar of fundamental period, $T = 0.1\text{sec}$, at a distance of 5 kpc with $E_0 = 100\text{TeV}$. Assuming a particle mass of 100 MeV (probably an exaggeration), one gets $z \approx 0.05$. Consequently, the first few terms of the expansion in eq. (10) are quite adequate.

It is worth remarking that the formulae for the absolute values of the Fourier coefficients are more useful from the point of view of extracting mass limits: in general, the phases of Fourier coefficients extracted from some observational data are more strongly affected by noise than the absolute values [43].

By looking at the expansion of the absolute value of the distortion factor, one can see that the velocity dispersion due to a finite rest mass of the carrier can only *broaden* the signal: higher Fourier coefficients are more strongly damped. (This can be shown in general. However, we are not going to give a general proof here, since the result is intuitively obvious.)

The formulae derived here allow us to estimate the rest mass of the carrier particle as follows.

We take ratios of the absolute values of the Fourier coefficients at the detector, averaged over the energy spectrum: this allows us to eliminate an overall normalization, which is detector dependent and carries a substantial uncertainty.

For the sake of definiteness, we consider the ratios of consecutive Fourier coefficients, although this is inessential and can be easily modified if necessary. We get:

$$\frac{m^2}{E_0^2} = \frac{2^{3/2}}{\omega d \sqrt{K(2n+1)}} \sqrt{\left| \frac{I_n^d}{I_{n+1}^d} \right|^2 - 1}. \quad (11)$$

In the last equation, I_n^d stands for the n^{th} Fourier coefficient observed at the detector, using eq. (5). (The estimate given by eq. (11) is accurate up to terms of $O(z^4)$.)

It is to be noted that K is a slowly varying function of its argument for "reasonable" values of the spectral index α , typically, $1.3 \lesssim \alpha \lesssim 2.5$. Hence, errors in the determination of the primary energies of the EAS associated with a given source do not influence the estimate of the rest mass substantially. As it can be seen from eq. (11), the only substantial uncertainty arises from that in determining the primary energy of the least energetic event. This is inevitable in any method of determining the mass of the primary from velocity dispersion.

The Fourier method systematizes the search for an upper limit on the mass of the UHE particles which cause extensive air showers. (Given a sufficiently large data sample, one may even hope to identify UHE neutron induced showers in the future.) The method is superior to previous estimates in that it utilizes all the information provided by the experimental data. (By contrast, eq. (1) utilizes, typically, two data points only.)

Knowledge of the energy averaged light curve at the source (deduced from muon poor showers) is essential in order to obtain a good determination of the mass of the primary. We believe that this is intuitively obvious: regardless of the

production mechanism, one expects the pulse to have a *finite width*. However, we saw a direct manifestation of this in a currently available data sample.

We tested the method here described on UHE data obtained during a burst of HER-X1 in 1986 [51]. In that reference, data on a total of 11 showers are published, making any separation of the showers according to their muon contents meaningless. Consequently, we made the assumption that the entire width of the signal originates from velocity dispersion. This hypothesis is identical to the one implicit in eq. (1); it means that one sets $I_{n+1}/I_n = 1$ in eq. (11). In this way, we obtained a slightly, but not dramatically lower mass limit than the one published in ref. [51], using eq. (1). Clearly, the natural width of the pulse is nonzero: without any knowledge of the latter, the upper limit on the mass becomes a very crude one.

4 Physics Results in Astro Particle Physics

We arrived at several distinct physics results in the area of particle physics. The first one is a modern recalculation of the galactic neutrino background. The second result is in an area where astroparticle physics is at its best: we place a limit on the deformation parameter of the Poincaré group by utilizing both the fact that UHE observations provide us with particle beams of energies not found elsewhere and the fact that the accelerator (in this case, an X-ray binary) is very far away. The third result concerns the search for “*new physics*” in EAS observations: we studied the development of “anomalous” air showers in order to obtain a picture of what are the physical observables to which current and planned measurements are sensitive. Finally, our last result is of an interdisciplinary nature; using the methods developed in the previous subsection, one sheds some light on the production mechanism of UHE particles in binary systems. We now proceed to describe these results.

4.1 The Galactic Neutrino Background

Galactic and extragalactic cosmic rays (mostly protons and perhaps heavier nuclei) interact with matter found in our galaxy. In those interactions a large number of mesons is produced, which, in turn, ultimately decay into charged leptons and neutrinos. The charged leptons interact electromagnetically with interstellar matter and magnetic fields, thus becoming an essentially random low energy leptonic background, practically unobservable. By contrast, the neutrinos interact weakly (even if at the relevant energies there is an increase of their interaction cross section, see [12]). Hence, they are likely to reach a neutrino telescope on Earth and induce events in it, as originally predicted by Markov [13]. Several estimates of this background exist, for a recent review see [14].

There are at least two reasons for recalculating the galactic neutrino background. First, all the works reviewed in ref. [14] are rather old ones. (The first

explicit estimate [15] was published in 1979). In the meantime, accelerator based data became available regarding inelastic cross sections and hadronic multiplicities, obviating the need for the extrapolation of the data over large energy ranges, as it was done in the older works. Second, in some of the works, the diffuse neutrino flux was expressed in terms of the gamma ray flux generated in a related process, typically using charge independence arguments. Again, the availability of high energy data permits a more explicit calculation. In view of the fact that there are several neutrino telescopes either currently taking data or in the planning stage, an explicit estimate of the neutrino background should be useful.

Our estimates differ from previous ones both in that in several respects they are more explicit than previous ones². and in that we use somewhat different assumptions about the elementary process. (In particular, we use as many accelerator based, high energy data available as practical.)

Due to the considerable uncertainties both in the behavior of hadronic cross sections beyond energies achievable by current accelerators and in our knowledge of the matter distribution in the Milky Way, we use a rather simple model in order to estimate neutrino production. In addition, the results may be influenced by the chemical composition of the primary cosmic radiation. Hence, our estimates are, by necessity, rather crude ones: we have to oversimplify many aspects of this rather complex process. Nevertheless, we believe that they serve a useful purpose in more than one way: first, they give a general idea about the kind of ν background to be expected in a neutrino detector; second, they exhibit some of the uncertainties associated with estimates of this kind. As a consequence, one hopes that as further data become available, the estimates given here will be improved upon.

It is now a generally accepted fact that the Milky Way (and, most probably, any galaxy) has to contain a certain amount of dark matter, not residing in stars. The exact nature of the dark matter is uncertain; however, some of it may be baryonic. If so, neutrinos are produced on it, hence, baryonic dark matter can be observed by the distortion of the angular distribution of the neutrino background.

In order to simplify matters, we assume that particle production in UHE interactions is one dimensional. This is a reasonably good assumption: at the energies of interest ($E_L \geq 100\text{TeV}$, say,) the angular spread in the laboratory is a fraction of a milliradian. We also neglect certain details of the strong interaction, such as leading particle effects *etc.* . (This is an acceptable approximation for the present purposes, since it introduces distortions in the resulting distributions of the order of 10 % or so.) If a primary proton of energy E_L hits a target and produces N secondaries, on the average each one of the latter has an energy E_L/N in the laboratory frame of reference. Further, we use simple analytic fits to the charged particle multiplicity, the total pp cross section and the energy distribution of the primary cosmic radiation.

We use the following fits to the charged multiplicity (N_{ch}) and total cross

²For instance, previous calculations have not given explicit expressions and/or curves for the angular distribution.

section (σ):

$$N_{ch} = 0.88 + 0.44 \ln s + 0.118(\ln s)^2, \quad (12)$$

$$\sigma = 62.4 - 7.68 \ln s + 0.66(\ln s)^2. \quad (13)$$

In both equations (12) and (13), s stands for the energy in the CMS squared, measured in $(\text{GeV})^2$; the cross section is measured in millibarns. Equation (12) is taken from the Particle Data Book [16], whereas equation (13) is a compromise between the fits of Block and Cahn [17] and that found in the Particle Data Book, omitting terms from both fits which decrease with s . Finally, we use a rough fit to the flux of the primary cosmic radiation with just two spectral indices. The fit was made to a variety of data, see, e.g. [18]. In essence, however, there is an agreement between the various measurements, thus the fit is not very sensitive to the sample of the data used. We express the flux in terms of the CMS energy squared in a nucleon - nucleon collision.

$$\frac{d\Phi}{ds} = \Phi_0 \left[s^{-\gamma_1} \Theta(s_1 - s) + s_1^{\gamma_2 - \gamma_1} s^{-\gamma_2} \Theta(s - s_1) \right] \quad (14)$$

The form given above gives a fair representation of the flux with the following choice of the parameters:

$$\Phi_0 = 36.4 \text{cm}^{-2} \text{sec}^{-1}, \gamma_1 = 2.6, \gamma_2 = 3.0, s_1 = 10^{6.8} \text{GeV}^2$$

The reader will notice that we did not include the flattening of the primary spectrum above $s \gtrsim 10^{18} \text{GeV}^2$. However, this causes a rather small error in our estimates, due to the smallness of the flux there. We assume the flux of primary cosmic radiation to be isotropic. This is well justified for energies below a few PeV in the laboratory; above that value there is some evidence for an anisotropy. However, the uncertainties existing in other input data (e.g. the matter distribution in the Milky Way) are comparable with any measured anisotropy. Thus, we feel justified in taking an isotropic primary cosmic radiation in the whole energy range considered.

In order to compute the production rate of UHE neutrinos, we need the distribution of baryonic matter in the Milky Way. Let us denote the baryon density in the galaxy by δ . Obviously, δ is a function of the position. However, we argue that for our purposes, it can be replaced by its average value, $\delta \approx 1 \text{cm}^{-3}$. The argument is the following.

We are interested in calculating the production of UHE neutrinos only. Apart from a relatively slowly varying background density arising from the presence of non-stellar material in our galaxy, the rapid, small scale variations of δ arise from the presence of stars in the Milky Way. Stars, however, are *not* good sources of UHE neutrinos, due to the fact that the hadrons produced in them interact several times before they decay. Therefore, the majority of the neutrinos emerging from the interaction of the primary cosmic radiation with a stellar object tend to be of low energy, unlikely to be observable by means of a neutrino telescope.

In order to get some quantitative estimates, we take a simple exponential model of the density distribution in a star, *viz.*

$$\rho(r) = \rho_0 e^{-r/h}, \quad (15)$$

where h stands for the scale height of the star and ρ_0 is the central density. The total mass of the star, of course, is given by

$$M = 4\pi \int_0^\infty \rho(r) r^2 dr. \quad (16)$$

The amount of stellar material "seen" by a particle incident upon the star at impact parameter b is given by the expression,

$$m(b) = \int_{-\infty}^\infty \rho \left((b^2 + x^2)^{1/2} \right) dx. \quad (17)$$

For an exponential distribution given by eq.(15), the expression in eq.(17) can be evaluated analytically. One gets:

$$m(b) = 2b\rho_0 K_1 \left(\frac{b}{h} \right), \quad (18)$$

where $K_1(x)$ is a standard MacDonald function. We further recall that — neglecting nuclear effects — the mean free path (mfp) of a reaction mediated by short range forces is given by a formula *approximately independent of the atomic number of the target*, (cf. [19]):

$$\lambda_i \approx \frac{1670}{\sigma}, \quad (19)$$

where λ is measured in g/cm² and σ in millibarns.

Before proceeding further, a brief remark is in order concerning equation (19). Strictly speaking, its right hand side would be independent of A , the atomic number of the target, in the case of a *small* elementary cross section σ , since in that case the nuclear cross section is proportional to A . According to theoretical calculations, cf. [20] the nuclear cross section (for light and medium heavy nuclei) is, to a good approximation, proportional to $A^{0.8}$. (This is in agreement with experimental data where they are available, see [21].) Hence, the right hand side of eq. (19) is proportional to $A^{0.2}$. For purposes of this estimate, such a weak dependence on A is negligible.

Somewhat arbitrarily, we now consider those impact parameters relevant from the point of view of UHE neutrino production where the amount of stellar matter penetrated by the primary is less than $3\lambda_i$. For smaller impact parameters cascading is substantial and the primary energy degenerates. Using eq.(18), we get that this requires

$$b \gtrsim 23.5h.$$

This value of the impact parameter was arrived at by assuming $\rho_0 \approx 150\text{g/cm}^3$ and $h \approx 8 \times 10^9\text{cm}$. (These parameters correspond to a Sun-like star; the average star in the Milky Way being about half that size, we actually overestimate the contribution of stellar matter to UHE neutrino production.) With this, one can easily estimate the amount of stellar matter contributing to UHE neutrino production. We integrate the density given by (15) outward from the minimal impact parameter quoted above and multiply the result by the number of stars in the Milky Way ($N_{\text{star}} \approx 10^{11}$). As a result, we get that the contribution of stellar matter — smeared over the entire Milky Way — corresponds to a density of about $10^{-6}\text{baryon/cm}^3$, *much smaller than the average density of interstellar matter*, which is about 1baryon/cm^3 .

This concludes our argument: *stars, in essence, do not contribute to the production of UHE neutrinos in proportion to their masses.* By the same token, we do not have to worry very much about the central bulge in the Milky Way either, [22]. The main contribution to the UHE neutrino production comes from the interaction of the primary cosmic radiation with non-stellar material (H, CO, dust, ...) present in the Milky Way: that component, in turn, has a remarkably uniform distribution out to the radius of the galactic disk, cf. [23], Ch. 3.

Bearing all this in mind, the angular distribution of UHE neutrinos can be fairly reliably computed by taking a rather simple-minded model of the Milky Way. We take a uniform matter distribution packed into an oblate rotational ellipsoid with semimajor axis $a \approx 15\text{kpc}$ and semiminor axis $b \approx 0.75\text{kpc}$. We further approximate the location of our solar system by placing it in the galactic plane at a distance of 10kpc from the galactic center.

We can now proceed to compute the angular distribution of neutrinos, assuming a constant baryon density of $1/\text{cm}^3$. Assuming that all hadrons produced decay rather than interact simplifies the calculation; however, such an assumption has to be justified. This is done later, now we proceed under this assumption.

The rate of charged hadron production in a volume element d^3x is given by the expression:

$$\frac{dN_h}{ds} = \frac{d\Phi}{ds} \sigma N_{ch} d^3x \delta. \quad (20)$$

Let further $\mathbf{x}$ stand for the vector pointing from the detector to the point of production under consideration. Assuming that all charged hadrons produced will decay into neutrinos, (*i.e.* $dN_h \approx dN_\nu$), the observed neutrino intensity at the detector is given by the formula:

$$\frac{dN_\nu}{dE_\nu} = \frac{1}{4\pi} \int_V d^3x \frac{dN_h}{ds} \frac{\cos \omega}{x^2} \frac{ds}{dE_\nu} \epsilon(\cos \omega) \quad (21)$$

Here, $\cos \omega$ stands for the angle between the normal of the plane of the detector and the vector $\mathbf{x}$; the integral is to be taken over the galactic volume. The function ϵ is the angular efficiency of the detector around the given direction. In what follows, we replace ϵ by a δ function and plot the intensity *per steradian* of solid angle.

In equating the number of hadrons with the number of neutrinos, we neglected muon neutrinos emerging from muon decay and electron neutrinos. Such an approximation introduces only a rather small error at the relevant energies. Without a detailed solution of the transport problem, one can estimate the fractional error thus committed as follows. We notice that at high energies in the laboratory, the energy of a decaying particle is, on the average, shared equally between the decay products. Similarly, neglecting leading particle effects in the hadronic interaction, the average energy of a secondary hadron in the laboratory is given in terms of the primary energy, E_p , by $E_h \approx E_p/(1.5N_{ch})$, since most of the secondaries are pions. As a consequence, given the energy of a muon neutrino, one needs a primary energy 3 times as high to produce it from muon decay rather than from the decay of a π or K meson. Therefore, the correction to the integral spectrum of UHE muon neutrinos, taking into account the ones coming from muon decay, is of the order of $3^{1-\gamma}$, where γ is some average spectral index of the primary spectrum. For spectral indices comparable to the fit given here, the correction is of the order of 40 %. In the approximation used in this paper, one can simply multiply all the distributions given by a factor of $1 + 3^{1-\gamma}$ in order to take into account the contribution of the muon neutrinos emerging from muon decay. Likewise, the distribution of electron neutrinos is obtained if one multiplies the distributions given by the factor $3^{1-\gamma}$.

The computation of the quantity ds/dE_ν requires some knowledge of the production and decay mechanisms. As stated earlier, we use a rather crude, but sufficiently representative model. We assume that if E_L is the primary energy of the incident primary proton in the rest frame of the detector, then each secondary hadron produced has an energy $E_h = E_L/N(s)$. Further, we assume that the energy of the emerging neutrino is given by $E_\nu = E_h/2$. Due to the fact that in the approximation we are using, the primary energy is distributed equally among the secondaries, we have

$$E_\nu \approx \frac{s}{9mN_{ch}(s)},$$

so that

$$dE_\nu = \frac{ds}{9mN_{ch}} \left(1 - \frac{d \ln N_{ch}}{d \ln s} \right) \quad (22)$$

We assumed that approximately 1/3 of the secondaries is neutral: the neutrals share the primary energy, but they are not an important source of neutrinos. The quantity m is some average hadron mass. One can, to a reasonable approximation, set it somewhere between the pion and kaon masses since those are the most copiously produced hadrons. The logarithmic derivative of the multiplicity in eq.(22) is a slowly varying quantity and above $\sqrt{s} \gtrsim 20\text{TeV}$ its presence amounts to a correction $\lesssim 10\%$.

We now have to justify our technical assumption, namely that in the interstellar space most of the hadrons produced will decay rather than interact. We do this by comparing the mfp for decay and interaction. The former has a more

rapid variation with energy than the latter:

$$\lambda_d = \frac{E_h \tau}{m_h}, \quad (23)$$

where m_h and τ stand for the mass and proper lifetime of the decaying hadron, respectively. (We use natural units, $\hbar = c = 1$.) By contrast, the variation of the interaction mfp with energy is logarithmical, which we neglect for the present purpose. As long as λ_d is smaller than the interaction mfp, the transport process is dominated by the decay of hadrons. Above a critical energy, defined by $\lambda_d(E_c) = \lambda_i$, interactions dominate over decays. The interaction mfp is given by $\lambda_i = 1/\delta\sigma$. Inserting $\delta \approx 1\text{cm}^{-3}$, and, for the sake of simplicity, a constant cross section, $\sigma \approx 100\text{mb}$, one gets $E_c \approx 2 \times 10^{21}\text{GeV}$ for pions. Hence, to a very good approximation, one may indeed neglect the interaction of the secondaries: all particles produced will decay.

For practical purposes, one is interested in the angular distribution of neutrinos above a certain energy (the threshold of the detector). The angular distribution depends on two angles. One of those was chosen an azimuthal angle, ϕ , in the galactic plane. The line of sight of the galactic center from the solar system is defined to have $\phi = 0$. The other angle, ψ , is, in essence a galactic latitude: it was chosen to be the angle between the line of sight of the detector and the galactic plane.

Figure 1 shows samples of the predicted angular distribution at various azimuthal angles, for two threshold energies, $E_{th} = 20\text{TeV}$ and $E_{th} = 100\text{TeV}$ in the laboratory system. The first of these thresholds corresponds, approximately, to the energy above which a DUMAND type underwater detector can certainly observe neutrinos with a high efficiency, cf.[24]; the second one corresponds to the approximate threshold of an EAS array at sea level. (The stated threshold of DUMAND is lower; we chose 20 TeV as an example in the region where the detector is highly efficient.) As it turns out, the effect of the logarithmic variation of the cross section and multiplicity on the distributions is not a very significant one. In a crude approximation, one can just scale up the result at 100 TeV by a factor of 23 in order to get the curves at a threshold of 20 TeV. In fact, neglecting the logarithmic dependence of the physical quantities, one can give a crude estimate of the flux above a threshold energy E_{th} by the approximate scaling law

$$\frac{N_\nu(> E_{th})}{N_\nu(> 100\text{TeV})} \approx \left(\frac{E_{th}}{100\text{TeV}} \right)^{-\gamma}, \quad (24)$$

where $\gamma \approx 1.8$ from a crude, single-exponent fit to the integral spectrum corresponding to eq.(14) over the entire energy range covered. Within the framework of the approximation scheme used in this work, eq.(24) holds at any angle of observation.

A final remark is in order here. Our results depend, to some extent, on the chemical composition of the primary cosmic radiation. This is due to the fact that

any measurement of the primary spectrum above a few hundred TeV measures the total energy per incident *particle*, not per incident nucleon. Therefore, if above some energy elements heavier than, say, H and He dominate, the energy per nucleon will be lower but the flux of nucleons, accordingly, higher. (This is a well known phenomenon. A qualitative discussion of its effect on galactic neutrino production is given in ref. [14], Ch. VIII.) We tested the effect of the chemical composition on our results by assuming that above $E_L = 1\text{PeV}$ the entire primary cosmic radiation consists of Fe nuclei. (This is a caricature of some recent observational data, [25], suitable for testing the effect of the chemical composition on the angular and energy distributions.) The neutrino spectrum is displayed in Fig. 2 for a threshold energy of 100 TeV. In Fig. 2a we assume that protons dominate over all the relevant energy range, whereas in Fig. 2b the same spectrum is displayed under the assumption of Fe dominance above 1PeV. The influence of these, physically rather different assumptions, on the angular distribution is, however, negligible. *Neutrino telescopes are not good tools in helping to resolve the question of chemical composition of the primary cosmic radiation.*

We now turn to describing the estimate of the effect of baryonic halo. It is intuitively obvious that if there exists baryonic matter in the halo, the angular distribution of background neutrinos will drop slower with increasing latitude than in the absence of such matter. The quantitative details depend on the specific assumptions made. Here we illustrate the behavior of the angular distribution under one specific assumption about the spectrum of the primary radiation. For the sake of simplicity, we assume that baryonic dark matter in the halo is distributed spherically, it has the same mass as the galaxy itself and the mass distribution is roughly uniform out to the galactic radius. Further, we assume a primary cosmic ray spectrum outside the galactic disk corresponding to the high energy part of the formula given in eq.(14). This corresponds to a two - source model of the primary cosmic radiation, see e.g. [26] and to the assumption that magnetic confinement of the primary cosmic radiation extends only to the galactic disk. With these — admittedly somewhat *ad hoc* — assumptions, we recalculated the angular distribution of the produced neutrinos. Figure 3 shows the contribution of the halo alone under the assumptions just stated. Figure 4 displays the total contribution (halo and disk) at the same azimuthal angles as in Fig.1. For comparison, we also plotted the angular distribution in the absence of a baryonic halo (the curves displayed in Fig. 1). It is clear that “looking” in directions near perpendicular to the galactic plane provides us with a good possibility of detecting a baryonic halo: the flux is both significantly higher and its distribution is noticeably flatter than in the absence of a baryonic halo.

Variations of these assumptions would produce results which differ from the ones shown here in quantitative aspects. Before engaging, however, into a more detailed study of this, it would be interesting to see whether such an effect exists at all. The *absence* of any substantial amount of baryonic dark matter in the halo is characterized by a rapid falloff of the neutrino intensity with increasing

latitude.

The galactic neutrino flux thus computed should be compared with some other sources of diffuse neutrino radiation. We just quote two of those; it is reasonable to expect that they are the main "competitors" to the galactic neutrino radiation and, hence, also to the search for either a baryonic halo or point sources.

- Atmospheric neutrinos. A rough extrapolation of the results of ref. [28] to the UHE region is in qualitative agreement with the calculations of Volkova [29]. However, similarly to the estimates of galactic neutrino production mentioned in the Introduction, Volkova's results are based on the extrapolation of old data to high energies. Therefore, all the available estimates of the flux of UHE atmospheric neutrinos are highly uncertain and it would be desirable to update them. It is probably reasonable to assume that above laboratory energies of the order of a few hundred TeV or so, the atmospheric neutrinos cease to dominate the observations.
- Active galactic nuclei (AGN). In a recent paper, Stecker *et al.*, ref. [30] gave an estimate of the diffuse neutrino radiation generated by AGN. According to their results, the background due to AGN dominates the neutrino sky above, approximately, $E_L = 100\text{TeV}$ or so and it is falling rather slowly with energy. It has to be pointed out, however, that one's lack of knowledge of details of the UHE processes taking place in AGN makes such estimates somewhat unreliable. For instance, the estimates of Sikora and Begelman [31] fall several orders of magnitude below the ones arrived at by Stecker *et al.* for energies higher than a few hundred TeV. It has also been emphasized that some presently existing observational results may already place a rather stringent upper limit on the neutrino radiation emitted by AGN, cf. Bergström *et al.*, ref. [32].

One has to conclude that the theoretical uncertainties in estimating the fluxes of the competing processes are still rather large. Probably a combination of further observations and theoretical work is needed in order to better clarify the situation.

In the absence of a substantial baryonic halo, one has the best chance of discovering distant neutrino point sources by looking in a direction approximately perpendicular to the galactic plane. In other directions (e.g. in the direction of the galactic center), the chance of discovery depends upon specific properties of the source. For instance, if the source is known to be a pulsed one, Fourier analyzing the data (or simply using Protheroe statistics) may enhance the signal to noise ratio.

Let us, finally, comment on the reliability of the calculation presented here. We have made the seemingly drastic simplification that the energy is distributed uniformly among all the secondaries. There is, however, some recent theoretical evidence that at ultrahigh energies this may be a general phenomenon, [33]. However, even if leading particle effects persist, a *few* high energy secondaries will not contribute significantly to the overall flux; consequently, the theoretical error on

our calculated flux is estimated to be about a factor of 2 or less, due to particle physics uncertainties. The same holds for uncertainties regarding the density of interstellar matter and its distribution. (For instance, we used a simplified target matter distribution, neglecting the presence of spiral arms, etc.) Overall, we estimate that our results are, probably, accurate to an order of magnitude. While this appears to be a more pessimistic estimate than the ones often given by the authors of previous computations (see [14] and references quoted there), in many of those papers the uncertainties involved in the extrapolation of particle physics data has not been taken into account. Assuming for the moment that our estimates can be taken at face value, one concludes that the galactic neutrino background is weak and, with the possible exception of very distant point sources, it does not represent a serious obstacle to the observation of point sources. By the same token, although the method described here appears to be the only known way of directly observing a baryonic halo around the Milky Way, such observations appear to be rather difficult ones.

4.2 Limits on the deformation of the Poincaré group

The special theory of relativity is one of the fundamental ingredients of any calculation or experiment in particle physics. Therefore, it is important to establish the accuracy to which it can be trusted.

This is a non-trivial problem. Practically all analyses of current particle physics experiments assume the validity of the special theory of relativity, *i.e.* the invariance of *local* physics under the Poincaré group: hence, by design, they cannot be used to test Poincaré invariance, unless, perhaps, that invariance fails in a dramatic and truly unexpected way. As a consequence, while one knows that the *special* theory of relativity is not applicable on length scales where gravitational effects are important, its validity at short distances is harder to test.

To be sure, it has been proposed several times that Poincaré invariance is violated at short distances; for some recent proposals see [34]. However, most of the proposals were introduced in an *ad hoc* fashion; the internal consistency of the schemes has never been investigated systematically. (There also exists a substantial number of older works in this area, involving a variety of ideas, such as the discretization of space-time, *etc.* ; it is not our purpose here to review them here.)

Due to developments in the theory of group deformations, the situation has changed. In some recent papers, [35], Lukierski *et. al.* developed a theory of the deformation of the Poincaré group. In order to accomplish the task, the authors of ref. [35] started from the DeSitter group. They performed a group deformation along the lines described in the classic papers on the deformation of simple groups, [36], [37], [38] and afterwards they performed a contraction to a deformed Poincaré group. While the procedure referred to is not a unique one, (see, *e.g.* [39]) it is internally consistent. Moreover, it is attractive from the physical point of view: the theory contains a characteristic mass, κ . The latter

determines the energies below which the special theory of relativity is valid to a good approximation; at energies above κ , significant deviations from the familiar relativistic kinematics are to be expected.

Existing and future observational data on UHE, distant events can place an interesting lower limit on the characteristic mass, κ . For this purpose one needs, in essence, only one of the results developed in refs. ([35]). In those papers it was shown that the quadratic Casimir invariant of the Poincaré group is deformed to the following expression:

$$C_\kappa = -\left(\mathbf{p}^2 + 2\kappa^2 \left(\cos \frac{E}{\kappa} - 1\right)\right). \quad (25)$$

In eq. (25) the quantities $\mathbf{p}$ and E have their usual meaning. The deformed Casimir invariant, eq. (25) goes over to the familiar quadratic form as $\kappa \rightarrow \infty$. However, at any $\kappa < \infty$, there is an observable difference between ordinary relativistic kinematics and the one described by eq. (25). Let us consider, in particular, a single particle state: in that case, $C_\kappa = m^2$, where m is related to the rest mass of the particle. On putting $\mathbf{p} = 0$, the rest energy is given by:

$$E^0 = \kappa \cos^{-1} \left(1 - \frac{m^2}{2\kappa^2}\right), \quad (26)$$

which agrees with the usual expression up to corrections of $O(m^4/\kappa^4)$. One can compute the group velocity (v) of the particle in the deformed Poincaré group. As within the framework of the usual (undeformed) Poincaré group, if the energy of the particle is much greater than m , the latter can be neglected. In this approximation one has:

$$\frac{1}{v} = \frac{d|\mathbf{p}|}{dE} \approx \cos \frac{E}{\kappa}. \quad (27)$$

If $E \ll \kappa$, the right-hand-side of eq. (27) can be expanded into a power series, retaining only the first few terms. One finds:

$$\frac{1}{v} \approx 1 - \frac{E^2}{8\kappa^2} + O\left(\left(\frac{E}{\kappa}\right)^4\right). \quad (28)$$

The deformed Poincaré group leads to the result that a massless particle is, in the usual sense of the word, always superluminal: what one accepts these days as the "speed of light" is only the *low frequency limit* of the group velocity.

One can now place a lower limit on the value of κ as follows. Suppose that a periodic UHE signal emitted by a distant source, (e.g. a pulsar) is observed.

The time delay a δ function pulse suffers between a source a distance d away from a (terrestrial) detector is given by the elementary formula:

$$\delta t = \frac{d}{v} \approx d \left(1 - \frac{E^2}{8\kappa^2} + \dots\right) \quad (29)$$

Any periodic signal can be expanded into a Fourier series. For a signal of fundamental period T and energy E , one has at the source:

$$I(E, t_s) = \sum_{n=-\infty}^{\infty} I_n^S(E) \exp in\omega t_s, \quad (30)$$

where $\omega = 2\pi/T$. The intensity is real, $I_{-n} = I_n^*$. At the detector, the Fourier coefficients differ from the ones at the source by an energy dependent phase shift, analogously to the time delay in eq. (29):

$$I_n^D(E) = I_n^S(E) \exp \left[i\omega n d \frac{E^2}{8\kappa^2} \right] \quad (31)$$

In arriving at eq. (31), we factored out a phase common to all Fourier coefficients and we omitted terms of $o(E^2/\kappa^2)$. We average eq. (31) over the energies, assuming, as usual that the spectrum of the source is given by a power law. (This averaging is needed because one does not know the energy dependence of the Fourier coefficients at the source.) The normalized spectrum, S , for a detector of nominal threshold energy E_{th} is:

$$S(E) = (\alpha - 1) \frac{1}{E_{th}} \left(\frac{E}{E_{th}} \right)^{-\alpha}. \quad (32)$$

Thus, the Fourier coefficients averaged over the spectrum are given by:

$$\langle I_n^D \rangle = \frac{\alpha - 1}{E_{th}^{1-\alpha}} \int_{E_{th}}^{\infty} dE I_n^S(E) E^{-\alpha} \exp \left[i\omega n d \frac{E^2}{8\kappa^2} \right] \quad (33)$$

As a consequence of the velocity dispersion, both the magnitude and the phase of $\langle I_n^D \rangle$ differ from I_n^S . Unfortunately, one has no information on the Fourier coefficients at the source. On assuming that the pulse at the source is proportional to a δ function at all energies — i.e. that the entire observed width of the pulse arises from velocity dispersion — one obtains a *lower limit* on κ . With the assumption $I_n(E) = I = \text{const}$, one has to the same accuracy as in eq. (28):

$$|I_n^D|^2 = |I|^2 \left\{ 1 - \left(\frac{\omega d n}{8} \right)^2 \left(\frac{E_{th}}{\kappa} \right)^2 \left[\frac{(\alpha - 1)^2}{(\alpha - 2)^2} - \frac{\alpha - 1}{\alpha - 3} \right] \right\} + \dots \quad (34)$$

and

$$\arg I_n^D = n \frac{\omega d E_{th}}{4\kappa} \frac{\alpha - 1}{\alpha - 2} + \dots \quad (35)$$

The unknown coefficient I is eliminated by taking ratios of the moduli of the Fourier coefficients. It is evident from these equations that in order to obtain good lower limits on κ ,

- one has to maximize the distance between the source and the detector,

- one has to use data at the highest energies available.

There was a burst observed in the binary system HER X1 in 1986. The distance of this system is about 5 kpc; particles of the highest energies were observed by the CYGNUS collaboration by means of detecting extensive air showers induced by them, *cf.* ref. [51]. The EAS of lowest energy, quoted in ref. [51] is 100 TeV, just above the stated threshold. Using the period-folded data from this experiment and the published fundamental period, ($T \approx 1.24$ sec), ref. [41], we determined the first few Fourier coefficients from the data, [42]. (Due to the fact that 11 showers were observed by the CYGNUS collaboration, the data allow for a meaningful determination of only the first few (≈ 5) complex Fourier coefficients. The coefficient I_0 , corresponding to a constant intensity, can always be subtracted.) The result is displayed in Table 1.

Table 1 : Fourier coefficients

n	$ I_n^D ^2$	$\arg I_n^D$
1	0.99	0.044
2	0.96	0.088
3	0.92	0.13
4	0.86	0.17
5	0.79	0.21

One observes that the phases of the Fourier coefficients determined from the data exhibit no noticeable trend; probably, they are considerably affected by “noise”, *i.e.* by the various errors in the data set. This is a general phenomenon [43]. For this reason, only the moduli of the Fourier coefficients were used in order to obtain a lower limit on κ . (One determines by inspection that the moduli contain less noise than the phases: they are decreasing almost linearly with n .)

Using the ratio, $|(I_2^D)/(I_1^D)|^2$ from eq. (34), one can compute a value of κ . As stated before, this is a lower limit on its value. Unfortunately, however, it turns out that the Taylor expansion quoted above is a very poorly convergent one: it can give a qualitative feeling about the relevant orders of magnitude at best. (This is, in essence, due to the fact that according to any estimate, the arguments of the incomplete Γ functions are of $O(1)$.) In order to do better, one has to resort to numerical techniques, taking advantage of the fact that the Fourier coefficients are expressible in terms of incomplete Γ functions. We got in this way

$$\kappa \gtrsim 3 \times 10^{12} \text{ GeV.} \quad (36)$$

In arriving at the result eq. (36), we used a value, $\alpha = 1.5$, obtained from a rough fit to the data. The quoted lower limit is rather insensitive to the precise

value of the spectral index. On taking other independent ratios of the moduli of the Fourier coefficients, one gets essentially the same lower limit on κ : the various ratios are certainly not inconsistent with each other within the accuracy to which the Fourier coefficients can be determined. (We believe that the best result is obtained by taking the ratio of the first two Fourier coefficients, since those are the ones which are the least affected by "noise" in the data. However, one has to develop a statistical treatment of the Fourier coefficients in order to utilize the full information available.)

It is worth remarking that in ref. [41], the authors quote a lower limit on the characteristic mass scale(s) of the *ad hoc* modifications of relativistic kinematics, [34]. They find a result which is equivalent in the present formulation to $\kappa \gtrsim 510^{11}$ GeV. The improvement represented by eq. (36) is due to the fact that by computing Fourier coefficients, one uses a larger amount of information contained in the data than by using the simpler estimates in ref. [41]. (In the approximation used in ref. [41], the entire estimate is dominated, in essence, by the data on the shower of highest energy. By contrast, the method outlined above uses the entire information available from the observations; hence, the sensitivity of the estimate to background contamination is also reduced.) In particular, we tested numerically the sensitivity of the lower limit on κ to the energies of the showers. It is obvious that the method used in [41] is virtually insensitive to the energy of the least energetic shower and so is ours. However, the cited uncertainty of the measurement of the primary energies (about a factor of 2), leads to a factor of 4 uncertainty in the value of κ in the method used in [4], whereas it leads to virtually no change in the lower limit in our computation. Thus, even though the numerical improvement in the lower limit on κ is not a dramatic one, the uncertainties due to background contamination and inaccuracies in the energy determination are substantially reduced.

It appears that deviations from Poincaré invariance are negligible at least up to energies which are fairly close to presumed grand unification masses. Presumably, as larger data samples become available, this limit can be further improved, or, better still, one may be able to estimate the energy where the special theory of relativity breaks down due to new physics at the unification scale. UHE astrophysical data are unique in this respect due to the high energies and large distances involved.

4.3 Simulation of "Anomalous" Air Showers.

The observation of extensive air showers (EAS) in the atmosphere associated with particle beams from pointlike astrophysical objects offers a unique avenue of investigation. The center-of-mass energy of a primary particle from one of these sources and a nucleon in the atmosphere can be orders of magnitude greater than that accessible at our present accelerators. The resultant EAS from such a collision is characterized by: (1) the physics of the initial interaction and (2) the

physics of the shower development. If the initial interaction is, in any way, similar to hadronic interactions, one expects multiplicities of the order of 100 or so at the relevant energies. On the average, the energy per particle will drop below substantially below the characteristic energy scale of new physics³. Therefore, it is expected that whatever new physics is present beyond the standard model, it will leave no footprint beyond the first interaction. Consequently, in describing the evolution of the shower, one has to model the first interaction according to the new physics to be explored and then evolve the shower by standard model physics. The simulation of EAS was carried out within the framework of a "generic" composite model of quarks and leptons. It should be kept in mind, however that there are many other possibilities to which a similar analysis applies. The primary aim in this investigation has been to assess which observables in an EAS are sensitive to changes in the initial interaction.

Soon after the discovery that pointlike objects were significant (detectable) sources of ultrahigh energy (UHE) cosmic rays, an anomalous muon content was measured in their associated EAS [1]. This is consistent with a scenario where UHE neutrinos interact strongly in the atmosphere because of some new physics [44]. We studied EAS initiated by an anomalously strong neutrino interaction. The argument for strong neutrino interactions is reviewed in subsection 4.3.1, a phenomenological model for the interaction is developed in subsection 4.3.2, the consequences of compositeness are explored in subsection 4.3.3, and the simulation and subsequent analysis of the EAS is described in subsections 4.3.4 and 4.3.5.

For the sake of definiteness, the results of simulated EAS are presented as they would be observed by the CYGNUS experiment at Los Alamos, although any detection level can be accommodated.

4.3.1 Strong Neutrino Interactions

A strong interaction between quarks and leptons, assuming the new physics resembles QCD, should possess two important features: (1) the new interaction becomes relevant when the transverse momentum (P_t) exchanged is the order of the scale of new physics Λ_{NP} , and (2) the magnitude of the cross section is much larger than the order of Λ_{NP}^{-2} . Feature (1) is a statement that approximate asymptotic freedom sets in only above the characteristic scale of the interaction; P_t is only a convenient kinematic variable. Feature (2) is not as easy to understand. In fact, a straightforward analysis of the effects of new physics suggests that the magnitude of the cross section should be of order Λ_{NP}^{-2} . However, this assumes that the dominate term in an effective Lagrangian of the new physics is a four-fermion interaction; this is true only far below the scale Λ_{NP} . If the new physics is a confining interaction with a rapid turn-on, then all length scales are relevant, not just long wavelengths, and a cross section of order Λ_{QCD}^{-2} can be gener-

³Typically, the flux of primaries with energies such that $\sqrt{s} \gg 1\text{TeV}$ is expected to be very small for any kind of point source.

ated. As an indication of the generality of these results, consider the features of a technicolor-induced interaction between quarks and leptons. The technifermions must be visible, i.e., have P_t of order Λ_{NP} , for an interaction to take place. A reasonable (quark-like) momentum distribution for the technifermions inside the quark and lepton will induce a rapid turn-on for the interaction. Finally, once the pseudo-threshold is crossed, the cross section will be $\gg \Lambda_{NP}^{-2}$, because no single scale dominates. An analysis of the unitarity of the Standard Model suggests that the scale of new physics is a few TeV, while the absence of strongly interacting neutrinos in our accelerators only places the same lower bound on this type of new physics.

The axis of a reconstructed EAS and the arrival time maintain the direction to the source and the timing of the source respectively. This allows an individual shower to be associated with a specific point source. A beam of photons from such a source should exhibit absorption on the thermal ensemble of 3°K background photons with a dip in the spectrum around a few PeV. The same pointlike objects that are sources of UHE photons produce a UHE neutrino flux that is up to an order of magnitude greater. Neutrinos, "ordinary" or not, should not be significantly absorbed at the center-of-mass energies involved. An analysis of the point source Cygnus X-3 suggests that an unabsorbed component of neutral particles reaches our atmosphere and interacts strongly. Because the muon excess can be explained by the production and decay of pions, this is interpreted as a signature of strong interactions between quarks and leptons, where the cross section of neutrinos on nucleons $\sigma_{\nu N}$ is .1-10 mb [46]. In this scenario, an UHE neutrino from a point source strikes a nucleon in the atmosphere with $\sqrt{S}/2 \geq \Lambda_{NP}$. The EAS develops like a proton-induced shower initiated deeper in the atmosphere than usual, but it maintains the directionality and timing of the point source as in a photon-induced shower.

4.3.2 The Phenomenology Of Anomalous Interactions

The model is inspired by the parton description of deep inelastic scattering that predated QCD. A neutrino scatters off a quark in a nucleon because of a strong interaction induced by new physics. The struck quark is assumed to carry all the momentum of the target nucleon: At the kinematic threshold for new physics, the relevant quark has a momentum fraction far out on the tail of its momentum distribution. The scattered "neutrino" is actually a jet of partons with color (anti-color) quantum numbers, because some anti-color (color) quantum numbers were lost to the scattering off the quark, which must evolve into a physical state. Since the time scale of new physics is of order Λ_{NP}^{-1} , which is $\ll \Lambda_{QCD}^{-1}$, it follows that the partons in the jet "quarkize," before they "hadronize." For reasons explained in section 4.3.3, the products of the strong interaction are ignored for the remainder of this investigation.

A "quarkization" procedure was adopted in the same spirit as the Field-Feynman approach to hadronization. In the Field-Feynman model, a hard quark pulls $q\bar{q}$ pairs out of the vacuum, producing a chain of such pairs. These pairs are identified as mesons carrying a fraction of the parent quark's energy, longitudinal momentum, or some other kinematic variable that is considered relevant; I choose longitudinal momentum (P_L) to be specific. The appearance of a meson with a fraction of the parent quark's P_L does not determine whether the meson was produced initially or far down the chain. As a result, the meson's P_L distribution function obeys an integral equation that is solved once the single meson P_L distribution is chosen.

The fundamental modification to this approach for a colored parton jet is that the "asymptotic" states are not color singlets. By conservation of quantum numbers for a leptonic jet that has lost some colored parton (and maybe other partons), it follows that the jet "dresses up" into $q\bar{q}$ pairs carrying momentum fraction z and $\bar{z}$.

Because the time scale $\Lambda_{NP}^{-1} \ll \Lambda_{QCD}^{-1}$, the quark and anti-quark are independent and should *not* be immediately paired into a meson. This result applies generally to any type of new physics that blurs the distinction between quarks and leptons and has a characteristic scale much larger than Λ_{QCD} : The introduction of partons simply simulates the flow of quantum numbers. The integral equation for the probability $D(z, \bar{z})$ of finding a $q\bar{q}$ pair with momentum fraction z and $\bar{z}$ is as follows:

$$D(z, \bar{z}) = f(z, \bar{z}) + \int \frac{dz'}{z'} \frac{d\bar{z}}{\bar{z}} f(z', \bar{z}) D\left(\frac{z}{1-z'-\bar{z}}, \frac{\bar{z}}{1-z'-\bar{z}}\right). \quad (37)$$

$f(z, \bar{z})$ is the single pair probability function. With the assumption that $f(z, \bar{z}) \propto (z\bar{z})^\mu (z + \bar{z})^\nu$, μ and $\nu > 0$, the solution of $D(z, \bar{z})$ is:

$$D(z, \bar{z}) = \frac{\Gamma(2\mu + \nu + 1)}{\Gamma(\mu)^2 \Gamma(\nu + 1)} (z\bar{z})^\mu E(z + \bar{z}). \quad (38)$$

This choice of $f(z, \bar{z})$ guarantees the correct phase behavior as $z, \bar{z} \rightarrow 0$. $E(x)$ is a power series in x , and the reader's eyes are spared from its exact form. Demanding that the probability function leads to a finite multiplicity of $q\bar{q}$ pairs puts further limitations on the allowed values of μ and ν . Values of $\mu = \frac{1}{2}$ and $\nu \geq 1$ give reasonable multiplicities. The subsequent analysis is insensitive to reasonable changes in these values. The final form chosen for $D(z, \bar{z})$ is:

$$D(z, \bar{z}) = \frac{4}{\pi} (z\bar{z})^{\frac{1}{2}} \frac{(1 - z - \bar{z})^3}{z + \bar{z}}. \quad (39)$$

Finally, up to this point, the effects of P_L have been ignored. Since the "dressing up" into quarks involves the new physics, the $q\bar{q}$ pair should have a relative P_L of order Λ_{NP} . It was assumed that the P_L distribution is Gaussian and centered on

Λ_{NP} . The subsequent analysis is insensitive to reasonable variations of the exact functional form of the P_t distribution.

Once $q\bar{q}$ pairs are produced, they must be hadronized. This was done with the Lund Monte Carlo program JETSET version 7.2 [47], because of its accessibility and simplicity. By sampling the probability distribution derived from $D(z, \bar{z})$ with a random number generator, $q\bar{q}$ pairs are generated until the energy of the jet is exhausted up to a small cutoff. Once some assumption is made about the energy and angular dependence of the new physics cross section, which is presented in subsection 4.3.3, the analysis can proceed. Each pair is assigned a P_t with respect to the jet in the jet frame, boosted to the pair's center-of-mass frame, and presented to JETSET. The final particle multiplicity, which is mildly dependent on the cutoff, is ≈ 119 per neutrino-nucleon interaction. The products of JETSET are then boosted to the lab frame and made available as an input to a standard EAS Monte Carlo program such as SHOWERSIM [48].

4.3.3 A Model Of Composite Quarks And Leptons

If quarks and leptons are composite and share constituents (henceforth referred to as "preons"), then, in high- P_t scattering processes, neutrinos reveal any colored structure they might possess. Historically, the deep inelastic scattering of electrons off a proton target exhibited the composite nature of protons. The differential cross section, with respect to the energy fraction of the scattered neutrino, follows from the assumptions that: (1) A colored preon in a neutrino with $P_t \geq \Lambda_{NP}$ scatters off a quark in a nucleon, (2) the quark carries all of the momentum of the nucleon, (3) the preon-quark total cross section is a constant over the kinematic range of interest, and (4) the fundamental interaction between fermions is vector-like (the general features are independent of this choice). The result is: [45]:

$$\frac{d\sigma}{d\zeta} = (1 - \zeta)\sigma_T \int_{\eta_{min}}^{\eta_{max}} \frac{d\eta}{\eta^2} \frac{1 + (1 - \eta)^2}{(\zeta + \eta - 1)}, \quad (40)$$

$$\begin{aligned} \eta_{min} &= \frac{1}{2}(1 - \sqrt{1 - \tau}) \\ \eta_{max} &= \frac{1}{2}(1 + \sqrt{1 - \tau}). \end{aligned}$$

ζ is the energy fraction of the scattered "neutrino", η is the inelasticity of the event, σ_T is the total preon-quark cross section, and $\tau = 4\Lambda_{NP}^2/S$ is a dimensionless quantity depending on the scale of new physics and the center-of-mass energy of the neutrino-nucleon system. An earlier analysis of this composite model suggests that σ_T is of order Λ_{QCD-2} [45], while Λ_{NP} , as explained in subsection 4.3.1, is a few TeV. The differential cross section for producing a $q\bar{q}$ pair with energies E_1 and E_2 is the convolution of the preon-quark differential cross section with $D(E_1, E_2)$ as derived in subsection 4.3.2:

$$\frac{d^2\sigma}{dE_1 dE_2} = \frac{16}{S\pi} \frac{(E_1 E_2)^{\frac{1}{2}}}{E_1 + E_2} \int_{\sqrt{\tau}}^1 \frac{d\zeta}{\zeta} \left(1 - \frac{2(E_1 + E_2)}{\sqrt{S}\zeta}\right)^3 \frac{d\sigma}{d\zeta}. \quad (41)$$

While we invoked a specific form of the initial interaction, the choice is almost irrelevant; kinematics alone determine that the momentum is evenly shared by the scattered neutrino and the quark-preon products in the center-of-mass frame, and they are strongly boosted into the lab frame. A detailed analysis of this form of the interaction, however, reveals one important feature: The products of the preon-quark interaction have a large invariant mass, and, as a result, are very soft. For this reason, this vertex of the interaction is ignored. From the point of view of a proponent of composite interactions, this can only reduce the magnitude of a possible signal.

4.3.4 EAS Simulation

EAS events were simulated with the Monte Carlo SHOWERSIM incorporating a library of EGS4 showers [49] to provide increased energy and timing resolution [50]. For the sake of definiteness, the EGS4 showers were generated for a single detection level of 800 g cm^{-2} , i.e., the depth at which the CYGNUS detector is located. However, if one is willing to forgo information about the energy and timing of the electron and photon component of an EAS, showers can be generated with SHOWERSIM alone down to sea-level; or, if one is willing to invest the time and resources, another library of EGS4 showers could be generated for a different detection level. In fact, this investigation did not exploit the timing information. We expect, however, to include it in the future.

The detector CYGNUS was modelled by an approximation to its 1986 configuration: a grid of 1 m^2 detectors spaced every 10 m over an area of 10^4 m^2 [51]. The CYGNUS trigger for recording an event during this time was that 20 or more detectors must be hit, and 8 or more of those with 3 or more particles, within a time window of 500 ns [52].

The average interaction depth λ in g cm^{-2} for a neutrino incident on the atmosphere is given by:

$$\lambda \approx 1700/\sigma_{\nu N}. \quad (42)$$

(The validity of this equation was discussed earlier in this report in connection with atmospheric neutrino production.) The range of .1-10 mb for $\sigma_{\nu N}$ from the Cygnus X-3 analysis allows for showers to begin either high in the atmosphere or underground. However, this is only an average interaction depth, so, even if the cross section is small, some showers will begin above ground. Regardless of where a shower begins, it must satisfy the CYGNUS trigger to be recorded. The depth of the first interaction was varied; the resulting electron density distribution was then used to determine whether the trigger was satisfied. The laboratory energy of the primary neutrino was chosen to be 12.5 PeV, which makes the new physics kinematically accessible if $\Lambda_{NP} \leq 2.5 \text{ TeV}$. This choice results in an average shower energy of 3.3 PeV; the remaining energy is lost to the soft hadronic jet from the preon-quark interaction. Finally, all showers were initiated at a zenith angle of 0° . This choice allows for easier analysis and interpretation of the results; if $\sigma_{\nu N}$ is

smaller than .1-10 mb, showers could still be detected from point sources at larger zenith angles. It was found that an EAS originating at 500 g cm^{-2} minimally satisfied the CYGNUS trigger. Computational constraints limited this analysis to a fixed interaction depth, primary energy, and zenith angle. For comparison, standard showers were also generated using SHOWERSIM, with a 3 PeV proton primary incident at the top of the atmosphere.

4.3.5 Results and Analysis

The lateral particle density, the average energy per particle, and the absolute particle number for electrons (e), photons (γ), muons (μ) and hadrons (π) were computed. These quantities are extractable from a standard EAS array such as CYGNUS. 3 PeV proton-induced showers are used for comparison. A fit was made to the electron density with the Nishimura-Kamata-Greisen (NKG) function to determine the shower age s , using a Molière radius at CYGNUS of 50 m. A neutrino-induced shower has an average shower age $s = .91$ (Fig.5). It is common lore that a proton-induced shower has $s \approx 1.25$ and the analysis of 3 PeV proton-induced showers resulted in $s = 1.24$. Therefore, the deeply initiated showers are markedly different from hadronic showers. The muon and hadron content of the neutrino-induced EAS, however, is indicative of a hadronic shower (Fig.6); however, the muon lateral distribution is narrower, and the photon-to-electron ratio is larger. In fact, the photon-to-electron ratio of the neutrino-induced EAS is similar to that of a 100 TeV proton primary; the total number and energy of particles is too high, however, to be misidentified as such. In principle, the muon lateral distribution for a ν -induced EAS can be distinguished from a p-induced EAS provided an array has efficient muon detectors with a coverage of $> 10^4 \text{ m}^2$. Scintillators are more sensitive to electrons than photons, but the photon density (and, hence, the photon-to-electron ratio) can be measured in at least two ways:

1. some scintillator can be covered with a thin lead sheet to allow the photons to convert,
2. a water-Cherenkov detector can be incorporated into the array.

With these techniques, an EAS array could distinguish the showers on the basis of the electron and photon densities.

The average energy per particle is another observable readily obtained by an EAS array. For the electromagnetic component of an EAS, this investigation shows that a ν -induced shower has a slightly higher average energy within 10m of the shower core than a p-induced shower. In contrast, the average energy per muon is consistently less than for a p-induced shower (Fig.7). The first observation is expected, since the ν -induced EAS originate much lower in the atmosphere, and the electromagnetic component has not scattered through as much material. The

second observation was unexpected and a consequence of the details of the initial interaction. (At present, it is not clear how robust such a feature is; this is a question to be investigated in the future.) For hadronic showers, a larger number of low energy muons is indicative of a primary with a larger mass number A . Qualitatively, the neutrino produces muons like a light nucleus.

The lesson is this: muons carry information about the initial interaction, while the electromagnetic component washes out some of the details.

The "smoking gun" for anomalous EAS is a shower age indicative of photon-induced shower and muon content indicative of a proton-induced shower. More subtle signals of an anomalously strong neutrino interaction are a narrow muon lateral distribution, a large number of lower energy muons, an electron-to-photon ratio like a lower energy EAS, and a high muon-to-electron ratio. The Monte Carlo results on the muon lateral distribution corroborate those of an analytic approach to the problem of ν -induced EAS [54]. The advantage of the Monte Carlo approach is that finer details, i.e., those dependent on energy and timing resolution, can be extracted. The disadvantage is the time commitment to simulate an EAS with millions of particles at the detection level. Therefore, the two-pronged approach to this problem advocated by us some time ago: *i.e.* using both analytical and Monte Carlo techniques appears to be the most reasonable one.

A careful measurement of the muon component of an EAS could prove fruitful in a search for new physics. Experiments that continue to associate an EAS with a point source on the basis of a muon-poor content will fail to find any such signal. With the technology available, it should be possible to extend this investigation to showers generated in rock and observed in underground proton decay experiments. It is important that these collaborations make a published statement of a muon anomaly in their data, if it indeed exists.

5 Search for New Physics in Heavy Quark Decays

A very substantial number of $t\bar{t}$ pairs will be produced at hadronic colliders in the near future. This will allow a precision measurement of the mass (m_t) of the top quark and the degree of polarization of the W -boson produced from the top quark decay. The radiative corrections arising from the electroweak [55, 56] and QCD [57, 58] interactions have been extensively studied. The electroweak radiative correction to the decay of the top quark has been shown to be around 1% if the G_F scheme [56] is used in the Born decay rate. The QCD correction, however, is around 10% for $100 \text{ GeV} < m_t < 200 \text{ GeV}$.

It is important to study QCD corrections to the decay of the t quark for at least two reasons:

- m_t is an undetermined parameter of the Standard Model which must be extracted from experiment [59], and

- the degree of polarization of the W -boson from top quark decay is sensitive to any new physics which might occur at the 100 GeV scale [60].

The determination of m_t depends on the jet definition and is affected by b -tagging. It is found, however, that the peak of the m_t distribution does not shift after including the QCD effects. The shape of the distribution, is, however, no longer of the Breit-Wigner form if b -tagging is used to discriminate a bottom quark jet from a gluon jet.

The distribution of m_{be} , the invariant mass of the bottom quark and lepton produced in the decay, is less sensitive to smearing effects of the detector; m_{be} is a measure of the polarization of the W -boson from top quark decay and therefore sensitive to m_t .

5.1 Radiative Corrections

5.1.1 Decay rate of a heavy top quark

The next to leading order radiative corrections to the decay process $t \rightarrow bW^+$ arising from either the electroweak [55, 56] or QCD [57, 58] interactions have appeared many times in the literature. The electroweak correction to the lowest order decay width $\Gamma_0(t \rightarrow bW^+)$ has been shown to be about 1% for $100 \text{ GeV} < m_t < 200 \text{ GeV}$ if the G_F scheme [56] is used. In this scheme, the tree level decay rate is

$$\Gamma_0(t \rightarrow bW^+) = \frac{G_F}{8\sqrt{2}\pi} \frac{(m_t^2 - M_W^2)^2}{m_t} \left(1 + \frac{2M_W^2}{m_t^2} \right), \quad (43)$$

where we have dropped terms of $\mathcal{O}(m_b^2/m_t^2)$ when considering the kinematics of a heavy top quark decay at the Born level. In terms of G_F , the W -boson mass M_W is obtained from

$$M_W^2 \left(1 - \frac{M_W^2}{M_Z^2} \right) = \frac{\pi\alpha}{\sqrt{2}G_F} \frac{1}{1 - \Delta r},$$

where the Z -boson mass M_Z is 91.17 GeV, the fine structure constant $\alpha = \frac{1}{137.036}$, Δr is defined by Sirlin [61] in the radiative corrections to the lifetime of the muon, and $G_F = 1.16637 \times 10^{-5} \text{ GeV}^{-2}$. For the purposes of the analysis, the following equation [56] has been used.

$$\begin{aligned} M_W &= 80.22 \text{ GeV} & \text{for } m_t &= 140 \text{ GeV}, \\ M_W &= 80.68 \text{ GeV} & \text{for } m_t &= 200 \text{ GeV}, \end{aligned}$$

The decay width $\Gamma_0(t \rightarrow bW^+)$ has been computed from eq.(43). The result is: $\Gamma_0 = 0.67 \text{ GeV}$ for $m_t = 140 \text{ GeV}$, and $\Gamma_0 = 2.44 \text{ GeV}$ for $m_t = 200 \text{ GeV}$. Compared with the QCD radiative correction, which is found to be around 10% for $100 \text{ GeV} < m_t < 200 \text{ GeV}$, the 1% electroweak radiative correction in the G_F scheme to

the lowest order decay width $\Gamma_0(t \rightarrow bW^+)$ is negligible and was ignored in the rest of the work.

The first step was a study of the kinematics of the decay process $t \rightarrow bW^+ \rightarrow be^+\nu_e$ when including the QCD radiative corrections. From the kinematics, one derives the correction to the determination of m_t and the degree of polarization of the W -boson from top quark decay. In considering the QCD radiative corrections to the decay process $t \rightarrow bW^+$, the work in reference [58] was followed. The inclusive decay rate Γ_{inc} is found to be

$$\begin{aligned}\Gamma^{(3)+(4)} &\equiv \Gamma_{inc}(t \rightarrow bW^+(g)) \\ &= \Gamma_0(t \rightarrow bW^+) \left\{ 1 - \frac{\alpha_s}{\pi} C_F \left[2 \text{Sp}(y) + \ln(y) \ln(1-y) + \frac{\pi^2}{3} \right] \right\} \\ &\quad - \frac{G_F m_t^3}{8\sqrt{2}\pi} \frac{\alpha_s}{\pi} C_F \left\{ (1-y-2y^2)y \ln(y) \right. \\ &\quad \left. + \frac{1}{2}(1-y)^2(5+4y) \ln(1-y) - \frac{1}{4}(1-y)(5+9y-6y^2) \right\}, \quad (44)\end{aligned}$$

with the color factor $C_F = \frac{4}{3}$ and $y \equiv M_W^2/m_t^2$. The Spence function is evaluated using

$$\text{Sp}(y) = - \int_0^y dt \frac{\ln(1-t)}{t} = \sum_{n=0}^{\infty} B_n \frac{z^{n+1}}{(n+1)!},$$

for $|y| < 1$ and $\text{Re}(y) < 0.5$, with $z = -\ln(1-y)$. B_n are the Bernoulli numbers, which can be found in reference [62]. In eq. (44), terms of $\mathcal{O}(m_b^2/m_t^2)$ have been dropped for a heavy top quark decay. In numerical calculations, the leading order running coupling α_s ($\simeq 0.1$) in the $\overline{\text{MS}}$ scheme was used.

From eq.(44), the radiative correction to the decay rate Γ_0 is found to be about -10%. The QCD corrected decay rate is

$$\begin{aligned}\Gamma_{inc}(t \rightarrow bW^+(g)) &= 0.61 \text{ GeV} \quad \text{for } m_t = 140 \text{ GeV}, \\ \Gamma_{inc}(t \rightarrow bW^+(g)) &= 2.20 \text{ GeV} \quad \text{for } m_t = 200 \text{ GeV}.\end{aligned}$$

The kinematic phase space of this decay process consists of two parts: a part where the emitted gluon is soft, and a part where the gluon is hard. This separation is accomplished by introducing a theoretical cutoff E_g^{th} on the energy of the gluon in the rest frame of the top quark. For $E_g^{th} < m_b$, the $\mathcal{O}(\alpha_s)$ decay width of a top quark with no hard gluon Γ^{nhg} is:

$$\begin{aligned}\Gamma^{(3)} &\equiv \Gamma^{nhg}(t \rightarrow bW^+) \\ &= \Gamma_0(t \rightarrow bW^+) \left\{ 1 - \left(\frac{\alpha_s}{\pi} \right) C_F \left[S(m_t, m_b, E_g^{th}) + \frac{3}{2(1+2y)} \ln(1-y) \right] \right\} \quad (45)\end{aligned}$$

where

$$S(m_t, m_b, E_g^{th}) = \left\{ \frac{5}{2} + 2\ln(2) + 2\ln\left(\frac{E_g^{th}}{m_t}\right) \right\} \left\{ 1 + \frac{1}{2}\ln\left[\frac{x}{(1-y)^2}\right] \right\} \\ - \frac{3}{2} + \frac{\pi^2}{6} + \text{Sp}(y) + 2[\ln(1-y)]^2 - \ln(x) \left[\frac{1}{2} + \ln(1-y) \right],$$

with $x \equiv m_b^2/m_t^2$, $y \equiv M_W^2/m_t^2$. The bottom quark mass m_b was kept to regularize the collinear divergent terms $\ln(x)$ when $m_b \rightarrow 0$. Of course, these divergent terms must necessarily be cancelled by those $\ln(x)$ terms arising from integrating the phase space of the gluon with $E_g > E_g^{th}$ from the top quark decay $t \rightarrow bW^+g$.

As shown in eq.(44), the inclusive decay rate $\Gamma^{(3)+(4)}$ is finite as $m_b \rightarrow 0$ and is independent of the theoretical cutoff E_g^{th} , which regularizes the soft divergence in $\Gamma^{(3)}$ as $E_g \rightarrow 0$. The decay width of a hard gluon event Γ^{hg} is given by

$$\Gamma^{hg}(t \rightarrow bW^+g) \equiv \Gamma^{(3)+(4)} - \Gamma^{(3)}, \quad (46)$$

where the soft divergence is regularized by E_g^{th} and the collinear divergence is absent by keeping m_b non-zero. To justify the choice of E_g^{th} , we compare eq.(46) with the result obtained by integrating the whole phase space for the decay process $t \rightarrow bW^+g$ with the gluon energy $E_g > E_g^{th}$. We perform this comparison in the next section.

5.1.2 Polarization of the W -boson from top quark decay

The most general operators for the coupling of $t \rightarrow bW^+$ are described by the effective interaction Lagrangian [60]:

$$\mathcal{L} = \frac{g}{\sqrt{2}} \left[W_\mu^- \bar{b} \gamma^\mu (f_1^L P_- + f_1^R P_+) t - \frac{1}{M_W} \partial_\nu W_\mu^- \bar{b} \sigma^{\mu\nu} (f_2^L P_- + f_2^R P_+) t \right] \\ + \frac{g}{\sqrt{2}} \left[W_\mu^+ \bar{t} \gamma^\mu (f_1^{L*} P_- + f_1^{R*} P_+) b - \frac{1}{M_W} \partial_\nu W_\mu^+ \bar{t} \sigma^{\mu\nu} (f_2^{R*} P_- + f_2^{L*} P_+) b \right], \quad (47)$$

where $P_\pm = \frac{1}{2}(1 \pm \gamma_5)$, $i\sigma^{\mu\nu} = -\frac{1}{2}[\gamma^\mu, \gamma^\nu]$, and $*$ denotes complex conjugation. In general, the form factors $f_1^{L,R}$ and $f_2^{L,R}$ can be complex. If the W -boson can be off-shell, there are additional factors, *viz.*

$$\partial^\mu W_\mu^- \bar{b} (f_3^L P_- + f_3^R P_+) t + \partial^\mu W_\mu^+ \bar{t} (f_3^{R*} P_- + f_3^{L*} P_+) b. \quad (48)$$

The form factors in (48) vanish for an on-shell W -boson or when the off-shell W -boson couples to massless on-shell fermions. Since we are considering the decay mode $t \rightarrow bW^+ \rightarrow b f_1 \bar{f}_2$, where f_1 and f_2 are light fermions, we shall not consider the form factors $f_3^{L,R}$ in this work since they vanish in the limit $m_f \rightarrow 0$. At tree level, the Standard Model form factors are

$$\begin{aligned} f_1^L &= 1 \\ f_1^R &= f_2^L = f_2^R = 0. \end{aligned}$$

With the QCD correction to $\mathcal{O}(\alpha_s)$ [58],

$$\begin{aligned} f_1^L &= 1 + f_1^{L'} \\ f_2^R &= C_F \frac{\alpha_s}{2\pi} \frac{1}{2\sqrt{y}} \ln \left(\frac{1}{1-y} \right), \end{aligned}$$

with $C_F = 4/3$ and $y \equiv M_W^2/m_t^2$. $f_1^{L'}$ is not infrared finite; both $f_1^{L'}$ and f_2^R are of $\mathcal{O}(\alpha_s)$. It follows from an approximate helicity conservation that $f_1^R = f_2^L = 0$ in the limit of $m_b = 0$. The probability for various W^+ polarization states from an unpolarized top quark decay in the limit $m_b \ll m_t$ is [60]

$$\begin{aligned} N_T &\equiv |\mathcal{M}(\lambda_W = -1)|^2 \sim |f_1^L + \frac{m_t}{M_W} f_2^R|^2, \\ N_L &\equiv |\mathcal{M}(\lambda_W = 0)|^2 \sim \frac{1}{2} |\frac{m_t}{M_W} f_1^L + f_2^R|^2, \end{aligned}$$

and

$$|\mathcal{M}(\lambda_W = +1)|^2 = 0,$$

where $\lambda_W = -1, 0, +1$ indicates the left-handed, longitudinal, and right-handed states of polarization of the W -boson. The fraction of longitudinally polarized W -bosons from top quark decay up to $\mathcal{O}(\alpha_s)$ is defined as

$$\begin{aligned} f_L \equiv \frac{N_L}{N_L + N_T} &= \frac{1}{1+2y} \left[1 - \frac{1}{\sqrt{y}} \text{Re}(f_2^R) \frac{4y(1-y)}{1+2y} \right], \\ &= \frac{1}{1+2y} \left[1 - C_F \frac{\alpha_s}{\pi} \left(\frac{1-y}{1+2y} \right) \ln \left(\frac{1}{1-y} \right) \right], \end{aligned} \quad (49)$$

for $0 < y < 1$, where $y \equiv M_W^2/m_t^2$. The result is that the QCD correction to the tree-level value of f_L , 0.604 (0.754), is -0.82 (-0.57)% for a 140 (200) GeV top quark. Because the correction is extremely small, this suggests that a good method to measure the top quark mass is to measure f_L [58, 60, 63]. However, in general, f_L will be slightly influenced by the kinematic cuts imposed on either the final state jet or charged lepton from the decay of a top quark produced in the lab frame.

As discussed in reference [60], the fraction of longitudinally polarized W -bosons f_L can be measured by:

- fitting the polar angle θ_e^* distribution of the e^+ in the rest frame of the W^+ with the z -axis defined as the moving direction of the W^+ in the rest frame of the top quark, or

- using the relation

$$\cos \theta_e^* \simeq \frac{2m_{be}^2}{m_t^2 - M_W^2} - 1, \quad (50)$$

where m_{be} is the invariant mass of b and e^+ .

As shown in eq. (49), the QCD correction to the Born level f_L is small, so the distribution of m_{be} will not be sensitive to the QCD corrections and provides a good method to determine m_t . The distribution of m_{be} does not change even after including the effects of detector resolution by smearing the 4-momenta of final state partons.

5.2 Event Generator

As discussed in subsection 5.1, a “theoretical cut”, E_g^{th} was introduced, in order to separate the hard gluon bremsstrahlung from the soft gluon bremsstrahlung events in the event generator. Any physical quantity is insensitive to an appropriate choice of E_g^{th} , which should be smaller than m_b as suggested by the formalism given in the previous subsection. The value $m_b = 5$ GeV was assumed.

The decay process

$$t \rightarrow be^+\nu_e(g)$$

was first studied in the rest frame of the top quark. Two event generators were used in this study: one (referred to as EG4 below) is to generate the phase space for the 4-body events $t \rightarrow bf_1\bar{f}_2g$, and the other (referred to as EG3 below) is to generate the phase space for the 3-body events $t \rightarrow bf_1\bar{f}_2$.

If the energy of the gluon in the rest frame of the top quark is bigger than E_g^{th} , this event is called a hard gluon bremsstrahlung event, and its rate $\Gamma^{(4)}$ is calculated numerically by the event generator EG4. For an event that contains a gluon with energy less than E_g^{th} , which we call a soft gluon bremsstrahlung event, we use the analytic result eq.(45) to calculate the decay rate $\Gamma^{(3)}$. The kinematics of this event is given by the event generator EG3. This event, therefore, includes the Born contribution, the one loop radiative correction and the soft gluon bremsstrahlung effect.

5.3 Results

Numerical results were obtained for both $m_t = 140$ GeV and 200 GeV. The strategy followed was a natural one: one first obtains the various distributions in the rest frame of the top quark; thereafter, the same distributions are boosted up into the appropriate frames corresponding to the various present and future accelerators (TEVATRON, LHC, SSC) where top production is expected to take place. Here we briefly summarize the results in the rest frame of the top quark, omitting the details of the distribution to be found at the various accelerators. They can be found in the original publication (S. Mrenna and C.P. Yuan, Phys.

Rev. D??, !!! (1982).) For each particular $\sqrt{s}$, appropriate kinematic cuts have been introduced and smearing due to a finite detector resolution has been taken into account.

The distributions were generated by means of the event generator EG4. A *jet* is defined as a cone of energy with radius $R = \sqrt{(\Delta\phi)^2 + (\Delta\eta)^2}$ centered along the moving direction of the parton jet, where ϕ is the azimuthal angle ($|\Delta\phi| < \pi$) and η is the pseudorapidity. The R -separation between partons i and j is then $\Delta R_{ij} = \sqrt{(\Delta\phi_{ij})^2 + (\Delta\eta_{ij})^2}$, where $\Delta\phi_{ij} \equiv \phi_i - \phi_j$ and $\Delta\eta_{ij} \equiv \eta_i - \eta_j$.

The events with gluon energy $E_g > 20$ GeV and R -separation $\Delta R_{bg} > 0.7$ now define a hard gluon bremsstrahlung event, which differs from the previous definition by the introduction of the R -separation cut. (Note that E_g^{th} , taken to be 1 GeV, is the theoretical cutoff in separating the two event generators EG4 and EG3. It should not be confused with the cutoff $E_g > 20$ GeV used in defining a jet.) The branching ratio for a hard gluon bremsstrahlung event turns out to be less sensitive to m_t as the cone size ΔR increases. The results for $\Delta R = 0.0$ should be compared with those in reference [64] in which the QCD virtual corrections to the total decay width of the top quark were not included. Since the QCD corrections reduce the total width for $t \rightarrow be^+\nu_e(g)$, the branching ratios are higher than those in reference [64] from 6 to 9% as m_t varies from 140 to 200 GeV.

In addition to the hard gluon bremsstrahlung events, there are three other types of events in the rest frame of the top quark. The 4-body events are clasified according to their topology as follows.

- (i) $\Delta R_{bg} < 0.7$ and $E_g > 20$ GeV,
- (ii) $\Delta R_{bg} < 0.7$ and $E_g < 20$ GeV,
- (iii) $\Delta R_{bg} > 0.7$ and $E_g < 20$ GeV,
- (iv) $\Delta R_{bg} > 0.7$ and $E_g > 20$ GeV.

To see the contribution of these different types of events to the determination of the top quark mass, one first generates a Breit-Wigner shaped m_t distribution with the QCD corrected width from eq.(44). For a 140 (200) GeV top quark mass, the QCD corrected width is 0.61 (2.20) GeV. The event generator EG4 was then used to generate the kinematics of the decay products b , e^+ , ν_e and g in the rest frame of the top quark and calculate m_t , m_{be} and $\cos\theta_e^*$. Both m_{be} and $\cos\theta_e^*$ measure the degree of polarization of the W -boson from the top quark decay.

As it turns out, most of the gluon bremsstrahlung, event types (i) and (ii), is reconstructed as part of the b jet. However, the low energy gluons that are out of the b cone, event type (iii), generate a tail below the peak of the m_t distribution.

After including the QCD radiative corrections to the decay process $t \rightarrow bW^+$, we find that the m_t distribution no longer has the Breit-Wigner shape. Instead, the m_t distribution has a long tail in the smaller invariant mass region. However,

the peak position of the m_t distribution remains the same even though the average value of m_t is smaller. (The *intrinsic* width of the top quark cannot be observed via the QCD processes after considering detector resolution.) The distribution of m_{be} , which is a measure of the polarization of the W -boson, is found to be insensitive to this smearing. Therefore, the m_{be} distribution remains a powerful tool to test the polarization of the W -boson from top quark decay and determine the mass of the top quark in the Standard Model. Measuring the m_{be} distribution alone, however, cannot determine the top quark mass [60, 72]; non-Standard Model values for the form factors in the Lagrangian (47) would upset the relationship between m_t and the m_{be} distribution [60]. But, once m_t is measured from the invariant mass of $b\bar{q}_1 q_2$, m_{be} can be used to test the gauge coupling of t - b - W .

The effects of b -tagging and jet definition on the determination of the top quark mass and the polarization of the W -boson were also discussed. Needless to say, one should study how hadronization affects the conclusions, which are based on a parton level calculation with exact QCD amplitudes. This is in contrast to the usual full event generators, e.g. ISAJET, where the virtual corrections from QCD loops are not included and the amount of final state gluon radiation depends on some arbitrary cutoff scale.

6 Conclusion

We summarized the progress made under the grant during a period of (approximately) one year. Progress made in two areas of research is summarized:

- Astro particle physics
- Accelerator based phenomenology

In astro particle physics, the development of the Fourier method for the analysis of periodic signals represents a new method by means of which future, large data samples obtained on point sources can be analyzed. Likewise, the computation of the neutrino background paves the way towards the particle physics use of future, large area ($\approx 1\text{km}^2$) neutrino detectors. The simulation of "anomalous" EAS, in essence, served as a first step towards understanding the pattern of shower observables sensitive to physics beyond the Standard Model. It is clear that further studies are needed: in particular, one wants to begin a more systematic exploration of the response of an EAS to variations of the parameters of the initial interaction.

The research on accelerator based phenomenology has been initiated in the same spirit. The effective Lagrangian approach within the framework of gauge theories has been proposed earlier [73]. It was clear, however, that any "low energy" search for new physics requires a detailed knowledge of the relevant standard model effects: radiative corrections generate terms in the effective action similar to the ones the new physics does (at energies below its characteristic scale). In

many ways, as pointed out in [73] and elsewhere, t decays offer a fertile ground for the search of the effects of new physics.

Figure Captions

Fig. 1 Variation of the neutrino integral flux with respect to the galactic latitude ψ . The upper set of curves correspond to $E_{th} = 20\text{TeV}$, the lower set to $E_{th} = 100\text{TeV}$. In each set, the curve which has the greatest value at $\psi = 0$ corresponds to azimuthal angle $\phi = 0$, the next lower curve is $\phi = \pi/4$, and the third curve is $\phi = \pi/2$.

Fig. 2 The neutrino integral flux for $E_{th} = 100\text{TeV}$. a) is the flux resulting from primary cosmic radiation dominated by protons; b) assumes cosmic radiation dominated by Fe nuclei. Both curves represent the flux looking toward the center of the galaxy.

Fig. 3 Angular variation of the integral flux ($E_{th} = 100\text{TeV}$) of neutrinos produced on a spherical halo of baryonic dark matter. The angles are as in Fig. 1

Fig. 4 Angular variation of the total neutrino integral flux (halo plus galactic disk; $E_{th} = 100\text{TeV}$). The top plot is for azimuthal angle $\phi = 0$, the middle plot for $\phi = \pi/4$, and the bottom plot $\phi = \pi/2$. The dashed curves are the same as in Fig. 1, included for comparison.

Fig. 5 ν -induced shower: radial density of photons, electrons, hadrons and muons.

Fig. 6 μ density for ν and p induced showers.

Fig. 7 Average energy per μ for ν induced showers, p induced showers injected at 0 gcm^{-2} and p induced showers injected at 410 gcm^{-2} .

Fig. 8 Scatterplot of the ratio of the number of muons to electrons. Squares = ν induced showers; circles = p-induced showers

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