

Neutrinos via Charm Decays in Astrophysical Sources

by

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Abstract

Acknowledgements

Abbreviations

AGN	Active Galactic Nucleus
CKM	Cabibbo, Kobayashi & Maskawa
CMB	Cosmic Microwave Background
DSA	Diffusive Shock Acceleration
EW	Electroweak Theory
FF	Fragmentation Function
GR	General Relativity
GZK	Greisen, Sazepin & Kusmin
LO	Leading Order
NLO	Next to Leading Order
PDF	Parton Distribution Function
PMNS	Pontecorvo, Maki, Nakagawa & Sakata
QCD	Quantum Chromodynamics
QED	Quantum Electrodynamics
QFT	Quantum Field Theory
QM	Quantum Mechanics
SM	Standard Model
SR	Special Relativity
UHECR	Ultrahigh Energy Cosmic Ray

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1 Introduction

2 Background

While traversing space, messenger particles are subjected to many different influences, both in terms of type and magnitude, before being detected at earth. Understanding this propagation requires a thorough comprehension of physical processes at all scales of application, from production in astrophysical sources, through interactions with the vast radiation and matter fields that fill the cosmos, to entering the solar system and finally the terrestrial atmosphere.

The following sections provide an incomplete overview of aspects relevant to the treatment of this complex topic, as well as references for further study on the various related subjects.

2.1 Particle Physics

In the modern view, interactions and categories of elementary particles are most accurately described by a construct called the *Standard Model* (SM) of particle physics. Underlying this formalism is the mathematical framework known as *Quantum Field Theory* (QFT) that combines and generalizes properties of *Quantum Mechanics* (QM) and *Special Relativity* (SR) to produce an extremely precise description of microscopic reality. Excitations in the associated fields are referred to as quanta and manifest as observable particles.

Since this work is mostly concerned with the statistical behaviour of large quantities instead of individual particle probabilistics, a more detailed explanation is omitted except for certain phenomena important to the justification of some later assumptions.

Fundamental Interactions

Electromagnetism is described by *Quantum Electrodynamics* (QED) and mediated by photons γ with Abelian $U(1)$ symmetry, whereas the weak force is carried by Z and W^\pm bosons related to the $SU(2)$ group. In terms of *Electroweak Theory* (EWT) all these particles actually result from symmetry breaking of virtual isospin and hypercharge fields, unifying $U(1) \times SU(2)$ under a single coherent structure. The previously massless gauge bosons then gain their masses via the Higgs mechanism, where only the γ is excluded and remains without mass.

Notably, the weak interaction exclusively couples to left handed fermions, violating parity.

Chirality, Helicity, Decay, Mixing, Quarks, Color Charge

Carriers of the strong force are named gluons and arise from the SU(3) symmetry group.

Weak, Electromagnetic, Electroweak, Higgs

Parity Violation, Mixing

Limits, Dark Matter, Dark Energy, Neutrino Oscillation, Quantum Gravity

Vector Bosons, Scalar Boson, Boson Spins, Fermion Spins

General Relativity

Leptons

Neutrinos, Decays

Hadrons

Confinement, Asymptotic Freedom, Residual Nuclear Force, Decays

2.2 Multimessenger Astronomy

Gravitational Waves

Cosmic Rays

Photons

Neutrinos

2.3 Astrophysical Sources

Magnetars

Active Galactic Nuclei

3 Methods

4 Results

5 Discussion

6 Conclusion & Outlook

Appendix

A Reference Frames

Depending on the application, energies in particle physics are either given as viewed from a suitable rest frame or independent from the choice of coordinate system altogether. One widely adopted formulation uses the Mandelstam variables

$$s = (p_1 + p_2)^2 \quad t = (p_1 - p_3)^2 \quad u = (p_1 - p_4)^2$$

to assign different channels in scattering processes via the squared momentum carried by the exchanged mediating particle. Implied in this context is the Minkowski inner product, making the above quantities manifestly Lorentz invariant.

When working with parametrizations defined for use in different subdisciplines it often becomes necessary to convert from center of mass energies \sqrt{s} to the energy E of a projectile in the target rest frame. With $E^2 = \mathbf{P}^2 c^2 + M^2 c^4$ as well as momenta $P = (E, \mathbf{P}c)$ and $p = (mc^2, 0)$ one finds

$$s = (P + p)^2 = (E + mc^2)^2 - \mathbf{P}^2 c^2 = 2Emc^2 + m^2 c^4 + M^2 c^4$$

for the invariant mass. This relation is typically approximated as $s = 2Emc^2$ at high energies.

B Cross Sections

By defining an effective area perpendicular to the velocity vectors of projectiles and targets, cross sections measure the probability of collision processes in particle physics. Due to depending on the strength of an interaction, these quantities generally scale with energy. Distinct from the integrated case, differential cross sections are usually given with respect to some independent variable such as angle or momentum of the particle.

Scattering

To model total cross sections in hadron proton scattering, this work uses the formula

$$\sigma_{hp} = H_h \ln^2(s/s_h) + P_h + R_h^1 (s_h/s)^{\eta_1} + R_h^2 (s_h/s)^{\eta_2} \quad (1)$$

as given in [1] for a universal analytic parametrization of the corresponding amplitudes.

All adjustable parameters are listed in table 1 together with relevant meson lifetimes for cooling. In this approach, the variable M relates to $H = \pi(\hbar c/M)^2$ and $s_h = (m_h + m_p + M)^2$ as an effective mass. Coefficients in (1) denote Heisenberg, Pomeranchuk and Regge terms which have some qualitative motivation, though the formula itself is primarily a quantitative result.

TABLE 1: Fits to the total inclusive scattering cross sections in hadron proton collisions. Parameters are taken from [1] with $M = 2.121$ GeV for $H = 0.272$ mb as the rate of growth. Both $\eta_1 = 0.447$ and $\eta_2 = 0.5486$ are dimensionless exponents. Decay times τ_h and rest masses m_h can be found in the particle listings [2] where the latter are given in natural units.

h	P_h / mb	R_h^1 / mb	R_h^2 / mb	τ_h / ns	m_h / GeV	s_h / GeV^2
p	34.41	13.07	7.39		0.938	15.98
π	18.75	9.56	1.767	26.03	0.140	10.23
K	16.36	4.29	3.408	12.38	0.494	12.62

Assuming a quasi universal ratio \mathcal{R} between elastic and total hadron cross sections, one obtains the inelastic cross section $\sigma_{\text{inel}} = (1 - \mathcal{R})\sigma_{\text{tot}}$ from $\sigma_{\text{el}} = \mathcal{R}\sigma_{\text{tot}}$ and $\sigma_{\text{el}} + \sigma_{\text{inel}} = \sigma_{\text{tot}}$ as a unitarity condition. Provided in [3] is the model independent parametrization

$$\mathcal{R}(s) = \frac{\sigma_{\text{el}}(s)}{\sigma_{\text{tot}}(s)} = \mathcal{A} \tanh(\gamma_1 - \gamma_2 \ln(s) + \gamma_3 \ln^2(s)) \quad (2)$$

with a constant asymptote \mathcal{A} at very high energies. Coefficients are given in table 2 for different physical settings. Both equations (1) and (2) use units of GeV^2 for the s variables.

TABLE 2: Almost model independent ratio of hadronic elastic and total scattering cross sections. Factors γ are taken from [3] for varying \mathcal{A} asymptotes.

\mathcal{A}	γ_1	γ_2	γ_3
1/2	0.466	0.0259	0.00177
1	0.2204	0.0111	0.00076

Reference [4] tests the asymptotic rise $\sigma(s) \propto \ln^2(s)$ derived in [5] as a theoretical upper bound and concludes that it is somewhat exceeded. Additionally, a ratio $\mathcal{A} = 1/3$ due to diffraction as opposed to the black disc limit $\mathcal{A} = 1/2$ from optical theorem predictions is suggested. Because parameters are only available in the latter case, all calculations of \mathcal{R} use function (2) as defined by an asymptote $\mathcal{A} = 1/2$ for this work. Data matching $\mathcal{A} = 1/3$ then implies underestimated values of $\sigma_{\text{inel}}(s)$ which should however not significantly influence the overall results.

Production

For charm quark production in proton air collisions, reference [6] gives

$$x_F \frac{d\sigma}{dx_F}(x_F, E_p) = ax_F^b (1 - x_F^m)^n$$

as the parametrized differential cross section with components

$$a = a_1 - a_2 \ln(E_p) \quad b = b_1 - b_2 \ln(E_p) \quad n = n_1 - n_2 \ln(E_p)$$

for which table 3 lists all necessary constants. Here proton energies E_p are defined as viewed by air nuclei at rest, while the Feynman scaling variable $x_F = p_c/p_s$ specifies magnitude ratios of produced charm quark longitudinal momentum to all available momentum in center mass coordinates of the colliding particles. Application of appendix A shows that this approximately fulfills $x_F = x_c$ where $x_c = E_c/E_p$ in the relevant energy ranges.

TABLE 3: Parametrization of the inclusive charm quark production differential cross section. Coefficients are calculated from [6] to write E_p in units of GeV without needing redundant conversion steps. The number $m = 1.2$ is a constant at all energies. For the application at hand, energy ranges beyond the given validity intervals might be used when mentioned in the text.

E_p / GeV	a_1 / μb	a_2 / μb	b_1	b_2	n_1	n_2
$10^4 - 10^8$	826	8411	0.197	0.016	1.061	0.107
$10^8 - 10^{11}$	403	2002	0.237	0.023	7.639	0.102

As in [7] it is assumed that the cross section scales linearly with nucleon number, yielding

$$\frac{d\sigma}{dx_c}(x_c, E_p) = A^{-1} \frac{d\sigma}{dx_F}(x_c, E_p)$$

for inclusive charm production in proton proton collisions. Approximating air as a gas mixture of roughly 75 % nitrogen and 25 % oxygen, one finds $A = 14.5$ for this scaling. Translation of charm quarks to charmed hadrons is achieved with a folding integral

$$\frac{d\sigma}{dx_h}(x_h, E_p) = \int_{x_h}^1 dz z^{-1} \frac{d\sigma}{dx_c}(x_c, E_p) D_c^h(z) \quad (3)$$

where $z = E_h/E_c$ and $x_h = E_h/E_p$ as well as $x_c = x_h/z$ are fractional energies. Limits for the integration follow from a basic inequality $E_h \leq E_c \leq E_p$ to incorporate kinematic constraints. Furthermore, the probability of observing any final state h originating from a c quark is encoded in a *Fragmentation Function* (FF) $D_c^h(z)$ dependent on the fraction of hadron to charm energy. Reference [8] addresses the connection between this concept and that of a *Parton Distribution Function* (PDF) among other things.

Where a PDF represents the probability density of finding a parton with given momentum in a color neutral particle, probabilities for color neutral states existing inside individual partons are given by the appropriate FF instead. The partons described here are either quarks or gluons, which can be free only asymptotically at high energies due to carrying color charges. In this limit, the running coupling of QCD is small enough for a power series expansion to be a sensible approach, leading to the definition of terms such as LO and NLO in reference to exponent order. There exist different factorization methods to separate these parts from the nonperturbative contributions contained in any PDF and FF for the confined constituents of hadrons. By fitting to existing data or perturbative results, models can extrapolate to low momentum fractions that have not yet been probed experimentally. A similar procedure has lead [9] to obtain

$$D_c^h(z) = \frac{N_h z(1-z)^2}{((1-z)^2 + \epsilon_h z)^2} \quad (4)$$

with parameters from e^+e^- data in table 4 as the charm hadron FF used throughout this work. It is important to note that such functions are invariant under charge conjugation, so that there is no differentiation between quark to particle or antiquark to antiparticle processes.

C Spectral Distributions

Constructing spectra dN_h/dE_h from proton injection requires folding of $F_{p \rightarrow h}$ as the hadronic distribution for a single pp interaction with the number of protons per energy interval given by the function dN_p/dE_p obtained from source specific modelling. An analogous approach can be applied to compute neutrino spectra dN_ν/dE_ν from dN_h/dE_h via distributions $F_{h \rightarrow \nu}$ formulated according to the involved decay modes.

Charm

Spectral distributions for charmed hadron production are calculated according to [10] through

$$F_{p \rightarrow h}(E_h, E_p) = E_p^{-1} \sigma_{pp}^{-1}(E_p) \frac{d\sigma}{dx_h}(x_h, E_p)$$

with $E_h = x_h E_p$ translating between variables. This formula can be understood as a normalization of (3) against proton energies and the inelastic pp cross section from appendix B to yield hadron numbers per unit energy.

To find neutrino spectra from charmed hadrons, the same approach as in [10] is used, which assumes an effective energy distribution approximated by three body decays for the semileptonic channel to a less massive pseudoscalar meson. By neglecting lepton masses, one obtains

$$\tilde{F}_{h \rightarrow \nu}(y) = D_h^{-1} \left(6b_h a_h^2 - 4a_h^3 - 12\lambda_h^3 a_h + 12\lambda_h^2 y - 6b_h y^2 + 4y^3 + 12\lambda_h^2 \ln((1-y)/\lambda_h) \right)$$

as a distribution with $y = E_\nu/E_h$ and $F_{h \rightarrow \nu}(E_\nu, E_h) = \tilde{F}_{h \rightarrow \nu}(y)/E_h$ for conversion.

Hadron specific coefficients for this equation are defined with the parameter $\lambda_h = \tilde{s}_h/m_h^2$ as

$$a_h = 1 - \lambda_h \quad b_h = 1 - 2\lambda_h \quad D_h = 1 - 8\lambda_h - 12\lambda_h^2 \ln(\lambda_h) + 8\lambda_h^3 - \lambda_h^4$$

where both s_h and m_h are listed in table 4 for all included charmed hadrons. The assumption of three body decays like $D^+ \rightarrow \bar{K}^0 e^+ \nu_e$ can be justified by consulting [2] for information on the relevant particles and comparing branching ratios, which indicate that purely leptonic modes are strongly suppressed. Hadronic channels such as $D^+ \rightarrow \pi^+ \pi^0$ or $D^+ \rightarrow K^- \pi^+ \pi^+$ are either very improbable as well or occur at significant rates but do not contribute many high energy neutrinos due to pions and kaons being subject to further cooling before decaying to leptons. By the same logic, secondary muon decay is neglected when determining the neutrino spectrum.

TABLE 4: Coefficients for charm hadron production, cooling and decay to neutrinos.

All parameters ϵ_h are taken from fits to LO QCD results via the FF as defined and described in [9] with normalizations N_h given by [10] to rescale the integration of (4) over $[0,1]$ to approximately match the fractions f_h provided in [11] from measurements. Effective masses $\sqrt{\tilde{s}_h}$ have been determined by [12] from fitting decay rates, except for D_s^+ which is copied from [10] because no original value could be found. Mean lifetimes τ_h and rest masses m_h are adopted from [2] in the particle listings. For mass type quantities, natural units are used.

h	N_h	ϵ_h	τ_h / fs	$\sqrt{\tilde{s}_h} / \text{GeV}$	m_h / GeV
D^0	0.577	0.101	410	0.67	1.86
D^+	0.238	0.104	1033	0.63	1.87
D_s^+	0.0327	0.0322	501	0.84	1.97
Λ_c^+	0.0067	0.00418	203	1.27	2.29

Pions & Kaons

By parametrizing event generator results, a neutral pion production spectrum of the form

$$\tilde{F}_\pi(x_\pi, E_p) = 4\alpha B x_\pi^{\alpha-1} \left(\frac{1 - x_\pi^\alpha}{1 + r x_\pi^\alpha (1 - x_\pi^\alpha)} \right)^4 \left((1 - x_\pi^\alpha)^{-1} + \frac{r(1 - 2x_\pi^\alpha)}{1 + r x_\pi^\alpha (1 - x_\pi^\alpha)} \right) \left(1 - \frac{m_\pi}{x_\pi E_p} \right)^{1/2}$$

is found in [13] with $m_\pi = 0.135 \text{ GeV}$ [2] translated to natural units and parameters

$$B = \tilde{B} + C \quad \alpha = \frac{\tilde{\alpha}}{\sqrt{C}} \quad r = \frac{\tilde{r}}{\sqrt{C}}$$

where a low energy cutoff is enforced via $E_\pi = x_\pi E_p$ in the mass term. From

$$C = c_1 - c_2 \ln(E_p) + c_3 \ln^2(E_p)$$

results a dependence on projectile energy for the shape of this distribution.

Coefficients are specified in table 5 and recalculated for E_p in GeV instead of TeV units. Under the assumption of a π^0 cross section approximately equal to the π^\pm average and with identical spectra for charged pions, it follows that $F_\pi = \tilde{F}_\pi/E_p$ should describe pion production regardless of charge reasonably well for the purpose of this work.

TABLE 5: Parametrized spectral distribution for neutral pion production. Factors are taken from [13] and converted to write E_p in units of GeV for c_k coefficients.

\tilde{B}	$\tilde{\alpha}$	\tilde{r}	c_1	c_2	c_3
0.25	0.98	2.6	1.515	0.206	0.075

For a convenient formulation of kaon production, references [14] and [15] indicate a constant ratio π/K at moderately high energies. Similar fractions are retrieved from multiplicities given in [16] and lead to $F_K/F_\pi = 0.12$ as a simplifying assumption, the validity of which cannot be guaranteed for the application at hand. Calculations of kaon spectra still employ this approach but are subject to considerable reservations as a result.

Decays of pions and kaons to neutrinos are approximated via the $h \rightarrow \mu^+ \nu_\mu$ two body channel with branching fractions of $F_\pi = 99.99\%$ and $F_K = 63.56\%$ given in the [2] particle listings. By decaying, muons produced in these processes can significantly impact the neutrino spectrum. Results from [10] suggest that this is particularly relevant for pions. Muonic three body decays of type $\mu^- \rightarrow e^- \bar{\nu}_e \nu_\mu$ as well as cooling factors depend on the polarization of participating leptons due to the nature of weak force coupling. This complicates computations and is thus omitted in service of restricting the present work to a manageable scope, though it should be remembered as an important caveat for the final results.

The remaining two body decays of ultrarelativistic hadrons h to leptons l obey a distribution

$$F_{h \rightarrow l}(E_l, E_h) = F_h E_h^{-1} (1 - \lambda_h)^{-1}$$

with $m_\nu = 0$ and $m_\mu = 0.106$ GeV [2] as well as $\lambda_h = m_\mu^2/m_h^2$ as a parameter. This formula is the same whether $l = \nu$ or $l = \mu$ because there is one muon for each neutrino. In addition, kinematic considerations lead to $E_\mu/E_h > \lambda_h$ and $E_\nu/E_h < 1 - \lambda_h$ for integral bounds $E_\mu < E_h < E_\mu/\lambda_h$ in the case of muons or $E_\nu/(1 - \lambda_h) < E_h < E_p$ when considering neutrinos.

D Cooling & Decay

Consider an infinitesimally thin slice of a target medium with particle number density n and volume $V = S dx$ where S and dx measure surface area and thickness, respectively. Accordingly, there exist $\tilde{N} = nV$ targets that each have effective interaction cross sections $\tilde{\sigma}$ with a total coverage of $\tilde{S} = \tilde{\sigma} \tilde{N}$ as viewed by an incident projectile.

The probability for stopping such a beam constituent then corresponds to the ratio $\mathcal{P} = \tilde{S}/S$ of both areas or explicitly $\mathcal{P} = \tilde{\sigma} n dx$ in case of dx as the covered distance. Expressing $\tilde{\sigma} = \kappa \sigma$ in terms of the inelastic scattering cross section σ and a dimensionless factor κ called inelasticity, one can identify a length scale $\lambda = (\kappa \sigma n)^{-1}$ as the mean free path between collisions. Multiplying with κ includes the ratio of remaining to initial energy, which is taken to be constant. From this follows a reduction in beam particles $dN = -N \lambda^{-1} dx$ proportional to N as the total projectile count and $\mathcal{P} = \lambda^{-1} dx$ for the reformulated probability which represents an ordinary differential equation of first order. The solution is found to follow an exponential law

$$N(x) = N_0 \exp\left(-\frac{x}{\lambda}\right)$$

where N particles remain over some distance x with N_0 as the initial amount. Furthermore,

$$P(x) = 1 - \exp\left(-\frac{x}{\lambda}\right)$$

gives the probability of a particle having been scattered after travelling x length units. Similar steps for time instead of distance lead to the well known equation

$$N(t) = N_0 \exp\left(-\frac{t}{\tau}\right)$$

describing exponential decay. It commonly appears in the context of radioactive materials but also applies to hadrons and leptons or more generally any quantity which decreases at a rate proportional to itself. Analogous to the previous case, a particle decays with probability

$$P(t) = 1 - \exp\left(-\frac{t}{\tau}\right)$$

before a time t has passed and for τ as the mean lifetime. Translating this from rest frame to laboratory coordinates defines the decay timescale $t_{\text{dec}} = \tau \Gamma$ with a Lorentz factor $\Gamma = E/m$ via projectile energy and invariant mass. This is equivalent to a characteristic decay length given by $\lambda_{\text{dec}} = v t_{\text{dec}}$ where the velocity $v = c$ can be set for highly relativistic particles. Additionally, mean free path and cooling distance $\lambda_{\text{cool}} = (\kappa \sigma n)^{-1}$ refer to exchangeable concepts. Particles lose energy in every collision, which is the same as reducing temperature from a thermodynamics perspective. Dividing by the speed of light c translates this expression to $t_{\text{cool}} = (\kappa \sigma n c)^{-1}$ as a cooling timescale. Analogously, the distance $\lambda_{\text{dec}} = c \tau E/m$ can be rewritten as $t_{\text{dec}} = \tau E/m$ in units of time. Substituting into the decay formula yields a cooling factor

$$C = 1 - \exp\left(-\frac{t_{\text{cool}}}{t_{\text{dec}}}\right) \quad (5)$$

which scales spectra from direct production to account for decay processes taking place after collisional energy losses have occurred. This is in some sense a core mechanism for the hypothesis that neutrinos from charm dominate pion and kaon contributions in high energy regimes. While longer lived particles experience significant cooling due to time dilation, charmed hadron decay is prompt in comparison. One requirement for the validity of (5) is that $\lambda_{\text{dec}} \ll d$ holds with a target field size d to ensure decays occur exclusive inside this region.

Appendix

E High Energy Cutoff

[CMB GZK \[17\]](#)

F Magnetic Field Scales

[18] criterion

G Stochastic Acceleration

Due to its wide applicability in different astrophysical scenarios, probabilistic collisions are often viewed as one of the more plausible mechanisms responsible for accelerating cosmic rays to high energies. The general case is described in [19] and supposes that for each collision, particles gain energy proportional to a constant factor η and remain in the region of acceleration with fixed probability ς on average. With initial conditions N_0 for the particle number and E_0 as the mean energy, this results in $N = N_0 \varsigma^k$ and $E = E_0 \eta^k$ after k collisions. Using $\ln x^k = k \ln x$ in

$$\frac{\ln(N/N_0)}{\ln(E/E_0)} = \frac{\ln(\varsigma)}{\ln(\eta)}$$

eliminates the exponent and by rearranging gives the relation

$$N = N_0 \left(\frac{E}{E_0} \right)^{\ln(\varsigma)/\ln(\eta)}$$

connecting energy and number of particles. This integrated spectrum incidentally follows a power law, which is an almost ubiquitous feature observed in cosmic ray physics. One obtains

$$\frac{dN}{dE} = \frac{N_0}{E_0} \left(\frac{E}{E_0} \right)^\alpha$$

for the differential spectrum where the spectral index

$$\alpha = \frac{\ln(\varsigma)}{\ln(\eta)} - 1$$

is constrained by $\ln(\varsigma)/\ln(\eta) < 0$ due to $\varsigma < 1$ and $\eta > 1$ as implied per the definitions.

The basic case of [DSA](#) considers strong shock fronts moving with velocity $\beta = v/c$ in a fully ionized gas. Requiring momentum isotropization without significant energy losses on both sides of the discontinuity results in $\ln(\varsigma)/\ln(\eta) = -1$ for a $dN/dE \propto E^{-2}$ spectral dependence that is discussed by [19] as well. A slightly steeper index around $\alpha = -2.5$ can be derived when nonlinear effects are accounted for to better reproduce measurements.

Energy gain increasing linearly with β leads this mechanism to be categorized as Fermi type acceleration of first order, whereas the originally proposed formulation scales like β^2 or as second order. Though shocks exceed the local speed of sound in the astrophysical medium, relativistic velocities are typically not achieved. Consequently, ratios $\beta \ll 1$ mean that lower order processes are much more efficient in reaching high particle energies.

H Pulsar Spindown

I Accretion Discs

[19]

Hydrostatic Equilibrium

Luminosity Limit

J Particle Theory

K Implementation

In order to calculate neutrino spectra from hadronic distributions, several integrals have to be computed. Discretizing this task allows the general case

$$F(x, y) = \int_{z_-}^{z_+} dz G(x, z) H(z, y)$$

to be rewritten as a Riemann sum. Assuming G and H are integrable over a given interval,

$$F_{ij} = \sum_k D_{kk} G_{ik} H_{kj}$$

converges to the exact solution for sufficiently small steps. Transforming variables

$$x \rightarrow x_i \quad y \rightarrow y_j \quad z \rightarrow z_k$$

and defining $D_{kk} = z_{k+1} - z_k$ leads to the above notation. It is easily shown how this expression in terms of indices translates to the product of corresponding matrices

$$\mathbf{F} = \mathbf{G} \mathbf{D} \mathbf{H}$$

as an equivalent formulation. Here the output $\mathbf{F} \in \mathbb{R}^{m \times n}$ is obtained from the inputs $\mathbf{G} \in \mathbb{R}^{m \times l}$ and $\mathbf{H} \in \mathbb{R}^{l \times n}$ as well as the square matrix $\mathbf{D} \in \mathbb{R}^{l \times l}$ that encodes all step sizes on its diagonal. These results enable a quick and efficient implementation of the required calculations as program code, where array arithmetic operations can greatly increase execution speed.¹

¹ In service of reproducibility, all implementations can be viewed in [this](#) repository.

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