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# Distinguishing neutrino mass hierarchies using dark matter annihilation signals at IceCube

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**Abstract.** We explore the possibility of distinguishing neutrino mass hierarchies through the neutrino signal from dark matter annihilation at neutrino telescopes. We consider a simple extension of the standard model where the neutrino masses and mixing angles are obtained via the type-II seesaw mechanism as an explicit example. We show that future extensions of IceCube neutrino telescope may detect the neutrino signal from DM annihilation at the Galactic Center and inside the Sun, and differentiate between the normal and inverted mass hierarchies, in this model.

**Keywords:** dark matter theory, neutrino theory

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### 1 Introduction

Although various lines of evidence support the existence of dark matter (DM) in the universe, its identity remains as a mystery. Among various proposed candidates, the weakly interacting massive particles (WIMPs) are promising ones that typically arises in many extensions of the standard model (SM) [1, 2]. Major direct, indirect, and collider experimental searches are currently underway to detect the particle nature of DM and determine its properties.

Just like DM, the origin of neutrino mass and mixing defines a very interesting area of investigation beyond the SM (BSM) [3, 4]. These are the two encouraging avenues of new physics and a large number of BSM scenarios have already been proposed to address them. It will be even more interesting to investigate models where the two sectors are connected. In fact, such a connection arises naturally in a class of models where DM is tied to the neutrino sector. For example, in type-II seesaw plus a singlet scalar scenario [5], it has been shown that the neutrino mass hierarchy can influence the DM annihilation rate to charged leptons as well as neutrinos and subsequently provide a possible explanation of the positron fraction excess observed at the AMS-02 experiment [6]. Similar studies have been done in the  $B - L$  extension of the minimal supersymmetric standard model (MSSM) where it is shown that the annihilation of the right-handed sneutrino DM produces neutrino final states with a distinct feature [7]. There have been studies on the possibility of discovering the  $B - L$  sneutrino DM at the LHC [8–10]. The connection between the DM and neutrinos in such scenarios can also be probed through DM direct detection [11] and some indirect detection experiments [12]. Therefore, a combination of direct, indirect and collider signatures could be useful to explore this connection (for a review, see [13]).

Our focus here is on the DM indirect detection signals as a means of testing the connection between the DM and neutrinos. Neutrino telescopes like IceCube are able to examine these models through the signal arising from the DM annihilation into neutrinos at the Galactic Center [14] and inside the Sun [15]. Unlike the photon flux, the neutrino flux arising from the DM annihilation at the Galactic Center has less astrophysical background uncertainty, which allows us to probe the exact nature of the DM more accurately. However, the model prediction for the photon fluxes coming from the DM annihilation at the Galactic Center should be consistent with the Fermi-LAT [16] data on the gamma-ray flux. Therefore, it will be interesting to combine the results from IceCube and Fermi-LAT to study models where the DM and the neutrino sectors are interconnected.

In this paper, we study these issues within the above mentioned extension of the SM [5, 6]. In this model, the light neutrino masses and mixing angles are generated using the well-known type-II seesaw mechanism [17–21] that introduces  $SU(2)_L$  triplet scalars. The neutral component of the triplet acquires a non-zero vacuum expectation value (vev) that generates tiny neutrino masses by breaking lepton number by two units and mixing among different neutrino flavors. Among the SM particles, these triplet scalars only couple to gauge bosons through gauge couplings and to leptons via Yukawa couplings. These Yukawa couplings are related to the triplet vev, which in turn determines different decay modes of these triplet scalars. It has been observed [22, 23] that the triplets will dominantly decay to a pair of leptons if the vev is less than 0.1 MeV, otherwise the gauge boson final state will become dominant. In order to accommodate a stable DM candidate, the model is augmented by a SM singlet scalar with a  $Z_2$  symmetry [5]. In the minimal case, the DM only couples to the Higgs and to the triplet scalars. For a vev smaller than 0.1 MeV, the annihilation of the DM produces a pair of triplets that will further decay to SM leptons (including neutrinos) with almost 100% branching fraction. Therefore, the flavor composition of the final states will depend upon the neutrino mass hierarchy that sets the Yukawa couplings. We exploit this feature and show how different neutrino mass hierarchies can be distinguished by using the muon signal arising from conversion of muon neutrinos from DM annihilation at the Galactic Center and inside the Sun at IceCube. We note that the photon signal from the Galactic Center does not discriminate amongst different hierarchies because the difference will not be significant.

The paper is organized as follows. In section 2, we discuss the model in brief. In section 3, we present the analysis of the neutrino signal from DM annihilation at the Galactic Center and inside the Sun for normal and inverted hierarchies. We also calculate the corresponding muon spectra at the detector and compare them with the background. Finally, in section 4 we present our discussion and conclusion.

## 2 Type-II seesaw with a scalar singlet DM candidate

The minimal version of the type-II seesaw model is an extension of the SM that includes a  $SU(2)_L$  scalar triplet  $\Delta$  with hypercharge  $Y = 2$ .

$$\Delta = \begin{pmatrix} \frac{\Delta^+}{\sqrt{2}} & \Delta^{++} \\ \Delta^0 & \frac{\Delta^-}{\sqrt{2}} \end{pmatrix}. \quad (2.1)$$

In order to accommodate a cold dark matter (CDM) candidate, we introduce a SM singlet real scalar  $D$ . The stability of the DM candidate is ensured by imposing a  $Z_2$  symmetry under which  $D$  is odd while all other SM particles and the triplet  $\Delta$  are even.

The scalar potential of the model, including the relevant terms for the DM candidate  $D$ , is given by [6]

$$\begin{aligned} \mathcal{V}(\Phi, \Delta) = & -m_\Phi^2(\Phi^\dagger \Phi) + \frac{\lambda}{2}(\Phi^\dagger \Phi)^2 + M_\Delta^2 \text{Tr}(\Delta^\dagger \Delta) + \frac{\lambda_1}{2} \left[ \text{Tr}(\Delta^\dagger \Delta) \right]^2 \\ & + \frac{\lambda_2}{2} \left( \left[ \text{Tr}(\Delta^\dagger \Delta) \right]^2 - \text{Tr} \left[ (\Delta^\dagger \Delta)^2 \right] \right) + \lambda_4(\Phi^\dagger \Phi) \text{Tr}(\Delta^\dagger \Delta) \\ & + \lambda_5 \Phi^\dagger [\Delta^\dagger, \Delta] \Phi + \left( \frac{\Lambda_6}{\sqrt{2}} \Phi^\dagger i\sigma_2 \Delta^\dagger \Phi + \text{h.c.} \right), \end{aligned} \quad (2.2)$$

$$\mathcal{V}_{\text{DM}}(\Phi, \Delta, D) = \frac{1}{2} m_D^2 D^2 + \lambda_D D^4 + \lambda_\Phi D^2(\Phi^\dagger \Phi) + \lambda_\Delta D^2 \text{Tr}(\Delta^\dagger \Delta), \quad (2.3)$$

where  $\Phi$  is the SM Higgs doublet. The couplings  $\lambda_i$  ( $i = 1 - 5$ ) can be taken as real without any loss of generality. Due to spontaneous symmetry breaking, the non-zero vev of the Higgs doublet generates a tadpole term through interaction involving the  $\Lambda_6$  coupling in eq. (2.2). This in turn induces a non-zero vev ( $v_\Delta$ ) for the neutral component of the triplet thereby breaking lepton number by two units. This triplet vev contributes to the SM gauge boson masses at the tree level which leads to a deviation in the electroweak  $\rho$  parameter. Hence,  $v_\Delta$  is constrained by the electroweak precision data that requires the  $\rho$  parameter to be close to the SM value of unity [24] which eventually demands  $v_\Delta < 2$  GeV [25]. In addition to this, current experimental bounds on lepton flavor violating processes put a lower bound on  $v_\Delta$  [26–28] as

$$v_\Delta M_\Delta \geq 150 \text{ eV GeV}. \quad (2.4)$$

It is to be mentioned here that we assume negligible mixing between the doublet and triplet scalars in our analysis<sup>1</sup> and with this assumption, the DM potential of eq. (2.3) can be expressed in terms of physical scalars ( $h, H^\pm, H^{\pm\pm}, H^0, A^0$ ) as

$$\begin{aligned} \mathcal{V}_{\text{DM}}(\Phi, \Delta, D) = & \frac{1}{2} m_{\text{DM}}^2 D^2 + \lambda_D D^4 + \lambda_\Phi v D^2 h + \frac{1}{2} \lambda_\Phi D^2 h^2 \\ & + \lambda_\Delta D^2 \left[ H^{++} H^{--} + H^+ H^- + \frac{1}{2} (H_0^2 + A_0^2) + \frac{1}{2} v_\Delta H_0 \right]. \end{aligned} \quad (2.5)$$

Here,  $m_{\text{DM}}^2 = m_D^2 + \lambda_\phi v^2 + \lambda_\Delta v_\Delta^2$  denotes the physical mass of the DM candidate. It is evident from eq. (2.5) that the DM candidate can annihilate to a pair of Higgs and to a pair of exotic triplet scalar particles through the coupling parameters  $\lambda_\phi$  and  $\lambda_\Delta$ . In the limit  $m_{\text{DM}} > m_\Delta$ , where,  $m_\Delta$  is the degenerate mass of all triplet scalars, the annihilation cross section of the DM candidate (non-relativistic) is expressed as

$$\langle \sigma v \rangle = \frac{1}{16\pi m_{\text{DM}}^2} \left[ \lambda_\phi^2 \left( 1 - \frac{m_h^2}{m_{\text{DM}}^2} \right) + 6\lambda_\Delta^2 \left( 1 - \frac{m_\Delta^2}{m_{\text{DM}}^2} \right) \right]. \quad (2.6)$$

This should yield the correct relic abundance of the DM particle that lies within the  $3\sigma$  limit of current Planck bound  $\Omega_{\text{DM}} h^2 = 0.1199 \pm 0.0027$  [30]. Now, the final state of DM annihilation will depend upon the branching fraction of the Higgs and the triplet scalar decay. The triplet scalars can couple to the SM gauge bosons through gauge coupling and to SM lepton doublets through Yukawa couplings.

$$\mathcal{L}_Y = \mathcal{L}_Y^{\text{SM}} - \frac{1}{\sqrt{2}} (Y_\Delta)_{ij} L_i^\top C i \sigma_2 \Delta L_j + \text{h.c.} \quad (2.7)$$

Here,  $C$  is the charge conjugation operator and  $L_i = (\nu_i, \ell_i)_L^T$  is the  $\text{SU}(2)_L$  lepton doublet with  $i$  being the three generation indices. Further, these Yukawa couplings can be obtained from the Majorana mass matrix of neutrinos that arises due to the non-zero triplet vev  $v_\Delta$ .

$$(M_\nu)_{ij} = v_\Delta (Y_\Delta)_{ij}, \quad (2.8)$$

$$Y_\Delta = \frac{M_\nu}{v_\Delta} = \frac{1}{v_\Delta} U^\top M_\nu^{\text{diag}} U, \quad (2.9)$$

where  $M_\nu^{\text{diag}} = \text{diag}(m_1, m_2, m_3)$  and  $U$  is the PMNS mixing matrix.

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<sup>1</sup>Detailed expressions for physical scalar mass eigenstates can be found in [29].

With the recent global analysis data (see eq. (A.6) in the appendix), and after using eq. (2.9), we obtain the following structure of the Yukawa coupling matrix for the normal hierarchy (NH) and inverted hierarchy (IH) scenarios

$$Y_\Delta = \frac{10^{-2} \text{ eV}}{v_\Delta} \times \begin{cases} \begin{pmatrix} 1.08 - 0.29i & -1.55 + 0.09i & 1.23 - 0.31i \\ -1.55 + 0.09i & 2.07 + 0.26i & -1.59 - 0.21i \\ 1.23 - 0.31i & -1.59 - 0.21i & 2.59 + 0.0i \end{pmatrix} & (\text{normal hierarchy}) \\ \begin{pmatrix} 3.84 + 0.34i & 1.21 - 0.13i & -1.39 - 0.94i \\ 1.21 - 0.13i & 2.97 - 0.35i & 1.98 - 0.65i \\ -1.39 - 0.94i & 1.98 - 0.65i & 2.66 + 0.01i \end{pmatrix} & (\text{inverted hierarchy}). \end{cases} \quad (2.10)$$

To find the Yukawa matrix for the degenerate case, we impose the 95% C.L. limit on the sum of all light neutrino masses  $\sum_i m_i < 0.23 \text{ eV}$  from Planck [30], which uses the Baryon Acoustic Oscillation (BAO) and Cosmic Microwave Background (CMB) data. By choosing  $m_1 \simeq m_2 \simeq m_3 = 0.07 \text{ eV}$ , we find

$$Y_\Delta = \frac{10^{-2} \text{ eV}}{v_\Delta} \times \begin{pmatrix} 6.80 - 0.06i & -0.13 - 0.04i & 0. - 1.65i \\ -0.13 - 0.04i & 6.91 - 0.03i & 0. - 1.10i \\ 0. - 1.65i & 0. - 1.10i & 6.71 + 0.10i \end{pmatrix} \quad (\text{degenerate case}). \quad (2.11)$$

However, as we will see later, the degenerate scenario does not offer any significant information in our study.

The triplet scalar decay to leptonic final states,  $H^{++} \rightarrow \ell^+ \ell^+$ ,  $H^+ \rightarrow \ell^+ \nu_\ell$ ,  $H_0/A_0 \rightarrow \nu \nu$  with almost 100% branching ratio when  $v_\Delta \leq 0.1 \text{ MeV}$ . On the other hand, for larger  $v_\Delta (> 0.1 \text{ MeV})$ , gauge boson final states become dominant with the decay modes  $H^{++} \rightarrow W^+ W^+$ ,  $H^+ \rightarrow W^+ Z$ ,  $H_0/A_0 \rightarrow ZZ/W^+ W^-$ .<sup>2</sup> For almost same order of  $\lambda_\phi$  and  $\lambda_\Delta$  within the range for producing correct relic density, the DM can annihilate to three different final states (depending upon the branching ratios i.e. on  $v_\Delta$  for triplets):

- (i) neutrino and charged lepton final states.
- (ii)  $W^\pm W^\mp$ ,  $ZZ$  final states mainly from triplet scalar decays;  $WW^*$ ,  $ZZ^*$  final states arise from the decay of the SM Higgs boson.
- (iii)  $b, \tau$  final states from the decay of the SM Higgs boson.

We choose two extreme values of  $v_\Delta$ , namely, 1 eV and 1 GeV, respectively to incorporate the effect of triplet scalars decay to only leptonic final state or only to non-leptonic final states. Our study is mainly focused on the study of neutrino flux coming from DM annihilation which can be obtained from all the above three final states. However, for triplet scalars decaying into leptonic final states includes the contribution of Yukawa couplings  $Y_\Delta$  which is a function of neutrino mass hierarchy and we expect to see the difference in the neutrino flux between the NH and IH scenarios. In the following sections, we will explicitly show how our analysis depends upon the neutrino mass hierarchy.

In this regard, it is to be noted that the triplet scalars can produce interesting signals at the colliders. The strongest limit on the scalar masses comes from the current searches at the LHC for the signature of doubly charged scalar where they can be pair produced

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<sup>2</sup>Throughout this analysis we assume degenerate triplet scalars.

via Drell-Yan and Vector Boson fusion and then decay to a pair of same-sign leptons. The experimental lower bound on  $m_{H^{\pm\pm}}$  has been set by the ATLAS experiment [31] from a pair of isolated lepton search for ( $H^{\pm\pm} \rightarrow e^\pm e^\pm$ ) and ( $H^{\pm\pm} \rightarrow \mu^\pm \mu^\pm$ ) decay modes at the center of mass energy 8 TeV. The 95% C.L. lower limit on the doubly charged Higgs mass for the same-sign isolated muons final state exclude mass range below 516 GeV. However, the experimental lower limit is based on the assumption that the doubly charged scalar decay to the dimoun channel with 100% branching ratio which is not the case in our scenario. In this model, the doubly charged scalar decay to dimuon channel with atmost 30% branching ratio and thus the lower limit on  $m_{H^{\pm\pm}}$  can be relaxed. Following figure 5 of reference [31], we set the degenerate mass of the triplets ( $m_\Delta$ ) as 400 GeV throughout our analysis.

### 3 Analysis

In this section, we demonstrate our findings on the DM annihilation at the Galactic Center and inside the Sun and the possibilities of using the respective neutrino signals to distinguish the NH and IH scenarios. For the purpose of our analysis, we choose to work with four benchmark points, shown in table 1, by fixing  $v_\Delta$  for two DM masses. We choose the DM mass to be greater than 400 GeV to get on-shell triplet scalars in the pair production processes of DM annihilation. Our choice of  $v_\Delta$  has already been justified in the previous section. We choose  $\lambda_\Delta$  and  $\lambda_\Phi$  couplings in the model (see eq. (2.3)) such that the correct DM abundance is obtained via thermal freeze-out with nominal value for thermally averaged annihilation cross section  $\langle \sigma_{\text{ann}} v \rangle \simeq 3 \times 10^{-26} \text{ cm}^3/\text{s}$ .  $\lambda_\Phi$  also enters into the direct detection cross section ( $\sigma_{\text{SI}}$ ) through the  $t$ -channel exchange of the SM Higgs. Values of  $\lambda_\Phi$  that satisfy thermal relic abundance yield a direct detection cross section well below the current limits from LUX experiment [32]. The choice of  $\lambda_\Delta$  and  $\lambda_\Phi$  and the corresponding relic density and direct detection cross section are also listed in table 1. The correct DM relic abundance can also be produced via non-thermal mechanisms [33–35] in which case a larger DM annihilation cross section is allowed. We will discuss the implications of a larger annihilaton cross section at the end of this section.

As mentioned earlier, the DM annihilation will give rise to neutrino fluxes of different flavors that will help in distinguishing the structure of Yukawa coupling and, hence, the neutrino mass hierarchy. The IceCube detects neutrinos by recording the Cherenkov light from relativistic charged particles in its volume. Muon neutrinos ( $\nu_\mu$ ) produce muon ( $\mu$ ) tracks via the charged current interactions in the detector. On the other hand, electron neutrinos ( $\nu_e$ ) and tau neutrinos ( $\nu_\tau$ ) result in hadronic and electromagnetic cascade events in the ice. Since the cascades are localized, they carry no directional information, and hence are not good enough for performing a meaningful DM search over the background. For this reason, we focus on  $\nu_\mu$  fluxes that arrive on the Earth. We also estimate the photon flux for all scenarios in our model. Cosmic ray showers create a muon background that can be controlled by selecting only the upward going events since muons are stopped in the Earth. This limits the observation of DM signal to the time when the source is below the horizon. With atmospheric muons thus eliminated, the most significant contribution to the remaining background comes from atmospheric neutrinos [36]. In addition to this, a portion of the detector may be used as a veto to observe the contained muon events with a conversion vertex inside the instrumented volume, as in the case of DeepCore array in IceCube. The veto procedure virtually eliminates the contribution to the background from atmospheric

Benchmarks	$v_\Delta$	$m_{\text{DM}}(\text{GeV})$	$m_\Delta(\text{GeV})$	$\lambda_\Delta$	$\lambda_\phi$	relic density	$\sigma_{\text{SI}}(\text{pb})$
BP1	1 eV	500	400	0.055	0.04	0.117	$2.25 \times 10^{-10}$
BP2	1 eV	700	400	0.075	0.05	0.122	$1.79 \times 10^{-10}$
BP3	1 GeV	500	400	0.055	0.04	0.117	$2.25 \times 10^{-10}$
BP4	1 GeV	700	400	0.075	0.05	0.122	$1.79 \times 10^{-10}$

**Table 1.** Benchmark Points (BPs) and the corresponding DM relic density and direct detection cross sections.

muons by selecting only contained vertices. This increases the potential observation time to the full year when the source is both above and below the horizon.

We use both contained and through-going muons in our analysis. For contained muon tracks, the vertex at which  $\nu_\mu$  is converted to a muon, is within the detector volume. Through-going muons represent those events that go through a surface inside the detector but may have been produced outside it. To acquire the muon spectra, we use three simple methods:

1. WimpSim [37] automatically provides muons converted from solar neutrinos.
2. DarkSUSY [38] gives muons from atmospheric neutrinos (which we average to make an isotropic approximation).
3. Finally, GENIE [39] provides muons converted from neutrinos coming from the Galactic Center.

GENIE provides the spectrum of contained muons produced from  $\nu_\mu$ s arriving at the detector. To find the total number of muon events, we use the total deep-inelastic charged-current cross sections from reference [40] to find the average  $\nu_\mu$ -to- $\mu$  and  $\bar{\nu}_\mu$ -to- $\bar{\mu}$  cross sections. Modeling the Galactic Center as an isotropic circle with radius of  $5^\circ$  and ignoring detector angular resolution, we obtain a  $5^\circ$  optimal cut, whereas the solar muons have a  $2^\circ$  optimal cut. To find the through-going muons, we propagate the muons using the parameters of table 4 of reference [41].

In order to generate the energy spectra of the SM particles that are produced at the annihilation point, we first create the model file for CalcHEP using FeynRules 2.0 [42] and then use CalcHEP 3.6.23 [43] and Pythia-6 [44] to get the spectra. These spectra are valid both at the Galactic Center and inside the Sun. In the following we present our result for these two distinct DM annihilation points.

### 3.1 Signals from DM annihilation in the Galactic Center

To calculate the fluxes near the Galactic Center, we run the indirect detection module of micrOMEGAsv4.1.8 [45] and measure the fluxes for angle of sight  $0^\circ \leq \Psi \leq 5^\circ$ . We use the Navarro-Frenk-White (NFW) [46] profile  $\rho(r) = \rho_0(r/r_s)^{-1}/(1+r/r_s)^2$  with  $r_s = 20$  kpc,  $\rho_0 = 0.4$  GeV/cm<sup>3</sup>. It is noteworthy that if the annihilation of our proposed DM candidate happens near the Galactic Center, then the final state neutrinos will encounter oscillations on their way to the Earth. In the appendix, we show the flavor dependent probability of neutrino flux that originates at the Galactic Center and reaches the detector surface on the Earth. Hence, the final neutrino fluxes will be given by eq. (A.4)–(A.7b). We have checked that varying all neutrino oscillation parameters within the  $3\sigma$  allowed range about their best

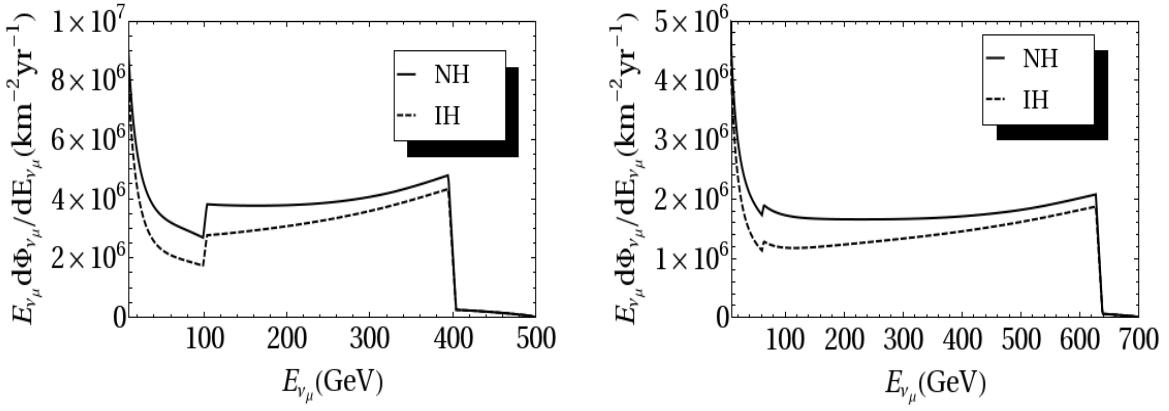
fit central value results in  $\mathcal{O}(3\text{--}4\%)$  change in the neutrino flux from DM annihilation at the Galactic Center. This is much smaller than the difference due to different mass hierarchies.

Following are the interesting features that are seen in figures 1, 2 and 3.

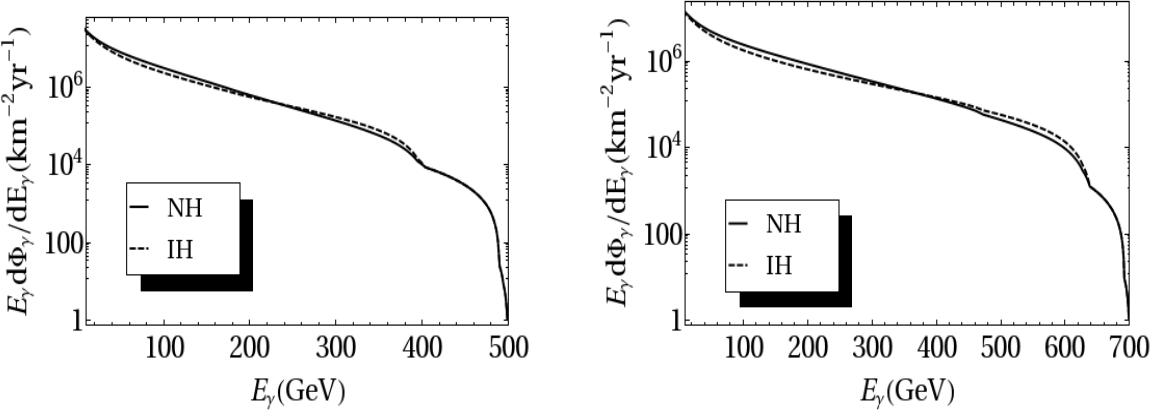
- Figure 1 displays the  $\nu_\mu$  fluxes for NH and IH for BP1 and BP2 which is for  $v_\Delta = 1\text{ eV}$ . As previously argued, at such low vev, the triplet scalars generate leptonic final state with almost 100% branching ratio and the DM annihilation to two on-shell triplet scalars mainly contribute to the annihilation cross section, hence to the  $\nu_\mu$  flux. Now, it is to be observed that for  $m_{\text{DM}} = 500\text{ GeV}$ , the flux rises at around 100 GeV and then falls near 400 GeV and similarly, the rise and fall occur at 60 GeV and 640 GeV for  $m_{\text{DM}} = 700\text{ GeV}$ . This is not surprising and can be understood by the kinematics of the DM annihilation. The triplets are produced on-shell with some boost that comes from the mass difference between the DM and the triplets. Thus the two-body final states render a box like feature with end points  $m_{\text{DM}}(1 \pm \beta)/2$  where  $\beta = \sqrt{1 - m_\Delta^2/m_{\text{DM}}^2}$ .
- We also see from figure 1 that there exists a significant difference between the NH and the IH scenarios. This is due to the fact that in the NH case, we get more taus ( $\tau$ ) from the triplet decays which produces  $\nu_\tau$  and the  $\nu_\tau$  further gets converted into  $\nu_\mu$ . For both the DM mass, the neutrino flux can be easily a factor of two higher in the NH scenario than the IH case. The neutrino flux due to degenerate case will be in between the flux generated for NH and IH cases. This is simply because in the degenerate case, we get less  $\tau$ s from the triplet decays compared to the NH case but more than the IH case. This is true for all our discussions to follow. Therefore, the degenerate case will not provide any new insight and so we restrain ourselves from doing any further remark on this case.
- In figure 2, we present the diffused photon fluxes for NH and IH cases. In our model, the photon flux arises from external charged lepton legs, final state radiation, and secondary decay of charged leptons that are directly produced from triplet decay. For both the benchmark points, we observe that the photon flux does not distinguish between NH and IH cases.<sup>3</sup>
- With the increase in  $v_\Delta$ , the branching ratio of the triplet scalars to leptons reduces and the decay modes with gauge boson final states open up and at sufficiently large  $v_\Delta$ ,  $\text{BR}(H^{\pm\pm} \rightarrow W^\pm W^\pm) \simeq 100\%$ , dominating over the leptonic final states. Hence, for  $v_\Delta = 1\text{ GeV}$ , the dependence on neutrino mass hierarchy gets washed away. Moreover, the contribution to neutrino fluxes coming from DM annihilating to Higgs will become more pronounced in this case. In figure 3, we show neutrino fluxes for BP3 and BP4 and evidently these do not distinguish between NH and IH cases. However, the shapes of neutrino spectra are different compared to the previous  $v_\Delta = 1\text{ eV}$  cases as shown in the figure 1. As expected, the flux is higher in the previous case because the leptonic final states originate from direct triplet decay. The photon flux, however, does not draw any discrimination be it between  $v_\Delta = 1\text{ eV}$  and  $1\text{ GeV}$  cases or between NH and IH scenarios.

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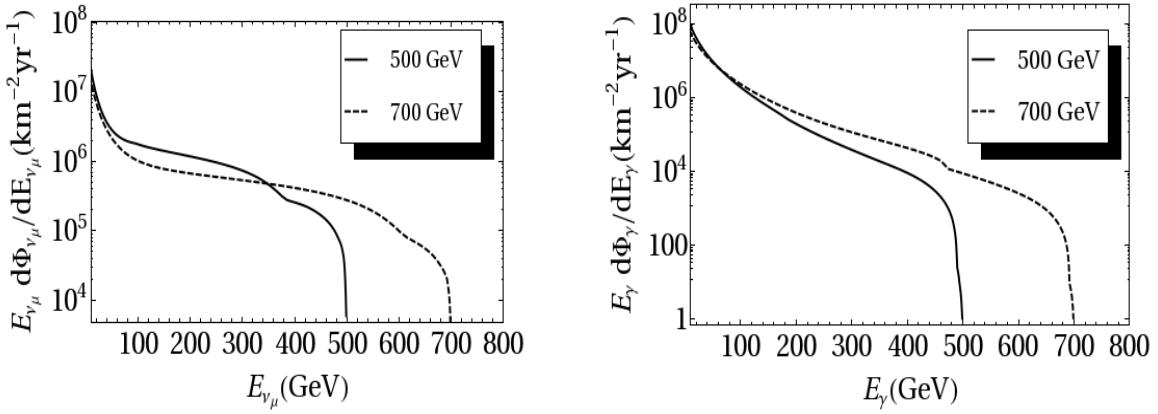
<sup>3</sup>The DM annihilation can also produce photon lines via charged scalar loop processes  $\text{DM} + \text{DM} \rightarrow \gamma\gamma$  and  $\text{DM} + \text{DM} \rightarrow Z\gamma$  with respective energies  $m_{\text{DM}}$  and  $m_{\text{DM}} \left(1 - \frac{m_Z^2}{m_{\text{DM}}^2}\right)$ . However, the line signal is highly suppressed relative to the diffused photon signal in our model, which is in agreement with the fact that that Fermi-LAT has not observed any photon line at such energies.



**Figure 1.** Spectra of  $\nu_\mu$  from DM annihilation at the Galactic Center for BP1 (left) with  $m_{\text{DM}} = 500 \text{ GeV}$  and BP2 (right) with  $m_{\text{DM}} = 700 \text{ GeV}$ .



**Figure 2.** Photon spectra from DM annihilation at the Galactic Center for BP1 (left) with  $m_{\text{DM}} = 500 \text{ GeV}$  and BP2 (right) with  $m_{\text{DM}} = 700 \text{ GeV}$ .



**Figure 3.** Spectra of  $\nu_\mu$  (left) and Photon (right) from DM annihilation at the Galactic Center for BP3 and BP4.

### 3.2 Signals from DM annihilations inside the Sun

DM particles that pass through the Sun lose energy due to the DM-nucleon scattering and get gravitationally trapped [47–53]. The total rate of DM annihilations  $\Gamma_{\text{ann}}$  inside the Sun follows [1]

$$\Gamma_{\text{ann}} = \frac{C}{2} \tanh^2(\sqrt{CA}t). \quad (3.1)$$

Here  $C$  is the capture rate of DM particles by the Sun, which is related to  $\sigma_{\text{SI}}$  in our model, and  $A$  is related to the DM annihilation cross section  $\sigma_{\text{ann}}$  (for details, see [1]). DM capture and annihilation reach equilibrium when  $t > \tau_{\text{eq}} \equiv (\sqrt{CA})^{-1}$ , in which case  $\Gamma_{\text{ann}} \approx C/2$ . For values of  $\sigma_{\text{SI}}$  given in table 1, and  $\langle \sigma_{\text{ann}} v \rangle = 3 \times 10^{-26} \text{ cm}^3 \text{ s}^{-1}$ , the equilibrium condition is achieved for the age of the Sun. We note that as long as equilibrium is established, the total rate of DM annihilations is controlled by the direct detection cross section  $\sigma_{\text{SI}}$ . Therefore, increasing  $\sigma_{\text{ann}}$  (which is allowed in scenarios of non-thermal DM) and keeping  $\sigma_{\text{SI}}$  fixed does not enhance the signal from DM annihilation inside the Sun. It only leads to a faster establishment of equilibrium between DM capture and annihilation.

DM annihilation final states that are most relevant for producing neutrinos inside the Sun are prompt neutrinos from  $H^0/A^0$  and  $H^+/H^-$  decays,  $W, Z$  and  $\tau$  particles from  $H^+/H^-$  and  $H^{++}/H^{--}$  decays. Other final states (like  $e$ ,  $\mu$ , and lighter quarks) lose their energy and stop almost immediately, which results in production of neutrinos at energies below detection threshold. The charged current interactions inside the Sun convert  $\nu_e$ ,  $\nu_\mu$ ,  $\nu_\tau$  to  $e$ ,  $\mu$ ,  $\tau$  respectively. Using the charged current neutrino-nucleon cross section [40], we find that neutrino absorption becomes important (i.e.,  $L_{\text{abs}} < R_{\text{C}}$ , with  $R_{\text{C}}$  being the core radius of the Sun) at energies  $E_\nu > 300 \text{ GeV}$ . For  $\nu_e$ , the flavor and mass eigenstates are the same deep inside the Sun, and hence absorption suppresses the flux of  $\nu_e$  at energies above 300 GeV. On the other hand,  $\nu_\tau$  absorption produces  $\tau$  particles that decay quickly before losing too much energy because of very short lifetime of  $\tau$ . This decay produces  $\nu_\tau$  at a lower energy, and hence this “regeneration” effect populates the spectrum at energies well below the DM mass. Since  $\nu_\mu$  and  $\nu_\tau$  undergo oscillations inside the Sun that is dominantly set by the atmospheric mass splitting  $\Delta m_{\text{atm}}^2$ , we have  $L_{\text{osc}} \propto E_\nu / \Delta m_{\text{atm}}^2$ . As long as  $L_{\text{abs}} \gtrsim L_{\text{osc}}/4$ , oscillations mix  $\nu_\mu$  and  $\nu_\tau$  efficiently and  $\nu_\mu$  final states also feel the regeneration effect. When  $L_{\text{abs}}$  drops below  $L_{\text{osc}}/4$ , which happens at  $E_\nu \sim 500 \text{ GeV}$ , oscillations cease to be effective and  $\nu_\mu$  gets absorbed similar to  $\nu_e$ . As a consequence, only the  $\nu_\tau$  final state retains a significant regeneration signature at energies above 500 GeV.<sup>4</sup>

For neutrinos of sufficiently low energies, vacuum flavor oscillations between the Sun and the Earth are averaged over half a year when the Sun is below the horizon at the south pole. In particular, below 100 GeV, the oscillation length set by the solar mass splitting  $\Delta m_{\text{sol}}^2$  is less than approximately 3 million kilometer change in the Earth-Sun distance over half a year. The situation changes at energies above 100 GeV where solar neutrino oscillations are not averaged out anymore. For monochromatic neutrinos of a single flavor, this significantly affects oscillations between  $\nu_e$  and  $\nu_\mu/\nu_\tau$ , and hence the  $\nu_\mu$  spectrum at the detector, at high energies. However, this is not an important effect in our model since DM annihilation produces different neutrino flavors with a continuous energy spectrum in this case.

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<sup>4</sup>Neutrinos also have neutral current interactions with matter inside the Sun that results in energy loss of the neutrinos from all flavors further shifting their spectra toward lower energies. However, the cross section for neutral current interactions is a factor of 3 smaller than that for charged current interactions, which makes them subdominant.

We use DarkSUSY 5.1.1 [38] to simulate production of neutrinos in the Sun, their propagation to the South Pole, and the interaction of  $\nu_\mu$  with ice at many different energies for all the SM channels. DarkSUSY does this by interpolating an older WimpSim [37] simulation. Here we run WimpSim 3.05 directly for the prompt neutrino channels, sample energies spaced at 10 GeV, and use oscillation parameters found in [54].

We then combine the energy spectra of all the channels at the production point with how they propagate, giving us the final neutrino spectra at the detector.<sup>5</sup> It should be noted that variation of all neutrino oscillation parameters within the  $3\sigma$  allowed range results in  $\mathcal{O}(3\text{--}4\%)$  change in the final neutrino spectra from the Sun, similar to the Galactic Center case, which is insignificant.

In figure 4 we show the spectra of  $\nu_\mu$  at the detector for  $v_\Delta = 1\text{ eV}$  (BP1, BP2) and for  $v_\Delta = 1\text{ GeV}$  (BP3) cases. It is seen that at energies below 400 GeV the NH scenario results in a larger number of  $\nu_\mu$  and eventually muons than the IH scenario. The reason is similar to the Galactic Center case and can be explained by the triplet scalar decay patterns. For  $v_\Delta = 1\text{ eV}$ , as mentioned above, triplet decays mainly produce  $\nu_\tau$ ,  $\nu_\mu$  and  $\tau$ ,  $\mu$  final states in the NH scenario. In the IH scenario, on the other hand, triplet decays mainly produce  $\nu_e$  and  $e$  final states. As far as charged lepton final states are concerned, only  $\tau$  decays before losing a significant fraction of energy and produces neutrinos. Regarding triplet decays to neutrinos,  $\nu_\tau$  (and to a lesser extent  $\nu_\mu$ ) is the most relevant flavor that survive at energies above 300 GeV due to regeneration effect. In consequence, for  $v_\Delta = 1\text{ eV}$ , a larger number of  $\nu_\mu$  arrive at the detector in the NH scenario. However, the neutrino spectra in the two scenarios do not differ when  $v_\Delta = 1\text{ GeV}$  as in earlier case. Comparing the neutrino signal from DM annihilation in the Sun (figure 4) and at the Galactic Center (figure 1), we notice some differences. First, the kinematic cuts in figure 1 are not prominent in figure 4. This is due to the fact that absorption/regeneration and down scattering inside the Sun erases sharp features in the spectra. More precisely, only a little bump develops at low energies for  $m_{\text{DM}} = 500\text{ GeV}$  due to the boosted triplet scalars, shown in right panel of figure 4, unlike the box-like structure that can clearly be observed in the Galactic Center neutrino flux. At higher energies, instead of a bump only the wiggles can be noticed which are the effect of neutrino oscillation. Furthermore, the difference between the spectra in NH and IH scenarios in figure 4 is smaller than that in figure 1. This is because stopping of muons and partial absorption of  $\nu_\mu$  inside the Sun decreases the overall flux of  $\nu_\mu$  arriving at the detector. Nevertheless, the NH scenario yields a larger flux than the IH scenario when  $v_\Delta = 1\text{ eV}$ .

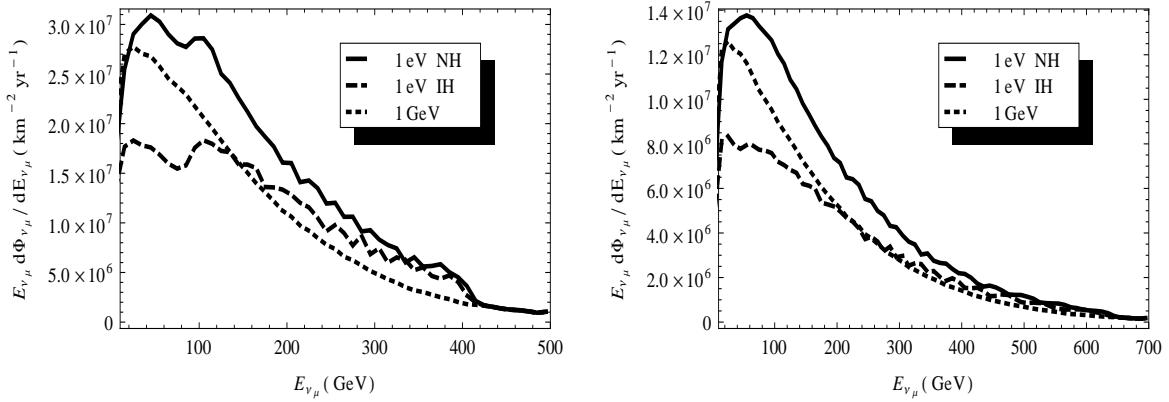
We emphasize that the neutrino signal from DM annihilation inside the Sun is complementary to that from the Galactic DM annihilation as they are set by different cross sections ( $\sigma_{\text{SI}}$  and  $\sigma_{\text{ann}}$  respectively). We see from figures 1 and 4 that for the model parameters given in table 1 the signal from the Sun is about an order of magnitude stronger than that from the Galactic Center. As we will discuss later, both of these signals can be enhanced further.

### 3.3 Muon spectra at the detector

We now present the muon spectra obtained from the conversion of  $\nu_\mu$  at the detector, which is the observed signal at neutrino telescopes. For simulation purpose, we assume a detector that has the same capability as the IceCube DeepCore array. In figure 5, we show the spectra of contained and through-going muons in the detector from DM annihilation at the Galactic

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<sup>5</sup>Since  $\bar{\nu}$ -nucleon cross section is about two times smaller than the  $\nu$ -nucleon cross section [55, 56], the number of absorbed  $\bar{\nu}$  inside the Sun is smaller. This results in a larger  $\bar{\nu}$  flux at the detector from DM annihilation inside the Sun.



**Figure 4.** Spectra of  $\nu_\mu$  from DM annihilation inside the Sun at the detector for  $m_{\text{DM}} = 500$  GeV (700 GeV) in the left (right) panel.

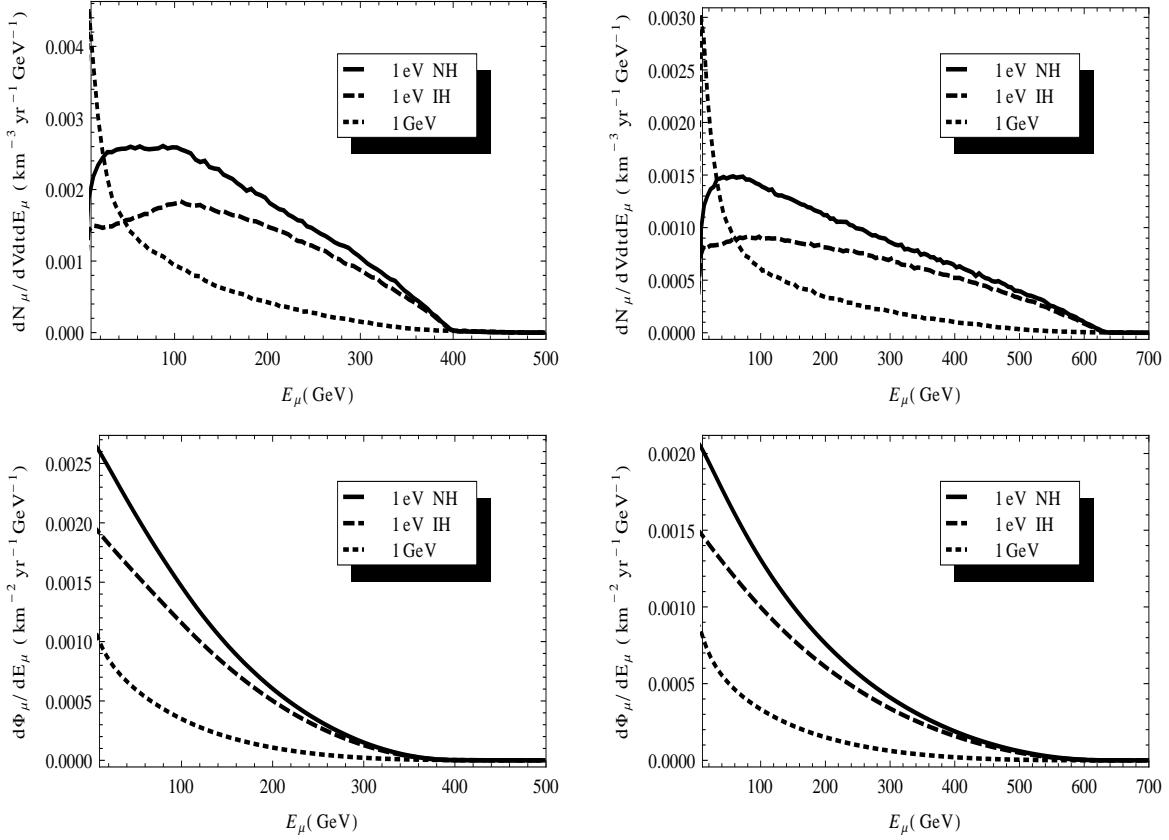
Center. In this case, the signal comes from a region around the Galactic Center that has  $5^\circ$  angular extension. Hence, in order to optimize the signal to background ratio, we have imposed a  $5^\circ$  angular cut on the muons relative to the center of the galaxy.

In figure 6, we show the spectra of contained and through-going muons from DM annihilation inside the Sun. Since the Sun is a point-like source, the optimal signal to background ratio is obtained for an angular cut of  $2^\circ$ , in this case. The peaks in both figures are due to the imposed angular cuts that eliminate muons produced from  $\nu_\mu$  conversion below a certain energy. For through-going muons the peak occurs at a lower energy, which can be understood by noting that the measured energy of a through-going muon is in general less than the actual energy at the production point.

As expected, the muon spectra have less features than the neutrino spectra shown in figures 1–4. Nevertheless, the difference between  $v_\Delta = 1$  eV and  $v_\Delta = 1$  GeV cases, as well as the NH and IH scenarios in the former case, are clearly visible in the muon spectra arising from both the Galactic Center and the Sun. Similar to the neutrino spectra, substantial deviation among the cases is seen in the signal arising from DM annihilation at the Galactic Center than in the Sun. However, we also note that the absolute value of the flux from the Galactic Center is smaller by an order of magnitude than that from the Sun.

The muon spectra shown in figures 5 and 6 are obtained assuming the nominal thermal freeze-out value of  $\langle \sigma_{\text{ann}} v \rangle$ . However, the Fermi-LAT data from Milky Way dwarf spheroidal galaxies [57] allows larger values on the DM annihilation cross section for the values of DM mass that we have considered here. One can see from figure 8 of [57] that for the  $b\bar{b}$  and  $W^\pm W^\mp$  final states  $\langle \sigma_{\text{ann}} v \rangle$  can be larger than the thermal freeze-out value by a factor of 4 (6) when  $m_{\text{DM}} = 500$  (700) GeV. In BP1 and BP2, where the distinction between the NH and IH scenarios is significant, these final states arise from DM annihilation to the SM Higgs that is controlled by the coupling  $\lambda_\Phi$  of eq. (2.3). As explained before,  $\lambda_\Phi$  also controls  $\sigma_{\text{SI}}$  in our model. Hence raising  $\langle \sigma_{\text{ann}} v \rangle$  by the above factors, due to a larger value of  $\lambda_\Phi$ , also results in an increase in  $\sigma_{\text{SI}}$  compared with that given in table 1, which is still consistent with the LUX bounds [32]. For the  $\tau^+\tau^-$  final state,  $\langle \sigma_{\text{ann}} v \rangle$  can be larger than the thermal freeze-out value by a factor of 10 (17) when  $m_{\text{DM}} = 500$  (700) GeV.<sup>6</sup> This final state arises due

<sup>6</sup>We have four-body annihilation final states in this case. Moreover, we have a combination of different final states instead of 100% of just one final state. Nevertheless, the Fermi-LAT limit provides a reasonable upper bound on  $\sigma_{\text{ann}}$  in this case too.

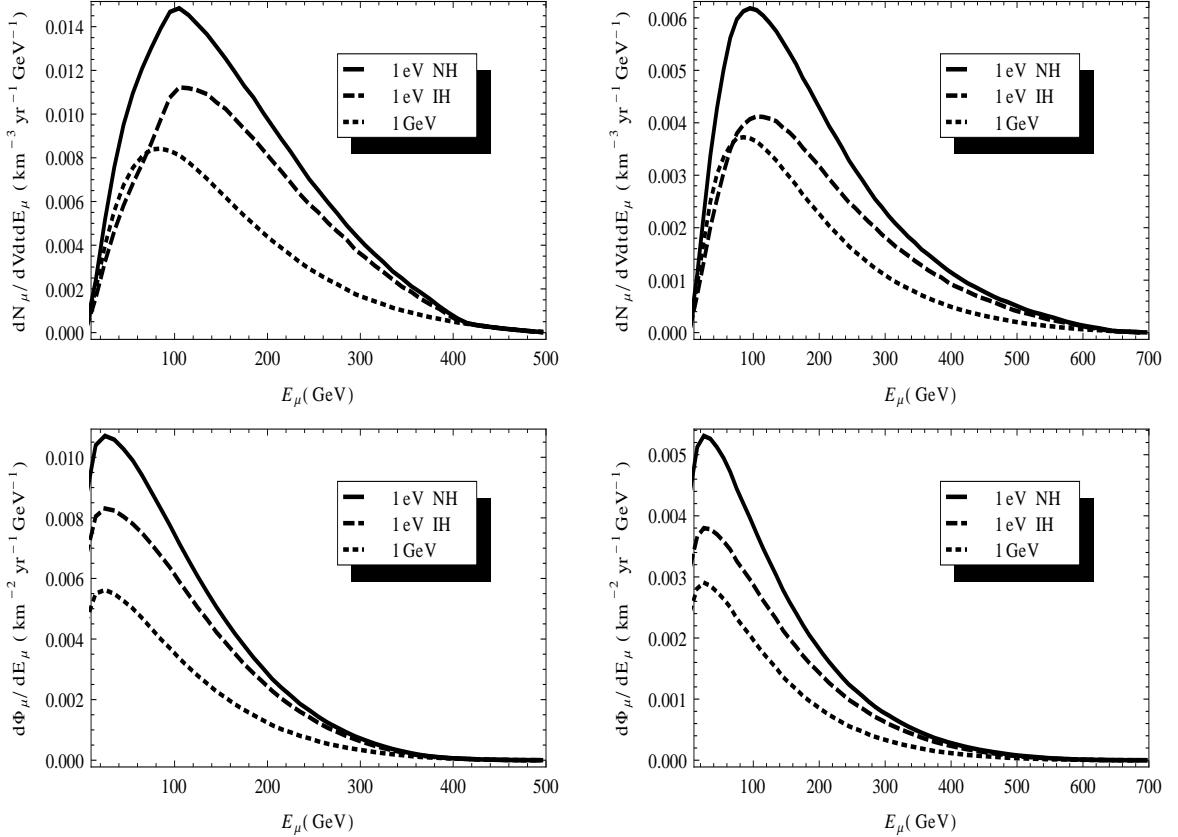


**Figure 5.** Spectra of muons from DM annihilation at the Galactic Center for  $m_{\text{DM}} = 500$  GeV (700 GeV) in the left (right) panel. Upper and lower panels show contained and through-going muon spectra with an angular cut of  $5^\circ$  respectively.

to DM annihilation to the triplet Higgs, which is controlled by the coupling  $\lambda_\Delta$  of eq. (2.3). Increasing  $\lambda_\Delta$  results in an increase in  $\sigma_{\text{ann}}$ , but does not affect  $\sigma_{\text{SI}}$ .

Therefore, by invoking non-thermal mechanisms for DM production in the early universe, we can obtain considerably larger neutrino fluxes from DM annihilation at the Galactic Center and inside the Sun. The resulting enhancement in the muon events, compared to figures 5 and 6, leads to a better prospect for detection of the neutrino signal against the atmospheric background. In figure 7, we show the spectra of contained muons due to DM annihilation at the Galactic Center (left panel) and inside the Sun (right panel) for annihilation cross sections that are just below the Fermi-LAT limits, where  $m_{\text{DM}} = 500$  GeV and the background from atmospheric neutrinos is also shown for comparison. Contained muons from DM annihilation inside the Sun provide the best detection opportunity. Within energy interval 100–400 GeV, the total number of contained muons for the NH and IH scenarios is 11 and 9 respectively (compared with 95 for the background). The maximum signal to background ratio, 17% and 15% for the NH and IH scenarios respectively, occurs at an energy of 255 GeV. For a neutrino telescope with the same capability as the IceCube DeepCore array,  $3\sigma$  discovery of NH and IH scenarios takes 8 and 12 years respectively. To distinguish the different neutrino mass hierarchies, one should go beyond the simple number count and perform a careful shape analysis, which is beyond the scope of this paper.

There are proposals for directly measuring the mass hierarchy by using atmospheric neutrinos in future extensions of neutrino telescopes such as PINGU (Precision IceCube Next



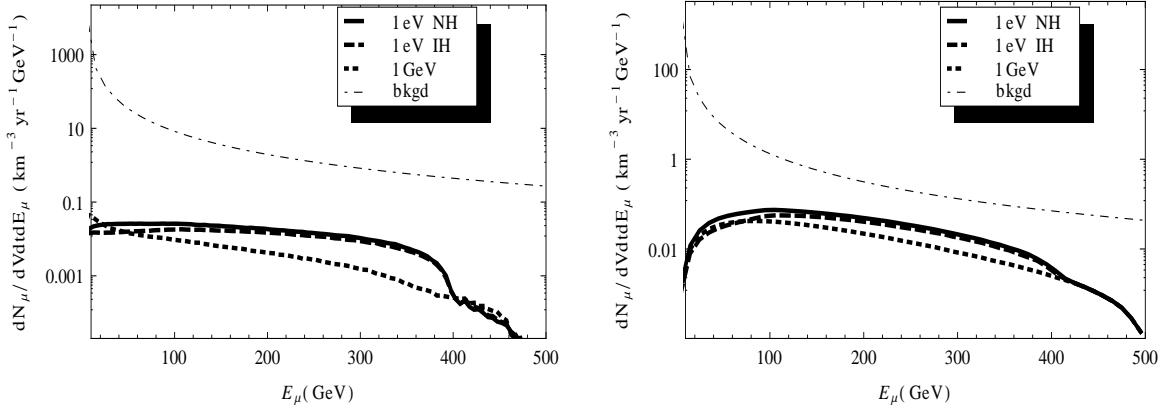
**Figure 6.** Spectra of muons from DM annihilation in the Sun for  $m_{\text{DM}} = 500 \text{ GeV}$  (700 GeV) in the left (right) panel. Upper and lower panels show contained and through-going muon spectra with an angular cut of  $2^\circ$  respectively.

Generation Upgrade) [58]. It is possible to exclude the wrong mass ordering by this method (as well as in neutrino beam experiments) at the  $3\sigma$  level within the next 10–15 years [59]. Our approach, which exploits an interesting connection between DM and neutrinos, is complementary to this direct method and makes another case for the future neutrino telescopes with much larger volume like IceCube-Gen2 [60].

#### 4 Summary and conclusions

In this paper, we have considered an extension of the SM that explains neutrino masses and mixing angles via the type-II seesaw mechanism and includes a singlet scalar as a viable DM candidate. The DM interacts with the triplet scalar that generates light neutrino masses thereby linking the two sectors. We have studied the neutrino signal from DM annihilation at the Galactic Center and inside the Sun in this model that could be detected at IceCube. Our main results are summarized as follows:

- In the region of the parameter space where the triplet scalar dominantly decays to leptonic final states, the flux of  $\nu_\mu$  from DM annihilation depends upon the neutrino mass hierarchy (normal, inverted or degenerate). The photon flux, on the other hand, is practically the same for different hierarchies.



**Figure 7.** Spectra of contained muons from DM annihilation at the Galactic Center (left) and inside the Sun (right) for  $m_{\text{DM}} = 500$  GeV. The value of  $\sigma_{\text{ann}}$  used to obtain the muon flux is just below the current bounds from Fermi-LAT. The top line in both figures is due to the background arising from the atmospheric neutrinos.

- The difference in the flux of  $\nu_\mu$  at the detector for NH and IH is visible. The NH scenario produces a larger flux because of its mass pattern. This holds for both signals from DM annihilation at the Galactic Center and inside the Sun.
- The difference is more significant in the neutrino signal arising from DM annihilation at the Galactic Center than that arises from DM annihilation inside the Sun. This is mainly because the interaction of neutrinos with matter inside the Sun results in a moderate attenuation of the final  $\nu_\mu$  flux at the detector.
- The same reason holds for the contained and through-going muons that are produced from the conversion of  $\nu_\mu$ s at the detector. The muon spectra have less features than the neutrino spectra. Nevertheless, these muons can be detected over the atmospheric background.
- For annihilation cross sections just below the Fermi-LAT limits, future extensions of IceCube are able to discover the neutrino signal and distinguish the different neutrino mass hierarchies with multiyear data.

The LHC in tandem with direct detection experiments, the IceCube and Fermi-LAT will be able to probe this minimal extension of the SM that utilizes a new Higgs sector. Finally, we would like to conclude by pointing out the fact that this analysis can be applied to other models that link DM to the neutrino sector.

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## A Oscillation probability calculation

The flavor eigenstate of neutrinos can be related to the mass eigenstate by the  $3 \times 3$  Pontecorvo-Maki-Nakagawa-Sakata (PMNS) mixing matrix.

$$\begin{pmatrix} \nu_e \\ \nu_\mu \\ \nu_\tau \end{pmatrix} = \begin{pmatrix} U_{e1} & U_{e2} & U_{e3} \\ U_{\mu 1} & U_{\mu 2} & U_{\mu 3} \\ U_{\tau 1} & U_{\tau 2} & U_{\tau 3} \end{pmatrix} \begin{pmatrix} \nu_1 \\ \nu_2 \\ \nu_3 \end{pmatrix}. \quad (\text{A.1})$$

The neutrinos from the Galactic Center encounter flavor oscillation while reaching the Earth. However, since it traverses very large distance with respect to its oscillation length, hence the oscillation term is averaged out. The probability of one neutrino flavor to oscillate into other becomes simply,

$$\langle P(\nu_\alpha \rightarrow \nu_\beta) \rangle = \sum_{i=1}^3 |U_{\alpha i}|^2 |U_{\beta i}|^2 \quad (\text{A.2a})$$

$$= XX^T \quad (\text{A.2b})$$

with

$$X = \begin{pmatrix} |U_{e1}|^2 & |U_{e2}|^2 & |U_{e3}|^2 \\ |U_{\mu 1}|^2 & |U_{\mu 2}|^2 & |U_{\mu 3}|^2 \\ |U_{\tau 1}|^2 & |U_{\tau 2}|^2 & |U_{\tau 3}|^2 \end{pmatrix}. \quad (\text{A.2c})$$

The PMNS matrix is usually parametrized in terms of the three mixing angles  $\theta_{12}; \theta_{23}; \theta_{13}$ , and one Dirac ( $\delta$ ) and two Majorana ( $\alpha_1; \alpha_2$ ) CP phases:

$$U = \begin{pmatrix} c_{12}c_{13} & s_{12}c_{13} & s_{13}e^{-i\delta} \\ -s_{12}c_{23} - c_{12}s_{23}s_{13}e^{i\delta} & c_{12}c_{23} - s_{12}s_{23}s_{13}e^{i\delta} & s_{23}c_{13} \\ s_{12}s_{23} - c_{12}c_{23}s_{13}e^{i\delta} & -c_{12}s_{23} - s_{12}c_{23}s_{13}e^{i\delta} & c_{23}c_{13} \end{pmatrix} \times \text{diag}(e^{i\alpha_1/2}, e^{i\alpha_2/2}, 1). \quad (\text{A.3})$$

The oscillated flux of the neutrinos at the detector at earth,

$$\begin{pmatrix} \Phi_{\nu_e} \\ \Phi_{\nu_\mu} \\ \Phi_{\nu_\tau} \end{pmatrix} = \begin{pmatrix} P_{11} & P_{12} & P_{13} \\ P_{21} & P_{22} & P_{23} \\ P_{31} & P_{32} & P_{33} \end{pmatrix} \begin{pmatrix} \Phi_{\nu_e}^0 \\ \Phi_{\nu_\mu}^0 \\ \Phi_{\nu_\tau}^0 \end{pmatrix}. \quad (\text{A.4})$$

From eq. (A.2b) and eq. (A.2c),

$$P_{11} = |U_{e1}|^4 + |U_{e2}|^4 + |U_{e3}|^4 \quad (\text{A.5a})$$

$$P_{12} = |U_{e1}|^2 |U_{\mu 1}|^2 + |U_{e2}|^2 |U_{\mu 2}|^2 + |U_{e3}|^2 |U_{\mu 3}|^2 = P_{21} \quad (\text{A.5b})$$

$$P_{13} = |U_{e1}|^2 |U_{\tau 1}|^2 + |U_{e2}|^2 |U_{\tau 2}|^2 + |U_{e3}|^2 |U_{\tau 3}|^2 = P_{31} \quad (\text{A.5c})$$

$$P_{22} = |U_{\mu 1}|^4 + |U_{\mu 2}|^4 + |U_{\mu 3}|^4 \quad (\text{A.5d})$$

$$P_{23} = |U_{\mu 1}|^2 |U_{\tau 1}|^2 + |U_{\mu 2}|^2 |U_{\tau 2}|^2 + |U_{\mu 3}|^2 |U_{\tau 3}|^2 = P_{32} \quad (\text{A.5e})$$

$$P_{33} = |U_{\tau 1}|^4 + |U_{\tau 2}|^4 + |U_{\tau 3}|^4. \quad (\text{A.5f})$$

We use the latest global analysis of the 3-neutrino oscillation data [54],

$$\begin{aligned} \Delta m_{21}^2 &= 7.60 \times 10^{-5} \text{ eV}^2, & \theta_{12} &= 34.6^\circ, \\ \Delta m_{31}^2 &= 2.48 \times 10^{-3} \text{ eV}^2 \text{ (NH)}, & \Delta m_{31}^2 &= 2.38 \times 10^{-3} \text{ eV}^2 \text{ (IH)}, \\ \delta &= 1.41\pi \text{ (NH)}, & \delta &= 1.48\pi \text{ (IH)}, \\ \theta_{23} &= 48.9^\circ \text{ (NH)}, & \theta_{23} &= 49.2^\circ \text{ (IH)}, \\ \theta_{13} &= 8.6^\circ \text{ (NH)} & \theta_{13} &= 8.7^\circ \text{ (IH)}. \end{aligned} \quad (\text{A.6})$$

$$\text{Normal Hierarchy: } P = \begin{pmatrix} 0.5380 & 0.1979 & 0.2640 \\ 0.1979 & 0.4240 & 0.3779 \\ 0.2640 & 0.3779 & 0.3579 \end{pmatrix} \quad (\text{A.7a})$$

$$\text{Inverted Hierarchy: } P = \begin{pmatrix} 0.5377 & 0.2009 & 0.2613 \\ 0.2009 & 0.4223 & 0.3766 \\ 0.2613 & 0.3766 & 0.3620 \end{pmatrix} \quad (\text{A.7b})$$

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