Observable features in (ultra)high energy neutrinos due to active-sterile secret interactions

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(Dated: July 16, 2020)

We consider the effects of active-sterile secret neutrino interactions, mediated by a new pseudoscalar particle, on high- and ultra high-energy neutrino fluxes. In particular, we focus on the case of 3 active and 1 sterile neutrino coupled by a flavor dependent interaction, extending the case of 1 active and 1 sterile neutrino we have recently examined. We find that, depending on the kind of interaction of sterile neutrino with the active sector, new regions of the parameter space for secret interactions are now allowed leading to interesting phenomenological implications on two benchmark fluxes we consider, namely an astrophysical power law flux, in the range below 100 PeV, and a cosmogenic flux, in the Ultrahigh energy range. First of all, the final active fluxes can present a measurable depletion observable in future experiments. Especially, in the case of only $\nu_\tau - \nu_e$ interaction, we find that the effects on the astrophysical power law flux can be so large to be already probed by the IceCube experiment. Moreover, we find intriguing features in the energy dependence of the flavor ratio.

Keywords: Ultra High Energy neutrinos, cosmogenic neutrinos, secret interactions

I. INTRODUCTION

High energy neutrinos can be produced by the interactions of High and Ultrahigh Energy Cosmic rays. Photodisintegration and hadron-hadron interactions of High Energy cosmic rays in astrophysical objects, such as Active Galactic Nuclei (AGN), Gamma-Ray Bursts (GRB), Starburst Galaxies (SBG), can produce high energy neutrinos with energy up to the order of 100 PeV. On the other hand, photodisintegration interaction of Ultrahigh Energy Cosmic rays (UHECRs) with the photons of the Cosmic Microwave Background (CMB) can give rise to cosmogenic neutrinos with energies in the range \(10^2 \div 10^{10}\) PeV.

Recently, in [2] we have studied the possibility that the cosmogenic neutrino flux suffers from a measurable depletion (called absorption effect) observable in future experiments, due to the presence of active-sterile secret neutrino interactions. In particular, in the scheme of 1 active neutrino and 1 sterile, we have shown that the absorption effect is maximal for energies around \(10^8 \div 10^9\) GeV, and it could be observed at experiments like GRAND [3] and ARIANNA [4].

In this paper we investigate the more general case 3 & 1 (3 active and 1 sterile neutrino $\nu_\tau$), where the interaction is possibly flavor dependent and mediated by a pseudoscalar particle $\varphi$. In comparison with our previous work, we examine in detail the constraints on the parameter space of the secret interaction, coming both from laboratory experiments and from cosmological observations. We find that new regions for the parameters are allowed compared to the safe one we had investigated in the previous paper. Extending the analysis to these new regions, we find again an absorption feature in the active neutrino fluxes, which can however occur over a much wider energy scale, reaching energies as low as \(10^6 \div 10^7\) GeV. Therefore, we consider two benchmark fluxes: an astrophysical power law flux, meant as a representative of the cumulative neutrinos from astrophysical objects, in the range below 100 PeV, and a cosmogenic flux, in the Ultrahigh energy range. We find that the effects on the astrophysical power law flux can be so important to be probed by the IceCube experiment [5]. More specifically, the effect of the interaction can produce a cutoff-like feature in the spectrum, which could potentially be an explanation for the lack of observed events above 10 PeV. Finally, we use the full 3 & 1 framework to generalize our predictions to the flavor structure modifications induced by the new interaction, finding interesting features in the energy dependence of the flavor ratio. This is especially important from a phenomenological point of view since flavor identification is performed differently in various energy ranges. In the IceCube energy range, flavor distinction can be partially achieved using the different topologies of cascade and track events. Further, tau neutrinos can give rise to a unique topology, the so called double bang events. Finally, electron antineutrinos should produce as a signature a peak in their interaction with the detector, due to the Glashow resonance in the process mediated by the $W^-$ boson. In the ultrahigh energy range flavor identification is more complicated: it can be partially achieved by distinguishing between the air showers produced by neutrinos of all flavors in the atmosphere, and the showers produced by skimming tau leptons. The latter originate from tau neutrinos passing through the Earth. This distinction therefore would in

\[ \text{energy Cosmic rays in astrophysical objects, such as Active Galactic Nuclei (AGN), Gamma-Ray Bursts (GRB), Starburst Galaxies (SBG), can produce high energy neutrinos with energy up to the order of } 100 \text{ PeV. On the other hand, photodisintegration interaction of Ultrahigh Energy Cosmic rays (UHECRs) with the photons of the Cosmic Microwave Background (CMB) can give rise to cosmogenic neutrinos with energies in the range } 10^2 \div 10^{10} \text{ PeV.} \]
principle allow an identification of the tau neutrino ratio in the ultrahigh energy range. We therefore reach the conclusion that a combined analysis of the energy and flavor structure of the astrophysical neutrino fluxes in the energy region above the PeV would allow to provide definite informations about the possibility of active-neutrino sterile interactions.

The outline of the paper is the following: in Section II we specify our model, emphasizing the role of the parameters. We then discuss in Section III the constraints, coming from laboratory experiments, cosmological observations and astrophysical data. A description of the benchmark fluxes we have analyzed, as well as of the methods used to describe the effects of the interaction, is provided in Section IV. In Section V we show our results and discuss it. Finally, in Section VI we come to our conclusions.

II. MODEL

In this section we describe the model of active-sterile neutrino interaction analyzed for this work. For definiteness we assume throughout the paper that neutrinos are described by Majorana spinors. We also consider just one sterile neutrino $\nu_s$, coupling with the active ones via a new interaction given by

$$\mathcal{L}_{SI} = \sum \lambda_\alpha \overline{\nu}_\alpha \gamma_5 \nu_s \varphi,$$  \hspace{1cm} (1)

where $\alpha = e, \mu, \tau$ and $\lambda_\alpha$ are dimensionless free couplings. The presence of $\gamma_5$ is connected with the Majorana nature of the neutrinos, since the scalar contraction $\overline{\nu}_\alpha \gamma_5 \nu_s$ would otherwise vanish. In order to maintain parity, the mediator of the interaction $\varphi$ is chosen as a pseudoscalar.

The interaction in eq. (1) is assumed to arise after the breaking of $SU_L(2)$ weak group, since it explicitly violates it. The study of a complete Standard Model Lagrangian is beyond the scope of the present paper, since we are only interested into the phenomenological consequences of the interaction (1). Nevertheless, it is interesting to observe that our interaction must be embedded into a more fundamental theory that will give rise to a $4 \times 4$ neutrino mass matrix in the bases $\nu_e, \nu_\mu, \nu_\tau, \nu_s$ diagonalized by a $4 \times 4$ unitarity matrix parametrized by three mixing angle between active-active states $\theta_{12}, \theta_{13}, \theta_{23}$ and three mixing angles between active-sterile states $\theta_{1s}, \theta_{2s}, \theta_{3s}$. For simplicity we assume here that $\theta_{is} \ll 1$, so small as to neglect its effects. Even though this is a simplification which restricts the space of parameters we explore, it allows us to disentangle the effects due to the interaction from the effects due to the active-sterile mixing.

The couplings $\lambda_\alpha$ are free parameters of the model, which means that we have an ample freedom of choice for our model. The most natural possibility is $\lambda_\alpha = \lambda_\mu = \lambda_\tau$, since it preserves lepton universality. However, also the case in which only $\lambda_\tau \not= 0$ is very interesting: even though it is not motivated by symmetry properties, we will see that it is only very weakly constrained by mesons decay. It can therefore lead to larger effects on the astrophysical fluxes without being excluded by present experiments. In our investigation we therefore consider these two benchmark cases.

As mentioned in [2], the cross section for the collision of sterile-active neutrinos exhibits a resonance in the $t$-channel. In fact, if a sterile neutrino with momentum $p$ collides with a fixed active neutrino, the resonance condition $t = M_\varphi^2$ gives the following expressions for the energy of the final sterile and of active neutrino:

$$E_s = \frac{m_\varphi^2 + m_i^2 - M_\varphi^2}{2m_i}, E_a = \sqrt{p^2 + m_\varphi^2 + m_i^2 - m_s^2 + M_\varphi^2},$$  \hspace{1cm} (2)

where $m_i$ is the mass of the $i$-th active neutrino, $M_\varphi$ is the mass of the pseudoscalar $\varphi$ and $m_s$ is the mass of the sterile neutrino. Since $m_i \ll m_s, M_\varphi$, from eq. (2) it follows that the resonance condition can be satisfied for positive energies if $m_s > M_\varphi$. If this condition is satisfied, the decay channel $\varphi \rightarrow \nu_i \nu$ is also kinematically suppressed. The amplitude for the process therefore depends critically on details of the model we have left unspecified. In fact, if the scalar mediator were completely stable, with no other decay channels, the $t$-resonance comes unregulated, giving rise to a non integrable pole in the differential cross section and a diverging total cross section. This situation is analogous to the divergence of the total cross section for Rutherford scattering. The regularization of this divergence in the case of a stable mediator particle depends, just as in the case of Rutherford scattering, on the transverse structure of the beam: if the radius of the beam is $a$, the cross section cannot exceed $\pi a^2$, so that the total cross section will saturate to this value. It is however uncommon for a particle to be completely stable, if this stability does not descend from some specific property or conservation law. Therefore, it is unlikely that our mediator should be completely stable and it may have other decay channels, giving rise to a finite total decay rate $\Gamma$. This decay rate regularizes the divergence. Obviously, this implies a dependence on a new parameter $\Gamma$ into our work for the region $m_s > M_\varphi$. In what follows, we adopt the choice that the dominant decay channel is the decay into two active neutrinos $\nu_i \nu$, if possible, the very small mixing angle. We will discuss in Section IV B the dependence of our results on this assumption.

III. CONSTRAINTS

In the simple extension of the Standard Model under consideration, we introduce two new species of matter:

\footnote{We do not include in this counting the CP violating phases.}
the scalar field $\varphi$ and the sterile neutrino $\nu_s$. Our model is then parametrized by the set

\[(\lambda_\alpha, M_\varphi, m_s).\] (3)

Since this model is in principle subject to a number of constraints from laboratory experiments, cosmology and astrophysics, leading to a restrictions of the free parameter space (3), we take into account the different constraints and we discuss them in the following.

A. Laboratory bounds

It is well known that mesons can decay leptonically as $M \rightarrow \nu_\ell \ell$, where $M$ represents a meson ($\pi^+, K^+, D^+$) and $\ell = e, \mu, \tau$ depending on the meson. The interaction given in eq. (1) opens the possibility of new leptonic decay channels: $M \rightarrow \nu_\ell \ell \varphi$ and $M \rightarrow \nu_\ell \ell \nu_s$. In this regard, we would like to remark that our assumption of small active-sterile mixing angle is crucial, since it prevents the appearance of new decay channels such as $M \rightarrow \nu_\ell \ell \varphi$.

Concerning the process $M \rightarrow \nu_\ell \ell \varphi$, it becomes possible only if the corresponding $\lambda_\alpha \neq 0$; moreover, it is kinematically allowed only if

\[m_s + M_\varphi \lesssim m_M - m_\ell,\] (4)

where $m_M$ is the mass of the decaying meson and $m_\ell$ the mass of charged lepton $\ell$. A lower limit on $m_s + M_\varphi$ arises from Big Bang Nucleosynthesis as discussed below. Using the relation (4), in table (I) we provide the maximal allowed values for $m_s + M_\varphi$, namely $m_M - m_\ell$. We observe that $\pi^+, K^+, D^+ \rightarrow e \nu_\ell \nu_s$ and $K^+, D^+ \rightarrow \mu \nu_\ell \nu_s$ have a large phase space available, while $\pi^+ \rightarrow \mu \nu_\ell \nu_s$ and $D^+ \rightarrow \tau \nu_\ell \nu_s$ are only marginally allowed. At last, the processes $\pi^+, K^+ \rightarrow \tau \nu_\ell \nu_s$ are not kinematically allowed. In the active sector, experimental bounds on meson decay provide limits on $\lambda_\ell$, see for example [6], where similar processes involving active neutrinos $\nu_\ell$ have been studied in details. The main difference between our case and the one studied in [3] is that the relation (3) is replaced with $M_\varphi \lesssim m_M - m_\ell$. In [3], the limit on $\lambda_\ell$ from $K^+ \rightarrow \mu \nu_s \nu_s$ has been found to be stronger with respect to the limit on $\lambda_\ell$ from $K^+ \rightarrow e \nu_s \nu_s$. This depends on the more accurate experimental data available for the former process compared to the latter. Moreover, for masses of the scalar field $\varphi$ smaller than about 300 MeV, the limits on $\lambda_{e,\mu}$ from $K^+$ decay are stronger compared to the one from $\pi^+$ and $D^+$ mesons decay. Concerning the $\lambda_\tau$ coupling, $\pi^+$ and $K^+$ can not provide information because of kinematics, see table (I). We have explicitly analyzed the rate of process $D^+ \rightarrow \tau \nu_\ell \nu_s$, finding that it is very small compared with the experimental bounds. Therefore, we find that for masses $M_\varphi$ and $m_s$ consistent with the cosmological constraints, $\lambda = 1$ is always allowed.

In addition to three-body decay $M \rightarrow \nu_\ell \ell \varphi$ discussed above, there is a further decay channel which occurs as a result of the new interaction: the four-body decay, $M \rightarrow \nu_\ell \ell \nu_s \nu_s$. In this case the kinematics only constrains the mass of the sterile neutrino. In particular, the decay is kinematically allowed if

\[2m_s \leq m_M - m_\ell.\] (5)

If $\lambda < 1$, the rate for four-body decay will be smaller by a factor of $\lambda^2$ compared to the three-body decay. We therefore expect that the four-body decay can relevantly change the exclusions in the parameter space only if $\lambda = 1$. Later we will explicitly verify the validity of the assertion.

We now discuss the explicit form of the decay rate of mesons through the new interactions. As a benchmark case we will discuss the Kaon decay into the muon channel: however, it is straightforward to obtain the decay rate for a different meson decay $M$ into the leptonic channel $\alpha$ by simply replacing in all the subsequent formulas $M_K$ by $m_M$ and $\mu$ by $\alpha$.

\[\begin{array}{lcr}
\text{Meson} & (m_s + M_\varphi)_{\text{max}} & \text{(MeV)} \\
\pi^+ & e\nu_\mu & 140 \\
\rightarrow & \mu\nu_\mu & 35 \\
\rightarrow & \tau\nu_\mu & - \\
K^+ & e\nu_\mu & 493 \\
\rightarrow & \mu\nu_\mu & 388 \\
\rightarrow & \tau\nu_\mu & - \\
D^+ & e\nu_\mu & 1870 \\
\rightarrow & \mu\nu_\mu & 1765 \\
\rightarrow & \tau\nu_\mu & 93 \\
\end{array}\]

Table I. New decay channels for light mesons induced by the interaction and relative maximal allowed values for $m_s + M_\varphi$. When the numerical value is missing it means that the corresponding process is kinematically forbidden.
We begin by examining the three body decay $K^+ \to \mu\nu_s\varphi$. In the limit of vanishing active neutrino masses, the decay rate is

$$d\Gamma_{K\to\mu\nu_s\varphi} = \frac{f_K^2 G_F^2 |\lambda_\mu|^2}{16 M_K (2\pi)^3} \int \frac{dE_p dE_k}{(M_\varphi^2 + m_\mu^2 - 2M_\varphi E_p)^2},$$

where

$$Q = 8M_\varphi^2[2(M_\varphi E_k - E_p E_k + p \cdot k) \times$$

$$+ (2E_p(M_\varphi - E_p) - M_\varphi E_p + m_\mu^2) -$$

$$(M_\varphi^2 + m_\mu^2 - 2M_\varphi E_p)(E_p E_k + p \cdot k)].$$

We have defined with $p, q, k$ and $P$ respectively the $\mu$, $\varphi$, $\nu_s$ and $K$ four-momenta, and in bold face their spatial three-momenta. $f_K$ is the Kaon decay form factor.

With an analogous notation for the four body decay $K^+ \to \mu\nu_s\nu_\tau\nu_\tau$, the decay rate can then be written:

$$d\Gamma_{K\to\mu\nu_s\nu_\tau\nu_\tau} = \frac{G_F^2 f_K^2 |\lambda_\mu|^2(\sum_\lambda |\lambda_\lambda|^2)}{(2\pi)^6 M}$$

$$\int \frac{|p|^2|p||q|^2|q||k|d\cos \theta d\phi}{E_p E_q E_k |p + q|} \frac{q \cdot lp \cdot k}{[(q + l)^2 - M_\varphi^2]^2},$$

where $p, q, l, k$ and $P$ denote the four-momenta respectively of $\mu, \nu_s, \nu_\tau, \nu_\tau$. For convenience, we also define $s = P - p$. In this case, there are 5 independent variables to parameterize the decay, which we chose to be $|p|, |q|, |k|$, the angle $\theta$ between $p$ and $q$, and the azimuthal angle $\phi$ between $k$ and the plane determined by $p$ and $q$. In case $p$ and $q$ were collinear, this should be interpreted as the azimuthal angle around the direction of $p$.

Both the processes $K^+ \to \mu\nu_\tau$ and $K^+ \to \mu\nu_s\nu_\tau\nu_\tau$ should be observed as $K \to \mu +$ missing energy. The closer Kaon decay process that is reported into PDG [7] is $K \to \mu\nu_\tau$ that can be used to constrain our processes. Therefore, we impose that the branching ratio to this channels should be smaller than $2.4 \times 10^{-6}$ [7].

As mentioned in section [1], a reasonable choice for a qualitative picture of the general case is to take $\lambda_\kappa = \lambda_\mu = \lambda_\tau = \lambda$. In Figure 2 we consider this case and we show the region excluded by Kaon decay in the $M_\varphi - m_s$ plane for various values of the coupling. From figure 2 we observe that if

$$\lambda \geq 0.01$$

and $(m_s$ or $M_\varphi) \gtrsim 30 \, MeV$, then the correction to Kaon decay is within the experimental bound.

The four-body decay channel only produces a bump in the right part of the exclusion contours, corresponding to a roughly horizontal line of exclusion that only constrains $m_s$, as expected from our previous considerations: just as we had deduced, this bump is only relevant for $\lambda = 1$.

The only case in which the results are drastically different from the choice of equal couplings for the three flavors is the one in which $\lambda_\tau \neq 0$ and the other two couplings are much smaller than it. In fact, we mentioned above, this case is unconstrained from meson physics and even for very small $\lambda_\tau \sim O(1)$ the only relevant bound in the $M_\varphi - m_s$ plane comes from Big Bang Nucleosynthesis as discussed below.

### B. Cosmological bounds

In addition to laboratory bounds, there could be additional constraints coming from cosmology at different epochs of the Universe. A first constraint comes from the Big Bang Nucleosynthesis (BBN) epoch and specifically from the number of relativistic degrees of freedom. In particular, the requirement is that there are no extra relativistic species (apart from the ones predicted from the Standard Model) at the moment of the BBN. This naturally happens if the new introduced species $\varphi$ and $\nu_s$ are non relativistic and in thermal equilibrium before and during BBN. Indeed, their distributions will be then Boltzmann suppressed by a factor of $e^{-m/\Delta_B}$, where $\Delta_B \approx 1 \, MeV$, and will not count as extra degree of freedom. Another constraint comes from the requirement that the new interaction does not affect the free-streaming nature (non interacting) of the active neutrinos at the time of the formation of the Cosmic Microwave Background (CMB).

A full discussion of the cosmological bounds would require explicit solution of the evolution equations for all the relevant species and is outside the scope of this work. We will limit ourselves to an order of magnitude estimate of the rates of the relevant processes to get a clear idea of which portion of the parameter space is effectively
constrained. Therefore, we did not distinguish between
the three different coupling $\lambda_\alpha$ but we used an effective
coupling $\lambda$. This latter have been chosen as the largest
coupling between the three and so the most relevant.

Big Bang Nucleosynthesis

As mentioned above, the requirements to not affect the
BBN yields are that the newly introduced species are non
relativistic at the time of the BBN and that they remain
in kinetic and chemical equilibrium at this epoch. The
first requirement is naturally met if both $M_\phi$ and $m_\alpha$
are chosen to be larger than about 10 MeV: in this way,
the Boltzmann factor is smaller than $10^{-4}$ and we can
safely assume that the species are non relativistic and
practically absent. Concerning the second requirement,
is necessary to compare the rates of the processes respon-
sible for the equilibrium with the rate of the expansion of
the Universe, in order to determine the temperature
of decoupling at which such processes become irrelevant.

- $\nu_\alpha \nu_\alpha \rightarrow \nu_\alpha \nu_\alpha$ and $\nu_\alpha \nu_\beta \rightarrow \nu_\beta \nu_\beta$: the cross section of
these processes, mediated by $\phi$, can be estimated as
$\lambda^4 T^2 / M_\phi^4$. If we assume a relativistic distribution
for the particles involved, the number density grows as
$n \sim T^3$. The decoupling temperature will be set by the
condition $na \sim \frac{M_\phi^4}{\lambda^2}$, where $M_\phi$ is the Planck
mass. Therefore, the decoupling temperature for sterile neutrinos is

$$T_{\nu_\alpha^{\prime} \nu_\alpha} = \left( \frac{M_\phi^4}{\lambda^2 M_{pl}^4} \right)^{1/3} \simeq 0.3 \text{MeV} \left( \frac{M_\phi}{1 \text{GeV}} \right)^{4/3} \left( \frac{\lambda}{0.01} \right)^{-4/3}$$

BBN will be left unchanged by the new interaction
if this temperature is smaller than the BBN
temperature, namely around 1 MeV. Therefore,
these processes are able to maintain both kinetic
and chemical equilibrium between active and ster-
elle neutrinos, and is therefore sufficient to guarantee
that sterile neutrinos become non relativistic before the
BBN.

- $\phi \phi \rightarrow \nu_\alpha \nu_\alpha$ or $\phi \phi \rightarrow \nu_\beta \nu_\beta$: the most efficient is
the first process mediated by active neutrino. The
cross section is estimated as $\frac{\lambda^4}{m_\phi^4}$, giving a decoupling
temperature

$$T_\phi = \frac{m_\phi^2}{\lambda^4 M_{pl}} \simeq 4 \times 10^{-28} \text{MeV} \left( \frac{m_\phi}{0.1 \text{eV}} \right)^2 \left( \frac{\lambda}{0.01} \right)^{-4}$$

As above, this temperature should be smaller than
1 MeV in order to keep the $\phi$ distribution as a
Maxwell-Boltzmann distribution throughout BBN,
so that the scalar particles have essentially dis-
appeared at the time of the nucleosynthesis. This
process is sufficient to maintain kinetic and chemical
equilibrium among the $\phi$ population, thereby
making them non relativistic before BBN.

In summary, for the parameter space which is of in-
terest to us, both scalar and sterile particles remain
in kinetic and chemical equilibrium throughout the primor-
dial nucleosynthesis. By taking them sufficiently massive,
namely

$$m_\alpha \gtrsim 10 \text{MeV} \quad \text{and} \quad M_\phi \gtrsim 10 \text{MeV}, \quad (10)$$

we can safely satisfy BBN limits, since the newly intro-
duced particles are so massive that they are Boltzmann
suppressed during BBN.

Cosmic Microwave Background

At the time of formation of the Cosmic Microwave
Background (CMB), sterile neutrinos and scalar parti-
cles have long disappeared. Active neutrinos, though,
can still secretly interact through the reactions
$\nu_\alpha \nu_\alpha^{\prime} \rightarrow \nu_\beta \nu_\beta$. There are in principle two ways in which this in-
teraction can proceed through the new interaction: either
the mixing angle between active and sterile neutrinos is
sufficiently large, so that the process $\nu_\alpha \nu_\alpha \rightarrow \nu_\beta \nu_\beta$ can be
converted via mixing to a four active neutrino process;
or the process can happen at next-to-leading order via
the box diagram. We will only analyze the latter process,
since we have assumed very small mixing angles between
active and sterile neutrinos. The cross section can be
estimated in order of magnitude as $\frac{\lambda^4}{m_\phi^4}$, so that the
decoupling temperature for this process is

$$T_{\nu_\alpha^{\prime} \nu_\alpha} = \left( \frac{M_\phi^8 m_\alpha^4}{\lambda^8 M_{pl}^4} \right)^{1/11} \simeq 10^5 \text{eV} \left( \frac{M_\phi}{10 \text{MeV}} \right)^{8/11} \left( \frac{m_\alpha}{10 \text{MeV}} \right)^{4/11} \lambda^{-8/11}.$$

In order to guarantee free-streaming active neutrinos at
CMB time, $T_{\text{dec}}$ has to be larger than the temperature
of CMB formation, around 1 eV. We have checked that
this is the case for all the parameter space we considered.

C. Astrophysical bounds

Another possible constraint we should take into ac-
count comes from the analysis of neutrino fluxes from
supernovae. In fact, since neutrinos in the supernovae
core have energies of order of tenth or hundredth of MeV,
they are sufficiently energetic to produce non relativistic
sterile neutrinos. If these sterile neutrinos interact suf-
fi ciently weakly with the active neutrinos in the core, they
could escape the supernova giving rise to an observable
energy loss. The conditions for this to happen are two:
in the first place, the mean free path of the sterile neutrino
inside the core, namely $(\sigma_{sa} n_\alpha)^{-1}$, with $\sigma_{sa}$ the cross sec-
tion and $n_\alpha$ the number density of active neutrinos in the
core, should be larger than the radius of the supernova
core, typically around 10 km. The cross section is evaluated for active neutrinos with typical energies of tenth of MeV and sterile neutrinos at rest. The number density $n_o$ can be estimated assuming a thermal distribution $f(E)$ of active neutrinos, with a typical temperature of tenth of MeV, as in [8].

The second condition to be verified is that sterile neutrinos should be copiously produced in the supernova core and that the energy injected into sterile neutrinos can exceed the threshold luminosity $L_s \approx 2 \times 10^{32}$ erg/s (namely $L_s \approx 8.2 \times 10^{36}$ MeV$^2$ in natural units, which we have used throughout this work) for the supernova SN 1987A [8]. We then estimate the luminosity in the proposed model as:

$$L_s = \int \frac{d\sigma_{e\nu_s\rightarrow e\nu_s}}{dE} E dE f(E', r) f(E'', r) dE' dE'' 4\pi r^2 dr,$$

where $f(E, r)$ is the distribution of active neutrinos inside the core of the supernova. The temperature profile $T(r)$ is taken from [8].

The model under consideration could be in conflict with SN 1987A data if both the above conditions would simultaneously met. We have numerically verified that this situation never occurs for all the parameter space we considered, with $m_s$ and $M$ larger than 10 MeV. Indeed, for large values of the coupling $\lambda$ the energy injected into sterile sector can easily exceed the threshold indicated above, but the interaction between sterile and active neutrinos is so strong that the mean free path is much shorter than the supernova dimensions. For small values of $\lambda$ we encounter the opposite situation where, even if sterile neutrinos are practically free to escape the supernova, they are produced in too small amounts to be observable. We can therefore deduce that the model we consider is not constrained by supernova data.

### IV. NEUTRINO FLUXES

Active-sterile neutrino interaction can become relevant at very different energy scales depending on the mass of the scalar mediator $\phi$: roughly we expect the energy scale at which the process of absorption over neutrinos from the Cosmic Neutrino Background (CNB) happens resonantly [7] at energies around $M_s^2/m_\nu$. For an active neutrino mass of 0.1 eV, we find that this energy scale can range from PeV to energies of order $10^4$ PeV in the selected parameter space. Close to the PeV scale the dominant source of neutrinos is expected to be constituted by galactic and extragalactic astrophysical sources, among which we mention Active Galactic Nuclei (AGN) and Gamma Ray Bursts (GRB). The details of the emitted neutrino spectra are sensitive to the physics of the sources. However, it is known that a good fit to the observed IceCube data in the region below the PeV is represented by a simple power law spectrum. Therefore, in this range of energy, we limit our discussion on the effect of the new interaction on a power law spectrum with parameters obtained by the fit to the IceCube data given in [7].

At higher energies, from 100 PeV, there are no experimental data on the neutrino flux. It is expected that a dominant source of neutrinos should have cosmogenic origin. On the other hand, recent studies have shown that a competing source of neutrinos could still be of astrophysical nature, provided for example by blazars [10] and Flat Spectrum Radio Quasars (FSRQ) [11].

The new interaction under consideration produces however effects which are qualitatively the same on all these fluxes. For simplicity, we stick to the treatment adopted in our recent paper [2] considering the effects of the new interaction on cosmogenic neutrino fluxes.

#### A. Without secret interaction

**Power Law**

We consider a collection of astrophysical neutrino sources, each one producing a power law spectrum in energy

$$g(E) = N E^{-\gamma},$$

where $g \equiv \phi_{\nu_e} + \phi_{\nu_\mu} + \phi_{\tau_e} + \phi_{\tau_\mu} + \phi_{\tau_\tau}$ and $\gamma$ is the spectral index. The IceCube analysis gives as best fit value for the throughgoing muons data set $\gamma = 2.28$ [12]. Due the similarity with the cosmogenic neutrino production, we found convenient to adopt the Star Forming rate $\rho(z)$ [13] for the cosmological evolution of these sources. The normalization $N$ is chosen in such a way to reproduce the best fit for the diffuse neutrino flux measured by the IceCube Collaboration in the throughgoing muons data sample. Therefore, the diffuse astrophysical spectrum is written as:

$$\frac{d\phi_\nu}{dE d\Omega} = \int \frac{dz'}{H(z')} \rho(z') g(E(1+z')).$$

We assume for definiteness a flavor structure at the source $(1:2:0)$, corresponding to pion beam sources.

Throughout our analysis we use the best fit values from the NuFit 3.2 global fit data for the active oscillation parameters [14], assuming normal neutrino mass ordering.

**Cosmogenic**

Cosmogenic neutrinos are produced by the scattering of high energy protons from the cosmic rays with the CMB photons. The production of cosmogenic neutrinos is quantitatively studied, for example, in [15]. In our

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3 Of course if the sterile mass is too large it can kinematically forbid the process: in determining the energy at which the absorption is most relevant one should take this factor into account.
previous work [2] we showed that their results can be reproduced by parameterizing the cosmogenic neutrino spectrum as:

\[
\frac{d\phi_\nu}{dE \, d\Omega} = \int \frac{dz'}{H(z')} \rho(z') f(E(1 + z')), \tag{13}
\]

where \( \rho(z) \) is the Star Forming rate [13]. We refer the reader to [2] for the method of determination of the function \( f(E) \) describing the energy spectrum.

For cosmogenic neutrinos we again assume a flavor structure at the source \((1:2:0)\).

### B. With secret interaction

Because of the secret interactions, active neutrinos can collide with active neutrinos from the CNB, producing sterile neutrinos and thereby causing a depletion of the flux observable at Earth. The transport equation for active neutrinos is in principle coupled to the transport equation for sterile neutrinos, since the secret interactions produce sterile neutrinos which can in turn collide with other CNB neutrinos to regenerate part of the original flux. The form of these equations has been given in [2], and we reproduce it here for the generalized multiflavor case. We define with \( \Phi_i(z, E) \) the flux of active neutrinos in the \( i \)th \((i = 1, 2, 3)\) mass eigenstate per unit energy interval per unit solid angle at a redshift \( z \), while \( \Phi_s(z, E) \) denotes the flux of sterile neutrinos, where in absence of mixing the sterile mass eigenstate is indicated with \( s \). The flux at Earth is connected with the flux at generic redshift by the relation

\[
\Phi_i(z, E) = \Phi(0, E).
\]

We write separate equations for the mass eigenstates because, as discussed in [2], the propagation is diagonal in the mass eigenstates, given the path between collisions much larger than the oscillation lengths. In other words, due to the very fast oscillations caused by mixing, in between two collisions, a neutrino decoheres to mass eigenstates.

The transport equations take the form:

\[
H(z)(1 + z) \left( \frac{\partial \Phi_i(z, E)}{\partial z} + \frac{\partial \Phi_i(z, E)}{\partial E} \frac{E}{1 + z} \right) = n(z) \sigma_i \Phi_i(z, E)
\]

\[
- \int dE' \Phi_s(z, E') \frac{d\sigma_{as}}{dE}(E' \rightarrow E)n(z)\]

\[
- \rho(z)(1 + z)f(E)\xi_i, \tag{14}
\]

where \( f(E) \) is the neutrino spectrum produced at the source and \( \xi_i \) is the fraction of neutrinos produced at the source in the \( i \)th mass eigenstate. Similarly, for the sterile flux we write:

\[
H(z)(1 + z) \left( \frac{\partial \Phi_s(z, E)}{\partial z} + \frac{\partial \Phi_s(z, E)}{\partial E} \frac{E}{1 + z} \right) = n(z) \sigma_s \Phi_s(z, E)
\]

\[
- \sum_i \int dE' \Phi_i(z, E') \frac{d\sigma_{as}}{dE}(E' \rightarrow E)n(z)
\]

\[
- \int dE' \Phi_s(z, E') \frac{d\sigma_{as}}{dE}(E' \rightarrow E)n(z). \tag{15}
\]

For convenience, we have denoted by \( \sigma_i \) and \( \sigma_s \) the cross sections for the collision of an \( i \)th mass eigenstate and a sterile neutrino, respectively, with a CNB neutrino. Correspondingly, \( \frac{d\sigma_{as}}{dE}(E' \rightarrow E) \) is the cross section for the production of a \( \beta \) neutrino with energy \( E \) after the collision of an \( \alpha \) neutrino with energy \( E' \) with a CNB neutrino.

If the regeneration processes play an important role, the task of determining the effect of the interaction is computationally expensive, since it requires the numerical solution of the system of four coupled partial integro-differential equations.

In our previous paper [2], we found that the regeneration was unimportant for a limited region of the parameter space, with masses of sterile neutrino and scalar mediator around 250 MeV. Here we have analyzed this question more thoroughly, taking in consideration a wider parameter space. We have adopted a perturbative approach in which the regeneration processes are treated as a perturbation and we have tested its validity \( a \) posteriori by comparing the perturbation induced by regeneration with the unperturbed flux.

We find that both cosmogenic and astrophysical fluxes are practically unaffected by regeneration. The physical reason behind this behavior is connected with the cosmological evolution of the sources, and in particular with the fact that the sources are distributed at various redshifts. In fact, when neutrinos produced at high redshifts, with \( z \gg 0.1 \), are severely suppressed due to the absorption on the CNB, neutrinos produced at low redshifts are only weakly absorbed. Thus the flux has always a component, produced at low redshift, which is roughly unabsorbed and which dominates against the small regenerated flux produced at high redshifts. The perturbative approach shows in fact that the corrections coming from regeneration, both for cosmogenic and astrophysical fluxes, are typically not larger than about 10%.

The negligible effect of regeneration is therefore connected with the presence of sources at small redshifts, masquerading the regenerated flux. Thus we expect that, for point-like sources localized at large redshifts, regeneration effects should instead be non negligible. Even though IceCube has identified so far a single realistic candidate of point-like astrophysical source, in the future one may expect noticeable improvements in this respect. Therefore, it might be interesting to have a qualitative idea of the effect of regeneration on the neutrino spectra from point-like sources. In Figure 3 we show the spectra expected at Earth for a generic source at two fixed red-
shift values $z$, namely 0.1 and 0.01, with an $E^{-2}$ reference spectrum. The effects of regeneration are, as expected, more important for larger redshifts of the source and can drastically change the results.

and astrophysical neutrino fluxes, we neglect the regeneration processes, so that the transport equation for the sole active neutrinos is given by

$$H(z)(1 + z) \left[ \frac{\partial \Phi_i}{\partial z} + \frac{\partial \Phi_i}{\partial E} \frac{E}{1 + z} \right] = n(z) \sigma_i(E) \Phi_i(E) + \rho(z) f(E) \xi_i$$

(16)

In this equation $\sigma_i$ is the cross section of the process $\nu_i + \nu_j \to \nu_\alpha + \nu_\beta$, namely

$$\sigma_i = \frac{1}{64\pi t_j^2} \sum_j \int_{t_1}^{t_2} |M_{ij \to \alpha \beta}|^2 (s,t) dt$$

(17)

in terms of the Mandelstam invariants $s = (p + l)^2$, $t = (p - k)^2$ and $u = (p - q)^2$ with $p$, $l$, $k$ and $q$ the momentum of the two active neutrinos and the two sterile neutrinos, respectively. Moreover

$$t_{1,2} = m^2 + m_s^2 - \frac{s}{2} \pm \sqrt{s^2 - 4m_s^2},$$

(18)

where $m_s$ is the mass of the active neutrino $\nu$ of CNB. The squared amplitude is given by

$$|M_{ij \to \alpha \beta}|^2 = |\sum_{\alpha, \beta} U_{\alpha j}^* U_{\beta i}^* \lambda_\alpha \lambda_\beta|^2$$

$$\times \left[ \frac{|t - (m - m_s)^2|^2}{(t - M_\varphi^2)^2 + \lambda^2 M_\varphi^2} + \frac{|u - (m - m_s)^2|^2}{(u - M_\varphi^2)^2 + \lambda^2 M_\varphi^2} \right]$$

$$- \frac{2[(t - M_\varphi^2)(u - M_\varphi^2) + \lambda^2 M_\varphi^2]}{[(t - M_\varphi^2)^2 + \lambda^2 M_\varphi^2][(u - M_\varphi^2)^2 + \lambda^2 M_\varphi^2]}$$

$$\times \left( \frac{(t - m^2 - m_s^2)^2}{4} + \frac{(u - m^2 - m_s^2)^2}{4} \right)$$

$$- \frac{s^2}{4} + s(m^2 + m_s^2 - m m_s) - 2m^2 m_s^2 \right] \right) \right]$$

(19)

where $\Gamma$ is the decay rate of the scalar mediator and $M_\varphi$ is its mass.

Therefore Eq.[16] contains only an absorption term and, for the astrophysical power law neutrino flux, it admits an analytical solution for the flux at Earth given by

$$\Phi_i(E) = \int_0^{+\infty} \frac{dz}{H(z)} \rho(z) g[E(1 + z)] \times \exp \left[ - \int_0^{z'} \frac{dz''}{H(z'')(1 + z'')} n(z'') \sigma_{\nu_i}[E(1 + z'')] \right] \xi_i$$

(20)

For cosmogenic neutrino fluxes the solution is identical with we consider the function $f(E)$ in place of $g(E)$.

V. RESULTS

A. Power law

We start our discussion of the results with the case of a power law astrophysical spectrum in the energy

Figure 3. Comparison between the spectra with pure absorption (dotted lines) and with both absorption and regeneration (dashed lines) for an $E^{-2}$ flux produced by a source at redshift 0.1 (top panel) and 0.01 (bottom panel). The thick line is the unabsorbed spectrum, reproduced for reference. The sterile and scalar mediator masses are fixed to the benchmark values of 10 MeV; the coupling $\lambda$ is chosen as 1 for the tau neutrinos.

A final comment on the regeneration effects is necessary in relation to the very small active-sterile mixing angle we adopted. In fact, as we mentioned above, in the regime $m_s > M_\varphi$, the mediator decay into an active and sterile neutrino is inhibited. Under these conditions we consider a mediator decay into two active neutrinos via a very small active-sterile mixing angle. We evaluated the dependence of negligible regeneration on the value of the mixing angle. In general large active-sterile mixing angles would lead to smaller regeneration effects, because they increase the decay rate of the mediator. On the other hand we find that even very small mixing angles still produce no relevant regeneration effects. Therefore, our conclusion that regeneration can be neglected for diffuse fluxes is still valid. Since the active-active cross section does not exhibit any t-channel resonance, it does not depend on the precise value of the mixing angle, making our results roughly independent of this assumption.

In the following, since we only deal with cosmogenic
range roughly below 100 PeV. In this region the effects of active-sterile interaction can be detected only if the sterile mass is sufficiently low that the process is not kinematically forbidden. As in the previous sections, we distinguish between the two possibilities: either $\lambda_e = \lambda_\mu = \lambda_\tau = \lambda_{af}$ (where $af$ denotes all flavors), or $\lambda_e = \lambda_\mu = 0$ and $\lambda_\tau \neq 0$.

In the first case, the Kaon decay strongly constrains the possible values of the coupling. In particular, we find that the optimal choice to have noticeable effects below 100 PeV is to have small sterile masses, large scalar masses and $\lambda_{af} = 1$. We take as benchmark values $m_s = 10$ MeV and $M_\varphi = 1$ GeV.

In the second case, in which the mediator only couples to tau neutrinos, the constraints from Kaon decay are irrelevant and we can also consider lower masses for $M_\varphi$. In order to maximize the effect in this energy range, we have chosen the benchmark values of $m_s = 15$ MeV, $M_\varphi = 10$ MeV and $\lambda_\tau = 1$, as represented in Fig. 4 where we show the neutrino spectra after the new interaction for both the choices of $\lambda$, together with the IceCube HESE data [16]. From this Figure we can infer that the second possibility is already testable using IceCube data while the first case is essentially undistinguishable from the power-law in the energy range probed by IceCube.

The effects of the new interaction can also cause significant changes in the flavor structure of the spectrum since the induced depletion acts differently on each flavor modifying the flavor ratio, namely the fraction of electron, muon and tau neutrino fluxes. Since the depletion is energy dependent, the result will be an energy dependent flavor ratio. In Figures 5 we show the flavor ratios as a function of the energy for the two cases, $\lambda_{af}$ and $\lambda_\tau$ respectively. We see that the case of only $\lambda_\tau$ has a threshold behavior with a sudden change of the flavor ratio. This change is quite relevant, especially when compared with the change in the case $\lambda_{af}$. We remind the reader that we assumed a flavor ratio at the source (1 : 2 : 0); at low energies, where the effects of the interaction are inactive, we recover the typical flavor structure (1 : 1 : 1) at the Earth as expected.

![Figure 4](image-url)  
**Figure 4.** Astrophysical all flavor neutrino power law spectra: the thick line is the flux with no interaction, while the dotted line corresponds to the case of $\lambda_{af} = 1$ and the dashed one denotes the case $\lambda_\tau = 1$, as described in the text. The experimental points are the IceCube HESE data.

![Figure 5](image-url)  
**Figure 5.** Flavor ratio at Earth as a function of the energy for the first benchmark case in the text ($m_s = 10$ MeV, $M_\varphi = 300$ MeV, $\lambda_e = \lambda_\mu = \lambda_\tau = 1$) (top panel) and second benchmark case in the text ($m_s = 15$ MeV, $M_\varphi = 10$ MeV, $\lambda_\tau = 1$) (bottom panel).

An interesting aspect revealed by Figure 4 is that the new interaction causes a cutoff-like feature in the spectrum in the range between 1 PeV and 10 PeV. In fact, the second case with only $\lambda_\tau$, shows a sudden drop of the flux at the energy at which the process $\nu_\alpha \nu_\beta \rightarrow \nu_\alpha \nu_s$ becomes kinematically allowed.

The effects of the interaction on the flavor structure of the spectrum, namely the modifications of the flavor...
ratio, can also be represented in the flavor triangle: we show this for case $\lambda_\tau = 1$, which has the largest effect, in Figure 6.

The red and the orange points correspond to an energy of $10^5$ GeV and $10^8$ GeV, respectively. The flavor sensitivity which has been forecasted for IceCube-Gen2 [17], in the case of pion beam sources with a flavor ratio $(1 : 2 : 0)$ at the source, has been shown as well. The triangle representation suggests the possibility that future experiments might be able to unveil a different flavor structure possibly caused by active-sterile secret interactions. It is worth noticing that this change induced by the interaction is also dominant with respect to the uncertainty due to the mixing parameters.

A fundamental feature of the change in flavor induced by secret interactions is that it has a unique energy dependence, which descends from the resonances and thresholds of the interaction. Since the data from IceCube-Gen2 might allow in the future to investigate the interplay between flavor and energy, this is a result which might be of experimental interest.

Figure 6. Flavor ratio, reproduced in the flavor triangle, for varying energy for the second case in the text ($m_s = 15$ MeV, $M_\phi = 10$ MeV, $\lambda_e = 1$). The red and orange points correspond to an energy of $10^5$ GeV and $10^8$ GeV, respectively. The forecasted sensitivity of IceCube-Gen2 is shown as well.

**B. Cosmogenic**

For the case of cosmogenic fluxes, which are relevant at higher energies, the absorption effect is most important for higher masses of the sterile neutrino and of the scalar mediator. In this part of the parameter space, the constraints from mesons decay are substantially irrelevant, so there is no need to distinguish between the two case studied above. We will therefore analyze as a single choice the case $\lambda_e = \lambda_\mu = \lambda_\tau = 1$, $m_s = 250$ MeV and $M_\phi = 250$ MeV. In Figure 7 we show the effect of the interaction on the cosmogenic flux. We can observe that the effect is maximal around $10^9 \div 10^{10}$ GeV. We address the reader to our previous paper [2] for more details also in relation to future experiments.

Figure 7. Cosmogenic all flavor neutrino power law spectra: the thick line is the flux with no interaction, while the dashed line correspond to the benchmark case described in the text.

As in the astrophysical neutrino case, also for the cosmogenic flux we analyze the flavor structure as a function of energy, as shown in Figures 8.

Figure 8. Flavor ratio at Earth as a function of the energy for the cosmogenic benchmark case in the text ($m_s = 250$ MeV, $M_\phi = 250$ MeV, $\lambda_e = \lambda_\mu = \lambda_\tau = 1$).
VI. CONCLUSIONS

In this work we have investigated the effects on high-and ultra high-energy active neutrino fluxes due to active-sterile secret interactions mediated by a new pseudoscalar particle. In particular, we extended our previous paper [2] considering three active neutrino flavors, leading to an ample freedom of choice for the couplings of the interactions. As a consequence, the laboratory constraints for the meson decays are more relaxed, allowing new regions of the parameter space otherwise forbidden. Active-sterile neutrino interactions become relevant at very different energy scales depending on the masses of the scalar mediator and of sterile neutrino. As already found in our previous paper, the final active fluxes can present a measurable depletion observable in future experiments. However, in this study we find that the flux depletion can also occur at lower energy, around the PeV, in the particular case of only $\nu_\tau \rightarrow \nu_s$ interaction. We adopted then two prescriptions for the neutrino flux, namely high-energy represented by a power-law, and ultra high-energy with cosmological origin (cosmogenic) in order to take into account the multi-scale energy phenomenology, varying from $10^6$ GeV up to $10^9$ GeV. Remarkably, when the depletion is around $10^6$ GeV, this effect could be very interesting for IceCube because it can produce a cutoff-like feature in the spectrum, which could potentially explain the lack of observed events above 10 PeV. For larger values of mediator and sterile masses, the depletion effect instead could be only visible at larger energy, around $10^8$ GeV, with future experiments like GRAND. Another interesting phenomenological aspect of active-sterile secret interactions is represented by the changing in the flavor ratio as a function of neutrino energy. This effect could be interesting for next generation of neutrino telescopes like IceCube-Gen2 or KM3NeT.

Acknowledgments: This work was partially supported by the research grant number 2017W4HA7S “NAT-NET: Neutrino and Astroparticle Theory Network” under the program PRIN 2017 funded by the Italian Ministero dell’Università e della Ricerca (MUR). The authors acknowledge partial support by the research project TAsP (Theoretical Astroparticle Physics) funded by the Instituto Nazionale di Fisica Nucleare (INFN).

[1] V. S. Berezinsky and G. T. Zatsepin, Phys. Lett. 28B, 423 (1969).
[2] D. F. Fiorillo, G. Miele, S. Morisi, and N. Saviano, Phys. Rev. D 101, 083024 (2020), arXiv:2002.10125 [hep-ph].
[3] J. Alvarez Muniz et al. (GRAND), Sci. China Phys. Mech. Astron. 63, 219501 (2020), arXiv:1810.09994 [astro-ph.HE].
[4] A. Anker et al., JCAP 03, 053 (2020), arXiv:1909.00840 [astro-ph.IM].
[5] M. Aartsen et al. (IceCube), Science 342, 124256 (2013), arXiv:1311.5238 [astro-ph.HE].
[6] J. M. Berryman, A. De Gouvea, K. J. Kelly, and Y. Zhang, Phys. Rev. D 97, 075030 (2018), arXiv:1802.00009 [hep-ph].
[7] T. et al. (Particle Data Group), Phys. Rev. D 98, 030001 (2018).
[8] L. Mastrototaro, A. Mirizzi, P. D. Serpico, and A. Esmaeil, JCAP 01, 010 (2020), arXiv:1910.10249 [hep-ph].
[9] D. R. Williams (IceCube), Int. J. Mod. Phys. Conf. Ser. 46, 1860048 (2018), arXiv:1909.05173 [astro-ph.HE].
[10] K. Murase, Y. Inoue, and C. D. Dermer, Phys. Rev. D 90, 023007 (2014), arXiv:1403.4089 [astro-ph.HE].
[11] C. Righi, A. Palladino, F. Tavecchio, and F. Vissani, (2020), arXiv:2003.08701 [astro-ph.HE].
[12] A. Schneider, PoS ICRC2019, 1004 (2020), arXiv:1907.11266 [astro-ph.HE].
[13] A. M. Hopkins and J. F. Beacom, Astrophys. J. 651, 142 (2009), arXiv:astro-ph/0601463.
[14] I. Esteban, M. Gonzalez-Garcia, M. Maltoni, I. Martinez-Soler, and T. Schwetz, JHEP 01, 087 (2017), arXiv:1611.01514 [hep-ph].
[15] M. Ahlers and F. Halzen, Phys. Rev. D 86, 083010 (2012), arXiv:1208.4181 [astro-ph.HE].
[16] E. Perotti, P. Blasi, F. Aharonian, G. Morlino, and P. Cristofari, Mon. Not. Roy. Astron. Soc. 493, 5880 (2020), arXiv:1911.06163 [astro-ph.HE].
[17] M. Bustamante and M. Ahlers, Phys. Rev. Lett. 122, 241101 (2019), arXiv:1901.10087 [astro-ph.HE].