

Superheavy supersymmetric dark matter as the origin of the KM3NeT ultrahigh energy signal

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We propose an explanation for the recently reported ultrahigh-energy neutrino signal at KM3NeT, which shows no clear association with known astrophysical sources. While decaying dark matter in the Galactic Center is a natural candidate, the observed arrival direction strongly suggests an extragalactic origin. We introduce a multicomponent dark matter scenario in which the components are part of a supermultiplet, with supersymmetry ensuring a nearly degenerate mass spectrum among the fields with different spins. In this setup, a cosmologically long-lived fermionic state decays into a slightly lighter bosonic dark matter state, producing a boosted neutrino spectrum with energy $E_\nu \sim 100$ PeV, determined by the mass difference. The heavy-to-light decay occurs at a cosmological redshift of $z \sim$ a few or higher, leading to an isotropic directional distribution of the signal.

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I. INTRODUCTION

The KM3NeT Collaboration has recently reported their investigation concerning the detection of an ultrahigh-energy neutrino event, designated as KM3-230213A [1]. The neutrino energy for this event is estimated to be $E_\nu = 220^{+2380}_{-146}$ PeV with a 90% confidence interval. This represents the most energetic neutrino observed to date. Multiple hypotheses have been proposed regarding the source of KM3-230213A [2–19]. Notably, no definitive associations with particular astrophysical point origins within galactic and extragalactic regions have been established so far [20,21].

We propose the decay of superheavy dark matter as a potential explanation for the observed signal at KM3NeT. Specifically, we focus on the distinctive characteristics of the signal:

- (i) The neutrino energy: $E_\nu \sim O(100)$ PeV.
- (ii) The signal originates from a direction nearly opposite to the Galactic Center.

Clearly, the second feature requires an explanation for the absence of a signal from the Galactic Center, which would naturally be expected in a conventional dark matter interpretation [22,23].

To address these challenges, we propose a multicomponent dark matter scenario with a heavy-to-light mass degeneracy, $\Delta M/M \ll 1$. Notably, this framework naturally emerges in a supersymmetric (SUSY) setting [24–26], where mass degeneracy is an intrinsic feature of supersymmetry. Moreover, when the SUSY-breaking scale is significantly smaller than the mass term, small mass splittings remain radiatively stable [27–31]. In this scenario, the heavier component decays into the lighter one, producing a boosted neutrino spectrum with energy $E_\nu \sim \Delta M \ll M$.

The near-degeneracy ensures that the dark matter produced from these decays remains nonrelativistic, preserving its role as cold dark matter. Furthermore, the lifetime of the heavier component can be easily controlled, allowing it to be safely set shorter than the age of the Universe. This guarantees that the resulting neutrino spectrum originates entirely from extragalactic sources. Given these characteristics, we refer to this scenario as *superheavy supersymmetric dark matter*.

This letter is organized as follows: In Sec. II, we discuss the conventional decaying dark matter (DM) scenario, where the expected signal from the Galactic Center is typically 510 times stronger than that from the opposite

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direction. In Sec. III, we present our two-component superheavy dark matter model and explain how it accounts for the KM3NeT signal. In Sec. IV, we analyze the expected signatures in neutrinos and associated gamma-rays, considering current observational constraints. Finally, in Sec. V, we summarize our scenario for the KM3NeT ultrahigh-energy (UHE) event and discuss its implications for UHE cosmic rays (UHECRs), UHE neutrinos (UHE ν), and future research directions.

II. CONVENTIONAL DARK MATTER SCENARIOS

In this section, we discuss the conventional decaying dark matter scenario. As we see in top panel of Fig. 1, the angular distribution of Galactic component of neutrinos

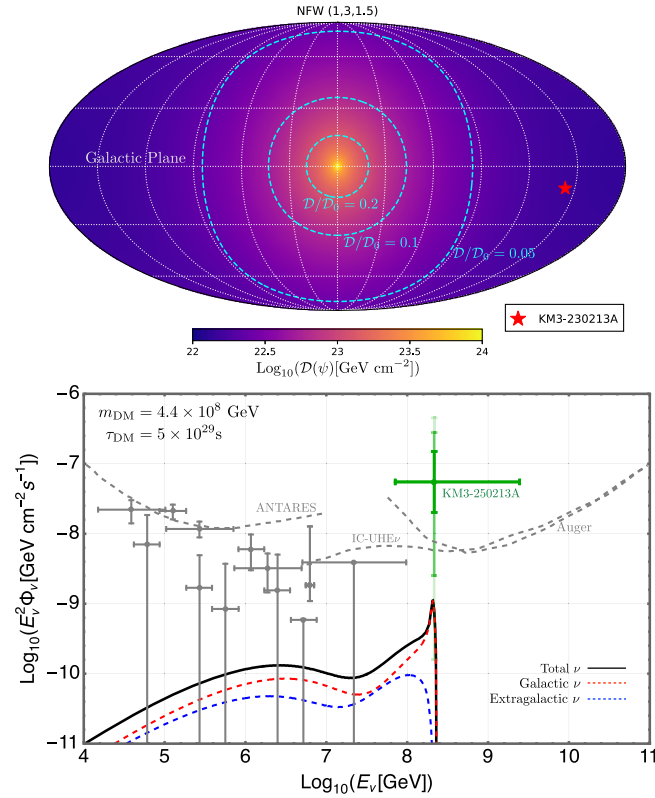


FIG. 1. The angular distribution of galactic component of neutrino flux from conventional decaying dark matter using NFW(1,3,1.5) profile (top) and the neutrino flux spectrum (bottom). In the top panel, we show the $\mathcal{D}$ -factor $\mathcal{D}(\psi) = \int \rho(l, \psi) dl$ on the Galactic coordinate map. Cyan dashed lines correspond to $\mathcal{D}(\psi)/\mathcal{D}(\psi_0) = 0.2, 0.1, 0.05$ from inner to outer contours where ψ is the angle between Galactic Center and the direction of line-of-sight. Here we consider the cutoff angle $\psi_0 = 1^\circ$ for the normalization. We overlap the direction of KM3-250213A neutrino origin as a red-star marker whose position is $(l, b) = (216.06^\circ, -11.13^\circ)$ in the Galactic coordinates. In the bottom panel, we show galactic (dashed red) and extragalactic (dashed blue) contribution to neutrino flux, and observational constraints on the diffuse flux from IceCube [32–35], Auger [36] and ANTARES [37] (gray lines).

from decaying DM is concentrated around the Galactic Center. Since the angle between the KM3NeT UHE ν 's direction (RA 94.3° , Dec -7.8°) and the Galactic Center (RA 266.4° , Dec -29.0°) is $\psi_{\text{KM3NeT}} = 142.5^\circ$, it is natural to see more energetic neutrinos at the Galactic Center. Even including extragalactic contribution, this anisotropy is not dramatically relaxed. For instance, let us define a ratio between energy-integrated fluxes for different two directions as

$$\mathcal{P} \equiv \frac{\int_{E_{\min}}^{E_{\max}} dE \frac{d^2\Phi}{dE d\Omega} \Big|_{\text{GC}}}{\int_{E_{\min}}^{E_{\max}} dE \frac{d^2\Phi}{dE d\Omega} \Big|_{\text{KM3}}} \quad (1)$$

We read the ratio $\mathcal{P}$ for three different DM profiles [NFW(1,3,1.5), NFW(1,3,1) and Isothermal) and two different energy windows ([72, 2600] PeV (90CL) and [170, 270] PeV (Narrow)] as

$$\begin{aligned} \mathcal{P}_{\text{NFW1315}}^{\text{Narrow}} &= 79.6, & \mathcal{P}_{\text{NFW1315}}^{90\text{CL}} &= 75.1, \\ \mathcal{P}_{\text{NFW131}}^{\text{Narrow}} &= 12.8, & \mathcal{P}_{\text{NFW131}}^{90\text{CL}} &= 9.0, \\ \mathcal{P}_{\text{Iso}}^{\text{Narrow}} &= 7.9, & \mathcal{P}_{\text{Iso}}^{90\text{CL}} &= 5.3, \end{aligned}$$

We use the cutoff angle $\psi_0 = 1^\circ$ to estimate GC flux for cuspy profiles. The parametrization NFW(α, β, γ) is defined as

$$\rho(r) = \rho_0 \cdot \left(\frac{r_s}{r}\right)^\gamma \left(1 + \left(\frac{r}{r_s}\right)^\alpha\right)^{\frac{\gamma-\beta}{\alpha}}. \quad (2)$$

As we see in bottom panel of Fig. 1, the spectral shape of flux is more narrow in the galactic contribution compared to the extragalactic flux, due to the redshift effect. In any cases, one expects the expected different number of events with a factor of 5–80 around $E_\nu \sim \mathcal{O}(100)$ PeV in the conventional decaying DM scenario. Nevertheless, since we have only one event within the $\mathcal{O}(100)$ PeV energy window, it obviously requires more enhanced statistics based on observational dataset with a larger effective area experiments. For now, let us keep this possibility, and proceed to construct a new scenario which gives the isotropic and extragalactic-dominant signals at high energies.

III. MODEL

A. Dark matter mass spectrum

The basic scheme of our model is shown in Fig. 2. We consider a dark matter supermultiplet, whose dominant mass term is given by the supersymmetric mass M . We allow the possibility that M can be as large as a grand unified scale, i.e., 10^{16} GeV. Since the SUSY breaking scale in our set-up is much smaller than M , it generically predict highly degenerated mass spectrum among dark matter components.

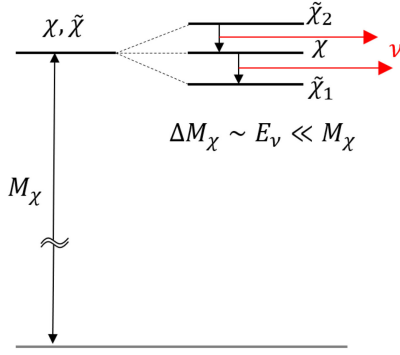


FIG. 2. Scheme of a model for superheavy supersymmetric dark matter. The dark matter supermultiplet possesses a large supersymmetric mass term, M_χ , with a small mass splitting induced by SUSY-breaking effects. This small mass splitting remains radiatively stable compared to the dark matter mass. Each component field contributes equally to the dark matter density through their supersymmetric production mechanism. The long-lived heavier component of the supermultiplet decays into the lightest one, emitting a highly energetic neutrino (directly, secondarily, or via oscillation) with energy $E_\nu \sim \Delta M_\chi \ll M_\chi$ at production time. Consequently, any late-time decay of χ and $\tilde{\chi}_2$ remains consistent with cosmological constraints on dark matter density.

Dark matter is given as the component field of the SM singlet chiral superfield X [38],

$$X = \tilde{\chi} + \sqrt{2}\theta\chi + \theta^2 F_\chi. \quad (3)$$

Here, $\tilde{\chi}$ is a complex scalar, χ is its fermionic superpartner, and F_χ is the auxiliary scalar field.

Since we consider superheavy dark matter with a mass much greater than the SUSY breaking scale, it is natural that the dominant mass contribution arises from the superpotential

$$W = \frac{1}{2}\mathcal{M}X^2, \quad (4)$$

where the spurion superfield $\mathcal{M}$ is given by

$$\mathcal{M} = M + \theta^2 F. \quad (5)$$

The scalar component of $\mathcal{M}$ provides a supersymmetric mass M for both the scalar and fermionic components, while the F -term contribution induces mass splitting among the scalars.

The spurion method efficiently captures the SUSY breaking effects in the mass spectrum of the supermultiplet as long as $M^2 \gg F$. Superheavy supersymmetric dark matter is realized for sufficiently large M . The quadratic potential for the scalar and fermion components is obtained as

$$V = \frac{1}{2}M(\chi\chi + \text{H.c.}) + M^2|\tilde{\chi}|^2 + \frac{1}{2}F(\tilde{\chi}^2 + \text{H.c.}). \quad (6)$$

The resulting mass splitting between the scalar and fermionic components is given by

$$M_{\tilde{\chi}_{1,2}} = M \pm \frac{F}{M}, \quad M_\chi = M. \quad (7)$$

Thus, the relative mass difference is

$$\frac{\Delta M_\chi}{M_\chi} \equiv \frac{M_\chi - M_{\tilde{\chi}_1}}{M_\chi} = \frac{M_{\tilde{\chi}_2} - M_\chi}{M_\chi} = \frac{F}{M^2} \ll 1. \quad (8)$$

This very degenerate mass spectrum remains radiatively stable as long as $F \ll M^2$, since F serves as the order parameter for SUSY breaking, ensuring that all other mass splittings are also proportional to F .

The lightest component of X (i.e., $\tilde{\chi}_1$) can serve as a stable dark matter candidate due to the $\mathbb{Z}_2$ symmetry, $X \rightarrow -X$. However, if the interaction between X and other fields is sufficiently suppressed, the heavier states χ and $\tilde{\chi}_2$ may also contribute to dark matter due to their long lifetimes. We can first think of the interaction that determines the dominant decay channel of heavier dark matter.

B. Neutrino flux from the decay of heavier dark matter components

Although the production mechanism of heavy dark matter in the early universe is not unique [39], it is reasonable to expect that the supersymmetric contribution dominates over the SUSY-breaking one, as dark matter is primarily produced through high-energy interactions exceeding the SUSY-breaking scale. Therefore, it is natural to assume that the initial abundances of $\tilde{\chi}_1$, χ , and $\tilde{\chi}_2$ are approximately equal, i.e., $\rho_{\tilde{\chi}_{1,2}} \sim \rho_\chi$.

The heavier components eventually decay into $\tilde{\chi}_1$. Due to the nearly degenerate mass spectrum, the resulting $\tilde{\chi}_1$ is not boosted and remains a component of cold dark matter. The key parameters governing the neutrino flux from these decays are the fraction of heavy components in the dark matter density, f_χ , the total decay rate of the heavier component, Γ_χ , and the decay rate of the heavier component into active neutrinos, given by $\text{Br}_\nu \cdot \Gamma_\chi$.

The isotropic extragalactic neutrino flux is then expressed as

$$\begin{aligned} \frac{d\Phi_\nu}{dE_\nu} &= \frac{\text{Br}_\nu \Gamma_\chi f_\chi \rho_{\text{DM}c}}{4\pi M_\chi H_0} \\ &\times \int_0^{z_{\text{max}}} \frac{dz}{1+z} \frac{e^{-\Gamma_\chi t(z)}}{\sqrt{\Omega_m(1+z)^3 + \Omega_\Lambda}} \frac{dN_\nu}{dE'_\nu} \end{aligned} \quad (9)$$

where $E'_\nu = (1+z)E_\nu$. The upper limit of the redshift integral, z_{max} , can, in principle, extend up to the neutrino decoupling redshift, $z_{\text{dec}} \simeq 10^{10}$, but for the parameter range of interest, the dominant contribution typically arises from $0 \leq z \lesssim 10$. The present-day dark matter density is

given by $\rho_{\text{DM}} = 1.15 \times 10^{-6} \text{ GeV cm}^{-3}$, and the Hubble radius is $c/H_0 = 1.37 \times 10^{28} \text{ cm}$. The matter and dark energy contributions to the total energy density are $\Omega_m = 0.315$ and $\Omega_\Lambda = 0.685$, respectively [40]. Since we consider the case where all heavier components decay before the present time, i.e., $\Gamma_\chi < H_0$, only the extragalactic contribution remains.

The necessary conditions for the branching fraction and mass splitting can be estimated from observations. The observed KM3NeT neutrino flux is given by

$$E_\nu^2 \frac{d\Phi_\nu}{dE_\nu} = 5.8_{-3.7}^{+10.1} \times 10^{-8} \left(\frac{\text{GeV}}{\text{cm}^2 \text{ sec sr}} \right) \quad (10)$$

for $E_\nu = (0.072 - 2.6) \times 10^9 \text{ GeV}$ [1]. Given the present dark matter density and the relation $E_\nu = \Delta M_\chi / (1 + z_\chi)$, where z_χ is the redshift at $t = 1/\Gamma_\chi$, the estimated flux for our model is

$$\begin{aligned} E_\nu^2 \frac{d\Phi_\nu}{dE_\nu} &\sim \frac{\text{Br}_\nu f_\chi}{4\pi M_\chi (1 + z_\chi)} \rho_{\text{DM}} c \\ &= 2.5 \times 10^{-8} \left(\frac{\text{GeV}}{\text{cm}^2 \text{ sec sr}} \right) \\ &\quad \times \left(\frac{\text{Br}_\nu f_\chi}{10^{-3}} \right) \left(\frac{\Delta M_\chi}{10^{-7} M_\chi} \right) \left(\frac{10}{1 + z_\chi} \right). \end{aligned} \quad (11)$$

A degenerate mass spectrum of this kind can naturally arise within supersymmetry, while the decay products and lifetimes depend on the specific model construction. The SUSY breaking mass $F/M \sim 10^8 - 10^9 \text{ GeV}$ is still much higher than the weak scale. This is consistent with the null result of SUSY searches [41–43]. The gravitino and superpartners of the Standard Model particles are heavy enough that they can decay much earlier than the BBN epoch if the R -parity violating terms exist. Therefore only X supermultiplet remains as a dark matter. A detailed discussion of R -parity violation (RPV) and its implications for BBN and proton stability is provided in the Appendix.

To allow for variations in the branching fraction into active neutrinos, we assume that the dominant decay channel of the dark matter supermultiplet is invisible. Introducing a singlet chiral superfield,

$$S = \tilde{\nu}_s + \sqrt{2}\theta\nu_s + F_s, \quad (12)$$

the superpotential governing the decay of the heavier dark matter components is

$$W_D = \frac{1}{2} \lambda_s S X^2, \quad (13)$$

where λ_s is a free parameter. For a mass much smaller than M_χ , this term induces the decays $\chi \rightarrow \tilde{\chi}_1 + \nu_s$ and

$\tilde{\chi}_2 \rightarrow \chi + \nu_s$. Assuming this decay channel dominates, the heavier dark matter decay rate is approximately

$$\Gamma_{\tilde{\chi}_2 \rightarrow \chi \nu_s} \simeq \Gamma_{\chi \rightarrow \tilde{\chi}_1 \nu_s} \simeq \frac{\lambda_s^2 M_\chi}{8\pi} \left(\frac{\Delta M_\chi}{M_\chi} \right)^2. \quad (14)$$

Note that the degeneracy provides a strong kinetic suppression for the decay rate.

We consider different possibilities for the decay into visible particles. The first case is that dark matter has a direct interaction with neutrinos. The corresponding minimal superpotential is

$$\Delta W_I = \frac{L H_u X^2}{2\Lambda}. \quad (15)$$

where Λ is assumed to be much larger than the Planck scale. This term leads to the component Lagrangian

$$\frac{\nu h \tilde{\chi} \chi}{\Lambda} + \text{H.c.} \quad (16)$$

which allows for three-body dark matter decay via the channels

$$\tilde{\chi}_2 \rightarrow \chi + \nu + h, \quad \chi \rightarrow \tilde{\chi}_1 + \nu + h \quad (17)$$

with the decay rate

$$\Gamma_{\tilde{\chi}_2 \rightarrow \chi \nu h} \simeq \Gamma_{\chi \rightarrow \tilde{\chi}_1 \nu h} \simeq \frac{M_\chi^3}{24\pi^3 \Lambda^2} \left(\frac{\Delta M_\chi}{M_\chi} \right)^4. \quad (18)$$

A similar decay channel to neutrinos from long-lived dark matter components has been considered without supersymmetry [44].

The second case is that the decay into active neutrinos is indirect. The simplest possibility is that the active neutrino is generated from the oscillation of the singlet fermion ν_s via

$$\Delta W_{II} = y_s L H_u S. \quad (19)$$

This leads to the decay $\chi \rightarrow \tilde{\chi}_1 + \nu_s$ and $\tilde{\chi}_2 \rightarrow \chi + \nu_s$, followed by $\nu_s \rightarrow \nu$ via oscillation [45], with a mixing angle determined by y_s and the masses of neutrinos, assuming the sterile neutrino mass is sufficiently small, on the order of eV. This ultimately results in a two-body decay spectrum

$$\chi \rightarrow \tilde{\chi}_1 + \nu \quad (20)$$

with a specific branching fraction into neutrinos, Br_ν .

C. Production mechanisms in the early Universe

Heavy dark matter with masses in the range $M \sim 10^{11} - 10^{16}$ GeV can be produced by several largely model-independent, nonthermal mechanisms. In particular, (i) gravitational particle production at the end of inflation and (ii) Hawking evaporation of light primordial black holes (PBHs) in the early Universe are both efficient over complementary mass intervals. Importantly, because these mechanisms depend primarily on the expansion history and on M rather than on specific dark-sector couplings, nearly degenerate components of a supermultiplet are produced in comparable amounts.¹

- (i) For a given inflationary Hubble scale H_{inf} , when $M \lesssim H_{\text{inf}}$ vacuum fluctuations and the nonadiabatic evolution across the end of inflation can generate a relic abundance controlled by the inflationary scale and the reheating history. A convenient parametric estimate is

$$\Omega_{\text{DM}} h^2 = \mathcal{C} \left(\frac{M}{10^{11} \text{ GeV}} \right)^2 \left(\frac{T_{\text{RH}}}{10^9 \text{ GeV}} \right), \quad (21)$$

where $\mathcal{C} = \mathcal{O}(1)$ encodes the mild dependence on the inflaton potential and the details of reheating (see, e.g., [46,47]). Gravitational production is particularly efficient for $M \lesssim 10^{14}$ GeV.

- (ii) Hawking evaporation provides another robust, non-thermal source of superheavy quanta. For heavier dark matter, $M \gtrsim 10^{14-16}$ GeV, evaporation of light PBHs can reproduce the observed dark-matter abundance. Using standard instantaneous-emission results around the evaporation epoch a_{ev} , one finds for the comoving number density

$$n_{\text{DM}}(a_{\text{ev}}) \simeq \left(\frac{1 \text{ g}}{M_{\text{PBH}}} \right)^4 \left(\frac{0.1 M_{\text{Pl}}}{M} \right)^2 \left(\frac{4\beta}{10^{-24}} \right) \text{ GeV}^3 \quad (22)$$

and the required initial PBH fraction $\beta \equiv \rho_{\text{PBH}}/\rho_{\text{tot}}$ (to obtain the CDM abundance today)

$$\beta \simeq 3.45 \times 10^{-20} \left(\frac{1 \text{ g}}{M_{\text{PBH}}} \right)^{1/2} \left(\frac{M}{10^{16} \text{ GeV}} \right), \quad (23)$$

valid for initial PBH masses $M_{\text{PBH}} \sim 1 \text{ g}$ to 10^8 g [48–51]. These relations illustrate that the viable β scales only with M_{PBH} and M , with no strong model

dependence in the dark sector. Recent studies have emphasized that this mechanism naturally yields the required dark-matter abundance in the above mass window [51,52].

In summary, for the exceedingly small couplings to the SM assumed in our scenario, thermal production is negligible, and nonthermal (semi)gravitational mechanisms dominate. Because their efficiencies primarily track M and the background cosmology, the nearly degenerate components of the supersymmetric multiplet are produced in similar proportions, consistent with our multicomponent framework.

IV. SIGNAL

As we notice in the previous section, in our model, there are three components for CDM at high scale M and the spectrum has small mass splittings characterized by a scale ΔM_χ . We consider two scenarios related to neutrino signals:

- (i) *Scenario I:* Heavier CDM components decay into lighter CDM, a active neutrino and SM Higgs via 3-body decay, $\chi_{\text{heavy}} \rightarrow \chi_{\text{light}} + \nu_\alpha + h$ ($\alpha = e, \mu, \tau$).
- (ii) *Scenario II:* Heavier CDM decay into lighter one and a light sterile neutrino via 2-body decay $\chi_{\text{heavy}} \rightarrow \chi_{\text{light}} + \nu_s$, and ν_s oscillates during its propagation with an active-sterile mixing θ_{as} .

Of course these two scenarios can simultaneously occur and contribute to neutrino signals in general.

In both scenarios, the processes of decays are $(\chi_{\text{heavy}}, \chi_{\text{light}}) = (\tilde{\chi}_2, \chi)$ and $(\chi, \tilde{\chi}_1)$. These two processes contribute to the neutrino signal similarly. For the case with relatively shorter lifetime compared to the age of our Universe $\tau_\chi (= \Gamma_\chi^{-1}) \lesssim \mathcal{O}(1)$ Gyr, because the dominant decaying region is very distant from the Earth ($\gtrsim \mathcal{O}(1)$ Gpc), we mainly focus on neutrinos and gamma-rays in this work.

Neutrinos. The mean-free-path of neutrinos in the extragalactic propagation is much larger than $\mathcal{O}(\text{Gpc})$ if the interaction between neutrinos cosmic target particles is dominated by electroweak interactions. The resulting neutrino flux at Earth is obtained by only redshifting the injection spectrum dN_ν/dE_ν at each redshift z .

The total neutrino spectrum is basically given by Eq. (9). The total neutrino spectrum dN_ν/dE_ν is given by

$$\frac{dN_\nu}{dE_\nu} = \left(\frac{dN_\nu}{dE_\nu} \right)_{\text{prim}} + \left(\frac{dN_\nu}{dE_\nu} \right)_{\text{sec}}. \quad (24)$$

where $(dN_\nu/dE_\nu)_{\text{prim}}$ is a monochromatic primary spectrum directly produced from hard decay processes ($\tilde{\chi} \rightarrow \chi \nu h$ or $\tilde{\chi} \rightarrow \chi \nu_s$) and $(dN_\nu/dE_\nu)_{\text{sec}}$ is a continuum tail of secondary neutrinos from electroweak/strong cascades, hadronization, and decays of soft mesons. For the evaluation of secondary spectrum, we use

¹In supersymmetric setups, superpotential couplings within a multiplet are aligned, and purely gravitational couplings are universal; in addition, our scenario assumes extremely small interactions with SM fields, so thermal production is negligible.

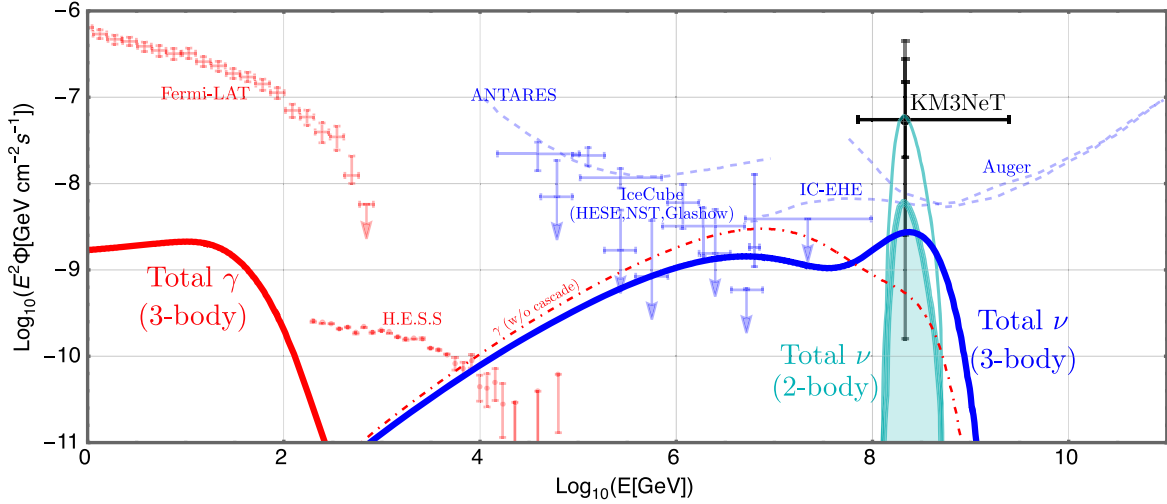


FIG. 3. The landscape of total neutrino and gamma-ray fluxes from scenario I ($\chi \rightarrow \tilde{\chi}_- + \nu_a + h$) and scenario II ($\chi \rightarrow \tilde{\chi}_- + \nu_s$ and $\nu_s \rightarrow \nu_\mu$ via oscillation). We show the total flux of neutrinos and gamma-rays in scenario I as thick blue and thick red lines respectively. The total flux of neutrinos in scenario II is shown as thick green line. In order to provide demonstrative examples, we choose the parameters as $M_\chi = 10^{16}$ GeV, $\Delta M_\chi = 3 \times 10^9$ GeV, $f_\chi = 0.6$, $\text{Br}_\nu = 10^{-2.6}$ for scenario I, and $M_\chi = 6 \times 10^{12}$ GeV, $\Delta M_\chi = 5 \times 10^8$ GeV, $f_\chi = 0.6$, $\text{Br}_\nu = 10^{-4}$ for scenario II. The lifetime of heavier component χ is always assumed to be $\tau_\chi \equiv \Gamma_\chi^{-1} = 1$ Gyr. For the details, see the main text. The upper limits on isotropic diffuse fluxes of gamma-rays (from H.E.S.S. [63] and Fermi-LAT [64]) and neutrinos (from ANTARES [37], IC-HESE [32], IC-NST [33], IC-Glashow [34], IC-EHE [35] and Auger [36]) are shown as faint red and blue lines, respectively. The desired neutrino flux based on KM3NeT effective area is indicated by black and gray lines [1]. In scenario I, gamma-ray flux without attenuation and electromagnetic cascade is shown as a dot-dashed red line for comparison. In scenario II, we also show the case with $\text{Br}_\nu = 10^{-3}$ as upper thin green line.

HDMSpectra [53].² We have both primary and secondary spectrum in scenario I ($\chi_h \rightarrow \chi_l + \nu + h$), and only primary spectrum in scenario II ($\chi_h \rightarrow \chi_l + \nu_s$).

Gamma-rays. Compared to neutrinos, the cosmological background is much less transparent to photons. For $\mathcal{O}(\text{PeV}) \lesssim E_\gamma \lesssim \mathcal{O}(\text{EeV})$, the mean-free-path is even less than $\mathcal{O}(1)$ Mpc [54,55]. Attenuation and electromagnetic cascade processes significantly affect its propagation [56]. The resulting gamma-ray flux at Earth is given by

$$\frac{d\Phi_\gamma^{\text{EG}}}{dE_\gamma} = \frac{\Gamma_\chi(\gamma) f_\chi \rho_{\text{DM}c}}{4\pi M_\chi H_0} \times \int_0^{z_{\text{max}}} \frac{dz}{1+z} \frac{e^{-\Gamma_\chi t(z)}}{\sqrt{\Omega_m(1+z)^3 + \Omega_\Lambda}} P\left(z, \frac{dN_\gamma^{\text{inj}}}{dE_\gamma}\right) \quad (25)$$

where $\Gamma_\chi(\gamma)$ is the decay rates for relevant decay process contributing to photon production and

²In order to estimate the signal flux in scenario I, we note that the mean energy of light particles $\langle E_{\nu,h} \rangle \approx \frac{\Delta M}{2}$ for the case with $|\mathcal{M}|^2 = \text{constant}$, and $\langle E_{\nu,h} \rangle \approx \frac{2\Delta M}{5}$ for the case with the effective operator (16) in the limit $\Delta M \ll M$. In Fig. 3, as an approximation we take $E_{\nu,h} = \frac{\Delta M}{2}$ for the energy of neutrino and Higgs boson at the rest frame of the heavy DM particle, to generate the flux of secondary photons and neutrinos from the decay of Higgs.

$dN_\gamma^{\text{obs}}/dE_\gamma \equiv P(z_0, dN_\gamma^{\text{inj}}/dE_\gamma)$ is the operation including the attenuation of high-energy gamma-rays through pair-production ($\gamma + \gamma_{\text{bkd}} \rightarrow e^+ + e^-$) and reproduction of down-scattered photons by inverse Compton scattering ($e^\pm + \gamma_{\text{bkd}} \rightarrow e^\pm + \gamma$) in the propagation from the injection position at the redshift $z = z_0$ to the Earth ($z = 0$). The cascade evolution can be checked by solving transport equation in the presence of CMB for microwave and extragalactic background light (EBL) for the infrared and optical components for each redshift value z . We adopt the EBL density at various redshift z from Ref. [57] for our calculation. In this work, the evaluation has been done by utilizing GammaCascade [58,59] which solve the transport equation semianalytically. Similar to neutrino case, the injection spectrum $dN_\gamma^{\text{inj}}/dE_\gamma$ was obtained by using HDMSpectra.

Signatures and observational constraints. Assuming $\tau_\chi = 1$ Gyr and $f_\chi = 0.6$, we consider one representative set of parameters for each scenarios as the following:

(i) Scenario I ($\chi_h \rightarrow \chi_l + \nu + h$)

$$\begin{aligned} M_\chi &= 10^{16} \text{ GeV}, \\ \Delta M_\chi &= 3 \times 10^9 \text{ GeV}, \\ \Lambda &= 1.44 \times 10^{31} \text{ GeV} \end{aligned} \quad (26)$$

(ii) Scenario II ($\chi_h \rightarrow \chi_l + \nu_s$)

$$\begin{aligned} M_\chi &= 6 \times 10^{12} \text{ GeV}, \\ \Delta M_\chi &= 5 \times 10^8 \text{ GeV}, \\ \lambda_s &= 1.12 \times 10^{-22} \end{aligned} \quad (27)$$

which correspond to $\text{Br}_\nu \simeq 10^{-2.6}$ for scenario I, and $\text{Br}_\nu \simeq 10^{-4} - 10^{-3}$ for scenario II.

For scenario I, although the scale Λ is much greater than the Planck scale, it can be induced by sub-Planckian physics with technically natural small couplings. An example could include a heavy right-handed neutrino and its superpartner with mass around m_N , having Yukawa interactions with a Dirac neutrino Yukawa y_ν and an additional Yukawa coupling y_N between heavy neutrino multiplet N and chiral multiplet X . The superpotential is written as

$$W = y_\nu L H_u N + y_N N X^2. \quad (28)$$

In this case the scale $\Lambda = m_N/(y_\nu y_N)$ appears from integrating out heavy field N . For instance, $y_\nu, y_N \sim 10^{-8}$ and $m_N \sim 10^{16} > M_\chi$ GeV can be a viable parameter choice in this scenario.

For this set of example parameters, our results are summarized in Fig. 3. In scenario II, we set our parameters $\text{Br}_\nu \lesssim 10^{-3}$ to satisfy the constraints on sterile-active mixing at $\mathcal{O}(\text{eV})$ mass [60], and $\Delta M_\chi/M_\chi \sim 10^{-4}$ is taken to be consistent with the Lyman-alpha forest constraints on the kick velocity for the daughter dark matter particle [61,62].

We show the expected signal of neutrinos and gamma-rays and compare them with existing isotropic diffuse neutrino/gamma-ray flux limits. One finds the relevant constraints on each energy region of the spectrum in the caption of the figure.

V. CONCLUSION AND DISCUSSION

In this paper, we have explored the superheavy supersymmetric dark matter (SSDM) scenario and its neutrino and gamma-ray signatures over a wide energy range. In the case of superheavy dark matter, supersymmetry naturally leads to a nearly degenerate mass spectrum among supermultiplet fields. The energy of neutrinos injected from the decay of heavier components carries direct information about the scale of SUSY breaking.

If neutrinos are produced via the direct decay of a heavier component, the process inevitably involves the production of Higgs bosons, leading to an additional gamma-ray signal. In contrast, if neutrinos originate from the oscillation of sterile neutrinos produced in the decay of a heavy

component, the resulting neutrino spectrum is expected to be sharper, with no accompanying gamma-ray signal.

Motivated by the recent detection of an ultrahigh-energy (UHE) neutrino event at $E_\nu \sim 220$ PeV by the KM3NeT Collaboration, we have analyzed this scenario in the context of SSDM. Our findings indicate that if the lifetime of the heavier component of dark matter is around $\mathcal{O}(1)$ Gyr, an extragalactic, nearly isotropic neutrino flux is robustly expected at energies around $E_\nu = \mathcal{O}(0.1 - 1)\Delta M_\chi$. Moreover, it is straightforward to consider scenarios in which the lifetime of heavier components is shorter, which could lead to distinctive observational consequences.

This scenario could have broader implications for the spectra of high-energy cosmic particles across various energy ranges, including ultrahigh-energy cosmic rays (UHECRs) at or above EeV scales and very-high-energy gamma rays in the TeV-PeV range. With upcoming neutrino and ultrahigh-energy cosmic ray observatories, such as IceCube-Gen2 [65] and AuperPrime [66], we anticipate that the nature of superheavy dark matter and the origins of ultrahigh-energy cosmic particles will be further unveiled in the near future [67–69].

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DATA AVAILABILITY

The data are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX: R-PARITY VIOLATION, BBN SAFETY, AND PROTON STABILITY

In our supersymmetric superheavy dark-matter scenario, R-parity violation (RPV) is introduced to ensure that the MSSM LSP decays before BBN, while not endangering proton stability:

- (i) The LSP decays promptly via RPV.
- (ii) We allow either B or L violation, but not both, thereby removing the dangerous dimension-4 p -decay operators. A heavy gravitino kinematically forbids exotic channels such as $p \rightarrow K^+ + \psi_{3/2}$.

In this appendix, we qualitatively show how these conditions can be satisfied.

1. BBN safety

In high-scale SUSY the LSP is typically heavy; even if a relatively light neutralino arises from a small μ -term, RPV-induced decays are fast enough provided the couplings are not extremely tiny. For trilinear RPV with an off-shell sfermion mediator, the dominant three-body width follows the standard estimate [70,71]

$$\Gamma_{\text{3body}}(\text{LSP} \rightarrow fff) \sim \frac{|\lambda_{\text{RPV}}|^2 m_{\text{LSP}}^5}{192\pi^3 \tilde{m}^4}, \quad (\text{A1})$$

where $\lambda_{\text{RPV}} \in \{\lambda, \lambda', \lambda''\}$ denote the RPV couplings associated with the superpotential terms LLE , LQD , and UDD , respectively; m_{LSP} is the LSP mass, and $\tilde{m}$ is a typical sfermion mass. Numerically,

$$\tau_{\text{LSP}} \approx 3.9 \times 10^{-14} \text{ sec} \times \left(\frac{10^{-7}}{\lambda_{\text{RPV}}}\right)^2 \left(\frac{\tilde{m}}{10^8 \text{ GeV}}\right)^4 \left(\frac{10^5 \text{ GeV}}{m_{\text{LSP}}}\right)^5. \quad (\text{A2})$$

Thus, for a heavy LSP (e.g., $m_{\text{LSP}} \gtrsim 10^5 \text{ GeV}$), the lifetime is far below 1 sec even for very small λ_{RPV} , ensuring BBN safety. If a lighter neutralino appears (e.g., $m_{\text{LSP}} \sim \text{TeV}$), Eq. (A1) implies that a moderately small λ_{RPV} still suffices for $\tau_{\text{LSP}} < 1 \text{ sec}$ when $\tilde{m}$ is large. Hence, across the relevant mass ranges, the MSSM LSP is not cosmologically long-lived in our setup.

2. Proton stability

Dimension-4 proton decay requires the simultaneous presence of B - and L -violating couplings [70,72]. We therefore adopt one of the two safe choices.

- (i) B -conserving, L -violating RPV: $\lambda, \lambda' \neq 0$, $\lambda'' = 0$.
 - (ii) L -conserving, B -violating RPV: $\lambda'' \neq 0$, $\lambda, \lambda' = 0$
- Either choice eliminates the dangerous tree-level operators for $p \rightarrow \text{meson} + \ell$. In the first case (B conserved) standard proton-decay modes are absent; in the second case (L conserved) they are also absent, though one must consider dinucleon decay and $n\bar{n}$ transitions. These are parametrically suppressed here by the large sfermion masses and remain compatible with bounds for the small λ'' needed only to ensure pre-BBN LSP decay (see, e.g., [73]); discrete symmetries such as baryon triality $\mathbb{Z}_3^B$ or lepton parity can enforce these patterns and forbid mixed B - and L -violating operators [74]. In our high-scale scenario the gravitino is very heavy ($m_{3/2} \gg m_p$), so proton-decay channels emitting a gravitino are kinematically forbidden. Thus neither RPV nor supergravity-induced operators open viable p -decay modes in our setup.

Consequently, our SUSY implementation is compatible with current proton-decay limits while meeting the cosmological requirement that all MSSM relics decay well before BBN. The only stable relic is the dark-sector state protected by its own Z_2 , which constitutes today's dark matter in our analysis.

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