



# Signals of very high energy neutralinos in future cosmic ray detectors

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## Abstract

“Top–down” models explain the observation of ultra-high energy cosmic rays (UHECR;  $E \gtrsim 5 \times 10^{19}$  eV) through the decay of very massive, long-lived “ $X$  particles”. If superparticles with masses near a TeV exist,  $X$  decays also lead to a significant flux of very energetic neutralinos, assumed to be the (stable or long-lived) lightest superparticles. There is a range of energies where neutrinos get absorbed in the Earth, but neutralinos can still traverse it. These neutralinos could in principle be detected. We calculate the detection rate in planned experiments such as OWL and EUSO. For bino-like neutralinos, which have been considered previously, we find detection rates below 1 event per Teraton of target and year in all cases; often the rates are much smaller. In contrast, if the neutralino is higgsino-like, more than ten events per year per Teraton might be observed, if the mass of the  $X$  particle is near its lower bound of  $\sim 10^{12}$  GeV.

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## 1. Introduction

The existence of ultra-high energy cosmic rays (UHECR), with  $E \gtrsim 5 \times 10^{19}$  eV, remains a mystery [1]. They have been detected by every major cosmic ray experiment, but we do not know of any astronomical objects that has sufficiently strong electromagnetic fields extending over a sufficiently large volume to accelerate charged particles to the required energies. Nor do we understand how these particles, once created, can reach us, given their energy loss through scattering on the cosmic microwave background [2].

One radical idea [3] is that UHECR originate from the decay of very massive, yet long-lived  $X$  particles. Since one starts with very energetic particles, which lose energy first through parton showering and fragmentation, and later while propagating through the universe, these class of models are known as “top–down” models. The most energetic CR event that has been observed to date has  $E \simeq 3 \times 10^{20}$  eV [4]. This implies a lower bound  $M_X \gtrsim 10^{12}$  GeV on the mass of the  $X$  particles. Since UHECR are observed today, the lifetime of  $X$  must be at least comparable to the age of the Universe. Several particle physics models containing candidates with sufficiently large mass and long lifetime have been suggested [1,5]. Ways to produce these particles in the very early universe are discussed in [3,6].

Models of this type can be made compatible with all existing data, including the first data from the Pierre Auger observatory [7]. However, in order to decisively test these models, one has to find predictions that allow to discriminate between top–down and the more conventional bottom–up [1] models. These two classes of models usually predict somewhat different spectra for photons and neutrinos at high energies, and/or different distributions of the arrival directions. However, distinguishing between UHE photons and protons is nontrivial. Gigaton class neutrino telescopes now under construction should see some very energetic neutrinos if these models are correct [8]; however, bottom–up models generically also lead to comparable neutrino fluxes. Anisotropies in the arrival direction can be expected [9], if  $X$  particles are distributed like (or even form the) dark matter in our galaxy; however, quantitative details depend on the distribution of matter near the galactic center, which is not well understood.

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These difficulties motivate the analysis of signals where bottom–up and top–down models make *qualitatively* different predictions. This may be possible if we postulate the existence of superparticles [10] at or near the electroweak energy scale. This assumption is quite natural in the given context, since supersymmetry is the only known way to stabilize the large hierarchy between  $M_X$  and the electroweak scale against radiative corrections.<sup>1</sup> Since  $M_X$  is much larger than the sparticle mass scale,  $X$  decays will produce large number of superparticles. This is true even if the primary decay of  $X$  only involves Standard Model (SM) particles; in this case superparticles will be produced in the subsequent parton shower [11,12]. If  $R$ -parity is (nearly) conserved, all these superparticles will decay into stable lightest superparticles (LSPs), assumed to be the lightest neutralino. In contrast, bottom–up models will produce a miniscule flux of superparticles. The reason is that the vast majority of UHE proton or photon interactions with matter only produces additional light particles (in particular, light mesons and baryons); the cross section for producing superparticles remains very small even at these energies.

This raises the question how one might observe these very energetic neutralinos. The crucial observation [13] is that there is a range of energies where neutrinos get absorbed in the Earth, whereas neutralinos can traverse it with little or no loss of energy. The reason for this difference is the smaller neutralino–nucleon scattering cross section, and/or the smaller neutralino energy loss per interaction [14]. Note that neutralino interactions always lead to a superparticle in the final state, which will decay back into a neutralino. An interaction will therefore not change the total neutralino flux, but will shift it to lower energies, where it is (even) more difficult to detect.

In this Letter we provide a detailed calculation of the neutralino event rates that one might expect in future cosmic ray detectors with very large target volumes, like OWL [15] and EUSO [16]. We improve on existing analyses [13,17,18] in several ways. We use neutralino spectra impinging on Earth calculated with the most complete code for  $X$  particle decays [12], where we analyze several different primary decay modes. We also carefully include the effects of neutralino propagation through the Earth, using the results of [14]. Our calculation of the event rates includes a cut on the visible energy deposited by a neutralino interaction; since this interaction again produces an invisible neutralino, the visible energy is usually significantly smaller than the energy of the incoming neutralino. Moreover, we investigate both bino- and higgsino-like neutralinos; the cross sections for the latter have also been computed in [14]. We find that higgsino-like neutralinos would in fact be much easier to detect; bino-like neutralinos most likely remain out of reach even for the planned EUSO and OWL missions. Finally, we calculate the neutrino background from the same model of  $X$  decays as the signal.

The remainder of this Letter is organized as follows. The calculation of the event rates is described in Section 2. In Section 3 we present numerical results, and Section 4 is devoted to a brief summary and some conclusions.

## 2. Calculation of event rates

Neutralinos are produced along with protons, photons, electrons and neutrinos at the location of  $X$  decays, following a prolonged parton shower [11,12]. We fix the normalization through the proton flux at  $10^{20}$  eV, which we take to be

$$E^3 F_p(E) = 1.6 \times 10^{24} \text{ eV}^2 \text{ m}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \quad (1)$$

at  $E = 10^{20}$  eV. This roughly corresponds to the flux observed by the HiReS experiment [20], which is somewhat smaller than that observed by AGASA [19]. Note, however, that we ignore the contribution of photons to the UHECR flux. This is phenomenologically motivated by the observation that UHECR events seem to be proton-like, rather than photon-like [21]. Normalizing to the sum of the proton and photon fluxes would obviously reduce the predicted neutralino flux, and hence the event rate; depending on the  $X$  decay model, the reduction factor would roughly lie between two and five. On the other hand, we ignore all propagation effects. If most  $X$  decays occur at significant distance from our galaxy, which may well be true if  $X$  particles are confined to topological defects, both the proton and photon fluxes might be depleted by propagation, while leaving the neutralino (and neutrino) flux essentially unchanged. The presence of significant propagation effects would therefore increase the predicted neutralino flux on Earth.

Neutralinos can interact with nucleons either through the exchange of a squark in the  $s$ -channel, or through the exchange of a  $Z^0$  or  $W^\pm$  gauge boson in the  $t$ -channel. In the following we treat these two contributions, which essentially do not interfere [14], in turn, before discussing the calculation of the neutrino-induced background.

As explained in [14,17,18], the  $s$ -channel contribution is dominated by the exchange of on-shell squarks. The event rate is given by:

$$\mathcal{N}_s = \sum_q \int_{E_{\min}}^{E_{\max}} dE_{\text{vis}} \int_{X_{\min}}^{X_{\max}} dX \int_0^{y_{\max,q}} dy \frac{1}{y} F_{\tilde{\chi}_1^0} \left( \frac{E_{\text{vis}}}{y}, X \right) \frac{d\sigma_s(E_{\text{vis}}/y, y)}{dy} \mathcal{V}. \quad (2)$$

<sup>1</sup> Note that “large” extra dimensions do not help here, since by construction the “fundamental scale” must be at least  $M_X$  in order to explain the observed UHECR; this is independent of the dimensionality of spacetime.

Here,  $F_{\tilde{\chi}_1^0}$  is the differential neutralino flux, which depends on the neutralino energy as well as the matter depth<sup>2</sup>  $X$ . The sum runs about all quark flavors  $q$ , and the first integration is over the visible energy  $E_{\text{vis}} = E_{\tilde{\chi}_{1,\text{in}}^0} - E_{\tilde{\chi}_{1,\text{out}}^0} = yE_{\tilde{\chi}_{1,\text{in}}^0}$ . The factor  $1/y$  appears because we integrate over the visible, rather than total, energy. The lower limit  $E_{\text{min}}$  on  $E_{\text{vis}}$  is determined by the energy sensitivity of the experiment, whereas the upper limit  $E_{\text{max}}$  is determined by kinematics,  $E_{\text{max}} \sim M_X/2$ ; however, after propagation through the Earth the neutralino flux at the highest kinematically allowed energy is very small. The lower bound on the column depth,  $X_{\text{min}} = 0.13 \times 10^6 \text{ GeV}^3$ , corresponds to an angular cut of about 5% on the signal, i.e. we only count events that emerge at least five degrees below the horizon; this cut greatly reduces the neutrino background.  $X_{\text{max}} = 2.398 \times 10^6 \text{ GeV}^3$  is the maximal earth column depth, corresponding to neutralinos that emerge vertically out of the Earth. The kinematic maximum of the scaling variable  $y$ , for 2-body decays  $\tilde{q} \rightarrow q + \tilde{\chi}_1^0$ , is  $y_{\text{max},q} = 1 - m_{\tilde{\chi}_1^0}^2/m_{\tilde{q}}^2$ . Since the maximal neutralino energy is finite, there should strictly speaking also be a non-vanishing lower bound on  $y$ ; note that we need the neutralino flux at  $E_{\tilde{\chi}_1^0} = E_{\text{vis}}/y$ . An explicit expression for the differential cross section  $d\sigma_s/dy$  can be found in [14]. Finally, the constant factor  $\mathcal{V}$  is given by

$$\mathcal{V} \equiv 2\pi V_{\text{eff}} \epsilon_{\text{DC}} t N_A \rho_w J_D. \quad (3)$$

Here,  $V_{\text{eff}}$  is the water equivalent (w.e.) effective volume,  $\epsilon_{\text{DC}}$  is the duty cycle (the fraction of time where the experiment can observe events),  $t$  is the observation time,  $N_A = 6.022 \times 10^{23} \text{ g}^{-1}$  is Avogadro's number,  $\rho_w = 10^6 \text{ g m}^{-3}$  is the density of water, and  $J_D = |d \cos \theta / dX|$  is the Jacobian for the transformation  $\cos \theta \rightarrow X(\cos \theta)$ .

The  $t$ -channel exchange diagrams predominantly lead to the production of heavier neutralinos or charginos in the final state [14], which we collectively denote by  $\tilde{\chi}_{\text{out}}$ . The visible energy therefore also depends on the  $\tilde{\chi}_{\text{out}}$  decay kinematics. The event rate can be written as:

$$\mathcal{N}_t = \int_{E_{\text{min}}}^{E_{\text{max}}} dE_{\text{vis}} \int_{X_{\text{min}}}^{X_{\text{max}}} dX \int_0^1 dy \frac{1}{y} F_{\tilde{\chi}_1^0} \left( \frac{E_{\text{vis}}}{y}, X \right) (G_{\tilde{\chi}_1^0}^{\text{NC}}(E_{\text{vis}}, y) + G_{\tilde{\chi}_1^0}^{\text{CC}}(E_{\text{vis}}, y)) \mathcal{V}. \quad (4)$$

Here we have written the contributions from charged and neutral currents separately. Each term is given by a convolution of a differential cross section for the production of  $\tilde{\chi}_{\text{out}}$  with the  $\tilde{\chi}_{\text{out}}$  decay spectrum. These convolutions are more easily written in terms of the variable  $z = E_{\tilde{\chi}_{1,\text{out}}^0} / E_{\tilde{\chi}_{1,\text{in}}^0} = 1 - y$ :

$$G_{\tilde{\chi}_1^0}^{\text{NC,CC}}(E_{\text{vis}}, y) = \int_z^{z_{1,\text{max}}} \frac{dz_1}{z_1} \frac{d\sigma_{t\tilde{\chi}}^{\text{NC,CC}}(E_{\text{vis}}/y, z_1)}{dz_1} \frac{1}{\Gamma} \frac{d\Gamma_{\tilde{\chi}_{\text{out}}}(z_1 E_{\text{vis}}/y, z_2 = (z/z_1))}{dz_2} \theta(z - z_{\text{min}}) \theta(z_{\text{max}} - z) \Big|_{z=1-y}. \quad (5)$$

Here  $z_1 = E_{\tilde{\chi}_{\text{out}}} / E_{\tilde{\chi}_{1,\text{in}}}$  describes the energy transfer from the incoming lightest neutralino to the heavier neutralino or chargino, and  $z_2 = E_{\tilde{\chi}_{1,\text{out}}} / E_{\tilde{\chi}_{\text{out}}}$  describes the energy transfer from this heavier neutralino or chargino to the lightest neutralino produced in its decay.  $z_2$  is chosen such that  $z \equiv z_1 z_2 = 1 - y$ . Explicit expressions for the differential cross sections, and for the limits  $z_{\text{min,max}}$ ,  $z_{1,\text{max}}$  in Eq. (5), can again be found in [14].<sup>3</sup> In principle one would need to include sums over  $\tilde{\chi}_{\text{out}}$  in Eq. (5). In practice, however, a single neutralino and a single chargino dominate neutral and charged current reactions, respectively [14].

The event rates (2) and (4) depend on the neutralino flux *after* propagation through the Earth. Of course, the propagation effects also depend on whether  $s$ - or  $t$ -channel exchange is dominant. We treat these effects through straightforward numerical integration of the transport equations, as described in [14].

The *background* is dominated by  $\nu_\tau$  scattering through  $t$ -channel exchange of  $W$  or  $Z$  bosons. At the relevant energies electron and muon neutrinos get absorbed efficiently in the Earth. However, since  $\nu_\tau$  interactions regenerate another  $\nu_\tau$ , albeit at lower energy,  $\tau$  neutrinos can always traverse the Earth, although their energy may be reduced drastically. Again treating charged and neutral current processes separately, the background rate can be written as

$$\mathcal{N}_\nu = \int_{E_{\text{min}}}^{E_{\text{max}}} dE_{\text{vis}} \int_{X_{\text{min}}}^{X_{\text{max}}} dX \int_0^1 dy \frac{1}{y} F_\nu \left( \frac{E_{\text{vis}}}{y}, X \right) \left( \frac{d\sigma_{t\nu}^{\text{NC}}(E_{\text{vis}}/y, y)}{dy} + N_\nu^{\text{CC}}(E_{\text{vis}}, y) \right) \mathcal{V}, \quad (6)$$

where  $y = 1 - E_{\nu,\text{in}}/E_{\nu,\text{out}}$ . In the case of NC scattering ( $Z$ -exchange) the entire visible energy results from the hadronic vertex. In case of CC scattering ( $W$ -exchange) we add the visible energy released in  $\tau$  decay to that produced at the hadronic vertex:

$$N_\nu^{\text{CC}}(E_{\text{vis}}, y) = \int_z^{z_{1,\text{max}}} \frac{dz_1}{z_1} \frac{d\sigma_\nu^{\text{CC}}(E_{\text{vis}}/y, z_1)}{dz_1} \frac{1}{\Gamma} \frac{d\Gamma(z_1(E_{\text{vis}}/y), z_2 = (z/z_1))}{dz_2} \theta(z - z_{\text{min}}) \theta(z_{\text{max}} - z) \Big|_{z=1-y}. \quad (7)$$

<sup>2</sup> Matter depth  $X$  is customarily given as a column depth, measured in  $\text{g cm}^{-2}$  or, in natural units, in  $\text{GeV}^3$ ; for the Earth,  $X \in [0, 2.398 \times 10^6 \text{ GeV}^3]$  [22].

<sup>3</sup> Note that the  $G_{\tilde{\chi}_1^0}^{\text{NC,CC}}$  of Eq. (5) are the integration kernels  $K_{\tilde{\chi}_1^0}^{\text{NC,CC}}$  of Ref. [14], multiplied with the total cross section for  $t$ -channel scattering.

This expression is formally very similar to Eq. (5), which also includes contributions to the visible energy from the decay of an unstable particle. This treatment is conservative since it ignores the fact that a  $\tau$  produced inside the target volume may decay outside of it. Moreover, if  $\tau$  production and decay both occur inside the target volume, it may be possible to use this “double bang” signature to remove these background events. On the other hand, we ignore the background from  $\tau$ s produced outside the target which decay inside the observed volume. This contribution should be smaller, since one would need higher neutrino energy to produce a given visible energy in this manner. Note that at the energies in question,  $\tau$  energy losses in rock or water are no longer negligible; this reduces the energy released in  $\tau$  decay even further. Recall that after propagation through the earth the  $\nu_\tau$  flux is a steeply falling function of energy.

The background rate (6) is proportional to the tau neutrino flux  $F_\nu$  emerging from the Earth. The  $\nu_\tau$  flux at the location of  $X$  decay is usually quite small [12]. However, due to near-maximal neutrino flavor mixing, the three neutrino fluxes impinging on Earth are very nearly equal, i.e. we take one third of the total neutrino flux, normalized according to Eq. (1), as estimate of the incoming  $\nu_\tau$  flux.

As mentioned above, tau neutrinos may lose much of their energy while traversing the Earth. We solve the corresponding transport equations using the methods of Ref. [14]. Since we are interested in very high energies, the tau leptons produced in CC  $\nu_\tau$  reactions may lose a significant fraction of their energy before decaying. We therefore modified the standard treatment [22] in order to at least crudely estimate the effects of  $\tau$  energy loss in matter. We do this by formally treating this energy loss as additional scattering. To this end, we modify the integration kernel in the transport equation for  $\nu_\tau$  as follows:

$$\frac{1}{\sigma(E_y)} \frac{d\sigma(E_y, z)}{dz} \rightarrow \int \frac{1}{\sigma(E_y)} \frac{d\sigma(E_y, z_1)}{dz_1} \frac{1}{L} \frac{dL(z_1 E_y, E'')}{dE''} dE'' \Big|_{z=E''/E}. \quad (8)$$

Here  $E_y = E/(1-y)$  is the energy of the incident neutrino that gives rise to a neutrino with energy  $E$  after the scattering, and the function  $dL(E_{\tau,\text{in}}, E_{\tau,\text{out}})/dE_{\tau,\text{out}}$  describes the  $\tau$  energy loss. We make the very simple ansatz [23]

$$\frac{dE_\tau}{dz} = -\beta\rho E_\tau \quad \text{with } \beta = 0.85 \times 10^{-6} \text{ cm}^2 \text{ g}^{-1} = \text{const.} \quad (9)$$

This implies  $E_\tau(z) = E_\tau(0)e^{-\beta\rho z}$ . We assume that all  $\tau$ s decay after traveling a distance  $z_{\text{dec}} = E_\tau c\tau_\tau/m_\tau$ , where  $\tau_\tau$  is the lifetime of the  $\tau$  lepton and  $c$  is the speed of light. Note that we estimate the average decay length from the  $\tau$  energy *after* propagation. This underestimates the decay length, and hence the effect of  $\tau$  energy loss. On the other hand, for  $E_{\nu_\tau} < 10^{10}$  GeV the ansatz (9) overestimates the energy loss [23]. Our approximation of a fixed decay length leads to

$$\frac{dL(E', E'')}{dE''} = \delta(E'' - E' \exp(-\kappa E'')), \quad (10)$$

with constant  $\kappa = \beta\rho c\tau_\tau/m_\tau$ . The integral over  $dL/dE''$ , which appears in Eq. (8), is then given by:

$$L = \int dE'' \delta(E'' - E' \exp(-\kappa E'')) = \frac{1}{1 + \kappa E' \exp(-\kappa E'')}, \quad (11)$$

where in the last expression  $E''$  has to be interpreted as a function of  $E'$ , as determined by the argument of the  $\delta$ -function. We can then evaluate the integral in Eq. (8):

$$\frac{1}{\sigma(E_y)} \frac{d\sigma(E_y, z)}{dz} \rightarrow (1 + \kappa z_1 E_y) \exp(\kappa z_1 E_y) \frac{1}{\sigma(E_y)} \frac{d\sigma(E_y, z')}{dz'} \Big|_{z'=z_1 \exp(\kappa z_1 E_y)}. \quad (12)$$

The obvious advantage of our simplified treatment is that it does not necessitate the numerical evaluation of additional integrals. This would have been very costly, since the length scales involved in  $\tau$  energy loss and decay (a few km for  $E_\tau \sim 10^8$  GeV) are very much shorter than the  $\nu_\tau$  interaction length in rock ( $\sim 10^3$  km for  $E_{\nu_\tau} = 10^8$  GeV) [23]. A more accurate treatment would therefore have required to use many more steps in  $X$  when integrating the transport equation; even with our simple treatment, or indeed without including the effects of  $\tau$  energy loss, calculating the  $\nu_\tau$  flux emerging from Earth takes up to several CPU days. On the other hand, our simplified treatment can only give us an indication of the size of effects due to  $\tau$  energy losses. We find that the effect on the  $\nu_\tau$  flux emerging from Earth is essentially negligible for  $E_{\nu_\tau} \lesssim 10^7$  GeV. This is also true for  $X \gtrsim 0.3X_{\text{max}}$ , since then the flux at  $E_{\nu_\tau} > 10^7$  GeV is negligible even if the  $\tau$  energy loss is ignored. However, it can reduce the  $\nu_\tau$  flux by a factor of two or more at large  $E_{\nu_\tau}$  and small  $X$ .

### 3. Results

We are now ready to present numerical results. Earlier estimates [13,17] have shown that one will need at least teraton scale targets in order to detect hadronic interactions of neutralinos in top-down models. Currently the only technology that might allow to monitor such large targets is optical observation from space [15,16]. Here one detects the light, either from Čerenkov radiation

or from fluorescence, emitted by very energetic showers in the atmosphere. The target is therefore quite thin: the neutralinos would have to interact either in the atmosphere itself, or just below it. One usually estimates an effective target thickness of 10 to 20 m w.e. A teraton target then results if one can monitor  $\mathcal{O}(10^6)$  km<sup>2</sup> simultaneously, which might be possible [15,16]. One drawback of this approach is that observations of this kind are only feasible on clear, moonless nights, leading to a duty cycle  $\epsilon_{\text{DC}}$  in Eq. (3) of only about 10%. In our numerical results we therefore take a target mass of 1 Tt,  $\epsilon_{\text{DC}} = 0.1$ , and assume an observation time of one year.

As shown in [13], the expected neutralino flux depends quite strongly on  $M_X$  as well as on the dominant  $X$  decay mode. Top-down models predict rather hard spectra, i.e.  $E^3$  times the flux increases with energy. Fixing the (proton) flux at  $E = 10^{20}$  eV therefore leads to smaller fluxes at  $E < 10^{20}$  eV as  $M_X$  is increased. Moreover, if  $M_X$  is not far from its lower bound of  $\sim 10^{12}$  GeV, much of the relevant neutralino flux is produced early in the parton cascade triggered by  $X$  decay, which is quite sensitive to the primary  $X$  decay mode. In contrast, if  $M_X \gg 10^{12}$  GeV, in the relevant energy range most LSPs originate quite late in the cascade; in that case the LSP spectrum is largely determined by the dynamics of the cascade itself, which only depends on Standard Model interactions, and is not very sensitive to the primary  $X$  decay mode(s).

Following Ref. [13] we therefore study scenarios with  $M_X = 10^{12}$  and  $10^{16}$  GeV, for four different primary  $X$  decay modes. In contrast to previous analyses [13,17,18] we calculate the event rates for both bino-like and higgsino-like neutralinos. As explained in Ref. [14] the former interact with hadronic matter almost exclusively through  $s$ -channel scattering, while the latter dominantly interact through  $t$ -channel diagrams.

Finally, we present results for two different values of the minimal visible energy  $E_{\text{min}}$ . Events with visible energy as “low” as  $10^6$  GeV might be observable via the Čerenkov light emitted by particles in the atmosphere with velocities exceeding the speed of light in air. On the other hand, the fluorescence signal (observed, e.g., by the HiReS experiment [20]) can probably only be seen for energies  $\gtrsim 10^9$  GeV. In all cases we require the event to come from an angle at least five degrees below the horizon. This greatly reduces the neutrino-induced background, as explained earlier.

We present results for higgsino- and bino-like neutralinos in Tables 1 and 2, respectively. We saw in Ref. [14] that the cross section for neutralino–nucleon scattering depends only weakly on details of the sparticle spectrum if  $\tilde{\chi}_1^0$  is higgsino-like. In Table 1 we therefore only show results for one scenario with higgsino-like LSP. It has an LSP mass of 300 GeV, with the second neutralino and first chargino, which are produced predominantly in NC and CC scattering respectively, having masses of 310 and 303 GeV, respectively; the near-degeneracy of these three states is a consequence of these states all being higgsino-like, which in turn follows if the LSP is a rather pure higgsino state.

As expected, we find much higher event rates for  $M_X = 10^{12}$  GeV than for  $M_X = 10^{16}$  GeV. In the former case we also see that the predicted event rate depends significantly on the primary  $X$  decay mode, again as expected. The decay into a lepton plus a slepton turns out to be most favorable. The reason is that this decay mode leads to a rather small number of protons produced per  $X$  decay, or, put differently, to a large ratio of the LSP and proton fluxes [12]. Since we normalize to the proton flux, this then leads to a rather large LSP flux. This decay mode also leads to the hardest  $\tilde{\chi}_1^0$  spectrum. Since the primary  $X$  decay only involves weakly interacting (s)particles, parton showering carries away a relatively small fraction of the energy of the original particles. The original slepton will then eventually decay into a very energetic neutralino. As a result, increasing the cut on  $E_{\text{vis}}$  by three orders of magnitude only reduces the predicted event rate by a factor of  $\sim 5$  in this case.

The second most favorable primary  $X$  decay mode is the one into five quarks and five squarks. Since we produce ten strongly interacting (s)particles already in the very first step, each of which initiates an extended QCD shower, the final multiplicity is very

Table 1

Predicted events rates per teraton and year (with duty cycle  $\epsilon_{\text{DC}} = 0.1$ ) for the scenario H2 of [14], where  $\tilde{\chi}_1^0$  is higgsino-like, and for the  $\nu_\tau$  induced background. Both signal and background depend on the mass  $M_X$  of the progenitor particle, as well as on the primary  $X$  decay mode. We show results for  $X$  decays into a first generation quark antiquark pair (“ $q\bar{q}$ ”), into a first generation quark squark pair (“ $q\tilde{q}$ ”), into a first generation lepton slepton pair (“ $l\tilde{l}$ ”), and into five quarks and five squarks (“ $5 \times q\tilde{q}$ ”). We only include events that emerge from an angle at least five degrees below the horizon

Event rates for higgsino-like $\tilde{\chi}_1^0$					
$M_X = 10^{12}$ GeV	$N_{\tilde{\chi}_1^0}$	$N_{\nu_\tau}$	$M_X = 10^{16}$ GeV	$N_{\tilde{\chi}_1^0}$	$N_{\nu_\tau}$
$E_{\text{vis}} \geq 10^6$ GeV			$E_{\text{vis}} \geq 10^6$ GeV		
$q\bar{q}$	0.56	0.44	$q\bar{q}$	0.033	0.050
$q\tilde{q}$	1.77	0.79	$q\tilde{q}$	0.024	0.035
$l\tilde{l}$	25.19	1.59	$l\tilde{l}$	0.022	0.033
$5 \times q\tilde{q}$	14.84	5.03	$5 \times q\tilde{q}$	0.038	0.055
$E_{\text{vis}} \geq 10^9$ GeV			$E_{\text{vis}} \geq 10^9$ GeV		
$q\bar{q}$	0.00883	0.00001	$q\bar{q}$	0.003187	0.000004
$q\tilde{q}$	0.08676	0.00001	$q\tilde{q}$	0.002685	0.000003
$l\tilde{l}$	4.09828	0.00002	$l\tilde{l}$	0.006532	0.000003
$5 \times q\tilde{q}$	0.17046	0.00005	$5 \times q\tilde{q}$	0.003668	0.000003



Table 2  
 Predicted event rates for bino-like LSP, for the same combinations of  $E_{\min}$ ,  $M_X$  and primary  $X$  decay mode as in Table 1. We show results for the three different mSUGRA scenarios of [14], with first generation squark masses of about 370 GeV (D1), 580 GeV (D2) and 1000 GeV (D3). The background is essentially the same as in Table 1

Event rates for bino-like $\tilde{\chi}_1^0$							
$M_X = 10^{12}$ GeV	$N_{D_1}$	$N_{D_2}$	$N_{D_3}$	$M_X = 10^{16}$ GeV	$N_{D_1}$	$N_{D_2}$	$N_{D_3}$
$E_{\text{vis}} \geq 10^6$ GeV				$E_{\text{vis}} \geq 10^6$ GeV			
$q\bar{q}$	0.0191	0.0192	0.0118	$q\bar{q}$	0.00095	0.00103	0.00075
$q\tilde{q}$	0.0471	0.0528	0.0388	$q\tilde{q}$	0.00070	0.00077	0.00055
$l\bar{l}$	0.3560	0.5376	0.5543	$l\bar{l}$	0.00079	0.00117	0.00062
$5 \times q\tilde{q}$	0.4567	0.4779	0.3051	$5 \times q\tilde{q}$	0.00113	0.00122	0.00088
$E_{\text{vis}} \geq 10^9$ GeV				$E_{\text{vis}} \geq 10^9$ GeV			
$q\tilde{q}$	0.00007	0.00070	0.00143	$q\tilde{q}$	0.000006	0.000058	0.000140
$q\tilde{q}$	0.00030	0.00314	0.00701	$q\tilde{q}$	0.000005	0.000047	0.000107
$l\bar{l}$	0.00567	0.06121	0.14800	$l\bar{l}$	0.000015	0.000149	0.000175
$5 \times q\tilde{q}$	0.00201	0.01982	0.03967	$5 \times q\tilde{q}$	0.000006	0.000067	0.000161

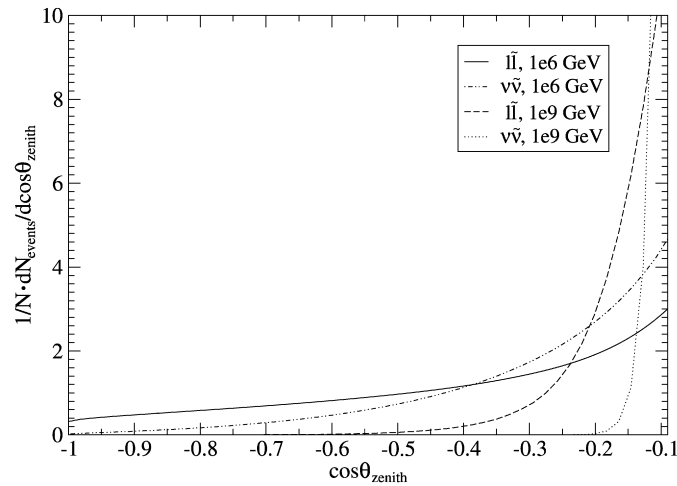


Fig. 1. Angular dependence of the signal from higgsino-like neutralinos from primary  $X \rightarrow l\bar{l}$  decays, and of the  $\nu_\tau$  induced background, for two different values of the lower limit on the visible energy.

large, but the fluxes are relatively soft. One then again needs a rather large normalization factor to reproduce the desired proton flux (1) at  $E = 10^{11}$  GeV. Since the  $\tilde{\chi}_1^0$  spectrum is quite soft, increasing  $E_{\min}$  from  $10^6$  to  $10^9$  GeV now reduces the predicted signal by nearly two orders of magnitude.

The worst case is  $X$  decay into SM quarks only. This gives a relatively hard proton spectrum. Moreover, superparticles are now only produced in the parton shower. This gives a small ratio of  $\tilde{\chi}_1^0$  to proton fluxes, and a relatively soft  $\tilde{\chi}_1^0$  spectrum. The fourth primary  $X$  decay we considered, into a quark and a squark, also leads to a relatively hard proton flux. However, since a superparticle is produced in the primary  $X$  decay, the  $\tilde{\chi}_1^0$  flux is larger, and significantly harder, than for  $X \rightarrow q\tilde{q}$  decays.

We see that at least three of the four cases might lead to observable signals if  $M_X$  is near its lower bound, and if visible energies around  $10^6$  GeV can be detected. Of course, at that energy one expects a huge number of ordinary CR induced events,  $\sim 1$  event per  $\text{km}^2$  and second or (including the duty cycle)  $\sim 3 \times 10^{11}$  events per year in an experiment observing  $10^5 \text{ km}^2$ , as required for a teraton-scale target mass [24]. One will therefore need an excellent discrimination against such down-going events in order to extract the signal of at best a handful events per year. To that end one may need to sharpen the angular cut somewhat. This may also be desired to further reduce the  $\nu_\tau$  induced background, which in this case is within an order of magnitude of the signal. Fig. 1 shows that for  $E_{\min} = 10^6$  GeV, imposing a stronger angular cut will not reduce the signal very much. This is in accord with the results of Ref. [14], which show large neutralino propagation effects only for LSP energies well beyond  $10^7$  GeV in this case. Note, however, that typically  $E_{\text{vis}} \lesssim 0.1 E_{\tilde{\chi}_1^0, \text{in}}$  for higgsino-like neutralino.

On the other hand, only the most favorable scenario remains observable if  $E_{\min}$  has to be increased to  $10^9$  GeV. On the positive side, the  $\nu_\tau$  induced background is now at least three orders of magnitude smaller than the signal, illustrating that the Earth can indeed be used as a filter. This is fortunate, since Fig. 1 shows that now the angular cut can be sharpened only at the cost of a significant reduction of the signal. However, in most cases one would need tens of Tt-yr to see a convincing signal even for

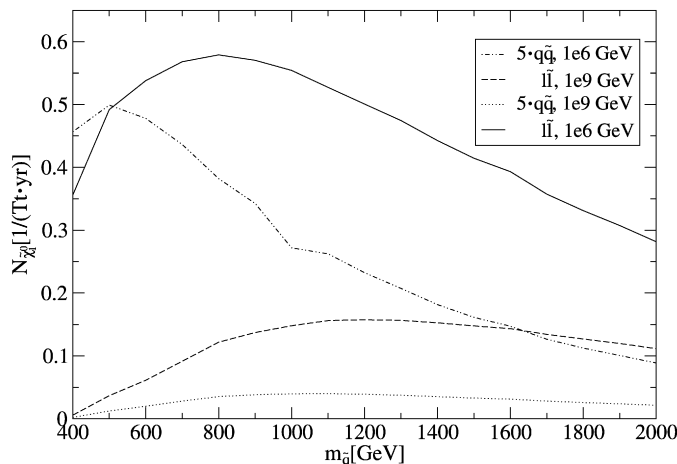


Fig. 2. Expected event rate due to bino-like neutralinos as function of the first generation squark mass, for two different primary  $X$  decay modes and two choices of the minimal visible energy  $E_{\min}$ . See the text for further details.

$M_X = 10^{12}$  GeV; for  $M_X = 10^{16}$  GeV and  $E_{\min} = 10^9$  GeV, one would need P<sub>tyr</sub> of target mass times observation time! This would require monitoring virtually the entire surface of the Earth. The neutralino flux from decays of such very heavy  $X$  particle would remain invisible to teraton scale detectors even for a threshold energy of  $10^6$  GeV. Note that in this case the predicted event rate is almost independent of the primary  $X$  decay mode. The reason is that now the entire relevant energy range satisfies  $x \equiv 2E/M_X \ll 1$ , where the spectrum is determined almost uniquely by the dynamics of the parton shower [12].

Table 2 shows event rates for bino-like neutralino. In this case the scattering cross section depends strongly on the squark mass [17,18,25]. We therefore show results for three different scenarios introduced in Ref. [14], with first generation squark masses near 370, 580 and 1000 GeV, respectively. We see that the event rate remains below one event per year and teraton in all cases. This result seems much less promising than that of earlier studies [13,17]. However, our rates are actually comparable to those of Ref. [17], once the differences in treatment are taken into account. To begin with, we assume that the  $X$  particles are distributed like Dark Matter, i.e. clump in our galaxy. Assuming a uniform distribution throughout the universe, as done in Ref. [17], increases the neutralino flux by about one order of magnitude [13]. The reason is that such a uniform distribution suppresses the proton flux due to the GZK effect. One therefore has to increase the normalization in order to match the observed flux. A more or less uniform distribution of  $X$  particles could be achieved only if they are bound to topological defects. However, scenarios with such defects, like cosmic strings, are nowadays quite tightly constrained by analyses of cosmic microwave background anisotropies [26]. Moreover, we quote events per year, whereas Ref. [17] finds about five events per lifetime of the experiment, taken to be three years. Finally, Ref. [17] applies a cut (of  $10^9$  GeV) on the total energy of the incident neutralino, whereas our cut is on the visible energy.

We note that for  $E_{\min} = 10^6$  GeV, the ten body decay mode and  $X \rightarrow \tilde{l}\tilde{l}$  decays now generally lead to similar event rates. The reason is that very energetic bino-like neutralinos lose energy considerably faster than higgsino-like neutralinos do: for rather light squarks the cross sections are comparable, but the energy loss per scattering is much larger for bino-like states, which produce a squark with  $m_{\tilde{q}} \gg m_{\tilde{\chi}_1^0}$ , than for higgsino-like states, which produce a heavier neutralino or chargino very close in mass to the LSP. The  $5 \times q\tilde{q}$  decay mode has a larger flux of softer neutralinos, which suffers less from propagation effects; for bino-like neutralinos this largely compensates the reduction of the rate due to the fact that the cross section is smaller at smaller LSP energy. However, if  $E_{\text{vis}} > 10^9$  GeV is required, even the relatively softer LSPs produced from the ten body decay mode will typically scatter several times before reaching the detector.  $X \rightarrow \tilde{l}\tilde{l}$  decays are then again more favorable, due to its initially much larger flux of very energetic neutralinos.

This brings us to a feature of our treatment which enhances the event rate compared to the numbers of Ref. [17]. In that analysis all neutralinos were discarded that interact even once before reaching the detector. This is not necessary, since this interaction will again yield a neutralino (from the decay of the produced squark), with typically about half the energy of the original LSP. Fig. 2 shows that this regeneration effect also leads to a much milder dependence of the final event rate on the cross section, and hence on the squark mass, than found in Ref. [17]. Increasing the squark mass reduces the cross section, and hence the event rate for given flux. However, it also reduces the effect of neutralino propagation through the Earth, i.e. it increases the flux. These two effects obviously tend to cancel. As a result the event rate as function of  $m_{\tilde{q}}$  shows a rather broad maximum, the location of which depends on the cut on  $E_{\text{vis}}$ . A lower  $E_{\text{vis}}$  means that softer neutralinos can contribute. Since the cross section increases with neutralino energy, softer neutralinos can tolerate lighter squarks before suffering significant propagation losses. As a result, at smaller  $E_{\min}$  the maximum rate occurs for smaller squark mass. This effect is less pronounced for primary  $X \rightarrow \tilde{l}\tilde{l}$  decays, since in this case the incident neutralino spectrum is in any case rather hard, even if no cut on  $E_{\text{vis}}$  is applied.

#### 4. Summary and conclusions

In this Letter we have calculated signal rates for the detection of very energetic neutralinos, as predicted by “top–down” models for the observed cosmic ray events at the highest energies. We use up-to-date calculations of the neutralino flux generated at the location of the decay of the superheavy particles, and of the effects due to propagation of the neutralinos through the Earth. We also for the first time treat the case of higgsino-like neutralino.

We conservatively assume that the progenitor “ $X$  particles” are distributed like Dark Matter, in which case most sources are “local”, i.e. effects of propagation through the interstellar or intergalactic medium are negligible. We then find detectable event rates in teraton scale experiments with duty cycle of  $\sim 10\%$ , typical for experiments based on optical methods, only if the following conditions are satisfied: the lightest neutralino must be a higgsino, rather than a bino;  $M_X$  must be rather close to its lower bound of  $\sim 10^{12}$  GeV; and the experiment must either be able to detect upgoing events with visible energy not much above  $10^6$  GeV, or most  $X$  particles undergo two-body decays involving at least one slepton and no strongly interacting (s)particle. The good news is that in all cases we studied the signal is at least several times larger than the  $\nu_\tau$  induced background, computed in the same  $X$  decay model. If  $M_X$  is near  $10^{16}$  GeV and the LSP is higgsino-like, or  $M_X \sim 10^{12}$  GeV and the LSP is bino-like, one will need  $\mathcal{O}(100)$  Tt-yr to collect a respectable event rate. In the worst case, with a bino-like LSP,  $M_X \sim 10^{16}$  GeV and a threshold of the visible energy near  $10^9$  GeV, one would observe less than one event per year even if one monitored the entire surface of the Earth! These numbers improve by about one order of magnitude if  $X$  particles are distributed more or less uniformly throughout the universe; this might be expected if they are confined to cosmic strings or similar topological defects. Recall, however, that scenarios with cosmic strings are constrained by observations of cosmic microwave anisotropies.

These numbers only include interactions of neutralinos with nuclei. It has been claimed in Ref. [18] that bino-like LSPs should lead to a detectable signal in Gt class experiments (like IceCube [27]) through resonant production of sleptons. However, they estimate the rates assuming a neutralino flux close to the upper bound on the neutrino flux; the kind of model we investigate here yields fluxes that are several orders of magnitude smaller than this. Moreover, the visible energy in such events is relatively small, since only the decay of the produced slepton contributes. At the relevant energies the Earth does not filter tau neutrinos very well; so even if one concentrates on upgoing events, the background in potentially realistic  $X$  decay models is several orders of magnitude larger than the signal.

Our overall conclusion is that next generation experiments, with effective target masses in the Tt range, would have to be lucky to observe a signal from neutralinos of “top–down” origin. Experiments with a relatively low energy threshold would stand a much better chance than those with high threshold. Unfortunately there are many reasonable  $X$  decay scenarios where the neutralino flux will remain invisible to such experiments. The goal of finding an experimentum crucis for top-down models may therefore remain elusive.

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