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Signatures of Collisional Annihilation in Ultra-High Energy Cosmic Rays

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Abstract

We first review the motivations and main aspects of the proposal that ultra high energy cosmic rays originate in collisional annihilation of superheavy dark matter particles in our galactic halo. We then emphasize that the various proposals for the origin of ultra high energy cosmic rays predict qualitatively different patterns of arrival directions for primary cosmic ray energies beyond the GZK bound. We expect that in about three years the anisotropy signal observed by the Pierre Auger observatory will tell us which proposal was correct.

I. INTRODUCTION

The existence of dark matter in the halos of galaxies has been established observationally through gravitational lensing and through the peculiar motions of galaxies and galaxy clusters, and on the theoretical side through the need of a sufficiently early start for structure formation and bounds on ordinary matter from the very successful theory of big bang nucleosynthesis. To understand the nature and composition of the dark matter is one of the most interesting challenges of modern science.

The existence of ultrahigh energy cosmic rays (UHECRs) with energies beyond the Greisen-Zatsepin-Kuzmin cutoff, $E > E_{GZK} \simeq 40\,\mathrm{EeV}$, and their lack of correlation with nearby active galactic nuclei constitute another major contemporary puzzle in science. The three main aspects of the UHECR puzzle are:
1. Scattering off cosmic microwave background photons limits the penetration depths of charged particles with energies $E > E_{GZK}$ to distances $< 100\,\mathrm{Mpc}$[1–6];
2. the distribution of arrival directions of UHECRs does not seem to favor any known astrophysical sources within the GZK cutoff length;
3. it seems extremely difficult to devise sufficiently efficient astrophysical acceleration mechanisms which could accelerate particles to energies $E > E_{GKZ}$, and at the same time overcome collisional and radiation losses.

Top-down models for the origin of ultrahigh energy cosmic rays combine both problems by proposing that ultrahigh energy cosmic rays arise in the decay[7–11] or collisional annihilation[12–14] of superheavy dark matter particles. We will use both the acronym SHDM and WIMPZILLA[1] for superheavy dark matter particles .

[1]The word WIMPZILLA was coined by Kolb, Chung and Riotto in their investigations of possible origins of superheavy dark matter particles[15].

SHDM decay can arise through direct decay of relic particles ("WIMPZILLA" decay) or through the annihilation of superheavy relic bound states ("WIMPZILLIUM" decay), but the underlying decay mechanism has no effect on the predicted pattern of UHECR arrival directions. This pattern differs strongly from the anisotropy pattern predicted by collisional WIMPZILLA annihilation. The reason for the difference is that the anisotropy pattern predicted by decay models is proportional to the WIMPZILLA or WIMPZILLIUM density $n_X(\mathbf{r})$, and is dominated by the smooth background halo. Collisional annihilation, on the other hand, predicts that the anisotropy pattern should be proportional to $n_X^2(\mathbf{r})$, and is constrained by unitarity limits on annihilation cross sections. Therefore collisional annihilation can only work in dense cores of dark matter substructure in the galactic halo. These dense cores were denoted as WIMPZILLA stars[14]. As a consequence, the collisional annihilation scenario predicts a pointlike source distribution with increasing density towards the galactic center.

All modern acceleration ("bottom-up") models for UHECR origin assume powerful extragalactic sources, and therefore predict that the pattern of observed arrival directions should not correlate with the galactic halo. These differences of predictions for UHECR arrival directions imply that a dedicated large statistics experiment like the Pierre Auger observatory will be able to identify the correct UHECR source model from its anisotropy signal.

Sec. II explains the origin of a mass window for the generation of relic superheavy dark matter particles during inflation. Sec. III summarizes and updates the calculation of the ultrahigh energy cosmic ray flux from collisional WIMPZILLA annihilation, and the origin of the unique pattern of arrival directions predicted by that model. In Sec. IV we compare the different anisotropy signatures predicted by all contemporary proposals for the origin of ultrahigh energy cosmic rays, and Sec. V contains our conclusions.

For comparison with experimental values, please note that $j(E) = dI/dE$ is the spectral flux, i.e. the differential flux per energy interval, and $J(E) = d^2I/dEd\Omega$ is the differential flux per energy interval and per unit of solid angle.

II. SUPERHEAVY DARK MATTER FROM INFLATION

Particle creation as a consequence of non-adiabatic expansion was discovered already in the late 1950s and 1960s, see[16] and references there. The effect was found to be negligible in radiation or dust dominated epochs; however, this mechanism for non-thermal particle creation was rediscovered and garnered much more interest after the necessity for inflation was realized. The effect is usually considered in terms of the Bogolubov transformation between in and out vacua in an expanding universe[15,17–20]. Particle creation during preheating after inflation can also arise as a consequence of a direct coupling between the inflaton and other matter fields[21,22].

Here we will expand on a simple discussion given in[14] to see how particle production in an inflationary universe can be understood by studying the evolution equations of weakly coupled scalar fields in the expanding universe. The reasoning outlined here is not intended to compete in any way with the traditional operator methods to study particle production from spatial expansion, but it may provide a complementary and helpful view.

If interactions with other matter fields can be neglected, a scalar field in a Friedmann–Robertson–Walker background with metric

$$\begin{aligned} ds^2 &= -dt^2 + a^2(t)\left(\frac{dr^2}{1-kr^2} + r^2 d\vartheta^2 + r^2\sin^2\vartheta\, d\varphi^2\right) \\ &= -dt^2 + a^2(t)\left(d\mathbf{x}^2 + \mathrm{k}\frac{(\mathbf{x}\cdot d\mathbf{x})^2}{1-\mathrm{k}\mathbf{x}^2}\right) \end{aligned} \tag{2.1}$$

satisfies

$$\ddot{\phi}(\mathbf{x},\mathrm{t}) + 3\frac{\dot{\mathrm{a}}(\mathrm{t})}{\mathrm{a}(\mathrm{t})}\dot{\phi}(\mathbf{x},\mathrm{t}) - \frac{1}{\mathrm{a}^2(\mathrm{t})}\left[\left(\delta^{\mathrm{ij}} - \mathrm{k}\mathrm{x}^{\mathrm{i}}\mathrm{x}^{\mathrm{j}}\right)\partial_{\mathrm{i}}\partial_{\mathrm{j}} - 3\mathrm{k}\mathrm{x}^{\mathrm{i}}\partial_{\mathrm{i}}\right]\phi(\mathbf{x},\mathrm{t}) + \mathrm{m}^2\phi(\mathbf{x},\mathrm{t}) = 0, \tag{2.2}$$

and the corresponding energy density per unit of comoving volume is

$$\begin{aligned}\varrho(\mathbf{x},\mathrm{t}) &= \sqrt{-g}\,T^{00} \\ &= \frac{a^3(t)}{2\sqrt{1-kr^2}}\left[\dot{\phi}^2(\mathbf{x},\mathrm{t})+\mathrm{m}^2\phi^2(\mathbf{x},\mathrm{t})+\frac{1}{\mathrm{a}^2(\mathrm{t})}\left(\delta^{\mathrm{ij}}-\mathrm{kx^ix^j}\right)\partial_{\mathrm{i}}\phi(\mathbf{x},\mathrm{t})\cdot\partial_{\mathrm{j}}\phi(\mathbf{x},\mathrm{t})\right].\end{aligned} \tag{2.3}$$

The violation of time translation invariance in (2.1) implies violation of energy conservation, of course:

$$\dot{\varrho}(\mathbf{x},\mathrm{t})-\mathrm{a}(\mathrm{t})\partial_{\mathrm{i}}\left[\dot{\phi}(\mathbf{x},\mathrm{t})\frac{\delta^{\mathrm{ij}}-\mathrm{kx^ix^j}}{\sqrt{1-\mathrm{kr}^2}}\partial_{\mathrm{j}}\phi(\mathbf{x},\mathrm{t})\right] \tag{2.4}$$

$$=\frac{a^2(t)\dot{a}(t)}{2\sqrt{1-kr^2}}\left[3\left(m^2\phi^2(\mathbf{x},\mathrm{t})-\dot{\phi}^2(\mathbf{x},\mathrm{t})\right)+\frac{1}{a^2(t)}\left(\delta^{ij}-kx^ix^j\right)\partial_i\phi(\mathbf{x},\mathrm{t})\cdot\partial_{\mathrm{j}}\phi(\mathbf{x},\mathrm{t})\right].$$

A priori the sign of the expression on the right hand side is indefinite, e.g. a rapidly evolving field of low mass and with small spatial fluctuations loses energy during spatial expansion as long as it evolves fast enough. However, Eq. (2.2) tells us that spatial fluctuations (which should never have been large anyway for the Friedmann–Robertson–Walker *ansatz* to work) are soon negligible in the inflationary expanding universe $a(t) \propto \exp(Ht)$. This leaves us with the simple equation

$$\ddot{\phi}(t)+3H\dot{\phi}(t)+m^2\phi(t)\simeq 0. \tag{2.5}$$

Approximately constant H during inflation then yields for the time evolution of the comoving energy density

$$\begin{aligned}\varrho(t) &\simeq \frac{1}{2}a^3(t)\left(\dot{\phi}^2(t)+m^2\phi^2(t)\right) \\ &\simeq A_+\exp\left(t\sqrt{9H^2-4m^2}\right)+A_-\exp\left(-t\sqrt{9H^2-4m^2}\right)+B.\end{aligned}$$

This implies a growing mode in the comoving energy density of weakly coupled scalars with $m < 1.5H \simeq 10^{14}$ GeV. What is special about the superheavy particles is that their comoving energy density is conserved after inflation, because the behavior of massive ($m > t^{-1}$) weakly coupled states in the subsequent radiation and dust dominated backgrounds preserves their energy. The asymptotic solution for weakly coupled massive scalars with $m > t^{-1}$ in such a background yields (with $\ell = 3$ for dust, $\ell = 4$ for radiation)

$$\phi(t) \propto t^{-3/\ell}\cos(mt+\varphi)\,,$$

$$\varrho(t) \propto a^3(t)t^{-6/\ell} \propto t^0,$$

and this implies in particular that the comoving density of massive particles freezes out at the end of inflation ($t \simeq 10^{-36}$ s) if

$$m > t^{-1} \simeq 10^{12}\,\mathrm{GeV}.$$

These considerations indicate a mass window for direct gravitational production of superheavy relic particles during inflation

$$10^{12}\,\mathrm{GeV} < m < 10^{14}\,\mathrm{GeV}.$$

After inflation intrinsically unstable matter of mass M will decay on time scales $\tau \lesssim M^2_{Planck}/M^3$, where the upper bound assumes that the particles couple to their decay products at least with gravitational strength. This means that *intrinsically unstable* superheavy particles should not be relic particles. This difficulty has motivated the collisional annihilation scenario for ultrahigh energy cosmic rays from superheavy dark matter[12], because the unitarity limit on reaction cross sections

$$\sigma_A v \lesssim \frac{4\pi\hbar^2}{M^2 v} \tag{2.6}$$

implies that superheavy dark matter without direct decay channels will still be around.

III. CALCULATION OF THE FLUX

We consider decay or annihilation of superheavy dark matter particles of mass $M_X \geq 10^{12}$ GeV. The spectral fluxes at our location $\mathbf{r}_\odot$ from decay or collisional annihilation of the dark matter particles of density $n_X(\mathbf{r})$ are then

$$j_d(E) = \frac{d\mathcal{N}(E, M_X)}{dE} \int d^3\mathbf{r}\, \frac{1}{4\pi|\mathbf{r}_\odot - \mathbf{r}|^2} \frac{\mathrm{n_X}(\mathbf{r})}{\tau_\mathrm{d}} \tag{3.1}$$

and

$$j_a(E) = \frac{d\mathcal{N}(E, 2M_X)}{dE} \int d^3\mathbf{r}\, \frac{\nu}{16\pi|\mathbf{r}_\odot - \mathbf{r}|^2} \mathrm{n_X}(\mathbf{r})^2 \langle \sigma_\mathrm{A} \mathrm{v} \rangle, \tag{3.2}$$

respectively. Here $d\mathcal{N}(E, E_{in})/dE = E_{in} d\mathcal{N}(x, E_{in})/dx$ is the number of particles in the energy interval $[E, E + dE]$ emerging from a decay or annihilation event of initial energy E_{in}. $x = E/E_{in}$ is a scaled energy variable. $d\mathcal{N}(E, E_{in})/dE$ is related to fragmentation functions via

$$\frac{d\mathcal{N}(E, E_{in})}{dE} = \sum_i \frac{1}{\sigma_A} \frac{d\sigma^{(i)}}{dE} = \frac{1}{E_{in}} \sum_i F^{(i)}(x, E_{in}^2),$$

where $F^{(i)}(x, E_{in}^2)$ is the differential number of particles of species i generated in the prescribed x-range in a decay or annihilation event with $s = E_{in}^2$. The factor ν in Eq. (3.2) equals 4 if the X-particles are Majorana particles, and 1 otherwise.

N-body simulations predict that usually about $f_{cl} \sim 5 - 10\%$ of dark matter halos should exist in substructure. Therefore in the decay models the ultrahigh energy cosmic ray flux will be dominated by the smooth background halo. Evaluating (3.1) e.g. for a Navarro-Frenk-White halo[23]

$$n_X(r) = 4n_X(r_s) r_s^3 / r(r + r_s)^2$$

yields

$$j_d(E) = 4 \frac{d\mathcal{N}(E, M_X)}{dE} \frac{n_X(r_s)}{\tau_d} \frac{r_s^3}{r_s^2 - r_\odot^2} \ln\left(\frac{r_s}{r_\odot}\right). \tag{3.3}$$

The scale radius r_s for the Milky Way is not well known, but for the whole range of say $5\,\mathrm{kpc} \leq \mathrm{r_s} \leq 50\,\mathrm{kpc}$ comparison between (3.3) and the ultrahigh energy cosmic ray flux from AGASA[24] yields that decaying superheavy dark matter should only make a small contribution to the galactic dark matter density. Of course, the problem is to explain a lifetime $\tau_d \geq 10^{17}$ s for unstable superheavy particles.

Eq. (3.2) can also be evaluated analytically for a Navarro-Frenk-White halo, but the unitarity limit for s-wave annihilation

$$\sigma_A v = \xi \times \frac{4\pi\hbar^2}{M_X^2 v} = \xi \times 4.40 \times 10^{-43}\,\mathrm{m}^3/\mathrm{s} \times \left(\frac{10^{12}\,\mathrm{GeV}}{M_X}\right)^2 \times \frac{100\,\mathrm{km/s}}{v},\ \xi \leq 1, \tag{3.4}$$

implies that any ultrahigh energy cosmic ray flux from collisional annihilation in the background halo is negligible. Therefore collisional WIMPZILLA annihilation can only make a noticeable contribution to the ultrahigh energy cosmic ray flux if it originates in dense cores of dark matter subclumps of the galactic halo[12]. These dense cores were denoted as WIMPZILLA stars in[14], and a simplified estimate of the cosmic ray flux from WIMPZILLA stars yields

$$j_a(E) \simeq \nu \frac{N_{cl}\overline{V}_{core}}{16\pi d^2} \frac{d\mathcal{N}(E, 2M_X)}{dE} n_X^2 \langle \sigma_A v \rangle \simeq \nu \frac{0.1 f_{cl} M_{halo}}{16\pi d^2 M_X} \frac{d\mathcal{N}(E, 2M_X)}{dE} n_X \langle \sigma_A v \rangle. \tag{3.5}$$

Here $d^{-2} = \langle r^{-2} \rangle$ is a mean inverse distance squared for our separation from galactic WIMPZILLA stars. In our estimate in[14] we had used $d \simeq 10\,\mathrm{kpc}$, which seems reasonable: e.g. our mean inverse distance squared to visible galactic substructures, the globular clusters, corresponds to $d = 7.28\,\mathrm{kpc}$. We will use the latter as a fiducial value for our calculations in the present paper, although we think that visible substructure will not trace dark matter substructure.

We have parameterized the unknown annihilation cross section already in terms of the unitarity bound (3.4). For the comparison with observed spectra we also parameterize the core density n_X of WIMPZILLA stars in terms of the solar density

$$n_X = \eta \times \frac{\varrho_\odot}{10^{12}\,\mathrm{GeV}} = \eta \times 7.89 \times 10^{17}\,\mathrm{m}^{-3}. \tag{3.6}$$

Insertion of Eqs. (3.4) and (3.6) into (3.5) determines the spectral current from collisional annihilation as a function of the unknown product $\nu\eta\xi$. For the rescaled spectral flux per steradian the result reads

$$E^3 J_a(E) \simeq \nu\eta\xi \frac{0.1 f_{cl} M_{halo} \hbar^2 c^4}{4\pi\,\mathrm{sr}\,\mathrm{d}^2 \mathrm{M}_\mathrm{X}^2 \mathrm{v}} x^3 \frac{d\mathcal{N}(x, 2M_X)}{dx} \varrho_\odot. \tag{3.7}$$

A fit of Eq. (3.7) to the ultrahigh energy cosmic ray spectrum as observed by AGASA[24] yields

$$\begin{aligned} \log(\nu\eta\xi) &= -2.95 \pm 0.28 - \log\left[\frac{f_{cl}}{0.05}\right] - \log\left[\frac{M_{halo}}{2\times 10^{12} M_\odot}\right] \\ &\quad + 2\log\left[\frac{d}{7.28\,\mathrm{kpc}}\right] + 2\log\left[\frac{M_X c^2}{10^{12}\,\mathrm{GeV}}\right] + \log\left[\frac{v}{100\,\mathrm{km/s}}\right] \end{aligned} \tag{3.8}$$

or

$$\nu\eta\xi \simeq 1.1 \times 10^{-3}. \tag{3.9}$$

It is reassuring for the collisional annihilation theory that the fit allows for small parameters η and ξ for the cross section and the typical core density.

Before we go on to discuss anisotropy patterns, we would also like to discuss the calculation of the flux under the assumption of two initial jets emerging from the decay or annihilation event. In that case the fragmentation functions $d\mathcal{N}(E, E_{in})/dE$ in the previous equations have to be replaced by jet fragmentation functions according to

$$d\mathcal{N}(E, E_{in})/dE \longrightarrow 2d\mathcal{N}(E, E_{jet})/dE = 2d\mathcal{N}(E, E_{in}/2)/dE.$$

Practically this makes a difference in fits to observed spectra, because all the fits use contemporary models or calculations for fragmentation functions which satisfy the continuity condition

$$\left.\frac{d\mathcal{N}(E, E')}{dE}\right|_{E=E'} = 0,$$

but do not vanish for $E = E'/2$. Therefore in all practical applications using one and the same model for fragmentation functions yields slightly different results for the two ways of calculation. This is because the direct calculation using Eqs. (3.1,3.2) allows for events where one particle may carry away almost all the energy from the decay or annihilation event, whereas in the model with two primary jets single particles may carry away at most one half of the energy from the decay or annihilation event.

With two primary jets the equations for fluxes from particle decay or annihilation become

$$\begin{aligned} j_d(E) &= \frac{d\mathcal{N}(E, M_X/2)}{dE} \int d^3\mathbf{r}\, \frac{1}{2\pi|\mathbf{r}_\odot - \mathbf{r}|^2} \frac{\mathrm{n_X}(\mathbf{r})}{\tau_\mathrm{d}} \\ &= 8\frac{d\mathcal{N}(E, M_X/2)}{dE} \frac{n_X(r_s)}{\tau_d} \frac{r_s^3}{r_s^2 - r_\odot^2} \ln\left(\frac{r_s}{r_\odot}\right), \end{aligned} \tag{3.10}$$

$$j_a(E) = \frac{d\mathcal{N}(E, M_X)}{dE} \int d^3\mathbf{r}\, \frac{\nu}{8\pi|\mathbf{r}_\odot - \mathbf{r}|^2} \mathrm{n_X}(\mathbf{r})^2 \langle \sigma_\mathrm{A} \mathrm{v}\rangle, \tag{3.11}$$

where $d\mathcal{N}(E,E_{jet})/dE = E_{jet}d\mathcal{N}(x,E_{jet})/dx$ is the number of particles in the energy interval $[E, E+dE]$ emerging from a jet of initial energy E_{jet}. The estimate of the cosmic ray flux from WIMPZILLA stars then yields

$$j_a(E) \simeq \nu\frac{N_{cl}\overline{V}_{core}}{8\pi d^2}\frac{d\mathcal{N}(E,M_X)}{dE}n_X^2\langle\sigma_A v\rangle \simeq \nu\frac{0.1f_{cl}M_{halo}}{8\pi d^2 M_X}\frac{d\mathcal{N}(E,M_X)}{dE}n_X\langle\sigma_A v\rangle$$
$$= \nu\eta\xi\frac{0.1f_{cl}M_{halo}\hbar^2}{2d^2M_X^4 v}\frac{d\mathcal{N}(E,M_X)}{dE}\varrho_\odot, \tag{3.12}$$

or

$$E^3J_a(E) \simeq \nu\eta\xi\frac{0.1f_{cl}M_{halo}\hbar^2c^4}{8\pi\,\mathrm{sr}\,\mathrm{d}^2\mathrm{M}_\mathrm{X}^2\mathrm{v}}x^3\frac{d\mathcal{N}(x,M_X)}{dx}\varrho_\odot. \tag{3.13}$$

The fit to the AGASA ultrahigh energy cosmic ray spectrum yields

$$\log(\nu\eta\xi) = -3.05\pm 0.26 - \log\left[\frac{f_{cl}}{0.05}\right] - \log\left[\frac{M_{halo}}{2\times 10^{12}M_\odot}\right]$$
$$+2\log\left[\frac{d}{7.28\,\mathrm{kpc}}\right] + 2\log\left[\frac{M_Xc^2}{10^{12}\,\mathrm{GeV}}\right] + \log\left[\frac{v}{100\,\mathrm{km/s}}\right] \tag{3.14}$$

or

$$\nu\eta\xi \simeq 8.9\times 10^{-4}. \tag{3.15}$$

The two fits (3.8) and (3.14) are shown in Fig. III.

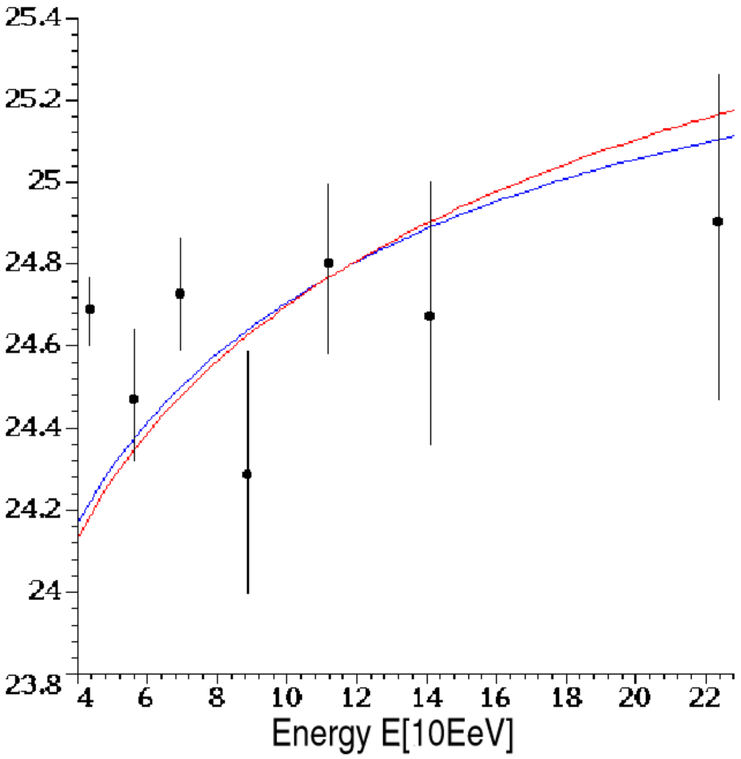

FIG. 1. The fits (3.8) (red) and (3.14) (blue) to the ultrahigh energy cosmic ray spectrum from AGASA. The y-axis is $\log(E^3J(E)/\mathrm{eV}^2\,\mathrm{m}^{-2}\,\mathrm{s}^{-1}\,\mathrm{sr}^{-1})$.

These fits have large potential uncertainties from the relatively small statistics of ultrahigh energy events observed in the era before the Pierre Auger observatory. However, the results (3.8,3.14) and the comparable results in[14,25] indicate that $\nu\eta\xi \sim 10^{-3}$ seems to be a good current estimate.

IV. EXPECTED ANISOTROPY PATTERNS BEYOND THE GZK BOUND

The fits (3.8,3.14) to the ultrahigh energy cosmic ray spectrum indicate that in the collisional annihilation scenario, cosmic rays above the Greisen-Zatsepin-Kuzmin bound should be dominated by the fragmentation products of WIMPZILLA annihilation. Here we want to emphasize that the anisotropy pattern observed by the Pierre Auger observatory will provide a crucial direct test of the different proposals for the origin of ultrahigh energy cosmic rays[25]: Bottom-up acceleration of charged particles in AGNs (see e.g.[26]) or gamma ray bursts[27] should naturally correlate with local superstructures if it is not connected to ultrahigh energy neutrinos traveling over cosmological distances. However, if ultrahigh energy neutrinos from the early universe play a role, they should create Z-bursts without a strong correlation to local superstructure[28]. For the top-down scenarios, on the other hand, WIMPZILLA or WIMPZILLIUM decay[7–11] should be dominated by the uniform component of our galactic halo, whereas collisional annihilation in WIMPZILLA stars would be dominated by relatively dense cores of halo substructure[12,14].

Origin of UHECRs	Expected anisotropy pattern
Collisional annihilation in WIMPZILLA stars	Pointlike sources with increasing density towards the galactic center. No direct correlation with galactic SN or SNRs.
WIMPZILLA or WIMPZILLIUM decay	Dominated by uniform increase towards galactic center.
Z-bursts	Approximately isotropic distribution with only weak correlation to structure within 150 Mpc.
Bottom-up acceleration	Correlation with local superstructure within 150 Mpc.

Once fully operational, the Pierre Auger observatory should see several hundred events per year with energies above the GZK bound. If the collisional annihilation scenario is correct, the anisotropy pattern in this energy range should include a large number of multiplets within the angular resolution of the observatory. The density of the multiplets will increase towards the galactic center.

V. CONCLUSIONS

Collisional annihilation of superheavy dark matter particles as a source model for ultrahigh energy cosmic rays has the advantage of naturally explaining the absence of a GZK cutoff in the spectrum without correlation to local AGNs, without the need of postulating an extremely powerful and efficient acceleration mechanism, and without the need to explain extremely long lifetimes of intrinsically unstable particles. It has the disadvantage of having to postulate formation of a few relatively dense cores of dark matter subhalos or subhalo remnants.

The two different top-down scenarios of decay or collisional annihilation of superheavy particles imply qualitatively different anisotropy signals in cosmic ray arrival directions, and these predictions also differ from the anisotropy signals predicted by the various bottom-up scenarios. For the collisional annihilation scenario the pattern of arrival directions should be dominated by a large number of multiplets within the angular resolution of the observatory. Based on current particle physics extrapolations for chemical composition of cosmic rays to very high energies, bottom-up acceleration scenarios enjoy much more popularity in the cosmic ray community. However, collisional WIMPZILLA annihilation is not an overly speculative theory, and based on its expected anisotropy pattern it should be confirmed or rejected within a few years. If we are right, the observational confirmation of collisional WIMPZILLA annihilation through the Pierre Auger observatory will open a window into the epoch of inflationary expansion of the universe and to an entirely new field of ultrahigh energy particle physics. It would also imply that the Pierre Auger observatory would provide us with an unprecedented view on galactic halo substructure, thus opening a new field of dark matter astronomy.

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