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Ultra-high energy cosmic rays from annihilation of superheavy dark matter

Pasquale Blasi [a,1], Rainer Dick [b,2], and Edward W. Kolb [a,c,3]

[a] *NASA/Fermilab Astrophysics Center, Fermi National Accelerator Laboratory, Batavia, Illinois 60510-0500*

[b] *Department of Physics and Engineering Physics, University of Saskatchewan, 116 Science Place, Saskatoon, SK S7N 5E2, Canada*

[c] *Department of Astronomy and Astrophysics, Enrico Fermi Institute, The University of Chicago, Chicago, Illinois 60637-1433*

Abstract

We consider the possibility that ultra-high energy cosmic rays originate from the annihilation of relic superheavy dark-matter particles. We find that a cross section of $\langle\sigma_A v\rangle \sim 10^{-26}\mathrm{cm}^2(M_X/10^{12}\,\mathrm{GeV})^{3/2}$ is required to account for the observed rate of super-GZK events if the superheavy dark matter follows a Navarro–Frenk–White density profile. We also calculate the possible signature from annihilation in sub-galactic clumps of dark matter and find that the signal from sub-galactic structures may dominate. Finally, we discuss the expected anisotropy in the arrival directions of the cosmic rays, which is a characteristic signature of this scenario.

Key words: Dark Matter, Ultra-High Energy Cosmic Rays
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1 Introduction

The nature of the dark matter and the origin of the ultra-high energy (UHE) cosmic rays are two of the most pressing issues in contemporary particle astrophysics. In this paper we explore the possibility that there is a common resolution of the two issues: dark matter is a relic superheavy dark-matter particle (WIMPZILLA), and the UHE cosmic rays are the annihilation products of WIMPZILLAS.

The dark matter puzzle results from the observation that visible matter can account for only a very small fraction of matter bound in large-scale structures. The evidence for dark matter is supported by mass estimates from gravitational lensing [1], by the peculiar velocities of large scale structures [2], by measurements of CMB anisotropy [3] and by measurements of the recession velocity of high-redshift supernovae [4]. Constraints from big-bang nucleosynthesis imply that the bulk of the dark matter cannot be baryonic, and most of the matter density of the universe must

1 blasi@carina.fnal.gov
2 rainer@sask.usask.ca
3 rocky@rigoletto.fnal.gov

arise from particles not accounted for by the Standard Model of particle physics [5].

The existence of UHE cosmic rays of energies above the Greisen-Zatsepin-Kuzmin cutoff [6], $E_{\rm GZK} \simeq 5 \times 10^{19}\,$eV, is a major puzzle because the cosmic microwave background constitutes an efficient obstacle for protons or nuclei of ultra-high energies to travel farther than a few dozen Mpc [6,7]. This suggests that the observed extremely energetic cosmic rays with $E > E_{\rm GZK}$ should originate in our cosmic neighborhood. Furthermore, the approximately isotropic distribution of arrival directions makes it difficult to imagine that nearby astrophysical sources are the accelerators of the observed UHE cosmic rays (but for the opposite point of view, see [8]).

An interesting possibility is that UHE cosmic rays are the products of the decay of some superheavy particle. This possibility has been proposed by Berezinsky, Kachelrieß and Vilenkin and by Kuzmin and Rubakov [9,10], see also [11] for a discussion of this in the framework of string/M-theory, and [12] for a discussion in the framework of topological defects. The superheavy particles must have masses $M_X \geq 10^{12}\,$GeV. Although this proposal circumvents some of the astronomical problems, there are two issues to address: Some cosmological mechanism must be found for producing particles of such large mass in the necessary abundance, and the lifetime of this very massive state must be in excess of $10^{20}\,$yr, if we want these particles to be both dark matter candidates and sources of UHE cosmic rays.

The simple assumption that dark matter is a thermal relic limits the maximum mass of the dark matter particle. The largest annihilation cross section in the early universe is expected to be roughly M_X^{-2}. This implies that very massive WIMPS have such a small annihilation cross section that their present abundance would be too large if the WIMPS are thermal relics. Thus, one predicts a maximum mass for a thermal WIMP, which turns out to be a few hundred TeV. While a thermal origin for WIMPS is the most common assumption, it is not the simplest possibility. It has been recently pointed out that dark particles might have never experienced local chemical equilibrium during the evolution of the universe, and that their mass may be in the range 10^{12} to 10^{19} GeV, much larger than the mass of thermal WIMPS [13–16]. Since these WIMPS would be much more massive than thermal WIMPS, such superheavy dark particles have been called WIMPZILLAS [16].

Since WIMPZILLAS are extremely massive, the challenge lies in creating very few of them. Several WIMPZILLA scenarios have been developed involving production during different stages of the evolution of the universe.

WIMPZILLAS may be created during bubble collisions if inflation is completed through a first-order phase transition [17,18]; at the preheating stage after the end of inflation with masses easily up to the Grand Unified scale of 10^{15}GeV [19] or even up to the Planck scale [20–22]; or during the reheating stage after inflation [15] with masses which may be as large as 2×10^3 times the reheat temperature.

WIMPZILLAS may also be generated in the transition between an inflationary and a matter-dominated (or radiation-dominated) universe due to the "nonadiabatic" expansion of the background spacetime acting on the vacuum quantum fluctuations. This mechanism was studied in details in Refs. [13,23,24]. The distinguishing feature of this mechanism is the capability of generating particles with mass of the order of the inflaton mass (usually much larger than the reheat temperature) even when the particles only interact extremely weakly (or not at all) with other particles, and do not couple to the inflaton.

The long lifetime required in the decay scenario for UHE cosmic rays is problematic if the UHE cosmic rays originate in single-particle decays. The lifetime problem can be illustrated in the decays of string or Kaluza–Klein dilatons. These particles can be decoupled from fermions [25] and therefore decay through dimension-five couplings to gauge fields, which in leading order are

$$\mathcal{L}_{I\Phi} = -\frac{\Phi}{4f}\, F^{\mu\nu}{}_{a} F_{\mu\nu}{}^{a}. \tag{1}$$

Here, f is a mass scale characterizing the strength of the coupling, and in Kaluza–Klein or string theory, it is of order of the reduced Planck mass $m_{\rm Pl} = (8\pi G_N)^{-1/2} \simeq 2.4 \times 10^{18}$ GeV. In heterotic string theory, *e.g.*, $f = m_{\rm Pl}/\sqrt{2} \simeq 1.7 \times 10^{18}$ GeV [26]. If the number of the vector fields is d_G, we find the lifetime estimated from the dilaton–vector coupling of Eq. (1)

$$\begin{aligned}\tau_\Phi &= \frac{32\pi f^2}{d_G m_\phi^3} \\ &= \frac{1.9\times 10^{-22}\,\mathrm{s}}{d_G}\left(\frac{10^{12}\,\mathrm{GeV}}{m_\phi}\right)^3 \\ &\quad \times \left(\frac{f}{1.7\times 10^{18}\,\mathrm{GeV}}\right)^2 .\end{aligned}$$

If there are direct decay channels through first order couplings, superheavy particles decay extremely fast, even if the coupling is only of gravitational strength and dimensionally suppressed. Superheavy relic particles with a sufficiently long lifetime require sub-gravitational couplings or exponential suppression of the decay mechanism due to wormhole effects [9], instantons [10], or magic from the brane world [27].

Motivated by the attractiveness of the decay scenario, in this paper we investigate the possibility that the observed UHE cosmic rays result from annihilation of relic superheavy dark matter.

Annihilation of dark matter in the halo has been a subject of much interest, with particular emphasis on possible neutralino signatures in the cosmic ray flux. A reasonable annihilation scenario for UHE cosmic rays is the production of two jets each of energy M_X, which then fragment into a (very) many-particle final state, including leading particles of mass comparable to M_X.

Assuming that the relic superheavy dark matter follows a Navarro–Frenk–White (NFW) density profile [28], in Sec. 2 we calculate the expected spectrum and the annihilation cross section required to account for the observed super-GZK events. We find that the necessary cross section exceeds the unitarity bound. We then discuss possibilities for circumventing the bound.

In Sec. 3 we consider the contribution from annihilation in sub-galactic clumps of dark matter. We find that this result is very sensitive to the assumed density profile of the subclumps.

In Sec. 4 we discuss the expected anisotropy in arrival direction of UHE cosmic rays if the source is superheavy relic particle annihilation.

Sec. 5 contains our conclusions.

2 Annihilation in the smooth component

In calculating the UHE cosmic ray flux from a smooth[4] superheavy dark matter distribution in the galactic halo, we assume a superheavy X-particle halo density spherically symmetric about the galactic

[4] We refer to a superheavy dark matter distribution as *smooth* if it can be described by a particle density $n_X(\mathbf{d})$ which decreases uniformly with distance from the galactic center.

center, $n_X(\mathbf{d}) = n_X(d)$, where d is the distance from the galactic center. We will assume $n_X(d)$ is given by a NFW profile [28]

$$n_X(d) = \frac{N_0}{d(d+d_s)^2}. \tag{2}$$

Navarro's estimate for the fiducial radius d_s for the Milky Way is of order 25 kpc [29]. Dehnen and Binney have examined a flattened NFW profile as a special case of a whole class of halo models and give a value $d_s = 21.8$ kpc [30]. We will use $d_s = 3d_\odot = 24$ kpc in our numerical estimates, where $d_\odot \simeq 8$ kpc is the distance of the solar system from the galactic center.

The dimensionless parameter N_0 may be found by requiring that the total mass of the Galaxy is $2\times 10^{12} M_\odot$, which gives

$$N_0 = \frac{8\times 10^{55}}{M_{12}} = \frac{2.8\times 10^{-9}}{M_{12}} \frac{\mathrm{kpc}^3}{\mathrm{cm}^3},$$

where $M_{12} = M_X/10^{12}$ GeV.

WIMPZILLA annihilation produces two jets, each of energy M_X, while decay of a WIMPZILLA produces two jets, each of energy $M_X/2$. The energy spectrum of observed UHE cosmic ray events from annihilation is

$$\mathcal{F} = 2\, \frac{d\mathcal{N}(E, E_{\rm jet} = M_X)}{dE} \langle \sigma_A v\rangle \times \int d^3\mathbf{d}\, \frac{n_X^2(d)}{4\pi\, |\mathbf{d}-\mathbf{d}_\odot|^2}. \tag{3}$$

Here, $d\mathcal{N}(E, E_{\rm jet})/dE$ is the fragmentation spectrum resulting from a jet of energy $E_{\rm jet}$. For comparison, the energy spectrum of observed UHE cosmic ray events from WIMPZILLA decay is

$$\mathcal{F} = 2\, \frac{d\mathcal{N}(E, E_{\rm jet} = M_X/2)}{dE} \Gamma_X \times \int d^3\mathbf{d}\, \frac{n_X(d)}{4\pi\, |\mathbf{d}-\mathbf{d}_\odot|^2},$$

where Γ_X is the decay width of the WIMPZILLA.

There are many discussions about the extrapolations and approximations to the fragmentation function $d\mathcal{N}(E, E_{\rm jet} = M_X)/dE$ (for a recent review see [31] and for a numerical approach, see [32]). In the results of Fig. 1, we use the fragmentation functions that are the MLLA limiting spectrum of ordinary QCD [33]. The salient features of the results can be understood by making the simple approximation that most of the content of the jet is pions, with a spectrum in terms of the usual variable $x \equiv E/E_{\rm jet}$ (of course, $0 \le x \le 1$),

$$\begin{aligned} \frac{d\mathcal{N}(x)}{dx} &= \frac{15}{16} x^{-3/2}(1-x)^2 \\ &\sim \frac{15}{16} x^{-3/2} \quad (x \ll 1). \end{aligned}$$

The UHE cosmic rays in this picture are photons resulting from the decay of neutral pions. The neutral pions are about one-third of the total number of pions in the jet, and each π^0 decay produces two pions.

Using this fragmentation approximation, the scaling of the flux with M_X can be found to be

$$\begin{aligned} \mathcal{F} &\propto M_X^{1/2} \langle \sigma_A\rangle M_X^{-2} \quad \text{(annihilation)} \\ &\propto M_X^{1/2} \Gamma_X M_X^{-1} \quad \text{(decay)}. \end{aligned} \tag{4}$$

The factor of $M_X^{1/2}$ is from the fragmentation function, and the factors of M_X^{-2} or M_X^{-1} arise from n_X^2 or n_X, respectively. Therefore, for a given $\mathcal{F}$, the necessary cross section scales as $M_X^{3/2}$ in the annihilation scenario and the necessary decay width scales as $M_X^{1/2}$ in the decay scenario.

Calculating the resulting UHE cosmic ray flux in the annihilation model, and comparing it to the similar calculation in the decay scenario, we obtain the results shown in Fig. 1.

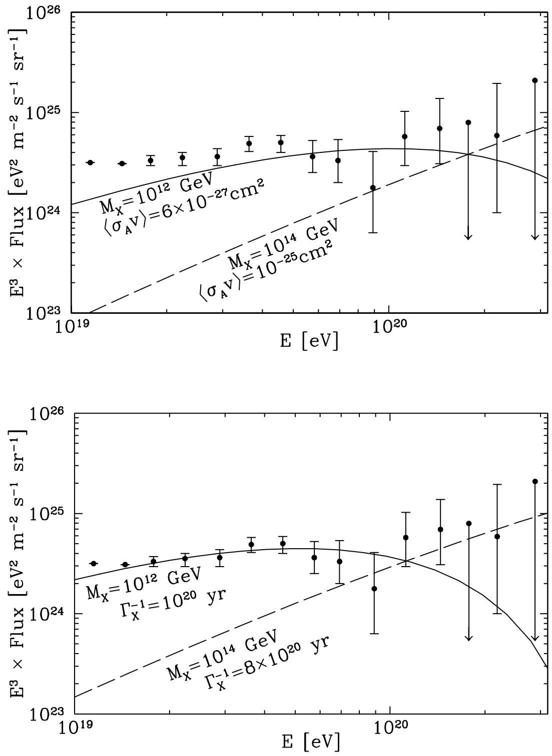

Fig. 1. UHE cosmic ray spectra from superheavy particle annihilation (upper panel) or decay (lower panel). For both figures the solid lines are for $M_X = 10^{12}\,\text{GeV}$ and the dashed lines are for $M_X = 10^{14}\,\text{GeV}$. For annihilation, the solid curve is for $\langle\sigma_A v\rangle = 6\times10^{-27}\,\text{cm}^2$ and the dashed curve corresponds to $\langle\sigma_A v\rangle = 10^{-25}\,\text{cm}^2$. In the decay case, the solid curve is for $\Gamma_X = 10^{20}\,\text{yr}$ and the dashed curve is for $\Gamma_X = 8\times10^{20}\,\text{yr}$.

The shape of the spectrum is determined by the mass of the WIMPZILLA and the overall normalization can be scaled by adjusting $\langle\sigma_A v\rangle$ or Γ_X. Clearly, in order to produce UHE cosmic rays in excess of $10^{20}\,\text{eV}$, M_X cannot be too much smaller than $10^{12}\,\text{GeV}$. In order to provide enough events to explain the observed UHE cosmic rays, $\langle\sigma_A v\rangle$ has to be in the range $10^{-25}\,\text{cm}^2$ to $10^{-27}\,\text{cm}^2$.

Similarly, in the decay scenario, the WIMPZILLA lifetime must be in the range $10^{20}\,\text{yr}$ to $10^{21}\,\text{yr}$.

We have already discussed the fact that a lifetime in this range for a particle so massive is rather unexpected (but, of course, not impossible). The required cross section in the annihilation scenario is also orders of magnitude larger than expected.

The necessary annihilation cross section is well in excess of the unitarity bound to the l-wave reaction cross section [34–36]:

$$\sigma_l v \leq \frac{4\pi}{M^2 v}(2l+1) \simeq \frac{1.5\times10^{-47}}{M_{12}^2} \times(2l+1)\frac{100\,\text{km s}^{-1}}{v}\ \text{cm}^2.$$

However, as emphasized by Hui [36], there are several ways to evade the bound. The most probable evasion in our scenario is possible finite size effects of the WIMPZILLA. The relevant limit in that case is $\sigma_A \lesssim 64\pi R^2$, where R is the "size" of either the particle or the range of the interaction. For the type of cross sections discussed above, we would require

$$R \gtrsim 4\times10^{-13}\left(\frac{\langle\sigma_A v\rangle}{10^{-26}\,\text{cm}^2}\right)^{1/2} \times\left(\frac{100\,\text{km s}^{-1}}{v}\right)^{1/2}\ \text{cm}.$$

This size is still uncomfortably large, but is of order of the length scale of the strong interactions. In the next section we will show that the signal from a clumped component of dark matter may dominate and the necessary cross section may be smaller and the requirements on R not as severe.

3 Annihilation in the clumped component

So far we have assumed that the galactic dark matter is smoothly distributed. In this section we will consider the contribution from inhomogeneities in the galactic distribution. We will first consider the signal from the sort of sub-galactic clumps (subclumps) predicted by numerical simulations. Then we will consider the possibility that we live within the core radius of such a subclump.

3.1 Individual Subclumps

We have modeled the smooth component of the dark matter by a NFW profile, Eq. (2). In addition to this smooth component, the dark matter may have a clumped component, as suggested by N-body simulations. Using the results of these simulations, the number of subclumps of mass M_{cl} per unit volume at distance d from the galactic center can be written in the form:

$$n_{cl}(d, M_{cl}) = n_{cl}^0 \left(\frac{M_{cl}}{M_H}\right)^{-\alpha} \times \left[1 + \left(\frac{d}{R_c^{cl}}\right)^2\right]^{-3/2},$$

where R_c^{cl} is the core radius of the subclump distribution in the galaxy, typically of order 10 to 20 kpc. The mass of the halo of the Galaxy is M_H. The normalization constant n_{cl}^0 can be calculated by requiring that the mass in the clumped component is a fraction ξ of the halo mass M_H. From simulations, $\xi \sim 10\%$. The power index, α, may also be found from simulations, with the result $\alpha \sim 1.9$. It is then easy to find

$$n_{cl}^0 \approx \frac{(2-\alpha)\xi}{4\pi (R_c^{cl})^2 R_H M_H \eta^{2-\alpha}},$$

where η is the ratio of the mass of the largest subclump to the halo mass, $\eta = M_{cl}(max)/M_H \sim 0.01 - 0.1$, and R_H is the radius of the halo.

Up to this point, all the discussion is independent of the density profile of the individual subclumps. Let us now discuss the effect of the density profile of individual subclumps.

3.1.1 Isothermal subclumps

First, we assume that the density profile of an individual subclump is an isothermal sphere, with radial profile

$$n_X(r) = n_H(d) \left(\frac{r}{r_0}\right)^{-2},$$

where r_0 is the radius of the individual subclump, defined as the radius at which the subclump density equals the density of the background halo at the distance d from the galactic center. If the subclump mass is M_{cl}, we can write

$$M_{cl} = 4\pi M_X n_H(d) \int_0^{r_0} dr\, r^2 (r/r_0)^{-2},$$

from which we obtain

$$r_0 = \left[\frac{M_{cl}}{4\pi M_X n_H(d)}\right]^{1/3}.$$

The rate of annihilation in an individual subclump at distance d from the galactic center is

$$\begin{aligned}\mathcal{R}(M_{cl}, d) = \langle\sigma_A v\rangle 4\pi \int_{R_{min}}^{r_0} dr\, r^2 \\ \times n_H^2(d)(r/r_0)^{-4} \\ + \langle\sigma_A v\rangle (4\pi/3) n_H^2(d) r_0^4 / R_{min},\end{aligned} \tag{5}$$

where R_{min} is the minimum radius of a subclump (the inner radius where the profile is flattened by efficient annihilations).

A reasonable upper limit to R_{min} can be obtained by requiring that the dynamical free-fall time of dark matter from the edge of the subclump is comparable with the annihilation rate at R_{min}. The dynamical free-fall time is

$$\tau_{ff} = r_0/v(r_0) = [4\pi G M_X n_H(d)]^{-1/2},$$

while the annihilation time scale is

$$\tau_A = [n(r)\langle\sigma_A v\rangle]^{-1} = \frac{1}{n_H(d)\langle\sigma_A v\rangle}\left(\frac{r}{r_0}\right)^2.$$

It is then easy to obtain the ratio

$$\frac{R_{min}}{r_0} = \langle\sigma_A v\rangle^{1/2}\left[\frac{n_H(d)}{4\pi G M_X}\right]^{1/4}. \qquad (6)$$

From Eqs. (6) and (6), we then obtain for the rate of annihilations in a subclump of mass M_{cl} at distance d from the galactic center:

$$\mathcal{R}(M_{cl}, d) = \frac{4}{3} M_{cl}\langle\sigma_A v\rangle^{1/2}(4\pi G)^{1/4} \times \left(\frac{n_H(d)}{M_X}\right)^{3/4}.$$

Now we can calculate the contribution of the clumped component in a similar manner as Eq. (3):

$$\mathcal{F} = 2\frac{d\mathcal{N}(E, E_{\rm jet} = M_X)}{dE} \times \int d^3\mathbf{d}\, \frac{1}{4\pi\,|\mathbf{d} - \mathbf{d}_\odot|^2} \times \int_{M_{min}}^{M_{max}} dM_{cl}\; n_{cl}(d, M_{cl})\mathcal{R}(M_{cl}, d). \qquad (7)$$

As before we will use an NFW profile for $n_H(d)$ with $d_s = 24\,\mathrm{kpc}$.

To understand the scaling of the events from subclumps, we can make the simple approximation that one is at the galactic center ($|\mathbf{d}_\odot| = 0$). In the limit that the contributions come from $d \ll R_c^{cl}$ and $d \ll d_s$, we find

$$\mathcal{F} \propto \langle\sigma_A v\rangle^{1/2} M_X^{-1}$$

Now turning to the location of the solar system in the Milky Way, using $R_c^{cl} = 15\,\mathrm{kpc}$, $\eta = 0.1$, and $\xi = 0.1$, we find

$$\frac{\mathcal{F}_{\rm clumped}}{\mathcal{F}_{\rm smooth}} \simeq 2\times 10^6 \left(\frac{M_X}{10^{12}\mathrm{GeV}}\right)^{1/2} \times \left(\frac{3\times 10^{-26}\,\mathrm{cm}^2}{\langle\sigma_A v\rangle}\right)^{1/2}.$$

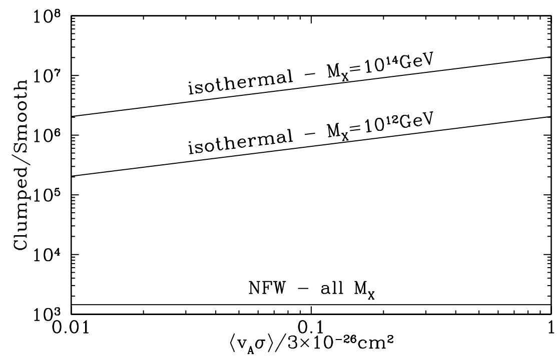

Fig. 2. The ratio of events from the subclump component to the smooth component assuming either an isothermal or a NFW profile for the subclumps.

The ratio of the fluxes from the subclump component and the smooth component are given in Fig. 2. Comparing Fig. 1 and Fig. 2, we see that for $M_X = 10^{12}$ GeV, to normalize the flux to the observed events would require a cross section 10^3 times smaller than the isotropic case if the WIMPZILLA subclumps have an isothermal profile.

3.1.2 NFW subclumps

Now, we take the density profile of the subclumps according to a NFW profile,

$$n_X(r) = n_H(d)\,\frac{r_0\,[r_0 + r_s]^2}{r\,[r + r_s]^2}, \qquad (8)$$

where r_s is the fiducial radius of the subclump and again r_0 is the radius of the subclump where $n_X(r) = n_H(d)$. We assume here that the ratio $r_0/r_s \gg 1$ is constant for all subclumps. The mass of the subclump can be written in terms of r_0 and r_s as

$$
\begin{aligned}
M_{cl} &= 4\pi M_X \int_0^{r_0} dr\, r^2 n_X(r) \\
&\simeq 4\pi M_X n_H(d) r_0^3 \ln\left(\frac{r_0}{r_s}\right).
\end{aligned}
$$

which provides the core radius r_0 for the subclump once we know r_s.

The rate of annihilations for one subclump is given by

$$
\begin{aligned}
\mathcal{R}(M_{cl}, d) = 4\pi\langle\sigma_A v\rangle \int_{R_{min}}^{r_0} dr\, r^2 n_X^2(r) \\
+(4\pi/3)R_{min}^3\langle\sigma_A v\rangle n_H^2(R_{min}),
\end{aligned}
$$

which becomes

$$
\mathcal{R}(M_{cl}, d) \simeq \frac{4\pi}{3}\langle\sigma_A v\rangle \frac{r_0^6}{r_s^3} n_H^2(d).
$$

To calculate $\mathcal{F}$ for NFW subclumps, we follow a procedure similar to Eq. 7, but with $\mathcal{R}(M_{cl}, d)$ given above. Again expressing the flux in terms of the flux from the smooth component, we find

$$
\frac{\mathcal{F}_{\text{clumped}}}{\mathcal{F}_{\text{smooth}}} \simeq 2 \times 10^3,
$$

independent of M_X or $\langle\sigma_A v\rangle$. Again, this result is represented in Fig. 2.

3.2 If we live in a subclump

Now consider the possibility that we live within the dense core of a subclump. The probability to be in the center of a subclump is given by the probability of being in a subclump, multiplied by the probability to be within the minimum radius, so $P_{\text{core}} \sim 0.1(R_{min}/r_0)^3$. For an isothermal subclump with mass $10^8 M_\odot$, and for $\langle\sigma_A v\rangle = 3 \times 10^{-27}\,\text{cm}^2$, we have $r_0 \sim 1\text{kpc}$ and $R_{min} \sim 10^{15}\text{cm}$, so that $P_{\text{core}} \sim 10^{-21}$. The probability is quite infinitesimal, but the possibility has interesting consequences.

The cases to consider are again that of an isothermal subclump and a NFW subclump, with the Earth in the core. In both cases it is relevant to calculate the value of the minimum radius, R_{min}, within which the efficiency of annihilations flattens the radial profile of dark matter. The reason for this is that the main contribution to the flux of UHE events comes from this region if the detector is within the core.

For an isothermal subclump,the flux is proportional to

$$
\begin{aligned}
\mathcal{F} = \frac{4}{3} n_H(d_\odot)^{11/12}\langle\sigma_A v\rangle^{-1/2} \\
\times M_{cl}^{1/3}(4\pi M_X)^{5/12} G^{3/4},
\end{aligned}
$$

where M_{cl} is the subclump mass and $n_H(d_\odot)$ is the typical density at the distance of the solar system from the galactic center. This flux should be compared with the total flux from a distribution of isothermal subclumps in the halo, given by Eq. (7).

One has to calculate R_{min} in the NFW case as well. As before, the way to do this is to require that the free-fall time scale equals the time scale for annihilation at R_{min}. We assume that the external radius of the subclump and its core radius,r_s, are related by $r_0/r_s = \delta \gg 1$. Then we can write:

$$
\begin{aligned}
R_{min} = \delta^{9/2}(GM_{cl})^{-1/2} r_s^{5/2} \\
\times n_H(d_\odot)\langle\sigma_A v\rangle.
\end{aligned}
$$

Note also that r_s is related to the total subclump mass by

$$r_s = \left[\frac{M_{cl}}{4\pi n_H(d_\odot)\delta^3 M_X f(\delta)} \right]^{1/3},$$

where $f(\delta) = \ln(1+\delta) - \delta/(1+\delta)$. The flux in this case is proportional to

$$\mathcal{F} \propto \left(\frac{R_{min}}{r_s} \right)^{-2} R_{min} \langle \sigma_A v \rangle n_{H,0}^2.$$

Here, $n_{H,0} \approx \delta^3 n_H(d_\odot)$. This number should be compared with the total flux from the distribution of subclumps with NFW profiles in the Galaxy.

Note that the fluxes are proportional to $\langle \sigma_A v \rangle^{-1/2}$ ($\langle \sigma_A v \rangle^{-1}$) for the isothermal (NFW) case. This means that decreasing the cross section actually helps increasing the flux. This is true provided the minimum radius, R_{min}, reaches the radius of (say) the atmosphere. This imposes a lower limit on the cross section, which is still above the unitarity limit.

4 Predicted signals

The annihilation rate is very sensitive to the local density. This implies that if the UHE cosmic rays originate from dark matter annihilation and dark matter in our galaxy is clumped, the arrival direction of UHE cosmic rays should reflect the dark matter distribution.

The first possibility we considered is that the dark matter distribution is smooth, and follows a NFW profile. In this case, the galactic center should be quite prominent. In Fig. 3 we illustrate the expected angular dependence of the arrival direction. The galactic center is prominent for the decay scenario where $\mathcal{F} \propto n_X$ [12], and even more prominent for the annihilation case where $\mathcal{F} \propto n_X^2$.

Now if we assume that there are dark matter subclumps within the galaxy, then the density in the subclumps would be larger than the ambient background dark-matter density, and events originating from subclumps will dominate the observed signal.

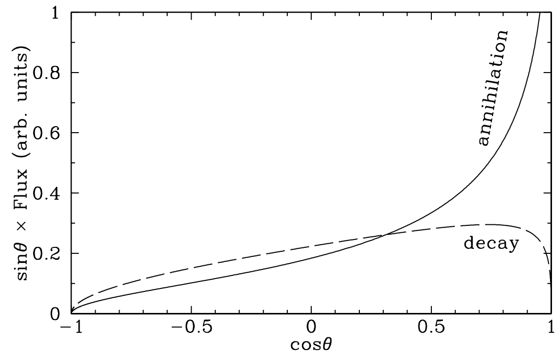

Fig. 3. The angular dependence of events originating from annihilation or decay of the smooth component of dark matter that follows a NFW profile. In this figure, $\cos\theta = 1$ corresponds to the direction of the galactic center.

The dominance of events from subclumps has two effects. The first effect is that a smaller annihilation cross section is required to account for the observed UHE flux. The second effect is that there will be a very large probability of detecting a nearby subclump.

In Fig. 4 we present histograms showing the number of occurrences of single and multiple events in arrival directions within a square degree. We generated many realizations of the expected subclump distribution, and calculated the flux from the subclumps that would result in about 100 detected events. We see that the *average* probabilities are quite reasonable, with most events in square degree areas with only one event, and a few pairs and triplets of events. However if we examine individual realizations containing about 100 events, we see that there is a large probability of a large number of events from

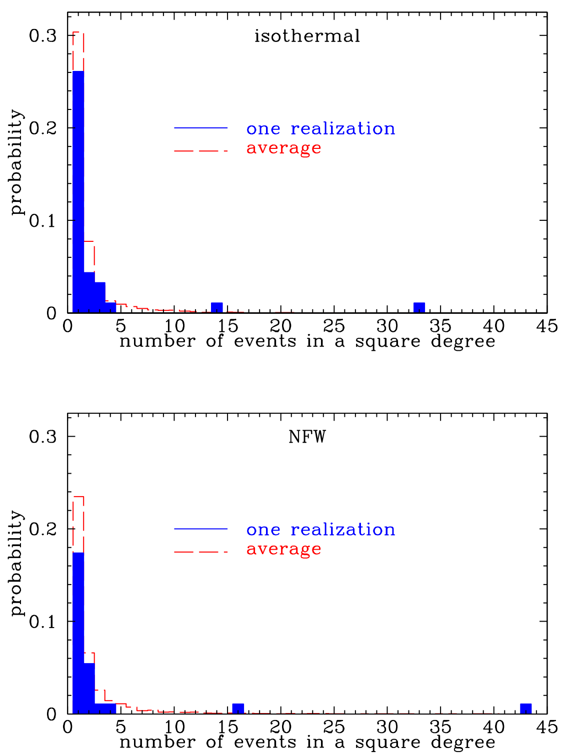

Fig. 4. The probability of finding a certain number of events in a one square degree area of the sky. The upper figure is for isothermal subclumps and the lower figure for NFW subclumps

single nearby subclumps. For instance, in the single realizations shown in Fig. 4, the isothermal subclumps result in a square degree bin with 14 events and a square degree bin with 33 events. In the typical NFW distribution, there is one bin with 16 events and another bin with 43 events. Therefore, if more than 100 events are observed with full sky coverage, the signature of subclumps will be unmistakable.

Finally, we remark on the characteristics of cosmic ray events originating from annihilation (or decay) of very nearby WIMPZILLAS, as might be expected if we live within the core radius of a subclump.

The process of jet production after the WIMPZILLA annihilation or decay is extremely complicated, but some general features can be understood as assuming that each elementary process of quark-antiquark generation or gluon radiation implies an average angular widening of the initial beam by the amount $\delta\alpha \sim \Lambda_{QCD}/E_{\rm jet}$, where $\Lambda_{QCD} \sim 0.3$ GeV and $E_{\rm jet} \sim M_X$. While the jets further fragments into other quarks and gluons, its opening angle is expected to make a random walk for a number of steps approximately given by $E_{\rm jet}/\Lambda_{QCD}$. Therefore, the ultimate opening angle should be $\alpha_{\rm jet} \sim (\Lambda_{QCD}/E_{\rm jet})^{1/2}$. We stress that this has to be intended as a rough estimate of the quantity $\alpha_{\rm jet}$, whose real value is determined by a variety of elementary particle physics processes.

If we adopt the expression obtained above for $\alpha_{\rm jet}$, the size of the hadronic shower at distance d from the production region can be estimated as:

$$l = \alpha_{\rm jet} d \sim 5 \times 10^8 M_{12}^{-1/2} d_{15}\,{\rm cm},$$

where $d_{15} = d/10^{15}$ cm. For simplicity, let us assume that the Earth is sitting in the center of a clump with an isothermal profile. In this case, most of the contribution to the local flux of UHE cosmic rays, as described in the previouse section, is provided by particles annihilating within R_{min}. From Eq. (6) we obtain for the minimum radius

$$R_{min} \approx 7 \times 10^{14} \frac{\langle \sigma_A v \rangle}{10^{-26}\,{\rm cm}^2} \times M_{12}^{-1/2} \left(\frac{M_{cl}}{10^8 M_\odot}\right)^{1/3} {\rm cm},$$

where $M_{cl,8}$ is the clump mass in units of $10^8 M_\odot$. Substituting into the expression for l we obtain

$$l < 3.5 \times 10^8 \frac{\langle \sigma_A v \rangle}{10^{-26}\,{\rm cm}^2}$$

$$\times M_{12}^{-1}\left(\frac{M_{cl}}{10^8 M_\odot}\right)^{1/3}\ \text{cm},$$

so that the size of the jets observed at the Earth can easily be smaller than the size of the atmosphere. This should result in the appearence of *anomalous* jets with many UHE cosmic rays arriving more or less at the same time, with a time spread of order $\Delta t \sim d\alpha_{\rm jet}/c$.

5 Conclusion

The origin of UHE cosmic rays from annihilation (or decay) of WIMPZILLAS has the attractive feature of the simplicity of generating ultrahigh energies: simple conversion of rest mass energy.

The decay scenario requires a lifetime of about 10^{20} years. This has proven to be a model building challenge. The annihilation scenario requires an uncomfortably large annihilation cross section, much larger than the unitarity limit. While the unitarity limit is a useful guideline, it is not sacrosanct.

In this paper we have examined the annihilation scenario. This proposal results in several striking predictions. While not discussed in this paper, like the decay scenario, the annihilation scenario suggests that the bulk of the UHE events are photons. The true characteristic of the annihilation scenario is the expected anisotropy in arrival direction. If the dark matter is smoothly distributed in the galaxy, the galactic center should be prominent. If the dark matter is clumped on sub-galactic scales, then the subclumps should be visible.

Thus, the annihilation scenario can be falsified by complete sky coverage. The Pierre Auger Observatory will be able to see the galactic center. By covering the southern sky, complete sky coverage should be able to pick out subclumps if they are present.

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