Extremely high energy cosmic rays from relic particle decays

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Abstract

The expected proton and neutrino fluxes from decays of massive metastable relic particles are calculated using the HERWIG QCD event generator. The predicted proton spectrum can account for the observed flux of extremely high energy cosmic rays beyond the Greisen-Zatsepin-Kuzmin cutoff, for a decaying particle mass of $\mathcal{O}(10^{12})$ GeV. The lifetime required is of $\mathcal{O}(10^{20})$ yr if such particles constitute all of the dark matter (with a proportionally shorter lifetime for a smaller contribution). Such values are plausible if the metastable particles are hadron-like bound states from the hidden sector of supersymmetry breaking which decay through non-renormalizable interactions. The expected ratio of the proton to neutrino flux is given as a diagnostic of the decaying particle model for the forthcoming Pierre Auger Project.

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I. INTRODUCTION

It has been known for some time that interactions on the 2.73 K blackbody cosmic microwave background (CMB) will severely degrade the energies of cosmic ray nucleons with energies beyond \( \sim 5 \times 10^{19} \text{ eV} \) — the Greisen-Zatsepin-Kuzmin (GZK) cutoff \([1]\). It was therefore very surprising when the Fly’s Eye atmospheric fluorescence detector reported the observation of an extremely high energy cosmic ray (EHECR) event with an energy of \((3.0 \pm 0.9) \times 10^{20} \text{ eV} \) \([2]\). This was followed by the detection of a \((1.7 - 2.6) \times 10^{20} \text{ eV} \) event by the AGASA air shower array \([3]\). These discoveries substantiated earlier claims from the Volcano Ranch \([4]\), Haverah Park \([5]\) and Yakutsk \([6]\) air shower arrays that cosmic rays do exist beyond the GZK cutoff. About a dozen such events are now known. Detailed accounts of the data may be found in recent reviews \([7]\).

In Figure 1 we show the EHECR spectrum for energies exceeding \(10^{18} \text{ eV} \) \([8]\); note that the fluxes have been multiplied by \(E^3\). It is believed that cosmic rays with energies up to \(\sim 5 \times 10^{18} \text{ eV} \), the so-called ‘ankle’, are predominantly of galactic origin, possibly accelerated by the Fermi mechanism in supernova remnants \([9]\). Above this energy, the spectrum flattens and the composition changes from being mostly heavy nuclei to mostly protons. Such a correlated change in the spectrum and composition was first established by the Fly’s Eye experiment \([2]\) and Figure 1 shows their suggested two-component fit to the data. The new component which dominates at energies beyond \(\sim 5 \times 10^{18} \text{ eV} \) is isotropic and therefore cannot possibly originate in the galactic disk \([10,11]\). However it also extends well beyond the GZK cutoff raising serious problems for hypothetical extragalactic sources. Because of the rapid energy degradation at these energies through photopion production on the CMB, such sources must exist within \(\sim 500 \text{ Mpc} \), in fact within \(\sim 50 \text{ Mpc} \) for the highest energy Fly’s Eye event \([12]\). For heavy nuclei, the energy loss is less severe according to a revised calculation \([13]\) so the range may extend up to \(\sim 100 \text{ Mpc} \). General arguments \([14,15]\) provide correlated constraints on the magnetic field strength and spatial extent of the region necessary to accelerate particles to such high energies and these requirements are barely met by likely astrophysical sites such as active galactic nuclei and the ‘hot spots’ of radio galaxies \([16]\). Moreover there are few such sources close to us and no definite correlations have been found between their locations and the arrival directions of the most energetic events \([17,10]\). It has been speculated that gamma-ray bursts which too are isotropically distributed, may be responsible for EHECRs \([18]\). However since these are at cosmological distances, one would expect to see the GZK cutoff in the cosmic ray spectrum contrary to observations (cf. ref. \([19]\)).

Some of the above arguments may be evaded if the EHECR events are due not to nucleons but neutral particles such as photons and neutrinos. Although high energy photons also suffer energy losses in traversing the CMB and the extragalactic radio background, there is no threshold effect which would cause a cutoff near the GZK value \([20]\). However the observed shower profile of the highest energy Fly’s Eye event \([2]\) argues against the primary being a photon since it would have interacted on the geomagnetic field and started cascading well before entering the atmosphere \([21]\). The observed events are also unlikely to be initiated by neutrinos as they all have incident angles of less than 40° from the zenith and thus too small a path length in the atmosphere for interactions \([22]\). This argument may be evaded if neutrinos become strongly interacting at high energies due to new physics beyond the Standard Model \([23,24]\), but such proposals are found not to be phenomenologically viable \([25]\) (although this is disputed \([26]\)). (Alternatively, the propagating high energy neutrinos could annihilate on the relic cosmic neutrino background, assumed to have a small mass of \(\mathcal{O}(0.1) \text{ eV} \), to make hadronic jets within the GZK zone \([27]\).)
Other exotic possibilities have been suggested, e.g. monopoles \[28\], stable supersymmetric hadrons \[29\] and loops of superconducting cosmic string (‘vortons’) \[30\]. However these possibilities have many phenomenological problems \[31,32\] and we do not discuss them further.

Thus one is encouraged to seek ‘top-down’ explanations for EHECRs in which they originate from the decay of massive particles, rather than being accelerated up from low energies. The most discussed models in this connection are based on the annihilation or collapse of topological defects such as cosmic strings or monopoles formed in the early universe \[33–36\]. When topological defects are destroyed their energy is released as massive gauge and Higgs bosons which are expected to have masses of $O(10^{16})$ GeV if such defects have formed at a GUT-symmetry breaking phase transition. The decays of such particles can generate cascades of high energy nucleons, $\gamma$-rays and neutrinos.

A more recent suggestion is that EHECRs arise from the decays of metastable particles with masses $m_X \sim 10^{13} - 10^{16}$ GeV which constitute a fraction of the dark matter \[37\]. These authors suggest that such particles can be produced during reheating following inflation or through the decay of hybrid topological defects such as monopoles connected by strings, or walls bounded by strings. The required metastability of the particle is ensured by an unspecified discrete symmetry which is violated by quantum gravity (wormhole) effects. Another suggestion is that the long lifetime is due to non-perturbative instanton effects \[38\]. In ref. \[39\], a candidate metastable particle is identified in a $SU(15)$ GUT.

A generic feature of these ‘top-down’ models is that the EHECR spectrum resulting from the decay cascade is essentially determined by particle physics considerations. Of course the subsequent propagation effects have astrophysical uncertainties but since the decays must occur relatively locally in order to evade the GZK cutoff \[37\], they are relatively unimportant. Thus although the proposal is speculative, it is possible, in principle, to make reliable calculations to confront with data. In this work we consider another possible candidate for a relic metastable massive particle \[40\] whose decays can give rise to the observed highest energy cosmic rays. First we discuss (§II) why this candidate, which arises from the hidden sector of supersymmetry breaking, is perhaps physically better motivated than the other suggested relics. We then undertake (§III) a detailed calculation of the decay cascade using a Monte Carlo event generator to simulate non-perturbative QCD effects. This allows us to obtain a more reliable estimate of the cosmic ray spectrum than has been possible in earlier work on both topological defect models \[34\] and a decaying particle model \[37\]. We confront our results with observations and identify the mass and abundance/lifetime required to fit the data. We conclude (§IV) with a summary of experimental tests of the decaying particle hypothesis.

II. MASSIVE, METASTABLE DARK MATTER FROM THE HIDDEN SECTOR

Soon after the discovery of the anomaly-free heterotic superstring theory in ten dimensions based on the gauge group $E_8 \times E_8$, it was pointed out \[41\] that in the physical low energy theory (where a grand unified $E_6$ or $O(10)$ group is broken by Wilson lines), the minimum value of magnetic charge is not the Dirac quantum $2\pi/e$ but an integral multiple thereof. Conversely, the
minimum electric charge is smaller than the electron charge $e$ by the same ratio. This was found to be a generic feature of all superstring models based on a level-one Kač-Moody algebra $^{[12]}$. In view of the severe experimental upper bounds on the relic abundance of fractional charges $^{[43]}$, this posed a potential embarrassment for superstring phenomenology $^{[44]}$.

A simple solution to the problem of fractional charges (with an obvious historical analogue in quarks and QCD) is to confine them and it was shown that this can be done in the hidden sector of supersymmetry breaking in the framework of the $SU(5) \otimes U(1)$ unification model $^{[45]}$. In this model, all fractionally charged states have charges $|Q_{em}| = \frac{1}{2}$ and are placed in $4$ or $6$ representations of a hidden $SU(4)$ gauge group which becomes strong at a scale $\Lambda_4 \sim 10^{12}$ GeV and in $10$ representations of a hidden $SO(10)$ group which becomes strong at a scale $\Lambda_{10} \sim 10^{15}$ GeV. This results in integer-charged particles — ‘cryptons’ — which may be 2-constituent mesons, 3-constituent baryons or 4-constituent ‘tetrons’ $^{[46]}$. Some of these mesons could be light (in analogy to the pion of QCD) but most of the states should be heavy with masses of order the confinement scale $\Lambda$. (Other possibilities for stable superstring relics have been discussed in ref. $^{[47]}$.)

The constituent fields have very few renormalizable ($N = 3$) superpotential interactions, so most of these states can only decay via higher-order ($N \geq 4$) superpotential terms. Generically, crypton lifetimes are expected to be $^{[48]}$

$$
\tau_X \simeq \frac{1}{m_X} \left( \frac{M}{m_X} \right)^{2(N-3)}, \quad m_X \sim \Lambda, \tag{1}
$$

where, $M \equiv M_P/\sqrt{8\pi} \simeq 2.4 \times 10^{18}$ GeV is the normalized Planck scale, giving

$$
\tau_4 \sim 10^{(12N_4-80)} \text{yr}, \quad \tau_{10} \sim 10^{(6N_{10}-65)} \text{yr}, \tag{2}
$$

for $SU(4)$ and $SO(10)$ bound states respectively. Thus $\tau_4 \gtrsim 1 \text{ sec} \,(10^{16} \text{ yr})$ for $N_4 \geq 6$ (8) and $\tau_{10} \gtrsim 1 \text{ sec} \,(10^{16} \text{ yr})$ for $N_{10} \geq 10$ (14). Detailed studies of the possible effects of decays of relic cryptons on primordial nucleosynthesis and the CMB spectrum $^{[48]}$, as well as on the diffuse $\gamma$-ray background $^{[48,49]}$ have established that such particles, if they survive as relics of the Big Bang, must either decay well before nucleosynthesis or have lifetimes longer than the age of the universe ($t_0 \sim 10^{10} \text{ yr}$). In the latter case, if such particles make an interesting contribution to the dark matter, their lifetime is further required to exceed $\sim 10^{16} \text{ yr}$ in order to respect experimental bounds on the flux of high energy neutrinos expected from their decays $^{[48,50]}$. It is seen from eq.(2) that these constraints favour $SU(4)$ mesons over their $SO(10)$ counterparts as possible constituents of the dark matter. It is then natural to contemplate the possibility that such cryptons with a mass of $m_4 \sim 10^{12}$ GeV and a lifetime $\tau_4 \gtrsim 10^{16} \text{ yr}$ are also responsible for the observed highest energy cosmic rays.

Recently the above discussion has been extended to other massive metastable particle candidates in superstring/M-theory $^{[40]}$. These authors discuss constructions with higher-level Kač-Moody algebras (necessary to accommodate adjoint Higgs representations in (unified) models other

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1The vacuum state of a physical theory in this scheme must be $M^4 \times K$ where $M^4$ is four-dimensional Minkowski space and $K$ is some compactified six-manifold. Such fractionally charged states exist because $K$ is not simply connected — these are states in which a closed string wraps around a non-contractible loop in $K$. 

4
than $SU(5) \otimes U(1)$) and note that similar metastable bound states occur in such models. They go on to consider other candidate particles in M-theory such as Kaluza-Klein states associated with extra dimensions but find that these are not as attractive, being either too heavy or too unstable. They suggest that although the $SU(5) \times U(1)$ model discussed above was constructed in the weak coupling limit, it may be elevated to an M-theory model in the strong coupling limit. The $SU(4)$ tetrons are then still the most likely candidates for massive metastable dark matter with the modification that the Planck scale $M$ in Eq. (1) may be replaced by a somewhat smaller scale.

The main reason why this possibility was not seriously entertained earlier concerns the expected relic abundance of such massive particles. If cryptoxs were maintained in chemical equilibrium in the early universe through self-annihilations, their present energy density is given by the usual ‘freeze-out’ calculation as inversely proportional to the (velocity-averaged) annihilation cross-section $\langle \sigma_{\text{ann}} v \rangle$. Estimating this to be $\langle \sigma_{\text{ann}} v \rangle \sim m_X^{-2}$ we see that equilibrium would have been established if the annihilation rate exceeded the Hubble expansion rate ($H \sim T^2/M$), i.e. at temperatures

$$T > T_{\text{dec}} \sim \frac{m_X}{\ln(M/m_X)}.$$  

(3)

The relic abundance is then simply estimated as the equilibrium value at decoupling:

$$\Omega_X \sim 10^{14} \left( \frac{m_X}{10^{12} \text{ GeV}} \right)^2.$$  

(4)

This is the basis for the conclusion that no stable relic particle may have a mass in excess of $\sim 10^5$ GeV without ‘overclosing’ the universe, i.e. contributing $\Omega_X > 1$ [51,17]. This does not necessarily apply to cryptoxs since a period of inflation should have diluted their abundance to essentially zero, along with monopoles and other such supermassive relics. If the reheating temperature following inflation is restricted to be $T_R \lesssim 10^9 - 10^{10}$ GeV in order not to produce too many gravitinos [52,53], cryptoxs would not have been generated afterwards.

However it has been recently recognized that in supersymmetric cosmology, there is likely to be a late stage of ‘thermal inflation’ [54] due to symmetry breaking along flat directions at intermediate scales [55,56]. This would adequately dilute the abundance of thermally generated gravitinos following inflation so the bound quoted above on the reheating temperature is no longer valid and the value of $T_R$ may be much higher. In that case cryptoxs even as massive as $10^{12}$ GeV may

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2 This was initially considered to be an ‘entropy crisis’ [55] since it would dilute any baryon asymmetry generated at the GUT scale. However there are now several plausible mechanisms for low temperature baryogenesis [57,58] which may operate after thermal inflation.

3 The vacuum energy $V(\phi)$ of the scalar ‘inflaton’ field is constrained to be $V^{1/4}/\epsilon^{1/4} \simeq 2.7 \times 10^{-2} M$ by the anisotropy in the CMB observed by COBE, where the slope parameter $\epsilon \equiv (M^2/2)(V'/V)^2$ is required to be $\ll 1$ to permit inflation to occur [53]. (The number of e-folds of expansion until the end of inflation is just $N = \int_0^{\phi_{\text{end}}} d\phi/M \sqrt{2\epsilon}$ and this should exceed $\sim 50 - 60$ in order to solve the flatness and homogeneity problems of the standard cosmology.) The reheat temperature $T_R$ can, in principle, have been as high as $\sim V^{1/4}$ although it is usually considerably smaller since the inflaton field is very weakly coupled in most inflationary models.
well have been brought back into thermal equilibrium during reheating after inflation and survived
with the huge relic abundance \( \Omega_X \). However thermal inflation would also have diluted this to an
acceptable level as was noted in ref. \[60\]; to obtain \( \Omega_X \sim 1 \), the number of e-folds of thermal
inflation required is just

\[
N_{\text{therm}} = \frac{1}{3} \ln(10^{14}) \simeq 11. \tag{5}
\]

This fits in well with the expectation that \( N_{\text{therm}} \sim \frac{1}{2} \ln(\Sigma / m_W) \sim 10 - 15 \) for the intermediate
scale \( \Sigma \) in the range \( (10^{-7} - 10^{-2})M \) \[54\]. Of course given the uncertainty in the value of \( N_{\text{therm}} \)
(and indeed the possibility that there may be more than one such epoch), \( \Omega_X \) could well have been
reduced to a negligibly small value.

Another possibility is that massive particles such as cryptons were never in thermal equilibrium
but were created with a cosmologically interesting abundance due to the varying gravitational field
during (primordial) inflation \[61,62\]. A cosmologically interesting relic abundance then arises for
\( m_X \sim (0.04 - 2)H \) where \( H \sim 10^{11} - 10^{14} \) GeV is the likely Hubble parameter during inflation
\[61\]. This is certainly very encouraging but it should be remembered that a later stage of thermal
inflation would dilute such an abundance to a negligible level, as discussed above.

It is clear that the relic abundance of massive particles such as cryptons will necessarily be
very uncertain given our ignorance of the thermal history of the universe prior to nucleosynthesis.
However as the above discussion illustrates, there are two complementary ways in which a
cosmologically interesting abundance may result so we may reasonably consider such particles as
candidates for the dark matter. We now move on to discuss whether relic cryptons can indeed be
the source of the EHECR by determining the expected spectrum of high energy particles from
their decays.

### III. COSMIC RAYS FROM MASSIVE PARTICLE DECAY

To calculate the expected flux of cosmic rays from decays of massive particles such as cryptons,
we must consider the contribution from both decaying particles in the halo of our Galaxy as well as
those elsewhere in the universe. Since such massive particles would behave as cold dark matter and
cluster efficiently in all gravitational potential wells, their abundance in our galactic halo would be
enhanced above their cosmological abundance by a factor

\[
f_{\cos} \equiv n_X^\text{halo} / n_X^\text{cos}. \tag{6}
\]

Note that \( \Omega_X = m_X n_X^\text{cos} / \rho_{\text{crit}} \) where \( \rho_{\text{crit}} \simeq 1.054 \times 10^{-4}h^2 \) GeV cm\(^{-3}\) is the critical density in
terms of the present Hubble parameter \( h \equiv H_0 / 100 \) km sec\(^{-1}\) Mpc\(^{-1}\). If for simplicity we assume
a spherical halo of uniform density,

\[
R_{\text{halo}} \sim 100 \text{ kpc}, \quad \rho_{\text{halo}} \sim 0.3 \text{ GeV cm}^{-3}, \tag{7}
\]

then \( f_{\cos} \sim 3 \times 10^4 h^{-2} \) and the number density of cryptons in the halo is

\[
n_X^\text{halo} \sim 3 \times 10^{-13} \text{ cm}^{-3} \left( \frac{f_{\cos} \Omega_X h^2}{3 \times 10^4} \right) \left( \frac{m_X}{10^{12} \text{ GeV}} \right)^{-1}. \tag{8}
\]
The actual density of dark matter in the halo must of course fall off as \( r^{-2} \) to account for the flat rotation curve of the disk but we do not consider it necessary at this stage to investigate realistic mass models. Thus the universal density of cryptons is smaller than the halo density by about the same numerical factor by which the distance to the horizon (\( \sim 3000h^{-1}\) Mpc) exceeds the halo radius, so the extragalactic contribution to the EHECR flux from decaying cryptons cannot exceed the halo contribution. In particular, the GZK cutoff scale for protons \[12\] or heavy nuclei \[13\] are all much smaller than the horizon distance, so only the halo contribution need be considered, as was emphasized in ref. \[37\]. Only for neutrinos would the extragalactic component be comparable in magnitude \[50\]. Henceforth we restrict ourselves to considering crypton decays in the halo alone.

Now the injection spectrum from particle decay is, to a good approximation,

\[
\Phi_i = \left| \frac{dN_X}{dt} \right| \frac{dN_i}{dE} = \frac{n_{X}^{\text{halo}}}{\tau_X} \frac{2}{m_X} \frac{dN_i}{dx} \quad (9)
\]

for lifetimes longer than the age of the universe (\( \tau_X \gg t_0 \)). Here

\[
x = \frac{E}{E_{\text{jet}}} = \frac{2E}{m_X} \quad (10)
\]

is a measure of particle energy (assuming 2-body decays) and the fragmentation function \( dN_i/dx \) is the average number of particles \( i \) released per decay, per unit interval of \( x \), at the value \( x \). The flux at Earth is then

\[
j_i(E) = \frac{1}{4\pi} R_{\text{halo}} \Phi_i(E). \quad (11)
\]

The final state particles which interest us most are ‘protons’ and neutrinos/antineutrinos where the former includes other nucleons, e.g. antiprotons and neutrons, since they all interact similarly in the Earth’s atmosphere. To compare with observations we multiply the fluxes by \( E^3 \) and define

\[
I_p(E) \equiv j_p(E)E^3 = \frac{1}{4\pi} \frac{n_{X}^{\text{halo}}}{\tau_X} R_{\text{halo}} \frac{2}{m_X} \frac{dN_p}{dx} E^3, \quad (12)
\]

\[
I_\nu(E) \equiv j_\nu(E)E^3 = \frac{1}{4\pi} \frac{n_{X}^{\text{halo}}}{\tau_X} R_{\text{halo}} \frac{2}{m_X} \frac{dN_\nu}{dx} E^3.
\]

For photons and electrons/positrons, propagation energy losses are substantial even within the halo and we do not attempt to determine these. However their injection spectra from particle decay are given by the same computation as for protons and neutrinos, to which we now turn.

\section*{A. Computing the fragmentation functions}

Heavy particles, whether GUT-scale bosons (in topological defect models), cryptons or other hypothetical massive particles, will decay into quarks and leptons. The quarks will hadronize producing jets of mostly pions with a small admixture of nucleons and antinucleons. The neutral pions will decay to give photons while charged pion decays will yield neutrinos and antineutrinos in addition to leptons. Thus the final spectrum of the decay produced particles will be essentially determined by the ‘fragmentation’ of quarks/gluons into hadrons. This is a non-perturbative QCD process and it has not been possible to calculate it by analytic means. Usually phenomenologically
motivated approximations are used to model experimental data on inclusive jet multiplicities and scaling violations [3].

So far, authors of proposals involving heavy particle decay, e.g. in the context of topological defect models [34], have employed a hadronic fragmentation function suggested by Hill [33]

\[
\frac{dN_h^{(Hill)}}{dx} = 0.08 \exp\left[2.6\sqrt{\ln(1/x)}(1-x)^2 \right] \frac{x}{x\sqrt{\ln(1/x)}}. 
\] (13)

It is further assumed that 3% of the hadronic jets from massive relic particle decays turn into nucleons, while the other 97% are pions which decay into photons and neutrinos. This was based on the leading logarithm approximation of QCD [64] applied to experimental data from PETRA on jet production in $e^+e^-$ collisions at tens of GeV. The estimated jet multiplicity from gluon fragmentation was convoluted with the gluon distribution to determine the total hadron yield; to estimate the spectrum, it was assumed that the first moment of the distribution is normalized to unity and the large $x$ behavior was guessed to be $(1-x)^2$ [33]. As we shall see, the Hill fragmentation function (13) significantly overestimates the yield of high $x$ final states from the decay of very massive particles and, moreover, photons and neutrinos are actually produced with a spectrum quite different from that of nucleons. Thus the decay spectra derived using eq.(13) for topological defect models [34] are inaccurate.

Subsequently, another form called the Modified Leading Logarithm Approximation (MLLA) which gives a better description of data at low $x$ has been proposed [63]; a gaussian approximation to this is

\[
\frac{dN_h^{(MLLA)}}{dx} = \frac{K_N}{x} \exp \left[ \ln^2\left(\frac{x}{x_{\text{max}}}\right) \right] \frac{2\sigma^2}{2\sigma^2}, 
\] (14)

where $K_N$ is a constant and

\[
2\sigma^2 = \frac{\pi}{2} \frac{2\pi}{3\alpha_S(s)} \simeq 0.09 \left( \ln \left( \frac{m_X^2}{\Lambda^2} \right) \right)^{3/2}, 
\] (15)

with $x_{\text{max}} = \sqrt{\Lambda/m_X}$ and $\Lambda = 0.234$ GeV. This fragmentation function is employed by the authors of ref. [37] to compute the spectrum from relic particle decays; they determine $K_N$ by requiring that the integral of $x dN_h^{(MLLA)}/dx$ over the range $x \in [0,1]$ be equal to the fraction of the energy transferred to hadrons. However this procedure is not exact as the form (14) is inapplicable for large $x$ and therefore cannot be normalized in this manner. Thus the shape of the cosmic ray spectrum computed [37] by this method for decaying particles is only reliable for small $x$ and its normalization uncertain.

Given the importance of determining the energy spectrum accurately, we decided to improve on these approximate formulations by using the standard tool employed by experimental high energy physicists to study QCD fragmentation, viz. a Monte Carlo event generator. Here the non-perturbative hadronization process is simulated on a computer by a well tested phenomenological model [66]. Although this requires extensive numerical calculations, it is the only means by which successful contact can be made between theory and experimental data. We chose the programme HERWIG [65] (Hadron Emission Reactions With Interfering Gluons) which incorporates the cluster model for hadronization and is based on a shower algorithm of the ‘jet calculus’ type [66]. To check
our results we also ran the JETSET programme \[14\] and found good agreement over the energy range where comparison was possible. However for the very high energies studied in this work, HERWIG proved to be more suitable for reasons of computing time \[68\]. Even so the calculations described here took several months on a Digital Alpha workstation.

For definiteness, we assume the heavy particles to decay into a quark-antiquark pair with unit branching ratio. The quark and antiquark, each carrying away energy \(m_X/2\), form jets which lead to the generation of many particles through cascading, hadronization and decays of some of the generated particles. This can be simulated by HERWIG via the annihilation process \(e^+e^- \rightarrow q\bar{q}\) with center-of-mass energy \(\sqrt{s} = m_X\), where \(q\) stands for all six kinematically allowed quark flavours. The event generator outputs kinematical details of all final state particles, e.g. protons, photons and leptons (electrons, positrons and neutrinos). We divided the \(x\)-range into 100 bins of width \(\Delta x = 0.01\). After each event simulation the number of protons, neutrinos, photons as well as electrons and positrons per energy bin was counted. We ran 10000 events for each of the masses \(m_X = 10^3, 10^5, 10^7, 10^9\) and \(10^{11}\) GeV. After all events had been run, the particle numbers in the bins are divided by the bin width and the number of events, in order to obtain the fragmentation functions \(dN_i/dx\). Apart from altering some relevant parameters in the computer code to allow it to run at the high energies studied here, we also switched off initial state radiation since it is not relevant for the present study. Unfortunately, it was not feasible to study the high \(x\) behaviour of the fragmentation functions for decaying particle masses higher than \(10^{11}\) GeV because of numerical convergence problems in the computer code. (Already for masses exceeding \(10^9\) GeV quadruple precision had to be used.) Hence we have had to extrapolate the fragmentation functions to high \(x\) for very heavy masses as described later.

First we show the proton fragmentation function obtained from the HERWIG runs in Figure 2 to illustrate that it depends on the decaying particle mass, contrary to the approximation (13) employed in previous work on topological defects \[34\]. Rather than being constant, it decreases with increasing \(m_X\) for \(x > \sim 0.1\), while at very low \(x\) it increases with increasing \(m_X\). The large fluctuations at \(x > \sim 0.5\) are due to the fact that relatively few particles are produced with such high energies despite the 10000 events per simulation. We note also that the shape differs significantly at high \(x\) from the approximation used in ref. \[37\].

In Figure 3 the fragmentation functions for protons, photons, neutrinos and electrons are compared for \(m_X = 10^{11}\) GeV. It is seen that at very low \(x\) there are more photons and neutrinos generated by the particle decay than electrons and protons. In the regime \(0.2 \lesssim x \lesssim 0.4\), photons, neutrinos and protons are generated with roughly equal abundances. However, for \(x > \approx 0.5\), photons and neutrinos again outnumber protons, in particular protons cut off at \(x \approx 0.75\) whereas neutrinos and photons are generated in the cascades with energies up to \(x \approx 0.95\). These differences will lead to different shapes of the expected fluxes \(I_i(E)\) as can be seen from eq. (12).

We now compare our proton fragmentation function with the commonly used Hill approximation \[33\] in Figure 4. Although his form provides a good fit for a low decaying particle mass, viz. \(m_X = 10^3\) GeV, it no longer does so for a high mass, viz. \(m_X = 10^{11}\) GeV. This is understandable given that the numerical co-efficients in eq. (13) were chosen to match relatively low energy collider data. However the functional form itself is well motivated and using our HERWIG runs we can determine new numerical co-efficients appropriate to heavier mass particles. Another advantage of the present approach is that the spectrum of neutrinos and photons is determined separately from that of the protons and not simply assumed to be proportional as in previous work \[34,35\].

To study the highest energy cosmic ray events we need to consider particle masses beyond...
10^{11} \text{ GeV} \) but this is difficult to do directly with HERWIG for technical reasons as mentioned earlier. We therefore resort to an extrapolation procedure as follows. For the range \( x \in [0, 0.2] \) the fragmentation functions are smooth and evolve monotonically with \( m_X \) so the fragmentation functions for a \( 10^{13} \text{ GeV} \) particle is obtained from simple linear extrapolation of the lower energy fragmentation functions in each individual energy bin. For \( x \in [0.2, 0.6] \) we first fit the calculated fragmentation functions to the form

\[
\frac{d N_{\text{fit1}}}{d x} = c_1 \frac{\exp[c_2 \sqrt{\ln(1/x)}(1 - x)^2]}{x \sqrt{\ln(1/x)}},
\]

for protons, and the form

\[
\frac{d N_{\text{fit2}}}{d x} = d_1 \frac{\exp[d_2 \ln(1/x)](1 - x)^2}{x \ln(1/x)},
\]

which proves more suitable for photons, neutrinos and electrons. The numerical co-efficients \( c_1, c_2 \) and \( d_1, d_2 \) are determined for particle masses less than \( 10^{11} \text{ GeV} \) by minimizing \( \chi^2 \) in the fit to the actual HERWIG runs. In Figures 5 and 6 we show these fits for \( x < \sim 0.6 \) to the proton and neutrino fragmentation functions corresponding to masses of \( 10^5 \) and \( 10^9 \text{ GeV} \). Then we determine the appropriate co-efficients for heavier masses by extrapolation. An example, for \( m_X = 10^{13} \text{ GeV} \), is shown in the figures. For \( x > \sim 0.6 \), statistical fluctuations become too severe so we extrapolate the fitting functions between the value at \( x = 0.55 \) and a cutoff which is taken to be \( x = 0.75 \) for protons and \( x = 0.95 \) for neutrinos, based on the observed behaviour for masses upto \( 10^{11} \text{ GeV} \) shown in Figures 2 and 3.

Finally, we mention the continuation of the proton fragmentation function for very low \( x \), viz. \( x \sim 0.015 \), which is relevant at high masses e.g. \( m_X = 10^{13} \text{ GeV} \). Since it proved impractical to have additional binning intervals at very small \( x \), we employ the fragmentation function (14) in this regime, normalized to our computations at \( x = 0.015 \).

**B. Comparison with observations**

With the fragmentation functions obtained above, we can now calculate the expected fluxes of protons and neutrinos from decays of particles such as cryptons in the halo. We normalize the calculated proton flux (12) to the observed cosmic ray flux at \( 10^{19} \text{ eV} \):

\[
\log_{10}[I_p(10^{19} \text{ eV})/\text{m}^{-2} \text{sec}^{-1} \text{sr}^{-1} \text{eV}^2] = 24.32.
\]

Note that the corresponding neutrino flux \( I_\nu(E) \) is then a *prediction* as the fragmentation function for neutrinos is computed independently.

The expected proton fluxes are shown in Figure 7. We see that a crypton with \( m_X = 10^{11} \text{ GeV} \) fits the flat power law well but cannot explain the events beyond \( 4 \times 10^{19} \text{ eV} \). Although this is easily achieved for \( m_X = 10^{13} \text{ GeV} \), the decays of such a massive particle would overproduce protons for \( E \gtrsim 3 \times 10^{19} \text{ eV} \). Thus a crypton with mass \( m_X \sim 10^{12} \text{ GeV} \) provides the best compromise although it too predicts a spectrum somewhat flatter than the one indicated observationally. (The reader is reminded that all differential fluxes have been multiplied by \( E^3 \) in eq. (12).)

An interesting signature for forthcoming experiments is the predicted ratio of the proton to neutrino flux [7]. In Figure 8 we compare the expected flux of protons and neutrinos for
\(m_X = 10^{12} \text{ GeV}. \) (We also show the photon flux to illustrate the difference from the prediction in ref. [37] but emphasize that this will be degraded through interactions with photon backgrounds during travel to Earth.) As can be seen, the neutrino flux exceeds the proton flux for \(10^{19} \text{ eV} \lesssim E \lesssim 10^{20} \text{ eV} \) and also for \(E \gtrsim 3 \times 10^{20} \text{ eV} \), as may have been anticipated from the comparison of their respective fragmentation functions. Thus the ratio \(I_p/I_\nu \) has a characteristic peak at about \(2 \times 10^{20} \text{ eV} \) as shown in Figure [3]. This could be a useful diagnostic of the decaying particle hypothesis for future experiments such as the Pierre Auger Project. Note that taking the extragalactic contribution into account would boost the neutrino flux by a factor of \(\sim 2 \) over that shown in the figures.

The abundance and lifetime of decaying particles such as cryptons are related through the spectrum normalization (18) as:

\[
\log_{10}(f_{\cos} \Omega_X h^2) = k + \log_{10}\left(\frac{\tau_X}{\text{yr}}\right)
\]

where \(k = -15.78, -16.13, -15.59 \) for crypton masses \(m_X = 10^{11}, 10^{12}, 10^{13} \text{ GeV} \) respectively. For a given crypton mass, a higher lifetime must be compensated for by a higher relic abundance, as illustrated in Figure [10]. So for example, if \(f_{\cos} \Omega_X h^2 \sim 1 \), cryptons with a mass of \(m_X = 10^{12} \text{ GeV} \) are required to have a lifetime of \(\tau_X \sim 10^{16} \text{ yr} \) if they are to explain the EHECR flux. If the enhancement in the halo is \(f_{\cos} \sim 3 \times 10^4 \) as expected for cold dark matter, then the lifetime may be increased to \(\sim 4 \times 10^{20} \text{ yr} \) if \(\Omega_X h^2 \sim 1 \); alternatively, for the same lifetime one could tolerate a lower relic abundance \(\Omega_X h^2 \sim 3 \times 10^{-5} \).

With regard to the fluxes of electrons and photons, both species would generate electromagnetic cascades on the prevalent radiation backgrounds through pair production and inverse Compton-scattering. A thorough analysis of such propagation effects and the resulting modifications of the injected photon and electron spectra has been performed [71]. It was found that the relic decaying particles with \(m_X \gtrsim 10^{14} \text{ GeV} \) would contribute excessively to the diffuse \(\gamma\)-ray background and are therefore ruled out. Hence, the mass range we favour, viz. \(10^{11} \lesssim m_X/\text{GeV} \lesssim 10^{13} \), does not lead to any conflict with observations. This conclusion is strengthened by the fact that according to our calculations the previous estimate [34] of the \(\gamma\)-ray flux from decaying particles was too high. Although the positrons released in the decays may be accumulated in the galactic halo, the astrophysical uncertainty in the containment time does not allow a restrictive constraint to be derived from limits on the positron flux in cosmic rays [72].

With regard to the neutrino background, the predicted flux at high energies is well below the upper limits derived from consideration of horizontal air showers [73], again because the decaying crypton mass is restricted to be less than about \(10^{13} \text{ GeV} \). It is also interesting to consider the flux at lower energies of \(\mathcal{O}(10^3) \text{ GeV} \) where experiments such the forthcoming ANTARES detector [74] will be most sensitive. As seen in Figure 8 the predicted neutrino flux dominates over the proton flux at low energies, thus the bulk of the energy released by the decaying cryptons ends up as neutrinos.\(^4\) Therefore we expect the neutrino flux at TeV energies to be at least \(\sim 10^8 \)

\(^4\)This expectation motivated the study undertaken earlier in which the abundance/lifetime of massive metastable relic particles was constrained using experimental limits on the high energy neutrino flux set by underground nucleon decay detectors and the Fly’s Eye experiment [30].
times larger than the EHECR flux at $10^{20}$ eV. Moreover the neutrinos should be well correlated in both time and arrival direction with the cosmic rays since the path length in the galactic halo is $\lesssim 100$ kpc. This is in contrast to the case of other suggested cosmologically distant sources such as gamma-ray bursts where the relative time delay can be up to $10^3$ yr [75].

**IV. CONCLUSIONS**

We have investigated the hypothesis that the highest energy cosmic rays, in particular those observed beyond the GZK cutoff, arise from the decay of massive metastable relic particles which constitute a fraction of the dark matter in the galactic halo. To simplify computations (using the HERWIG Monte Carlo event generator) we have considered only decays into $q\bar{q}$ pairs with unit branching ratio. Comparison with experimental data indicates that a decaying particle mass of $O(10^{12})$ GeV is required to fit the spectral shape while the absolute flux requires a lifetime of $O(10^{20})$ yr if such particles contribute the critical density. The predicted decay spectra may be somewhat altered if 3-body decays and other final states (e.g. supersymmetric particles [76]) are considered. However our conclusions regarding the preferred mass and relic abundance/lifetime of the decaying particle are unlikely to be affected. In particular it would appear that the approximations used to calculate the particle spectra in previous studies of decaying topological defects [34] and hypothetical massive particles [77] were not sufficiently accurate. Our work indicates that the topological defect model is disfavoured unless the mass of the decaying gauge bosons is less than about $10^{13}$ GeV, which is well below the unification scale of $\sim 10^{16}$ GeV. (A similar conclusion is arrived at by independent arguments in refs. [77,78].) By contrast, cryptons from the hidden sector of supersymmetry breaking have a mass of the required order, as well as a decay lifetime which is naturally suppressed. However their relic abundance is difficult to estimate reliably, although we have argued that it may be cosmologically interesting.

The primary intention of this work is to attempt to quantify the decaying particle hypothesis in a manner which is of interest to experimentalists. We have therefore computed the expected neutrino to proton ratio as a function of energy since this is an important test of competing hypotheses for forthcoming experiments, in particular the Pierre Auger project [83]. Of course our cleanest prediction is that the cosmic ray spectrum should cut off just below the mass of the decaying crypton, at $\sim 3 \times 10^{20}$ eV. Moreover, with sufficient event statistics it should be possible to identify the small anisotropy which should result from the distribution of the decaying particles in the Galactic halo [79]. Thus although the hypothesis investigated here is very speculative, it is nevertheless testable. Perhaps Nature has indeed been kind to us and provided a spectacular cosmic signature of physics well beyond the Standard Model.

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FIG. 1. The high energy cosmic ray spectrum beyond the ‘ankle’. (Note that the differential flux has been multiplied by $E^3$.) The data shown are from AGASA, stereo Fly’s Eye, Haverah Park and Yakutsk and are AGASA-normalized [8]. The highest energy monocular Fly’s Eye event at $3 \times 10^{19}$ eV is also shown. A fit to the spectrum from the superposition of a steeply falling and a flatter power law (dashed lines) is indicated [2].
FIG. 2. Fragmentation function of protons, for decaying particle masses $m_X = 10^3$, $10^7$, and $10^{11}$ GeV, computed using the event generator HERWIG.
FIG. 3. Fragmentation functions of photons, neutrinos, and electrons compared to that of protons for a decaying particle mass of $m_X = 10^{11}$ GeV.
FIG. 4. Comparison of the computed proton fragmentation function with the leading-log approximation of eq. (13), normalized to the HERWIG computation at $x = 0.042$. The upper solid line refers to a decaying particle mass of $m_X = 10^3$ GeV and the lower one to $m_X = 10^{11}$ GeV.
FIG. 5. Extrapolation of the computed proton fragmentation function to higher decaying particle masses. The upper two solid lines are HERWIG results for decaying particles of mass $m_X = 10^5$ and $10^9$ GeV while the dashed lines are the best fitting functions according to eq. (16). (The corresponding $\chi^2$-values are also indicated.) The lower solid line is the fragmentation function for $m_X = 10^{13}$ GeV obtained from extrapolating the fitting parameters.
FIG. 6. Same as Figure 5 for the case of neutrinos with the fitting function now given by eq. (17).
FIG. 7. Expected proton flux for decaying particle masses $m_X = 10^{11}$, $10^{12}$ and $10^{13}$ GeV compared with observations. The theoretical spectra are normalized at $10^{19}$ eV to the flat component (dashed line) suggested by the Fly’s Eye data.
FIG. 8. Expected fluxes of protons, neutrinos and photons from decays of a decaying particle with mass $m_X = 10^{12}$ GeV (normalized as in Figure 7) compared with the observations. Note that the photon flux will be degraded through interactions with the CMB during travel to Earth and is shown for illustrative purposes only.
FIG. 9. The ratio of the proton to neutrino flux for the decaying particle mass $m_X = 10^{12} \text{ GeV}$.
FIG. 10. The relic decaying particle abundance versus lifetime for various masses, as required by the flux normalization to the observations.