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Are ultrahigh energy cosmic rays signals of supersymmetry?

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We investigate the suggestion that cosmic ray particles of energy larger than the Greisen–Zatsepin–Kuzmin cutoff are not nucleons, but stable, massive, supersymmetric hadrons, S^0s, which are *uds*-gluino bound states. Because the propagation range of S^0s through the cosmic background radiation is significantly longer than the range of nucleons, S^0s can originate from sources at cosmological distances.

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I. INTRODUCTION

Detection of cosmic rays of energies above 10^{20} eV [1,2] has raised yet unsettled questions regarding their origin and composition. The first problem is that it is difficult to imagine any astrophysical site for the cosmic accelerator (for a review, see Ref. [3]). The Larmour relation for a particle of charge Z, $(E/10^{18}\,\mathrm{eV}) = Z(R/\,\mathrm{kpc})(|\vec{B}|/\mu\mathrm{Gauss})$, sets the scales for the required size, R, and magnetic field strength, $|\vec{B}|$, of the accelerator. One would expect any sources with sufficient $R|\vec{B}|$ to accelerate particles to ultrahigh energies to appear quite unusual in other regards.

A second issue is the composition of the observed cosmic rays. The shower profile of the highest energy events are inconsistent with their identification as photons, and suggests that they are hadrons [4]. The problem is that the propagation of neutrons, protons, or nuclei over astrophysical distances is strongly affected by the existence of the cosmic background radiation (CBR). Above threshold, cosmic-ray nucleons lose energy by photoproduction of pions, $N\gamma \to N\pi$, resulting in the Greisen–Zatsepin–Kuzmin (GZK) cutoff in the maximum energy of cosmic-ray nucleons. If the primary is a heavy nucleus, then it will be photo-disintegrated by scattering with CBR photons. This suggests that if the primary is a nucleon or a nucleus, the sources must be nearby (less than about 50 Mpc).

Since UHE cosmic rays should be largely unaffected by intergalactic or galactic magnetic fields, by measuring the incident direction of the cosmic ray it should be possible to trace back and identify the source. The relatively short range of nucleons or nuclei is a severe problem, because unusual sources such as quasars and Seyfert galaxies typically are beyond the "range" expected for nucleons or nuclei. Possible candidate sources within 10° of the ultrahigh-energy (UHE) cosmic ray observed by the Fly's Eye [2] were studied in Ref. [5].[1] The quasar 3C 147 and the Seyfert galaxy MCG 8-11-11 are attractive candidates. Lying within the 1σ error box of the primary's incoming direction, the quasar 3C 147 has a large radio luminosity ($7.9 \times 10^{44}\,\mathrm{erg\,s^{-1}}$) and an X-ray luminosity of about the same order of magnitude, indicative of a large number of strongly accelerated electrons in the region. It also produces a large Faraday rotation, with a rotation measure $\mathrm{RM} = -1510 \pm 50\,\mathrm{rad\,m^{-2}}$, indicative of a large magnetic field over large distances. It is also interesting that this source is within the error box of a UHE event seen by the Yakutsk detector. However, it lies at a red shift of about $z = 0.545$, well beyond $z < 0.0125$ adopted in Ref. [5] as the GZK cutoff.

Just outside the 2σ error box of the primary's incoming direction is the Seyfert galaxy MCG 8-11-11. It is also unusual, with large X-ray and low-energy gamma-ray luminosities ($4.6 \times 10^{44}\,\mathrm{erg\,s^{-1}}$ in the $20 - 100\,\mathrm{keV}$ region and $7 \times 10^{46}\,\mathrm{erg\,s^{-1}}$ in the $0.09 - 3\,\mathrm{MeV}$ region). At a redshift of $z = 0.0205$, it is much closer than 3C 147, but still too distant for the flux to be consistent with the observed proton flux at lower energies [5], and beyond the GZK cutoff.

Briefly stated, the problem is that there are no known candidate astronomical sources within the range of protons, neutrons, or nuclei. In this paper we propose that the answer to this cosmic-ray conundrum is that UHE cosmic rays are not protons, neutrons, or nuclei, but a new species of particle we denote as the uhecron, U. The meager information we have about the cosmic ray events allows us to assemble a profile for the properties of the uhecron:

1) The uhecron interacts strongly: Although there are only a handful of UHE events, the observed shower developments suggest a strongly interacting primary.

2) The uhecron is stable or very long lived: Clearly if the particle originates from

[1] Ten degrees is the estimated maximum deflection angle due to magnetic fields for a proton of this energy.

cosmological distance, it must be stable, or at least remarkably long lived, with $\tau \gtrsim 10^6\mathrm{s}(m_U/3\,\mathrm{GeV})(L/1\,\mathrm{Gpc})$ where L is the distance to the source.

3) The uhecron is massive, with mass greater than about 2 GeV: If the cosmic ray is massive, the threshold energy for pion production increases, and the energy lost per scattering on a CBR photon will decrease. We will go into the details of energy loss in the paper, but it is simple to understand the reason from simple kinematics. In $U\gamma \to U\pi$, the threshold for pion production is $s_{\rm min} = m_U^2 + m_\pi^2 + 2m_U m_\pi$. In the cosmic-ray frame where the U has energy $E_U \gg m_U$ and the photon has energy $E_\gamma \sim 3T$ (where $T = 2.4\times 10^{-4}\,\mathrm{eV}$ is the temperature of the CBR), $s \simeq m_U^2 + 4E_\gamma E_U$. Thus, the threshold for pion production, $s \geq s_{\rm min}$, results in the limit $E_U \gtrsim m_\pi m_U/2E_\gamma$. For $E_\gamma = 3T$, and identifying the uhecron as a proton, the threshold is $E_U \approx 10^{20}\,\mathrm{eV}$. Of course the actual threshold is more involved because there is a distribution in photon energy and scattering angle, but the obvious lesson is that if the mass of the primary is increased, the threshold for pion production increases, and the corresponding GZK cutoff will increase with the mass of the cosmic ray. Furthermore, since the fractional energy loss per collision will be approximately m_π/m_U, a massive uhecron will lose energy via pion-photoproduction at a slower rate than that of a lighter particle. Another potential bonus if the cosmic ray is not a neutron or a proton is that the cross section for $U\gamma \to U\pi$ near the threshold may not be strongly enhanced by a resonance such as $\Delta(1232)$, as when the U is a nucleon. Although there may well be a resonance in the $U\pi$ channel, it might not have the strength or be as near the pion-photoproduction threshold as the Δ in the pion-nucleon channel.

4) We will assume the uhecron is electrically neutral: Although not as crucial a requirement as the first three, there are three advantages if the uhecron is neutral. The first is that it will not lose energy through e^+e^- pair production off the CBR photons. Another advantage of a neutral particle is that since it will be unaffected by intergalactic and galactic magnetic fields, its straight trajectory will directly point

back to its source, and thus it will be able to account for the common direction of the two highest energy events. Thirdly, there will be no energy losses due to synchrotron or bremsstrahlung radiation. Of course since a neutral particle will not be accelerated by normal electromagnetic mechanisms, it is necessary to provide at least a plausibility argument that they can be produced near the source, for instance as secondaries in collisions induced by high-energy protons.

In this paper we analyze the possibility that a supersymmetric baryon S^0 (uds-gluino bound state whose mass is expected to be in the range 1.9 to 2.3 GeV—see below) is the uhecron rather than a proton, as first proposed in Ref. [6]. The S^0 has strong interactions, it can be stable, it is more massive than the nucleon, and it is neutral with vanishing magnetic moment [6]. If UHE cosmic rays are S^0s, we will show that their range is at least an order of magnitude greater than that of a proton, putting MCG 8-11-11 (and possibly even 3C 147) within range of the Fly's Eye event.

II. PRODUCTION OF UHE S^0s

We first address the question of whether there is a plausible scenario to produce UHE S^0s. This is a tricky question, since there is no clear consensus on the acceleration mechanism even if the primary particle is a proton. Here we simply assume that somehow UHE protons are produced, and ask if there is some way to turn a beam of UHE protons into UHE S^0s.

Assuming that there exists an astrophysical accelerator that can accelerate protons to energies above 10^{21} eV, one can envisage a plausible scenario of S^0 production through proton collision with hadronic matter surrounding the accelerator. A p-nucleon collision will result in the production of R_ps, the uud-gluino states whose mass is about 200 MeV above that of the S^0. The R_p decays to an S^0 and a π^+,[2] with the S^0 receiving a momentum fraction of about $(m_{S^0}/m_{R_p})^2$. From a triple Regge model of the collision, one estimates that the distribution of the produced R_ps as a function of the outgoing momentum fraction x is $d\sigma/dx \sim (1-x)^{1-2\alpha}(s')^{\alpha_P-1}$ as x approaches unity. Here, $s' = (1-x)s$, and α is the Regge intercept of the SUSY-partner of the Pomeron. Hence, one has $\alpha = \alpha_P - 1/2 = \epsilon + 1/2$, where $\epsilon \approx 0.1$ is the amount the pomeron trajectory is above 1 at high energies. Thus we parameterize the S^0 production cross section in a p-nucleon collision as $d\sigma/dx = AE_p^\epsilon$ and x is the ratio of the S^0 energy to the incident energy. Parameterizing the high energy proton flux from the cosmic accelerator as $dN_p/dE_p = BE_p^{-\gamma}$, we have a final S^0 flux of $dN_{S^0}/dE = \kappa nLABE^{-\gamma+\epsilon}$, where nL is the matter column density with which the proton interacts to produce an R_p and κ is of order 1 (for $\gamma = 2$, $\kappa = 0.4$). Note that the produced S^0s are distributed according to a spectrum that is a bit flatter than the high energy proton spectrum. However a disadvantage to this indirect method of producing S^0s is the suppression factor of about AE^ϵ/σ_{pp}, where σ_{pp} is the proton-proton total diffractive cross section.[3] This may disfavor extremely distant candidates like 3C 147 as a source, since the required particle flux for the detected flux on Earth already pushes the luminosity limit. Assuming that the 3.2×10^{20} eV event of Fly's Eye came during its exposure to 3C 147, the resulting time-averaged flux is $11\,\mathrm{eV\,cm^{-2}\,s^{-1}}$, which is greater than the X-ray luminosity of 3C 147 [5].

Another possible mechanism of high energy S^0 production is the direct acceleration of charged light SUSY hadrons (mass around 2 GeV), such as R_p and R_Ω, whose lifetime

[2]The decay $R_p \to S^0\pi$ was the subject of an experimental search [7]. However the sensitivity was insufficient in the mass range of interest ($m(R_p) = 2.1 - 2.5$ GeV) for a signal to have been expected.

[3]Here we have placed the constraint that the column density satisfy the condition $nL\sigma_{pp} \sim 1$ since the proton will otherwise lose too much energy before producing an S^0.

is about 10^{-11} s [6]. If the accelerated SUSY hadrons have a sufficient time-dilation factor, whatever electromagnetic mechanism accelerates the protons may also be able to accelerate the high energy SUSY hadrons. Then, one can imagine that the high energy tail of the hadronic plasma which gets accelerated by some electromagnetic mechanism will consist of a statistical mixture of all light strong-interaction-stable charged hadrons. Then the flux of the resulting S^0 will have the same spectrum as the protons, differing in amplitude by a factor of order unity, which depends on the amount of SUSY hadrons making up the statistical mixture. Unfortunately, conventional shock wave acceleration mechanisms would require a too long time scale for this mechanism to be feasible (e.g., Ref. [8]). However, some electromagnetic "one push" mechanisms similar to the one involving electric fields around pulsars [3,9] may allow this kind of acceleration if the electric field can be large enough.

A somewhat remote possibility is that there may be gravitational acceleration mechanisms which would not work for a charged particle (because of radiation energy losses and magnetic confinement) but would work for a neutral, zero magnetic moment particle such as an S^0. For example, if S^0s exist in the high energy tail of the distribution of accreting mass near a black hole (either by being gravitationally pulled in themselves or by being produced by a proton collision), they may have enough energy to escape with a large energy. A charged particle, on the other hand, will not be able to escape due to radiation losses. Unfortunately, this scenario may run into low flux problems due to its reliance on the tail of an energy distribution.

Finally, decay of long-lived superheavy relics of the big bang would produce all light particles present in the low-energy world, including the S^0.

III. PROPAGATION OF UHE COSMIC RAYS

To calculate the energy loss due to the primary's interaction with the CBR, we follow the continuous, mean energy loss approximation used in Ref. [10]. In this approximation we smooth over the discrete nature of the scattering processes, neglecting the stochastic nature of the energy loss, to write a continuous differential equation for the time evolution of the primary energy of a single particle. The proper interpretation of our result is the mean energy of an ensemble of primaries traveling through the CBR. We shall now delineate the construction of the differential equation.

For an ultrahigh energy proton (near 10^{20} eV in CBR frame[4]), three main mechanisms contribute to the depletion of the particle's energy: pion-photoproduction, e^+e^- pair production, and the cosmological redshift of the momentum. Pion-photoproduction consists of the reactions $p\gamma \to \pi^0 p$ and $p\gamma \to \pi^+ n$. Pion-photoproduction, which proceeds by excitation of a resonance, is the strongest source of energy loss for energies above about 10^{20} eV, while below about $10^{19.5}$ eV, e^+e^- pair production dominates. For the scattering processes (pion-photoproduction and e^+e^- pair production), the mean change in the proton energy (E_p) per unit time (in the CBR frame) is

$$\frac{dE_p(\text{scatter})}{dt} = -\sum_{\text{events}} (\text{mean event rate}) \times \Delta E \tag{1}$$

where the sum is over distinct scattering events with an energy loss of ΔE per event. The mean event rate is given by

$$\text{mean event rate} = \frac{1}{\gamma}\frac{d\sigma}{d\xi} f(E_\gamma) dE_\gamma d\xi \tag{2}$$

where $\gamma = E_p/m_p$ is necessary to convert from the event rate in the proton frame (proton's rest frame), where we perform the calculation, to the CBR frame, $d\sigma/d\xi$ is

[4] Let this be the frame in which CBR has an isotropic distribution.

the differential cross section in the proton frame,[5] and f is the number of photons per energy per volume in the proton frame. To obtain f we start with the isotropic Planck distribution and then boost it with the velocity parameter β to the proton frame

$$n(E_\gamma, \theta) = \frac{1}{(2\pi)^3}\left[\frac{2E_\gamma^2}{\exp[\gamma E_\gamma(1+\beta\cos\theta)/T]-1}\right] \tag{3}$$

where θ is the angle that the photon direction makes with respect to the boost direction. Integrating Eq. (3) over the solid angle[6] and taking the ultrarelativistic limit, we find

$$f = \frac{E_\gamma T}{2\pi^2\gamma}\ln\left[\frac{1}{1-\exp(-E_\gamma/2\gamma T)}\right]. \tag{4}$$

For ΔE, the energy loss per event in the CBR frame, we can write

$$\Delta E(\cos\theta, p_r) = \gamma m_p\left[1+\frac{\beta p_r}{m_p}\cos\theta - \sqrt{1+\left(\frac{p_r}{m_p}\right)^2}\right] \tag{5}$$

where p_r, which may depend on E_γ and $\cos\theta$, is the recoil momentum of the proton and θ is the angle between the incoming photon direction and the outgoing proton direction. Putting all these together, the energy loss rate due to scattering given by Eq. (1) becomes

$$\begin{aligned}\frac{dE_p(\text{scatter})}{dt} &= -\gamma^{-1}\int dE_\gamma f(E_\gamma)\\ &\quad\times\sum_i\int d\xi_i\frac{d\sigma_i}{d\xi_i}(E_\gamma,\xi_i)\,\Delta E(\cos\theta(E_\gamma,\xi_i), p_r(E_\gamma,\xi_i))\end{aligned} \tag{6}$$

where only functions yet to be specified are the recoil momentum and the differential cross section (for each type of reaction i).

For the reaction involving the production of a single pion, the recoil momentum of the protons in the proton frame can be expressed as

[5]The differential $d\xi$ is $dQd\eta$ (Q and η are defined below) for the e^+e^- pair production while it is $d\cos\theta$ for the pion-photoproduction.

[6]The exact angular integration range is unimportant as long as the range encompasses $\cos\theta = -1$ (where the photon distribution is strongly peaked in the ultrarelativistic limit) since we will be taking the ultrarelativistic limit.

$$p_r(E_\gamma, \cos\theta) = \frac{2q^2 E_\gamma \cos\theta \pm (E_\gamma + m_p)\sqrt{4E_\gamma^2 m_p^2 \cos^2\theta - 4m_\pi^2 m_p (E_\gamma + m_p) + m_\pi^4}}{2\left[(E_\gamma + m_p)^2 - E_\gamma^2 \cos^2\theta\right]} \tag{7}$$

where $q^2 = m_p(m_p + E_\gamma) - m_\pi^2/2$. When the photon energy E_γ is approximately at threshold energy of $m_\pi + m_\pi^2/2m_p$ and the proton recoils in the direction $\theta = 0$, the recoil momentum is about m_π. The recoil momentum is a double valued function, where the negative branch corresponds to the situation where most of the photon's incoming momentum is absorbed by the pion going out in the direction of the incoming photon. Thus, since the positive branch will be more effective in retarding the proton (in the CBR frame), we will neglect the negative branch to obtain a conservative estimate of the "cutoff" distance. It is possible to work out the kinematics for multipion production, but for our purpose of making a reasonably conservative estimate, it is adequate to use Eq. (7) as the recoil momentum even for multipion production.[7]

The pion-photoproduction cross section has been estimated by assuming that the s-wave contribution dominates, which would certainly be true near the threshold of the production. The cross section is taken to be a sum of a Breit-Wigner piece and two non-resonant pieces:

$$\begin{aligned}
\sigma(\text{pion}) &= 2\sigma_{1\pi}\, \Theta\!\left(E_\gamma - m_\pi - \frac{m_\pi^2}{2m_p}\right) + 2\sigma_{\text{multipion}} \\
\sigma_{1\pi} &= \frac{4\pi}{p_{cm}^2}\left[\frac{m_\Delta^2 \Gamma(\Delta \to \gamma p)\Gamma(\Delta \to \pi P)}{(m_\Delta^2 - s)^2 + m_\Delta^2 \Gamma_{\text{tot}}^2}\right] + \sigma_{\text{nonres}} \\
\Gamma(\Delta \to Xp) &= \frac{p_{cm}^X \omega_X}{8 m_\Delta \sqrt{s}} \\
\Gamma_{\text{tot}} &= \frac{p_{\text{cm}}^\pi}{\sqrt{s}} \frac{2m_\Delta^2 \overline{\Gamma}_{\text{tot}}}{\sqrt{[m_\Delta^2 - (m_\pi + m_p)^2]\,[m_\Delta^2 - (m_p - m_\pi)^2]}}
\end{aligned}$$

[7]For example, one can easily verify that the maximum proton recoil during one pion production is greater than the maximum proton recoil during two pion production.

$$\begin{aligned}\sigma_{\text{nonres}} &= \frac{1}{16\pi s}\frac{\sqrt{[s-(m_p+m_\pi)^2]\,[s-(m_p-m_\pi)^2]}}{(s-m_p^2)}|\mathcal{M}(p\gamma\to\pi p)|^2 \\ \sigma_{\text{multipion}} &= a\tanh\left(\frac{E_\gamma - E_{\text{multi}}}{m_\pi}\right)\Theta(E_\gamma - E_{\text{multi}}) \end{aligned} \qquad (8)$$

where ω_X is defined through $4\pi\omega_X \equiv \int d\Omega|\mathcal{M}(\Delta\to Xp)|^2$, $\mathcal{M}$ denotes an invariant amplitude, the center of momentum momentum is given as usual by

$$p_{cm}^X = \sqrt{\frac{[s-(m_p+m_X)^2]\,[s-(m_p-m_X)^2]}{4s}}, \qquad (9)$$

and $\sigma_{\text{multipion}}$ is a crude approximation[8] for the contribution from the multipion production whose threshold is at $E_{\text{multi}} = 2(m_\pi + m_\pi^2/m_p)$. The σ_π component of the cross section is fit[9] to the $p\gamma\to n\pi^0$ data of Ref. [11], while the amplitude a for $\sigma_{\text{multipion}}$ is estimated from the $p\gamma\to Xp$ data for energies $E_\gamma \gtrsim 0.6\,\text{GeV}$. The numerical values of the parameters resulting from the fit are $(\omega_\gamma\omega_\pi) = 0.086\,\text{GeV}^4$, $|\mathcal{M}(p\gamma\to\pi p)| = 0.018$, $\overline{\Gamma}_{\text{tot}} = 0.111\,\text{GeV}$, $m_\Delta = 1.23\,\text{GeV}$, and $a = 0.2\text{mb}$. The factor of 2 multiplying σ_π accounts for the two reactions $p\gamma\to\pi^0 p$ and $p\gamma\to\pi^+ n$, since a neutron behaves, to first approximation, just like the proton. For example, the dominant pion-photoproduction reactions involving neutrons are $n\gamma\to\pi^0 n$ and $n\gamma\to\pi^- p$ which have similar cross sections as the analogous equations for protons. Thus, we are really estimating the energy loss of a nucleon, and not just a proton.

Taking the $p\gamma\to e^+e^-p$ differential cross section from Ref. [12] (as done in Ref. [10]),

[8] The functional form was chosen to account for the shape of the cross section given in Ref. [11].

[9] The fit is qualitatively good, but only tolerable quantitatively. The fit to the data in the range between 0.212 GeV and 0.4 GeV resulted in a reduced $\chi^2_{16} \sim 50$ (due to relatively small error bars). This is sufficient for our purposes since our results should depend mainly upon the gross features of the cross section.

we use[10]

$$
\begin{aligned}
\frac{d\sigma(\text{pair})}{dQd\eta} &= \Theta(E_\gamma - 2m_e) \times \frac{4\alpha^3}{E_\gamma^2}\frac{1}{Q^2}\left\{\ln\left(\frac{1-w}{1+w}\right)\left[\left(1-\frac{E_\gamma^2}{4m_e^2\eta^2}\right)\right.\right. \\
&\times\left.\left(1-\frac{1}{4\eta^2}+\frac{1}{2\eta Q}-\frac{1}{8Q^2\eta^2}-\frac{Q}{\eta}+\frac{Q^2}{2\eta^2}\right)+\frac{E_\gamma^2}{8m_e^2\eta^4}\right] \\
&\left.+w\left[\left(1-\frac{E_\gamma^2}{4m_e^2\eta^2}\right)\left(1-\frac{1}{4\eta^2}+\frac{1}{2\eta Q}\right)+\frac{1}{\eta^2}\left(1-\frac{E_\gamma^2}{2m_e^2\eta^2}\right)(-2Q\eta+Q^2)\right]\right\}, \quad (10)
\end{aligned}
$$

where $w = [1 - 1/(2Q\eta - Q^2)]^{1/2}$. The recoil momentum is contained in $Q = p_r/2m_e$, and the photon energy is contained in $\eta = E_\gamma \cos\theta/2m_e$.

The final ingredient in our energy loss formula is the redshift due to Hubble expansion. We assume a matter-dominated, flat FRW universe with no cosmological constant. Thus, the cosmological scale factor is proportional to $t^{2/3}$. The energy loss for relativistic particles (such as our high energy proton) due to redshift is then given by

$$
\frac{dE_p(\text{redshift})}{dt} = -\frac{2E_p}{3t}. \tag{11}
$$

Furthermore, note that the expansion of the universe causes the temperature to vary with time as $t^{-2/3}$.

Adding Eqs. (6) and (11), we have the proton energy loss equation

$$
\frac{dE_p}{dt} = \frac{dE_p(\text{scatter})}{dt} + \frac{dE_p(\text{redshift})}{dt}\ , \tag{12}
$$

whose integration from some initial cosmological time t_i to the present time t_0 gives the present energy of the proton that was injected with energy E_i at time t_i. Note that we are interested in plotting $E_p(t_0)$ as a function of $t_0 - t_i$ with t_0 fixed, which is not equivalent to fixing t_i and varying t_0 because there is no time translational invariance in an FRW universe. Note also that we need to set the Hubble parameter h (where the Hubble constant is $100h\,\text{km}\,\text{s}^{-1}\,\text{Mpc}^{-1}$) in our calculation because the conversion

[10] We ignore that n does not pair produce e^+e^-. However, this has consequences only for energies below about $10^{19.5}$ eV.

between time and the redshift depends on h. To show the degree of sensitivity of our results to h we will calculate the energy loss for $h = 0.5$ and $h = 0.8$.

Now, suppose the primary cosmic ray is an S^0 instead of a proton. The e^+e^- pair production will be absent (to the level of our approximation) because of the neutrality of S^0. Furthermore, the mass splitting between S^0 and any one of the nearby resonances that can be excited in a γS^0 interaction is larger than the proton-Δ mass splitting, leading to a further increase in the attenuation length of the primary. Perhaps most importantly, the mass of S^0 being about two times that of the proton increases the attenuation length significantly because of two effects. One obvious effect is seen in Eq. (7), where the fractional energy loss per collision to leading approximation is proportional to p_r/m_p while p_r has a maximum value of about m_π when E_γ is at the pion-photoproduction threshold. Replacement of $m_p \to m_{S^0}$ obviously leads to a smaller energy loss per collision. The second effect is seen in Eqs. (4) and (6), where for the bulk of the photon energy integration region, a decrease in γ (in the exponent) resulting from an increase in the primary's mass suppresses the photon number. In fact, it is easy to show that if we treat the cross section to be a constant, the pion-photoproduction contribution to the right hand side of Eq. (6) can be roughly approximated as

$$\frac{dE_p(\pi)}{dt} \approx -\frac{m_\pi^2 T^2 \sigma}{\pi^2} \exp(-y/2) \left(1 + \frac{3}{y} + \frac{4}{y^2}\right) \tag{13}$$

where $y = m_\pi m_p/(E_p T)$, clearly showing a significant increase in the attenuation length as m_p is replaced by m_{S^0}.

The relevant resonances for the $S^0\gamma$ collisions are spin-1 R_Λ and R_Σ [6] (whose constituents are those of the usual Λ and Σ baryons, but in a color octet state, coupled to a gluino [13]). There are two R-baryon flavor octets with $J = 1$. Neglecting the mixing between the states, the states with quarks contributing spin 3/2 have masses of about $385 - 460\,\mathrm{MeV}$ above that of the S^0 and the states with quarks contributing spin 1/2

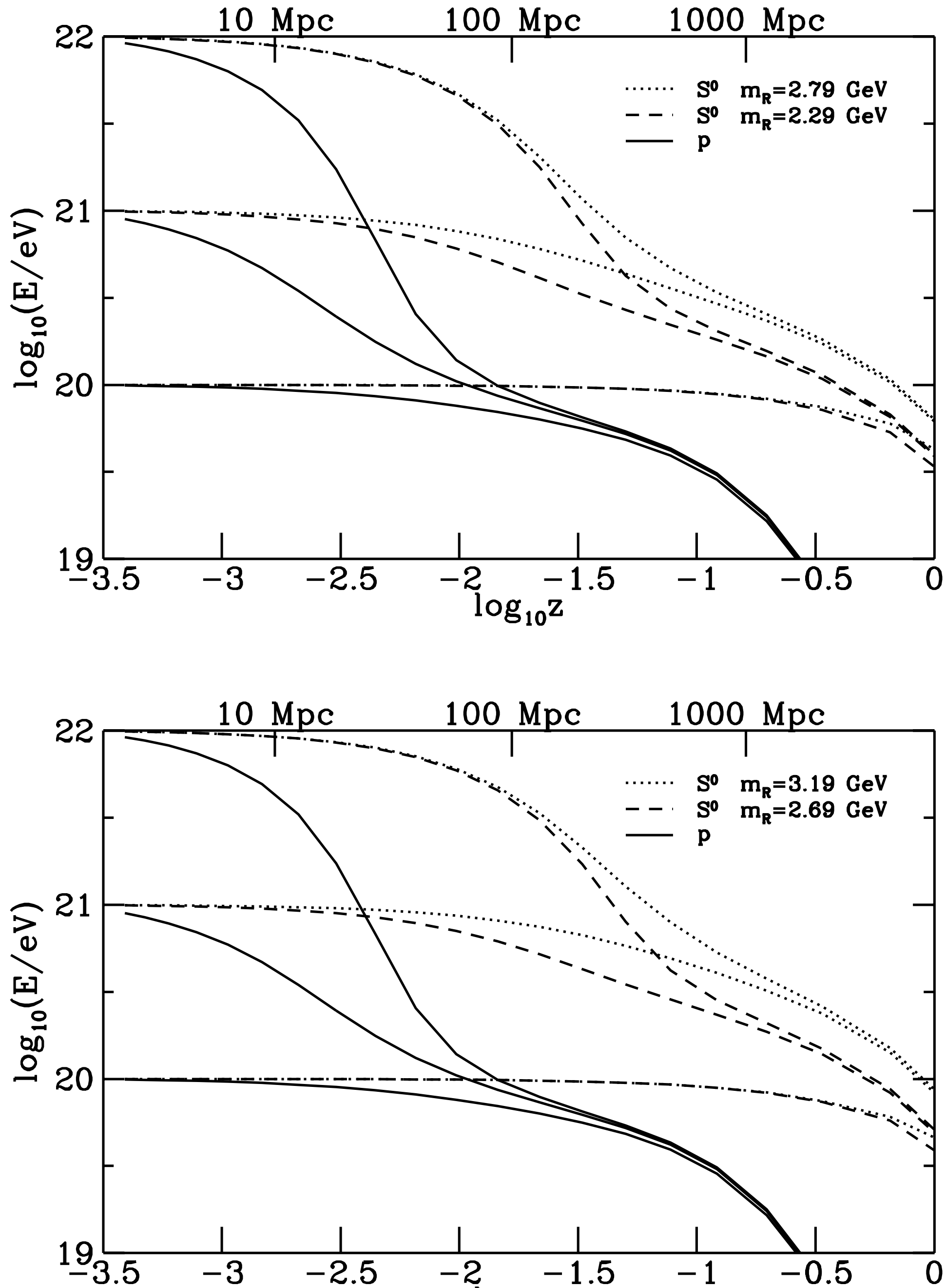
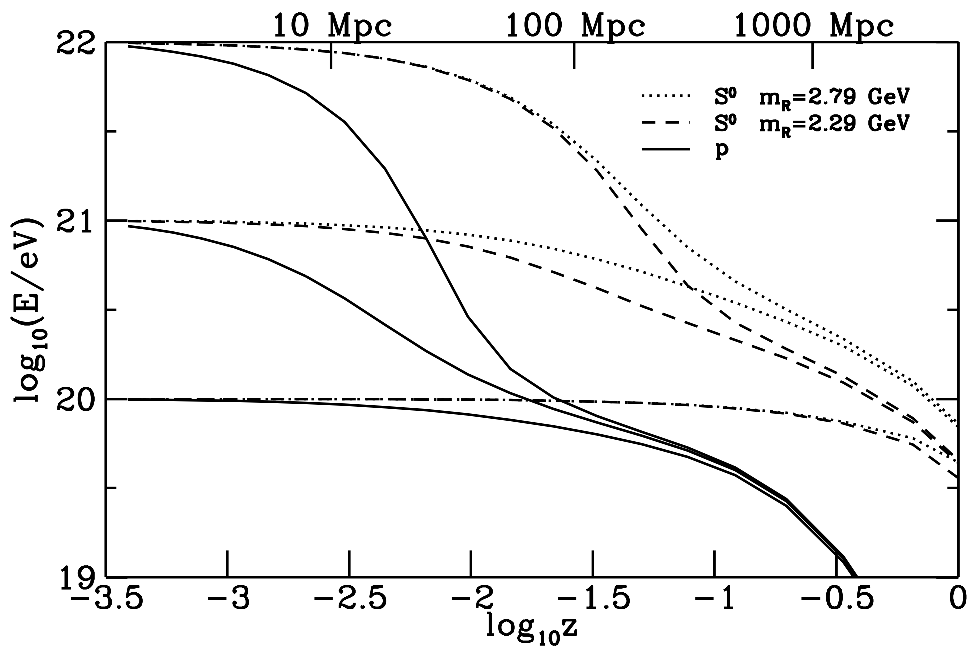


Fig. 1: The figures show the primary particle's energy as it would be observed on Earth today if it were injected with various energies (10^{22} eV, 10^{21} eV, and 10^{20} eV) at various redshifts. The distances correspond to luminosity distances. The mass of S^0 is 1.9 GeV in the upper plot while it is 2.3 GeV in the lower plot. Here, the Hubble constant has been set to $50\,\mathrm{km\,s^{-1}\,Mpc^{-1}}$.



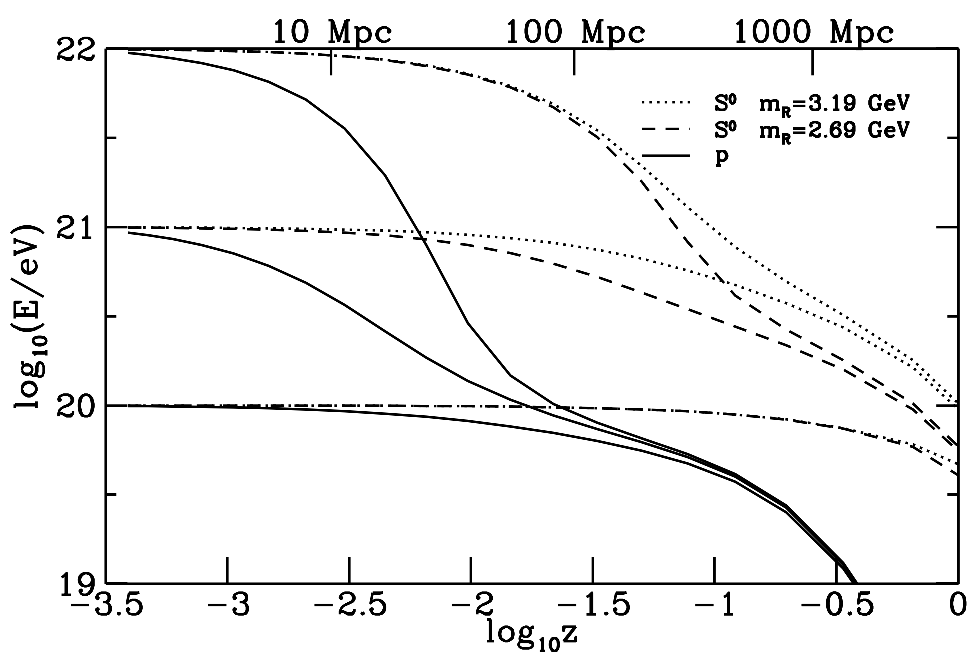


Fig. 2: Same as Fig. 1 except with the Hubble constant equal to $80\,\mathrm{km\,s^{-1}\,Mpc^{-1}}$.

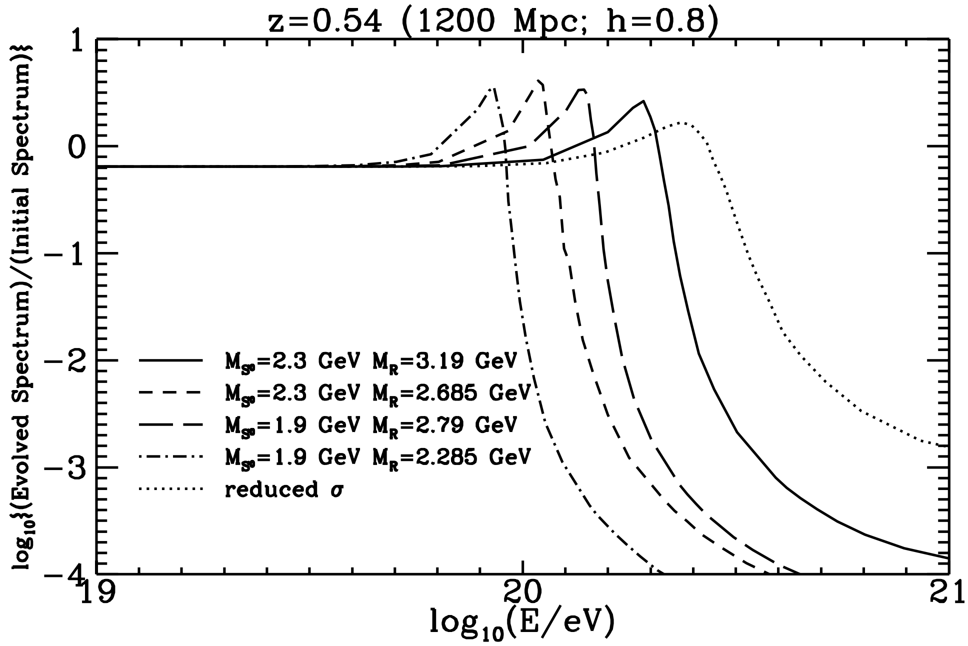


Fig. 3: An initial S^0 injection spectrum having a power law form of E^{-2} is evolved through the particle's interaction with the CBR during its 1200 Mpc travel to Earth. The masses of the S^0 and its associated resonance are shown. The curve labeled "reduced σ" has the same mass parameters as the solid curve except with our conservative estimation of the total cross section reduced by a factor of half.

have masses of about $815 - 890$ MeV above that of the S^0. If we take the mass of R^0 to be about 1.7 GeV and require that the photino ($\tilde{\gamma}$) is a significant dark matter component, then according to Ref. [14], $m_{\tilde{\gamma}}$ lies in the range 0.9 to 1.4 GeV. If we assume that S^0 is minimally stable, we have $m_{S^0} \approx m_p + m_{\tilde{\gamma}}$ resulting in an m_{S^0} in the range 1.9 to 2.3 GeV. The other resonance parameters are fixed at the same values as those for the protons.

In Fig. 1, we show the proton energy and the S^0 energy today (with $h = 0.5$) if it had been injected at a redshift z (or equivalently from the corresponding distance[11]) with an energy of 10^{22} eV, 10^{21} eV, and 10^{20} eV. To explore the interesting mass range, we have

[11] Marked are the luminosity distances $d_L = H_0^{-1} q_0^{-2} \left[z q_0 + (q_0 - 1)(\sqrt{2 q_0 z + 1} - 1) \right]$ where the deceleration parameter q_0 is 1/2 in our $\Omega_0 = 1$ universe.

set the S^0 mass to 1.9 GeV in the upper plot while we have set it to 2.3 GeV in the lower plot. For the cosmic rays arriving with 10^{20} eV, the distance is increased by more than thirty times, while for those arriving with $10^{19.5}$ eV, the distance is increased by about fifteen times. In Fig. 2, we recalculate the energies with $h = 0.8$.

Using the mean energy approximation, we can also calculate the evolved spectrum of the primary S^0 spectrum observed on Earth given the initial spectrum at the source (where all the particles are injected at one time). With the source at $z = 0.54$ (the source distance for 3C 147) and the initial spectrum having a power law behavior of E^{-2}, the evolved spectrum is shown in Fig. 3. We see that even though there is significant attenuation for the S^0 number at 3×10^{20} eV for most of the cases shown, when the overall cross section (which was originally estimated quite conservatively) is reduced by a factor of half, the bump lies very close to the Fly's Eye event. Moreover, taking the Fly's Eye event energy to be 2.3×10^{20} eV which is within 1σ error range, we see that an S^0 can easily account for the Fly's Eye event. For sources such as MCG 8-11-11, S^0s clearly can account for the observed event without upsetting the proton flux at lower energies.

IV. CONCLUSIONS

In conclusion, we have shown that if S^0s are the primary cosmic ray particles at energies above the GZK cutoff, they can propagate at least fifteen to thirty times longer through the CBR than the nucleons, for the same amount of total energy loss. Thus, if there exists an acceleration mechanism which can generate an adequate high-energy spectrum, S^0s can serve as messengers of the phenomena which produce them, allowing MCG 8-11-11 Seyfert galaxy or 3C 147 quasar to be viable sources for these ultrahigh energy cosmic rays.

Although much of the relevant hadronic physics in the atmospheric shower development will be similar to that for the proton primaries, some subtle signatures of an S^0 primary are still expected. Because an S^0 is expected to have a cross section on nucleons or nuclei somewhere between 1/10 and 4/3 of the p-p cross section, the depth of the shower maximum may be a bit larger than that due to the proton. Furthermore, because it is about twice as massive as the proton, it deposits its energy a bit more slowly than a proton, broadening the distribution of the shower. There may be further signatures in the shower development associated with the different branching fractions to mesons, but we leave that numerical study for the future.

A prediction of this scenario which can be investigated after a large number of UHE events have been accumulated is that UHE cosmic rays primaries point to their sources. If there are a limited number of sources, multiple UHE events should come from the same direction. Also, the UHE cosmic-ray spectrum from each source should exhibit a distinct energy dependence with a cutoff (larger than the GZK cutoff) at an energy which depends on the distance to the source. The systematics of the spectrum in principle could reveal information about both the S^0 mass and the primary spectrum of the source accelerator.

We note that the proposal that the uhecron is a new, massive, stable hadron is more general than the S^0 hypothesis. If the uhecron is produced in "fixed-target" mode by the collision of a high-energy proton beam with a stationary proton, then at 10^{21} eV primary proton energy, the available center-of-mass energy is $\sqrt{2m_pE_p} \simeq 1400\,\mathrm{TeV}$. Thus, the uhecron could conceivably have a mass of several hundred TeV before its production would start to be kinematically suppressed. Extensions of the standard model often predict new colored particles in the 1 to 10 TeV range. In some instances the lightest of the new particles would be stable either because of an accidental symmetry or because of a new conserved quantum number. The new, massive, colored particle would form a color singlet by binding to light quarks. Such a particle would propagate through the CBR without significant energy loss, because the threshold energy for pion production in inelastic collisions is proportional to its mass. The possibility that UHE cosmic rays could be be a window to mass scales beyond the reach of accelerators is exciting, however present observations may already rule out such a possibility. The very massive hadron should have a distinctive signature in the shower development in the atmosphere. Although it is strongly interacting, its fractional energy loss *per collision* in the atmosphere is only of order $1\,\mathrm{GeV}/M$, where M is the mass of the heavy hadron.[12] Thus if the observed energy deposition spectrum of the UHE events is indeed typical of a nucleon or nucleus, as present evidence suggests, we cannot identify the uhecron with a very massive stable hadron.

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[12]It is the momentum of the light partons that is redistributed in a hadronic collision. In the infinite momentum frame for the heavy hadron, the light partons have the same velocity as the heavy parton, but their mass is only of order Λ_{QCD}. Therefore, the fractional momentum carried by light partons is of order $1\,\mathrm{GeV}/M$. Of course a hard collision with the heavy quark would produce a large fractional energy loss, but the cross section for such a collision is small, of order α_s^2/E^2.

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