# Are ultrahigh energy cosmic rays a signal for supersymmetry?

Daniel J. H. Chung*
*Department of Physics and Enrico Fermi Institute, The University of Chicago, Chicago, Illinois 60637*
*and NASA/Fermilab Astrophysics Center, Fermi National Accelerator Laboratory, Batavia, Illinois 60510*

Glennys R. Farrar†
*Department of Physics and Astronomy, Rutgers University, Piscataway, New Jersey 08855*

Edward W. Kolb‡
*NASA/Fermilab Astrophysics Center, Fermi National Accelerator Laboratory, Batavia, Illinois 60510*
*and Department of Astronomy and Astrophysics and Enrico Fermi Institute, The University of Chicago, Chicago, Illinois 60637*
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We investigate the possibility that cosmic rays of energy larger than the Greisen-Zatsepin-Kuzmin cutoff are not nucleons, but a new stable, massive, hadron that appears in many extensions of the standard model. We focus primarily on the $S^0$, a $uds$-gluino bound state. The range of the $S^0$ through the cosmic background radiation is significantly longer than the range of nucleons, and therefore $S^0$'s can originate from sources at cosmological distances. [S0556-2821(98)04808-5]

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## I. INTRODUCTION

The 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 $Z$, $(E/10^{18}\ \mathrm{eV}) = Z(R/\mathrm{kpc})(|\vec{B}|/\mu\mathrm{G})$, 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 event [2] is consistent with its identification as a hadron but not as a photon [4]. Ultrahigh-energy[1] (UHE) events observed in air shower arrays have a muonic composition indicative of hadrons [1]. The problem is that the propagation of hadrons—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, $N\gamma \rightarrow N\pi$, 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. Indeed, even photons of such high energies have a mean free path of less than 10 Mpc due to scattering from cosmic background radiation (CBR) and radio photons [5]. Thus unless the primary is a neutrino, the sources must be nearby (less than about 50 Mpc). This would present a severe problem, because unusual sources such as quasars and Seyfert galaxies typically are beyond this range.

However, the primary cannot be a neutrino because the neutrino interaction probability in the atmosphere is very small. This would imply an implausibly large primary flux, and worse yet, would imply that the depths of first scattering would be uniformly distributed in column density, contrary to observation. The suggestion that the neutrino cross section grows to a hadronic size at UHE [4] has recently been shown to be inconsistent with unitarity and constraints from lower energy particle physics [6].

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. Possible candidate sources within 10° of the UHE cosmic ray observed by the Fly's Eye [2] were studied in Ref. [5].[2] The quasar 3C 147 and the Seyfert galaxy MCG 8-11-11 are attractive candidates. Lying within the $1\sigma$ error box of the primary's incoming direction, the quasar 3C 147 has a large radio luminosity ($7.9\times10^{44}$ 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 rotation measure RM $=-1510\pm50$ rad m$^{-2}$, indicative of a large magnetic field over large distances. It is noteworthy that this source is within the error box of a UHE event seen by the Yakutsk detector. However, 3C 147 lies at a redshift of about $z=0.545$, well beyond $z<0.0125$ adopted in Ref. [5] as the distance upper limit for the source of UHE proton primaries.

*Electronic mail: djchung@yukawa.uchicago.edu
†Electronic mail: farrar@farrar.rutgers.edu
‡Electronic mail: rocky@rigoletto.fnal.gov

[1] We use the term ultrahigh energy to mean energies beyond the GZK cutoff (discussed below) which can be taken to be $10^{19.6}$ eV.

[2] Ten degrees is taken as the extreme possible deflection angle due to magnetic fields for a proton of this energy.

 

Just outside the $2\sigma$ 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 ($4.6\times10^{44}$ erg s$^{-1}$ in the 20–100 keV region and $7\times10^{46}$ erg s$^{-1}$ in the 0.09–3 MeV region). At a redshift of $z=0.0205$, it is much closer than 3C 147, but it is still too distant for the flux to be consistent with the observed proton flux at lower energies [5].

Briefly stated, the problem is that there are no known candidate astronomical sources within the range of protons, neutrons, nuclei, or even photons. Yet there are good candidate sources at 100–1000 Mpc. In this paper we propose that the answer to this cosmic-ray conundrum may be that UHE cosmic rays are not known particles 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 development and muonic content suggests a strongly interacting primary.

(2) The uhecron is stable or very long lived: Clearly if the particle originates from 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 later in the paper, but this general feature can be understood from simple kinematics. In $U\gamma\rightarrow U\pi$, the threshold for pion production is $s_{\min}=m_U^2+m_\pi^2+2m_Um_\pi$. 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 $T=2.4\times10^{-4}$ eV is the temperature of the CBR), $s\simeq m_U^2+4E_\gamma E_U$. Thus, the threshold for pion production, $s\geqslant s_{\min}$, results in the limit $E_U\gtrsim m_\pi m_U/(2E_\gamma)$. More generally, the threshold for producing a resonance of mass $M_R=M_U+\Delta$ is $E_U=\Delta m_U/(2E_\gamma)$. For $E_\gamma=3T$, and if the uhecron is the proton, the threshold for pion photoproduction is $E_U\approx10^{20}$ 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 will be of order $m_\pi/m_U$, a massive uhecron will lose energy via pion-photoproduction at a slower rate than a lighter particle. Another potential bonus if the cosmic ray is not a neutron or a proton is that the cross section for $U\gamma\rightarrow U\pi$ near 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 $\Delta(1232)$ is in the pion-nucleon channel.

(4) We will assume that 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 because it will be unaffected by intergalactic and galactic magnetic fields, its arrival direction on the sky will point back to its source. Third, there will be no energy losses due to synchrotron or bremsstrahlung radiation. Of course because neutral particles 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, they may be produced 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 1.9–2.3 GeV—see below) is the uhecron instead of the proton, as first proposed in Ref. [7]. 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 [7]. Remarkably, this particle is not experimentally excluded. The light gluino required in this scenario would have escaped detection. Experimental limits and signatures are discussed in [7] and the reviews of Farrar [8,9].

If UHE cosmic rays are $S^0$'s, 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.

While the main thrust of this paper is an investigation into the scenario where the $S^0$ is the uhecron, most of our analysis can also be applied to the case where the uhecron is much more massive than assumed for the $S^0$. Extensions of the standard model often predict new heavy, e.g., multi-TeV, colored particles which in some instances have a conserved or almost-conserved quantum number. Bound to light quarks these form heavy hadrons, the lightest of which would be stable or quasistable. Such a particle would propagate through the CBR without significant energy loss because the threshold energy for inelastic collisions is proportional to its mass. Some mechanisms for uhecron production discussed below would be applicable for a new very massive hadron. However, such a particle probably would not be an acceptable candidate for the uhecron because its interaction in the atmosphere is quite different from that of nucleons, nuclei, or an $S^0$. Although it is strongly interacting, its fractional energy loss per collision in the Earth's atmosphere is only of order ($1\ \mathrm{GeV}/M$), where $M$ is the mass of the heavy hadron.[3] Thus if the uhecron energy deposition spectrum is indeed typical of a nucleon or nucleus, as present evidence suggests, we cannot identify the uhecron with a very massive stable hadron. The maximum uhecron mass consistent with observed shower properties is presently under investigation [10].

## II. PRODUCTION OF UHE $S^0$'s

We first address the question of whether there is a plausible scenario to produce UHE $S^0$'s. This is a tricky question,

[3]In the infinite momentum frame for the heavy hadron, this is the fractional momentum carried by light partons since they have the same velocity as the heavy parton, but their mass is of order $\Lambda_{QCD}$. It is the momentum of these light partons which is redistributed in a hadronic collision. 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: $\approx\alpha_s^2/E^2$.

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 UHE protons into UHE $S^0$'s. Our intent is not to establish the viability of any particular mechanism but to see that finding a satisfactory mechanism is not dramatically more difficult than it is for protons.

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 $R_p$'s, the $uud$-gluino state whose mass is about 200 MeV above the $S^0$. The $R_p$ decays to an $S^0$ and a $\pi^+$,[4] 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 $R_p$'s 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 $\alpha$ is the Regge intercept of the supersymmetric (SUSY) partner of the Pomeron. Thus, $\alpha=\alpha_P-1/2=\epsilon+1/2$, where $\epsilon \approx 0.1$ is the amount the Pomeron trajectory is above 1 at high energies. Hence, we parametrize the $S^0$ production cross section in a $p$-nucleon collision as $d\sigma/dx = AE_p^{\epsilon}$; $x$ is the ratio of the $S^0$ energy to the incident energy. Parametrizing 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 $\kappa$ is of order 1 (for $\gamma=2$, $\kappa=0.4$). Note that the produced $S^0$'s are distributed according to a spectrum that is a bit flatter than the high energy proton spectrum.

A disadvantage of this ''beam-dump'' $S^0$ production mechanism is the suppression factor of about $AE^{\epsilon}/\sigma_{pp}$, where $\sigma_{pp}$ is the proton-proton total diffractive cross section. This suppression could be of order $10^{-1}$–$10^{-2}$ for typical energies. However the produced $S^0$'s enjoy a compensating advantage. The large column densities characteristic of most candidate acceleration regions makes it hard to avoid energy degradation of protons before they escape. That is, $L(n_p\sigma_{pN}+n_e\sigma_{pe}+n_\gamma\sigma_{p\gamma})$ may be much greater than unity. By contrast, $S^0$'s may escape with little or no energy loss. Their electromagnetic interactions are negligible, and analogy with glueball wave functions suggests that $\sigma_{S^0N}$ could be as small as $10^{-1}\sigma_{pN}$ [7]. Thus the emerging $S^0$ and nucleon fluxes could be of the same order of magnitude. This would be necessary for a very distant source such as 3C 147 to be acceptable, since the required particle flux for the detected flux on Earth already pushes its luminosity limit. Assuming that the $3.2\times10^{20}$ eV event of Fly's Eye came during its exposure to 3C 147, the resulting time-averaged flux is 11 eV cm$^{-2}$ s$^{-1}$, which is greater than the x-ray luminosity of 3C 147 [5].

In connection with the ''beam dump'' mechanism, we note that it is possible to have at the source a nucleon flux significantly greater than the $S^0$ flux, and yet at Earth still have a large enough $S^0$ flux to account for the high energy end of the spectrum without being inconsistent with the rest of the observed cosmic ray spectrum. To see that this is possible, suppose as an illustrative example that the $S^0$ spectrum for energies above $10^{20}$ eV is a smooth extrapolation of the proton spectrum at energies below the GZK cutoff; i.e., if $J_p(E)=AE^{-3}$ for $E<10^{19.6}$ eV, then $J_{S^0}(E)=AE^{-3}$ for $E>10^{20}$ eV. Denoting the $S^0$-to-proton suppression factor by $\eta$, the proton flux for $E>10^{20}$ eV is then $J_p(E)=\eta^{-1}AE^{-3}$. Protons of energy greater than the GZK cutoff (here taken to be $10^{19.6}$ eV) will bunch up in the decade in energy below the GZK cutoff [11,12]. The total number in the pileup region will receive a contribution from protons from the source above the GZK cutoff as well as those originally in the pileup region. With $\eta=10^{-2}$, there will be equal contributions from the pileup protons and the protons originally below the source. The statistics of the number of events with energy above $10^{18.5}$ eV is too poor to exclude this scenario; indeed there is some indication of a bump in the spectrum in this region [1].

Note that even for a point source as far away as 1200 Mpc (e.g. 3C 147), the required flux of high energy protons at the accelerator is not unacceptable. For instance extrapolating the spectrum as $7.36\times10^{18}E^{-2.7}/(\text{eV m}^2\text{ sr s})^{-1}$ and using our pessimistic efficiency for $S^0$ production (factor of 1/100) requires the high energy proton luminosity of the source to be $\sim10^{47}$ ergs/s. This is indeed a high value, but not impossible.

Another possible mechanism of high energy $S^0$ production is the direct acceleration of charged light SUSY hadrons (mass around 2–3 GeV), such as $R_p$ and $R_\Omega$, whose lifetime is about $2\times10^{-10}$–$2\times10^{-11}$ sec [7]. Because of the large time-dilation factor ($E/m\approx10^{11}$), 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. In that case the flux of the resulting $S^0$ will have the same spectrum as the protons, differing in magnitude by a factor of order unity, which depends on the amount of SUSY hadrons making up the statistical mixture. Conventional shock wave acceleration mechanisms probably require a too long time scale for this mechanism to be feasible (e.g., Ref. [14]). However, some electromagnetic ''one push'' mechanisms similar to the one involving electric fields around pulsars [15] may allow this kind of acceleration if the electric field can be large enough. It is certainly tantalizing that the time scale of the short time structure of pulsars and gamma ray bursts is consistent with the scale implied by the time-dilated lifetime of charged $R$-baryons.

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 $S^0$'s 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 col-

[4]The decay $R_p \to S^0\pi$ was the subject of an experimental search [13]. However the sensitivity was insufficient in the mass and lifetime range of interest [$m(R_p)=2.1$–$2.5$ GeV, $\tau(R_p)=2\times10^{-10}$–$2\times10^{-11}$; see [7]] for a signal to have been expected.

lision), they may be able 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.

A final possibility is the decay of long-lived superheavy relics of the big bang, which would produce all light particles present in the low-energy world, including the $S^0$. For instance if such relics decay via quarks which then fragment, as in models such as Ref. [16], the $S^0$/nucleon ratio is probably in the range $10^{-1}$–$10^{-2}$ based on a factor of about 10 suppression in producing a 4-constituent rather than 3-constitutent object, and possibly some additional suppression due to the $S^0$'s higher mass.[5]

Of the scenarios considered above, only the last two are conceivably relevant for a superheavy (0.1–1000 TeV) uhecron. Although the energy in $p$-nucleon collisions ($\sqrt{s} = \sqrt{2E_p m_p} \sim 10^3$ TeV for a primary proton energy of $10^{21}$ eV) is sufficient for superheavy particle production, the production cross section is too small for the ''beam dump'' mechanism to be efficient.[6] Also, the direct acceleration mechanism is not useful for a superheavy uhecron unless it is itself charged or is produced in the decay of a sufficiently long-lived charged progenitor. Even if a sufficient density of superheavy hadrons could be generated in spite of the small production cross section, the time scale required for the early stages of acceleration could be too long since it is proportional to $\beta^2$. This leaves the decay of a superheavy relic (either a particle or cosmic defect) as the most promising source of uhecrons if their masses are greater than tens of GeV.

## 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 Refs. [12] and [19]. 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[7]), 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 $p\gamma \rightarrow \pi^0 p$ and $p\gamma \rightarrow \pi^+ n$. 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

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

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

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

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 the differential cross section in the proton frame,[8] 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 $\beta$ to the proton frame

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

where $\theta$ is the angle that the photon direction makes with respect to the boost direction. Integrating Eq. (3) over the solid angle[9] and taking the ultrarelativistic limit, we find

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

For $\Delta E$, the energy loss per event in the CBR frame, we can write

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

where $p_r$, which may depend on $E_\gamma$ and $\cos\theta$, is the recoil momentum of the proton and $\theta$ is the angle between the

[5]After our work was completed, Ref. [17] appeared with an estimate of the production of gluino-hadrons from the decay of cosmic necklaces. Note that their pessimism regarding the light gluino scenario is mostly based on arguments which have been rebutted in the literature (see for example Refs. [18] and [9]).

[6]The cross section is proportional to the initial parton density at $x \sim M_U/\sqrt{s}$ times the parton-level cross section, which scales as $M_U^{-2}$.

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

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

[9]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.

incoming photon direction and the outgoing proton direction. Putting all these together, the energy loss rate due to scattering given by Eq. (1) becomes

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

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

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

where $q^2=m_p(m_p+E_\gamma)-m_\pi^2/2$. When the photon energy $E_\gamma$ is approximately at the threshold energy of $m_\pi+m_\pi^2/2m_p$ and the proton recoils in the direction $\theta=0$, the recoil momentum is about $m_\pi$. 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.[10]

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:

$$\sigma(\text{pion})=2\sigma_{1\pi}\Theta\left(E_\gamma-m_\pi-\frac{m_\pi^2}{2m_p}\right)+2\sigma_{\text{multipion}}$$

$$\sigma_{1\pi}=\frac{4\pi}{p_{\text{c.m.}}^2}\left[\frac{m_\Delta^2\Gamma(\Delta\to\gamma p)\Gamma(\Delta\to\pi P)}{(m_\Delta^2-s)^2+m_\Delta^2\Gamma_{\text{tot}}^2}\right]+\sigma_{\text{nonres}}$$

$$\Gamma(\Delta\to Xp)=\frac{p_{\text{c.m.}}^X\omega_X}{8m_\Delta\sqrt{s}}$$

$$\Gamma_{\text{tot}}=\frac{p_{\text{c.m.}}^\pi}{\sqrt{s}}\,\frac{2m_\Delta^2\overline{\Gamma}_{\text{tot}}}{\sqrt{[m_\Delta^2-(m_\pi+m_p)^2][m_\Delta^2-(m_p-m_\pi)^2]}}$$

$$\sigma_{\text{nonres}}=\frac{1}{16\pi s}\,\frac{\sqrt{[s-(m_p+m_\pi)^2][s-(m_p-m_\pi)^2]}}{(s-m_p^2)}\times|\mathcal{M}(p\gamma\to\pi p)|^2$$

$$\sigma_{\text{multipion}}=a\tanh\left(\frac{E_\gamma-E_{\text{multi}}}{m_\pi}\right)\Theta(E_\gamma-E_{\text{multi}}) \quad (8)$$

where $\omega_X$ 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 is given as usual by

$$p_{\text{c.m.}}^X=\sqrt{\frac{[s-(m_p+m_X)^2][s-(m_p-m_X)^2]}{4s}}, \quad (9)$$

and $\sigma_{\text{multipion}}$ is a crude approximation[11] for the contribution from the multipion production whose threshold is at $E_{\text{multi}}=2(m_\pi+m_\pi^2/m_p)$. The $\sigma_\pi$ component of the cross section is fit[12] to the $p\gamma\to n\pi^0$ data of Ref. [20], while the amplitude $a$ for $\sigma_{\text{multipion}}$ is estimated from the $p\gamma\to Xp$ data for energies $E_\gamma\gtrsim 0.6$ GeV. The numerical values of the parameters resulting from the fit are $(\omega_\gamma\omega_\pi)=0.086\ \text{GeV}^4$, $|\mathcal{M}(p\gamma\to\pi p)|=0.018$, $\overline{\Gamma}_{\text{tot}}=0.111$ GeV, $m_\Delta=1.23$ GeV, and $a=0.2$ mb. The factor of 2 multiplying $\sigma_\pi$ accounts for the two reactions $p\gamma\to\pi^0p$ and $p\gamma\to\pi^+n$, since a neutron behaves, to a first approximation, just like the proton. For example, the dominant pion-photoproduction reactions involving neutrons are $n\gamma\to\pi^0n$ 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. [21] (as done in Ref. [12]), we use[13]

[10]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.

[11]The functional form was chosen to account for the shape of the cross section given in Ref. [20].

[12]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 $\chi_{16}^2\sim 50$ (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.

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

$$\frac{d\sigma(\text{pair})}{dQ\,d\eta}=\Theta(E_\gamma-2m_e)\frac{4\alpha^3}{E_\gamma^2}\frac{1}{Q^2}\Bigg\{\ln\left(\frac{1-w}{1+w}\right)\Bigg[\left(1-\frac{E_\gamma^2}{4m_e^2\eta^2}\right)\times\left(1-\frac{1}{4\eta^2}+\frac{1}{2\eta Q}-\frac{1}{8Q^2\eta^2}-\frac{Q}{\eta}+\frac{Q^2}{2\eta^2}\right)+\frac{E_\gamma^2}{8m_e^2\eta^4}\Bigg]$$
$$+w\Bigg[\left(1-\frac{E_\gamma^2}{4m_e^2\eta^2}\right)\left(1-\frac{1}{4\eta^2}+\frac{1}{2\eta Q}\right)+\frac{1}{\eta^2}\left(1-\frac{E_\gamma^2}{2m_e^2\eta^2}\right)(-2Q\eta+Q^2)\Bigg]\Bigg\}, \tag{10}$$

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 the Hubble expansion. We assume a matter-dominated, flat Friedmann-Robertson-Walker (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 a FRW universe. Note also that we need to set the Hubble parameter $h$ (where the Hubble constant is $100h$ km s$^{-1}$ Mpc$^{-1}$) in our calculation because the conversion 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 $\gamma S^0$ interaction is larger than the proton-$\Delta$ mass splitting, leading to a further increase in the attenuation length of the primary. Perhaps most importantly, the mass of $S^0$ being about 2 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 the leading approximation is proportional to $p_r/m_p$ while $p_r$ has a maximum value of about $m_\pi$. Replacement of $m_p\rightarrow m_{S^0}$ 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 $\gamma$ (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

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

where $y=m_\pi m_p/(E_pT)$, 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 $R_\Lambda$ and $R_\Sigma$ [7] (whose constituents are those of the usual $\Lambda$ and $\Sigma$ baryons, but in a color octet state, coupled to a gluino [22]). 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 MeV above that of the $S^0$ and the states with quarks contributing spin 1/2 have masses of about 815–890 MeV above that of the $S^0$. If we require that the photino be a significant dark matter component so $1.3\leq M_{R^0}/m_\gamma\leq 1.6$ according to Ref. [23], and take the mass of $R^0$ to be about $1.6-1.8$ GeV as expected, then $m_\gamma$ lies in the range 0.9–1.3 GeV. If we assume that $S^0$ is minimally stable, we have $m_{S^0}\approx m_p+m_\gamma$, resulting in $m_{S^0}$ in the range 1.9–2.3 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[14]) with an energy of $10^{22}$ eV, $10^{21}$ eV, and $10^{20}$ eV. To explore the interesting mass range, we have 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 30 times, while for those arriving with $10^{19.5}$ eV, the distance is increased by about 15 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 $3\times10^{20}$ eV for most of the cases shown, when the overall

[14]Marked are the luminosity distances $d_L=H_0^{-1}q_0^{-2}[zq_0+(q_0-1)(\sqrt{2q_0z+1}-1)]$ where the deceleration parameter $q_0$ is 1/2 in our $\Omega_0=1$ universe.

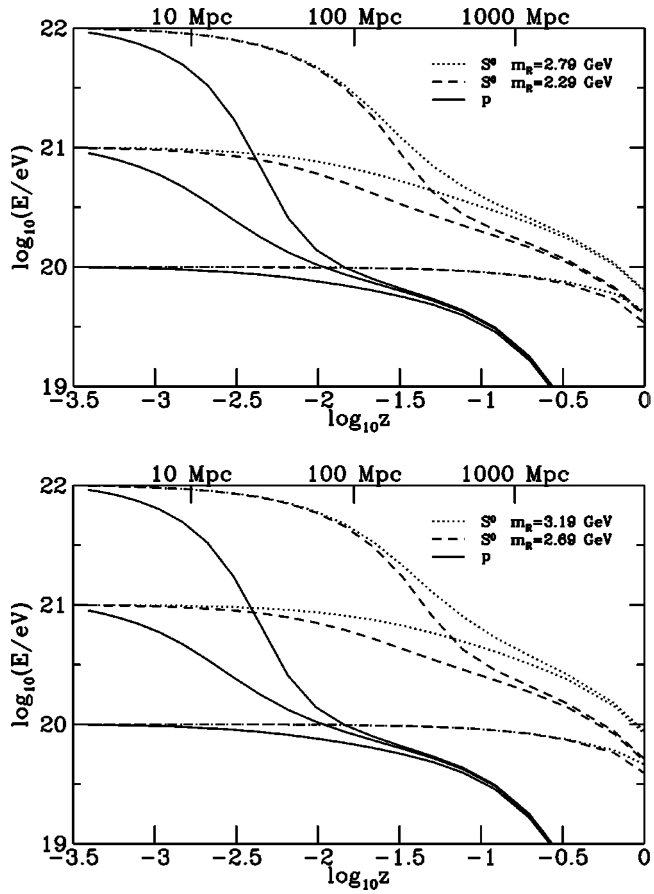


FIG. 1. Primary particle's energy as it would be observed on Earth today if it were injected with various energies ($10^{22}$ eV, $10^{21}$ eV, and $10^{20}$ eV) at various redshifts. The distances correspond to luminosity distances. The mass of $S^0$ is 1.9 GeV in the upper plot while it is 2.3 GeV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

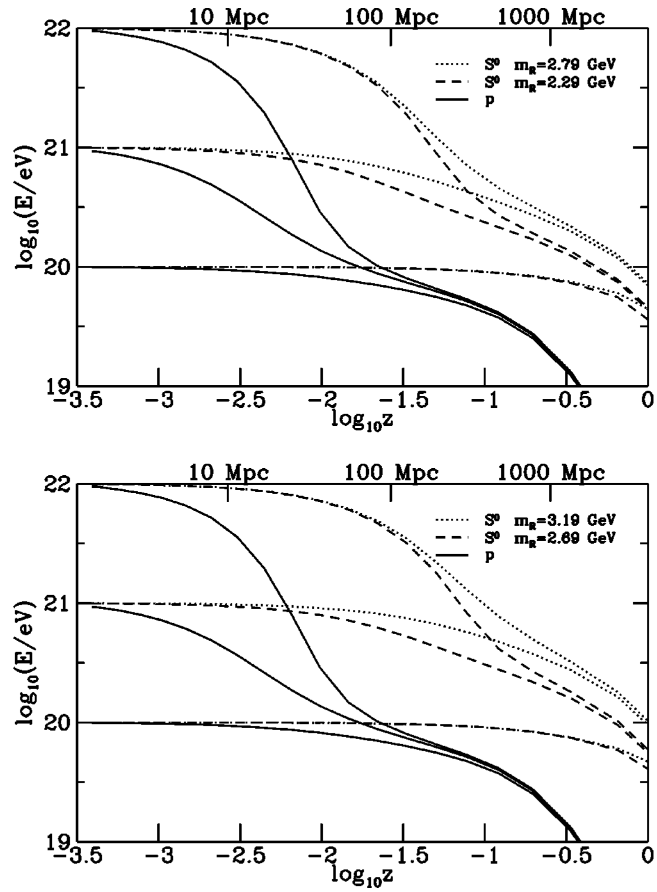


FIG. 2. Same as Fig. 1 except with the Hubble constant equal to 80 km s$^{-1}$ Mpc$^{-1}$.

cross section (which was originally estimated quite conservatively) is reduced by a factor of one-half, the bump lies very close to the Fly's Eye event. Moreover, taking the Fly Eye's event energy to be $2.3\times10^{20}$ eV which is within a $1\sigma$ error range, we see that the $S^0$ can easily account for the Fly Eye's event. For sources such as MCG 8-11-11, $S^0$ clearly can account for the observed event without upsetting the proton flux at lower energies.

## IV. CONCLUSIONS

We have considered the suggestion that the very long-lived or stable new hadron called $S^0$, a $uds$-gluino bound state predicted in some supersymmetric models, can account for the primary cosmic ray particles at energies above the GZK cutoff. We noted ways that conventional acceleration mechanisms might result in acceptable fluxes of high energy $S^0$'s. We also found that the $S^0$ can propagate at least 15–30 times longer through the CBR than do nucleons, for the same amount of total energy loss. Thus, if $S^0$'s exist and there exists an acceleration mechanism which can generate an adequate high-energy spectrum, $S^0$'s can serve as messengers of the phenomena which produce them, allowing the 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 maxi-

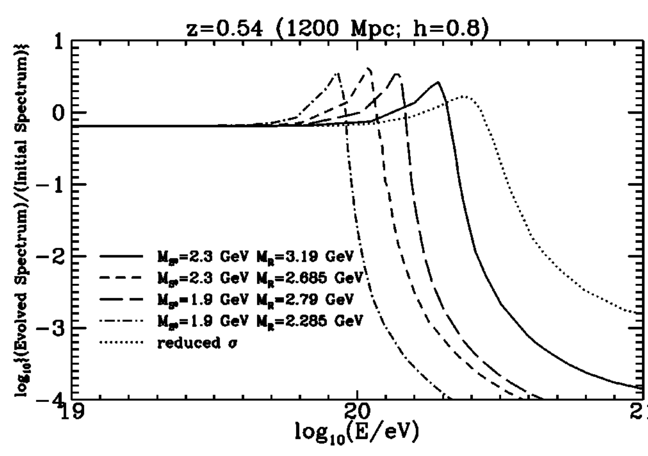


FIG. 3. An initial $S^0$ injection spectrum having a power law form of $E^{-2}$ is evolved through the particle's interaction with the CBR during its 1200 Mpc travel to Earth. The masses of the $S^0$ and its associated resonance are shown. The curve labeled reduced $\sigma$ has the same mass parameters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

mum 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 ray 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 masses of supersymmetric particles and the primary spectrum of the source accelerator.

We noted that the mass range for a new hadron which can account for the observed properties of UHE cosmic ray events is limited: it must be at least 2 GeV in order to evade the GZK bound, yet small enough that the atmospheric shower it produces will mimic an ordinary hadronic shower.

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[1] J. Linsley, Phys. Rev. Lett. **10**, 146 (1963); N. Hayashida *et al.*, *ibid.* **73**, 3491 (1994); D. J. Bird *et al.*, Astrophys. J. **424**, 491 (1994).
[2] D. J. Bird *et al.*, Astrophys. J. **441**, 144 (1995).
[3] P. L. Bierman, J. Phys. G **23**, 1 (1997).
[4] F. Halzen *et al.*, Astropart. Phys. **3**, 151 (1995).
[5] J. W. Elbert and P. Sommers, Astrophys. J. **441**, 151 (1995); G. Sigl *et al.*, Astropart. Phys. **2**, 401 (1994).
[6] G. Burdman *et al.*, hep-ph/9709399.
[7] G. R. Farrar, Phys. Rev. Lett. **76**, 4111 (1996).
[8] G. R. Farrar, hep-ph/9710277.
[9] G. Farrar, Phys. Rev. D **51**, 3904 (1995).
[10] I. F. Albuquerque, G. R. Farrar, and E. W. Kolb (in preparation).
[11] C. T. Hill and D. N. Schramm, Phys. Rev. D **31**, 564 (1985).
[12] J. Geddes, T. Quinn, and R. Wald, astro-ph/9506009.
[13] I. Albuquerque *et al.*, Phys. Rev. Lett. **78**, 3252 (1997).
[14] J. P. Rachen and P. L. Biermann, Astron. Astrophys. **272**, 161 (1993).
[15] K. S. Cheng, C. Ho, and M. Ruderman, Astrophys. J. **300**, 500 (1986); Ref. [3].
[16] V. A. Kuzmin and V. A. Rubakov, astro-ph/9709187.
[17] V. Berezinsky and M. Kachelreiss, hep-ph/9709485.
[18] R. Plaga, Phys. Rev. D **51**, 6504 (1995).
[19] J. W. Cronin, in *Cosmic Rays Above $10^{19}$ eV*, Proceedings of the International Workshop, Paris, France, 1992, edited by J. Cronin *et al.* [Nucl. Phys. B (Proc. Suppl.) **28B**, 234 (1992)].
[20] *Total Cross-sections for Reactions of High Energy Particles*, edited by H. Schopper (Springer-Verlag, New York, 1988).
[21] R. Jost, J. M. Luttinger, and M. Slotnick, Phys. Rev. **80**, 189 (1950).
[22] F. Bucella, G. R. Farrar, A. Pugliese, Phys. Lett. **153B**, 311 (1985).
[23] D. J. H. Chung, G. R. Farrar, and E. W. Kolb, Phys. Rev. D **56**, 6096 (1997).