Probing Split Supersymmetry with Cosmic Rays

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Abstract
A striking aspect of the recently proposed split supersymmetry is the existence of heavy gluinos which are metastable because of the very heavy squarks which mediate their decay. In this paper we correlate the expected flux of these particles with the accompanying neutrino flux produced in inelastic $pp$ collisions in distant astrophysical sources. We show that an event rate at the Pierre Auger Observatory of approximately $1 \, \text{yr}^{-1}$ for gluino masses of about 500 GeV is consistent with existing limits on neutrino fluxes. Such an event rate requires powerful cosmic ray engines able to accelerate particles up to extreme energies, somewhat above $5 \times 10^{13}$ GeV. The extremely low inelasticity of the gluino-containing hadrons in their collisions with the air molecules makes possible a distinct characterization of the showers induced in the atmosphere. Should such anomalous events be observed, we show that their cosmogenic origin, in concert with the requirement that they reach the Earth before decay, leads to a lower bound on their proper lifetime of the order of 100 years, and consequently, to a lower bound on the scale of supersymmetry breaking, $\Lambda_{\text{SUSY}} > 2.6 \times 10^{11}$ GeV. Obtaining such a bound is not possible in collider experiments.

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I. GENERAL IDEA

The standard model (SM) of particle physics has had outstanding success in describing all physical phenomena up to energies $\sim 500$ GeV [1]. Nonetheless, there is a general consensus on that it is not a fundamental theory of nature (apart from the fact that the SM does not include gravity): with no new physics between the energy scale of electroweak unification ($M_W \sim 10^2$ GeV) and the vicinity of the Planck mass ($M_{Pl} \sim 10^{19}$ GeV) the higgs mass must be fine-tuned to an accuracy of order $(M_W/M_{Pl})^2$ to accommodate this enormous desert. The leading contender for the elaboration of the desert has been the supersymmetric extension of the SM [2]. Supersymmetry (SUSY) posits a “complete democracy” between integral and half-integral spins, implying the existence of many as-yet-undiscovered superpartners. Thus, if SUSY can serve as a theory of low energy interactions, it must be a broken symmetry. The most common assumption is that the minimal low energy effective supersymmetric theory (MSSM) has a breaking scale of order $\Lambda_{SUSY} \sim 1$ TeV, thus avoiding ’t Hooft naturalness problem with the higgs mass.

MSSM has a concrete advantage in embedding the SM in a grand unified theory: the supersymmetric beta functions for extrapolating the measured strengths of the strong, electromagnetic, and weak couplings lead to convergence at a unified energy value of the order $M_{GUT} \sim 10^{16}$ GeV [3]. The model, however, is not free of problems. In particular, dimension four $R$-parity violating couplings in the superpotential yield unacceptably large proton decay rates and neutrino masses. This can be readily solved by imposing $R$-parity conservation, which as a byproduct ensures the stability of the lightest SUSY particle, making it a possible candidate for cold dark matter [4]. However, there are other problems in the MSSM: dimension five operators tending to generate excessive proton decay; new CP-violating phases which require suppression for agreement with limits on electric dipole moment limits; and excessive flavor violations, due e.g., to the absence of a complete flavor degeneracy in the Kähler potential of minimal supergravity [5].

Of course, the fine-tuning involved in accommodating the above constraints is miniscule in comparison to that required in generating a cosmological constant that satisfies ’t Hooft naturalness. Recent experimental data [6] strongly indicate that the universe is expanding in an accelerating phase, with an effective de Sitter constant $H$ that nearly saturates the upper bound given by the present-day value of the Hubble constant, i.e., $H \lesssim H_0 \sim 10^{-33}$ eV. According to the Einstein field equations, $H^2$ provides a measure of the scalar curvature of the space and is related to the vacuum energy density, $\epsilon_4$, according to $M_{Pl}^2 H^2 \sim \epsilon_4$. However, the “natural” value of $\epsilon_4$ coming from the zero-point energies of known elementary particles is found to be at least $\epsilon_4 \sim \Lambda_{SUSY}^4$, yielding $H \gtrsim 10^{-3}$ eV. The failure of ’t Hooft naturalness then centers on the following question: why is the vacuum energy determined by the Einstein field equations 60 orders of magnitude smaller than any “natural” cut-off scale in effective field theory of particle interactions, but not zero? Nowadays, the only existing framework which can address aspects of this question is the anthropic approach [7].

Very recently Arkani–Hamed and Dimopoulos (AD) have looked at SUSY from a different angle [8]. In their model, the scale of SUSY breaking is pushed to a very high energy (say, $\Lambda_{SUSY} \sim 10^{13}$ GeV) and a higgs mass of order TeV is recovered by invoking fine tuning. For this breaking scale, the bosonic superpartners are heavy, while the extra fermions retain TeV-scale masses thanks to protection by chiral symmetry. (We follow Giudice and Romanino [9] in adopting the designation “split SUSY” for the AD model). This scenario preserves the achievements of the MSSM while resolving the problems mentioned above. In particular,
analyses of one loop [10] and two loops [9] running of the RG equations, show that the AD scenario preserves unification of couplings. Moreover, aside from the light higgs tuning, the other flaws inherent to the MSSM elegantly disappear when the scalar superpartners decouple.

The AD model can be discussed in the same anthropic framework adopted for examination of the cosmological constant. Recent investigations in String Theory have applied a statistical approach to the enormous “landscape” of vacua present in the theory [11]. Among this vast number of metastable vacua, there can be small subset $O(10^{40})$ exhibiting low scale SUSY breaking, a TeV-scale higgs, as well as the remaining traditional MSSM physics [12]. However, the fine tuning required to achieve a small cosmological constant implies the need of a huge number of vacua, far more than the $O(10^{40})$ characterizing low-scale SUSY breaking [13]. Remarkably, if one posits high-scale SUSY breaking and superpartners widely separated in mass, $O(10^{200})$ vacua become available, enough to fine tune both the cosmological constant and the higgs mass.

It is therefore instructive to explore how drastically the AD scenario can change the phenomenology of conventional MSSM. Prospects for probing split SUSY at the LHC [14] as well as in dark matter searches [15] have been recently developed. In what follows we show that cosmic ray data may also provide important information about the AD scenario. As a principal result of this paper, we will delineate conditions under which one can set a lower bound on the SUSY breaking scale in a region of parameter space far beyond that probed at the LHC.

An intriguing prediction in this scenario, which represents a radical departure from the MSSM, is the longevity of the gluino [16]. As mentioned above, in split SUSY the squarks are very massive and so gluino decay via virtual squarks becomes strongly suppressed, yielding a $\tilde{g}$ lifetime of the order of [8]

$$\tau_0 \approx \frac{64\pi^3\Lambda_{SUSY}^4}{M_{\tilde{g}}^5} \approx 10^7 \left(\frac{\text{TeV}}{M_{\tilde{g}}}\right)^5 \left(\frac{\Lambda_{SUSY}}{10^{13} \text{ GeV}}\right)^4 \text{yr}$$

where $M_{\tilde{g}}$ is the gluino mass. Very strong limits on heavy isotope abundance in turn require the gluino to decay on Gyr time scales, leading to an upper bound for the scale of SUSY breaking $O(10^{13})$ GeV [8]. Because of the large mass of the gluino, the threshold for inelastic scattering on the cosmic microwave background is $\gtrsim 10^{14}$ GeV [17], allowing ultrahigh energy gluino-containing hadrons (“$G$’s”) to reach us unimpeded from cosmological distances.

In this work we study the possibility of $G$-detection with cosmic ray observatories. To this end, in Sec. II we discuss the main characteristics of cascades induced by $G$-hadrons and estimate the sensitivity of the Pierre Auger Observatory (PAO). After that, in Sec. III, we correlate the expected flux of $G$-hadrons with the accompanying neutrino flux produced in inelastic $pp$ collisions in distant astrophysical sources. We show that an event rate $\approx 1 \text{ yr}^{-1}$ at PAO for gluino masses of about 500 GeV is consistent with existing limits on neutrino fluxes. The actual observation of a few $G$-events will then directly imply a lower bound on $\Lambda_{SUSY}$. The details of this interesting possibility are presented in Sec. IV. Section V contains our conclusions.
II. CHARACTERISTICS OF AIR SHOWERS INITIATED BY G–HADRONS

The interaction of a high energy cosmic ray in the upper atmosphere gives rise to a roughly conical cascade of particles that reaches the Earth’s surface in the form of a giant “saucer”, traveling at nearly the speed of light. In the case of proton– or nucleus–induced cascades, the leading particle and other high energy hadrons (mostly pions) in the shower core readily cascade to lower energies as they interact with the air molecules. Because of the prompt decay of neutral pions, 1/3 of the energy in each of these interactions transits into energetic γ–rays. Electromagnetic subshowers are then initiated: the high energy photons produce pairs that lose energy by bremsstrahlung and ionization before annihilation into a new photon at lower energy. Eventually, the average energy per particle drops below a critical energy $\epsilon_0 \sim 86$ MeV at which point ionization takes over from bremsstrahlung and pair production as the dominant energy loss mechanism. The changeover from radiation losses to ionization losses depopulates the shower, and defines $X_{\text{max}}$, the longitudinal coordinate of maximum multiplicity.

The number of muons (and neutrinos) does not increase linearly with energy, because at higher energy more generations are required to cool the pions to the point where they are likely to decay before interaction. Production of extra generations results in a larger fraction of the energy being lost to the electromagnetic cascade, and hence a smaller fraction of the original energy being delivered to the $\pi^\pm$. The electrons, positrons and photons are thus the most prolific constituents in the thin disk of particles showering towards the ground, and most of the energy (about 90%) is dissipated in the electromagnetic cascade.

By the time they reach the ground, relatively vertical showers have evolved fronts with a curvature radius of a few km, and far from the shower core their constituent particles are well spread over time, typically of the order of a few microseconds. For such a shower both the muon component and a large portion of the electromagnetic component survive to reach the ground. For inclined showers the electromagnetic component is absorbed long before reaching the ground, as it has passed through the equivalent of several vertical atmospheres: 2 at a zenith angle $\theta = 60^\circ$, 3 at $70^\circ$, and 6 at $80^\circ$. In these showers, only high energy muons created in the first few generations of particles survive past 2 equivalent vertical atmospheres. The rate of energy attenuation for muons is much smaller than it is for electrons, thus the shape of the resulting shower front is very flat (with curvature radius above 100 km), and its time extension is very short (less than 50 ns) [18].

The energy lost by a $G$ during collision with nucleons is primarily through hard scattering [19]. This implies a fractional energy loss per collision, $K_{\text{inel}} \approx 1 \text{ GeV}/M_G$. A heuristic justification for this result is as follows: consider the inclusive process $GN \rightarrow GX$. It is a simple kinematic exercise to show that the minimum momentum transfer for $M_X^2, M_G^2 \ll s$ is given by $|t_{\text{min}}| \sim M_G^2, M_X^2/s^2$, where $\sqrt{s}$ is the center-of-mass energy of the $GN$ collision. Kinematics also determine the fractional energy loss, $K_{\text{inel}} = M_X^2/s$. Combining these two results we obtain that $K_{\text{inel}} \approx |t_{\text{min}}|^{1/2}/M_G$. The hard scattering restriction in QCD requires $|t_{\text{min}}| \gtrsim 1 \text{ GeV}^2$, so that for a fiducial $M_G = 500$ GeV, we obtain $K_{\text{inel}} \approx 0.002$. Note that, for $M_G = 50$ GeV, our formula agrees with the result in [19].

In the case of a $G$-induced shower, the very low inelasticity of $G$-air interactions implies the leading particle retains most of its energy all the way to the ground, while the secondary particles promptly cascade to low energies as for any other air shower. This results in an ensemble of mini-showers strung along the trajectory of the leading particle. Since the typical distance between mini-showers is about 10 times smaller than the extent of a
single longitudinal profile, it is not possible to resolve the individual showers experimentally. Instead one observes a smooth envelope encompassing all the mini-showers, which extends from the first interaction all the way to the ground. Monte Carlo simulations have been performed which confirm this phenomenological description (see Fig. 3 in Ref. [19]). The $G$-showers indeed present a distinct profile: (1) there is only a few percent probability for its $X_{\text{max}}$ to be mistaken for that of a proton shower [20] (2) the flatness of the longitudinal development is unique to the extremely low inelasticity of the scattering, and can be easily isolated from background.

The characteristics of $G$-induced showers observed at the ground should also be distinct from those characteristic of proton and nucleus induced showers. Each mini-shower generates a bundle of muons which survive to the ground. Since each bundle is produced at a different slant-depth, the muon component of the shower front exhibits a much more pronounced curvature than what one would expect for a shower of the same energy interacting only in the upper atmosphere. The difference between a proton/nucleus and a $G$ would thus be much more evident in the case of inclined showers for which there is much less electromagnetic contamination. Specifically, a highly inclined $G$-induced shower produces many muon bundles, and so should exhibit a much more curved shower front than a proton/nucleus shower with the same energy and zenith angle.

The experimentally interesting region to search for $G$ events then lies between $70^\circ < \theta < 90^\circ$. The reduction of the solid angle acceptance to larger than $70^\circ$ eliminates the hadronic background. Moreover, there needs to be sufficient pathlength for the $G$, with its low inelasticity, to lose sufficient energy. The mean free path for ultra-high energy $G$‘s is about 30 g cm$^{-2}$ [19], and the pathlength of the atmosphere at 70$^\circ$ is about 100 times this [18], allowing more than 20% of the energy to evolve in the air shower.

With this in mind, the $G$-event rate for a given cosmic ray experiment is found to be

$$
\frac{dN}{dt} = \int_{E_{G,\text{min}}}^{E_{G,\text{max}}} J_G(E_G) A(E_G) dE_G, \tag{2}
$$

where $A(E_G)$ is the hadronic aperture for $\theta > 70^\circ$.

There are two major techniques which can be employed in detecting cosmic ray air showers. The most commonly used detection method involves sampling the shower front at a given altitude using an array of sensors spread over a large area. Sensors, such as plastic scintillators or Čerenkov detectors are used to infer the particle density and the relative arrival times of the shower front at different locations. The muon content is usually sought either by exploiting the signal timing in the surface sensors or by employing dedicated detectors which are shielded from the electromagnetic shower component. Inferring the primary energy from energy deposits at the ground is not completely straightforward, and involves proper modeling of both the detector response and the physics of the first few cascade generations. Another highly successful air shower detection method involves measurement of the longitudinal development of the cascade by sensing the fluorescence light produced via interactions of the charged particles in the atmosphere. Excited nitrogen molecules fluoresce producing radiation in the 300 - 400 nm ultraviolet range, to which the atmosphere is quite transparent. The shower development appears as a rapidly moving spot of light whose angular motion depends on both the distance and the orientation of the shower axis. The integrated light signal is proportional to the total energy deposited in the atmosphere. Fluorescence observations can only be made on clear moonless nights, yielding a duty cycle of about 10%.
Over the next few years, the best observations of the extreme end of the cosmic ray spectrum will be made by the PAO [21], which is actually comprised of two sub-observatories. The Southern site is currently operational and in the process of growing to its final size of $S \approx 3000$ km$^2$. Another site is planned for the Northern hemisphere. The PAO works in a hybrid mode, and when complete, each site will contain 24 fluorescence detectors overlooking a ground array of 1600 water Cherenkov detectors. During clear, dark nights, events are simultaneously observed by fluorescence light and particle detectors, allowing powerful reconstruction and cross-calibration techniques. Simultaneous observations of showers using two distinct detector methods will also help to control the systematic errors that have plagued cosmic ray experiments to date.

For showers at inclination $\theta > 70^\circ$ and energy $E_{\text{sh}} > 10^9$ GeV the probability of detecting an event falling within the physical area is roughly 1. Note that, from the previous discussion, $E_G \gtrsim 5 E_{\text{sh}} \approx 5 \times 10^9$ GeV. The total aperture (2 sites) of the surface array is found to be

$$A_{\text{PAO}} = \int_{70^\circ}^{90^\circ} S \cos \theta \, d\Omega \approx 2200 \, \text{km}^2 \, \text{sr},$$

(3)

where $S \cos \theta$ is the projected surface of the array in the shower plane, and $d\Omega$ is the acceptance solid angle. From Eqs. (2) and (3), one notes that an event rate of $\approx 1 \, \text{yr}^{-1}$ requires an integrated flux

$$\int_{E_G,\text{min}}^{E_G,\text{max}} J_G(E_G) \, dE_G \approx 1.4 \times 10^{-21} \, \text{cm}^{-2} \, \text{s}^{-1} \, \text{sr}^{-1}.$$

(4)

Requiring that events also trigger the fluorescence detectors increases this flux by a factor of 10.

Detailed characteristics of ultra-high energy cosmic ray sources are largely unknown [22], rendering a direct calculation of the expected $G$-flux speculative. In the next section we assess the viability of this flux by comparing with existing limits on gamma ray and neutrino fluxes.

III. PRODUCTION OF $G$–HADRONS IN ASTROPHYSICAL SOURCES

Among the non-thermal sources in the universe, radio–loud active galactic nuclei (AGNs) seem to be the most important energetically. There are other interesting powerful sources, like gamma ray bursts and cluster of galaxies; however, the non-thermal energy release in these astrophysical processes does not come close that of AGNs. At radio frequencies, where very large baseline interferometers can resolve the emission regions at milliarcsecond scale, many of radio–loud AGNs exhibit compact jets of relativistic plasma which are remarkably well collimated, with opening angles about a few degrees or less. The AGNs come in various disguises according to the orientation of their radio jets axes and characteristics of the circum-nuclear matter in their host galaxies. The most extreme versions are Fanaroff Riley radio-galaxies with the radio jet axes almost in the plane of the sky and blazars with the radio jet axes pointing close to the line of sight to the observer, yielding a significant flux enhancement because of Doppler boosting.

A total of 66 blazars have been detected to date as GeV $\gamma$-ray sources by the Energetic Gamma Ray Experiment Telescope (EGRET) on board of the Compton Gamma Ray Observatory (CGRO) [23]. In addition, observations from ground-based Čerenkov telescopes
indicated that at least in 2 of these sources the $\gamma$-ray spectrum can be traced to more than a TeV observed photon energy [24]. The non-thermal emission of these powerful objects indicates a double-peak structure in the overall spectral energy distribution. The first component (from radio to $X$-rays) is generally interpreted as being due to synchrotron radiation from a population of non-thermal electrons, whereas the second component ($\gamma$-rays) is explained either through inverse Compton scattering of the same electron population with the various seed photon fields traversed by the jet [25], or by the decay of neutral pions produced when the highly relativistic baryonic outflow collides with diffuse gas targets moving across the jet [26].

In this work we focus on the “relativistic jet meets target” scenario, in which $G$-hadrons can be produced in collisions of ultra-high energy protons in the jet with those in surrounding gas. In the course of these collisions pions are produced which, on decay, give rise to a flux of photons and neutrinos. In what follows, we estimate the relative probabilities for production of $G$’s and high energy neutrinos. We can then assess whether existing limits on the flux of high energy energy neutrinos and EGRET data from low energy gamma rays are consistent with the $G$ flux given in Eq. (4).

In order to specify detection criteria, the following kinematic analysis is relevant. The average energy of the produced $G$ in the target system is given by

$$E_{G}^{\text{lab}} \simeq \sqrt{\frac{E_{p}^{\text{lab}}}{2 M_p} E_{G}^{\text{cm}}},$$

where $E_{p}^{\text{lab}}$ is the energy of the high energy proton in the jet, and $E_{G}^{\text{cm}}$ is the $G$ energy in the center-of-mass of the $pp$ collision. Full acceptance at PAO requires a minimum energy $E_{G,\text{min}}^{\text{lab}} \simeq 5 \times 10^9$ GeV for the showering $G$ hadron. From Eq. (5), this determines a minimum energy for the high energy proton:

$$E_{p}^{\text{lab}} > 2 \times 10^{14} \text{ GeV} \left(\frac{500 \text{ GeV}}{M_G}\right)^2 \left(\frac{4 M_G^2}{\hat{s}}\right),$$

where we have made use of the fact that, on the average, $E_{G}^{\text{cm}} = \sqrt{\hat{s}}/2$, where $\hat{s}$ is the square of the energy in the center-of-mass of the parton-parton collision producing the $G$’s.

It is immediately apparent from Eq. (6) that ultra-high energy sources are required in order to produce $G$’s which can generate extensive air showers. The required energy can be reduced by restricting production to large $\hat{s}$. In order that the maximum energy at the source does not exceed $10^{14}$ GeV, a speculative number sometimes used in the literature [27], we take $\hat{s} \geq 16 M_G^2$, which via Eq. (6) leads to $E_{p,\text{min}}^{\text{lab}} \approx 5 \times 10^{13}$ GeV.

We now turn to evaluating the consistency of the required $G$ flux and its accompanying pion flux with existing limits on neutrino and gamma ray fluxes. The required relationship is

$$\int J_{\nu}(E_{\nu}) \, dE_{\nu} = \frac{2}{3} \frac{\sigma_{\text{inel}}}{\sigma_{\text{pp} \rightarrow G}(\hat{s}_{\text{min}})} \frac{\langle N_{\nu} \rangle}{N_G} \int J_{G}(E_{G}) \, dE_{G},$$

where $J_{\nu}$ is the neutrino flux, $\sigma_{\text{inel}}$ and $\sigma_{\text{pp} \rightarrow G}$ are the total $pp$ inelastic and inclusive $pp \rightarrow G$ cross sections, and $N_{\nu}$ and $N_G = 2$ are the neutrino and gluino multiplicities per collision. The factor of 2/3 accounts for the fact that only charged pions contribute to neutrino production. To get our estimates we adopt $\sigma_{\text{inel}} \sim 130 \text{ mb}$ [28].
The inclusive $pp \rightarrow G$ production cross section has been evaluated [29] using CTEQ5L parton distribution functions [30]. For $M_G = 500$ GeV, and integrated over all values of $\hat{s}$, the cross section can be conveniently parametrized as

$$\sigma_{pp\rightarrow G} = 1.17 \times 10^{-43} \left( \frac{E_p}{\text{GeV}} \right)^{1.0565} \text{cm}^2. \quad (8)$$

When the condition $\hat{s} \geq 16M_G^2$ is imposed on the integration, we have estimated that the cross section given in Eq. (8) is decreased by a factor of 2. The dependence on $M_G$ can be roughly described by the scaling behavior $e^{-0.007(M_G-500 \text{ GeV})}$.

The $E_G$ integration interval in Eq. (7) is very narrow, as discussed above. Since we are working at the very high energy end of the cosmic ray spectrum, we will examine the consistency of Eq. (7) with bounds on the neutrino flux at the highest energy. Inserting Eq. (4), into Eq. (7), we obtain (with $N_G = 2$) the high energy neutrino flux accompanying $G$-production

$$J_\nu(E_\nu) \Delta E_\nu \approx 1200 \langle N_\nu \rangle \int J_G(E_G) \ dE_G \quad (9)$$

$$\approx \langle N_\nu \rangle 1.7 \times 10^{-18} \text{ cm}^2 \text{ s}^{-1} \text{ sr}^{-1} \quad (10)$$

for an event rate of 1 yr$^{-1}$. At the end of the spectrum, we may approximately consider only the most energetic of the secondary $\pi^\pm$, carrying about 8\% of the primary energy [31]. This entails a neutrino multiplicity $\langle N_\nu \rangle = 3$/event, each carrying 1/4 of the pion energy. Thus $\Delta E_\nu \sim \langle E_\nu \rangle \approx 0.02 \langle E_p \rangle = 1.5 \times 10^{12}$ GeV. The expected neutrino flux (all flavors) at this energy is then

$$J_\nu(1.5 \times 10^{12} \text{ GeV}) \approx 4.4 \times 10^{-30} \text{ GeV}^{-1} \text{ cm}^2 \text{ s}^{-1} \text{ sr}^{-1}. \quad (11)$$

An upper bound for the neutrino flux in this energy range has been obtained through the absence of radio signals originating in the Moon’s rim (GLUE [32]) or in the Greenland ice sheet (FORTE satellite [33]). The flux associated with 1 $G$ event/yr at PAO, given in Eq. (11), is comfortably below these limits. The RICE experiment [34] (detecting electron neutrino-induced radio Cerenkov radiation in the polar ice cap) has reported upper bounds on the $\nu_e + \bar{\nu}_e$ flux for energies up to $10^{12}$ GeV. A slight extrapolation of their result to $1.5 \times 10^{12}$ GeV gives a (3-flavor) upper bound which is a factor of 2.5 lower than the flux in Eq. (11). The flux in Eq. (11) is also a factor of 3 lower than the model-independent bound found in [35] from the absence of horizontal air showers. Finally, extrapolation of the ultra-high energy neutrino intensity given in Eq. (11) down to lower energies, assuming $J_\nu(E_\nu) \propto E_\nu^{-2}$ (more on this below), leads to a neutrino flux which is in agreement with all upper limits on $J_\nu(E_\nu)$ reported by the AMANDA Collaboration [36].

There is an additional bound (known as the cascade limit) coming from EGRET’s observation of GeV gamma rays [37]. The relation to neutrinos is as follows: isotopically symmetric triplets of $\pi^+$, $\pi^-$, and $\pi^0$ produced at high energy sources yield 3 $\nu$’s per charged pion and 2 $\gamma$’s per $\pi^0$, with energies $E_\nu = E_\pi/4$ and $E_\gamma = E_\pi/2$, respectively. ($E_\pi$ is the pion energy.) The integrated neutrino and gamma ray energies then satisfy [38]

$$\int_{E_\pi,\text{min}/4}^{E_\pi,\text{max}/4} E_\nu \ J_\nu(E_\nu) \ dE_\nu = \frac{3}{2} \int_{E_\pi,\text{min}/2}^{E_\pi,\text{max}/2} E_\gamma \ J_\gamma(E_\gamma) \ dE_\gamma. \quad (12)$$
Normalization to EGRET data [37],
\[ \int_{0.1 \text{ GeV}}^{\infty} J_\gamma(E') \, dE' = 1.45 \times 10^{-5} \text{ cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1}, \]
with a spectrum \( J_\gamma(E') = C_\gamma E'^{-2} \), implies \( C_\gamma = 1.45 \times 10^{-6} \text{ GeV cm}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \). Using Eq. (12) we obtain \( C_\nu = 3/2C_\gamma \) for the normalization constant of the accompanying neutrino spectrum, again on the assumption of a spectrum \( J_\nu \propto E_\nu^{-2} \). At face value this gives a flux at \( 1.5 \times 10^{12} \text{ GeV} \) which is nearly a factor of 5 smaller than the flux in Eq. (11). However, this represents a very large extrapolation: even logarithmic corrections to the \( E^{-2} \) spectrum could result in sizeable deviations at very high energies. To see how such corrections might arise, we note that for a primary high energy proton with energy \( E_p \), the resulting pion spectrum is expected to obey a modified Feynman scaling in the central rapidity region,
\[ \frac{dN_\pi}{dE_\pi}(E_\pi) \approx C(E_p)/E_\pi, \]
where the Lorentz dilated \( C(E_p) \) is generically a function which grows as a power of \( \ln(E_\pi) \), falling to zero at the cutoff \( E_\pi = 0.08 \, E_{p,\text{max}} \). Hence we suspend judgement with respect to the constraint imposed by the cascade bound.

**IV. \( \Lambda_{\text{SUSY}} \) WRITTEN IN THE SKY?**

We have described conditions under which \( G \)-hadrons could be produced in astrophysical sources in sufficient abundance to allow detection in an air shower array. Moreover, the complete longitudinal and shower-front curvature profiles of \( G \)-hadron cascades would uniquely differentiate them from proton and nucleus backgrounds. In this last part of the paper we examine what we can learn about SUSY if such events are actually observed.

We have seen that limits on heavy isotope abundance place an *upper* bound on \( \Lambda_{\text{SUSY}} \). Detection of \( G \)-hadrons, presumed to originate at cosmological distances, will place a *lower* bound on the proper lifetime \( \tau_0 \) of the \( G \). With the use of Eq. (1) this can be translated into a lower bound on the scale of SUSY breaking. The argument can be specified as follows. In the presence of decay, the integration of the \( G \)-flux over source distances \( r \) out to the horizon \( R \approx 3 \text{ Gpc} \) is modified by inclusion of a damping factor
\[ f = \int_0^R e^{-r/[c\tau(E_G)]} \, dr = \frac{c\tau(E_G)}{R} \left( 1 - e^{-R/[c\tau(E_G)]} \right), \]
where the Lorentz dilated \( G \) lifetime is given in terms of the proper lifetime \( \tau(E) = \tau_0 \, E_G/M_G \), with \( E_G \approx 5 \times 10^9 \text{ GeV} \). If a few \( G \) events are seen during the lifetime of PAO, and the observation of \( >10^{12} \text{ GeV} \) neutrinos merits the expectation of 1 \( G \) event/yr, then \( f \) cannot be too small, say \( f \gtrsim 0.1 \). This implies that \( c\tau(E_G)/R \gtrsim 0.1 \), and places a bound
\[ \tau_0 \gtrsim 100 \text{ yr} \]
on the proper lifetime. From Eq.(1), one would then obtain a lower limit on the SUSY breaking scale,

$$\Lambda_{\text{SUSY}} > 2.6 \times 10^{11} \text{ GeV}. \quad (18)$$

In conjunction with the upper limit from isotope abundance, $$\Lambda_{\text{SUSY}} < 10^{13} \text{ GeV}$$, we can then adduce the remarkable result that the observation of a few $G$ events during the operating life of PAO can fix the scale of high energy SUSY breaking.

V. SUMMARY

In this work we have examined under which conditions metastable $G$-hadrons can be synthesized in powerful astrophysical environments and can be detected on Earth. We show, that if cosmic ray sources are able to accelerate protons somewhat above $5 \times 10^{13} \text{ GeV}$, about 0.5-1 $G$-hadron induced shower per year could be detected at PAO. Additionally, $G$-cascades provide a signal which is easily differentiated from background. These numbers are for $M_G = 500 \text{ GeV}$: the event rate could be doubled for $M_G = 400 \text{ GeV}$.

Should some of these very distinctive showers be observed, a lower bound on $\Lambda_{\text{SUSY}}$ can be deduced in a manner unavailable at colliders. The combination of such a lower bound with the upper bound imposed by the scarcity of heavy isotopes fixes the scale of SUSY breaking to a relatively narrow window, $2.6 \times 10^{11} \text{ GeV} < \Lambda_{\text{SUSY}} < 10^{13} \text{ GeV}$. This result, if true, combined with the expected low mass of the higgs would provide strong support for a finely-tuned universe.

Note added: After this paper was completed, a work with related interesting ideas for detection of long-lived gluinos at colliders and at IceCube appeared [41]. For observation at IceCube, the gluino production takes place in cosmic ray collisions in the Earth’s atmosphere. The steeply falling cosmic ray luminosity above $10^9 \text{ GeV}$ allows sufficient $G$ production for a measurable signal at IceCube only for $M_G \lesssim 150 \text{ GeV}$, which is very close to the present bounds from Tevatron [42]. In examining production at extraterrestrial sources, the present work can substantially extend the possible range of $M_G$ able to be probed by cosmic ray experiments.

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