ULTRA HIGH ENERGY COSMIC RAYS FROM DECAYING SUPERHEAVY PARTICLES *

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Decaying superheavy particles can be produced by Topological Defects or, in case they are quasi-stable, as relics from the early Universe. The decays of these particles can be the sources of observed Ultra High Energy Cosmic Rays ($E \sim 10^{10} - 10^{12}$ GeV). The Topological Defects as the UHE CR sources are critically reviewed and cosmic necklaces and monopole-antimonomope pairs are identified as most plausible sources. The relic superheavy particles are shown to be clustering in the halo and their decays produce UHE CR without GZK cutoff. The Lightest Supersymmetric Particles with Ultra High Energies are naturally produced in the cascades accompanying the decays of superheavy particles. These particles are discussed as UHE carriers in the Universe.

I. INTRODUCTION

The observation of cosmic ray particles with energies higher than $10^{11}$ GeV [1] gives a serious challenge to the known mechanisms of acceleration. The shock acceleration in different astrophysical objects typically gives maximal energy of accelerated protons less than $(1 - 3) \cdot 10^{10}$ GeV [2]. The unipolar induction can provide the maximal energy $1 \cdot 10^{11}$ GeV only for the extreme values of the parameters [3]. Much attention has recently been given to acceleration by ultrarelativistic shocks [4], [5]. The particles here can gain a tremendous increase in energy, equal to $\Gamma^2$, at a single reflection, where $\Gamma$ is the Lorentz factor of the shock. However, it is known (see e.g. the simulation for pulsar relativistic wind in [6]) that particles entering the shock region are captured there or at least have a small probability to escape.

Topological defects, TD, (for a review see [7]) can naturally produce particles of ultrahigh energies (UHE). The pioneering observation of this possibility was made by Hill, Schramm and Walker [8] (for a general analysis of TD as UHE CR sources see [8] and for a review [10]).

In many cases TD become unstable and decompose to constituent fields, superheavy gauge and Higgs bosons (X-particles), which then decay producing UHE CR. It could happen, for example, when two segments of ordinary string, or monopole and antimonomopole touch each other, when electrical current in superconducting string reaches the critical value and in some other cases.

In most cases the problem with UHE CR from TD is not the maximal energy, but the fluxes. One very general reason for the low fluxes consists in the large distance between TD. A dimension scale for this distance is the Hubble distance $H_0^{-1}$. However, in some rather exceptional cases this dimensional scale is multiplied to a small dimensionless value $r$. If a distance between TD is larger than UHE proton attenuation length (due to the GZK effect [11]), then the flux at UHE is typically exponential suppressed.

Ordinary cosmic strings can produce particles when a loop annihilate into double line [12]. The produced UHE CR flux is strongly reduced due to the fact that a loop oscillates, and in the process of a collapse the two incoming parts of a loop touch each other in one point producing thus the smaller loops, instead of two-line annihilation. However, this idea was recently revived due to recent work [13]. It is argued there that the energy loss of the long strings is dominated by production of very small loops with the size down to the width of a string, which immediately annihilate into superheavy particles. A problem with this scenario is too large distance between strings (of order of the Hubble distance). For a distance between an observer and a string being the same, the observed spectrum of UHE CR has an exponential cutoff at energy $E \sim 3 \cdot 10^{19}$ eV.

Superheavy particles can be also produced when two segments of string come into close contact, as in cusp events [14]. This process was studied later by Gill and Kibble [15], and they concluded that the resulting cosmic ray flux is far too small. An interesting possibility suggested by Brandenberger [14] is the cusp “evaporation” on cosmic strings. When the distance between two segments of the cusp becomes of the order of the string width, the cusp

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may “annihilate” turning into high energy particles, which are boosted by a very large Lorentz factor of the cusp [14]. However, the resulting UHE CR flux is considerably smaller than one observed [10].

Superconducting strings [17] appear to be much better suited for particle production. Moving through cosmic magnetic fields, such strings develop electric currents and copiously produce charged heavy particles when the current reaches certain critical value. The CR flux produced by superconducting strings is affected by some model-dependent string parameters and by the history and spatial distribution of cosmic magnetic fields. Models considered so far failed to account for the observed flux [15].

Monopole-antimonopole pairs \((\bar{M}M)\) can form bound states and eventually annihilate into UHE particles [19], [20]. For an appropriate choice of the monopole density \(n_M\), this model is consistent with observations; however, the required (low) value of \(n_M\) implies fine-tuning. In the first phase transition \(G \to H \times U(1)\) in the early Universe the monopoles are produced with too high density. It must then be diluted by inflation to very low density, precisely tuned to the observed UHE CR flux.

Monopole-string networks can be formed in the early Universe in the sequence of symmetry breaking

\[
G \to H \times U(1) \to H \times Z_N. 
\]

For \(N \geq 3\) an infinite network of monopoles connected by strings is formed. The magnetic fluxes of monopoles in the network are channeled into strings that connect them. The monopoles typically have additional unconfined magnetic and chromo-magnetic charges. When strings shrink the monopoles are pulled by them and are accelerated. The accelerated monopoles produce extremely high energy gluons, which then fragment into UHE hadrons [2]. The produced flux is too small to explain UHE CR observation [22].

Cosmic necklaces are TD which are formed in a sequence of symmetry breaking given by Eq. (1) when \(N = 2\). The first phase transition produces monopoles, and at the second phase transition each monopole gets attached to two strings, with its magnetic flux channeled along the strings. The resulting necklaces resemble “ordinary” cosmic strings with monopoles playing the role of beads. Necklaces can evolve in such way that a distance between monopoles diminishes and in the end all monopoles annihilate with the neighboring antimonopoles [23].

The produced UHE CR flux [21] is close to the observed one and the shape of the spectrum resembles one observed. The distance between necklaces can be much smaller than attenuation length of UHE protons.

Superheavy relic particles can be sources of UHE CR [24, 25]. In this scenario Cold Dark Matter (CDM) have a small admixture of long-lived superheavy particles. These particles must be heavy, \(m_X > 10^{12} \text{ GeV}\), long-lived \(\tau_X > t_0\), where \(t_0\) is the age of the Universe, and weakly interacting. The required life-time can be provided if this particle has (almost) conserved quantum number broken very weakly due to warmhole [25] or instanton [22] effects. Several mechanisms for production of such particles in the early Universe were identified. Like other forms of non-dissipative CDM, X-particles must accumulate in the halo of our Galaxy [25] and thus they produce UHE CR without GZK cutoff and without appreciable anisotropy.

The UHE carriers produced at the decay of superheavy relic particles or from TD, can be be nucleons, photons and neutrinos or neutralinos [25]. Production of neutralinos occurs in particle cascade, which originates at the decay of superheavy X-particle in close analogy to QCD cascade [25]. Though flux of UHE neutralino is of the same order as neutrino flux, its detection is more problematic because of smaller cross-section. The other particles discussed as carrier of UHE signal are gluino [26, 27], in case it is Lightest Supersymmetric Particle (LSP), and heavy monopole [20, 27].

In this paper I will present the results obtained in our joint works with M.Kachelriess and A.Vilenkin about necklaces and superheavy relic particles as possible sources of UHE CR. Neutralino and gluino as the carriers of UHE signal will be also shortly discussed.

II. NECKLACES

Necklaces produced in a sequence of symmetry breaking \(G \to H \times U(1) \to H \times Z_2\) form the infinite necklaces having the shape or random walks and a distribution of closed loops. Each monopole in a necklace is attached to two strings.

The monopole mass \(m\) and the string tension \(\mu\) are determined by the corresponding symmetry breaking scales, \(\eta_s\) and \(\eta_m\) (\(\eta_m > \eta_s\)): \(m \sim 4\pi \eta_m / e\), \(\mu \sim 2\pi \eta_s^2\). Here, \(e\) is the gauge coupling. The mass per unit length of string is equal to its tension, \(\mu\). Each string attached to a monopole pulls it with a force \(F \sim \mu\) in the direction of the string. The monopole radius \(\delta_m\) and the string thickness \(\delta_s\) are typically of the order \(\delta_m \sim (\eta_m)^{-1}\), \(\delta_s \sim (\eta_s)^{-1}\).

An important quantity for the necklace evolution is the dimensionless ratio

\[
\eta = \frac{\eta_s}{\eta_m}. 
\]
\[ r = m/\mu d, \] (2)

We expect the necklaces to evolve in a scaling regime. If \( \xi \) is the characteristic length scale of the network, equal to the typical separation of long strings and to their characteristic curvature radius, then the force per unit length of string is \( f \sim \mu/\xi \), and the acceleration is \( a \sim (r + 1)^{-1}\xi^{-1} \). We assume that \( \xi \) changes on a Hubble time scale \( \sim t \). Then the typical distance travelled by long strings in time \( t \) should be \( \sim \xi \), so that the strings have enough time to intercommute in a Hubble time. This gives \( at^2 \sim \xi \), or

\[ \xi \sim (r + 1)^{-1/2}t. \] (3)

The typical string velocity is \( v \sim (r + 1)^{-1/2} \).

It is argued in Ref. (23) that \( r(t) \) is driven towards large value \( r \gg 1 \). However, for \( r \geq 10^6 \) the characteristic velocity of the network falls down below the virial velocity, and the necklaces will be trapped by gravitational clustering of the matter. This may change dramatically the evolution of network. One possible interesting effect for UHE CR is the enhancement of necklace space density within Local Supercluster – a desirable effect as far absence of the GZK cutoff is concerned. However, we restrict our consideration by the case \( r < 10^6 \). The distance between necklaces is still small enough, \( \xi \gtrsim 3 \text{ Mpc} \), to assume their uniform distribution, when calculating the UHE CR flux.

Self-intersections of long necklaces result in copious production of closed loops. For \( r \gtrsim 1 \) the motion of loops is not periodic, so loop self-intersections should be frequent and their fragmentation into smaller loops very efficient. A loop of size \( \ell \) typically disintegrates on a timescale \( \tau \sim r^{-1/2}t \). All monopoles trapped in the loop must, of course, annihilate in the end.

Annihilating \( MM \) pairs decay into Higgs and gauge bosons, which we shall refer to collectively as \( X \)-particles. The rate of \( X \)-particle production is easy to estimate if we note that infinite necklaces lose a substantial fraction of their length to closed loops in a Hubble time. The string length per unit volume is \( \sim \xi^{-2} \), and the monopole rest energy released per unit volume per unit time is \( r\mu/\xi^2 t \). Hence, we can write

\[ n_X \sim r^2 \mu/(t^2 m_X), \] (4)

where \( m_X \sim e\eta_m \) is the \( X \)-particle mass.

\( X \)-particles emitted by annihilating monopoles decay into hadrons, photons and neutrinos, the latter two components are produced through decays of pions.

The diffuse flux of ultra-high energy protons can be evaluated as

\[ I_p(E) = \frac{c n_X}{4\pi m_X} \int_0^{t_0} dt \ W_N(m_X, x_0) \frac{dE_p(E, t)}{dE} \] (5)

where \( dn_X/dt \) is given by Eq. (4), \( E \) is an energy of proton at observation and \( dE_p(E, t) \) is its energy at generation at cosmological epoch \( t \), \( x_0 = E_0/E \) and \( W_N(m_X, x) \) is the fragmentation function of \( X \)-particle into nucleons of energy \( E = x m_X \). The value of \( dE_p/dE \) can be calculated from the energy losses of a proton on microwave background radiation (e.g. see 3). In Eq. (5) the recoil protons are taken into account, while in Ref. (23) their contribution was neglected.

The fragmentation function \( W_N(m_X, x) \) is calculated using the decay of \( X \)-particle into QCD partons (quark, gluons and their supersymmetric partners) with the consequent development of the parton cascade. The cascade in this case is identical to one initiated by \( e^+e^- \) -annihilation. We have used the fragmentation function in the gaussian form as obtained in MLLA approximation in 34 and 35.

In our calculations the UHE proton flux is fully determined by only two parameters, \( r^2 \mu \) and \( m_X \). The former is restricted by low energy diffuse gamma-radiation. It results from e-m cascades initiated by high energy photons and electrons produced in the decays of \( X \)-particles.

The cascade energy density predicted in our model is

\[ \omega_{\text{cas}} = \frac{1}{2} f_\pi r^2 \mu \int_0^{t_0} dt \ \frac{1}{t^2 (1 + z)^2} = \frac{3}{4} f_\pi r^2 \mu \frac{t_0^2}{t_0}, \] (6)

where \( t_0 \) is the age of the Universe (here and below we use \( h = 0.75 \)), \( z \) is the redshift and \( f_\pi \sim 1 \) is the fraction of energy transferred to pions. In Eq. (1) we took into account that half of the energy of pions is transferred to photons and electrons. The observational bound on the cascade density, for the kind of sources we are considering here, is \( \omega_{\text{cas}} \lesssim 10^{-5} \text{ eV/cm}^3 \). This gives a bound on the parameter \( r^2 \mu \).
In numerical calculations we used $r^2 \mu = 1 \times 10^{28} \text{ GeV}^2$, which results in $\omega_{\text{max}} = 5.6 \cdot 10^{-6} \text{ eV/cm}^3$, somewhat below the observational limit. Now we are left with one free parameter, $m_X$, which we fix at $1 \cdot 10^{14} \text{ GeV}$. The maximum energy of protons is then $E_{\text{max}} \sim 10^{13} \text{ GeV}$. The calculated proton flux is presented in Fig.1, together with a summary of observational data taken from ref. [37].

Let us now turn to the calculations of UHE gamma-ray flux from the decays of X-particles. The dominant channel is given by the decays of neutral pions. The flux can be readily calculated as

$$I_{\gamma}(E) = \frac{1}{4\pi} \hat{n}_X \lambda_{\gamma}(E) N_{\gamma}(E),$$

where $\hat{n}_X$ is given by Eq.(4), $\lambda_{\gamma}(E)$ is the absorption length of a photon with energy $E$ due to $e^+e^-$ pair production on background radiation and $N_{\gamma}(E)$ is the number of photons with energy $E$ produced per one decay of X-particle. The latter is given by

$$N_{\gamma}(E) = \frac{2}{m_X} \int_{E/m_X}^{1} \frac{dx}{x} W_{\pi^0}(m_X, x)$$

where $W_{\pi^0}(m_X, x)$ is the fragmentation function of X-particles into $\pi^0$ pions.

At energy $E > 1 \cdot 10^{10} \text{ GeV}$ the dominant contribution to the gamma-ray absorption comes from the radio background. The significance of this process was first noticed in [38] (see also book [3]). New calculations for this absorption were recently done [39]. We have used the absorption lengths from this work.

When evaluating the flux at $E > 1\cdot 10^{10} \text{ GeV}$ we neglected cascading of a primary photon, because pair production and inverse compton scattering occur at these energies on radio background, and thus at each collision the energy of a cascade particle is halved. Moreover, assuming an intergalactic magnetic field $H \geq 1 \cdot 10^{-8}$, the secondary electrons and positrons lose their energy mainly due to synchrotron radiation and the emitted photons escape from the considered energy interval [10].

The calculated flux of gamma radiation is presented in Fig.1 by the curve labelled $\gamma$. One can see that at $E \sim 1 \cdot 10^{11} \text{ GeV}$ the gamma ray flux is considerably lower than that of protons. This is mainly due to the difference in the attenuation lengths for protons ($110 \text{ Mpc}$) and photons ($2.6 \text{ Mpc}$ [39] and $2.2 \text{ Mpc}$ [38]). At higher energy the attenuation length for protons dramatically decreases ($13.4 \text{ Mpc at } E = 1 \cdot 10^{12} \text{ GeV}$) and the fluxes of protons and photons become comparable.

A requirement for the models explaining the observed UHE events is that the distance between sources must be smaller than the attenuation length. Otherwise the flux at the corresponding energy would be exponentially suppressed. This imposes a severe constraint on the possible sources. For example, in the case of protons with energy $E \sim (2 - 3) \cdot 10^{11} \text{ GeV}$ the proton attenuation length is $19 \text{ Mpc}$. If protons propagate rectilinearly, there should be several sources inside this radius; otherwise all particles would arrive from the same direction. If particles are strongly deflected in extragalactic magnetic fields, the distance to the source should be even smaller. Therefore, the sources of the observed events at the highest energy must be at a distance $R \lesssim 15 \text{ Mpc}$ in the case or protons.

In our model the distance between sources, given by Eq.(4), satisfies this condition for $r > 3 \cdot 10^{4}$. This is in contrast to other potential sources, including superconducting cosmic strings and powerful astronomical sources such as AGN, for which this condition imposes severe restrictions.

The difficulty is even more pronounced in the case of UHE photons. These particles propagate rectilinearly and their absorption length is shorter: $2 - 4 \text{ Mpc at } E \sim 3 \cdot 10^{11} \text{ GeV}$. It is rather unrealistic to expect several powerful astronomical sources at such short distances. This condition is very restrictive for topological defects as well. The necklace model is rather exceptional regarding this aspect.

### III. UHE CR FROM RELIC QUASISTABLE PARTICLES

This possibility was recognized recently in Refs([4,23]).

Our main assumption is that Cold Dark Matter (CDM) has a small admixture of long-lived supermassive X-particles. Since, apart from very small scales, fluctuations grow identically in all components of CDM, the fraction of X-particles, $\xi_X$, is expected to be the same in all structures. In particular, $\xi_X$ is the same in the halo of our Galaxy and in the extragalactic space. Thus the halo density of X-particles is enhanced in comparison with the extragalactic density. The decays of these particles produce UHE CR, whose flux is dominated by the halo component, and therefore has no GZK cutoff. Moreover, the potentially dangerous e-m cascade radiation is suppressed.
First, we address the elementary-particle and cosmological aspects of a superheavy long-living particle. Can the relic density of such particles be as high as required by observations of UHE CR? And can they have a lifetime comparable or larger than the age of the Universe?

Let us assume that $X$-particle is a neutral fermion which belongs to a representation of the $SU(2) \times U(1)$ group. We assume also that the stability of $X$-particles is protected by a discrete symmetry which is the remnant of a gauge symmetry and is respected by all interactions except quantum gravity through wormhole effects. In other words, our particle is very similar to a very heavy neutralino with a conserved quantum number, $R'$, being the direct analogue of $R$-parity (see [32] and the references therein). Thus, one can assume that the decay of $X$-particle occurs due to dimension 5 operators, inversely proportional to the Planck mass $m_{Pl}$ and additionally suppressed by a factor $\exp(-S)$, where $S$ is the action of a wormhole which absorbs $R'$-charge. As an example one can consider a term

$$\mathcal{L} \sim \frac{1}{m_{Pl}} \bar{\Psi} \nu \phi \exp(-S),$$

where $\Psi$ describes $X$-particle, and $\phi$ is a $SU(2)$ scalar with vacuum expectation value $v_{EW} = 250$ GeV. After spontaneous symmetry breaking the term $\bar{\Psi} \nu \phi$ results in the mixing of $X$-particle and neutrino, and the lifetime due to $X \to \nu + q + \bar{q}$, e.g., is given by

$$\tau_X \sim \frac{192(2\pi)^3}{(G_F v_{EW})^2} m_X^{-2S},$$

where $G_F$ is the Fermi constant. The lifetime $\tau_X > t_0$ for $X$-particle with $m_X \geq 10^{13}$ GeV needs $S > 44$. This value is within the range of the allowed values as discussed in Ref. [32].

Let us now turn to the cosmological production of $X$-particles with $m_X \geq 10^{13}$ GeV. Several mechanisms were identified in [24], including thermal production at the reheating stage, production through the decay of inflaton field at the end of the "pre-heating" period following inflation, and through the decay of hybrid topological defects, such as monopoles connected by strings or walls bounded by strings.

For the thermal production, temperatures comparable to $m_X$ are needed. In the case of a heavy decaying gravitino, the reheating temperature $T_R$ (which is the highest temperature relevant for our problem) is severely limited to value below $10^8 - 10^{10}$ GeV, depending on the gravitino mass (see Ref. [32] and references therein). On the other hand, in models with dynamically broken supersymmetry, the lightest supersymmetric particle is the gravitino. Gravitinos with mass $m_{3/2} \leq 1$ keV interact relatively strongly with the thermal bath, thus decoupling relatively late, and can be the CDM particle [33]. In this scenario all phenomenological constraints on $T_R$ (including the decay of the second lightest supersymmetric particle) disappear and one can assume $T_R \sim 10^{11} - 10^{12}$ GeV. In this range of temperatures, $X$-particles are not in thermal equilibrium. If $T_R < m_X$, the density $n_X$ of $X$-particles produced during the reheating phase at time $t_R$ due to $a + \bar{a} \to X + \bar{X}$ is easily estimated as

$$n_X(t_R) \sim N_a n_a^2 \sigma_X t_R \exp(-2m_X/T_R),$$

where $N_a$ is the number of flavors which participate in the production of $X$-particles, $n_a$ is the density of $a$-particles and $\sigma_X$ is the production cross-section. The density of $X$-particles at the present epoch can be found by the standard procedure of calculating the ratio $n_X/s$, where $s$ is the entropy density. Then for $m_X = 1 \cdot 10^{13}$ GeV and $\xi_X$ in the wide range of values $10^{-8} - 10^{-4}$, the required reheating temperature is $T_R \sim 3 \cdot 10^{11}$ GeV.

In the second scenario mentioned above, non-equilibrium inflaton decay, $X$-particles are usually overproduced and a second period of inflation is needed to suppress their density.

Finally, $X$-particles could be produced by TD such as strings or textures. Particle production occurs at string intersections or in collapsing texture knots. The evolution of defects is scale invariant, and roughly a constant number of particles $\nu$ is produced per horizon volume $t^3$ per Hubble time $t$. ($\nu \sim 1$ for textures and $\nu \gg 1$ for strings.) The main contribution to to the $X$-particle density is given by the earliest epoch, soon after defect formation, and we find $\xi_X \sim 10^{-6} \nu(m_X/10^{13} \text{ GeV})(T_f/10^{10} \text{ GeV})^3$, where $T_f$ is the defect formation temperature. Defects of energy scale $\eta \geq m_X$ could be formed at a phase transition at or slightly before the end of inflation. In the former case, $T_f \sim T_R$, while in the latter case defects should be considered as "formed" when their typical separation becomes smaller than $t$ (hence $T_f < T_R$). It should be noted that early evolution of defects may be affected by friction; our estimate of $\xi_X$ will then have to be modified. X particles can also be produced by hybrid topological defects: monopoles connected by strings or walls bound by strings. The required values of $n_X/s$ can be obtained for a wide range of defect parameters.

The decays of $X$-particles result in the production of nucleons with a spectrum $W_X(m_X, x)$, where $m_X$ is the mass of the $X$-particle and $x = E/m_X$. The flux of nucleons $(p, \bar{p}, n, \bar{n})$ from the halo and extragalactic space can be calculated as
where index $i$ runs through $h$ (halo) and $e x$ (extragalactic), $R_i$ is the size of the halo $R_h$, or the attenuation length of UHE protons due to their collisions with microwave photons, $\lambda_p(E)$, for the halo case and extragalactic case, respectively. We shall assume $m_X n_X^h = \xi_X \rho_{CDM}^h$ and $m_X n_X^{ex} = \xi_X \Omega_{CDM} \rho_{cr}$, where $\xi_X$ describes the fraction of $X$-particles in CDM, $\Omega_{CDM}$ is the CDM density in the units of the critical density $\rho_{cr}$, and $\rho_{CDM}^h \approx 0.3 \text{ GeV/cm}^3$ is the CDM density in the halo. We shall use the following values for these parameters: a large DM halo with $R_h = 100 \text{ kpc}$ (a smaller halo with $R_h = 50 \text{ kpc}$ is possible, too), $\Omega_{CDM} h^2 = 0.2$, the mass of $X$-particle in the range $10^{13} \text{ GeV} < m_X < 10^{16} \text{ GeV}$, the fraction of $X$-particles $\xi_X \ll 1$ and $\tau_X \gg t_0$, where $t_0$ is the age of the Universe. The two last parameters are convolved in the flux calculations in a single parameter $r_X = \xi_X t_0 / \tau_X$. For $W_N(m_X, x)$ we shall use like in the previous section the QCD fragmentation function in MLLA approximation. For the attenuation length of UHE protons due to their interactions with microwave photons, we use the values given in the book [3].

The high energy photon flux is produced mainly due to decays of neutral pions and can be calculated for the halo case as

$$I_\gamma(E) = \frac{1}{4\pi} \frac{n_X}{\tau_X} R_h N_\gamma(E), \quad \text{(13)}$$

where $N_\gamma(E)$ is the number of photons with energy $E$ produced per decay of one $X$-particle, which is given by Eq.[3].

For the calculation of the extragalactic gamma-ray flux, it is enough to replace the size of the halo, $R_h$, by the absorption length of a photon, $\lambda_\gamma(E)$. The main photon absorption process is $e^+e^-$-production on background radiation and, at $E > 1 \cdot 10^{10} \text{ GeV}$, on the radio background. The neutrino flux calculation is similar.

Before discussing the obtained results, we consider the astrophysical constraints.

The most stringent constraint comes from electromagnetic cascade radiation, discussed in the previous section. In the present case this constraint is weaker, because the low-energy extragalactic nucleon flux is $\sim$ 4 times smaller than that one from the Galactic halo (see Fig. 2). Thus the cascade radiation is suppressed by the same factor.

The cascade energy density calculated by integration over cosmological epochs (with the dominant contribution given by the present epoch $z = 0$) yields in our case

$$\omega_{cas} = \frac{1}{5} r_X \frac{\Omega_{CDM} \rho_{cr}}{H_0 t_0} = 6.3 \cdot 10^2 r_X f_\pi \text{ eV/cm}^3.$$

To fit the UHE CR observational data by nucleons from halo, we need $r_X = 5 \cdot 10^{-11}$. Thus the cascade energy density is $\omega_{cas} = 3.2 \cdot 10^{-8} f_\pi \text{ eV/cm}^3$, well below the observational bound.

Let us now discuss the obtained results. The fluxes shown in Fig. 2 are obtained for $R_h = 100 \text{ kpc}$, $m_X = 1 \cdot 10^{13} \text{ GeV}$ and $r_X = \xi_X t_0 / \tau_X = 5 \cdot 10^{-11}$. This ratio $r_X$ allows very small $\xi_X$ and $\tau_X > t_0$. The fluxes near the maximum energy $E_{max} = 5 \cdot 10^{12} \text{ GeV}$ were only roughly estimated (dotted lines on the graph).

It is easy to verify that the extragalactic nucleon flux at $E \leq 3 \cdot 10^9 \text{ GeV}$ is suppressed by a factor $\sim 4$ and by a much larger factor at higher energies due to energy losses. The flux of extragalactic photons is suppressed even stronger, because the attenuation length for photons (due to absorption on radio-radiation) is much smaller than for nucleons (see Ref. [33]). This flux is not shown in the graph. The flux of high energy gamma-radiation from the halo is by a factor 7 higher than that of nucleons and the neutrino flux, given in the Fig.2 as the sum of the dominant halo component and subdominant extragalactic one, is twice higher than the gamma-ray flux.

The spectrum of the observed EAS is formed due to fluxes of gamma-rays and nucleons. The gamma-ray contribution to this spectrum is rather complicated. In contrast to low energies, the photon-induced showers at $E > 10^9 \text{ GeV}$ have the low-energy muon component as abundant as that for nucleon-induced showers [12]. However, the shower production by the photons is suppressed by the LPM effect [33] and by absorption in geomagnetic field (for recent calculations and discussion see [14] and references therein).

We wish to note that the excess of the gamma-ray flux over the nucleon flux from the halo is an unavoidable feature of this model. It follows from the more effective production of pions than nucleons in the QCD cascades from the decay of $X$-particle.

The signature of our model might be the signal from the Virgo cluster. The virial mass of the Virgo cluster is $M_{\text{Virgo}} \sim 1 \cdot 10^{15} M_\odot$ and the distance to it $R = 20 \text{ Mpc}$. If UHE protons (and antiprotons) propagate rectilinearly from this source (which could be the case for $E_p \sim 10^{11} - 10^{12} \text{ GeV}$), their flux is given by

$$F_{p,p}^{\text{Virgo}} = \frac{M_{\text{Virgo}}}{t_0 R^2 m_X^2} W_N(m_X, x). \quad \text{(15)}$$
The ratio of this flux to the diffuse flux from the half hemisphere is $6.4 \cdot 10^{-3}$. This signature becomes less pronounced at smaller energies, when protons can be strongly deflected by intergalactic magnetic fields.

IV. LSP IS UHE CARRIER

LSP is the Lightest Supersymmetric Particle. It can be stable if R-parity is strictly conserved or unstable if R-parity is violated. To be able to reach the Earth from most remote regions in the Universe, the LSP must have lifetime longer than $\tau_{\text{LSP}} \gtrsim t_0/\Gamma$, where $t_0$ is the age of the Universe and $\Gamma = E/m_{\text{LSP}}$ is the Lorentz-factor of the LSP. In case $m_{\text{LSP}} \sim 100$ GeV, $\tau_{\text{LSP}} > 1$ yr.

Theoretically the best motivated candidates for LSP are the neutralino and gravitino. We shall not consider the latter, because it is practically undetectable as UHE particle.

In all elaborated SUSY models the gluino is not the LSP. Only, if the dimension-three SUSY breaking terms are set to zero by hand, gluino with mass $m_{\tilde{g}} = \mathcal{O}(1 \text{ GeV})$ can be the LSP \cite{26}. There is some controversy if the low-mass window $1 \text{ GeV} \lesssim m_{\tilde{g}} \lesssim 4 \text{ GeV}$ for the gluino is still allowed \cite{46,47}. Nevertheless, we shall study the production of high-energy gluinos and their interaction with matter being inspired by the recent suggestion \cite{26} (see also \cite{27}), that the atmospheric showers observed at the highest energies can be produced by colorless hadrons containing gluinos. We shall refer to any of such hadron as $\tilde{g}$-hadron. Light gluinos as UHE particles with energy $E \gtrsim 10^{16}$ eV were considered in some detail in the literature in connection with Cyg X-3 \cite{48,49}. Additionally, we consider heavy gluinos with $m_{\tilde{g}} \gtrsim 150$ GeV \cite{27}.

UHE LSP are most naturally produced at the decays of unstable superheavy particles, either from TD or as the relic ones \cite{28}.

The QCD parton cascade is not a unique cascade process. A cascade multiplication of partons at the decay of superheavy particle appears whenever a probability of production of extra parton has the terms $\alpha \ln Q^2$ and $\alpha \ln^2 Q^2$, where $Q$ is a maximum of parton transverse momentum, i.e. $m_{\chi}$ in our case. Regardless of smallness of $\alpha$, the cascade develops as long as $\alpha \ln Q^2 \gtrsim 1$. Therefore, for extremely large $Q^2$ we are interested in, a cascade develops due to parton multiplication through $SU(2) \times U(1)$ interactions as well. Like in QCD, the account of diagrams with $\alpha \ln Q^2$ gives the Leading Logarithm Approximation to the cascade fragmentation function.

For each next generation of cascade particles the virtuality of partons $q^2$ diminishes. When $q^2 \gg m_{\text{SUSY}}^2$, where $m_{\text{SUSY}}$ is a typical mass of supersymmetric particles, the number of supersymmetric partons in the cascade is the same as their ordinary partners. At $q^2 < m_{\text{SUSY}}^2$ the supersymmetric particles are not produced any more and the remaining particles decay producing the LSP. In Ref.\cite{28} a simple Monte Carlo simulation for SUSU cascading was performed and the spectrum of emitted LSP was calculated. LSP take away a considerable fraction of the total energy ($\sim 40\%$).

The fluxes of UHE LSP are shown in Fig. 3 for the case of their production in cosmic necklaces (see section II). When the LSP is neutralino, the flux is somewhat lower than neutrino flux. The neutralino-nucleon cross-section, $\sigma_{\chi N}$, is also smaller than that for neutrino. For the theoretically favorable masses of supersymmetric particles, $\sigma_{\chi N} \sim 10^{-34} \text{ cm}^2$ at extremely high energies. If the the masses of squarks are near their experimental bound, $M_{L,R} \sim 180$ GeV, the cross-section is 60 times higher.

**Gluino as the LSP** is another phenomenological option. Let us discuss shortly the status of the gluino as LSP.

In all elaborated SUSY models the gluino is not LSP, and this possibility is considered on purely phenomenological basis. Accelerator experiments give the lower limit on the gluino mass as $m_{\tilde{g}} \gtrsim 150$ GeV \cite{28}. The upper limit of the gluino mass is given by cosmological and astrophysical constraints, as was recently discussed in \cite{27}. In this work it was shown that if the gluino provides the dark matter observed in our galaxy, the signal from gluino annihilation and the abundance of anomalous heavy nuclei is too high. Since we are not interested in the case when gluino is DM particle, we can use these arguments to obtain an upper limit for the gluino mass. Calculating the relic density of gluinos (similar as in \cite{27}) and using the condition $\Omega_{\tilde{g}} \ll \Omega_{\text{CDM}}$, we obtained $m_{\tilde{g}} \ll 9 \text{ TeV}$.

Now we come to very interesting argument against existence of a light stable or quasistable gluino \cite{24}. It is plausible that the glueballino ($\tilde{g}g$) is the lightest hadronic state of gluino \cite{38,43}. However, **gluebarino**, i.e. the bound state of gluino and three quarks, is almost stable because baryon number is extremely weakly violated. In Ref.\cite{38} it was argued that the lightest gluebarino is the neutral state ($\tilde{q}u\bar{d}$). These charged gluebarinos are produced by cosmic rays in the earth atmosphere \cite{50}, and light gluino as LSP is excluded by the search for heavy hydrogen or by proton decay experiments (in case of quasistable gluino). In the case that the lightest gluebarino is neutral, see \cite{38}, the arguments of \cite{24} still work if a neutral glueballino forms a bound state with the nuclei. Thus, a light gluino is disfavored.

The situation is different if the gluino is heavy, $m_{\tilde{g}} \gtrsim 150$ GeV. This gluino can be unstable due to weak R-parity violation \cite{24} and have a lifetime $\tau_{\tilde{g}} \gtrsim 1$ yr, i.e. long enough to be UHE carrier (see beginning of this section).
Then the calculated relic density at the time of decay is not in conflict with the cascade nucleosynthesis and all cosmologically produced $\tilde{g}$-hadrons decayed up to the present time. Moreover, the production of these gluinos by cosmic rays in the atmosphere is ineffective because of their large mass.

Glueballino, or more generally $\tilde{g}$-hadron, looses its energy while propagating from a source to the Earth. The dominant energy loss of the $\tilde{g}$-hadron is due to pion production in collisions with microwave photons. Pion production effectively starts at the same Lorentz-factor as in the case of the proton. This implies that the energy of the GZK cutoff is a factor $m_{\tilde{g}}/m_p$ higher than in case of the proton. The attenuation length also increases because the fraction of energy lost near the threshold is small, $\mu/m_{\tilde{g}}$, where $\mu$ is a pion mass. Therefore, even for light $\tilde{g}$-hadrons, $m_{\tilde{g}} \gtrsim 2\text{ GeV}$, the steepening of the spectrum is less pronounced than for protons.

The spectrum of $\tilde{g}$-hadrons from the cosmic necklaces accounted for absorption in intergalactic space, is shown in Fig. 3.

A very light UHE $\tilde{g}$-hadron interacts with the nucleons in the atmosphere similarly to UHE proton. The cross-section is reduced only due to the radius of $\tilde{g}$-hadron and is of order of $\sim 1\text{ mb}$ [49]. In case of very heavy $\tilde{g}$-hadron the total cross-section can be of the same order of magnitude, but the cross-section with the large energy transfer, relevant for the detection in the atmosphere, is very small [28]. This is due to the fact that interaction of gluino in case of large energy transfer is characterized by large $Q^2$ and thus interaction is a deep inelastic QCD scattering.

Thus, only UHE gluino from low-mass window $1\text{ GeV} \leq m_{\tilde{g}} \leq 4\text{ GeV}$ could be a candidate for observed UHE particles, but it is disfavored by the arguments given above.

V. CONCLUSIONS

Topological Defects naturally produce particles with extremely high energies, much in excess of what is presently observed. However, the fluxes from most known TD are too small. So far only necklaces [23] and monopole-antimonopole pairs [24] can provide the observed flux of UHE CR.

Another promising sources of UHE CR are relic superheavy particles [24,25]. These particles should be clustering in the halo of our Galaxy [25], and thus UHE CR produced at their decays do not have the GZK cutoff. The signatures of this model are dominance of photons in the primary flux and Virgo cluster as a possible discrete source.

Apart from protons, photons and neutrinos the UHE carriers can be neutralinos [28], gluino [26–28] and monopoles [29,27]. While neutralino is a natural candidate for the Lightest Supersymmetric Particle (LSP) in SUSY models, gluino can be considered as LSP only phenomenologically. LSP are naturally produced in the parton cascade at the decay of superheavy X-particles. In case of neutralino both fluxes and cross-sections for interaction is somewhat lower than for neutrino. In case of gluinos the fluxes are comparable with that of neutralinos, but cross-sections for the production of observed extensive air showers are large enough only for light gluinos. These are disfavored, especially if the charged gluebarino is lighter than the neutral one [50].

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FIG. 1. Predicted proton (p) and gamma-ray (γ) fluxes from necklaces in comparison with experimental data.
FIG. 2: Predicted fluxes from decaying X-particles: nucleons (p, n, n) from the halo (curve $I^\text{halo}_N$), extragalactic protons (curve $I^\text{extra}_p$), photons from the halo (curve $I^\text{halo}_\gamma$), and neutrinos from the halo and the extragalactic space (curve $I^\text{tot}_\nu$).
FIG. 3: Predicted fluxes from cosmic necklaces with $r^2 \mu = 2 \times 10^{27}$ GeV$^2$: neutrinos (curve $I_\nu$), neutralinos (curve $I_\chi$) and $\gamma$-hadrons with mass $m_\gamma = 2.5$ GeV, $m_\chi = 2.0$ GeV and $m_\chi = 1.5$ GeV.