Ultrahigh Energy Cosmic Rays without Greisen-Zatsepin-Kuzmin Cutoff

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We study the decays of metastable superheavy particles as the source of ultrahigh energy cosmic rays (UHE CR). These particles are assumed to constitute a tiny fraction ξ_X of cold dark matter in the Universe. The UHE CR fluxes produced at the decays of X particles are calculated. The dominant contribution is given by fluxes of photons and nucleons from the halo of our Galaxy and thus does not exhibit the GZK cutoff. The extragalactic components of UHE CR are suppressed by the smaller extragalactic density of X particles and, hence, the cascade limit is relaxed. We discuss the spectrum of produced extensive air showers and a signal from a Virgo cluster as signatures of this model. [S0031-9007(97)04670-X]

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The observations of ultrahigh energy cosmic rays (UHE CR) reveal the presence of a new, isotropic component at energies $E \geq 1 \times 10^{10}$ GeV (for a review, see Ref. [1]). This component is thought to be extragalactic, since the galactic magnetic field cannot isotropize the particles of such energies. On the other hand, the observation of particles of the highest energies, especially of the two events with energies $2{-}3 \times 10^{11}$ GeV [2], contradicts the Greisen-Zatsepin-Kuzmin (GZK) cutoff [3] at $E \sim 3 \times 10^{10}$ GeV, which is the signature of extragalactic UHE CR. All known extragalactic sources of UHE CR, such as active galactic nuclei [4], topological defects [5], or the Local Supercluster [6], result in a well pronounced GZK cutoff, although in some cases the cutoff energy is shifted closer to 1×10^{11} GeV [6]. UHE neutrinos [7] could give a spectrum without cutoff, but the neutrino fluxes and the neutrino-nucleon cross section are not large enough to render the neutrino a realistic candidate for the UHE CR events.

In this Letter, we propose a scenario in which the UHE CR spectrum has no GZK cutoff and is nearly isotropic. 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 ξ_X is expected to be the same in all structures. In particular, ξ_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 cascade radiation [18] is suppressed. Long-lived massive relic particles were already discussed in the literature as a source of high energy neutrino radiation [9]. However, in our case the particles must be much heavier ($m_X \sim 10^{13}{-}10^{16}$ GeV).

The plan of our paper is as follows. First, we take a phenomenological approach and treat the density n_X of X particles and their lifetime τ_X as free parameters fixed only by the requirement that the observed UHE CR flux is reproduced. We calculate the fluxes of nucleons, photons, and neutrinos, considering the production of cascade radiation, positrons, antiprotons, and radio fluxes as constraints. We then discuss how the required properties of X particles can be realized.

The decays of X particles result in the production of nucleons with a spectrum $W_N(m_X, x)$, where m_X is the mass of the X particles and $x = E/m_X$. The flux of nucleons $(p, \overline{p}, n, \overline{n})$ from the halo and extragalactic space can be calculated as

$$I_N^i(E) = \frac{1}{4\pi}\,\frac{n_X^i}{\tau_X}\,R_i\,\frac{1}{m_X}\,W_N(m_X, x)\,, \tag{1}$$

where index i runs through h (halo) and ex (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 $m_X n_X^h = \xi_X \rho_{\rm CDM}^h$ and $m_X n_X^{\rm ex} = \xi_X \Omega_{\rm CDM} \rho_{\rm cr}$, where ξ_X describes the fraction of X particles in CDM, $\Omega_{\rm CDM}$ is the CDM density in units of the critical density $\rho_{\rm cr}$, and $\rho_{\rm CDM}^h \approx 0.3$ GeV/cm^3 is the CDM density in the halo. We shall use the following values for these parameters: a large DM halo with $R_h = 100$ kpc (a smaller halo with $R_h = 50$ kpc is possible, too), $\Omega_{\rm CDM} h^2 = 0.2$, $h = 0.6$, the mass of X particle in the range 10^{13} GeV $< m_X < 10^{16}$ GeV, and 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 $r_X = \xi_X t_0/\tau_X$. Following [10], we shall use at small X the QCD fragmentation function in the modified leading logarithm approximation (see [11])

$$W_N(m_X, x) \approx \frac{K_N}{x}\,\exp\!\left(-\frac{\ln^2 x/x_m}{2\sigma^2}\right), \tag{2}$$

where

$$2\sigma^2 = (1/6)\,(\ln m_X/\Lambda)^{3/2},$$

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$x = E/m_X$, $x_m = (\Lambda/m_X)^{1/2}$, and $\Lambda = 0.234$ GeV. The normalization constant K_N is found from energy conservation as

$$K_N \int_0^1 dx\, x W_N(m_X, x) = f_N\,,$$

where f_N is the fraction of energy transferred to nucleons. Using Z^0 decay as a guide, we assume $f_N \approx 0.05 f_\pi$, where f_π is the corresponding fraction for pions (e^+e^--collider LEP gives 0.027 for $p\overline{p}$ only). For the attenuation length of UHE protons due to their interactions with microwave photons, we use the values given in the book [12].

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

$$I_\gamma^h(E) = \frac{1}{4\pi}\frac{n_X}{\tau_X} R_h N_\gamma(E)\,, \tag{3}$$

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

$$N_\gamma(E) = \frac{2K_{\pi^0}}{m_X}\int_{E/m_X}^1 \frac{dx}{x^2}\exp\left(-\frac{\ln^2 x/x_m}{2\sigma^2}\right). \tag{4}$$

The normalization constant K_{π^0} is again found from the condition that neutral pions take away the fraction $f_\pi/3$ of the total energy m_X.

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 $E > 1 \times 10^{10}$ GeV, on the radio background. The neutrino flux calculation is similar.

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

The most stringent constraint comes from electromagnetic cascade radiation, which is initiated by high-energy photons and electrons from pion decays and is developing due to interaction with low energy background photons. The relevant calculations were performed in Ref. [8]. In our case, this constraint is weaker, because the low-energy extragalactic nucleon flux is $\sim$4 times smaller than that from the Galactic halo (see Fig. 1). Thus, the cascade radiation is suppressed by the same factor.

The relevant parameter which characterizes the flux of cascade radiation is the total energy density of cascade radiation ω_{cas}. The observation of the low-energy diffuse gamma-ray flux results in the limit $\omega_{\text{cas}} < 1 \times 10^{-5}$–$1 \times 10^{-6}$ eV/cm^3 [8]. In our case, the cascade energy density calculated by integration over cosmological epochs (with the dominant contribution given by the present epoch $z = 0$) yields

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

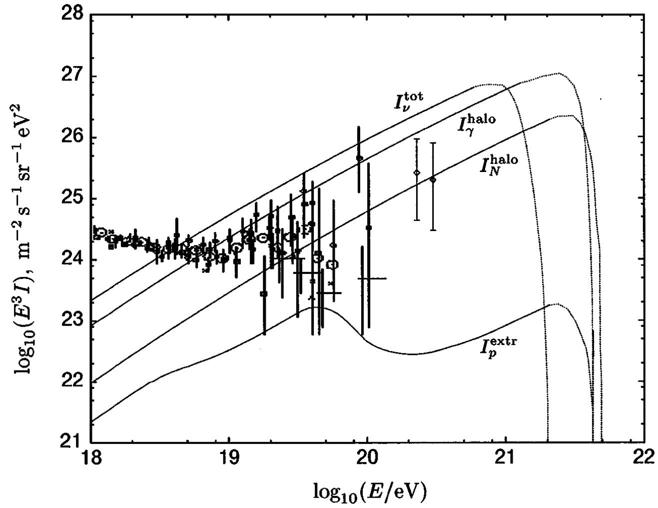

FIG. 1. Predicted fluxes from decaying X particles: nucleons $(p, \overline{p}, n, \overline{n})$ from the halo (curve I_N^{halo}), extragalactic protons (curve I_p^{extr}), photons from the halo (curve I_γ^{halo}), and neutrinos from the halo and the extragalactic space (curve I_ν^{tot}). The data points are based on the compilation made in Ref. [22].

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

The other constraints come from the observed fluxes of positrons and antiprotons in our Galaxy and from the extragalactic component of the radio flux. We performed detailed calculations which will be published elsewhere. In all cases, these constraints are satisfied and they are weaker than that due to cascade gamma radiation.

Now we address the elementary-particle and cosmological aspects of a supermassive, long-lived particle. Can the relic density of such particles be as high as required in our calculations? And can they have a lifetime comparable or larger than the age of the Universe?

Let us assume that the 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 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 analog of R parity (see [13] and the references therein). Thus, one can assume that the decay of X particles 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

$$\mathcal{L} \sim \frac{1}{m_{\text{Pl}}}\overline{\Psi}\nu\phi\phi\exp(-S)\,, \tag{6}$$

where Ψ describes X particles, and ϕ is a SU(2) scalar with a vacuum expectation value $v_{EW} = 250$ GeV. After spontaneous symmetry breaking the term (6) results in the mixing of the X particle and the neutrino, and the lifetime due to $X \to \nu + q + \overline{q}$, e.g., is given by

$$\tau_X \sim \frac{192(2\pi)^3}{(G_F v_{EW}^2)^2} \frac{m_{\rm Pl}^2}{m_X^3} e^{2S}, \tag{7}$$

where G_F is the Fermi constant. The lifetime $\tau_X > t_0$ for X particles with $m_X \geq 10^{13}$ GeV requires $S > 44$. This value is within the range of the allowed values as discussed in Ref. [14].

Let us now turn to the cosmological production of X particles with $m_X \geq 10^{13}$ GeV. Several mechanisms can be considered, including thermal production at the reheating stage, production through the decay of the inflation field at the end of the "preheating" period following inflation, and through interactions and decays of various topological defects.

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 values below 10^8–10^{10} GeV, depending on the gravitino mass (see Ref. [15] 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 [16]. 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 $a + \overline{a} \to X + \overline{X}$ is easily estimated as

$$n_X(t_R) \sim N_a n_a^2 \sigma_X t_R \exp(-2m_X/T_R), \tag{8}$$

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 σ_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 $m_X = 1 \times 10^{13}$ GeV and ξ_X in the wide range of values 10^{-8}–10^{-4}, the required reheating temperature is $T_R \sim 3 \times 10^{11}$ GeV.

In the second scenario mentioned above, nonequilibrium 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 topological defects, such as strings or textures (for a review of defects, see [17]). 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 ν 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 the X-particle density is given by the earliest epoch, soon after the 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 $\eta \gtrsim m_X$ 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 ξ_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.

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

It is easy to verify that the extragalactic nucleon flux at $E \leq 3 \times 10^9$ GeV is suppressed by a factor of $\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. [18]). This flux is not shown in the graph. The flux of high energy gamma radiation from the halo is by a factor of 7 higher than that of nucleons, and the neutrino flux (given in Fig. 1 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 extensive air showers (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 $E > 10^9$ GeV have the low-energy muon component as abundant as that for nucleon-induced showers [19]. However, the shower production by the photons is suppressed by the Landau-Pomeranchuk-Migdal (LPM) effect [20] and by absorption in the geomagnetic field (for recent calculations and discussion, see [8,21] and references therein). These effects are energy dependent. The LPM effect starts at 10^9–10^{10} GeV and it almost fully suppresses the production of "normal" EAS at $E_\gamma \geq 1 \times 10^{12}$ GeV, when maximum EAS reach the see level practically for all zenith angles [8]. The calculation of the spectrum of EAS is outside the scope of this paper and the normalization of the halo nucleon spectrum by observational data at $E \sim 2 \times 10^{11}$ GeV in Fig. 1 has an illustrative character. The general tendency of greater suppression of photon-induced showers with increase of energy might improve the agreement between calculated and observed spectra.

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

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

$$F_{p,\overline{p}}^{\mathrm{Virgo}} = r_X \frac{M_{\mathrm{Virgo}}}{t_0 R^2 m_X^2} W_N(m_X, x) . \tag{9}$$

The ratio of this flux to the diffuse flux from the hemisphere is 6.4×10^{-3}. This signature becomes less pronounced at smaller energies, when protons can be strongly deflected by intergalactic magnetic fields.

This signature, as well as our prediction of the absence of the GZK cutoff up to very high energies, can be tested by the future Auger detector [23].

When our work was in progress, we learned that a similar idea was put forward by V. A. Kuzmin and V. A. Rubakov [24]. The main difference is that we take into account the radiation from the galactic halo, which is the main issue of our Letter, while the above authors limited their consideration to the extragalactic component. We are grateful to V. A. Kuzmin and V. A. Rubakov for interesting discussions. M. K. was supported by the Alexander von Humboldt-Stiftung and A. V. was supported in part by the National Science Foundation.

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