Cygnus X-3: a cosmic source of glueballinos

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The nature of the 10^{15}–10^{16} eV neutral primary radiation recorded for the galactic source Cygnus X-3 is discussed. It is argued that these particles cannot be γ-rays, neutrinos, or photinos. It is suggested that they are the bound states of a gluino and a gluon (glueballino) with a mass of the order of 2–4 GeV. The properties of glueballinos are discussed, and it is shown that the above hypothesis is not inconsistent with data deduced from accelerator experiments and those obtained in cosmology.

I. INTRODUCTION

The flux of 10^{14}–10^{16} eV neutral radiation recorded from Cygnus X-3 (Ref. 1) is usually interpreted[2] as consisting of γ-rays. However, this natural interpretation encounters two difficulties. The first is that the atmospheric showers in which the Cygnus X-3 radiation has been recorded contain a relatively large number of muons, whereas showers due to γ-rays should be depleted in muons. The number of muons in such showers in the direction of Cygnus X-3 was measured in Ref. 3 and it was found that it was the same as the usual number in nuclear showers: N_μ (Cyg X-3)$/N_\mu(p) \geqslant 0.8$. On the other hand, in showers initiated by γ-rays, the corresponding ratio is much smaller: $N_\mu(\gamma)/N_\mu(p) < 0.1$ (Ref. 4). The Akeno installation[5] has been used in a search for radiation from Cygnus X-3 in muon-depleted showers, i.e., showers generated by γ-rays. Of the 18 recorded events in the direction of Cygnus X-3, seven were classified in Ref. 5 as belonging to the signal. The events were distributed within a very wide phase range (0.45–0.8). If these events can be interpreted as a real effect and not simply as the upper limit, the flux corresponding to energies $E \geqslant 10^{15}$ eV is $(1.1 \pm 0.4) \times 10^{-14}$ cm^{-2} s^{-1}, whereas the flux at the same energy in the phase interval 0.55–0.65, measured by the Haverah Park installation in the same year (1984), was 7×10^{-14} cm^{-2} s^{-1} (see Ref. 1e). This means that the γ-ray flux recorded by the Akeno system was several times smaller than the total flux.

The other difficulty involves the expected absorption of γ-rays with energy $E \simeq 2 \times 10^{15}$ eV in the Galaxy as a result of interactions with primordial photons: $\gamma + \gamma_p \to e^+ + e^-$ (Ref. 6). If we accept the current lower limit for the distance to Cygnus X-3, namely, $r \geqslant 13$ kpc, the absorption of γ-rays should lead to a reduction in the flux by a factor of 5 at $E \sim 2 \times 10^{15}$ eV. This valley has not been seen in the spectrum of Cygnus X-3.

Two other known particles, namely, the neutron and the neutrino, again cannot be the primary particles incident on the Earth from Cygnus X-3. The lifetime of neutrons with energies of 10^{15}–10^{16} eV is too short to allow them to reach the Earth (at $E = 10^{15}$ eV, the time taken to traverse the distance between Cygnus X-3 and the Earth is about 2000 neutron half-lives).

If we suppose that showers with energies in excess of 10^{15} eV are generated by neutrinos incident on the Earth, Cygnus X-3 will have to produce a flux of about 10^{-7} cm^{-2} s^{-1} on the Earth in order to account for the observed ground-level flux of extensive atmospheric showers (AES) of 7×10^{-14} cm^{-2} s^{-1}. (We are assuming that, at $E = 10^{15}$ eV, the νN interaction cross section is $\sigma_N \simeq 10^{-33}$ cm^2. This estimate takes into account the fact that, in this energy range, the linear increase in the cross section with energy is no longer valid.) If this neutrino flux were produced in the source by known mechanisms, i.e., $p + N \to \pi + X$, $\pi \to \mu\nu$, and $\mu \to e\bar{\nu}\nu$ reactions, it would require the emission of $L \sim 10^{44}\ \Omega/4\pi$ erg s^{-1} by accelerated protons in the Cygnus X-3 source, where Ω is the solid angle into which the protons radiate. A luminosity of this order is out of the question.

The neutrinos incident on the Earth take part in the reactions $\nu + N \to \mu + X$ (if they are muon neutrinos) or $\nu + N \to \tau + X$ and $\tau \to \nu_\tau \mu \nu_\mu$ (if $\nu = \nu_\tau$) and create high-energy muons underground, which can be recorded by underground installations (IMB,[7] NUSEX,[8] Soudan,[9] and the Baksan scintillation-counter telescope[10]). A neutrino flux of the above order should give rise to a high muon counting rate in the IMB and Baksan installations in the horizontal direction (3000 muons per annum for $\nu = \nu_\mu$ and 600 muons per annum for $\nu = \nu_\tau$ in the IMB installation), which, of course, would have been noticed. If $\nu = \nu_\mu$, the resulting muon flux due to the neutrions would exceed by a factor of 40 the NUSEX muon flux from Cygnus X-3, and should give rise to a weak dependence of the counting rate on the zenith angle, i.e., on the thickness of the ground traversed by the neutrinos, whereas experiment[8] shows that the muon flux decreases with increasing zenith angle. The conflict between data obtained in underground experiments and the consequences of the neutrino hypothesis could be avoided if the flux incident on the Earth were to consists of electron neutrinos alone. However this is difficult to imagine, even when possible neutrino oscillations are taken into account, because the source must generate both muon and electron neutrinos in comparable amounts.

Stenger[11a] has put forward the qualitative hypothesis that high-energy photinos from Cygnus X-3 could explain both the muon flux recorded by the Sudan and NUSEX detectors and the extensive atmospheric showers at 10^{15}–10^{16} eV. However, it can be shown that this hypothesis requires a proton luminosity that is too high for a galactic source ($L_p \sim 10^{44}$ erg·s^{-1}, Ref. 11b). Moreover, a muon neutrino

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flux would unavoidably accompany the photino flux, and would lead to a disagreement with underground detector experiments analogous to those discussed above.

The main reason for the contradiction that has arisen is that the photino-nucleon cross section is small, namely, $\sigma_{\tilde{\gamma}N} \simeq 10^{-32}\,(40\ \mathrm{GeV}/m_{\tilde{q}})^2$ (Ref. 12), where $m_{\tilde{q}}$ is the mass of the scalar quark ($m_{\tilde{q}} > 20$ GeV, Ref. 13), so that the shower imitation probability in the atmosphere is $W \sim 10^{-5}$ instead of the $W \sim 1$ for a photon.

In this paper, we consider the hypothesis that the primary AES particles with energies of 10^{14}–10^{16} eV that arrive on the Earth from Cygnus X-3 are the bound states of the gluino and gluon, to which we shall refer as glueballinos.

The idea of the glueballino flux from Cygnus X-3 was put forward by us in Ref. 14 and independently by Auriemma *et al.*[15] The present paper is a somewhat extended version of the preprint in Ref. 14.

2. BASIC IDEA

Let us suppose that the neutral particles with energies of 10^{14}–10^{16} eV that arrive on the Earth from Cygnus X-3 are stable or quasistable bound states of the gluino $\tilde{g}$ and gluon g (glueballinos $\tilde{G} = \tilde{g}g$). To ensure that the glueballino does not decay during its flight, the gluino must be the lightest (with the exception, possibly, of the gravitino) supersymmetric particle. In particular, the gluino must be lighter than the photino $m_{\tilde{g}} < m_{\tilde{\gamma}}$) and, moreover, $\tilde{G}$ must be lighter than $\tilde{g}\bar{q}q$ (see below for a discussion of this point). The mass of the gluino can be chosen on the basis of the following considerations. On the one hand, we must have a high enough cross section for the creation of gluon pairs in hadronic collisions, which sets an upper bound for m_g ($m_{\tilde{g}} < 3$–5 GeV). On the other hand, $m_{\tilde{g}} > 1$ GeV because the gluino production cross section for $m_{\tilde{g}} \lesssim 1$ GeV is greater than the cross section for the production of charm, and it seems unlikely that such a light glueballino would have remained unnoticed in laboratory experiments. The calculations given below refer to the case $m_{\tilde{g}} = 3$ GeV.

The idea underlying our hypothesis can readily be understood if we assume, as a crude approximation (more accurate estimates will be given below), that the cross section for the production of the gluino in *p-p* interactions and the cross section for the gluino-nucleon interaction can be characterized by $\sigma \sim 1$ mbarn. The range of the glueballino in matter is then $x \sim 1000\ \mathrm{g\cdot cm^{-2}}$, and the probability that a shower will be initiated in the atmosphere is $W \sim 1$. If the thickness of the "target" in Cygnus X-3 is $x \sim 1000\ \mathrm{g\cdot cm^{-2}}$, the glueballinos will freely leave the source, but photons will be strongly absorbed by it. The luminosity of the source will therefore be greater than that in the usual "photon model" by the factor

$$\sigma(pp \to \pi^0 X)/\sigma(pp \to \tilde{g}\tilde{g}X) \sim 10$$

(a more accurate estimate will be given below), i.e., by a relatively small factor.

3. THE PROPERTIES OF LIGHT GLUEBALLINOS

Let us first consider the status of the light, long-lived (stable or quasistable) gluino. This question has already been discussed[16] from another standpoint. To ensure that a glueballino with $E \sim 10^5$–10^6 GeV will traverse the distance $r = 13$ kpc without decaying, the gluino must be either a stable particle or a long-lived particle with a lifetime $\tau > 10^6\,(m_{\tilde{g}}/\mathrm{GeV})$s. Such a long lifetime excludes the possibility of the decay of the gluino to any known particles other than the gravitino: the gluino must be either the lightest supersymmetric particle or it can be heavier only than the gravitino. This possibility does not arise in the simplest supersymmetric models: the mass $m_{3/2}$ of the gravitino in such models is of the order of the mass m_W of the W boson and the gluino mass is greater than the photino mass. However, neither of these conditions is essential. For example, the inequality $m_{\tilde{g}} > m_{\tilde{\gamma}}$ may not be necessary if the supersymmetric theory is described not by a simple Lie group, but by a product of such groups, so that the photino and the gluino are gauge fields of different groups. Next, the global supersymmetry-breaking parameter may not be related to the gravitino mass $m_{3/2}$ (Ref. 17), in which case $m_{3/2} \ll m_W$ becomes admissible.

If the gravitino is the lightest supersymmetric particle, the gluino may decay into a gluon and a gravitino. The Hamiltonian for the decay interaction has the form[18]

$$H = \frac{1}{4M}\bar{\xi}^a \gamma_\mu \sigma_{\lambda\nu} \psi_\mu G_{\nu\lambda}{}^a + \mathrm{H.c.}, \tag{1}$$

where $M = M_{pl}/(8\pi)^{1/2} = 2.4 \times 10^{16}$GeV, ξ^a is the gluino field (Majorana spinor), ψ_μ is the gravitino field, and $G^a_{\nu\lambda}$ is the gluino field ($a = 1,\ldots,8$). The decay width of the gluino, calculated with this Hamiltonian, is

$$\Gamma = (2\pi M^2 m_{\tilde{g}}^3)^{-1}(m_{\tilde{g}}^2 - m_{3/2}^2)^3(1 + m_{\tilde{g}}^2/3m_{3/2}^2). \tag{2}$$

Here, we note that Γ tends to infinity as $m_{3/2} \to 0$. The reason for this is as follows. It is readily verified that, for unbroken symmetry, when the gluino is massless, the current $\eta_\mu = \bar{\xi}^a \gamma_\mu \sigma_{\nu\lambda} G^a_{\nu\lambda}/4M$ corresponding to the emission of a gravitino in (1) is conserved ($\partial_\mu \eta_\mu = 0$), so that no longitudinal gravitino is emitted and the theory does not contain infrared divergences. When spontaneous symmetry breaking ensures that the gluino acquires mass, the current η_μ ceases to be conserved ($\partial_\mu \eta_\mu \sim m_g$), the longitudinal gravitinos begin to be emitted, and the theory acquires infrared singularities in $m_{3/2}$ that are proportional to $m_{\tilde{g}}$.

When $m_{\tilde{g}} = 3$ GeV and $m_{3/2} \geqslant 10$ MeV, the lifetime of the gluino is $\tau \gtrsim 1\ y$. A gluino with this lifetime and $E \geqslant 10^5$ GeV would not decay to any appreciable extent along its path between Cygnus X-3 and the Earth.

There are cosmological limitations on the light gravitino considered here ($m_{3/2} > 10$ meV). In noninflationary models, gravitino masses in the range $1\ \mathrm{keV} < m_{3/2} < 10^4$ GeV are forbidden[19] because the generation of gravitinos at the early stages of the hot Universe leads to energy densities above the critical value. Exponential expansion in the inflationary scenario gives rise to a lower density of these gravitinos, but repeated heating of the Universe at the end of the inflationary stage leads to the regeneration of gravitinos (see Ref. 18 and the references cited therein). The temperature T_r after repeated heating must not be too high because, otherwise, the gravitinos or their decay products would give rise to effects in conflict with observations. The strongest limitation on the maximum admissible temperature in repeated heating is obtained for the unstable gravitino as a consequence of the splitting of deuterium nuclei by photons from gravitino decays. For an unstable gravitino with mass $m_{3/2} \sim 100$ GeV, this gives rise to a serious problem because the temperature in the case of repeated heating must be less than 10^9–10^{10} GeV (Ref. 18), and such a low temperature will not ensure the necessary baryon excess due to the decays of heavy Higgs particles. For the light stable gravitino, there are less stringent restrictions on the maximum temperature in repeated heating. The most stringent limitation is obtained for nonrelativistic gravitinos from the nuclear fusion reaction producing helium: $T_r < 5 \times 10^{17}$ GeV for $m_{3/2} = 2$ GeV ($T_r^{\max}$ increases with decreasing $m_{3/2}$). Thus, the admissible temperature in repeated heating allows the generation of a baryon asymmetry due to the decay of heavy Higgs particles.

For a gravitino with $m_{3/2} < 10$ MeV the decay of the gluino along the path between Cygnus X-3 and the Earth is significant, and our hypothesis will not work in this case.

The gluino forms colorless hadrons with gluons and quarks. It is expected that the lightest of these states is the bound state of the gluino and gluon. In particular, it is lighter than the state $\tilde{g}\bar{q}q$. The argument for this is that the dimension of the operator $G^a_{\mu\nu}\sigma_{\mu\nu}\xi^a$ (G^a is the gluon operator and ξ^a is the gluino operator) corresponding to the $\tilde{g}g$ system is less than the dimension of the operator $\bar{q}\lambda^a\gamma_\mu q\gamma_\mu\xi^a$ corresponding to the system $\tilde{g}\bar{q}q$. Within the framework of the operator expansion approach (the QCD sum rules), this relation must lead to[1] $m_{\tilde{g}} < m_{\tilde{g}\bar{q}q}$. It will be seen from the discussion presented below that the existence of a light stable (or quasistable) gluino will not be in conflict with existing experimental data if there are no bound states of the glueballinos with nucleons or nuclei. We shall assume that this is actually so. This situation will arise, in particular, if there are repulsive forces between the glueballino and the nucleon at low energies. There are some arguments in favor of the latter possibility. Consider the $\tilde{G}N$ forward-scattering amplitude $T(s)$ at high energies. Since $\tilde{G}$ is a flavor singlet, amplitude $T(s)$ is determined by the exchange of P and P' poles with the framework of the Regge approach. The signs of the residues at the P and P' poles are the same for all known particles, which is usually explained by saying that the coupling constants between P and P' are proportional to the matrix elements of the energy-momentum tensor. We shall suppose that this is so for the glueballino, i.e., the signs of the interaction constants between glueballino and P and P' are the same. We shall take $T(s)$ in the form

$$T(s) = -\sum_{i=P,\,P'} \frac{1+\exp(-i\pi\alpha_i)}{\sin\pi\alpha_i} s^{\alpha_i}\beta_{i\tilde{G}}\beta_{iN}, \tag{3}$$

where α_i are the intercepts, $\alpha_P = 1$, $\alpha_{P'} = 1/2$, and β_i are the residues of the P and P' trajectories. For large s,

$$\operatorname{Im} T(s) = \beta_{P\tilde{G}}\beta_{PN}s > 0$$

and

$$\operatorname{Re} T(s) = -\frac{1+\cos\pi\alpha_{P'}}{\sin\pi\alpha_{P'}}\beta_{P'\tilde{G}}\beta_{P'N}. \tag{4}$$

In accordance with the foregoing, it follows from (4) that

$$\operatorname{Re} T(s) < 0. \tag{5}$$

If this inequality could be extrapolated to low energies, it would imply a positive $\tilde{G}N$ potential, i.e., a repulsion. Of course, this discussion merely shows that the assumed absence of resonances in the $\tilde{G}N$ system is not in conflict with the behavior of the $\tilde{G}N$ scattering cross sections at high energies. We note that our discussion, as given above, is not valid for the system $\tilde{g}\bar{q}q$ scattering may involve the exchange of other Regge poles, apart form P and P'.

We now turn to another case in which the gravitino is much heavier than the gluino, i.e., the case of the stable gluino. Stable gluinos should be created in the Universe at temperature $T \gtrsim m_{\tilde{g}}$ and should exist at present in the form of primordial particles. Their density in the present Universe can be calculated as the primordial concentration of stable neutral hadrons,[21] which is $n \sim 10^{-11}\, n_B$ independently of the mass, where n_B is the density of baryons. If the glueballinos were to form bound states with nucleons or nuclei, this concentration would exceed the upper bound for the density of anomalous isotopes, discussed in Ref.22. Since there are no such states, it is very difficult to detect glueballinos with densities $n \sim 10^{-11}\, n_B$.

Let us now estimate the cross section for the interaction between glueballinos and nucleons at high energies. We shall compare the $\tilde{G}N$ and πn (or ρN) scattering cross sections and will assume, as is often done in estimates of scattering cross sections for charmed particles, that

$$\sigma(\tilde{G}N)/\sigma(\pi N) \approx (r_{\tilde{G}}/r_\pi)^2, \tag{6}$$

where $r_{\tilde{G}}$ and r_π are the glueballino and pion radii, respectively. We shall use the constituent model to estimate $r_{\tilde{G}}/r_\pi$. We then have

$$r_{\tilde{G}}/r_\pi \approx \alpha_{sq}(m_q/2)/\alpha_{sg}\mu, \tag{7}$$

where α_{sq} and α_{sg} are the strong interaction constants in the $\bar{q}q$ and $\tilde{g}g$ systems, $\alpha_{sq}/\alpha_{sg} = 4/9$, $m_q = 350$ MeV is the mass of a constituent u or d quark, and $\mu = m_g m_{\tilde{g}}/(m_g + m_{\tilde{g}})$ is the reduced mass of the gluino and the constituent gluon. Assuming that $m_{\tilde{g}} = 3$ GeV and $m_g = 0.6$ GeV, we find from (7) that $r_{\tilde{G}}/r_\pi \simeq 1/7$, so

$$\sigma(\tilde{G}N)/\sigma(\pi N) = {}^1/_{50}, \tag{8}$$

that[2] $\sigma(\tilde{G}N) \simeq 1/2$ mbarn. It is, of course, important to bear in mind that the above estimate is very crude and can be criticized on a number of counts. For example, when we derived (7), we did not take into account the fact that α_s was a function of distance, i.e., the ratio $\alpha_{sq}(r_\pi)/\alpha_{sg}(r_{\tilde{G}})$ must be greater than 4/9. This effect may make the right-hand side of (8) several times greater. However, the true $\tilde{G}N$ scattering cross section can hardly exceed a few millibarns. We also not that the low $\tilde{G}N$ scattering cross section is an argument in favor of the absence of resonances in the $\tilde{G}N$ system.

The fact that the gluino in the glueballino takes practically no part in processes with high energy loss is an important feature of glueballino-nucleon scattering. To verify this, it will suffice to note that, when the gluino is scattered, the square of the miminum transferred momentum is given by

$$Q^2_{min} = m^2_{\tilde{g}}(E'/E + E/E' - 2), \tag{9}$$

where E and E' are the energies of the gluino before and after scattering, respectively. For the heavy gluino ($m_{\tilde{g}} \simeq 3$ GeV) and the high energy loss $(E - E')/E \sim 1/2$, the momentum $Q^2_{\min}$ turns out to be high ($Q^2_{\min} \sim m^2_{\tilde{g}}$), which leads to a small gluino scattering cross section. It is readily seen that these considerations will not apply to the lighter gluon ($m_g \simeq 0.6$ GeV) in the glueballino. Thus, it may be considered in a crude approximation that the inelastic interaction between the glueballino and the nucleon involves the participation of only the gluon in the glueballino. If we use this to estimate the cross section for $\tilde{G}N$ scattering, and assume that the hadron cross sections are proportional to the number of active constituent components of the hadron, we find that $\sigma(NN):\sigma(\pi N):\sigma(\tilde{G}N) = 3:2:1$ and $\sigma(\tilde{G}N) \sim 10$ mbarn. This seems to be the upper limit. We have seen that the scattering of the gluino in the glueballino cannot lead to a large energy loss. On the other hand, the scattering of the constituent gluon cannot ensure a large energy loss either. The fact is that, if the gluon is scattered through a large angle, the loss of longitudinal momentum by the glueballino will be $\Delta p_{\parallel}/p_{\parallel} \sim m_g/(m_{\tilde{g}} + m_g) < 0.2$. We may therefore suppose that the high-energy glueballino loses about 10–20% of its energy in the $\tilde{G}N$ collision.

We must now consider the upper bound for the mass of quasistable (stable) gluino, which can be deduced from accelerator experiments. First of all, the fact that the glueballino is neutral is of cardinal importance from this point of view. If the lowest bound state of the light gluino and the gluon or quark were charged (like, for example, the system $\tilde{g}\bar{q}q$), it would have been discovered long ago.

It is quite clear that the restrictions deduced from the beam-dump experiment and the experiment involving the search for unbalanced transversed jets in colliding-beam systems have been completely removed: in the former case, the gluino does not generate penetrating particles and, in the second, there are no jets that are balanced in $p_{\perp}$.

The glueballinos should be most noticeable in the experiment of Gustafson *et al.*[23] In this experiment, the proton beam was incident on a target in the form of short pulses, and secondary particles were recorded at a distance l in the detector of thickness $x = 900$ g cm^{-2}. The velocity and the Lorentz factor γ of a secondary particle could be determined by measuring the time of flight Δt, and the energy released in the detector δE was assumed equal to the particle energy E. The mass of the particle was found from the formula $m = [2c(\delta E)^2 \Delta t/l]^{1/2}$. The analysis reported in Ref. 23 was confined to events with $m \geqslant 2$ GeV because the neutron flux was too high for $m < 2$ GeV. If the detector in Ref. 23 were to intecept glueballinos, then for $\sigma(\tilde{G}N) \simeq 3$ mbarn, it should have recorded about two collisions in the detector, and for $\Delta E/E \simeq 0.2$ the energy release was $\delta E = 0.36E$. For this reason, the mass obtained in the experiment was smaller by a factor of three than the true value, i.e., the glueballino fell into the region of the neutron background.

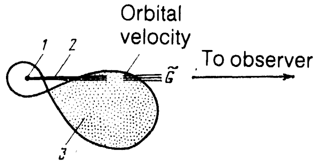

FIG. 1. Schematic diagram illustrating the model of Cygnus X-3. The pulsar 1 emits a beam of accelerated protons 2 in the direction of the observer. As the binary system rotates, the gas target low-density part of the Roche cavity of the massive companion 3) cuts the beam periodically, and the glueballinos are created in this region.

4. GENERATION OF GLUEBALLINOS IN CYGNUS X-3

The so-called hidden neutrino source, suggested in Ref. 24 and illustrated in Fig. 1, will now be adopted as a working model. The massive component ($M \sim 10 M_{\odot}$) filling the Roche cavity forms a binary system with the active pulsar which generates the beam of accelerated protons pointing toward the observer. Gluino generation ($p + p \to \tilde{g} + \tilde{g} + X$) occurs in the peripheral parts of the Roche cavity when the thickness of matter along the line of sight is of the order of the range of the generated particles. The number of protons with kinetic energy E emitted per second by the pulsar is

$$N_p(E)\,dE = (\gamma - 1)\gamma\left(\frac{E}{E_0} + 1\right)^{-(\gamma+1)} \frac{L_p}{E_0}\frac{dE}{E_0}, \tag{10}$$

where γ is the exponent of the integrated spectrum, E_0 is the normalizing energy, and L_p is the luminosity of the pulsar due to accelerated protons with energies $E > E_0$. From now on, we take $\gamma = 1.1$, $E_0 = 1$ GeV, and E and L_p in units of GeV and GeV/c, respectively. The flux of glueballinos with energy E emitted by a source that is transparent to glueballinos can readily be deduced by analogy with the generation of neutrinos[25]

$$F_{\tilde{g}}(E) = \frac{\gamma(\gamma - 1)}{1 - \alpha^{\gamma}} \varphi_{\tilde{g}}(E) L_p E^{-(\gamma+1)} w, \tag{11}$$

where w is the fraction of time per period in which the glueballinos are emitted, i.e., when the proton beam falls on the target, $\alpha \simeq 0.5$ is the fraction of energy retained by the proton in an inlastic p, p collision, and $\varphi_{\tilde{g}}(E)$ is the gluino yield given by

$$\varphi_{\tilde{g}}(E) = 2\int_0^1 dx\, x^{\gamma} \frac{1}{\sigma_{in}} \frac{d\sigma(E/x, x)}{dx}, \tag{12}$$

where $\sigma_{in} \simeq 40$ mbarn is the inelastic p, p scattering cross section, $x = E/E_p$, and $d\sigma(E,x)/dx$ is the inclusive cross section for the production of gluinos in the $p + p \to \tilde{g} + X$ reaction. By analogy with Ref. 26, our calculations were performed in the fusion model, i.e., we took into account the reactions

$$g + g \to \tilde{g} + \tilde{g}, \quad \bar{q} + q \to \tilde{g} + \tilde{g}, \quad q + g \to \tilde{g} + \tilde{q}.$$

The nucleon structure functions determined in Ref. 27 were employed. The $p + p \to \tilde{g} + X$ cross section and the yield $\varphi_g(E)$ for $m_g = 3$ GeV are shown in Fig. 2. In particular, when $E = 10^6$ GeV, $\varphi_g = 6 \times 10^{-3}$. For comparison, we reproduce the yields for γ-rays (from neutral pion generation)

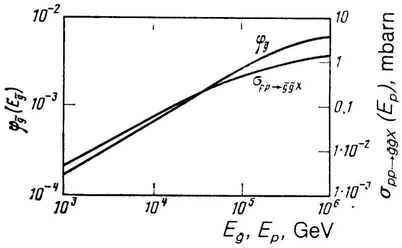

FIG. 2. $pp \to \tilde{g}\tilde{g}X$ reaction cross section (right-hand scale) as a function of the proton energy in the laboratory system, and the gluino yield as a function of the gluino energy.

and the neutrinos $\nu_\mu + \bar{\nu}_\mu$ (from the generation of charged pions and kaons): $\varphi_\gamma \simeq \varphi_{\nu\mu} + \varphi_{\bar{\nu}\mu} \simeq 0.12$ (Ref. 25). We note that we have calculated the yields for gluino production in the fusion model. In the case of charm, it is known[28] that the observed charmed hadron production cross section is several times greater than the theoretical value obtained in the fusion model, and the fraction of energy taken up by the charmed hadron is found to be much greater. We shall now assume (and, in our view, this is a very conservative assumption) that the yield $\varphi_{\tilde{g}}$ is greater by a factor of 2–3 than the prediction of the fusion mode, so that $\varphi_g/\varphi_\gamma \sim 0.1$.

To ensure that the glueballino flux exceeds the photon flux from the source, the range $x_{\tilde{G}}$ of the glueballinos in matter must be appreciably greater than the photon range $x_{\rm rad}$, and the target thickness x must be greater than the photon range. The smallness of $\varphi_{\tilde{g}}/\varphi$ is balanced in the ratio of the glueballino to photon fluxes by the relative absorption factor $k = \exp[x/x_{\rm rad} - x/x_{\tilde{G}}]$.

For the above glueballino-nucleon cross sections and the energy loss in $\tilde{G}N$ collisions, the glueballino range is much greater than the photon range. For example, when $\sigma(\tilde{G}N) \simeq 3$ mbarn and $\Delta E/E \simeq 0.2$, the glueballino range corresponding to the loss of half the energy is $x_{\tilde{G}} \simeq 1500$ g·cm^{-2} as compared with the photon range in hydrogen $x_{\rm rad} \simeq 63$ g·cm^{-2}. One can therefore readily imagine the conditions in the source under which the photons are completely absorbed, but practically all the glueballinos created there will escape from the source. The source luminosity due to high-energy protons must be greater than the "photon model" value by only an order of magnitude. By analogy with the usual assumption employed in the photon model, the absolute luminosity can be reduced by assuming that the emission of the proton beam (and, consquently, the glueballino beam) occurs within a relatively small solid angle Ω.

5. ATMOSPHERIC SHOWERS DUE TO CYGNUS X-3

When they enter the Earth's atmosphere, the glueballinos generate a nuclear cascade as a result of their strong interaction. The particular feature of the development of this cascade in the gradual injection of energy (two-fourfold, depending on the cross section) along the cascade development length. This produces a spreading of the shower maximum and superposition of electromagnetic cascades of different age at sea level.

The muon component of the shower in the reigon of low-energy muons (of the order of a few GeV) is approximately the same as in the usual nuclear showers, which is in agreement with observations.[1,3] It is thus clear that the glueballino hypothesis removes both difficulties encountered when it is assumed that the flux from Cygnus X-3 consists of γ-rays.

Showers containing 1000 GeV or less cannot be explained in terms of glueballinos. If we impose the condition that the difference between times of arrival of particles with energies E_1 and E_2 due to the different velocities does not spread the peak[1] on the phase histogram of duration $\tau \simeq 0.1T$, where $T = 4.8$ h is the source period, we find that the primary-particle mass is

$$m < E_2 \left[\frac{cT}{5r(E_2^2/E_1^2 - 1)} \right]^{1/2}, \qquad (13)$$

where $r = 13$ kpc is the assumed distance to the source. When $E_1 = 10^3$ GeV and $E_2 = 1.5 \times 10^3$ GeV, we find that $m < 68$ MeV ($m < 51$ MeV as $E_2 \to \infty$), whereas, for $E_1 = 10^6$ GeV and $E_2 = 1.5 \times 10^6$ GeV, result is $m < 68$ GeV ($m < 51$ GeV as $E_2 \to \infty$). It seems to us exceedingly likely that particles with $E \lesssim 1000$ GeV are γ-rays emitted by electrons in the magnetosphere of the pulsar in Cygnus X-3 because this radiation is observed for the pulsars in the Crab and Vela. Radiation with $E \sim 1000$ GeV should then be periodic, owing to the intrinsic rotation of the pulsar.

A separate and possibly unrelated problem is presented by high-energy muons recorded by the underground detectors Soudan[8] and NUSEX[9]). When these are described in terms of our hypothesis, it must be remembered that high-energy muons are created by glueballinos with a greater probability than by nucleons When it interacts with a nucleon, the glueballino may be peripherally excited and may undergo a transition to the system $\tilde{g}\bar{q}q$. The most probable decay of this system is into a glueballino and a pair of pseudoscalar mesons. Since the glueballino is a flavor singlet, this leads to an increase in the yield of leading kaons (which carry off 0.05–0.1 of the energy) as compared with NN collisions. On the other hand, charged kaons have shorter lifetimes and therefore produce muons more effectively than pions. This can be used to deduce the following estimate for the flux of muons with $E_\mu > 5$ GeV, generated by the glueballinos:

$$j_\mu(E > E_\mu) \approx \frac{0{,}64}{4} n_K \frac{\varepsilon_K}{2E_\mu} j_{\tilde{G}}(E > 20E_\mu) + \frac{1}{4} n_\pi \frac{\varepsilon_\pi}{1{,}25 E_\mu} j_{\tilde{G}}(E > 12{,}5E_\mu), \qquad (14)$$

where $n_K = 2$ is the K-meson multiplicity, 1/4 is the probability of a $\tilde{G}N$ collision with the creation of a K^+K^- pair, $\varepsilon_K \sim 1000$ GeV is the critical energy of a charged K meson, at which the decay range of a K meson is equal to the nuclear interaction length in the atmosphere, and $E_K = 2E_\mu = 10^4$ GeV is the energy of a K meson creating a muon of the required energy. Using $j_{\tilde{G}}(E > 1 \times 10^5 \text{ GeV}) = 2 \times 10^{-12}$ cm^{-2} s^{-1}, we find from (14) that j_μ $(E > 5 \times 10^5$ GeV$) \sim 1 \times 10^{-13}$ cm^{-2}s^{-1}, which is lower by a factor of 30 than the flux observed in the NUSEX experiment for $E > 5000$ GeV.

The glueballino hypothesis thus enables to explain the atmospheric showers observed for $E \sim 10^6$ GeV in the direction of Cygnus X-3, which characterized by the same muon component as ordinary nuclear-electromagnetic showers. The necessary luminosity L_p of the source due to acceleration protons must exceed the photo-model luminosity by only an order of magnitude. The impurity of high-energy photons depends on the source model or, specifically, on the time spent by the thin part of the target with $x = x_{\rm rad}$ in the line of sight. Although the glueballinos give rise to an enhanced muon flux and $E_\mu > 1$ TeV in showers, this mechanism does not appear to be sufficient to explain the high-energy muon flux observed in the NUSEX[8] and Soudan[9] experiments.

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[1)] Chanowitz and Sharpe,[20] whose work we unfortunately omitted from our preprint,[14] have calculated the $\tilde{g}g$ and $\tilde{g}\bar{q}q$ masses in the bag model. The masses obtained in Ref. 20 for certain particular values of the parameters are consistent with $m_{\tilde{g}g} < m_{\tilde{g}\bar{q}q}$. We note that the name glueballino, used here for $\tilde{g}g$, was proposed previously for this system in Ref. 20. We are indebted to M. Chanowitz for sending us a copy of Ref. 20.

[2)] Our discussion is confirmed by an analogous estimate of the $J/\psi,N$ scattering cross section. We have $r_{J/\psi}/r_\pi \simeq m_q/m_c \simeq 1/4$, so that $\sigma(J/\psi,N) \sim 1/16\sigma(\pi,N) \sim 2$ mbarn, in agreement with experiment.

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