Diffuse Cosmic Gamma-Ray Background as a Probe of Cosmological Gravitino Regeneration and Decay

Keith A. Olive
Astrophysics Theory Group, Fermilab, Batavia, Illinois 60510

and

Joseph Silk
Department of Astronomy, University of California, Berkeley, California 94720
(Received 22 March 1985)

We predict the presence of a spectral feature in the isotropic cosmic gamma-ray background associated with gravitino decays at high red shifts. With a gravitino abundance that falls in the relatively narrow range expected for thermally regenerated gravitinos following an inflationary epoch in the very early universe, gravitinos of mass several gigaelectronvolts are found to yield an appreciable flux of 1–10-MeV diffuse gamma rays.

PACS numbers: 98.70.Vc, 11.30.Pb, 14.80.Pb, 95.30.Cq

Decays of long-lived particles can have significant astrophysical implications. Decay products may produce far-ultraviolet background photons,[1] and even yield a uniform cosmological density of dark matter.[2] There is one class of particle, common to most phenomenological supersymmetry theories, whose long-lived decay has a significant effect that appears to have been hitherto overlooked. We shall show below that gravitino decays are likely to lead to an appreciable diffuse cosmic gamma-ray background. Decays of primordial gravitinos can also present a severe embarrassment to cosmological models, via effects on primordial nucleosynthesis and microwave background distortions.[3–8] One promising solution is that an epoch of inflation diluted the primordial gravitino abundance to an acceptable level.[9] Gravitinos are regenerated during the reheating process after inflation, and hence the reheating will be subject to the constraints from gravitino decays. Remarkably in this case, we find that the predicted flux of diffuse cosmic gamma rays resulting from gravitino decays in the early universe is similar to the observed gamma-ray background for a wide range of gravitino masses, provided that the maximum temperature to which the universe reheats after inflation lies in a acceptable range.

Despite all the uncertainties regarding the model-dependent particle spectrum in locally supersymmetric theories, there are at least two things which are common to all: (1) There is one stable supersymmetric (SUSY) particle and (2) one somewhat long-lived SUSY particle. By a long-lived SUSY particle (LSP), we mean one for which its decay rate is proportional to the gravitational constant $G_N = M_p^{-2}$. The gravitino is an example of such a SUSY particle and, depending on its mass, may indeed be very long lived.

The identity of the LSP is of course the subject of debate. In a wide class of models, the LSP is either the photino[10,11] or Higgs fermion (or mixed state of the two),[12] while other candidates are the scalar neutrino[13] and the SUSY partner of the axion.[14] In this Letter, we will assume that the LSP is the photino. Provided that the gravitino is not the LSP, then it will be unstable to decay into lighter SUSY particles. For simplicity, we will take the photino to be the only SUSY particle lighter than the gravitino. Other models in which, for example, both photinos and gluinos are lighter than the gravitino will not drastically affect our ensuing discussion. The decay rate for gravitino→ photon+photino has been calculated[11,6] and can be expressed as $\Gamma = 8\pi \alpha m_{3/2}^3 / M_P^2$ where the "coupling" α depends on the photino mass: $\alpha = (32\pi)^{-1}[1-(m_{\tilde{\gamma}}/m_{3/2})^2]^3$. Hence a gravitino with mass $m_{3/2} < 100$ GeV will have a lifetime $\tau > 4\times 10^8$ s. (The actual values of α differ in Refs. 6 and 11. We have chosen the value of Ref. 6 and treat α as a parameter to cover for this uncertainty, in addition to uncertainties coming from higher-order corrections.)

Because of their late decays into photons, the abundance of gravitinos (with respect to photons, for example) becomes very important. Various limits from entropy production,[4] overall mass density of the decay products,[11,12,8] distortions of the microwave background,[5,6,8] and the destruction of primordial deuterium place limits on the gravitino abundance Y which we define as $Y = n_{3/2}/n_\gamma$ where n_γ is the number density of photons today. The strongest of these limits implies that[7,8] $Y < 6\times 10^{-12}/m_{3/2}$. In a standard (noninflationary) model, one would expect that initially (before decoupling) $Y = 1$ and that through the annihilation of all other particle species the abundance would be brought down to $Y \sim 1/N_D \sim 10^{-2}$ today where N_D is the total number of degrees of freedom at gravitino decoupling. This leads to a catastrophe if gravitinos are long lived, since the expansion of the universe becomes dominated by nonrelativistic particles far too early. Inflation resolves[9] this problem by diluting the primordial gravitino abundance to a negligible level. However, after inflation, the Universe will reheat to some temperature T_R, during which additional gravitinos will be produced. One can estimate the secondary production of gravitinos as[12] $Y \simeq (\alpha/\sqrt{N})T_R/M_p \sim 2\times10^{-3}T_R/M_p$ for the gauge coupling $\alpha_G \sim \frac{1}{25}$ and $N \sim 300$. This agrees very closely with a more exact calculation[6] which yields $Y = 2.8\times10^{-3}T_R/M_p$. Thus the final abundance Y of gravitinos will be determined by T_R, and the limits on Y become limits on T_R.

Many models of supersymmetric inflation[15] generally predict a low value for T_R. In these models, the superpotential f (which generates the scalar potential for inflation) contains only one parameter whose value differs from $O(1)$. For example $f(\phi) = \mu^2 g(\phi)$ where all couplings in $g(\phi)$ are $O(1)$. The scalar potential is then $V \sim \mu^4$ and the mass of ϕ is $M_\phi \sim \mu^2 M_p$, etc. More importantly, μ also determines both the magnitude of density fluctuations produced during inflation[16] and the reheating temperature T_R (Ref. 15). In this type of model, the density fluctuations are given by $\delta\rho/\rho \sim 10^3\mu^2 \sim 10^{-4}$–$10^5$ so that $\mu^2 \sim 10^{-7}$–10^{-8} according to limits on the isotropy of the microwave background.[17] The reheating temperature, on the other hand, is given by $T_R \simeq M_\phi^{3/2}/M_p^{1/2} \simeq 0.2\mu^3 M_p \simeq (6\times10^{-12}$–$2\times10^{-13})M_p$ implying that the abundance of gravitinos is $Y \simeq 10^{-14}$–4×10^{-16}, below the limit given earlier.

Given the abundance Y of gravitinos and their decay rate, one can determine the time or red shift of decay. The age of the Universe at a given (recent, i.e., matter-dominated, regime) red shift ($>>1$) can be expressed as $t = t_0(1+z)^{-3/2}$, where the present age of the Universe is taken to be $t_0 = 2\times10^{17}h^{-1}\Omega^{-1/2}$ s, the Hubble parameter is defined as $H_0 = 100h$ km Mpc^{-1} s^{-1}, and $\Omega = \rho/\rho_c$, where the critical density $\rho_c = 1.88$ g/cm^3. The quantity $1+z$ is related to the background temperature by $T/T_0 = 1+z$ where $T_0 \simeq 2.7$ K is the present temperature. Gravitinos will decay when $\Gamma = t_D^{-1}$ or when

$$\alpha m_{3/2}^3 = 2\times10^{-5}h(1+z_D)^{3/2}\Omega^{1/2}\ \text{GeV}^3, \qquad (1)$$

where z_D is the red shift of decay. In the rest of the gravitino, if we define the energy of the photon produced in the decay to be $m_{3/2}/\gamma$ (with $\gamma \geqslant 2$) and β to be its energy today in megaelectronvolts, then

$$\beta = 10^3\gamma^{-1}m_{3/2}(1+z_D)^{-1}, \qquad (2)$$

where we are using units such that all masses are given in gigaelectronvolts.

It is now relatively straightforward to calculate the spectrum of γ's coming from the decay. At any time t the number of γ's is $n_\gamma^{(D)} = Yn(1-e^{-\Gamma t})$. If we define $x = \Gamma t = (1+z_D)/(1+z)$, we can write the differential flux as

$$\Phi^{(D)} = \frac{dn^{(D)}}{dE} = \frac{3}{8\pi}Yn_\gamma\frac{x^{1/2}}{\beta}\exp(-x^{3/2}). \qquad (3)$$

In this notation the photon energy is $E = x\beta$. The produced flux of γ's must now be compared to the observed flux at $E \sim 1$ MeV,

$$\Phi \simeq 1.2\times10^{-2}x^{-3}\beta^{-3}\ \text{cm}^{-2}\ \text{s}^{-1}\ \text{sr}^{-1}\ \text{MeV}^{-1}. \qquad (4)$$

One can check that the ratio of Eqs. (3) and (4) is maximized at $x = (\frac{7}{3})^{2/3} \simeq 1.76$. At this value of x the γ-ray flux produced by gravitino decays is comparable to the observed diffuse γ-ray flux for $Y \simeq 1.2\times10^{-14}\beta^{-2}$. We argue below that the $\sim$1-MeV γ-ray bump is due to gravitino decays. As one can see, for β in the range 1–10, the abundance of gravitinos agrees remarkably with the estimated abundance given by inflation and the predicted magnitude of density perturbations. From Eqs. (1) and (2) the inferred gravitino mass is $m_{3/2} \sim 1$–10 GeV. Unfortunately, there is no accepted source for the origin of the diffuse γ-ray background flux in the vicinity of 1 MeV, or for that matter at any other energy. However, it is likely that distant discrete sources make up the continuous part of the spectrum. Moreover, the quality of the existing data is such that one can do little more than model these data as the superposition of a continuous source with a feature of the general form given by Eq. (3). Hence, the best that we can do is to use the data to constrain the gravitino parameters which we will now go on to do.

A limit to the decay red shift from which an observable gamma flux can originate may be inferred from the absorption of the energetic decay photons by ambient intergalactic matter. In the energy range of interest for decays to yield gamma rays at present, the relevant absorption process near z_d is pair production. The pair-production cross section is given by

$$\sigma_{pp} \approx \tfrac{28}{9}\alpha_f r_e^2[\ln(2E_\gamma/m_0c^2) - \tfrac{109}{42}] \sim \alpha_f\sigma_T, \qquad (5)$$

where $\sigma_T = \frac{8}{3}\pi r_e^2$ (r_e is the classical electron radius) is the Thompson cross section. The absorption probability is then approximately given by

$$\tau = \int_0^{z_D}\alpha_f\sigma_T c(H_0^{-1}/\Omega^{1/2})n_b(1+z)^{1/2}dz$$
$$= 3.3\times10^{-4}(\Omega_b/\Omega^{1/2})h[(1+z_D)^{3/2}-1], \qquad (6)$$

and α_f is the fine-structure constant, n_b is the baryon density, and Ω_b is the fraction of Ω in the form of baryons. Requiring $\tau < 1$ in order to observe the decay photons, we then constrain $1+z_D \leqslant 200\Omega_b^{-2/3}\times\Omega^{1/3}h^{-2/3}$.

The requirement that the decay photons survive until today enables us now to set limits on β, $m_{3/2}$, and Y in regimes where gravitino decays contribute to the γ-

ray background. Now from Eq. (1) and the definition of β we can write $1+z_D=730h^{2/3}(\alpha^{2/3}\beta^2\gamma^2)^{-1}\Omega^{1/3}$. We obtain from the upper limit on $1+z_D$

$$1.9(\Omega_b h^2/\alpha)^{1/3}\gamma^{-1}<\beta, \tag{7a}$$

$$m_{3/2}<0.4\Omega_b^{-1/3}\alpha^{-1/3}\Omega^{1/3}, \tag{7b}$$

and

$$Y<3.3\times 10^{-15}(\Omega_b h^2/\alpha)^{-2/3}\gamma^2. \tag{7c}$$

These constraints are summarized in Fig. 1 where we have plotted the allowable regions in the β-$m_{3/2}$ plane. The limits from $\tau<1$ show the excluded region in the lower right-hand corner for $h=\frac{1}{2}$ and $h=1$. All results plotted in the figure have assumed $\gamma=2$. The line perpendicular to the limits represents a range of values about $\alpha=1/32\pi$. Depending on the value of the photino mass, α may be much smaller (one must then readjust the limits for $\gamma>2$). Of course in order to explain the feature in the observed γ-ray background, we need β to be around 1–2 MeV [cf. Eq. (3)].

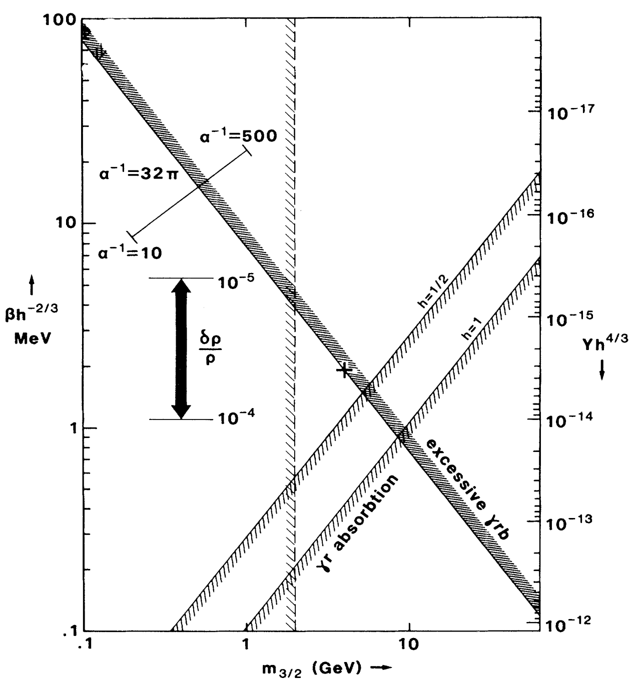

FIG. 1. Constraints in the β-$m_{3/2}$ plane, due to photon absorption (cross-hatched lines); and to threshold for anitproton production from $\tilde{\gamma}$ annihilation (Ref. 20) (dashed-hatched line). The diagonal line corresponds to our predicted γ-ray yield for $\alpha=1/32\pi$ and the scaling to other values of α is indicated. This determines an upper limit on Y to be read on the vertical scale on the right. Also indicated is the expected range for Y from density perturbations $\delta\rho/\rho \sim 10^{-4}$–$10^{-5}$. The cross represents a particular example. We have taken $\Omega_b h^2=0.01$ from big-bang nucleosynthesis (Ref. 25) and $\Omega=1$.

Photinos themselves may also play an important role in cosmology. For example, photinos with a mass[19] in the range[10–12] 1–2 GeV would provide a sufficient energy density so as to make $\Omega \doteq 1$. This would make photinos a prime candidate for the dark matter of the Universe. Annihilations of photinos with mass around 2 GeV in the halos of galaxies have also been suggested as a source for low-energy cosmic-ray antiprotons.[20] We have drawn a dashed vertical line in Fig. 1 to represent the cutoff in the gravitino mass for which this effect is still possible. Our results with regard to the γ-ray background, however, apply to both sides of this line. We can now see that our expected range for Y corresponds rather well to the values of β needed to explain a possible feature in the γ-ray background spectrum[18] near 1 MeV, and falls within the limits due to photon absorption. All γ-ray experiments flown to date in the range 1–20 MeV have confirmed the existence of a bump in the spectrum of diffuse γ's, and no satisfactory explanation exists.[20] The gravitino mass needed for the proposed decay mechanism is fairly unrestricted, although larger masses would require smaller couplings α.

Because of the absence of observed supersymmetric particles such as the SUSY electron, most standard supersymmetric theories would require $m_{3/2}>20$ GeV. This is, however, a tree-level estimate which is subject to gravitational radiative corrections.[21] A mass on the order of a few gigaelectronvolts is not unreasonable. There is also a class of supersymmetric theories called SU(N, 1) no-scale models[22] which are particularly attractive in that the only scale put in by hand is the Planck scale. The weak scale is then determined through radiative corrections. In these models, however, the gravitino mass is not necessarily related to the weak scale or the mass splittings between particles and SUSY particles. Hence $m_{3/2}$ is arbitrary, leaving the horizontal axis in the figure completely unconstrained. The cross on the figure represents a concrete example. For $m_{3/2}=4$ GeV, $\alpha=1/32\pi$, and $m_{\tilde{\gamma}}\lesssim 2$ GeV ($\gamma=2$), decays occur at a red shift $1+z_D \simeq 10^3\Omega^{-1/3}h^{-2/3}$ and a contribution comparable to the observed γ-ray background will appear at $\beta \simeq 2h^{2/3}\Omega^{1/3}$ ($E\simeq 3.5h^{2/3}\Omega^{1/3}$ MeV, for $x=1.76$). The required gravitino abundance lies in the expected range, $Y\simeq 3\times 10^{-15}h^{-4/3}\Omega^{-2/3}$, and we have therefore arrived at a specific mechanism for explaining the γ-ray background spectum near 1 MeV.

Before concluding, we would like to point out that the gravitino is not the unique slowly decaying particle. Nearly all models of supergravity employ what is known as a hidden sector in order to break local supersymmetry. In these models there may be two or more scalar fields which couple only gravitationally to our standard low-energy world. In the event that there are several scalars (which may also interact only gravitationally among themselves) it may be possible to produce photons of various energies.[23]

In summary, a plausible range of parameters (maximum reheating temperature, α, and mass $m_{3/2}$) results in gravitinos decaying at high red shift ($z \sim 1000$) and producing a possibly detectable diffuse gamma-ray flux near 1–10 MeV. A most intriguing possibility is that a possible feature in the diffuse gamma-ray spectrum near 1 MeV could actually be the signature of such decays. Far more exotic interpretations of this feature can be found in the literature.[24] Unfortunately, the spectral region near 1 MeV is notoriously difficult to observe, and the reality of a spectral feature is not unambiguous, nor is the isotropy of the diffuse gamma rays well known in this region.

This research has been supported in part by U.S. Department of Energy Grant No. 84-ER40161 at Berkeley.

[1]R. Kimble, C. S. Bowyer, and P. Jakobsen, Phys. Rev. Lett. **46**, 80 (1981); P. Salati and J. C. Wallet, Phys. Lett. **144B**, 61 (1984); N. Cabibbo, G. R. Farrar, and L. Maiani, Phys. Lett. **105B**, 155 (1981); D. W. Sciama, Mon. Not. Roy. Astron. Soc. **198**, 1 (1982).

[2]M. Davis, M. Lecar, C. Pryor, and E. Witten, Astrophys. J. **250**, 423 (1981); P. Hut and S. D. M. White, Nature, (London) **310**, 637 (1984); M. Turner, G. Steigman, and L. Krauss, Phys. Rev. Lett. **52**, 2090 (1984); G. Gelmini, D. N. Schramm, and J. W. F. Valle, Phys. Lett. B. (to be published); K. A. Olive, D. Seckel, and E. Vishniac; Astrophys. J. (to be published); K. A. Olive, D. N. Schramm, and M. Srednicki, Nucl. Phys. B. (to be published).

[3]S. Weinberg, Phys. Rev. Lett. **48**, 1303 (1982).

[4]D. V. Nanopoulos, K. A. Olive, and M. Srednicki, Phys. Lett. **127B**, 30 (1983).

[5]M. Yu. Khlopov and A. D. Linde, Phys. Lett **138B**, 265 (1984).

[6]J. Ellis, J. Kim, and D. V. Nanopoulos, Phys. Lett. **145B**, 181 (1984).

[7]J. Ellis, D. V. Nanopoulos, and S. Sarkar, Nucl. Phys. **B259**, 175 (1985).

[8]R. Juszkiewicz, J. Silk, and A. Stebbins, to be published.

[9]J. Ellis, A. D. Linde, and D. V. Nanopoulos, Phys. Lett. **118B**, 59 (1982).

[10]H. Goldberg, Phys. Rev. Lett. **50**, 1419 (1983).

[11]L. M. Krauss, Nucl. Phys. **B227**, 556 (1983).

[12]J. Ellis, J. Hagelin, D. V. Nanopoulos, K. A. Olive, and M. Srednicki, Nucl. Phys. **B238**, 453 (1984).

[13]L. E. Ibanez, Phys. Lett. **137B**, 160 (1984); J. Hagelin, G. Kane, and S. Raby, Nucl. Phys. **B241**, 638 (1984).

[14]J. Kim, Phys. Lett. **136B**, 150 (1984); J. Kim, A. Masiero, and D. V. Nanopoulos, Phys. Lett. **139B**, 346 (1984).

[15]J. Ellis, D. V. Nanopoulos, K. A. Olive, and K. Tamvakis, Nucl. Phys. **B221**, 224 (1983); D. V. Nanopoulos, K. A. Olive, M. Srednicki, and K. Tamvakis, Phys. Lett. **123B** 41 (1983); D. V. Nanopoulos, K. A. Olive, M. Srednicki, Phys. Lett. **127B**, 30 (1983); R. Holman, P. Ramond, and G. G. Ross, Phys. Lett. **137B**, 343 (1984).

[16]A. Guth and S. Y. Pi, Phys. Rev. Lett. **49**, 1110 (1982); S. W. Hawking, Phys. Lett. **115B**, 295 (1982); J. M. Bardeen, P. S. Steinhardt, and M. S. Turner, Phys. Rev. D **28**, 679 (1983); A. A. Starobinski, Phys. Lett. **117B**, 175 (1982); J. Ellis, D. V. Nanopoulos, K. A. Olive, and K. Tamvakis, Phys. Lett. **120B**, 331 (1983).

[17]J. Silk, in *Proceedings of the Inner Space/Outer Space Workshop,* edited by E. Kolb, M. S. Turner, D. Lindley, K. A. Olive, and D. Seckel (Univ. of Chicago Press, Chicago, 1985).

[18]V. Schonfelder, F. Graml, and F. P. Penningsfeld, Astrophys. J **240**, 350 (1980).

[19]A photino mass of ~ 100 eV would also yield $\Omega = 1$ and correspond to setting $\gamma = 2$.

[20]J. Silk and M. Srednicki, Phys. Rev. Lett. **53**, 624 (1984).

[21]R. Barbieri and S. Cicotti, Z. Phys. C **17**, 183 (1983); M. Srednicki and S. Theisen, to be published.

[22]E. Cremmer, S. Ferrara, C. Kounas, and D. V. Nanopoulos, Phys. Lett. **133B**, 287 (1983); J. Ellis, C. Kounas, and D. V. Nanopoulos, Nucl. Phys. **B241**, 406 (1984). J. Ellis, K. Enquist, and D. V. Nanopoulos, Phys. Lett. **147B**, 99 (1984).

[23]However, hidden sectors generally present other cosmological difficulties: c.f. C. D. Coughlan, W. Fishler, E. W. Kolb, S. Raby, and G. G. Ross, Phys. Lett. **131B**, 59 (1983).

[24]Attempts to interpret the diffuse gamma-ray feature near 1 MeV have been made by F. W. Stecker, Astrophys. J. **157**, 507 (1969), involving primordial cosmic-ray interactions, and by F. W. Stecker, D. L. Morgan, and J. Bredecamp, Phys. Rev. Lett. **27**, 1469 (1971), and F. W. Stecker, Nature (London) **273**, 493 (1978), who discussed matter-antimatter annihilation in a baryon-symmetric universe. Alternatives to these diffuse production mechanisms involve essentially arbitrary extrapolation to γ-ray energies of the contribution of active galaxy x-ray sources, as discussed in Ref. 18.

[25]J. Yang, M. S. Turner, G. Steigman, D. N. Schramm, and K. A. Olive, Astrophys. J **281**, 413 (1984).