Cosmic-Ray Antiprotons as a Probe of a Photino-Dominated Universe

Joseph Silk
Astronomy Department, University of California, Berkeley, California 94720, and Institute for Theoretical Physics, University of California, Santa Barbara, California 93106

and

Mark Srednicki
Physics Department, University of California, Santa Barbara, California 93106
(Received 8 June 1984)

Observational tests of the hypothesis that the universe is flat and dominated by dark matter in the form of massive photinos include the production of significant fluxes of cosmic rays and gamma rays in our galactic halo. Specification of the cosmological photino density and the masses of scalar quarks and leptons determines the present annihilation rate. The predicted number of low-energy cosmic-ray antiprotons is comparable to the observed flux.

PACS numbers: 94.40.Cn, 11.30.Pb, 14.80.Pb, 98.80.Bp

One of the most important predictions of inflationary cosmology is that the mean density Ω, in units of the critical Einstein–de Sitter value $\rho_{cr} = 3H_0^2/8\pi G$ (where H_0 is Hubble's constant), is unity.[1] Primordial nucleosynthesis constraints[2] limit the baryonic contribution Ω_b to be at most 0.2, and several nonbaryonic weakly interacting particle candidates have been proposed to make up the discrepancy. Viable scenarios for galaxy formation and clustering appear to require that the nonbaryonic matter be cold at epochs when the horizon first contained a mass comparable to that of a galaxy.[3] One of the more plausible particle candidates is the stable massive photino. We show here that there are noteworthy observational tests of the hypothesis that there is a critical cosmological density of gigaelectronvolt photinos.

Our derivation proceeds as follows. Specification of Ω suffices to fix the photino annihilation rate Γ, and specification of the various scalar quark and lepton masses then fixes the photino mass $m_{\tilde{\gamma}}$.[4,5] The low-energy limit of Γ then yields the annihilation rate in our galactic halo, which is assumed to consist of photinos. Accretion of photinos into galactic halos is unavoidable in a photino-dominated universe. We then find that observable annihilation products of $\sim$3-GeV photinos in our halo include γ rays, cosmic-ray positrons, and, most significantly, low-energy cosmic-ray antiprotons. Following Gunn *et al.*[6] and Stecker,[7] Sciama[8] proposed that gamma rays would be produced in observable amounts by photino annihilations. Zel'dovich *et al.*[9] examined cosmic-ray production as a constraint on the mass of a stable, heavy, neutral lepton, but the possibility of antiparticle production at a significant level has not previously been considered.

To commence, we adopt a minimal supersymmetric model in which the lightest supersymmetric particle is a photino (with a negligible admixture of Higgs fermion). The thermally averaged annihilation cross section for $\tilde{\gamma}\tilde{\gamma} \to \bar{f}f$, where f is any quark or lepton, is[4,5]

$$\langle \sigma v \rangle = 8\pi\alpha^2 Q_f^4 r_f \left[\frac{m_f^2}{m_{sf1}^4} + \left(\frac{2T}{m_{\tilde{\gamma}}} \right) \left(\frac{m_{\tilde{\gamma}}^2}{m_{sf2}^4} - \frac{m_f^2}{m_{sf3}^4} \right) \right], \tag{1}$$

where Q_f and m_f are the charge and mass of f, $m_{sf1,2,3}$ are various functions[5] of the masses of the two scalar partners of f, and $r_f = (1 - m_f^2/m_{\tilde{\gamma}}^2)^{1/2}$. To obtain the total annihilation cross section, one should sum over all quarks and leptons with $m_f < m_{\tilde{\gamma}}$. By following standard methods,[10] this cross section can be used to compute a relic photino mass density. Requiring that this equal the critical density (for $H_0 = 50$ km/s$\cdot$Mpc) gives a value for $\langle \sigma v \rangle$ at the freeze-out temperature T_*, when the annihilation rate per particle was equal to the expansion rate.[11] (For values of $\langle \sigma v \rangle$ of interest here, $T_* \approx \frac{1}{20} m_{\tilde{\gamma}}$.)[4,5] This needed value of $\langle \sigma v \rangle_*$ is plotted versus $m_{\tilde{\gamma}}$ as the solid line in Fig. 1, with the subsequent heating of the relic photon background by other species of particles[12] taken into account. We see that we need a total $\langle \sigma v \rangle$ of about 10^{-26} cm^3/s. Also plotted in Fig. 1 for $m_{sf1,2,3} = 50$ GeV is $\langle \sigma v \rangle_*$ as given by Eq. (1) (dashed line). The intersection of the two curves shows that scalar masses of 50 GeV and a photino mass of $\sim$2.8 GeV would yield a critical density of relic photinos. Decreasing (increasing) the scalar masses would move the dashed curve up (down).

Given $m_{\tilde{\gamma}}$ and m_{sf1}, we can compute the annihila-

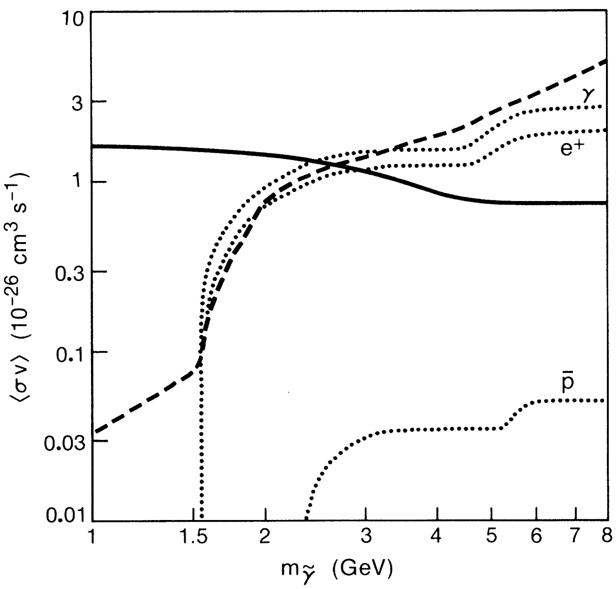

FIG. 1. The photino annihilation cross section necessary to have $\Omega = 1$ with $h = \frac{1}{2}$, as a function of photino mass (solid line); the annihilation cross section at the freeze-out temperature for 50-GeV quark and lepton scalars (dashed line); and the production cross sections for antiprotons, positrons, and gamma rays with 50-GeV scalars (dotted lines).

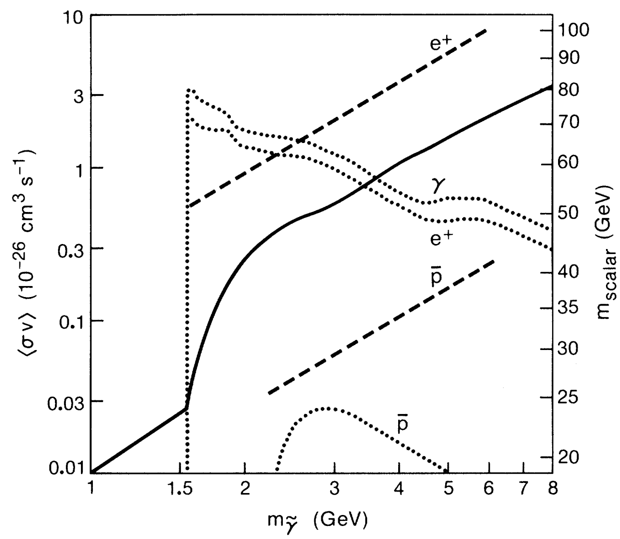

FIG. 2. The scalar quark and lepton masses (taken to be equal) as a function of photino mass needed to have $\Omega = 1$ with $h = \frac{1}{2}$ (solid line); the antiproton, positron, and gamma-ray production cross sections with these scalar masses (dotted lines); and the cross sections necessary to reproduce the observed antiproton and positron fluxes with our halo parameters (dashed lines).

tion rate today. For $m_{\tilde{\gamma}} \lesssim 1.5$ GeV, the annihilation rate is negligible. For 1.8 GeV $\lesssim m_{\tilde{\gamma}} \lesssim 5$ GeV, the primary annihilation products are τ leptons and charmed quarks. For $m_{\tilde{\gamma}} \gtrsim 5$ GeV, the bottom-quark mass, bottom quarks are also produced. These particles will then decay to ordinary particles. The annihilation products of most interest are e^+, $\bar{p}$, and γ. The total number of each type of particle produced per $\tilde{\gamma}\tilde{\gamma}$ annihilation can be estimated from experimental results[13] on $e^+e^- \to \tau^+\tau^-$, $\bar{c}c$, and $\bar{b}b$. We estimate the following yields of stable particles for each of the relevant channels:

$$\tilde{\gamma}\tilde{\gamma} \to \tau^+\tau^- \to 1.5(e^+e^-) + 0(\bar{p}p) + 1\gamma + 5.5\nu,$$
$$\tilde{\gamma}\tilde{\gamma} \to \bar{c}c \to 4(e^+e^-) + 0.2(\bar{p}p) + 7\gamma + 22\nu, \quad (2)$$
$$\tilde{\gamma}\tilde{\gamma} \to \bar{b}b \to 7.5(e^+e^-) + 0.3(\bar{p}p) + 13\gamma + 41\nu.$$

With these numbers, the total production cross sections for e^+, $\bar{p}$, and γ are plotted in Figs. 1 and 2. These results are sensitive to the *relative* masses of scalar quarks and leptons. Lighter scalar quarks and heavier scalar leptons would cause the e^+, $\bar{p}$, and γ lines in Fig. 2 to move up, but not by more than a factor of 3.

The energy distributions of the particles are harder to estimate. We guess that each γ has about twice the energy of each $e^\pm$ or ν ($\pi^0 \to \gamma\gamma$ vs $\pi^\pm \to e^\pm\nu\nu\nu$), and that the average kinetic energy of a p or $\bar{p}$ is about the same as a typical γ energy. Decay products of τ leptons (which account for 70% of all annihilation events below the b quark threshold of ~ 5 GeV and 58% above it, if all scalar masses are equal) have higher average energies than c or b quark decay products; we estimate $\epsilon_\gamma \simeq 0.2 m_{\tilde{\gamma}}$ for γ's from τ decay. Decay products of c and b quarks have a much softer spectrum; we estimate $\epsilon_\gamma \simeq 0.09 m_{\tilde{\gamma}}$ for γ's from c quark decay and $\epsilon_\gamma \simeq 0.05 m_{\tilde{\gamma}}$ for γ's from b quark decay.

Consider first the isotropic gamma-ray background. We integrate the volume emissivity $\epsilon_\gamma n_{\tilde{\gamma}}^2 \langle\sigma v\rangle_\gamma$ (where $n_{\tilde{\gamma}}$ is the average photino number density and $\langle\sigma v\rangle_\gamma$ is the production cross section for a single γ) over the red-shift range $1 \leq 1+z \leq \epsilon_\gamma/\epsilon_\gamma^0$, where $\epsilon_\gamma \simeq f_\gamma m_{\tilde{\gamma}}$, $f_\gamma \simeq 0.2$, and ϵ_γ^0 is the observed, red-shifted gamma-ray energy.

The resulting isotropic γ-ray background flux is[6,7]

$$F_\gamma = \int_0^{\epsilon_\gamma/\epsilon_\gamma^0} \frac{c}{4\pi H_0} dz\,(1+z)^{1/2} \rho_{\rm cr}^2 \langle\sigma v\rangle_\gamma m_{\tilde{\gamma}}^{-2}$$
$$= 6.3 \times 10^{-10} \langle\sigma v\rangle_{26} \Omega_{\tilde{\gamma}}^2 h^3 m_3^{-2} (\epsilon_\gamma/\epsilon_\gamma^0)^{3/2}$$
$$({\rm cm}^2\ {\rm s\ sr})^{-1}, \quad (3)$$

where $\langle \sigma v\rangle_{26} = \langle \sigma v\rangle_\gamma/(10^{-26}\ \mathrm{cm}^3\ \mathrm{s}^{-1})$ and $m_3 = m_{\tilde\gamma}/(3\ \mathrm{GeV})$. This is some four orders of magnitude (with $m_{\tilde\gamma} \approx 3$ GeV) below the observed isotropic γ-ray flux above 100 MeV, which amounts to about[14] $0.9\times10^{-5}\ (\mathrm{cm}^2\ \mathrm{s\ sr})^{-1}$.

However, cold dark matter cannot be prevented from accreting into our galactic halo. The most natural assumption is that our halo is made of photinos. The enhancement in flux from our halo relative to the isotropic background amounts to a factor $\sim(\rho_h/\rho)^2(a/ct_0)$, where ρ_h is the halo density within the uniform-density halo core, of radius a. For a more reliable estimate,[6] the halo density distribution is specified by an analytic fit to a self-consistent dynamical model, defined by measured rotation velocity $v_{200} \equiv v/(200\ \mathrm{km\ s}^{-1})$ and core radius $a_{10} \equiv a/(10\ \mathrm{kpc})$, with ρ_h given by

$$\rho_h(r) = (v^2/4\pi G a^2)(1+r^2/a^2)^{-1}; \qquad (4)$$

hence the photino density in the halo core is

$$n_0 = 0.09 v_{200}^2 a_{10}^{-2} m_3^{-1}\ \mathrm{cm}^{-3}.$$

More sophisticated models[15] yield a similar density for $\rho_h(r)$ in the solar neighborhood. The halo γ-ray flux can now be expressed as

$$F_\gamma^h = \langle \sigma v\rangle_\gamma n_0^2 a \int_0^\infty dx\ (x^2+a^2)^{-2}/4\pi$$
$$= 1.7\times10^{-7}\langle \sigma v\rangle_{26} v_{200}^4 a_{10}^{-3} m_3^{-2}\ (\mathrm{cm}^2\ \mathrm{s\ sr})^{-1}. \qquad (5)$$

This cannot exceed the observed high-latitude γ-ray flux, whence we infer the weak bound

$$\langle \sigma v\rangle_{26} \lesssim 53 m_3^2 a_{10}^3 v_{200}^{-4}. \qquad (6)$$

A photino with mass ~ 3 GeV yields an integral γ-ray flux above 100 MeV that is at least an order of magnitude below that observed.

However, a substantial fraction of the γ-ray energy flux will be in rather energetic gamma rays (from the τ decay channel) with energy $\epsilon_\gamma \approx 0.2 m_{\tilde\gamma}$. This may lead to a more significant prediction if $m_{\tilde\gamma} \gtrsim 3$ GeV, since the observed isotropic gamma-ray spectrum appears to be very steep above 100 MeV. In fact, if the number flux decreases as E^{-2} per unit of energy interval, as the data[14] suggest, over the range 100 MeV $\lesssim \epsilon \lesssim 1$ GeV, we infer that

$$\langle \sigma v\rangle_{26} \lesssim 2.1 a_{10}^3 v_{200}^{-4} \qquad (7)$$

for $1.5 \lesssim m_{\tilde\gamma} \lesssim 5$ GeV.

Cosmic-ray observations provide a more important constraint on halo photinos than do γ-ray observations. Cosmic rays are trapped in the halo by diffusion against magnetic field irregularities, and thus the observable flux is boosted by a considerable factor. The cosmic-ray diffusion time in the galactic disk is known to be of order 10^7 yr, but is likely to be considerably longer in the halo. The diffusion coefficient in most cosmic-ray confinement models[16] lies in the range $(1\text{–}10)\times10^{28}$ $\mathrm{cm}^2\ \mathrm{s}^{-1}$. Writing $K = 10^{29}K_{29}\ \mathrm{cm}^2\ \mathrm{s}^{-1}$, we infer a confinement time scale $t_h \sim (a^2/3)K \sim 10^8 a_{10}^2 K_{29}^{-1}$ yr. Hence the annihilation fluxes of cosmic rays are enhanced over the gamma-ray flux by a factor of order $(\langle\sigma v\rangle/\langle\sigma v\rangle_\gamma)ct_h/a$. We focus on the cosmic-ray antiprotons and positrons to obtain the most sensitive limits. Observations of p and e^+ in the range 0.1–1 GeV are subject to uncertainty because of solar modulation; however, the $\bar p/p$ and e^+/e^- ratios are reasonably well known.

There is some indication that primary cosmic-ray antiprotons may already have been observed below the kinematic threshold for secondary production of about 2 GeV. In the demodulated energy range of approximately 0.6–1.2 GeV, the observed $\bar p$ flux is[17]

$$F_{\bar p} \approx 3\times10^{-6}\ (\mathrm{cm}^2\ \mathrm{s\ sr})^{-1} \quad (0.6 \lesssim E \lesssim 1.2\ \mathrm{GeV}). \qquad (8)$$

A secondary origin cannot be excluded for the antiprotons below 1 GeV, but the models become extremely contrived.[18] The antiproton flux predicted by halo photino annihilation is [with $\langle\sigma v\rangle_{27} = \langle\sigma v\rangle_{\bar p}/(10^{-27}\ \mathrm{cm}^3\ \mathrm{s}^{-1})$]

$$F_{\bar p}^h = 5\times10^{-6}\langle\sigma v\rangle_{27} v_{200}^4 f_{0.1} a_{10}^{-2} K_{29}^{-1} m_3^{-2}\ (\mathrm{cm}^2\ \mathrm{s\ sr})^{-1}, \qquad (9)$$

where the factor $f \equiv 0.1 f_{0.1}$ corrects for the fact that only a fraction f of the annihilation $\bar p$ produced with a mean energy of about $0.09 m_{\tilde\gamma}$ lies in the observed, demodulated energy range. Clearly, a reasonable choice of parameters, and $K_{29} \lesssim 1$, suffices to produce the observed low-energy $\bar p$ flux below 1 GeV.

The predicted positron flux is given by an expression similar to (9) but with a production cross section approximately 30 times higher than $\langle\sigma v\rangle_{\bar p}$ (Fig. 1). Now the observed flux is[18] $e^+/(e^+ + e^-) \approx 0.2$ over $0.1 \lesssim E \lesssim 1$ GeV, while the absolute electron cosmic-ray density can be inferred from the synchrotron emissivity of the galaxy. The interstellar e^+ flux is approximately[19]

$$F_{e^+} = 1\times10^{-3}\ (\mathrm{cm}^2\ \mathrm{s\ sr})^{-1} \quad (0.2 \lesssim E \lesssim 1\ \mathrm{GeV}). \qquad (10)$$

The predicted flux of positrons amounts to [taking $\langle \sigma v \rangle_{26} = \langle \sigma v \rangle_{e^+} / (10^{-26}\,\mathrm{cm}^3\,\mathrm{s}^{-1})$]

$$F^h_{e^+} = 5\times 10^{-4} \langle \sigma v \rangle_{26} v^4_{200} a_{10}^{-2} K_{29}^{-1} m_3^{-2} \quad (\mathrm{cm}^2\,\mathrm{s\,sr})^{-1}, \qquad (11)$$

and is comparable to that observed.

In summary, we have found that reasonable values of the masses of scalar leptons and quarks can lead to a critical cosmological density of photinos and to predictions of observable fluxes of cosmic-ray antiprotons and positrons, and high-energy gamma rays. For scalar masses of ~ 50 GeV and a photino mass of ~ 3 GeV, plausible assumptions about the parameters of our dark halo yield an antiproton flux comparable to that observed below the threshold for secondary production in standard cosmic-ray propagation models.

We would like to thank R. M. Barnett and M. Witherell for discussions of charm particle production, and B. Achterberg for discussions of cosmic-ray diffusion. This work was supported in part by the Department of Energy under Contract No. AT0 3-82ER 40069, in part by the National Science Foundation through Grant No. PHY77-27084, supplemented by funds from NASA, and in part by the National Science Foundation through Grant No. PHY83-13324.

Note added.—That massive photinos make a good dark matter candidate was pointed out by S. Weinberg [Phys. Rev. Lett. **50**, 387 (1983)].

[1]A. Guth, Phys. Rev. D **23**, 347 (1981).

[2]J. Yang, M. Turner, G. Steigman, D. Schramm, and K. Olive, Astrophys. J. (to be published).

[3]M. Davis, C. Frenk, and S. White, to be published.

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

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

[6]J. Gunn, B. Lee, I. Lerche, D. Schramm, and G. Steigman, Astrophys. J. **233**, 1015 (1978).

[7]F. W. Stecker, Astrophys. J. **233**, 1032 (1978).

[8]D. Sciama, to be published.

[9]Ya. B. Zel'dovich, A. A. Klypin, M. Yu. Khlopov, and V. M. Chechetkin, Yad. Fiz. **31**, 1286 (1980) [Sov. J. Nucl. Phys. **31**, 664 (1980)].

[10]G. Steigman, Annu. Rev. Nucl. Part. Sci. **29**, 313 (1979).

[11]B. Lee and S. Weinberg, Phys. Rev. Lett. **39**, 165 (1977).

[12]K. Olive, D. Schramm, and G. Steigman, Nucl. Phys. **B180**, 497 (1981).

[13]M. Roos *et al.* (Particle Data Group), Phys. Lett. **111B**, 1 (1982); G. S. Abrams *et al.*, Phys. Rev. Lett. **44**, 10 (1980); J. L. Siegrist *et al.*, Phys. Rev. D **26**, 969 (1982); M. W. Coles *et al.*, Phys. Rev. D **26**, 2190 (1982); M. S. Alam *et al.*, Phys. Rev. Lett. **51**, 1143 (1983), and references therein.

[14]B. C. Fichtel, G. Simpson, and D. Thompson, Astrophys. J. **222**, 833 (1978).

[15]J. Bahcall and R. Soneira, Astrophys. J. Suppl. **44**, 73 (1980).

[16]V. Ginzburg and V. Ptuskin, Rev. Mod. Phys. **48**, 161 (1976).

[17]A. Buffington, S. Schindler, and C. Pennypacker, Astrophys. J. **248**, 1179 (1981).

[18]For example, see R. Protheroe, Astrophys. J. **251**, 387 (1982); P. Kiraly, J. Szabelski, J. Wdowczyk, and A. W. Wolfendale, Nature **293**, 120 (1981); F. Stecker, R. Prothero, and D. Kazanas, Astrophys. Space Sci. **96**, 171 (1983); L. C. Tan and L. K. Ng, Astrophys. J. **269**, 751 (1983).

[19]For example, see R. Protheroe, Astrophys. J. **254**, 391 (1982).