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Distinctive positron feature from particle dark-matter annihilations in the galactic halo

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If the dark matter in our galactic halo consists of weakly interacting massive particles (WIMP's) heavier than the $W^{\pm}$ boson which have a significant annihilation branch into $W^{\pm}$ and Z^0 pairs, e.g., a Higgsino-like neutralino, a very distinctive feature in the cosmic-ray positron spectrum arises from W^+ and Z^0 decays. Because of inherent astrophysical uncertainties such a signal is by no means guaranteed even if heavy WIMP's *do* comprise the galactic halo. However, the positron signature is virtually a "smoking gun" for particle dark matter in the halo and thus worthy of note.

I. INTRODUCTION

The nature of the ubiquitous dark matter known to exist throughout the Universe is a most urgent issue in both cosmology and particle physics.[1] The stringent nucleosynthesis constraint to the mass density contributed by baryons, $\Omega_B \lesssim 0.12$,[2] and the compelling arguments for the flat Einstein–de Sitter cosmology ($\Omega=1$) based upon structure formation and inflation, provide ample motivation for the hypothesis that the dark matter is composed of relic elementary particles. (Here Ω is the ratio of the total mass density to the critical mass density, and Ω_B is the fraction of critical mass density contributed by baryons.) The neutralino, a linear combination of the supersymmetric partners of the photon, Z^0 boson, and neutral Higgs bosons, is a very promising particle dark-matter candidate:[3,4] In large regions of the parameter space for the minimal supersymmetric extension of the standard model,[5] the relic abundance of neutralinos provides closure density. The case for neutralino dark matter is so compelling that major experimental efforts are underway to design and build ultra-low-background detectors that are sensitive to the small energy that is deposited when a neutralino elastically scatters with ordinary matter. However, the operation of such detectors is still a long way off.[6]

Others have suggested that neutralino dark matter in the halo could be detected by its annihilation products, including antiprotons,[7] γ rays,[8] and positrons.[9] A continuum spectrum of such particles is produced by the annihilations of neutralinos and hadronization of the annihilation products, and the cosmic-ray flux of such particles has been used to constrain the neutralino parameter space. However, it seems unlikely to us that a case for the existence of neutralino dark matter in the halo could ever be made on the basis of continuum annihilation products: The astrophysical uncertainties involving the origin and propagation of the conventional sources of such cosmic rays are too great.

Along similar lines, some have suggested using large underground detectors, such as Kamiokande II, Irvine-Michigan-Brookhaven (IMB), MACRO, and Frejus,[10] to search for high-energy neutrinos produced by the annihilation of neutralinos that have accumulated in the Sun or Earth. The prospects for indirect detection of neutralinos by this means are more promising since a competing background is not a problem. However, the problem here is rate: The predicted fluxes are generally very small.

Several authors have suggested that particle dark matter in the halo might be detected through a narrow γ-ray[11] or positron[12] line. While such a feature is virtually a "smoking gun" for particle dark matter in the halo, the annihilation rates for both these lines are discouragingly small—at least for a conventional neutralino. The direct annihilation of neutralinos into $e^{\pm}$ pairs is strongly suppressed—by a factor of $m_e^2/m_{\tilde{\chi}}^2$—for reasons having to do with chirality,[13] making the positron-line signature a long shot. Similarly, the rate for annihilation into a final state that contains at least one monoenergetic photon is expected to be very small.[14]

For a neutralino that is heavier[15] than the $W^{\pm}$ boson, there is another mechanism for producing high-energy positrons: The annihilation of two neutralinos into a $W^{\pm}$ or Z^0 pair, followed by the direct decay of the W^+ into a positron and an electron neutrino (11% decay branching ratio) or decay of the Z^0 to an $e^{\pm}$ pair (3% decay branching ratio). Although the resulting positron spectrum is not quite as distinctive as that from the line radiation discussed in Ref. 12, it can be much more easily distinguished from conventional cosmic-ray sources (the background) than the continuum positron radiation considered in Ref. 9. Positrons from direct gauge-boson decays have an average energy of half the neutralino mass, with a spectrum that drops off sharply at energies slightly higher or lower. In contrast, the "continuum positrons" are produced in the cascades that result from fermion and Higgs-boson final states; their energies are much lower and there is no sharp drop-off in the spectrum.

A continuous spectrum of positrons will also be produced in the cascade following gauge-boson decays to quarks, muons, and τ leptons, but since their energies are generally lower than those from direct gauge-boson decays, the positron peak at half the neutralino mass and sharp drop-off near the neutralino mass is preserved. In fact, we find that inclusion of the continuum radiation may result in a second peak at much lower energies (about $\frac{1}{20}$ of the neutralino mass); however, since hadronization and decay of quarks is very complicated, this second peak may be less prominent than our calculations indicate. In any case, we are confident that the continuum radiation does not wash out the distinctive positron feature produced by direct decays.

As we will see, neutralinos that are nearly a pure Higgsino state provide the best candidate for producing an observable feature in the cosmic-ray positron spectrum. Although we will focus on this case, it is also conceivable that a mixed Higgsino-gaugino neutralino, or some other heavy dark-matter candidate that has a significant annihilation branch to $W^{\pm}$- or Z^0-boson pairs, such as a heavy Majorana neutrino, could equally well produce a positron signature. Because of the many inherent astrophysical uncertainties, nonobservation of such a feature in the positron spectrum cannot be used to constrain the properties of particle dark-matter candidates.

Our point is that under reasonable but optimistic assumptions a positron signature can arise and may be observable. Since this signature is so striking and since its discovery would be of such enormous importance, this point seems worthy of note despite the fact that less optimistic assumptions about the uncertainties would make detection difficult.

In the next section we address the positron feature that arises directly from gauge-boson decays ($W^+ \to e^+ + \nu_e$, $Z^0 \to e^+ + e^-$). In Sec. III we consider the continuous spectrum of positrons that result from the hadronization and decays of quarks, muons, and τ leptons that are also produced by gauge-boson decays, and in Sec. IV we summarize our results and add some concluding remarks.

II. POSITRONS FROM W^+ AND Z^0 DECAY

To determine the cosmic-ray positron energy spectrum $d\mathcal{F}_+/dE$ produced by WIMP annihilations into $W^{\pm}$ and Z^0 pairs, we need to know the source distributions $f(\epsilon)$ of positrons from W^+ and Z^0 decays as a function of positron energy ϵ and the Green's function $G(E,\epsilon)$ for cosmic-ray positrons during their confinement and propagation through the interstellar medium. Given $f(\epsilon)$ and $G(E,\epsilon)$, the positron spectrum we observe at energy E is

$$\frac{d\mathcal{F}_+}{dE} = \int G(E,\epsilon) f(\epsilon) d\epsilon \; ; \tag{1}$$

the units of $d\mathcal{F}_+/dE$ are $\text{cm}^{-2}\,\text{sr}^{-1}\,\text{GeV}^{-1}\,\text{sec}^{-1}$.

In Ref. 12 a simple homogeneous model of cosmic-ray propagation is developed which includes energy loss due to Compton scattering off the cosmic microwave background and synchrotron radiation in a magnetic field of strength 3×10^{-6} G. The energy-loss time scale for a positron of energy ϵ is $\tau_{\Delta E} \equiv \epsilon/(d\epsilon/dt) \simeq 0.66$ G yr $/(\epsilon/\text{GeV})$. The Green's function for this model is the steady-state differential-energy flux produced by a δ-function source of positrons with energy ϵ and strength a (in units of $\text{cm}^{-3}\,\text{sec}^{-1}$) and is given by

$$G(E,\epsilon) = 8\times10^{26}\frac{a}{E^2}\exp\left[\frac{90}{\tau_7}\left[\frac{1}{\epsilon}-\frac{1}{E}\right]\right]\theta(\epsilon-E)\ \text{cm}^{-2}\,\text{sec}^{-1}\,\text{sr}^{-1}\,\text{GeV}^{-1}\ , \tag{2}$$

where τ_7 is the containment time for positrons in units of 10^7 yr, ϵ and E are given in GeV, and $\theta(\epsilon-E)$ is the Heaviside function. The source strength of positrons from WIMP annihilations is

$$a = 1.5\times10^{-27}\langle\rho_{0.4}^2\rangle\langle\sigma\beta\rangle_8 m_{\tilde{\chi}}^{-2}\ \text{cm}^{-3}\,\text{sec}^{-1}\ , \tag{3}$$

where $\langle\rho_{0.4}^2\rangle$ is the average of the square of the mass density of neutralinos in the halo (in units of 0.4 GeV, cm^{-3}), and $m_{\tilde{\chi}}$ is to be given in GeV. The quantity $\langle\sigma\beta\rangle_8$ is the cross section for annihilation of neutralinos into $W^{\pm}$ pairs (in units of $10^{-8}\ \text{GeV}^{-2}$) and may be obtained from Ref. 4.

The Green's function $G(E,\epsilon)$ rises to a maximum at an energy ϵ and drops sharply to zero—in this model positrons can only lose energy. The δ-function source is broadened due to energy loss and has a characteristic width $\Delta E/\epsilon = \tau/\tau_{\Delta E} \simeq \epsilon\tau_7/(90\ \text{GeV})$. The width of the broadened line is controlled by the ratio of the energy-loss time $\tau_{\Delta E}$ to the containment time τ: An increase in the magnetic-field strength (which decreases $\tau_{\Delta E}$) or an increase in the containment time will broaden the line and vice versa.

Since little is known with certainty about cosmic-ray propagation at higher energies, it is possible, if not probable, that the containment time τ varies with energy. For example, if the containment time varies inversely with positron energy, $\tau(\epsilon) = \tau(E_0)E_0/\epsilon$, then the Green's function is

$$G(E,\epsilon) = 6\times10^{26}\frac{a}{\epsilon^2}\left[\frac{E}{\epsilon}\right]^{r-2} \times\theta(\epsilon-E)\ \text{cm}^{-2}\,\text{sec}^{-1}\,\text{sr}^{-1}\,\text{GeV}^{-1}\ , \tag{4}$$

where $r = \tau_{\Delta E}(\epsilon)/\tau(\epsilon) = (66\ \text{GeV})/(\tau_7 E_0)$ and is independent of ϵ, and we have set $\tau(20\ \text{GeV}) = 10^7\tau_7$ yr. Because the containment time decreases with energy, the flux in this model decreases relative to background more quickly than the flux in the previous model.

The source distribution $f_D(\epsilon)$ of positrons from direct decays $W^+ \to e^+\nu$ and $Z^0 \to e^{\pm}$ is easily obtained from kinematics. In the galactic halo WIMP's move with velocites of order $10^{-3}c$, and so their annihilations occur nearly at rest. To be concrete, let us consider neutralino annihilation into a $W^{\pm}$ pair. After the two neutralinos annihilate the outgoing $W^{\pm}$ bosons move with a velocity $\beta=(1-1/\gamma^2)^{1/2}$, where the Lorentz factor $\gamma=m_{\tilde{\chi}}/m_W$, $m_{\tilde{\chi}}$ is the neutralino mass, and $m_W \simeq 80$ GeV is the $W^{\pm}$ mass. If in the rest frame of the W^+ the positron is emitted at an angle ϕ from the direction of motion of the W^+, the energy of the positron is $m_{\tilde{\chi}}(1+\beta\cos\phi)$. Provided the W^+ boson is unpolarized (which is the case in most CP-conserving theories), its rest-frame decay will be isotropic, and so the energies of the positrons produced will be uniformly distributed (in energy) from $m_{\tilde{\chi}}(1-\beta)/2$ to $m_{\tilde{\chi}}(1+\beta)/2$. In the case of direct neutralino annihilation into an $e^{\pm}$ pair, the energy of the positron is $m_{\tilde{\chi}}$—a positron line; in this case there is a positron "rectangle" centered at an energy of $m_{\tilde{\chi}}/2$ with a width $m_{\tilde{\chi}}\beta$, and the source distribution is

$$f_D(\epsilon)=\frac{B_{W\to e\nu}}{\beta m_{\tilde{\chi}}}\theta\left[\epsilon-\frac{m_{\tilde{\chi}}(1-\beta)}{2}\right] \times\theta\left[\frac{m_{\tilde{\chi}}(1+\beta)}{2}-\epsilon\right], \tag{5}$$

where $B_{W\to e\nu}\simeq 0.11$ is the branching ratio for W^+ decay to a positron.

For positron energies greater than about 90 GeV, cosmic-ray propagation smears out the source distribution by order unity. Thus the positron spectrum produced by W^+ (or Z^0) decays—which has a natural width of order $\beta m_{\tilde{\chi}}$—is not affected much more than a positron line. Moreover, the sharp feature associated with the maximum positron energy, $E(\max) =(1+\beta)m_{\tilde{\chi}}/2$ remains intact.

It is likely that for many WIMP's whose mass exceeds that of the Z^0 boson ($m_Z \simeq 90$ GeV) the cross section for annihilation to a Z^0 pair is comparable to that for annihilation into a $W^{\pm}$ pair. If this is the case—as it is for the neutralino—then a similar discussion applies to positrons from Z^0 decays. Although the Z^0 decays to $e^{\pm}$ only about 3% of the time, a positron is produced by each Z^0, and so $B_{Z\to e^{\pm}}=0.06$.

The cosmic-ray electron flux has been measured up to energies of about 2 TeV; for energies greater than 10 GeV, the cosmic-ray electron flux is given to within a factor of 2 by[16]

$$\frac{d\mathcal{F}_-}{dE}\simeq 0.07(E/\mathrm{GeV})^{-3.3}\ \mathrm{cm}^{-2}\ \mathrm{sr}^{-1}\ \mathrm{GeV}^{-1}\ \mathrm{sec}^{-1}\ . \tag{6}$$

At energies greater than a few GeV, the positron flux is roughly 3%–5% that of the electron flux, which is consistent with estimates from "conventional sources."[17] The dominant conventional source is believed to be the interaction of primary cosmic-ray protons and helium nuclei with nuclei in the interstellar medium, which produces $\pi^{\pm}$ mesons whose decays ultimately produce positrons. For reference estimates for the differential positron flux divided by the sum of the electron plus positron differential fluxes are consistent with[18]

$$\frac{d\mathcal{F}_+/dE}{d\mathcal{F}_+/dE+d\mathcal{F}_-/dE}=0.02+0.10\left[\frac{E}{\mathrm{GeV}}\right]^{-0.5}, \tag{7}$$

which we will use as our estimate for the conventional source "background" to the WIMP-produced positron signal.

Now that we have an idea of the cosmic-ray positron flux expected from halo WIMP annihilations into $W^{\pm}$ and Z^0 pairs and from conventional sources, let us discuss the theoretical expectations for $\langle\sigma\beta\rangle_8$ and $\langle\rho_{0.4}^2\rangle$, whose values determine the source amplitude a. First, we note that the most promising supersymmetric models are ones in which the neutralino is nearly a pure-Higgsino state and where either the squark masses are significantly greater than the neutralino mass or the top quark is not much heavier than the current lower bound. The reason is simple: In these models the neutralino annihilates almost entirely to gauge bosons.[4]

In models where the neutralino is almost a pure-gaugino state, the annihilation branching ratio into $W^{\pm}$ or Z^0 pairs is very small (although the amount of continuum positron radiation from decays of other annihilation products may be comparable to or larger than that expected from conventional sources; see Ref. 9). Models where the neutralino is a mixed state may or may not produce a distinctive positron signature depending on the branching ratio to gauge-boson final states.

Since a Higgsino-like neutralino offers the most favorable case (and simplifies the analysis), we will focus on Higgsinos (although one should keep in mind that mixed-state neutralinos may also be of interest). We should point out that the regions of parameter space in which the neutralino is almost purely Higgsino are quite large (see Fig. 1) and that the squark mass may very well be large enough so that annihilation to fermions is negligible for Higgsinos; thus the models we are considering may be quite general.

In the case that the neutralino is a pure Higgsino state, the cross section times relative velocity for neutralino-neutralino annihilation into a $W^{\pm}$ pair (as the relative velocity approaches zero) is[4]

$$\langle\sigma\beta\rangle_{WW}=\left[1-\frac{m_W^2}{m_{\tilde{\chi}}^2}\right]\frac{G_F^2 m_W^4 m_{\tilde{\chi}}^2}{\pi(2m_{\tilde{\chi}}^2-m_W^2)^2}, \tag{8}$$

and that for annihilation into a Z^0 pair is

$$\langle\sigma\beta\rangle_{ZZ}=\left[1-\frac{m_Z^2}{m_{\tilde{\chi}}^2}\right]\frac{G_F^2 m_Z^4 m_{\tilde{\chi}}^2}{2\pi(2m_{\tilde{\chi}}^2-m_Z^2)^2}. \tag{9}$$

Here G_F is the Fermi coupling constant which, for heavy neutralino annihilations, should be taken to be about 1.07 times its value as measured in low-energy experiments, because of the running of the weak-coupling constant. Above threshold, the cross section increases rapidly and reaches a maximum of 2.5×10^{-8} GeV^{-2} (1.6×10^{-8} GeV^{-2}) at about 110 GeV (120 GeV) for $W^{\pm}$ (Z^0Z^0), re-

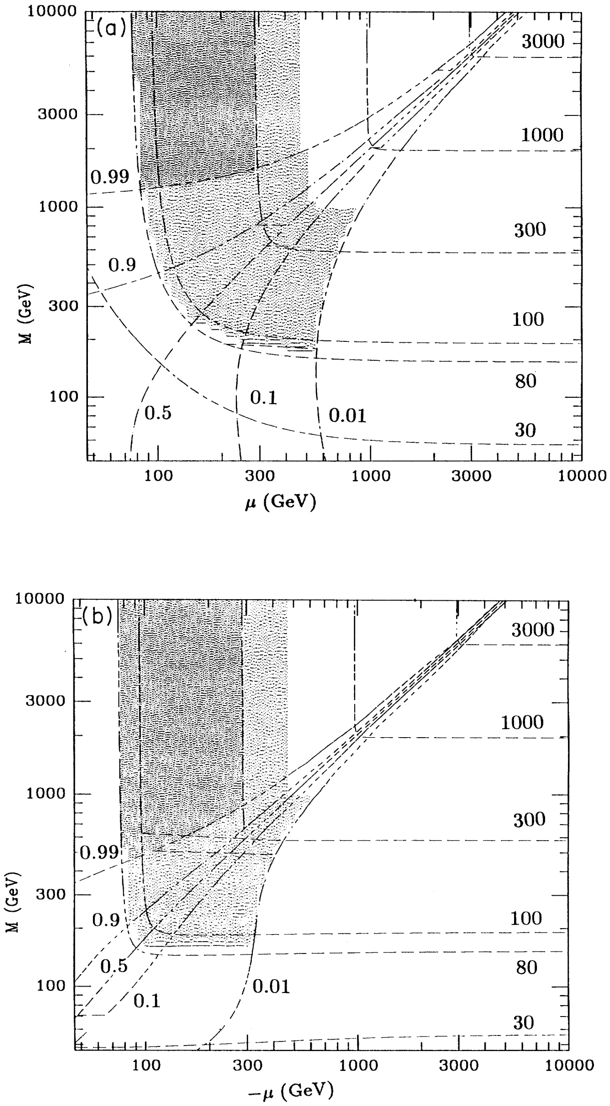

FIG. 1. Higgsino fraction and mass of the neutralino in the M-μ plane for $\tan\beta=2$, for positive and negative μ. In the dark-shaded region the Higgsino fraction is greater than 0.99, and the neutralino mass is between 80 and 300 GeV—the most favorable case for producing a distinctive positron feature. In the light-shaded region the neutralino is either a mixed Higgsino-gaugino state with mass between 80 and 300 GeV or is a Higgsino state with mass between 300 and 500 GeV—in this region a distinctive positron feature is possible.

spectively. For very large neutralino masses $m_{\tilde{\chi}} \gg m_W$, the cross section varies inversely with the neutralino mass squared.

The relic neutralino abundance is inversely proportional to the annihilation cross section evaluated at a relative velocity of about half the speed of light. (Although the annihilation cross sections given above are only strictly valid in the nonrelativistic limit, the qualitative behavior at large relative velocities is similar.) The relic abundance of a heavy neutralino has been computed in Ref. 4 (see Fig. 17). Above the threshold for the $W^{\pm}$ annihilation channel, the relic abundance drops from $\Omega_{\tilde{\chi}}h^2 \sim 1$ to a minimum of $\Omega_{\tilde{\chi}}h^2 \simeq 0.006$ for $m_{\tilde{\chi}} \simeq 110$ GeV; for neutralino masses $m_{\tilde{\chi}} \gg m_W$, the relic abundance is roughly $\Omega_{\tilde{\chi}}h^2 \simeq 2.5\times10^{-3}(m_{\tilde{\chi}}/100\ \text{GeV})^2$ (where the present value of the Hubble constant is $H_0=100h$ km sec^{-1} Mpc^{-1}). Based upon the age of the Universe, Ωh^2 must be less than 1; thus a neutralino with $\Omega_{\tilde{\chi}}h^2 \gtrsim 1$ is cosmologically unacceptable, which excludes a Higgsino-like neutralino more massive than about 2 TeV. We also note that for fixed annihilation cross section, the relic abundance could be smaller (e.g., if there was significant entropy production) or larger (e.g., if the expansion rate at early times was larger) than in the canonical case.[19]

Taking $0.4 \lesssim h \lesssim 1$, we can infer that for a flat universe with $\Omega_B \sim 0.1$ and $\Omega_{\tilde{\chi}} \sim 0.9$, the quantity $\Omega_{\tilde{\chi}}h^2 \simeq 0.1-0.9$, which (keeping in mind the comment above) would fix the annihilation cross section to within a factor of about 10. However, it could be that nature is indeed supersymmetric and $\Omega_{\tilde{\chi}}$ is not unity—either because $\Omega \neq 1$ or because some othe relic accounts for the bulk of the mass density. (As Griest[3] has emphasized, if low-energy supersymmetry is realized in nature, the neutralino abundance is very likely to be significant—greater than a percent or so of critical density.) Since relic neutralinos behave like cold dark matter, they will find their way into the halos of galaxies, whether or not they contribute the critical density. The amount of material known to exist in the halos of spiral galaxies could contribute as little as $\Omega_{\text{halo}} \simeq 0.03$; thus, even in the least favorable case, $\Omega_{\tilde{\chi}}h^2 \sim 0.006$, Higgsino-like neutralinos could comprise galactic halos. Since their relic abundance can always be sufficient to account for the halo dark matter, we shall assume that Higgsinos comprise the galactic halo.

Next, consider the spatial average of the Higgsino mass density squared, $\langle \rho_{0.4}^2 \rangle$. Rotational velocities in our own Galaxy constrain the halo mass interior to our position and allow us to determine the local halo density $\rho_{\text{halo}} \simeq 0.4$ GeV cm^{-3}, with an uncertainty of about a factor of 2. If the halo density interior to our position were constant (e.g., if the core radius of the halo is comparable to or greater than our distance from the galactic center), the uncertainty in the quantity $\langle \rho_{0.4}^2 \rangle$ would be a factor of 4. However, if the halo density increases rapidly toward the center of the Galaxy (e.g., if there is a bulge population of WIMP's or if the halo core radius is small),[20] then there can be a large enhancement in the value of $\langle \rho_{0.4}^2 \rangle$.[21] In particular, suppose the halo density is of the form

$$\rho_{\text{halo}}(r)=\rho_{\text{local}}(R^2+r_{\text{core}}^2)/(r^2+r_{\text{core}}^2)\ , \tag{10}$$

where $\rho_{\text{local}} \simeq 0.4$ GeV cm^{-3} is the local halo density, $R \simeq 8-10$ kpc is the distance from the solar system to the galactic center, and r_{core} is the core radius. If $R/r_{\text{core}} \gg 1$, then $\langle \rho_{0.4}^2 \rangle \simeq \pi R/12 r_{\text{core}} \gg 1$, which represents a significant enhancement.[21]

Finally, as discussed in more detail in Ref. 21, there are other astrophysical uncertainties, including the confinement time of positrons and the number of positrons produced by halo-WIMP annihilations that find their way into the cosmic-ray positron confinement volume, that could increase (or decrease) our estimate of the positron flux associated with halo-WIMP annihilations. It is probably fair to say that because of the various irreducible uncertainties our estimates have a factor of 100—perhaps even a factor of 1000—uncertainty.

In computing the positron flux from neutralino annihilations, we have computed $\langle\sigma\beta\rangle_8$ from Eqs. (8) and (9) and included positrons from both W^+ and Z^0 decays. To make things interesting we have increased the source amplitude a by a factor of 10 over the canonical value in Eqs. (2) and (4) (which is equivalent to setting $\langle\rho_{0.4}^2\rangle=10$). In Fig. 2 we show the cosmic-ray positron fraction for neutralino masses $m_{\tilde\chi}=81$, 90, 100, 120, 300, and 500 GeV, including a "background" from conventional sources; cf. Eq. (7). We show our results for cosmic-ray propagation models with an energy-independent confinement time (solid curves; $\tau_7=1$) and with an energy-dependent confinement time [broken curves; $\tau(E_0=20\text{ GeV})=10^7$ yr].

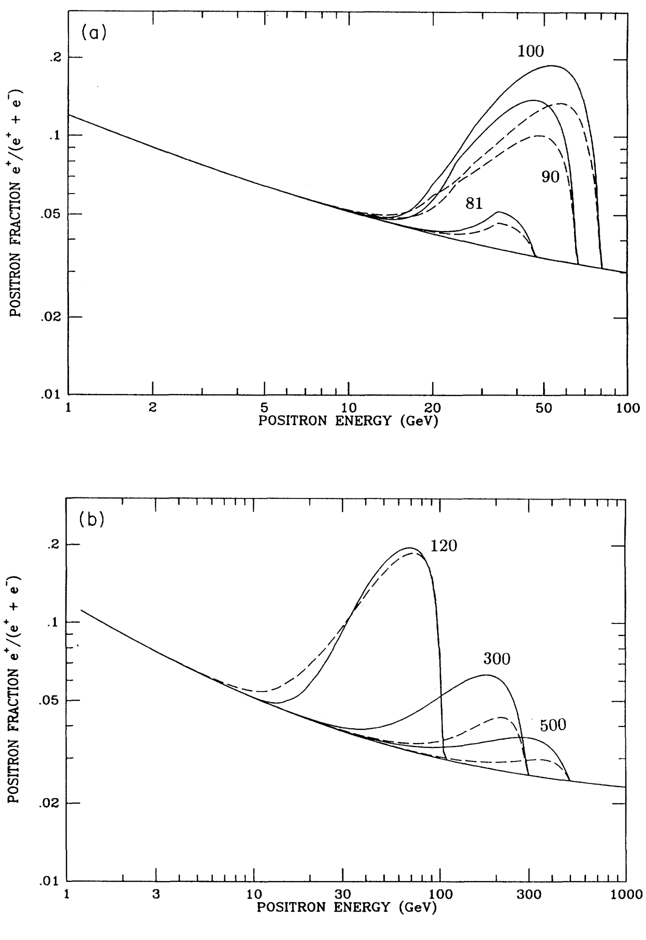

FIG. 2. Differential positron flux divided by the sum of the differential electron and positron fluxes as a function of energy for cosmic-ray propagation models where the positron confinement time is assumed to be energy independent (solid curves and $\tau_7=1$) and where the positron confinement time is assumed to be energy dependent [broken curves and $\tau(E_0=20\text{ GeV})=10^7$ yr]. In (a) the neutralino masses are $m_{\tilde\chi}=81$, 90, and 100 GeV, and in (b) $m_{\tilde\chi}=120$, 300, and 500 GeV. In both we have included a "positron" background from conventional sources given by Eq. (7) and boosted the source amplitude over the canonical value by a factor 10.

As seen in Fig. 2 the feature in the positron fraction for a Higgsino of mass 80 GeV to nearly 500 GeV can be very distinctive and qualitatively similar to that of a positron line. The strength of the feature grows rapidly as the neutralino mass increases above the $W^\pm$ annihilation threshold and has its maximum strength for a neutralino mass of about 120 GeV where the annihilation cross section is maximum. The sharp drop-off in the positron flux at energies near the neutralino mass is particularly distinctive and could be easily distinguished from a background of "smooth" conventional sources, providing a "smoking-gun" signal for particle dark matter in the galactic halo. This should be contrasted with the rather smooth enhancement in the positron flux expected from cascades of other annihilation products.[9]

III. CONTINUUM POSITRON RADIATION

In addition to the positrons produced by the direct decays of gauge bosons (typical energy of about $m_{\tilde\chi}/2$), there will be a continuous spectrum of positrons[9] of lower energy. These positrons are produced as secondary decay products (e.g., $W^+\to\tau^+\to e^+$, $W^+\to b\to e^+$, etc.) and from the decays of pions that are produced in hadronic decays of the $W^\pm$ and Z^0 bosons. The number of pion-produced positrons outnumbers those from direct gauge-boson decays by a factor of about 10, and so one might worry that these positrons could wash out the prominent feature discussed in the previous section. As we shall see, because their energies are much smaller, they do not.

The calculation of the continuum cosmic-ray positron flux is more difficult and more uncertain since the distribution of positrons from the cascade following hadronization and decay of the gauge-boson decay products must be modeled and approximations are made. We ask the reader to keep this in mind as we try to estimate the continuum positron radiation.

Positrons from secondary decays result from $W^\pm$ decays to muons (which always decay to positrons), to τ leptons (which decay to positrons 18% of the time), and to b and c quarks (which decay to positrons 13% of the time). The typical energy of these positrons is about $m_{\tilde\chi}/6$. As noted by Rudaz and Stecker,[9] the distribution of positron energies ϵ from the decay of a μ^+, τ^+, or b or c quark with energy E_f may be approximated by a step function $\theta(kE_f-\epsilon)/(kE_f)$ where $k=0.7$. Integration over the distribution of quark and lepton energies, which is flat from $m_{\tilde\chi}(1-\beta)/2$ to $m_{\tilde\chi}(1+\beta)/2$, like that for direct positrons, gives the source distribution from secondary-decay positrons:

$$f_S(\epsilon)=[B_{W\to\mu}+0.18B_{W\to\tau}+0.13(B_{W\to c}+B_{W\to b})]g(\epsilon)\ , \tag{11}$$

where

$$g(\epsilon)=\begin{cases}\ln[(1+\beta)/(1-\beta)]/km_{\tilde\chi}\beta & \text{for } \epsilon\le km_{\tilde\chi}(1-\beta)/2\ ,\\ \ln[km_{\tilde\chi}(1+\beta)/2\epsilon]/km_{\tilde\chi}\beta & \text{for } km_{\tilde\chi}(1-\beta)/2\le\epsilon\le km_{\tilde\chi}(1+\beta)/2\ ,\\ 0 \text{ for } \epsilon>km_{\tilde\chi}(1+\beta)/2. \end{cases} \tag{12}$$

The quantities B_i are branching ratios for W decay to the various fermion channels. In addition, there are positrons from the tertiary decays of gauge boson (e.g., $W^+\to\tau^+\to\mu^+\to e^+$), but these positrons have lower energies. Since the positron distribution at low energies is dominated by pion decays, we will not consider tertiary decays.

The hadronization of quarks from gauge-boson decays results in a shower of charged pions, which eventually decay to positrons ($\pi^+\to\mu^+\to e^+$). By integrating Rudaz and Stecker's expression for the energy spectrum of pions produced by quarks of energy E_f over the quark energy distribution [which is flat for $m_{\tilde\chi}(1-\beta)/2$ to $m_{\tilde\chi}(1+\beta)/2$] and taking the positron energy to be $\frac{1}{4}$ of the pion energy, we obtain the source distribution of pion-produced positrons:

$$f_\pi(\epsilon)=\frac{B_{W\to\text{hadrons}}}{m_{\tilde\chi}\beta}\int_{m_{\tilde\chi}(1-\beta)/2}^{m_{\tilde\chi}(1+\beta)/2}\left[93\exp\left[-\frac{68\epsilon}{E_f}\right]+56\exp\left[-\frac{27.6\epsilon}{E_f}\right]\right]dE_f\ , \tag{13}$$

where $B_{W\to\text{hadrons}}\simeq\frac{2}{3}$ is the hadronic branching ratio for W decay. In deriving Eq. (13) we multiplied Rudaz and Stecker's distribution function by 2 since a pair of quarks can come from the decay of either gauge boson.

The continuum positron spectrum is then obtained by the convolution of the sources $f_S(\epsilon)$ and $f_\pi(\epsilon)$ with the Green's function [Eq. (2) or (4)]. The complete positron spectrum from decays of gauge bosons produced by halo annihilations of a Higgsino of mass 120 GeV is shown in Fig. 3. Note that the peak associated with the direct decays of W^+ and Z^0 bosons remains quite prominent and is not washed out by the continuum positron radiation. Moreover, there appears to be a second, less prominent peak centered at an energy of about $m_{\tilde\chi}/20$, which could provide a signature for heavy dark matter in the halo that annihilates primarily into gauge bosons. We caution the reader that hadronization and decay of quarks is quite complicated, and it could well be that low-energy peak is much less pronounced than our calculations suggest.

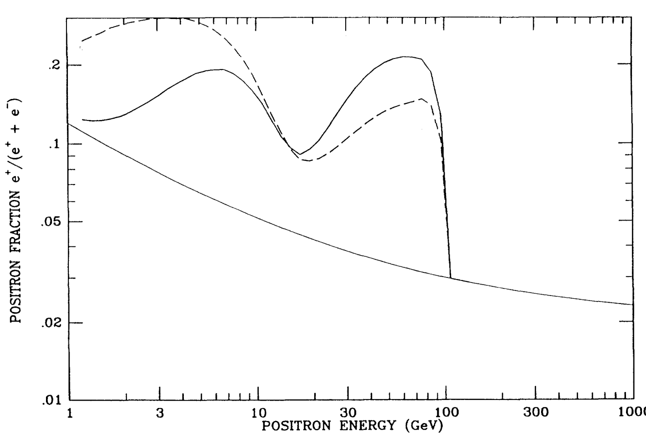

FIG. 3. Differential positron flux divided by the sum of the differential electron and positron fluxes as a function of energy for a neutralino of mass 120 GeV, for models of cosmic-ray propagation where the positron confinement time is constant (solid curve) and where it decreases with energy (broken curve). In addition to the positrons produced by the direct decays of the gauge bosons, we have included the "continuum positron radiation" resulting from the other decay modes of the gauge bosons.

IV. CONCLUDING REMARKS

Since the composition of the ubiquitous dark matter in the Universe is of such great importance to both particle physics and cosmology, any and all avenues that can lead to the *discovery* of its constituents must be pursued. Here we have emphasized the distinctive feature in the cosmic-ray positron spectrum that arises from halo WIMP annihilations into $W^\pm$ and Z^0 pairs followed by W^+ or Z^0 decay into an energetic positron of energy around half the WIMP mass.

We have shown that with somewhat optimistic assumptions regarding the inherent astrophysical and particle-physics uncertainties a very distinctive feature arises in the positron spectrum. We reiterate that even if WIMP's *do* make up the galactic halo, because of the same uncertainties, there is no guarantee that a positron signal would be observable—and therefore it is not possible to use nonobservation of such a signal to rule out dark-matter candidates.

We have also checked to make sure that additional, lower-energy positrons produced by other gauge-boson decays do not wash out this feature; in fact, it appears that they lead to another feature at an energy of about $m_{\tilde\chi}/20$. However, we are quick to remind the reader of the uncertainties and approximations made in calculating the flux due to the continuum positron radiation.

While we have restricted our quantitative analysis to a Higgsino-like neutralino, we stress the generality of our results: Any WIMP heavier than the mass of the $W^\pm$ boson which has a significant annihilation branch into $W^\pm$ or Z^0 pairs could produce such a feature in the positron flux. One such possibility is a heavy Majorana neutrino.

In fact, it is also possible that a similar signal could arise from neutralinos lighter than the $W^{\pm}$: In the minimal supersymmetric extension of the standard model,[5] there is always a neutral Higgs boson (denoted as H_2^0) which is lighter than the Z^0. In some models where the neutralino mass is less than m_W but greater than $(m_Z+m_{H_2^0})/2$, the neutralino can annihilate predominantly into $Z^0H_2^0$ (see the Appendix of Ref. 22). While the Higgs-boson decays produce few positrons (because the relevant coupling is proportional to the electron mass squared), the Z^0 decay will produce many, and the analysis here applies.

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