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Evidence for “sterile neutrino” dark matter?

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ABSTRACT

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I show that it may be possible to explain the present evidence for a gamma-ray emission from the galactic halo as due to halo WIMP annihilations. Not only the intensity and spatial pattern of the halo emission can be matched but also the relic density of the candidate WIMP can be in the cosmologically interesting domain. After a model-independent analysis to learn about the properties of a suitable candidate, I present a working model: a sterile neutrino in a model with an extended Higgs sector. Two examples indicate the existence of an interesting region in the model parameter space where present observational and experimental constraints are satisfied and the gamma-ray emission is reproduced.

1. Introduction

A sophisticated analysis of EGRET data (Dixon et al. 1998) has found evidence for gamma-ray emission from the galactic halo. Filtering the data with a wavelet expansion, Dixon et al. have subtracted an isotropic extra-galactic component and expected contributions from cosmic ray interactions with the interstellar gas and from inverse Compton of ambient photons by cosmic ray electrons, and they have produced a map of the intensity distribution of the residual gamma-ray emission. Besides a few “point” sources, they find an excess in the central region extending somewhat North of the galactic plane, and a weaker emission from regions in the galactic halo. They mention an astrophysical interpretation for this halo emission: inverse Compton by cosmic ray electrons distributed on larger scales than those commonly discussed and with anomalously hard energy spectrum.

I find it intriguing that the angular distribution of the halo emission resembles that expected from pair annihilation of dark matter WIMPs in the galactic halo (Gunn et al. 1978, Turner 1986), and moreover, that the gamma-ray intensity is similar to that expected from annihilations of a thermal relic with present mass density 0.1–0.2 of the critical density (Gondolo 1998a). Namely, the emission at $b \geq 20^\circ$ is approximately constant at a given angular distance from the galactic center – except for a region around $(b, l) = (60^\circ, 45^\circ)$, correlated to the position of the Moon, and a region around $(b, l) = (190^\circ, -30^\circ)$, where there is a local cloud (b and l are the galactic latitude and longitude, respectively). Dixon et al. (1998) argue against the possibility of WIMP annihilations on the base that direct annihilation of neutralinos into photons would give too low a gamma-ray signal. However most of the photons from WIMP annihilations are usually not produced directly but come from the decay of neutral pions generated in the particle cascades following annihilation.

Presently preferred dark matter candidates tend to give a gamma emission which is too low even in the continuum.^a In a search for a suitable candidate it is worth examining the impact of various constraints on the properties of the candidate. I will later introduce a working model, so to make the discussion more concrete.

2. A candidate WIMP can be a thermal relic.

It is useful to use the WIMP mass m_χ and its annihilation cross section (times relative velocity at $v = 0$) σv as parameters in the discussion. The requirement to approximately match the Dixon et al. maps selects a band in the $\sigma v - m_\chi$ plane (band marked “halo γ 's” in fig. 1). Another band is selected by the requirement that the WIMP is a thermal relic from the early universe, with a relic abundance in the range $0.025 < \Omega h^2 < 1$ (band marked “ $0.025 < \Omega h^2 < 1$ ” in fig. 1). The intersection of the two bands defines the interesting region.

The gamma-ray band is obtained as follows. The gamma-ray intensity from WIMP annihilations in the galactic halo from direction with galactic longitude b and galactic latitude l is given by

$$\phi_\gamma(b, l, >E) = n_\gamma(>E) \frac{\sigma v}{4\pi m_\chi^2} \int \rho_\chi^2 dl. \quad (1)$$

$\phi_\gamma(b, l, >E)$ is in photons/(cm² s sr), $n_\gamma(>E)$ is the number of photons with energy above E generated per WIMP annihilation, σv is the WIMP annihilation cross section times relative velocity, m_χ is the WIMP mass, and ρ_χ is the WIMP mass density in the halo. The integral in eq. (1) is along the line of sight in direction (b, l) . This integral gives the angular dependence of the gamma-ray flux and depends on details of the dark halo model, which are not well known. For a canonical halo model, $\rho(r) = \rho_{\text{loc}}(r_c^2 + R^2)/(r_c^2 + r^2)$, where ρ_{loc} is the WIMP mass density in the solar neighborhood, r_c is the halo core radius, R is the distance of the Sun from the galactic center, and r is the galactocentric distance. In this case, the integral depends only on the angle ψ between the direction of observation and the galactic center,

$$\int \rho_\chi^2 dl = \rho_{\text{loc}}^2 R \frac{x}{2(x - c^2)^{3/2}} \left[\frac{\pi}{2} + \arctan \frac{c}{\sqrt{x - c^2}} + \frac{c\sqrt{x - c^2}}{x} \right]. \quad (2)$$

Here $x = 1 + (r_c/R)^2$ and $c = \cos \psi = \cos b \cos l$. From dynamical studies one finds $\rho_{\text{loc}} = 0.3\text{--}0.5$ GeV/cm³, $R = 8\text{--}8.5$ kpc, and $r_c = 2\text{--}8$ kpc. It is interesting to notice that the gamma-ray intensities at $\psi = 40^\circ$ and $\psi = 60^\circ$ are approximately in the ratio 2:1 as on the Dixon et al. maps. Assuming WIMP annihilation into quark-antiquark and lepton-antilepton pairs, fixing $n_\gamma(>1\text{GeV})$ with the Lund Monte-Carlo, varying

^aAttempts to increase the flux by clumpiness in the halo (Berström, Edsjö and Ullio 1998) tend to produce either too many antiprotons or too few photons below 1 GeV, in contrast to the Dixon et al. maps (see section 4).

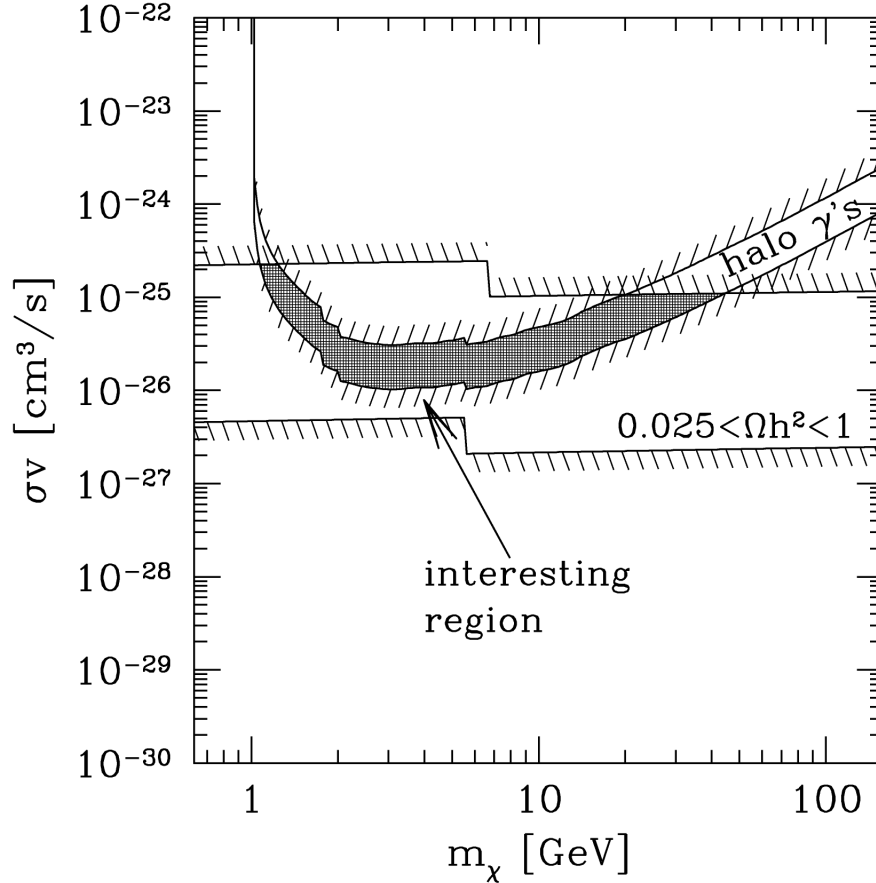


Fig. 1. Region in the annihilation cross section versus mass plane where the gamma-ray halo may be explained by annihilations of a thermal dark matter relic.

the halo parameters in the range given above, and matching the observed intensity to eq. (1) to within 20%, I obtain the required WIMP annihilation cross section as a function of the WIMP mass. This is the band marked “halo γ 's” in fig. 1.

The second band in fig. 1 comes from the requirement that the WIMP relic density Ωh^2 be in the cosmologically interesting range 0.025–1. (h is the Hubble constant in units of 100km/s/Mpc.) The WIMP relic density is related to the WIMP annihilation cross section through the approximate relation (Kolb & Turner 1990, Gondolo & Gelmini 1991)

$$\sigma v = \frac{2.0 \times 10^{-27} \text{cm}^3/\text{s}}{g g_*^{1/2} x_f \Omega h^2}, \quad (3)$$

where I have assumed that the annihilation cross section is dominated by the $v = 0$ term both in the galactic halo and at freeze-out (s-wave annihilation). The freeze-out temperature $x_f m$ can be obtained solving $x_f^{-1} + \frac{1}{2} \ln(g_*/x_f) = 80.4 + \ln(g m_\chi \sigma v)$, with m_χ in GeV and σv in cm^3/s . For a Majorana WIMP $g = 2$, for a Dirac WIMP $g = 4$. g_* is the effective number of relativistic degrees of freedom at freeze-out: $g_* \simeq 81$

before and $g_* \simeq 16$ after the QCD quark-hadron phase transition. Letting Ωh^2 vary in the above range gives the relic density band in fig. 1.

The two bands intersect for WIMP masses between 1.2 and 50 GeV. For example, $\Omega = 0.2$ and $H = 60$ km/s/Mpc give the required cross section of 3×10^{-26} cm³/s at $m_\chi = 5$ GeV. (This case was presented in Gondolo 1998a.)

3. Constraints on candidates that couple to the Z boson.

I assume in this section that $\chi\chi$ annihilation and χ -nucleon scattering are dominated by Z boson exchange. In this case, the annihilation cross section reads

$$\sigma v = \frac{G_F^2}{\pi} \sum_f \beta_f \left[m_f^2 (a_\chi^2 a_f^2 + v_\chi^2 v_f^2 - 2v_\chi^2 a_f^2) + 2m_\chi^2 v_\chi^2 (a_f^2 + v_f^2) \right], \quad (4)$$

where $\beta_f = (1 - m_f^2/m_\chi^2)^{1/2}$, $a_f = T_f$ and $v_f = T_f - 2e_f \sin^2 \theta_W$ are the usual axial and vector couplings of fermion f to the Z boson, and a_χ and v_χ are the analogous quantities for the χ -Z coupling.

The χ -nucleon scattering cross section $\sigma_{\chi N}$, which is limited by negative direct dark matter searches, is related to the annihilation cross section by crossing symmetry. The experimental bound on $\sigma_{\chi N}$ depends on the WIMP mass and on the spin-dependent or spin-independent character of the interaction (for a recent compilation of limits see Bernabei et al. 1998). For Z boson exchange, the spin-dependent and spin-independent χ -nucleon cross sections read

$$\sigma_{\chi N}^{\text{SD}} = \frac{6\mu_{\chi N}^2}{\pi} G_F^2 a_\chi^2 \left[(\Delta u - \Delta d)^2 + (\Delta s)^2 \right], \quad (5)$$

and

$$\sigma_{\chi N}^{\text{SI}} = \frac{\mu_{\chi N}^2}{4\pi} G_F^2 v_\chi^2 \left[(1 - 4 \sin^2 \theta_W)^2 + 1 \right]. \quad (6)$$

Here $\mu_{\chi N} = m_\chi m_N / (m_\chi + m_N)$ is the reduced χ -nucleon mass, and Δq is the quark q contribution to the spin of the proton. (From neutron and hyperon decay, $\Delta u - \Delta d = 1.2573 \pm 0.0028$, while Δs is uncertain: $\Delta s = 0$ in the naive quark model, $\Delta s = -0.11 \pm 0.03$ from deep inelastic data, and $\Delta s = -0.15 \pm 0.09$ from elastic $\nu p \rightarrow \nu p$ data.)

The coupling of the χ to the Z boson also gives a contribution to the Z boson decay width, if $m_\chi < m_Z$. Namely,

$$\Gamma(Z \rightarrow \chi\bar{\chi}) = \frac{G_F m_Z}{6\sqrt{2}\pi} \left[a_\chi^2 m_Z^2 \beta_\chi^3 + \beta_\chi v_\chi^2 (m_Z^2 + 2m_\chi^2) \right], \quad (7)$$

with $\beta_\chi = (1 - 4m_\chi^2/m_Z^2)^{1/2}$. The experimental limit is $\Gamma(Z \rightarrow \chi\bar{\chi}) < 5$ MeV (Barnett et al. 1996).

Once a relation between a_χ and v_χ is specified, the experimental bounds on $\sigma_{\chi\mathcal{N}}$ and $\Gamma(Z \rightarrow \chi\bar{\chi})$ translate into a bound on σv . Fig. 2 plots these bounds for two cases: a Dirac particle with $v_\chi = a_\chi$ ($V - A$ interaction), and a Majorana particle with $v_\chi = 0$ (axial interaction). For Dirac particles, the interesting region is not fully excluded by these constraints, but for Majorana particles it is. Hence the impact of these constraints is model dependent.

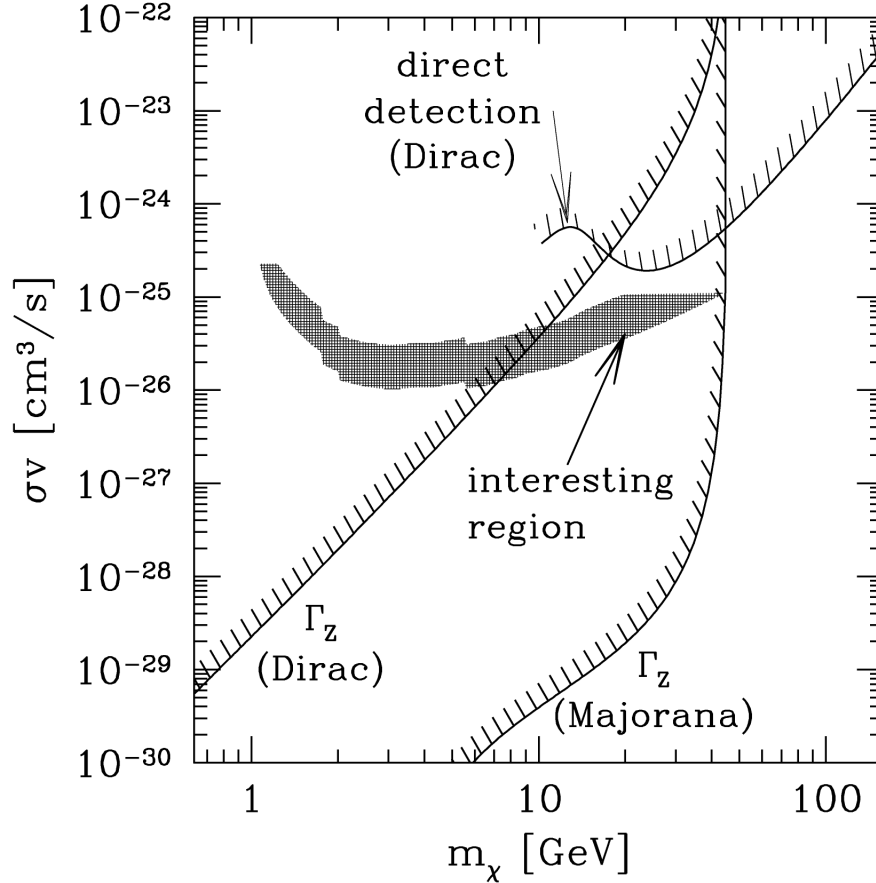


Fig. 2. Constraints on the interesting region for a particle coupled to the Z boson.

4. Constraints from dark matter searches.

A candidate WIMP has to satisfy constraints from negative dark matter searches. In this section I consider indirect detection through production of rare cosmic rays (antiprotons) and through neutrino production in the Sun and the Earth, and direct detection through elastic scattering off nuclei in a laboratory detector.

The most important constraint comes from the measured flux of cosmic ray antiprotons. If the gamma-rays are produced in jets originated by quarks, there is an

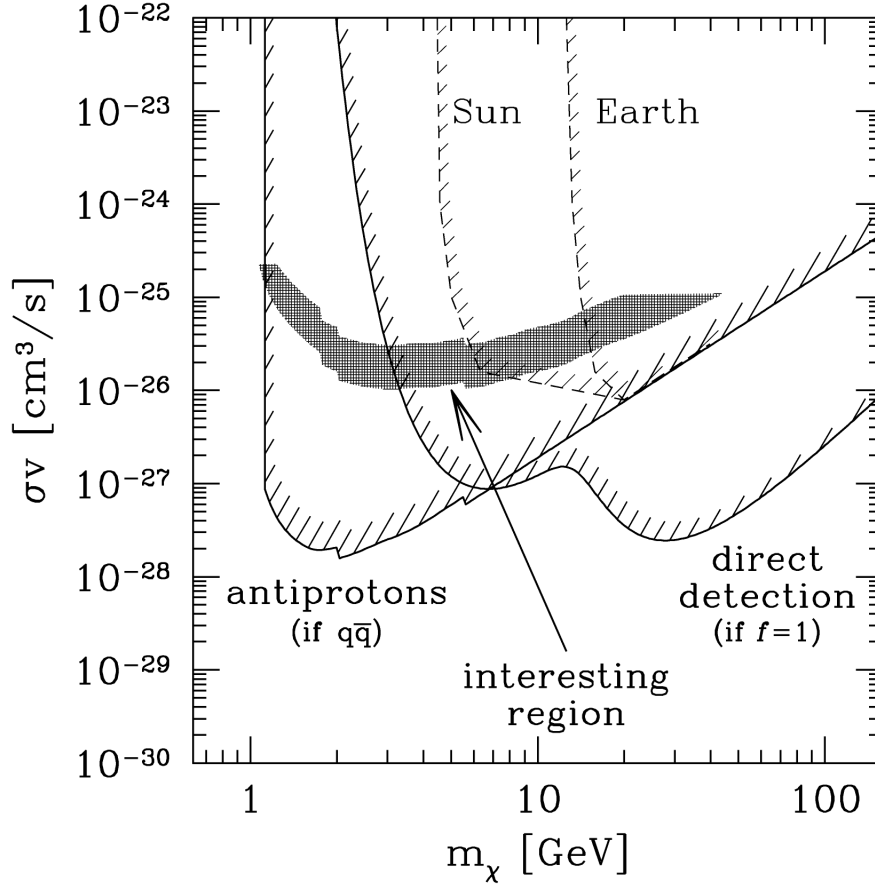


Fig. 3. Constraints on the interesting region from dark matter searches.

associated production of antiprotons. The ratio of antiproton and gamma-ray fluxes is independent of the WIMP annihilation cross section and of the local mass density, and the relative number of antiprotons and photons per annihilation is fixed by the physics of jets. The antiproton flux at a \bar{p} kinetic energy T at the top of the atmosphere is

$$\phi_{\bar{p}}(T) = \frac{dN_{\bar{p}}}{dT} \frac{\sigma v}{4\pi m_{\chi}^2} \rho_{\text{loc}}^2 v_{\bar{p}} t_{\text{cont}} \mu, \quad (8)$$

where $dN_{\bar{p}}/dT$ is the antiproton spectrum per annihilation, t_{cont} is the \bar{p} containment time, which in the diffusion model of Chardonnet et al. (1996) is $t_{\text{cont}} \simeq (1 + p/3\text{GeV})^{-0.65} \times 10^{15}$ s, and $\mu = [T(T + 2m_{\bar{p}})] / [(T + \Delta)(T + \Delta + 2m_{\bar{p}})]$ takes into account solar modulation. Fig. 3 shows the bound obtained by taking $\Delta = 600$ MeV and imposing $\phi_{\bar{p}}(150\text{--}300\text{MeV}) < 3 \times 10^{-6}$ $\bar{p}/\text{cm}^2/\text{s}/\text{sr}/\text{GeV}$ (Moiseev et al. 1997). For WIMPs heavier than ~ 10 GeV, a related constraint on the gamma-ray flux between 300 and 1000 MeV comes from the approximate relation $n_{\gamma}(300\text{--}1000\text{MeV}) \simeq dN_{\bar{p}}(150\text{--}300\text{MeV})/dT \times 10\text{GeV}$. This gives $\phi_{\gamma}(300\text{--}1000\text{MeV}) \lesssim 1.3 \times 10^7$ pho-

tons/cm²/s/sr if photons originate in quark jets with $m_\chi \gtrsim 10$ GeV. Compared with the observed value of $\sim 8 \times 10^7$ photons/cm²/s/sr, it implies too low a gamma-ray flux below 1 GeV. However, these bounds from antiprotons are uncertain because of uncertainties in the antiproton propagation in the galaxy and in the effect of the solar modulation. Moreover, these bounds depend sensitively on the relative branching ratios into the various annihilation channels, and for example can be avoided if the dominant decay channel is leptonic (see next section).

The χ fermions can also accumulate in the Sun and in the Earth, annihilate therein and produce GeV neutrinos. Accumulation is not efficient for WIMPs lighter than the evaporation mass, which is ~ 12 GeV for the Earth and ~ 3 GeV for the Sun. The curves marked “Earth” and “Sun” in fig. 3 show approximate constraints obtained from the experimental upper bounds on the flux of through-going muons in the Baksan detector (Suvorova 1997).

Another limit comes from direct searches through the crossing relation between annihilation and scattering cross sections. This relation can be written as

$$\mu_{\chi\mathcal{N}}^2 \sigma v = f m_\chi^2 c \sigma_{\chi\mathcal{N}}, \quad (9)$$

where c is the speed of light and f accounts for the details of the interactions and of the nuclear structure. The “direct detection” bound in fig. 3 assumes $f = 1$, but actually f can range from 0 to infinity. For example, consider a WIMP that couples to quarks and leptons only through exchange of scalar bosons, and consider annihilation through s-channel exchange and scattering through t-channel exchange. For kinematical reasons, if the scalar is CP-even, the annihilation cross section σv at $v = 0$ vanishes and the scattering cross section at small velocity is finite, so $f = 0$. On the other hand, if the scalar is CP-odd, it is the scattering cross section that vanishes and the annihilation cross section is finite, so f is infinite. Choosing the relative strength of CP-odd and CP-even exchange, the magnitude of the annihilation and scattering cross sections can be tuned. An explicit example is given in the next section.

5. A candidate WIMP can be a sterile neutrino.

The analysis in the previous paragraphs may have been instructive, but the conclusions are so model-dependent that the discussion would be clearer if done in a specific model. Here I show that a particle with the required properties is a sterile right-handed neutrino ν_s in a model with an extended Higgs sector.

The Higgs sector contains two Higgs doublets H_1 and H_2 and a Higgs singlet N . I assume the following Higgs potential

$$V_{\text{Higgs}} = \lambda_1 \left(H_1^\dagger H_1 - v_1^2 \right)^2 + \lambda_2 \left(H_2^\dagger H_2 - v_2^2 \right)^2 + \lambda_3 \left| N^2 - v_N^2 \right|^2 + \lambda_4 |H_1 H_2 - v_1 v_2|^2 + \lambda_5 \left| H_1 H_2 - v_1 v_2 + N^2 - v_N^2 \right|^2 + \lambda_6 \left| H_1^\dagger H_2 \right|^2. \quad (10)$$

To fix the notation, $H_1 H_2 = H_1^0 H_2^0 - H_1^- H_2^+$. Taking all the λ 's real and positive guarantees that the absolute minimum of the Higgs potential is at $\langle H_1^0 \rangle = v_1$, $\langle H_2^0 \rangle = v_2$, and $\langle N \rangle = v_N$, with the vacuum expectation values of all the other fields vanishing.

Through the Yukawa terms

$$\mathcal{L}_{\text{Yukawa}} = f_d Q H_1 d_R + f_u Q H_2 u_R + f_e L H_1 e_R + h \overline{\nu_{sR}^c} \nu_{sR} N, \quad (11)$$

H_1 gives masses to the up-type quarks and the charged leptons, H_2 gives masses to the down-type quarks, and N gives a Majorana mass to the right-handed neutrino. There is no mixing of the right-handed neutrino with ordinary neutrinos, otherwise the new neutrino would have decayed in the early universe and would not be in the galactic halo at present.^b

There are 2 would-be Goldstone bosons – a charged one G^\pm and a neutral one G^0 – and 6 physical Higgs bosons: a charged one H^\pm , three neutral “scalars” S_i ($i = 1, \dots, 3$), and two neutral “pseudoscalars” P_i ($i = 1, 2$).

The physical charged Higgs field $H^\pm = H^{\mp*} \sin \beta + H_2^\pm \cos \beta$ has squared mass

$$m_{H^\pm}^2 = \lambda_6 v^2. \quad (12)$$

As usual, $v = \sqrt{v_1^2 + v_2^2}$ and $\tan \beta = v_2/v_1$.

The squared mass matrix of the two physical pseudoscalar neutral Higgs fields reads

$$\mathcal{M}_P^2 = \begin{pmatrix} (\lambda_4 + \lambda_5)v^2 & 2\lambda_5 v v_N \\ 2\lambda_5 v v_N & 4(\lambda_3 + \lambda_4)v_N^2 \end{pmatrix} \quad (13)$$

in the basis $\sqrt{2}(\sin \beta \text{Im} H_1^0 + \cos \beta \text{Im} H_2^0, \text{Im} N)$. The orthogonal combination of fields is the would-be Goldstone boson G^0 . Let U^P be the unitary matrix that diagonalizes \mathcal{M}_P^2 , namely $P_i = U_{i1}^P A + U_{i2}^P N_I$ where $i = 1, 2$.

The scalar Higgs boson mass matrix is

$$\mathcal{M}_S^2 = \begin{pmatrix} 4\lambda_1 v_1^2 + (\lambda_4 + \lambda_5)v_2^2 & (\lambda_4 + \lambda_5)v_1 v_2 & 2\lambda_5 v_2 v_N \\ (\lambda_4 + \lambda_5)v_1 v_2 & 4\lambda_2 v_2^2 + (\lambda_4 + \lambda_5)v_1^2 & 2\lambda_5 v_1 v_N \\ 2\lambda_5 v_2 v_N & 2\lambda_5 v_1 v_N & 4(\lambda_3 + \lambda_5)v_N^2 \end{pmatrix} \quad (14)$$

in the basis $\sqrt{2}(\text{Re} H_1^0, \text{Re} H_2^0, \text{Re} N)$. Its mass eigenstates S_i ($i = 1, \dots, 3$) are obtained as $S_i = \sqrt{2}(U_{i1}^S \text{Re} H_1^0 + U_{i2}^S \text{Re} H_2^0 + U_{i3}^S \text{Re} N)$.

The original Higgs fields can be expressed in terms of the physical fields as

$$H_1^0 = v_1 + \sqrt{\frac{1}{2}} \left(U_{i1}^S S_i + i U_{i1}^P P_i \sin \beta \right), \quad (15)$$

$$H_2^0 = v_2 + \sqrt{\frac{1}{2}} \left(U_{i2}^S S_i + i U_{i1}^P P_i \cos \beta \right), \quad (16)$$

$$N = v_N + \sqrt{\frac{1}{2}} \left(U_{i3}^S S_i + i U_{i2}^P P_i \right). \quad (17)$$

^bF. Vissani has kindly pointed out that a particle that does not mix with ordinary neutrinos should not be called a “neutrino.” I keep this name for lack of a better one.

This gives the interactions of the Higgs bosons with the quarks, the leptons and the right-handed neutrino,

$$\mathcal{L}_{\text{int}} = \frac{gm_{\nu_s}}{2m_W} \frac{v}{v_N} \left[U_{i3}^S S_i \bar{\nu}_s \nu_s + i U_{i2}^P P_i \bar{\nu}_s \gamma_5 \nu_s \right] \quad (18)$$

$$+ \frac{gm_u}{2m_W} \left[\frac{U_{i2}^S}{\sin \beta} S_i \bar{u} u + i \cot \beta U_{i1}^P P_i \bar{u} \gamma_5 u \right] \quad (19)$$

$$+ \frac{gm_d}{2m_W} \left[\frac{U_{i1}^S}{\cos \beta} S_i \bar{d} d + i \tan \beta U_{i1}^P P_i \bar{d} \gamma_5 d \right] \quad (20)$$

It is then easy to work out the annihilation cross section,

$$\sigma v = \frac{G_F^2 m_\nu^4}{\pi} \frac{v^2}{v_N^2} \left[\sum_{k=1}^2 \frac{U_{k2}^P U_{k1}^P}{m_{P_k}^2 - 4m_\nu^2} \right]^2 \sum_f c_f \beta_f m_f^2 \kappa_f^2, \quad (21)$$

and the scattering cross section off nucleons,

$$\sigma_{\nu N} = \frac{G_F^2}{\pi} \frac{2m_N^4 m_\nu^4}{(m_N + m_\nu)^2} \frac{v^2}{v_N^2} \left[\sum_{j=1}^3 \frac{U_{j3}^S}{m_{S_j}^2} \left(\frac{k_d U_{j1}^S}{\cos \beta} + \frac{k_u U_{j2}^S}{\sin \beta} \right) \right]^2. \quad (22)$$

Here $c_f = 3$ for quarks, $c_f = 1$ for leptons, $\kappa_f = \cot \beta$ for up-type quarks, and $\kappa_f = \tan \beta$ for down-type quarks and leptons. Moreover, $k_d = \langle m_d \bar{d} d + m_s \bar{s} s + m_b \bar{b} b \rangle = 0.21$ and $k_u = \langle m_u \bar{u} u + m_c \bar{c} c + m_t \bar{t} t \rangle = 0.15$.

Table 1 lists important quantities for two interesting cases: a 3 GeV neutrino and a 7 GeV neutrino.

Table 1. Two interesting cases.

	“3 GeV”	“7 GeV”	Experimental/ Observational Values
m_ν [GeV]	3	7	
$\phi_\gamma(56^\circ)$ [photons(>1GeV)/cm ² /s/sr]	7.3×10^{-7}	7.4×10^{-7}	$\sim 7.6 \times 10^{-7}$
Ω ($H = 60$ km/s/Mpc)	0.11	0.12	~ 0.2
$\phi_{\bar{p}}(200\text{MeV})$ [\bar{p} /cm ² /s/sr/GeV]	~ 0	7×10^{-6}	$\lesssim 3 \times 10^{-6}$
$\sigma_{\chi N}$ [pb]	1.2×10^{-3}	1.3×10^{-4}	see text
m_{P_1}, m_{P_2} [GeV]	5, 155	4.2, 156	
$m_{S_1}, m_{S_2}, m_{S_3}$ [GeV]	5, 155, 220	4.2, 156, 220	
m_{H^\pm} [GeV]	110	110	$\gtrsim 53$
$\Gamma(Z \rightarrow f\bar{f}P, S)/\Gamma(Z \rightarrow f\bar{f})$	1.3×10^{-6}	6.3×10^{-6}	$\gtrsim 10^{-4}$
“ $\sin^2(\beta - \alpha)$ ” (Zh)	2.3×10^{-7}	2.3×10^{-6}	$\gtrsim 10^{-2}$
“ $\cos^2(\beta - \alpha)$ ” (hA)	6×10^{-7}	6×10^{-6}	$\lesssim 0.3$

For the “1 GeV” model, the parameters are $h = 0.9$, $v/v_N = 52$, $\tan \beta = v_2/v_1 = 40$, and $\lambda_1, \dots, \lambda_6 = g^2$ (the square of the SU(2) coupling constant). The gamma-ray

intensity is close to the observed one and the relic density is in the cosmologically interesting range. There is practically no antiproton flux from annihilations because only the $\tau^+\tau^-$ channel is effective. The scattering cross section off nucleons is compatible with the present bound of 2×10^{-2} pb. Constraints on the Higgs sector are discussed below.

For the “7 GeV” model, the parameters are $h = 0.8$, $v/v_N = 20$, $\tan \beta = 33$, and $\lambda_1, \dots, \lambda_6 = g^2$. The gamma-ray intensity is closed to the observed one, and the relic density is in the cosmologically interesting range. The antiproton flux is slightly higher than the measured one, but, given the big uncertainties in its estimation, it is compatible with it. The scattering cross section is smaller than the present limit of 3×10^{-4} pb. The Higgs sector is discussed in the following.

Two Higgs bosons (S_1 and P_1) are light and one has to worry about their production at colliders. But this production is very suppressed because the light Higgs bosons are mostly singlets. First, the LEP bounds on the search of Higgs bosons in two-doublet models are satisfied. Since there is an additional Higgs singlet, $\sin^2(\beta - \alpha)$ in the bound from $e^+e^- \rightarrow Zh$ must be replaced with $|U_{11}^S \cos \beta + U_{12}^S \sin \beta|^2$, and $\cos^2(\beta - \alpha)$ in the bound from $e^+e^- \rightarrow hA$ must be replaced with $|U_{12}^P U_{12}^S - U_{11}^P U_{11}^S|^2$. In the allowed region, the reinterpreted $\sin^2(\beta - \alpha)$ and $\cos^2(\beta - \alpha)$ are smaller than 5×10^{-5} , and are not excluded by accelerator searches (Decamp et al. 1992). Secondly, Higgs bremsstrahlung from final state leptons and quarks in Z decays is suppressed by the small mixing between singlet and doublet Higgs bosons. In the present model,

$$\frac{\Gamma(Z \rightarrow f\bar{f}X)}{\Gamma(Z \rightarrow f\bar{f})} = \frac{\sqrt{2}G_F}{4\pi^2} g\left(\frac{m_X}{m_Z}\right) c_f m_f^2 A_{Xf}^2, \quad (23)$$

where $A_{Pif} = U_{i1}^P \cot \beta$, $A_{Sif} = U_{i1}^S / \cos \beta$ for up-type quarks, $A_{Pif} = U_{i1}^P \tan \beta$, $A_{Sif} = U_{i2}^S / \sin \beta$ for down-type quarks and leptons, $c_f = 3$ for quarks and $c_f = 1$ for leptons. The function $g(y)$ comes from the phase-space integration, and is $g(y) \cong 1$ at $m_X = 12$ GeV. The values obtained in the two examples (see table 1) are lower than the LEP constraint on two leptons + two jets production $\Gamma(Z \rightarrow f\bar{f}X)/\Gamma(Z \rightarrow f\bar{f}) < 1.5 \times 10^{-4}$ at $m_X = 12$ GeV (Decamp et al. 1992).

I conclude that these two examples satisfy the experimental and observational constraints considered. A more general scan of the model parameter space is under way (Gondolo 1998b).

6. Conclusions

I have shown that it may be possible to explain the Dixon et al. gamma-ray emission from the galactic halo as due to halo WIMP annihilations. Not only the intensity and spatial pattern of the halo emission can be matched but also the relic density of the candidate WIMP can be in the cosmologically interesting domain. After a model-independent analysis to learn about the properties of a suitable candidate,

I have presented a working model: a sterile neutrino in a model with an extended Higgs sector. Two examples in parameter space indicate the existence of an interesting region in the model parameter space in which present observational and experimental constraints are satisfied and the gamma-ray emission is reproduced.

7. Acknowledgments

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8. References

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