

Cosmic Ray Anomalies and Dark Matter Annihilation to Muons via a Higgs Portal Hidden Sector

Kazunori Kohri,^{*} John McDonald,[†] and Narendra Sahu[‡]

*Cosmology and Astroparticle Physics Group,
University of Lancaster, Lancaster LA1 4YB, UK*

Abstract

Annihilating dark matter (DM) models based on a scalar hidden sector with Higgs portal-like couplings to the Standard Model are considered as a possible explanation for recently observed cosmic ray excesses. Two versions of the model are studied, one with non-thermal DM as the origin of the boost factor and one with Sommerfeld enhancement. In the case of non-thermal DM, four hidden sector scalars which transform under a $U(1)_X$ symmetry are added. The heaviest scalars decouple and later decay to DM scalars, so providing the boost factor necessary to explain the present DM annihilation rate. The mass of the annihilating scalars is limited to $\lesssim 600$ GeV for the model to remain perturbative. $U(1)_X$ breaking to Z_2 at the electroweak transition mixes light $O(100)$ MeV hidden sector scalars with the Higgs. The DM scalars annihilate to these light scalars, which subsequently decay to two $\mu^+\mu^-$ pairs via Higgs mixing, so generating a positron excess without antiprotons. Decay to $\mu^+\mu^-$ rather than e^+e^- is necessary to ensure a fast enough light scalar decay rate to evade light scalar domination at nucleosynthesis. In the version with Sommerfeld enhancement only three new scalars are necessary. TeV scale DM masses can be accommodated, allowing both the higher energy electron plus positron excess and the lower energy PAMELA positron excess to be explained. DM annihilates to $2\mu^+\mu^-$ pairs as in the non-thermal model. This annihilation mode may be favoured by recent observations of the electron plus positron excess by FERMI and HESS.

PACS numbers: 12.60.Jv, 98.80.Cq, 95.35.+d

^{*}Electronic address: k.kohri@lancaster.ac.uk

[†]Electronic address: j.mcdonald@lancaster.ac.uk

[‡]Electronic address: n.sahu@lancaster.ac.uk

I. INTRODUCTION

Recent results from the satellite experiment PAMELA [1] indicate an excess of positrons at 10-100 GeV as compared with the expected galactic background, confirming the earlier results from HEAT [2] and AMS [3]. Surprisingly, PAMELA did not find any antiproton excess below 100 GeV as compared with the galactic background [4]. Evidence was obtained from the balloon experiments ATIC [5] and PPB-BETS [6] of an excess electron plus positron flux as compared with the galactic background in the energy range 100-800 GeV. These results have recently been reconsidered by FERMI [7] and HESS [8], which do not confirm the large excess and spectral features observed by ATIC and PPB-BETS. However, HESS does not rule out the possibility of an electron plus positron excess, although there is no indication of structure in the electron plus positron spectrum [8], while FERMI observes a flattening of the electron spectrum relative to that predicted by a conventional diffusive model for the background, which suggests new physics, although again no prominent spectral features are observed [7]. Therefore an electron plus positron excess remains a possibility. These results raise the exciting prospect that the positron and the electron plus positron excesses could be attributed to annihilation of dark matter (DM) particles¹. If DM annihilation is the explanation for the positron excess at lower energies and the possible electron plus positron excess at higher energies, then the annihilation rate of DM at present should be larger than that expected from the canonical thermal relic annihilation cross-section in the case of a smooth distribution of DM in the galaxy ($\approx 3 \times 10^{-26}$ cm³/s) [14]. This is the boost factor². The origin of the boost factor could be astrophysical, because of the merger of sub-structures, or entirely from particle physics, or a combination of the two.

A popular method to achieve the boost factor is Sommerfeld enhancement of the DM annihilation cross-section [14–17]. This typically requires the introduction of new light bosons of mass $M_B \sim \alpha M_{DM}$ in order to mediate a force between the DM particles, where M_{DM} is the DM particle mass and α is the interaction's fine-structure constant. (An exception is discussed in [15], where the enhancement is mediated by electroweak interactions.) An alternative approach is to use non-thermal production of DM to accommodate a large annihilation cross-section, usually via

¹ Nearby astrophysical sources [9–11] and decaying DM [12, 13] are also possible explanations.

² Different authors define the boost factor in different ways, with some reserving this term for the astrophysical boost due to clumpy DM. We will use it to refer to the total enhancement of the DM annihilation rate.

heavy particle decay³ [18]. In the following we will consider both possibilities⁴.

In addition, in order to produce positrons without a sizable amount of antiprotons, a mechanism to allow DM to annihilate primarily to leptons is required. One approach is to introduce new ‘leptophilic’ couplings of the DM particles to leptons [23]. In this paper we will instead consider a DM sector which interacts with the SM via generic (non-leptophilic) couplings. Our goal is to determine whether such non-leptophilic models can account for the observed cosmic ray excesses and to obtain the necessary conditions on their masses, couplings and field content. Our analysis of the ingredients required to construct successful non-leptophilic models may then guide the construction of more complete models which can explain the necessary features.

Our models are based on a SM singlet scalar sector interacting with the SM via Higgs portal-like interactions. Adding a scalar DM particle S is a particularly simple way to extend the SM to account for DM [24, 25]. The scalar is typically stabilised by either a discrete Z_2 or $U(1)$ symmetry. It interacts with the SM sector via the coupling $S^\dagger SH^\dagger H$ (which has come to be known as the ‘Higgs portal’ [26]), which is the only renormalizable coupling of the S to SM particles. Several DM models based on this type of coupling have been proposed [27]. However, this coupling alone cannot account for DM annihilation primarily to leptons, nor can it account for the boost factor. Here we extend the symmetry of the DM particle to a sector of SM singlet scalar fields. The DM sector is composed of the DM scalar S and additional scalar fields χ_i , all of which carry non-trivial charges under a symmetry $U(1)_X$. In the model with non-thermal DM, three scalars $\chi_i, i = 1, 2, 3$, are introduced. The heaviest scalar χ_1 populates the number density of DM, so providing the boost factor, while the lightest scalar χ_3 ensures an annihilation channel of DM to two $\mu^+\mu^-$ pairs. The role of χ_2 is to mix χ_3 with the Higgs via its vacuum expectation value (VEV), inducing its decay to leptonic final states. In order to ensure that χ_3 decays before dominating at nucleosynthesis, $U(1)_X$ must be broken to a discrete Z_2 symmetry which maintains the stability of S . This occurs spontaneously at the electroweak phase transition, when χ_2 acquires a VEV. $U(1)_X$ breaking also mixes the lightest scalar χ_3 with the Higgs boson, providing a mechanism for leptonic annihilation of S to $2\mu^+\mu^-$ pairs⁵. In this version of the model the mass of S is constrained by perturbativity to be less than approximately 600 GeV. In the version of the model with Sommerfeld enhancement, the light scalar χ_3 which mixes with the Higgs is also used to mediate the enhancement. This

³ We note that in SUSY models a natural alternative is Q-ball decay [19, 20].

⁴ Another possibility, annihilation close to a pole, has been considered in [21, 22].

⁵ The Higgs mixing mechanism was first described in [16].

model requires only two scalars, χ_2 and χ_3 , and can accommodate a larger range of S mass. DM again annihilates to $2\mu^+\mu^-$ pairs via Higgs mixing as in the model with non-thermal DM. Annihilation to muons appears to be favoured by recent data from FERMI and HESS [28, 29], with annihilation to 4μ being favoured by the analysis of [29].

Our paper is organised as follows. In Section II we present and discuss a Higgs portal model with non-thermal production of DM as the source of the boost factor. In Section III we discuss a version of the model with Sommerfeld enhancement in place of non-thermal production. In Section IV we present our conclusions.

II. A HIGGS PORTAL MODEL WITH NON-THERMAL DARK MATTER

A. Overview

We extend the SM by adding a dark sector composed of a singlet DM scalar S of mass M_S and three additional scalar fields χ_i , $i = 1, 2, 3$ with masses m_i , $i = 1, 2, 3$, such that $m_1 \gg m_2 \gg m_3$. We impose a symmetry $U(1)_X$, under which the fields χ_i carry a charge $+1$ and S carries a charge $+3/2$. In order to avoid a Goldstone boson from $U(1)_X$ breaking we will consider this to be a gauge symmetry. The gauge interaction will not have any significant effect on the cosmological evolution of the model, only contributing to the already rapid annihilation and scattering between the hidden sector scalars. The SM fields are neutral under $U(1)_X$. The dark sector fields interact with the SM via Higgs portal-type couplings to the Higgs bilinear $H^\dagger H$. $U(1)_X$ will be broken at the electroweak (EW) phase transition to a surviving Z_2 symmetry under which S is odd while rest of the fields, including the SM fields, are even. Since S is odd under the surviving Z_2 symmetry, it is stable and a candidate for DM.

Our model is based on generic couplings of the gauge singlet scalars to $H^\dagger H$. The renormalizable couplings of the scalar sector of the Lagrangian are given by

$$\begin{aligned} \mathcal{L} \supseteq & m_i^2 \chi_i^\dagger \chi_i + M_S^2 S^\dagger S + m_H^2 H^\dagger H \\ & + \lambda_S (S^\dagger S)^2 + \lambda_H (H^\dagger H)^2 + \gamma S^\dagger S H^\dagger H \\ & + \eta_{ijkl} \chi_i^\dagger \chi_j \chi_k^\dagger \chi_l + (\eta_S)_{ij} \chi_i^\dagger \chi_j S^\dagger S + (\eta_H)_{ij} \chi_i^\dagger \chi_j H^\dagger H, \end{aligned} \quad (1)$$

where H is the SM Higgs and $i = 1, 2, 3$. We assume that all couplings are real and that all particles in the dark sector have positive masses squared ($m_i^2, M_S^2 > 0$). As usual $m_H^2 < 0$ so that H acquires

a vacuum expectation value (VEV). If $(\eta_H)_{22} < 0$ and m_2^2 is sufficiently small then χ_2 gains a $U(1)_X$ -breaking VEV when the electroweak (EW) phase transition occurs.

The cosmological evolution of the model can be summarized as follows. χ_1 is assumed to have the largest mass in the hidden sector. S has a mass $\gtrsim 100$ GeV if its annihilation is to account for the PAMELA observations and so freezes-out at a temperature $\gtrsim 5$ GeV. A key requirement of the model is that χ_1 decay occurs sufficiently long after the S density has frozen out of thermal equilibrium that it can boost the S relic density. χ_1 can decay to $\chi_i^\dagger \chi_j \chi_k$, $\chi_i S^\dagger$ or $\chi_i H^\dagger H$ ($i, j, k \neq 1$). After the EW phase transition, when χ_2 acquires a VEV, χ_1 also decays through the two body processes: $\chi_1 \rightarrow \chi_j^\dagger \chi_k$, with $j, k \neq 1$ and $\chi_1 \rightarrow S^\dagger S, hh, \chi_j h$ with $j \neq 1$, where h is the physical Higgs scalar. We will see that late decay of χ_1 requires these couplings to be very small, of order 10^{-10} . Due to $U(1)_X$ -breaking, χ_3 mixes with the physical Higgs scalar h and decays to SM fermions via the Yukawa couplings. If the χ_3 mass is in the range $2m_\mu$ to $2m_{\pi^0}$ (212-270 MeV) then it decays predominantly to $\mu^+ \mu^-$ pairs. For larger χ_3 mass, χ_3 decay to π^0 pairs produces photons while χ_3 decay to nucleon-antinucleon pairs produces antiprotons. Whether the photon flux from pion decay is excluded depends on the nature of the DM halo, with cuspy NFW halos excluded but cored isothermal halos still likely to be consistent with present bounds⁶ [30, 31]. In the following we will consider the χ_3 mass to be in the range $2m_\mu$ to $2m_{\pi^0}$, although the upper bound may be increased to $2m_{proton}$ if the photon flux from pion decay is within observational limits. This range of χ_3 mass requires that the couplings of χ_3 to H and χ_2 are less than 10^{-6} . The lifetime for χ_3 decay to $\mu^+ \mu^-$ is short compared with the time at which nucleosynthesis begins, so the relic χ_3 density, which would otherwise dominate at nucleosynthesis, safely decays away. This would not be true for decay to $e^+ e^-$, which would apply if the χ_3 mass was less than $2m_\mu$. The present S density annihilates primarily to χ_3 pairs which promptly decay to muons. The subsequent decay $\mu^+ \rightarrow e^+ + \nu_e + \bar{\nu}_\mu$ then accounts for the positron flux without any antiproton flux.

⁶ A possible problem with photons from pion decay was previously noted in the context of an axion decay model for cosmic ray anomalies [13].

B. Electroweak Phase Transition and Spontaneous Breaking of $U(1)_X$

After the EW phase transition H develops a VEV, which triggers a VEV for χ_2 . The VEV of H and χ_2 also induce a VEV for χ_1 through the couplings η_{1222} and $(\eta_H)_{12}$

$$\langle \chi_1 \rangle \approx -\frac{\eta_{1222} u^3 + (\eta_H)_{12} uv^2}{m_1^2} \quad (2)$$

where $\langle \chi_2 \rangle = u$ and $\langle H \rangle = v$. As we will show, η_{1222} and $(\eta_H)_{12}$ are required to be no larger than $O(10^{-10})$ in order to ensure the late decay of χ_1 . Therefore with $u \sim v \sim 100$ GeV and $m_1 \sim O(1)$ TeV we find that $\langle \chi_1 \rangle$ is negligibly small, $O(100)$ eV. Similarly, χ_3 also gains a VEV

$$\langle \chi_3 \rangle \approx -\frac{\eta_{2333} u^3 + (\eta_H)_{23} uv^2}{m_3^2}. \quad (3)$$

We will show later that the mass of the lightest mass eigenstate χ_3' must be less than $O(1)$ GeV in order to ensure that it will decay primarily to leptons. This is most easily understood if all the terms in the mass matrix are less than 1 GeV, which in turn requires that all the couplings of χ_3 to χ_2 and H are less than $O(10^{-6})$. (Larger entries in the mass matrix are possible but would require sufficient cancellation between the contributions to the lightest mass eigenvalue.) Since m_3 is also no larger than $O(1)$ GeV, this means that $\langle \chi_3 \rangle \lesssim u, v$. Although a value for $\langle \chi_3 \rangle$ which is comparable to $\langle H \rangle$ and $\langle \chi_2 \rangle$ is possible, this will not qualitatively alter our results from the case where $\langle \chi_3 \rangle \ll u, v$, since it will only alter the admixtures of χ_2 and H in the lightest mass eigenstate by $O(1)$ factors. Therefore, to simplify the analysis we will set $\langle \chi_3 \rangle = 0$ in the following and consider

$$\langle \chi_1 \rangle = 0, \quad \langle S \rangle = 0, \quad \langle H \rangle = v, \quad \langle \chi_2 \rangle = u, \quad \langle \chi_3 \rangle = 0, \quad (4)$$

with u and v obtained by minimizing the scalar potential

$$V = m_H^2 v^2 + m_2^2 u^2 + \eta_{2222} u^4 + \lambda_H v^4 + (\eta_H)_{22} u^2 v^2. \quad (5)$$

We assume that $(\eta_H)_{22}$ and m_H^2 are negative with all other terms positive. Vacuum stability requires that $(\eta_H)_{22} > -2\sqrt{\eta_{2222}\lambda_H}$. Minimizing Eq. (5) gives

$$u = \sqrt{\frac{(2\lambda_H m_2^2 - (\eta_H)_{22} m_H^2)}{((\eta_H)_{22}^2 - 4\eta_{2222}\lambda_H)}} \quad (6)$$

and

$$v = \sqrt{\frac{(2\eta_{2222}m_H^2 - (\eta_H)_{22}m_2^2)}{((\eta_H)_{22}^2 - 4\eta_{2222}\lambda_H)}}. \quad (7)$$

In the following we will assume that η_{2222} , $|(\eta_H)_{22}|$ are ~ 0.1 and that $m_2 \lesssim 100$ GeV, therefore $u \sim v$.

The χ_2 expectation value breaks the $U(1)_X$ symmetry to a discrete symmetry, under which the scalars transform as $\phi \rightarrow e^{i2\pi Q}\phi$, where Q is the $U(1)_X$ charge of ϕ . Thus χ_i ($Q = 1$) transforms as $\chi_i \rightarrow \chi_i$, while S ($Q = 3/2$) transforms as $S \rightarrow -S$. Therefore S is stable due to the residual discrete symmetry.

Once χ_2 gains a $U(1)_X$ -breaking expectation value, the Higgs h (defined by $Re(H^0) \equiv (h + v)/\sqrt{2}$) mixes with the real parts of χ_1 , χ_2 and χ_3 . The χ_2 expectation value is assumed to be $\sim v$, so $\chi_2 - h$ mixing will be large, which should have significant consequences for Higgs phenomenology. $\chi_3 - h$ mixing provides the mechanism for DM to annihilate primarily to lepton final states. For simplicity we will consider only the mixing of the Higgs with χ_3 , which is responsible for the important physics. In the basis spanned by $\sqrt{2}Re\chi_3$ and h , the effective mass squared matrix is given by

$$\mathcal{M}^2 = \begin{pmatrix} m_3^2 + 6\eta_{2233}u^2 + (\eta_H)_{33}v^2 & (\eta_H)_{23uv} \\ (\eta_H)_{23uv} & 2\lambda_H v^2 + (\eta_H)_{22}u^2 \end{pmatrix}. \quad (8)$$

In this we have assumed that η_{ijkl} is independent of the order of i, j, k, l . Diagonalising this gives mass eigenstates χ_3' and h' with

$$M_{\chi_3'}^2 \approx m_3^2 + 6\eta_{2233}u^2 + (\eta_H)_{33}v^2 - \frac{((\eta_H)_{23uv})^2}{M_{h'}^2} \quad (9)$$

and

$$M_{h'}^2 \approx 2\lambda_H v^2 + (\eta_H)_{22}u^2, \quad (10)$$

where we assume $\chi_3 - h$ mixing is small. The $\chi_3 - h$ mixing angle is

$$\beta \approx \frac{(\eta_H)_{23uv}}{M_{h'}^2}. \quad (11)$$

In order to have DM annihilation to muon final states we require that $M_{\chi_3'}$ is in the range 212-270 MeV. Therefore we require that $m_3 \lesssim 300$ MeV, $\eta_{2233} \lesssim 10^{-6}$, $(\eta_H)_{33} \lesssim 10^{-6}$ and $(\eta_H)_{23} \lesssim 10^{-6}$, assuming that $u \sim v \sim M_{h'} \sim 100$ GeV. Although for simplicity we considered only the mixing between χ_3 and h , in general χ_3 will mix with χ_2 and χ_1 in addition to χ_3 . The corresponding

couplings will also be constrained by the requirement that the light eigenstate mass is $\sim O(100)$ MeV, therefore the additional mixings will not change the model qualitatively. This illustrates an important feature of generic Higgs portal models for cosmic ray excesses: some couplings must be strongly suppressed. We will show that suppressed couplings are also necessary to produce the boost factor via non-thermal production of DM.

C. Non-Thermal Production of S Dark Matter

The S density is due to out-of-equilibrium decay of χ_1 . This must occur at a sufficiently low temperature that the S scalars can have a boosted annihilation cross-section without annihilating away after being produced by χ_1 decay. An initial density produced by χ_1 decay at temperature T_{decay} will annihilate down to a density at T_{decay} given by

$$n_S(T_{decay}) = \frac{H(T_{decay})}{(\sigma|v_{rel}|)_S}, \quad (12)$$

where n_S is the number density and $(\sigma|v_{rel}|)_S$ is the annihilation cross-section times relative velocity (which is T independent for the case of annihilating scalars). Eq. (12) is true if the initial S number density from χ_1 decay is larger than $n_S(T_{decay})$. Since $n_S \propto g(T)T^3$ while $H \propto g(T)^{1/2}T^2$, where $g(T)$ is the effective number of relativistic degrees of freedom, Eq. (12) implies that the present ratio of the S number density from χ_1 decay to that from thermal freeze-out (which is given by Eq. (12) with T_{decay} replaced by the S freeze-out temperature T_S) is

$$\frac{n_{S\,decay}}{n_{S\,th}} = \left(\frac{g(T_S)}{g(T_{decay})} \right)^{1/2} \frac{T_S}{T_{decay}}. \quad (13)$$

Since the annihilation cross-section is enhanced by a factor B over that which accounts for observed DM via a thermal relic density, it follows that the thermal relic S density is smaller than that observed by a factor B . So in order to account for the observed DM via χ_1 decay we must require that $n_{S\,decay} \approx B \times n_{S\,th}$. Therefore

$$T_{decay} = \left(\frac{g(T_S)}{g(T_{decay})} \right)^{1/2} \frac{T_S}{B}. \quad (14)$$

In this we have neglected the logarithmic dependence of the freeze-out temperature on the annihilation cross-section and therefore treated T_S as a constant. T_S is related to the S mass by $T_S = M_S/z_S$ with $z_S \approx 20$ [33]. For example, if $m_S \sim 400$ GeV then $T_S \sim 20$ GeV. Since the observed positron and electron excess requires that $B \sim 10^3$, we would then require that $T_{decay} \sim 20$ MeV.

This very low decay temperature is difficult to achieve via particle decay, as it implies a very long-lived particle with lifetime $\tau \sim H^{-1} \sim 10^{-3}s$. χ_1 can decay to S pairs via the three-body decays $\chi_1 \rightarrow \chi_2 S^\dagger S$ and $\chi_3' S^\dagger S$. The decay rate is given by

$$\Gamma_{\chi_1} \approx \frac{\lambda^2}{128\pi^3} M_{\chi_1'} , \quad (15)$$

where $\lambda^2 = (\eta_S)_{12}^2 + (\eta_S)_{13}^2$. In order to have a late enough χ_1 decay we require that $\Gamma_{\chi_1} < H$ at $T = O(10)$ MeV, which implies that $\lambda \lesssim 10^{-10}$. χ_1 can also decay to S pairs via the two-body decay $\chi_1 \rightarrow S^\dagger S$ once χ_2 gains a VEV. The decay rate is given by

$$\Gamma_{\chi_1} \approx \frac{(\eta_S)_{12}^2}{16\pi} \frac{u^2}{M_{\chi_1'}} . \quad (16)$$

With $u \sim 100$ GeV and $M_{\chi_1'} \sim 1$ TeV, this decay rate as a function of $(\eta_S)_{12}$ is comparable to Eq.(15) as a function of λ . Therefore we also require that $(\eta_S)_{12} \lesssim 10^{-10}$. (Similar constraints apply to other χ_1 two-body decay modes.)

The above assumes that the initial S density from χ_1 decay is larger than that given in Eq. (12). This requires that $n_{\chi_1}(T_{decay}) > n_S(T_{decay})/2$, since each χ_1 decay produces 2 S . $n_{\chi_1}(T_{decay})$ is given by

$$n_{\chi_1}(T_{decay}) = \frac{g(T_{decay}) T_{decay}^3}{g(T_{\chi_1}) T_{\chi_1}^3} \frac{H(T_{\chi_1})}{(\sigma|v_{rel}|)_{\chi_1}} , \quad (17)$$

where T_{χ_1} is the χ_1 freeze-out temperature. The condition $n_{\chi_1}(T_{decay}) > n_S(T_{decay})/2$ then implies that

$$\frac{g(T_{decay})^{1/2}}{g(T_{\chi_1})^{1/2}} \frac{T_{decay}}{T_{\chi_1}} > \frac{1}{2} \frac{(\sigma|v_{rel}|)_{\chi_1}}{(\sigma|v_{rel}|)_S} . \quad (18)$$

This translates into an upper bound on B ,

$$B < B_0 \equiv \frac{2(\sigma|v_{rel}|)_S z_{\chi_1} M_S}{(\sigma|v_{rel}|)_{\chi_1} z_S M_{\chi_1}} \frac{g(T_S)^{1/2}}{g(T_{\chi_1})^{1/2}} . \quad (19)$$

Therefore if $B < B_0 \approx (\sigma|v_{rel}|)_S/(\sigma|v_{rel}|)_{\chi_1}$ (assuming that $M_{\chi_1} \sim M_S$ and $z_{\chi_1} \sim z_S$) then the required χ_1 decay temperature will be given by Eq. (14). The cross-section times relative velocity for non-relativistic χ_1 pair annihilation to S and H is given by

$$(\sigma|v_{rel}|)_{\chi_1} = \frac{1}{32\pi M_{\chi_1}^2} [(\eta_S)_{11}^2 + (\eta_H)_{11}^2] . \quad (20)$$

The annihilation cross-section times relative velocity for non-relativistic S is given by

$$(\sigma|v_{rel}|)_S = \frac{1}{32\pi M_S^2} [(\eta_S)_{ij}^2 + \gamma^2] , \quad i = 2, 3 . \quad (21)$$

With $B \sim 10^2 - 10^3$, Eq. (19) is therefore satisfied if $(\eta_H)_{11}, (\eta_S)_{11} \lesssim 10^{-2}$, assuming that $\gamma \sim 0.1$. If this is not satisfied and $B > B_0$, then the boost factor is given by B_0 rather than B . In this case the S density comes directly from χ_1 decay without subsequent annihilations. An even lower χ_1 decay temperature and smaller $(\eta_S)_{12}$ would then be necessary in order to account for the observed DM density.

D. Dark Matter Annihilation Rate

The dominant S annihilation mode is assumed to be to χ_3' pairs. In this case the annihilation cross-section times relative velocity is given by

$$(\sigma|v_{rel}|)_S = \frac{(\eta_S)_{33}^2}{32\pi M_S^2}. \quad (22)$$

In order to account for the cosmic ray excesses, the annihilation cross-section times relative velocity necessary to account for thermal relic DM, $(\sigma|v_{rel}|) \approx 3 \times 10^{-26} \text{cm}^3/\text{s} \equiv 2.6 \times 10^{-9} \text{GeV}^{-2}$, must be boosted by $B \sim 10^2 - 10^3$ for DM masses in the range $\text{O}(100)\text{GeV} - \text{O}(1)\text{TeV}$. This requires that

$$(\eta_S)_{33} \approx (5 - 16) \times \left(\frac{M_S}{1 \text{ TeV}} \right). \quad (23)$$

Therefore for the theory to remain perturbative ($(\eta_S)_{33} \lesssim 3$) we require that $M_S \lesssim 600$ (190) GeV for $B = 10^2$ (10^3). Thus while the model can account for the positron excess in the range 1-100 GeV observed by PAMELA, an electron plus positron excess at energies up to $\text{O}(1)\text{TeV}$ cannot be explained if the model is to remain perturbative, in which case an alternative explanation for the electron plus positron excess is required, most likely astrophysical. This conclusion is likely to apply rather generally to models which do not have Sommerfeld enhancement of the annihilation cross-section. (However, this does not exclude the possibility of a large annihilation cross-section due to strong coupling between S and χ_3 , which is only constrained by unitarity [32].)

E. Leptonic Final States via Dark Matter Annihilation to χ_3'

In order to account for the positron excess without an accompanying antiproton flux, the S annihilations at present should proceed primarily through leptonic decay channels. In our model this is achieved through a mixture of χ_3 -Higgs mixing and kinematics. $U(1)_X$ breaking due to $\langle \chi_2 \rangle$ causes the real part of χ_3 to mix with h . If the dominant S annihilation mode is $S^\dagger S \rightarrow \chi_3'^\dagger \chi_3'$

and if the χ'_3 mass is in the range $2m_\mu < M_{\chi'_3} < 2m_{\pi^0}$, then the mixing of χ_3 with h leads to the decay $\chi'_3 \rightarrow \mu^+\mu^-$ via the muon Yukawa coupling. This is illustrated in Figure 1, treating the mixing as a mass insertion. Thus S annihilation will produce a 4μ final state via the process shown in Figure 2.

The decay rate for $\chi'_3 \rightarrow \mu^+\mu^-$ is given by

$$\Gamma_{\chi'_3} = \frac{\beta^2 Y_\mu^2}{8\pi} M_{\chi'_3} \equiv \frac{(\eta_H)_{23}^2 Y_\mu^2}{8\pi} \left(\frac{uv}{M_h^2} \right)^2 M_{\chi'_3}, \quad (24)$$

where Y_μ is the Yukawa coupling of the SM Higgs to $\mu^+\mu^-$. This gives for the lifetime of χ'_3

$$\begin{aligned} \tau_{\chi'_3} \equiv \frac{1}{\Gamma_{\chi'_3}} &= 4 \times 10^{-4} \left(\frac{10^{-6}}{(\eta_H)_{23}} \right)^2 \left(\frac{6.07 \times 10^{-4}}{Y_\mu} \right)^2 \\ &\times \left(\frac{M_h}{150 \text{ GeV}} \right)^4 \left(\frac{200 \text{ MeV}}{M_{\chi'_3}} \right) \text{ s}, \end{aligned} \quad (25)$$

where we have used $u = 100$ GeV and $v = 174$ GeV. The short lifetime of χ'_3 ensures that the thermally produced χ'_3 will decay well before the onset of nucleosynthesis at $O(1)$ s. This is essential, as the χ'_3 will freeze-out while relativistic (since there are no annihilation channels for χ'_3 once $T \lesssim M_{\chi_2}, M_{h'}$) and so they will dominate the energy density at nucleosynthesis. The decay rate to e^+e^- is suppressed by $Y_e^2/Y_\mu^2 \sim 10^{-5}$, leading to a lifetime ~ 10 s. Thus for χ'_3 to decay before nucleosynthesis, *decay to muon pairs must be kinematically allowed*⁷.

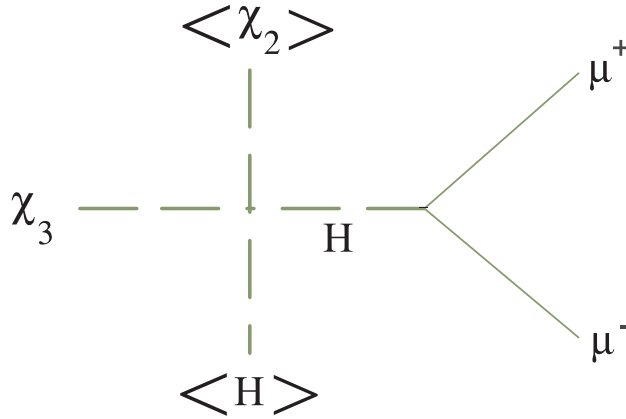


FIG. 1: χ'_3 decay into $\mu^+\mu^-$ via χ_3 -Higgs mixing and the muon Yukawa coupling.

⁷ We note that the mechanism described in [16], which is based on a single scalar field ϕ and which gives a decay rate equivalent to that here but with $\langle \chi_2 \rangle$ replaced by $\langle \chi_3 \rangle \sim 200$ MeV (where χ_3 is the equivalent of ϕ in [16]), results in a lifetime which is too long and so ϕ domination at nucleosynthesis.

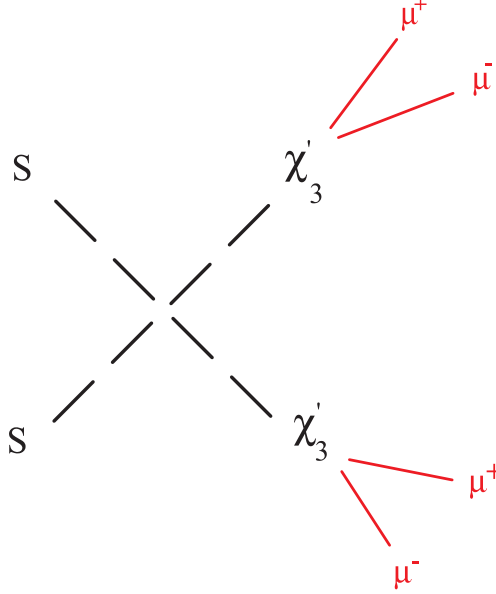


FIG. 2: The primary S annihilation process to 4μ via χ_3' decay.

The previous discussion applies to the decay of the real part of χ_3 , which mixes with the physical Higgs. The imaginary part of χ_3 does not mix with the Higgs and is therefore stable. However, so long as the χ_3 self-coupling η_{3333} is large, the χ_3 scalars will maintain a thermal equilibrium with each other even after they have decoupled from thermal equilibrium with other particles. Since the imaginary part of χ_3 is generally heavier than the real part after the latter mixes with the Higgs, the imaginary part of χ_3 will annihilate to the lighter real part once $T < M_{\chi_3}$, thus ensuring that the entire χ_3 density can decay to muons prior to nucleosynthesis.

F. Sub-dominant S annihilation to Higgs pairs

We have so far considered the annihilation $S^\dagger S \rightarrow \chi_3' \chi_3'$ via the quartic coupling $(\eta_S)_{33}$, which primarily produces $\mu^+ \mu^-$ pairs. However, it is also possible to have $S^\dagger S \rightarrow H^\dagger H$ via the coupling⁸ γ . The branching ratio to Higgs pairs $B_{S^\dagger S \rightarrow H^\dagger H} \approx \gamma^2 / (\eta_S)_{33}^2$ should be small enough that the production of Higgs pairs does not result in a large antiproton signal. This requires that $B_{S^\dagger S \rightarrow H^\dagger H} \lesssim 0.1$. However, this still allows a significant coupling to Higgs pairs, which can contribute a small antiproton component to the cosmic rays from DM annihilation. The $S^\dagger S H^\dagger H$

⁸ We are assuming that the S mass is sufficiently large that we can approximately calculate the annihilation rate in the $\langle H \rangle \rightarrow 0$ limit.

coupling also mediates the coupling of S to nucleons, which may allow direct detection of DM. These possibilities are distinctive features of Higgs portal models which distinguish them from models based on purely leptophilic couplings.

G. Positron Excesses from $S^\dagger S$ Annihilation

The annihilation of $S^\dagger S$ pairs will give rise to mostly μ^\pm pairs which finally decay to e^\pm and neutrinos. The electrons and positrons from $S^\dagger S$ annihilation then travel under the influence of the galactic magnetic field and therefore the motion of e^\pm is expected to be a random walk. As a result a fraction of e^\pm flux will reach the solar system.

The positron flux in the vicinity of the solar system can be obtained by solving the diffusion equation [15, 34, 35]

$$\frac{\partial}{\partial t} f_{e^+}(E, \vec{r}, t) = K_{e^+}(E) \nabla^2 f_{e^+}(E, \vec{r}, t) + \frac{\partial}{\partial t} [b(E) f_{e^+}(E, \vec{r}, t)] + Q(E, \vec{r}), \quad (26)$$

where $f_{e^+}(E, \vec{r}, t)$ is the number density of positrons per unit energy E , $K_{e^+}(E)$ is the diffusion constant, $b(E)$ is the energy-loss rate and $Q(E, \vec{r})$ is the positron source term. The positron source term $Q(E, \vec{r})$ from $S^\dagger S$ annihilation is given by

$$Q(E, \vec{r}) = n_S^2(\vec{r}) \sigma_S |v_{\text{rel}}| \frac{dN_{e^+}}{dE}. \quad (27)$$

In the above equation the fragmentation function dN_{e^+}/dE represents the number of positrons with energy E which are produced from the annihilation of $S^\dagger S$.

We assume that the positrons are in a steady state, i.e. $\partial f_{e^+}/\partial t = 0$. Then from Eq. (26), the positron flux in the vicinity of the solar system can be given in a semi-analytical form [15, 34, 35]

$$\Phi_{e^+}(E, \vec{r}_\odot) = \frac{v_{e^+}}{4\pi b(E)} (n_S)_\odot^2 \sigma_S |v_{\text{rel}}| \int_E^{M_S} dE' \frac{dN_{e^+}}{dE'} I(\lambda_D(E, E')), \quad (28)$$

where $\lambda_D(E, E')$ is the diffusion length from energy E' to energy E and $I(\lambda_D(E, E'))$ is the halo function which is independent of particle physics. An analogous solution for electron flux can also be obtained.

Positrons in our galaxy are not only produced by $S^\dagger S$ annihilation but also by the scattering of cosmic-ray protons with the interstellar medium [36]. Thus the positrons produced from the latter sources can act as background for the positrons produced from the annihilation of $S^\dagger S$.

The background fluxes [36] of primary and secondary electrons and secondary positrons can be parameterized as [37]:

$$\begin{aligned}
\Phi_{\text{prim, e}^-}^{\text{bkg}} &= \frac{0.16\varepsilon^{-1.1}}{1 + 11\varepsilon^{0.9} + 3.2\varepsilon^{2.15}} \text{GeV}^{-1} \text{cm}^{-2} \text{s}^{-1} \text{sr}^{-1} \\
\Phi_{\text{sec, e}^-}^{\text{bkg}} &= \frac{0.70\varepsilon^{0.7}}{1 + 11\varepsilon^{1.5} + 600\varepsilon^{2.9} + 580\varepsilon^{4.2}} \text{GeV}^{-1} \text{cm}^{-2} \text{s}^{-1} \text{sr}^{-1} \\
\Phi_{\text{sec, e}^+}^{\text{bkg}} &= \frac{4.5\varepsilon^{0.7}}{1 + 650\varepsilon^{2.3} + 1500\varepsilon^{4.2}} \text{GeV}^{-1} \text{cm}^{-2} \text{s}^{-1} \text{sr}^{-1}, \tag{29}
\end{aligned}$$

where the dimensionless parameter $\varepsilon=E/(1 \text{ GeV})$. The net positron flux in the galactic medium is then given by

$$(\Phi_{e^+})_{\text{Gal}} = (\Phi_{e^+})_{\text{bkg}} + \Phi_{e^+}(E, \vec{r}_{\odot}). \tag{30}$$

The second term in the above equation is given by Eq. (28), which depends on various factors: $b(E)$, $\lambda_D(E, E')$, $I(\lambda_D(E, E'))$, v_{e^+} , $(n_S)_{\odot}$ and the injection spectrum dN_{e^+}/dE' . The energy loss term $b(E)$ (due to inverse Compton scattering and synchrotron radiation due to the galactic magnetic field) is determined by the photon density and the strength of magnetic field. Its value is taken to be $b(E) = 10^{-16}\varepsilon^2 \text{GeV s}^{-1}$ [37]. The number density of S DM in the solar system is given by

$$(n_S)_{\odot} = \frac{\rho_{\odot}}{M_S}, \tag{31}$$

where $\rho_{\odot} \approx 0.3 \text{GeV}/\text{cm}^3$. In the energy range we are interested in, the value of v_{e^+} is taken approximately to be c , the velocity of light. The values of diffusion length $\lambda_D(E, E')$ and the corresponding halo function $I(\lambda_D(E, E'))$ are based on astrophysical assumptions [15, 34, 35]. By considering different heights of the galactic plane and different DM halo profiles the results may vary slightly. In the following we take the height of the galactic plane to be $\lesssim 4 \text{ kpc}$, which is referred to as the "MED" model [15, 34, 35], and we have used the NFW DM halo profile [38],

$$\rho(r) = \rho_{\odot} \left(\frac{r_{\odot}}{r} \right) \left(\frac{1 + \left(\frac{r_{\odot}}{r_s} \right)}{1 + \left(\frac{r}{r_s} \right)} \right)^2, \tag{32}$$

to determine the halo function $I(\lambda_D(E, E'))$, where $r_s \approx 20 \text{kpc}$ and $r_{\odot} \approx 8.5 \text{kpc}$. (We find that our results are not strongly sensitive to the halo profile.) In Figure 3, plotted using DARKSUSY [39], the positron fraction from $S^{\dagger}S$ annihilation is compared with the data from AMS, HEAT and PAMELA for the case of $M_S = 600 \text{ GeV}$, showing that a good fit is obtained in this case.

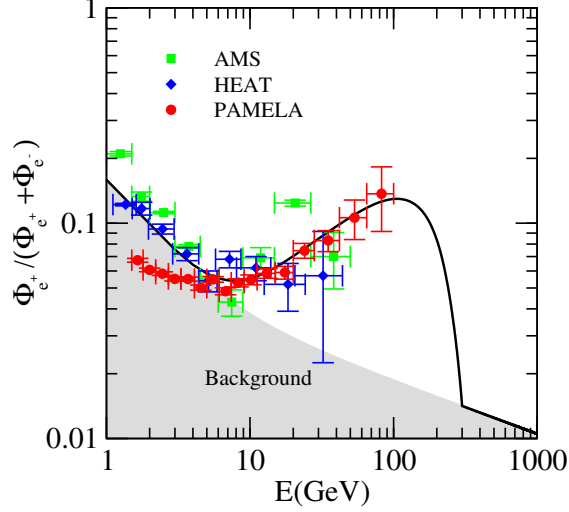


FIG. 3: Positron fraction from $S^\dagger S \rightarrow 2\mu^+\mu^-$ at $M_S = 600$ GeV. We have used the annihilation cross-section $\langle\sigma|v|\rangle = 4.5 \times 10^{-23} \text{cm}^3/\text{s}$.

H. Nucleosynthesis, Gamma-Ray and CMB Constraints on Enhanced $S^\dagger S$ Annihilation

So far we have considered large annihilation cross sections of the order of $10^{-23} \text{cm}^3 \text{s}^{-1}$ in order to fit the excess of the observed cosmic-ray electron fraction. This value is approximately $10^2 - 10^3$ times larger than the canonical value of the annihilation cross section for thermal relic DM ($\simeq 3 \times 10^{-26} \text{cm}^3 \text{s}^{-1}$). Therefore we have to check if this value is consistent with other cosmological and astrophysical constraints, in particular those from nucleosynthesis and due to gamma-rays from the galactic centre (GC) and halo. We will consider S annihilation primarily to χ'_3 pairs, but we will include the possibility of a small but significant branching ratio to Higgs pairs.

First we shall discuss constraints which come from BBN [40–43]. Even after the freeze-out of $S^\dagger S$ annihilations, a small amount of S pairs continue to annihilate. In our model, the $S^\dagger S$ pair dominantly annihilates into $\mu^+\mu^-$ pairs with some fraction into $H^\dagger H$.

The photodissociation of D and ^4He is severely constrained by observational value of $^3\text{He}/\text{D}$. According to [42] we have a constraint on the annihilation cross section into $\mu^+\mu^-$ pairs,

$$\langle\sigma v\rangle < \langle\sigma v\rangle_{S^\dagger S \rightarrow 2\mu^+\mu^-}^{\text{photo}} = \frac{1.0 \times 10^{-21}}{B_{S^\dagger S \rightarrow 2\mu^+\mu^-}} \text{cm}^3 \text{s}^{-1} \left(\frac{E_{\text{vis}}/M_S}{0.7} \right)^{-1} \left(\frac{M_S}{1 \text{ TeV}} \right), \quad (33)$$

where E_{vis}/M_S represents the fraction of the total energy $2M_S$ which goes into visible energy E_{vis} i.e. charged particles and photons. $B_{S^\dagger S \rightarrow 2\mu^+\mu^-}$ is the branching ratio for $S^\dagger S$ annihilation into 2

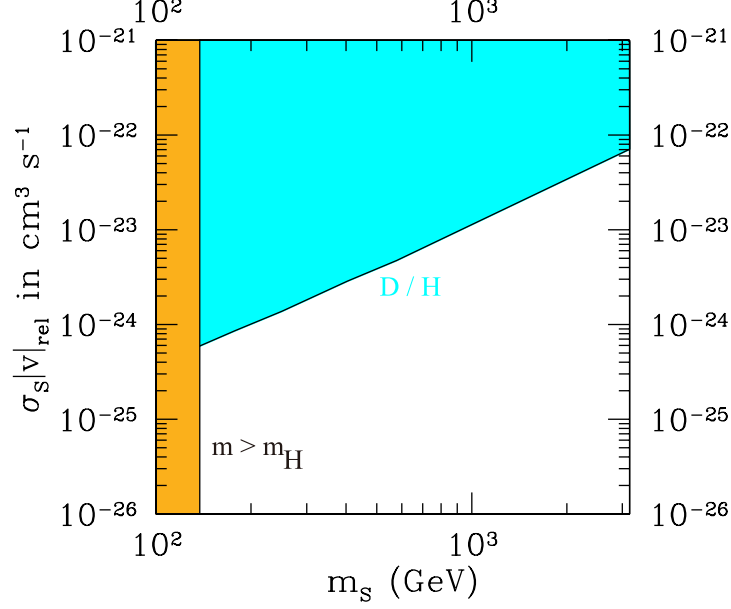


FIG. 4: Upper bounds from BBN on the annihilation cross section of $S^\dagger S$ into a Higgs $H^\dagger H$ pair as a function of the DM mass, where the branching ratio is normalized to $B_{S^\dagger S \rightarrow H^\dagger H} = 1$. Here we have assumed the mass of Higgs boson is 130 GeV. The name of the light element used for the constraint is written near each line. The vertical band at the left side indicates the region which is not kinematically allowed.

$\mu^+ \mu^-$ pairs. In case of the muon decay, $E_{\text{vis}} \sim 0.7 M_S$.

In addition, in order to limit the branching ratio to Higgs pairs, we have calculated the constraints on the cross section to $H^\dagger H$ which follow from photodissociation and hadron emission. For photodissociation we find [42]

$$\langle \sigma v \rangle < \langle \sigma v \rangle_{S^\dagger S \rightarrow H^\dagger H}^{\text{photo}} = \frac{7.0 \times 10^{-22}}{B_{S^\dagger S \rightarrow H^\dagger H}} \text{cm}^3 \text{s}^{-1} \left(\frac{E_{\text{vis}}/M_S}{1.0} \right)^{-1} \left(\frac{M_S}{1 \text{ TeV}} \right), \quad (34)$$

where PYTHIA [44] gives $E_{\text{vis}} \sim 1.0 M_S$ and $B_{S^\dagger S \rightarrow H^\dagger H}$ is the branching ratio into $H^\dagger H$. (In the low energy limit this becomes the total branching ratio to W , Z and h .) In the current case $B_{S^\dagger S \rightarrow H^\dagger H} = 1 - B_{S^\dagger S \rightarrow 2\mu^+ \mu^-}$. The dominant upper bound comes from the smaller of $\langle \sigma v \rangle_{S^\dagger S \rightarrow 2\mu^+ \mu^-}^{\text{photo}}$ and $\langle \sigma v \rangle_{S^\dagger S \rightarrow H^\dagger H}^{\text{photo}}$. These bounds are generally compatible with the range of values required to account for the cosmic ray excesses, $\langle \sigma v \rangle \sim 3 \times 10^{-24} - 3 \times 10^{-23} \text{cm}^3 \text{s}^{-1}$ for boost factors $10^2 - 10^3$.

The most severe bound on hadron emission comes from the overproduction of deuterium by the

destruction of ${}^4\text{He}$. This process is constrained by observational D/H [42]. From this we obtain the hadron emission constraint

$$\langle\sigma v\rangle < \langle\sigma v\rangle_{S^\dagger S \rightarrow H^\dagger H}^{\text{had}} = \frac{1.3 \times 10^{-23}}{B_{S^\dagger S \rightarrow H^\dagger H}} \text{cm}^3 \text{s}^{-1} \left(\frac{N_n}{1.0}\right)^{-1} \left(\frac{M_S}{1 \text{ TeV}}\right)^{1.5}, \quad (35)$$

with N_n the number of emitted neutrons per single annihilation. In the case of $H^\dagger H$ emission, N_n is approximately 1.0, which is obtained using PYTHIA [44]. This is again consistent with the range of $\langle\sigma v\rangle$ required to account for the observed cosmic ray excesses. We have plotted the results for annihilation into $H^\dagger H$ in Figure 4, with the normalization $B_{S^\dagger S \rightarrow H^\dagger H} = 1$.

In summary, for the range of S mass which is compatible with perturbative couplings, the boost factor required to account for the positron and/or electron plus positron excess via annihilation to muons is compatible with present BBN constraints.

We next consider constraints from gamma-rays. A possible gamma-ray signal from the GC due to DM annihilation has been extensively studied as it could provide a good method to study the nature of DM astrophysically. So far the HESS group has reported that power-law signals were observed from the GC [45, 46] for $200 \text{ GeV} \lesssim E_\gamma \lesssim 700 \text{ GeV}$. Quite recently the FERMI satellite group also reported their preliminary result for the signals observed from the galactic mid-latitude ($10^\circ < |b| < 20^\circ$) for $200 \text{ MeV} \lesssim E_\gamma \lesssim 10 \text{ GeV}$. When we adopt a cuspy profile of the galaxy, such as the NFW profile, the gamma-ray signal from muon emission can exceed the observed signal. However, if we take a milder profile such as the cored isothermal profile, then for the moment DM annihilation is not constrained by the current observations [30, 31, 47]⁹. To clarify the dependence of the DM constraints on the density profile, we need more accurate data on the diffuse gamma-ray background, which will be provided by FERMI in the near future.

In addition, there are CMB constraints on the enhanced $S^\dagger S$ annihilation cross-section. It has been shown in ref. [48, 49] that energetic particles from rapid $S^\dagger S$ annihilation can reionize neutral hydrogen at the last scattering surface, leaving an imprint on the CMB. The analysis of [48] concludes that current data from WMAP5 imposes a 2- σ upper bound on the $S^\dagger S$ annihilation cross-section which is given by

$$\langle\sigma v\rangle_{S^\dagger S \rightarrow \chi_3^\dagger \chi_3} < \frac{3.6 \times 10^{-24} \text{cm}^3/\text{s}}{f} \left(\frac{M_S}{1 \text{ TeV}}\right), \quad (36)$$

⁹ Note that the positron and electron plus positron signals will not change even if we used the cored isothermal profile because local annihilation within 1 kpc dominates the production of electrons and positrons with $\gtrsim 10 \text{ GeV}$ energies.

where f is in the range $0.2 - 0.3$ for annihilation to $2\mu^+\mu^-$ pairs. Thus a boost factor of $O(1000)$ is marginally allowed by the current data.

Finally, we briefly comment of the possibility of neutrino signals from the GC. Detecting such neutrino signals in the future might be useful to distinguish the Higgs portal DM model from others, since muon neutrinos are produced by the decay of the $\mu^+\mu^-$ pairs coming from DM annihilation and subsequent χ_3' decay. So far Super-K has reported upper bounds on the up-going muon flux coming from neutrinos emitted from the GC [50]. We can compare the theoretical prediction of the neutrino flux in our model with this Super-K upper bound. According to the discussion of Ref. [51], our model is presently allowed since neutrinos are not produced directly but indirectly through the decay of the charged leptons and possibly mesons. It is expected that future neutrino experiments such as KM3NeT [52, 53] or IceCube DeepCore [54, 55] will be able to detect the up-going muons induced by the neutrinos emitted from the GC.

III. A SOMMERFELD ENHANCED VERSION OF THE MODEL

In the previous section we studied the conditions for a successful Higgs portal model with non-thermal production of DM. In this section we will replace non-thermal production with Sommerfeld enhancement of the annihilation cross-section as the source of the boost factor. The main difference between the two models is the reduced number of hidden sector scalar fields, since χ_1 is no longer needed as the source of the non-thermal DM density. This will also eliminate the most heavily suppressed $O(10^{-10})$ couplings, which were necessary to ensure the late decay of χ_1 . The S DM annihilation to χ_3 pairs and subsequent χ_3 decay to $\mu^+\mu^-$ pairs is unchanged from the non-thermal scenario.

Since in the non-thermal model there must exist a light scalar χ_3 if we wish to avoid leptophilic couplings, it is natural to ask whether we can eliminate χ_1 and consider instead thermal DM with a Sommerfeld enhanced annihilation cross-section, with the attractive force mediated by χ_3 -exchange. The correct thermal relic density of S DM is obtained if the $S^\dagger S \chi_3^\dagger \chi_3$ coupling is in the range $0.1-1$ for $M_S \sim 0.1 - 1$ TeV [24]. If we then consider the coupling $(\eta_S)_{23}$ in Eq.(1) and introduce $\langle \chi_2 \rangle$, we obtain the interaction

$$(\eta_S)_{23} \langle \chi_2 \rangle \chi_3 S^\dagger S + \text{h.c.} . \quad (37)$$

This interaction can produce the required long-range force between S particles via χ_3 exchange.

The condition for a Sommerfeld enhanced annihilation rate is $M_{\chi_3} \lesssim \alpha M_S$, where $\alpha = \lambda^2/4\pi$ and the effective coupling from χ_3 exchange is $\lambda \approx (\eta_S)_{23} \langle \chi_2 \rangle / M_S$. Therefore

$$M_{\chi_3} \lesssim 1 \text{ GeV} (\eta_S)_{23}^2 \left(\frac{\langle \chi_2 \rangle}{100 \text{ GeV}} \right)^2 \left(\frac{1 \text{ TeV}}{M_S} \right). \quad (38)$$

Since $M_{\chi_3} \sim 200 \text{ MeV}$ in our model, this will be satisfied if $(\eta_S)_{23} \gtrsim 0.4$ when $M_S \sim 1 \text{ TeV}$.

Therefore, in addition to simplifying the model by eliminating χ_1 , Sommerfeld enhancement permits larger DM masses, $M_S \sim 1 \text{ TeV}$. This may be significant in light of recent analyses [28, 29] of the new FERMI and HESS electron plus positron data, which favour DM particles with TeV scale masses which annihilate to muons (with the case of annihilation to 4μ being favoured by the analysis of [29]). Since the Higgs portal models generally predict that DM annihilates to two $\mu^+\mu^-$ pairs via decay of the primary χ_3 pair, a Sommerfeld enhanced version of the Higgs portal model, in contrast with the non-thermal model, could provide an explanation for both the higher energy electron plus positron excess and the lower energy PAMELA positron excess.

IV. CONCLUSIONS

We have considered two DM models for cosmic ray excesses which are based on Higgs portal-type couplings of a scalar DM sector to the SM, one with non-thermal DM as the explanation of the boost factor and the other with thermal DM and Sommerfeld enhancement of the annihilation cross-section.

In the case of the model with non-thermal production of DM, the DM scalar mass must be less than about 600 GeV if the model is to remain perturbative. Therefore if this model is correct then the PAMELA positron excess can be explained by DM annihilation but the higher energy electron plus positron flux suggested by FERMI and HESS must have a different explanation. This is likely to be true of most models without Sommerfeld enhancement. Non-thermal production of DM is possible via quartic scalar couplings. However, the couplings leading to decay of the heavy scalar which produces the DM density must be highly suppressed, $\lesssim 10^{-10}$, in order to ensure that the heavy particle decays well after the DM particle freezes-out.

A successful model must also account for DM annihilation to primarily leptonic states. If we do not wish to introduce DM which couples preferentially to leptons then the only way to achieve this is kinematically, by ensuring that DM annihilates to unstable final states which are too light to subsequently decay to hadrons. Our model can generate the required decay process via mixing

of the χ_3 scalar of the hidden sector with the Higgs, leading to the decay of χ_3 primarily to $\mu^+\mu^-$ via the muon Yukawa coupling if its mass is in the range $2m_\mu - 2m_{\pi^0}$ (212 – 270 MeV). The small χ_3 mass requires that the quartic scalar couplings of χ_3 to the Higgs and to χ_2 are $\lesssim 10^{-6}$. The $\mu^+\mu^-$ final state is essential if we require that the χ_3 density decays prior to nucleosynthesis (which χ_3 would otherwise dominate) but does not decay to pions or nucleons, which would produce a potentially dangerous photon or antiproton flux. This is a clear prediction of the Higgs portal model, which applies equally to the Sommerfeld enhanced version.

We conclude that quartic couplings of a relatively simple scalar DM sector can achieve the required enhancement of the annihilation rate and leptonic final states, but appropriate mixtures of strongly suppressed and unsuppressed quartic couplings and large and small mass terms are required. In the absence of symmetries or dynamical effects which can explain them, such hierarchies would appear unnatural. It is therefore to be hoped that the pattern of masses and couplings can be understood in terms of the symmetries or dynamics of a complete theory, for which the present model is the low energy effective theory.

A significant feature of the Higgs portal model, which can distinguish it from those with purely leptophilic annihilation modes, is that there can be a significant coupling of DM to Higgs pairs. This could produce a non-negligible antiproton component in the cosmic rays from DM annihilation if the annihilation process $S^\dagger S \rightarrow H^\dagger H$ is not too suppressed relative to the dominant process $S^\dagger S \rightarrow \chi_3^\dagger \chi_3$. The $S^\dagger S H^\dagger H$ coupling may also allow direct detection of DM.

Constraints from BBN are important for the model with non-thermal DM, since the annihilation rate is large at all temperatures. We found that both the muon and Higgs final states are consistent with an annihilation cross section as large as $10^{-23} \text{ cm}^3 \text{ s}^{-1}$ for $M_S \lesssim 600 \text{ GeV}$. The model is also consistent with the gamma-ray signal from the galactic centre and from the diffuse gamma-ray background in the case of a cored isothermal halo profile, but not in the case of a cuspy NFW profile.

In the Sommerfeld enhanced version of the model, the low mass χ_3 scalar which accounts for leptonic DM annihilation also mediates the force responsible for the Sommerfeld enhancement. In this case we can reduce the number of additional scalars by one, since χ_1 is no longer needed to produce the DM non-thermally, which also eliminates the most highly suppressed couplings. This version of the model can accommodate a TeV scale DM particle, allowing it to explain the electron plus positron excess suggested by FERMI and HESS as well as the positron excess observed by PAMELA. Exactly as in the non-thermal model, DM annihilates to χ_3 pairs which subsequently

decay via Higgs mixing to $2\mu^+\mu^-$ pairs. This may be significant, as recent analyses suggest that the new FERMI and HESS electron plus positron data favours TeV scale DM particles annihilating to muons [28, 29], with annihilation to 4μ via intermediate decaying scalars being favoured by [29].

The Higgs portal models considered here should have phenomenological signals due to the coupling of the DM sector to the Higgs bilinear and the mixing of the Higgs with the SM singlet χ_2 [56]. If S or χ_2 are light enough then they may be produced via Higgs decay at the LHC. The mass eigenstate Higgs boson is also expected to have a large singlet component, with consequences for Higgs phenomenology. These features may not be unique to our model, but they would provide indirect support for it. In addition, the muon neutrinos produced by the decay of the $\mu^+\mu^-$ pair from DM annihilation may be detectable via upward-moving muons at future neutrino experiments.

Acknowledgement

The research at Lancaster is supported by EU grant MRTN-CT-2004-503369 and the Marie Curie Research and Training Network "UniverseNet" (MRTN-CT-2006-035863).

-
- [1] O. Adriani *et al.*, arXiv:0810.4995 [astro-ph].
 - [2] S. W. Barwick *et al.* [HEAT Collaboration], *Astrophys. J.* **482**, L191 (1997) [arXiv:astro-ph/9703192]; J. J. Beatty *et al.*, *Phys. Rev. Lett.* **93** (2004) 241102 [arXiv:astro-ph/0412230].
 - [3] M. Aguilar *et al.* [AMS-01 Collaboration], *Phys. Lett. B* **646**, 145 (2007) [arXiv:astro-ph/0703154].
 - [4] O. Adriani *et al.*, arXiv:0810.4994 [astro-ph].
 - [5] J. Chang *et al.*, *Nature* **456**, 362 (2008).
 - [6] S. Torii *et al.*, arXiv:0809.0760 [astro-ph].
 - [7] F. Collaboration, arXiv:0905.0025 [astro-ph.HE].
 - [8] H. E. S. Aharonian, arXiv:0905.0105 [astro-ph.HE].
 - [9] D. Hooper, A. Stebbins and K. M. Zurek, arXiv:0812.3202 [hep-ph].

- [10] D. Hooper, P. Blasi and P. D. Serpico, *JCAP* **0901**, 025 (2009) [arXiv:0810.1527 [astro-ph]] ; S. Profumo, arXiv:0812.4457 [astro-ph].
- [11] H. Yuksel, M. D. Kistler, and T. Stanev, arXiv:0810.2784; K. Ioka, arXiv:0812.4851; N. J. Shaviv, E. Nakar, and T. Piran, arXiv:0902.0376; E. Borriello, A. Cuoco and G. Miele, arXiv:0903.1852 [astro-ph.GA]; P. Blasi, arXiv:0903.2794; N. Kawanaka, K. Ioka and M. M. Nojiri, arXiv:0903.3782 [astro-ph.HE]; Y. Fujita, K. Kohri, R. Yamazaki and K. Ioka, arXiv:0903.5298 [astro-ph.HE]; P. Blasi and P. D. Serpico, arXiv:0904.0871.
- [12] S. Shirai, F. Takahashi and T. T. Yanagida, arXiv:0905.0388 [hep-ph]; C. H. Chen, C. Q. Geng and D. V. Zhuridov, arXiv:0905.0652 [hep-ph].
- [13] Y. Nomura and J. Thaler, arXiv:0810.5397 [hep-ph].
- [14] M. Cirelli, M. Kadastik, M. Raidal and A. Strumia, arXiv:0809.2409 [hep-ph].
- [15] J. Hisano, S. Matsumoto, O. Saito and M. Senami, *Phys. Rev. D* **73**, 055004 (2006) [arXiv:hep-ph/0511118].
- [16] N. Arkani-Hamed, D. P. Finkbeiner, T. R. Slatyer and N. Weiner, *Phys. Rev. D* **79**, 015014 (2009) [arXiv:0810.0713 [hep-ph]].
- [17] M. Pospelov and A. Ritz, *Phys. Lett. B* **671** (2009) 391 [arXiv:0810.1502 [hep-ph]].
- [18] T. Moroi and L. Randall, *Nucl. Phys. B* **570**, 455 (2000) [arXiv:hep-ph/9906527]; M. Fairbairn and J. Zupan, arXiv:0810.4147 [hep-ph].
- [19] J. McDonald, arXiv:0904.0969 [hep-ph].
- [20] K. Enqvist and J. McDonald, *Phys. Lett. B* **440**, 59 (1998) [arXiv:hep-ph/9807269]; M. Fujii and K. Hamaguchi, *Phys. Lett. B* **525**, 143 (2002) [arXiv:hep-ph/0110072]; *Phys. Rev. D* **66**, 083501 (2002) [arXiv:hep-ph/0205044]; M. Fujii and M. Ibe, *Phys. Rev. D* **69**, 035006 (2004) [arXiv:hep-ph/0308118].
- [21] D. Feldman, Z. Liu and P. Nath, *Phys. Rev. D* **79**, 063509 (2009) [arXiv:0810.5762 [hep-ph]].
- [22] M. Ibe, H. Murayama and T. T. Yanagida, *Phys. Rev. D* **79**, 095009 (2009) [arXiv:0812.0072 [hep-ph]]; W. L. Guo and Y. L. Wu, *Phys. Rev. D* **79**, 055012 (2009) [arXiv:0901.1450 [hep-ph]]; M. Ibe, Y. Nakayama, H. Murayama and T. T. Yanagida, *JHEP* **0904**, 087 (2009) [arXiv:0902.2914 [hep-ph]]. F. Y. Cyr-Racine, S. Profumo and K. Sigurdson, arXiv:0904.3933 [astro-ph.CO].
- [23] P. J. Fox and E. Poppitz, arXiv:0811.0399 [hep-ph]. H. S. Goh, L. J. Hall and P. Kumar, arXiv:0902.0814 [hep-ph].
- [24] J. McDonald, *Phys. Rev. D* **50**, 3637 (1994).

- [25] V. Silveira and A. Zee, Phys. Lett. B **161**, 136 (1985); C. P. Burgess, M. Pospelov and T. ter Veldhuis, Nucl. Phys. B **619**, 709 (2001) [arXiv:hep-ph/0011335].
- [26] B. Patt and F. Wilczek, arXiv:hep-ph/0605188.
- [27] M. C. Bento, O. Bertolami, R. Rosenfeld and L. Teodoro, Phys. Rev. D **62**, 041302 (2000) [arXiv:astro-ph/0003350]; M. C. Bento, O. Bertolami and R. Rosenfeld, Phys. Lett. B **518**, 276 (2001) [arXiv:hep-ph/0103340]; J. McDonald, Phys. Rev. Lett. **88**, 091304 (2002); L. Lopez Honorez, E. Nezri, J. F. Oliver and M. H. G. Tytgat, JCAP **0702**, 028 (2007) [arXiv:hep-ph/0612275]; N. Sahu and U. Sarkar, Phys. Rev. D **76**, 045014 (2007) [arXiv:hep-ph/0701062]; J. McDonald, N. Sahu and U. Sarkar, JCAP **0804**, 037 (2008) [arXiv:0711.4820 [hep-ph]]; J. March-Russell, S. M. West, D. Cumberbatch and D. Hooper, JHEP **0807**, 058 (2008) [arXiv:0801.3440 [hep-ph]]; J. McDonald and N. Sahu, JCAP **0806**, 026 (2008) [arXiv:0802.3847 [hep-ph]]; Y. G. Kim, K. Y. Lee and S. Shin, JHEP **0805** (2008) 100 [arXiv:0803.2932 [hep-ph]]; D. G. Cerdeno, C. Munoz and O. Seto, arXiv:0807.3029 [hep-ph].
- [28] L. Bergstrom, J. Edsjo and G. Zaharijas, arXiv:0905.0333 [astro-ph.HE].
- [29] P. Meade, M. Papucci, A. Strumia and T. Volansky, arXiv:0905.0480 [hep-ph].
- [30] G. Bertone, M. Cirelli, A. Strumia and M. Taoso, arXiv:0811.3744 [astro-ph].
- [31] M. Kawasaki, K. Kohri and K. Nakayama, arXiv:0904.3626 [astro-ph.CO].
- [32] K. Griest and M. Kamionkowski, Phys. Rev. Lett. **64**, 615 (1990).
- [33] B. W. Lee and S. Weinberg, Phys. Rev. Lett. **39**, 165 (1977).
- [34] T. Delahaye, R. Lineros, F. Donato, N. Fornengo and P. Salati, Phys. Rev. D **77**, 063527 (2008) [arXiv:0712.2312 [astro-ph]].
- [35] M. Cirelli, R. Franceschini and A. Strumia, Nucl. Phys. B **800**, 204 (2008) [arXiv:0802.3378 [hep-ph]].
- [36] I. V. Moskalenko and A. W. Strong, Astrophys. J. **493**, 694 (1998) [arXiv:astro-ph/9710124].
- [37] E. A. Baltz and J. Edsjo, Phys. Rev. D **59**, 023511 (1999) [arXiv:astro-ph/9808243].
- [38] J.F. Navarro, C.S. Frenk and S.D.M. White, Astrophys.J.**462**, 563 (1996) .
- [39] P. Gondolo, J. Edsjo, L. Bergstrom, P. Ullio, M. Schelke and E.A. Baltz, [astro-ph/0406204]. Dark-SUSY homepage <http://www.physto.se/~edsjo/darksusy/>, JCAP **0407** (2004) 008.
- [40] K. Jedamzik, Phys. Rev. D **70**, 083510 (2004).
- [41] J. Hisano, M. Kawasaki, K. Kohri and K. Nakayama, arXiv:0810.1892 [hep-ph].
- [42] J. Hisano, M. Kawasaki, K. Kohri, T. Moroi and K. Nakayama, arXiv:0901.3582 [hep-ph].

- [43] M. Kawasaki, K. Kohri and T. Moroi, Phys. Lett. B **625**, 7 (2005); Phys. Rev. D **71**, 083502 (2005).
- [44] T. Sjostrand, S. Mrenna and P. Skands, JHEP **0605**, 026 (2006).
- [45] F. Aharonian *et al.* [The HESS Collaboration], Astron. Astrophys. **425**, L13 (2004);
- [46] F. Aharonian *et al.* [H.E.S.S. Collaboration], Phys. Rev. Lett. **97**, 221102 (2006) [Erratum-ibid. **97**, 249901 (2006)].
- [47] P. Salucci, A. Lapi, C. Tonini, G. Gentile, I. Yegorova and U. Klein, Mon. Not. Roy. Astron. Soc. **378**, 41 (2007) [arXiv:astro-ph/0703115]; M. Persic, P. Salucci and F. Stel, Mon. Not. Roy. Astron. Soc. **281**, 27 (1996) [arXiv:astro-ph/9506004].
- [48] T. R. Slatyer, N. Padmanabhan and D. P. Finkbeiner, Phys. Rev. D **80**, 043526 (2009) [arXiv:0906.1197 [astro-ph.CO]].
- [49] T. Kanzaki, M. Kawasaki and K. Nakayama, arXiv:0907.3985 [astro-ph.CO].
- [50] S. Desai *et al.* [Super-Kamiokande Collaboration], Phys. Rev. D **70**, 083523 (2004) [Erratum-ibid. D **70**, 109901 (2004)]; Astropart. Phys. **29**, 42 (2008).
- [51] J. Hisano, M. Kawasaki, K. Kohri and K. Nakayama, arXiv:0812.0219 [hep-ph].
- [52] Kappes, A. et al. 2007, ApJ, 656, 870
- [53] Kappes, A. & Consortium, f. t. K. 2007, arXiv:0711.0563
- [54] Achterberg, A. et al. 2007, PRD, 76, 042008 [Erratum: PRD, 77, 089904 (2008)]
- [55] D. F. Cowen [IceCube Collaboration], J. Phys. Conf. Ser. **110**, 062005 (2008).
- [56] V. Barger, P. Langacker, M. McCaskey, M. J. Ramsey-Musolf and G. Shaughnessy, Phys. Rev. D **77**, 035005 (2008) [arXiv:0706.4311 [hep-ph]].