

Likely dominance of WIMP annihilation to fermion pair +W/Z

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Abstract: Arguably, the most popular candidate for Dark Matter (DM) is a massive, stable, Majorana fermion. However, annihilation of Majorana DM to two fermions features a helicity-suppressed s -wave rate. The process radiating a gauge boson via electroweak (EW) and electromagnetic (EM) bremsstrahlung removes this s -wave suppression, and is likely to be the dominant modes of gauge-singlet Majorana DM annihilation. Given their enhanced annihilation rate with radiated W and Z gauge bosons, and the subsequent dominant W/Z decays via hadronic channels, Majorana DM tends to produce more antiprotons than positrons. This result contrasts with observations, thereby presenting a challenge to model building with Majorana DM.

Keywords: icrc2013, Dark Majorana annihilation, indirect Dark Matter detection, electroweak bremsstrahlung.

1 Introduction

The “WIMP miracle” (WIMP, for Weakly Interacting Massive Particle) for Dark Matter (DM) physics is that $\sigma_{\text{ann}} \sim (\frac{\alpha_{\text{weak}}}{4\pi})^2 / M_{\text{WIMP}}^2 \sim \text{pb} \times (\frac{200 \text{ GeV}}{M_{\text{WIMP}}})^2$ yields the critical thermally-averaged annihilation “rate” $\langle v\sigma \rangle \sim 3 \times 10^{-26} \text{ cm}^3/\text{s}$ to effect a DM abundance in agreement with what is observed today, $\Omega_{\text{DM}} \sim 0.24$. Here, v is the DM Møller velocity. Among the many candidates for DM, the most common particle-type is the Majorana fermion (χ). A splendid example is the neutralino of supersymmetry models; being a partner to a massless gauge-boson in the unbroken SUSY theory, the neutralino has two independent spin states, not four. Thus, χ is a Majorana fermion. Here we focus not on iso-vector DM like the neutralino, but rather on gauge-singlet Majorana DM. (Note that for gauge-singlet DM, the decay to $W^+ W^-$ and ZZ modes is very suppressed by quantum number conservation.)

Majorana-ness means that in the annihilation process to a fermion pair, $\chi\chi \rightarrow f\bar{f}$, the two initial state χ 's contribute two amplitudes, the usual one and a crossed- χ diagram with a relative minus sign (deducible from Fig. 1) when the external gauge boson lines are removed). The process is mediated by a t - and u -channel exchange of a virtual scalar particle which we label η . For fermionic dark matter, Fierz transformations effect the natural projection of $2 \rightarrow 2$ t - and u -exchange processes into partial waves. When the fermion currents (connected fermion lines) are Fierz rearranged into “charge retention” order (a χ line and a light fermion line), the result is an axial vector coupling in the s -channel, plus corrections proportional to the χ 's velocity, v . This comes about because for each t -channel diagram (the first three shown in Fig. 1), there is an accompanying u -channel diagram (the last three shown in Fig. 1), obtained by interchanging the momentum and spin of the two Majorana fermions. The relative sign between the t - and u -channel amplitudes is (-1) in accord with Fermi statistics. Therefore, one obtains an elegant simplification: V and T couplings in the Fierz bilinears of the χ -current are zero to $\mathcal{O}(v)$, and so ignorable. Furthermore, the chiral nature of the fermion couplings offers no S or P terms after Fierzing. Thus, after Fierzing, only the axial vector coupling of the χ -current is significant. It has been known for many years [1] that the spin/orbital angular momentum of the s -

channel axial-vector coupling requires a helicity flip of one of the produced fermions in the $L = 0$ (“ s -wave”) amplitude. Thus, the s -wave amplitude is suppressed $\sim (m_f/M_\chi)$. The $L = 1$ (“ p -wave”) amplitude does not require a helicity flip. However, on general grounds, the L^{th} partial wave amplitude is suppressed as v^L , and so the rate from the p -wave is suppressed as $\langle v^2 \rangle$.

Standard statistical mechanics simulations reveal that decoupling occurs at $T_{\text{dec}}/M_\chi \sim \frac{1}{6} \langle v^2 \rangle \sim 1/20$ to $1/50$; the resulting $\langle v^2 \rangle$ is therefore ~ 0.1 to 0.3 , which implies only a mild p -wave suppression at DM decoupling in the Early Universe. Today, $v \sim 300 \text{ km/s} \sim 10^{-3} c$ in galactic halos, so the p -wave contribution is highly suppressed by $\sim 10^{-6}$, and only the s -wave contribution is expected to be significant. However, as we have mentioned, in Majorana fermion DM models, the s -wave annihilation into a fermion pair $\chi\chi \rightarrow f\bar{f}$ is helicity suppressed by the factor $(m_f/M_\chi)^2$. These two suppressions are quite general: The helicity-suppression of the s -wave amplitude applies to all Majorana fermion $\chi\chi \rightarrow f\bar{f}$, while the velocity-suppressed p -wave applies to all DM $\rightarrow f\bar{f}$, Majorana or otherwise.

It is becoming increasingly appreciated that if the light fermion pair is produced in association with a gauge boson, then the spins of the final state can match the s -wave angular momentum requirement without a helicity flip. Thus, there is no s -wave mass-suppression factor for the $2 \rightarrow 3$ process $\chi\chi \rightarrow f\bar{f} + \text{gauge-boson}$. This unsuppressed s -wave was first calculated in 1989 for the photon bremsstrahlung reaction $\chi\chi \rightarrow f\bar{f}\gamma$ [2, 3], where the photon is radiated from one of the external particle legs (final state radiation, FSR) or from a the virtual mediator particle η (internal bremsstrahlung, IB). On the face of it, the radiative rate is down by the usual QED coupling factor of $\alpha/4\pi \sim 10^{-3}$. However, and significantly, photon bremsstrahlung can lift the helicity suppression of the s -wave process, which more than compensates for the extra coupling factor. (And if the dark matter annihilates to colored fermions, radiation of a gluon would also lift the helicity suppression.)

The importance of electroweak radiative corrections to dark matter annihilation was recognized more recently. Electroweak bremsstrahlung was investigated first in the context of cosmic rays [4, 5, 6, 7, 8, 9, 10]; the possibility of W/Z bremsstrahlung to lift initial-state velocity and final-

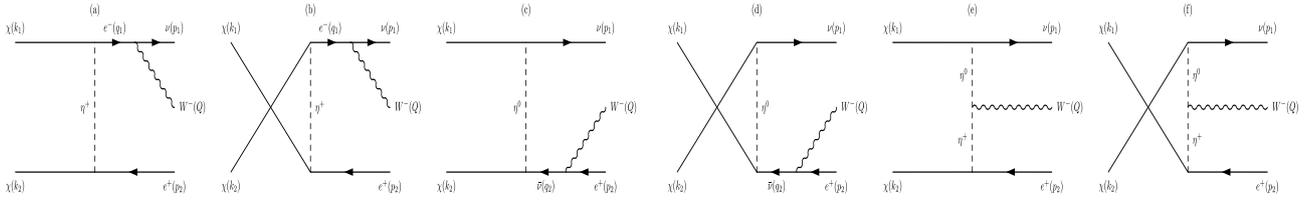


Figure 1: The t - (left three) and u - (right three) channel Feynman diagrams for $\chi\chi \rightarrow e^+ \nu W^-$. All fermion momenta in the diagrams flow with the arrow except p_2 and q_2 , with $q_1 = p_1 + Q$, $q_2 = p_2 + Q$.

state helicity suppressions was alluded to in Refs.[6, 9], and the unsuppressed s -wave in the context of $\chi\chi \rightarrow f\bar{f} + W/Z$ was noted in [11, 12]. Some signature channels for indirect detection of DM χ 's were calculated in Ref. [13], where the point was made that even if the $\chi\chi \rightarrow 2$ process is tuned to suppress direct $p\bar{p}$ production, the $2 \rightarrow 3$ EW process will necessarily produce an antiproton signal via the production and hadronic decay of the W and Z . As a result, the experimental upper limit on the cosmic antiproton flux provides a meaningful constraint on Majorana DM [13]. The dominance of the helicity-unsuppressed $f\bar{f}W/Z$ channel has also been elaborated upon in Refs. [14, 15].

The diagrams contributing to EW bremsstrahlung are shown in Fig. 1. There are a few important distinctions between electromagnetic (EM) and electroweak (EW) bremsstrahlung. An obvious one is that EM bremsstrahlung produces just photons, whereas EW bremsstrahlung and subsequent decay of the W or Z leads to leptons (including positrons), hadrons (including antiprotons) and gamma rays, offering correlated ‘‘multi-messenger’’ signals for indirect dark matter searches. This is an important result for future DM indirect searches. Another distinction between EW and EM bremsstrahlung is that in the former the massive W/Z 's have a longitudinal mode not available to the photon. The rate of longitudinal W_L emission is proportional to the mass-squared splitting $M_{\eta^\pm}^2 - M_{\eta^0}^2$ of the two intermediate scalars in the IB graph, and may even exceed the radiation rate of the transverse W_T [16].

2 Helicity and v^{2L} Suppressions Explored

The discrete symmetries C , P , and T , and angular momentum conservation, place constraints on fermion pair $\bar{\chi}\chi$ states [6, 11], as summarized in Table (1). A fermion pair can have total spin $S = 1$ in the symmetric state or $S = 0$ in the antisymmetric state. The parity of the two-fermion state is $P = (-)^{L+1}$, where L is the orbital angular momentum of the pair. Since the annihilation rate is suppressed as v^{2L} , only the $L = 0$ partial wave gives an unsuppressed rate for DM in today's Universe. The general rule for charge-conjugation (particle-antiparticle exchange) is $C = (-)^{L+S}$. Two identical fermions comprise a Majorana pair, and so a Majorana pair is even under charge-conjugation, and from $C = (-)^{L+S}$ one infers that L and S must be either both even, or both odd for a Majorana pair.

Consider the $L \leq 2$ states. In spectroscopic notation $(^{2S+1})L_J$ and spin-parity notation (J^{PC}), the vector 3S_1 (1^{--}), C -odd axial vector 1P_1 (1^{+-}), and assorted 3D_J (J^{--}) states are all C -odd and therefore disallowed for a Majorana pair. The pseudo-scalar 1S_0 (0^{-+}), scalar 3P_0 (0^{++}), axial vector 3P_1 (1^{++}), C -even tensor 3P_2 (2^{++}), and pseudo-tensor 1D_2 (2^{-+}) are all C -even and therefore

allowed. In particular, the sole $L = 0$ state, with no v^{2L} suppression, is the pseudo-scalar 1S_0 (0^{-+}).

At threshold, defined by $s = (2M_\chi)^2$ or $v = 0$, the orbital angular momentum L is necessarily zero. With two identical Majorana fermions, the two-particle wave function must be antisymmetric under particle interchange. Since $L = 0$ at threshold, the $\chi\chi$ spatial wave function is even, and the wave function must be antisymmetrized in its spin. The antisymmetric spin wave function is the $S = 0$ state. Thus, the only contributing partial wave at threshold is the 1S_0 state. We have just seen that this is also the only state with no v^{2L} suppression, so one may expect an unsuppressed Majorana annihilation rate at threshold if and only if there is a 1S_0 partial wave.

Finally, one may also invoke CP invariance, the rule $CP = (-)^{S+1}$, and the fact that $S = 0$ and 1 are the only possibilities for a pair of spin 1/2 particles, to deduce that total spin S is conserved in any Dirac or Majorana annihilation to an $f\bar{f}$ final state.

What does this mean for a two-Majorana initial state which annihilates to a two-fermion final state? The implications are best recognized after a Fierz transformation of the two mixed fermion bilinears to ‘‘charge-retention’’ order, i.e., to a χ -bilinear and an f -bilinear (many details may be found in an appendix of [11]). The Majorana pair couples only to the C -even basis fermion bilinears: the pseudo-scalar, the scalar, and the axial vector. If the Fierz'd bilinears contain a pseudo-scalar, there is no suppression of the rate. Otherwise, there is a v^2 rate suppression. If the Fierz'd bilinears contain an axial vector piece, it offers a $(m_f/M_\chi)^2$ -suppressed s -wave contribution (unless accompanied by a radiated W or Z or γ , as we shall show).¹

To address the question of which products of Fierz'd currents are suppressed and which are not, we set v^2 to zero in the χ -current, and m_f^2 to zero in the fermion current. In Table 2 we give the results for the product of all standard Dirac bilinears. (The derivation of these results is outlined in an appendix of [11].) Read across rows of this table to discover that the only unsuppressed s -channel products of bilinears for the $2 \rightarrow 2$ process are those of the pseudo-scalar, vector, and tensor. For Majorana DM, the vector and tensor bilinears are disallowed by charge-conjugation arguments and one is left with just the unsuppressed pseudo-scalar. However, the couplings of chiral fermions to scalar η contains no pseudoscalar (or scalar) in its Fierz transformation.

Reference to Table (2) reveals that the only component of the axial current which is non-vanishing in the s -wave

1. The axial-vector is an $L = 1$ mode, and we have seen that this mode elicits a v^2 suppression in the rate. However, the exchange particle is typically off-shell and so has an $L = 0$ timelike pseudo-scalar piece, with unsuppressed velocity-dependence, but still with $(m_f/M_\chi)^2$ helicity-suppression.

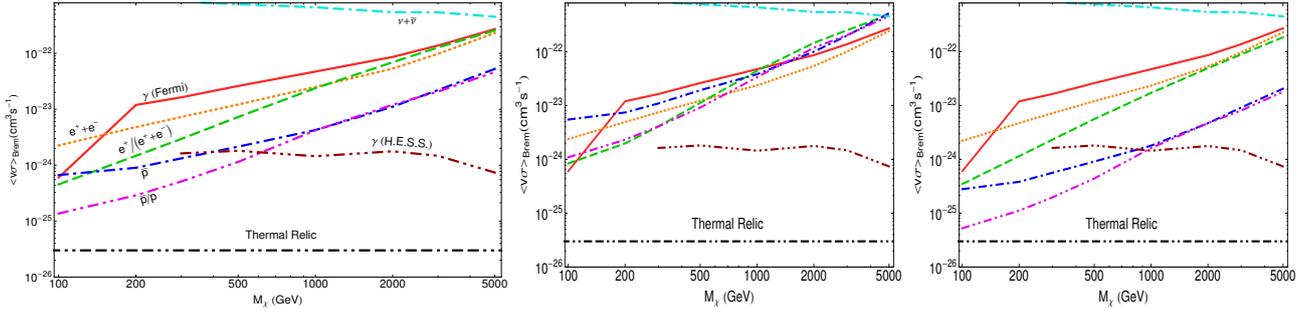


Figure 2: Upper limits on $\langle \nu \sigma \rangle_{\text{Brem}}$ using (left panel) the ‘med’ diffusion parameter set. Shown are constraints based on the Fermi extragalactic background light (solid, red), $e^+ + e^-$ flux (dots, orange), $e^+/(e^+ + e^-)$ ratio (dashes, green), \bar{p} flux (dot-dashes, blue), \bar{p}/p ratio (dot-dot-dashes, magenta), HESS γ 's (dot-dot-dot-dashes, maroon), and ν 's (dot-dash-dashes, cyan). Also shown is the expected cross section for thermal relic dark matter, $3 \times 10^{-26} \text{ cm}^3/\text{s}$. As for left panel, but using the ‘min’ (center) and ‘max’ (right) diffusion parameter sets.

L	S	$P = (-)^{L+1}$	$C = (-)^{L+S}$	$2S+1L_J$	J^{PC}	Name	Dirac Op	ν^{2L}
C-even states								
0	0	–	+	1S_0	0^{-+}	pseudo-scalar	$i\gamma_5$	ν^0
1	1	+	+	3P_0	0^{++}	scalar	1	ν^2
1	1	+	+	3P_1	1^{++}	axial-vector	$\gamma_5 \gamma^k$	ν^2
0	0	–	+	1S_0	0^{-+}		$\gamma_5 \gamma^0$	ν^0
C-odd states (unavailable for Majorana pair)								
0	1	–	–	3S_1	1^{--}	vector	γ^k	ν^0
1	0	+	–	1P_1	1^{+-}		γ^0	ν^2
1	0	+	–	1P_1	1^{+-}	tensor	σ^{jk}	ν^2
0	1	–	–	3S_1	1^{--}		σ^{0k}	ν^0
Redundant C-even states unavailable for Majorana pair								
0	0	–	+	1S_0	0^{-+}	pseudo-tensor	$\gamma_5 \sigma^{jk}$	ν^0
1	1	+	+	3P_1	1^{++}		$\gamma_5 \sigma^{0k}$	ν^2

Table 1: Decomposition of fermion bilinear currents into s -channel partial waves. Note that the relation $\gamma_5 \sigma^{\mu\nu} = \frac{i}{2} \epsilon^{\mu\nu\alpha\beta} \sigma^{\alpha\beta}$, implies that (i) the pseudo-tensor does not couple to Majorana fermions, and that (ii) $P(\gamma_5 \sigma^{\mu\nu}) = P(\sigma^{\mu\nu})$, and $C(\gamma_5 \sigma^{\mu\nu}) = -C(\sigma^{\mu\nu})$.

($\nu = 0$ limit) is $\bar{\chi} \gamma_5 \gamma^0 \chi$. However, there is no corresponding non-vanishing current $\bar{\Psi} \gamma_5 \gamma^0 \Psi$ or $\bar{\Psi} \gamma^0 \Psi$ in the Table. Thus, the s -wave amplitude must be helicity suppressed. Table (2) also reveals that a coupling of $\bar{\chi} \gamma_5 \gamma^0 \chi$ to $\bar{\Psi} \gamma_5 \gamma^i \Psi$ or $\bar{\Psi} \gamma^i \Psi$ is helicity un-suppressed, but requires a spin flip from parallel spinors in the initial state to antiparallel spinors in the final state. Such a direct coupling would violate Lorentz invariance. To the rescue comes gauge boson emission, which alters the fermion-pair spin state by one unit of helicity, and couples $\bar{\chi} \gamma_5 \gamma^0 \chi$ to $\bar{\Psi} \gamma_5 \gamma^i \Psi$ and $\bar{\Psi} \gamma^i \Psi$. An un-suppressed s -wave amplitude will be the result.

We now explain why the t - or u -channel scalar exchange with opposite fermion chiralities at the vertices is so common. It follows from a single popular assumption, namely that the dark matter is a gauge-singlet Majorana fermion. As a consequence of this assumption, annihilation to SM fermions, which are $SU(2)$ doublets or singlets, requires either an s -channel singlet boson, or a t - or u -channel singlet or doublet scalar that couples to χ - f . In the first instance, there is no symmetry to forbid a new force between SM fermions, a disfavored possibility. In the second instance, unitarity fixes the second vertex as the hermitian adjoint of the first. One gets chiral-opposites for the two vertices

of the t - or u -channel, and after a Fierz transformation, no unsuppressed pseudoscalar term. So either the fermion bilinear is suppressed by m_f in the s -wave or the χ bilinear is suppressed by ν in the p -wave.

3 W/Z -strahlung and Cosmic Signatures

Recent data from several experiments reveal an excess of an astrophysical positron flux at energies up to $\mathcal{O}(\text{TeV})$, while no excess of antiprotons is seen. In fact, the observed \bar{p} flux is well reproduced by standard astrophysical processes. While new astrophysical sources are thought to ultimately be the mechanism behind this excess (see, e.g., [17]), DM annihilation in the Galactic halo has been advanced as an alternative explanation. Many of the popular models proposed to explain the positron excess invoke Majorana DM.

Ref. [13] presented the spectra of stable annihilation products produced via DM annihilation including $\gamma/W/Z$ -bremsstrahlung. After accounting for propagation through the Galaxy, we set upper bounds on the annihilation cross section via a comparison with observational data. Fig. 2 collects our upper limits on the bremsstrahlung rate $\langle \nu \sigma \rangle_{\text{Brem}}$.

s-channel bilinear $\bar{\Psi}\Gamma_D\Psi$		v = 0 limit (projects out pure s-wave)		M = 0 limit (reveals helicity suppression)	
		parallel spinors	antiparallel spinors	parallel spinors	antiparallel spinors
scalar	$\bar{\Psi}\Psi$	0	0	\sqrt{s}	0
pseudo-scalar	$\bar{\Psi}i\gamma_5\Psi$	$-2iM$	0	$-i\sqrt{s}$	0
axial-vector	$\bar{\Psi}\gamma_5\gamma^0\Psi$	$2M$	0	0	0
	$\bar{\Psi}\gamma_5\gamma^j\Psi$	0	0	0	$\sqrt{s}(\pm\delta_{j1} - i\delta_{j2})$
vector	$\bar{\Psi}\gamma^0\Psi$	0	0	0	0
	$\bar{\Psi}\gamma^j\Psi$	$\mp 2M\delta_{j3}$	$-2M(\delta_{j1} \mp i\delta_{j2})$	0	$-\sqrt{s}(\delta_{j1} \mp i\delta_{j2})$
tensor	$\bar{\Psi}\sigma^{0j}\Psi$	$\mp 2iM\delta_{j3}$	$-2iM(\delta_{j1} \pm \delta_{j2})$	$-i\sqrt{s}\delta_{j3}$	0
	$\bar{\Psi}\sigma^{jk}\Psi$	0	0	$\pm\sqrt{s}\delta_{j1}\delta_{k2}$	0
pseudo-tensor	$\bar{\Psi}\gamma_5\sigma^{0j}\Psi$	0	0	$\pm i\sqrt{s}\delta_{j3}$	0
	$\bar{\Psi}\gamma_5\sigma^{jk}\Psi$	$\mp 2M\delta_{j1}\delta_{k2}$	$-2M(\delta_{j2}\delta_{k3} \mp i\delta_{j3}\delta_{k1})$	$-\sqrt{s}\delta_{j1}\delta_{k2}$	0

Table 2: Extreme non-relativistic and extreme relativistic limits for s-channel fermion bilinears. For a term with an initial-state DM bilinear and a final-state SM bilinear to remain unsuppressed, the DM bilinear must have a nonzero entry in the appropriate cell of the “v = 0 limit” columns, and the SM bilinear must have a non-zero term in the appropriate cell of the “M = 0 limit” columns. Recall that antiparallel spinors correspond to parallel particle spins (and antiparallel particle helicities for the $M = 0$ current), and vice versa. Amplitudes are shown for $\bar{u}\Gamma_D v = [\bar{v}\Gamma_D u]^*$. The two-fold \pm ambiguities reflect the two-fold spin assignments for parallel spins, and separately, for antiparallel spins.

While our analysis techniques are conservative, there are large astrophysical uncertainties in the propagation of charged particles through galactic magnetic fields, and in the DM density profile which probably contains substructure. Consequently, our constraints are illustrative, but not robust. We assumed astrophysical propagation parameters [18] which are consistent with a ‘median’ \bar{p} flux [19]. However, by assuming alternate parameters, e.g. from the ‘max’ or ‘min’ \bar{p} flux scenarios, our results may be strengthened or weakened by up to an order of magnitude, as shown in panels of Fig. 2. Our conclusions hold in all cases considered except for the extreme “min” choice.

For the “med” parameter set, the constraint from the antiproton ratio is stronger than that from the positron data by a factor of ~ 5 . Therefore, if the observed positron fraction were attributed to the bremsstrahlung process, then the same process would overproduce antiprotons by about a factor of five, and thereby preclude a sizable Majorana fermion DM contribution.

4 Conclusions

If DM is Majorana in nature, then its $2 \rightarrow 2$ annihilation to SM fermions is suppressed due to helicity considerations. However, both electroweak and photon bremsstrahlung lift this suppression, thereby becoming the dominant channels for DM annihilation (EW exceeding EM if the DM mass exceeds ~ 150 GeV).

Unsuppressed production and subsequent decay of the emitted W and Z gauge bosons will produce fluxes of hadrons, including \bar{p} 's, in addition to e^- 's, e^+ 's, ν 's, and γ 's. Importantly, we find that the null \bar{p} data make it difficult for helicity-suppressed Majorana DM annihilation to two fermions to source the reported cosmic e^+ excesses. The obstacle is the copious production and subsequent hadronic decay of the EW gauge bosons, which leads to a significant antiproton flux.

Acknowledgment: This work was partially supported by U.S. DOE award DE-FG05-85ER40226.

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