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# Minimal asymmetric dark matter

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ABSTRACT

In the early Universe, any particle carrying a conserved quantum number and in chemical equilibrium with the thermal bath will unavoidably inherit a particle–antiparticle asymmetry. A new particle of this type, if stable, would represent a candidate for asymmetric dark matter (DM) with an asymmetry directly related to the baryon asymmetry. We study this possibility for a minimal DM sector constituted by just one (generic) $SU(2)_L$ multiplet $\chi$ carrying hypercharge, assuming that at temperatures above the electroweak phase transition an effective operator enforces chemical equilibrium between $\chi$ and the Higgs boson. We argue that limits from DM direct detection searches severely constrain this scenario, leaving as the only possibilities scalar or fermion multiplets with hypercharge $y = 1$, preferentially quintuplets or larger $SU(2)$ representations, and with a mass in the few TeV range.



## 1. Introduction

The existence of dark matter (DM) is a well established fact, confirmed by a plethora of observations including the most recent cosmological surveys [1]. However, so far all evidences for DM come solely from gravitational effects, and its nature remains yet to be understood. If DM is constituted by new fundamental particles, the most compelling question is perhaps which other types of interactions these particles can have with ordinary matter, which could allow its 'discovery' via non-gravitational effects. But the little we know about DM brings about other puzzles, and one of the most intriguing ones is *why is the DM energy density so close to the energy density of baryons:* $\Omega_{DM}/\Omega_B \approx 5.5$ [1]?

In recent years, numerous models and constructions have been put forward in the attempt to explain this puzzle. Two main classes of models have been studied in the literature so far: asymmetric DM (ADM) [2–10] with all its variants, and WIMP-based schemes, as for example the ones proposed in [11,12] (see [13–16] for recent reviews). These constructions usually rely on new symmetries (for instance, in order to transfer the asymmetry) and/or extended hidden sectors. It should be remarked, however, that symmetries can just explain why the *number densities* are comparable: $n_{DM}/n_B \approx \mathcal{O}(1)$, while the numerical coincidence is in the *energy densities:* $\rho_{DM}/\rho_B \equiv (m_{DM}n_{DM})/(m_N n_B) \approx \mathcal{O}(1)$. In most cases a suitable value for $m_{DM}$ is chosen in order to reproduce the observations, which means that the coincidence is not really explained. Models in which an explanation is provided for the ratio of energy densities do exist, but often rely on unusual scenarios [17,18].

In this paper we investigate whether it is possible to relate the baryon and dark matter number densities using just the gauge symmetries of the standard model (SM). Our framework assumes a minimal ADM (MADM) sector but is otherwise fairly model-independent. We assume that at temperatures well above the temperature $T_{EW}$ of the electroweak (EW) phase transition, a CP asymmetry is generated in the thermal bath (the origin of this asymmetry is not relevant for us). At sufficiently low temperatures ($T \lesssim 10^6$ GeV) the rates of all SM interactions become faster than the Universe expansion rate, and chemical equilibrium is enforced among all SM particle species, that are thus characterized by numerically similar density asymmetries. We introduce in this scenario a new $SU(2)_L$ multiplet $\chi$ carrying hypercharge, whose lightest (neutral) component is rendered stable by a matter parity. An effective operator ensures that at $T \gtrsim T_{EW}$ $\chi$ is in chemical equilibrium with the Higgs multiplet, and thus it inherits an asymmetry which, after the symmetric component has annihilated away, is at the origin of its present relic density. We will show that limits from DM searches via direct detection (DD) experiments, together with the requirement that the effective interaction goes out of equilibrium before hypercharge symmetry gets spontaneously broken, render this scenario quite constrained. We find that the only viable MADM candidates are fermion or scalar multiplets with

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hypercharge $y = 1$.[1] An important difference with respect to other ADM scenarios is that in our case, while the DM relic density is indeed inherited from an initial asymmetry, DM is no more asymmetric in the present cosmological era. This is because when the Higgs field acquires a vacuum expectation value (vev) and hypercharge symmetry gets broken, the same operator responsible for the asymmetry transfer generates a splitting between the two real degrees of freedom $\chi^0_{1,2}$ of the neutral component of the complex multiplet. DM corresponds to the lightest state $\chi^0_1$ (a real scalar or a Majorana fermion) which can well undergo self annihilation and produce indirect detection signals. This is clearly different from the cases in which the present-day DM population is still characterized by an asymmetry in a conserved quantum number, and no signal from DM annihilation is expected.[2]

## 2. Minimal asymmetric dark matter

The particle content of our DM scenario is that of the SM augmented with an $SU(2)_L$ multiplet containing a neutral component which accounts for the DM. In this respect it might resemble the minimal DM (MDM) scenario proposed in [21,22]. However, while MDM considers self-conjugate multiplets with zero hypercharge, we require non-zero hypercharge to ensure that the DM multiplets are not self-conjugate, and can thus carry a particle–antiparticle asymmetry. This implies that the phenomenology of MADM is genuinely different from that of MDM.

As usual, in order to enforce DM stability, we need to impose a parity symmetry under which $\chi$ is the only odd field. A second important requirement is DM neutrality. An $SU(2)_L$ multiplet $\chi$ of weak isospin $t$ and hypercharge $y$ (without loss of generality we take $y$ to be positive) can contain an electrically neutral component if $t = y + k$, with $k$ a non-negative integer.[3] For the minimal case $k = 0$, for which the multiplet has the lowest dimension, the electrically neutral component corresponds to the lowest weight $t_3 = -y$, while for non-minimal multiplets with $k > 0$ the lowest weights are negatively charged. In all cases we need to ensure the neutral component remains the lightest one within the multiplet. A mass splitting between the charged and neutral component of $\chi$ can be generated after EW symmetry breaking by any type of $\chi$ couplings to the Higgs involving $\chi$ bilinears that are not by themselves invariant under $SU(2)_L$. For scalars there is always such a renormalizable operator:

$$\mathcal{O}^{\vec{t}} = \lambda_v \left(\chi^\dagger \vec{t}\,\chi\right)\left(\phi^\dagger \frac{\vec{\tau}}{2}\phi\right), \tag{1}$$

where $\vec{t}$ are the $SU(2)$ matrices in the representation in which $\chi$ transforms and $\vec{\tau}$ are the Pauli matrices. After EW symmetry breaking $\mathcal{O}^{\vec{t}}$ induces a mass difference between two $\chi$ components of isospin eigenvalues $t_3$ and $t'_3$ given by:

$$\delta m^v = -(t_3 - t'_3)\frac{\lambda_v v^2}{4m_\chi} \approx -151\,(t_3 - t'_3)\,\lambda_{0.02}^{1\,\mathrm{TeV}}\ \mathrm{MeV}, \tag{2}$$

where $v = \langle\phi\rangle = (2\sqrt{2}G_F)^{-1/2} \approx 174$ GeV is the Higgs vev and we have defined $\lambda_{0.02}^{1\,\mathrm{TeV}} = \frac{\lambda_v}{0.02}\,\frac{1\,\mathrm{TeV}}{m_\chi}$. For fermions the operator corresponding to Eq. (1) is of dimension five, and the result (2) holds with the replacement $1/m_\chi \to 1/\Lambda$. Neutral-charged mass splittings receive also contributions from gauge boson loops. We obtain:

$$\begin{aligned}\delta m^{\alpha_2} &= \frac{\alpha_2}{2}\,(t_3 - t'_3)\left\{(t_3 + t'_3)\left(M_W - c_W^2 M_Z\right) + 2y s_W^2 M_Z\right\}\\ &= 152\,(t_3 - t'_3)\,\{1.1(t_3 + t'_3) + 4.6y\}\ \mathrm{MeV},\end{aligned} \tag{3}$$

with $\alpha_2 = \frac{g^2}{4\pi}$ and $(c_W)\,s_W$ the (co)sine of the weak mixing angle. Eq. (3) agrees with the result given in [21] and holds for scalars as well. We see that for minimal multiplets (those with $t_3^{\min} = -y$) $\delta m^{\alpha_2}$ and (for $\lambda_v < 0$) also $\delta m^v$ shift the mass of the charged components above the mass of the neutral one (e.g. for a scalar triplet with $y = 1$ and reference values of the parameters, the mass splittings between the $Q = +1$ and $Q = 0$ components are $\delta m^{\alpha_2} \sim 540$ MeV and, for negative $\lambda_v$, $\delta m^v \sim 151$ MeV).

For non-minimal multiplets ($t_3^{\min} = -(y + k)$) states of weight $-(y + l)$ with $1 \le l \le k$ are all (negatively) charged. Among them, loop corrections would make heavier than the $t_3 = -y$ neutral state only those with $l > 2.3\,y$. Since $y = 1/2$ is the minimum hypercharge value allowing for a neutral component in the multiplet, and since by definition a charged state with $l = 1$ is always present, if loop induced mass splittings were dominant, all non-minimal multiplets would remain excluded as DM candidates. However, including the tree level contribution $\delta m^v$ allows to evade this conclusion. We find that the neutral state is always the lightest one for positive values of $\lambda_{0.02}^{1\,\mathrm{TeV}}$ falling within the interval:

$$\lambda_{0.02}^{1\,\mathrm{TeV}} = 2.5y \pm 1.1\,. \tag{4}$$

Let us now assume that a (non-Hermitian) effective operator of dimension $d \ge 4$ mediates an interaction between a pair of $\chi$ particles and the Higgs field $\phi$. Since the hypercharge of the Higgs is $y(\phi) = -1/2$ this operator takes the form

$$\mathcal{O}^\phi = \frac{1}{\Lambda^{4y-x}}\,\chi\chi\phi^{4y}\,, \tag{5}$$

where $x = 1\,(2)$ if $\chi$ is a fermion (boson), $y = y(\chi)$ is the hypercharge of the $\chi$ particle, and $\Lambda$ is the scale where the effective operator is generated.[4] The operator $\mathcal{O}^\phi$ plays two roles:

- At $T > T_{EW}$: $\mathcal{O}^\phi$ can enforce chemical equilibrium between $\phi$ and $\chi$, communicating the asymmetry present in the thermal bath to the DM sector.
- At $T < T_{EW}$: $\mathcal{O}^\phi$ generates a mass splitting between the two real degrees of freedom $\chi^0_{1,2}$ of the neutral component of the complex multiplet:

$$\delta m_0^x = \frac{v^{4y}}{\Lambda^{4y-x}} \tag{6}$$

where for fermions ($x = 1$) $\delta m_0 \equiv m_{\chi_2^0} - m_{\chi_1^0}$ while for bosons ($x = 2$) $\delta m_0^2 \equiv m^2_{\chi_2^0} - m^2_{\chi_1^0} \approx 2m_\chi\,\delta m_0$.

Let us comment on the previous two points. For definiteness, in the first point we have assumed that some baryogenesis mechanism generates an asymmetry in the SM sector, which is then communicated to the $\chi$ sector via the operator $\mathcal{O}^\phi$. We stress however, that the opposite possibility is also viable. The main difference would simply be that the fundamental asymmetry is no

[1] Subleading contributions from an asymmetric DM component to $\Omega_{DM}$ are however possible also in other cases.

[2] Even in this case, if the symmetric component has been only partially annihilated away, indirect detection signals, although accordingly suppressed, might still be detectable, see e.g. [19,20].

[3] Electric charge is defined here as $Q = t_3 + y$ with $t_3$ the diagonal generator of weak isospin.

[4] In Eq. (5) we have implicitly absorbed in the scale $\Lambda$ an overall coupling $\lambda$ multiplying the effective operator. Of course, any bound derived on $\Lambda$ should be then understood as a bound on $\Lambda/\lambda^{1/(4y-x)}$. For the fermion doublet case ($x = 1$, $y = 1/2$) this operator was already used in [23] to relate DM and the baryon asymmetry.

more in the SM $B-L$ charge, that remains exactly conserved and with vanishing asymmetry, but in a global hypercharge asymmetry of the SM particles, which is exactly compensated by an equal in size and opposite in sign asymmetry in the $\chi$ sector [24,25].

Note that the requirement of gauge invariance allows to write other operators suitable to enforce chemical equilibrium between $\chi$ and the SM particles. Of course, operators of higher dimension are not relevant and can be neglected, however, for integer $y$'s, the operator

$$\mathcal{O}^{e_R} = \frac{1}{\Lambda^{3y-x}} \chi\chi (e_R e_R)^y \,, \tag{7}$$

where $e_R$ is any of the SM $SU(2)_L$ singlet leptons (with $y(e_R) = -1$) is allowed, and its dimension is $y$ units lower than the dimension of $\mathcal{O}^\phi$ in Eq. (5). Motivated by minimality, one could assume that the ultra-violet realization of the model is such that operators of this type are either forbidden, or that they are suppressed by additional powers of $\Lambda$ with respect to naive power counting. However, for completeness, in Section 2.2 we will briefly comment on the effects of $\mathcal{O}^{e_R}$ in the case of $y=1$ multiplets, which include the interesting cases of fermion and scalar triplets.

The second role played by $\mathcal{O}^\phi$ after EW symmetry breaking is also of fundamental importance: the lightest new particle $\chi_1^0$ (a real scalar or a Majorana fermion) does not couple to the $Z$ boson, but virtual $Z$ exchange can mediate the inelastic transition $\chi_1^0 \to \chi_2^0$. In order to evade the stringent limits imposed by direct searches for DM scatterings off nuclei, we need to ensure that in most cases the kinetic energy of the incoming DM particle will not suffice to trigger the inelastic scattering, so that the rate of events gets kinematically suppressed below the observable level. This implies a lower limit on the mass splitting:

$$\delta m_0 = 2m_\chi \left(\frac{v}{\Lambda}\right)^{4y} \left(\frac{\Lambda}{2m_\chi}\right)^x \gtrsim \delta m^{\min} \,. \tag{8}$$

Values of $\delta m^{\min}$ have been derived in [26] for different DM masses and different hypercharges $y$. In the DM mass range relevant for us they can be roughly parameterized as $\delta m^{\min} \sim (1+0.2y)\times 175$ keV for $m_\chi$ of order few TeV.

In order for DM to originate from the asymmetry present in the primordial plasma, the following steps are required to occur in sequential order of decreasing temperature:

1. At some temperature $T \gg T_{EW}$ the effective operator $\mathcal{O}^\phi$ mediates in-equilibrium reactions feeding an asymmetry between the SM sector and the $\chi$ sector.
2. At a certain temperature $T_a > T_{EW}$ the rate of these reactions drops below the Hubble rate $H$, and the $\chi$ sector gets chemically decoupled from the thermal bath. The relevant effective Lagrangian at $T_a$ is then characterized by a global $U(1)_\chi$ symmetry corresponding to rephasing of the $\chi$ field. The quantity $Y_{\Delta\chi} \equiv Y_\chi - Y_{\bar\chi}$ (where $Y_\chi = n_\chi/s$, with $s$ the entropy density) is associated to the $U(1)_\chi$ global charge, and it remains conserved.
3. The annihilation $\chi\bar\chi \to SM$ that proceeds, for example, via (unsuppressed) gauge interactions, continues to erase the symmetric DM component until it freezes out at a temperature $T_s < T_a$. After $U(1)$ hypercharge symmetry is spontaneously broken at $T_{EW}$ no conserved charge remains associated with the $\chi$ neutral members. To avoid that the surviving ADM component will restart annihilating away via e.g. $\chi_1^0\chi_1^0 \to W^+W^-(ZZ)$ mediated by $t$-channel exchange of $\chi^\pm$ ($\chi_2^0$), we need to require $T_s > T_{EW}$. If at $T_s$ $Y_{\bar\chi} \ll Y_{\Delta\chi} \approx Y_\chi$, then the present DM relic abundance is dominated by the initial $\chi$ asymmetry.
4. At some temperature $T_d < T_{EW}$, which depends on $m_\chi$ and on the charged/neutral mass splitting $\delta m_\chi$, $\chi^\pm$ will decay to the lighter neutral states. Later on (but still safely before Big Bang Nucleosynthesis), also $\chi_2^0 \to \chi_1^0$ decays occur. Eventually, at $T \ll T_d$ we will have $Y_{\chi_1^0} = Y_{\Delta\chi}$ and the present DM energy density then is given by $\rho_{DM} = s\, m_\chi Y_{\Delta\chi}$.

Let us note that if the annihilation of the symmetric part $\chi\bar\chi \to SM$ proceeds mainly via gauge interactions, freeze out occurs around $T_s \sim m_\chi/25$. The requirement $T_s > T_{EW}$ (point 3.) then implies $m_\chi \gtrsim 25\, T_{EW}$.[5] Estimating precisely the SM value of $T_{EW}$ is a hard task, and for relatively large Higgs masses $> 100$ GeV only a few studies exist [27–29]. In particular, for a Higgs mass $\sim 125$ GeV, Ref. [29] quotes $T_{EW} \sim 130$ GeV. Due to the large uncertainties involved in these estimates we will conservatively impose the condition $T_{EW} > 100$ GeV which yields the lower limit $m_\chi \gtrsim 2.5$ TeV. As regards the freeze out of the interactions mediated by the effective operator $\mathcal{O}^\phi$, we will take it to be $T_a \sim m_\chi/10 > T_{EW}$. This value results in a Boltzmann suppression that yields a MADM relic asymmetry in the ballpark to account for $\Omega_{DM}$.

### 2.1. Constraints from chemical decoupling

We now discuss, in a general way, the conditions under which $\chi$ can provide a DM candidate with a relic density originating from the same primordial asymmetry giving rise to the cosmological matter/antimatter asymmetry.

The operator $\mathcal{O}^\phi$ in Eq. (5) induces two types of reactions which can maintain $\chi$ in chemical equilibrium with the thermal bath: $s$-channel annihilation $\chi\chi \leftrightarrow \phi^{4y}$, and inelastic $t$-channel scattering[6] $\chi\phi^* \leftrightarrow \chi^*\phi^{4y-1}$. We define $T_a$ as the temperature at which chemical equilibrium cannot be any longer maintained, which happens when the rates for both these reactions

$$\Gamma_{\chi\chi} = n_\chi^0 \langle\sigma|v|\rangle_{\chi\chi} \,, \tag{9}$$

$$\Gamma_{\chi\phi} = n_\phi^0 \langle\sigma|v|\rangle_{\chi\phi} \,, \tag{10}$$

become slower than the Hubble expansion rate:

$$\Gamma_{\chi\chi},\ \Gamma_{\chi\phi} \lesssim H(T_a) \,. \tag{11}$$

After decoupling, the relic abundance of DM remains approximately fixed. Chemical decoupling of $\chi$ always occurs in the non-relativistic limit $T_a < m_\chi$ while the requirement $T_a > T_{EW}$ implies that a relativistic number density is the one appropriate for the Higgs boson. Thus, the appropriate equilibrium number densities for Eqs. (9)–(10) are:

$$n_\chi^0 = g_\chi \left(\frac{m_\chi T}{2\pi}\right)^{3/2} e^{-m_\chi/T} \,, \tag{12}$$

$$n_\phi^0 = g_\phi \frac{\zeta(3) T^3}{\pi^2} \,, \tag{13}$$

with $g_\chi$ and $g_\phi$ the respective numbers of degrees of freedom and $\zeta(3) \approx 1.2$. The thermally averaged cross sections for the two processes can be estimated as:

[5] For scalars, annihilation can also proceed via renormalizable operators like $\mathcal{O}^{\vec{t}}$ and $\lambda_s(\chi^\dagger\chi)(\phi^\dagger\phi)$. For particularly large couplings $\lambda_{v,s} > 1$ they could be dominant and yield $T_s < m_\chi/25$.

[6] We thank S. Tulin for pointing out to us the relevance of the $t$-channel reactions.

$$\langle\sigma|v|\rangle_{\chi\chi} \sim \eta_{PS}^{(n)}\, m_\chi^{-2}\left(\frac{m_\chi}{\Lambda}\right)^{2(4y-x)}, \tag{14}$$

$$\langle\sigma|v|\rangle_{\chi\phi} \sim \langle\sigma|v|\rangle_{\chi\chi}\left(\frac{T}{m_\chi}\right)^{4(2y-1)}, \tag{15}$$

where $\eta_{PS}^{(n)}$ is a $n=4y$ body phase space numerical factor. The temperature dependent multiplicative factor for the $t$-channel process (15) arises because while for $s$-channel annihilation the available phase space for the final states is determined by $m_\chi$, for the $t$-channel is determined by the $\phi^*$ momentum, which is of order $T$. We can now check by direct comparison which process, for the different cases, is the relevant one to maintain chemical equilibrium down to $T_a$. The condition $\Gamma_{\chi\chi} > \Gamma_{\chi\phi}$ is satisfied when

$$\frac{z}{\log z} \lesssim 4(2y-1)+\frac{3}{2}, \tag{16}$$

where we have defined $z = m_\chi/T$. For $y = 1/2$ this inequality is never satisfied, so that the relevant processes enforcing chemical equilibrium are the $t$-channel scatterings. For $y=1$ $\Gamma_{\chi\chi}$ dominates as long as $z \lesssim 15$. Since, as mentioned above, the correct DM relic density is obtained if chemical decoupling occurs around $z_a \sim 10$, for $y = 1$ $s$-channel annihilation is the most relevant process. Finally, for $y > 1$ the decoupling temperature is always determined by $\Gamma_{\chi\chi}$, and $t$-channel scatterings can be safely neglected. We thus need to consider separately the case $y = 1/2$ (scalar and fermion DM doublets belong to this class) from the cases with $y \geq 1$ (scalar and fermion triplets belong to this class). Let us start from the latter case.

To evaluate $\Gamma_{\chi\chi}$ let us first estimate the value of $n_\chi^0$ in Eq. (9) which would yield a correct DM relic density. Before the EWPT chemical equilibrium between the Higgs and the DM multiplet imposes the condition:

$$\frac{\Delta n_\chi}{n_\chi^0} = -2y\frac{\Delta n_\phi}{n_\phi^0}, \tag{17}$$

where $\Delta n_\chi = n_\chi - n_{\bar\chi}$ and $\Delta n_\phi = n_\phi - n_{\bar\phi}$, and the minus sign follows from requiring consistency between the hypercharge assignments $y(\chi) > 0$, $y(\phi) < 0$ and hypercharge conservation. By normalizing both asymmetries to the entropy density, Eq. (17) can be rewritten as:

$$\frac{Y_{\Delta\chi}}{Y_{\Delta\phi}} = -2y\frac{n_\chi^0}{n_\phi^0}. \tag{18}$$

Assuming the SM content of relativistic particles, the Higgs asymmetry is related to $\Delta_{B-L}$ by $Y_{\Delta\phi} = -\frac{8}{79}Y_{\Delta_{B-L}}$ [30]. We further have $Y_{\Delta_{B-L}} = \frac{79}{28}Y_{\Delta B}$ [31] so that:

$$\frac{Y_{\Delta\chi}}{Y_{\Delta\phi}} = -\frac{7}{2}\frac{Y_{\Delta\chi}}{Y_{\Delta B}} = -\frac{7}{2}\omega\frac{m_B}{m_\chi}, \tag{19}$$

where we have defined $\omega \equiv \frac{\Omega_{DM}}{\Omega_B}$ and $m_B \approx 1$ GeV is the nucleon mass. Putting together Eq. (18) and Eq. (19) we obtain:

$$n_\chi^0 = \frac{7\,\omega}{4y}\frac{m_B}{m_\chi}n_\phi^0. \tag{20}$$

By means of Eq. (14) and Eq. (20) the condition $\Gamma_{\chi\chi} \lesssim H(T_a)$ translates to:

$$m_\chi^{-2}\, z_a^{-1}\left(\frac{m_\chi}{\Lambda}\right)^{2(4y-x)} \lesssim \frac{4\pi^3}{21\zeta(3)}\sqrt{\frac{\pi g_*}{5}}\frac{y}{\omega\eta_{PS}^{(n)}}\frac{1}{M_P m_B} = 6.1\frac{y}{\eta_{PS}^{(n)}}\times 10^{-19}\ \text{GeV}^{-2} \tag{21}$$

where $M_P = 1.2\times 10^{19}$ GeV is the Planck mass, and we have used $H = \left(4\pi^3 g_*/45\right)^{1/2} T_a{}^2/M_P$ for the Hubble parameter with $g_* = 106.75$ the number of relativistic degrees of freedom, and $\omega \approx 5.5$ from cosmological observations [1]. In the numerical analysis we have adopted for the phase space factor the parametrization $\eta_{PS}^{(n)} = 1/[4\pi(3^3\times 2^{11}\times\pi^4)^{2y-1}]$ which reproduces correctly 3-body and 4-body phase space when particle multiplicities and identical particle final states are accounted for.

For $y = 1/2$ the condition $\Gamma_{\chi\phi} \lesssim H(T_a)$ yields instead:

$$m_\chi^{-1} z_a^{-1}\left(\frac{m_\chi}{\Lambda}\right)^{2(2-x)} < 5.9\times 10^{-16}\ \text{GeV}^{-1}. \tag{22}$$

*2.2. DM multiplets with different hypercharges*

For each value of the hypercharge $y(\chi)$, Eq. (8) and Eq. (21) (or Eq. (22) if $y = 1/2$) provide strong constraints on the viable parameter space. Another constraint that we will use follows from the requirement that the effective operator (5) provides a consistent description of the interaction enforcing chemical equilibrium, which requires $m_\chi < \Lambda$. Let us now study a few cases.

For a fermion multiplet ($x = 1$) with hypercharge $y = 1$ (the minimal choice is a complex $SU(2)$ triplet) the two constraints (8) and (21) yield:

$$\Lambda \lesssim \left(\frac{v^4}{\delta m^{\min}}\right)^{1/3} \approx 17\ \text{TeV}, \tag{23}$$

$$m_\chi \approx 10\left(\frac{\Lambda}{17\ \text{TeV}}\right)^{3/2}\left(\frac{z_a}{10}\right)^{1/4}\ \text{TeV}, \tag{24}$$

where in the first equation we have used $\delta m^{\min} \approx 200$ keV. With $z_a \sim 10$ and taking into account the limit on the cutoff scale (23) we obtain $m_\chi \lesssim 10$ TeV which, for $z_s \sim 25$, is completely compatible with the requirement $T_s \gtrsim T_{EW}$. Therefore a complex $SU(2)_L$ fermion multiplet with $y = 1$ can be a viable MADM candidate. The relatively low value of $\Lambda$ implies that dimension five operators yield a rather large charged/neutral tree level mass difference $\delta m^v \sim 1$ GeV which dominates over the loop contributions, and is also much larger than the splitting $\delta m^0$ between the two neutral states $\chi_{1,2}$. (In the presence of the transfer operator $\mathcal{O}^{e_R}$ Eq. (7) a $y = 1$ fermion multiplet would still be a viable MADM candidate within the narrower window 2.5 TeV $\lesssim m_\chi \lesssim$ 6.7 TeV.)

The results for fermions in our MADM scenario for hypercharges $y = 1, \frac{3}{2}, 2$ are depicted in Fig. 1 (corresponding to $\mathcal{O}^\phi$ operators respectively of dimension 7, 9, 11). The horizontal dashed line gives the lower limit on the freeze-out temperature for $\bar\chi\chi$ annihilation $T_s \sim m_\chi/25 > 100$ GeV and the gray region below is then excluded. The thick black line bisecting the figure selects the region $m_\chi > \Lambda$ (in gray) which must be excluded because the description of the asymmetry transfer via the effective operator $\mathcal{O}^\phi$ breaks down. The regions on the right of the three vertical lines corresponding respectively to $y = 1, \frac{3}{2}, 2$, delimit the values of $\Lambda$ that give a too large suppression of the $\chi_2^0 - \chi_1^0$ mass difference ($\delta m \lesssim 200$ keV) so that a signal would have been observed in DD experiments.

The results for $\chi$ contributing dominantly to the DM of the Universe are obtained from Eq. (21). The width of the band corresponds to varying the fraction of the relic abundance $f \equiv \Omega_\chi/\Omega_{DM}$ from 50% to 100%. As we have discussed above, for a fermion multiplet with $y = 1$ there is a region up to $m_\chi \approx 10$ TeV and $\Lambda \approx 17$ TeV in which the $\chi$ relic density generated via an initial $\chi$–$\bar\chi$ asymmetry can account for the dominant amount of DM, while respecting the other bounds. The case $y = \frac{3}{2}$ corresponds to the red band and $y = 2$ to the green band. In both cases the entire

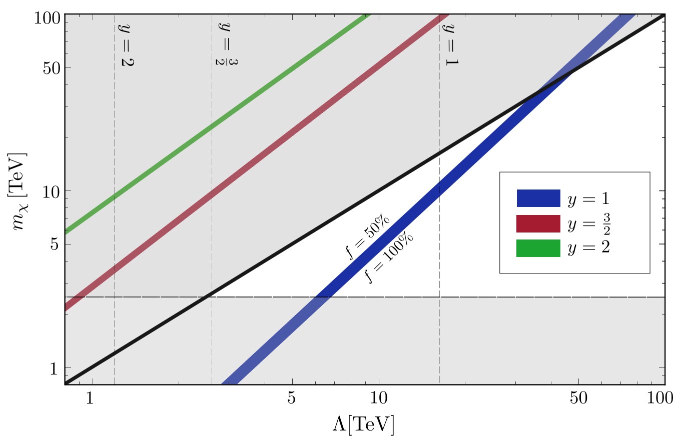


**Fig. 1.** Constraints on the parameter space $m_\chi$ *vs.* $\Lambda$ for fermion DM with hypercharge $y = 1$ (blue), $\frac{3}{2}$ (red) and 2 (green). The width of the bands correspond to varying $f = \Omega_\chi/\Omega_{DM}$ in the interval $0.5 < f < 1.0$. The region above the line bisecting the figure corresponds to $m_\chi > \Lambda$ for which the effective operator description (5) breaks down. This excludes the $y > 1$ lines. The region below the dashed line is excluded by the requirement that all relevant reactions freeze-out above $T_{EW}$. The regions at the right of the vertical lines labeled $y = 1, \frac{3}{2}, 2$ are excluded by DD experiments. (For interpretation of the references to color in this figure, the reader is referred to the web version of this article.)

parameter space selected lies in the $m_\chi > \Lambda$ half-plane, where the effective field theory description of the asymmetry transfer cannot be applied. A first conclusion is that the case of a fermion multiplet with $y = 1$ can be viable for a certain mass range, while for hypercharges $y > 1$ the MADM scenario is not viable, or more precisely the possible contribution of an asymmetry to the DM relic density cannot be relevant. For a fermion multiplet with hypercharge $y = 1/2$ (the minimal choice is an $SU(2)_L$ doublet) Eq. (8) should be used together with Eq. (22). The first condition yields the upper limit $\Lambda \lesssim 1.5 \times 10^5$ TeV, while Eq. (22) implies the lower bound $\Lambda \gtrsim 4.1 \times 10^5 \left(\frac{T_a}{100\text{ GeV}}\right)^{1/2}$ TeV. Given that $T_a > T_{EW} \gtrsim 100$ GeV the two bounds are in conflict, and we can conclude that for fermion doublets (and more generically for fermion multiplets with hypercharge $y = 1/2$) the MADM scenario is not viable.

For a scalar multiplet ($x = 2$) with hypercharge $y = 1$ the two constraints (8) and (21) yield:

$$\Lambda \lesssim \frac{v^2}{\left(2m_\chi \delta m^{\min}\right)^{1/2}}\,, \tag{25}$$

$$m_\chi \approx \frac{\Lambda^2}{m_B}\left(6.1 \times 10^{-14} z_a\right)^{1/2}. \tag{26}$$

The value of $m_\chi$ is maximized by saturating the inequality (25) in which case solving the system gives:

$$\Lambda \approx 18 \left(\frac{10}{z_a}\right)^{1/8} \text{ TeV}\,, \tag{27}$$

$$m_\chi \lesssim 6.7 \left(\frac{z_a}{10}\right)^{1/4} \text{ TeV}\,. \tag{28}$$

The last equation then allows for 2.5 TeV $\lesssim m_\chi \lesssim$ 6.7 TeV and shows that values of $m_\chi$ large enough to ensure that chemical equilibrium reactions and annihilation processes freeze out before $T_{EW}$ are possible in a rather large window. (Asymmetry equilibration via the transfer operator $\mathcal{O}^{e_R}$ Eq. (7) would instead imply $m_\chi \lesssim 1.3$ TeV, and would render the $y = 1$ scalar case not viable.) Fig. 2 depicts the results for scalar DM. Graphical conventions are the same as in Fig. 1. We see from the picture that the only case in which a DM asymmetry can give sizable contributions to $\Omega_{DM}$ is for $y = 1$. For higher values of the hypercharge the bands lie in the $m_\chi > \Lambda$ half-plane, and the corresponding MADM possibilities are therefore ruled out.

For a scalar multiplet with hypercharge $y = 1/2$ (e.g. a scalar doublet), the operator $\mathcal{O}^\phi$ is of dimension four (renormalizable). Thus there is no cutoff $\Lambda$ in the model and $m_\chi$ is the only new scale, a feature that is unique to this case. In order to pin down the values of $m_\chi$ it is convenient to keep explicit the coupling constant $\lambda$ of the transfer operator. The constraint from DD, Eq. (8), implies the upper bound $\frac{m_\chi}{\lambda} \lesssim 8 \times 10^4$ TeV, while the chemical freeze-out condition (22) yields the lower limit $\frac{m_\chi}{\lambda} \gtrsim 4.1 \times 10^5 \left(\frac{T_a}{100\text{ GeV}}\right)^{1/2}$ TeV.[7] The conflict between these two bounds leaves no window in parameter space where scalar doublets can work as MADM.

## 3. Mass limits from symmetric annihilation

We have seen in the previous sections that the bounds on the MADM parameter space from (i) limits on nucleon recoils signals via tree-level $Z$ boson exchange and (ii) constraints from the freeze-out temperature of asymmetry transfer and annihilation processes, select as the only possibilities multiplets with hypercharge $y = 1$. The minimal dimension of the corresponding representations are $SU(2)_L$ triplets, and the next-to-minimal are quintuplets. An important issue that should be discussed in more detail is which ranges of masses are allowed by the requirement that $\chi\bar{\chi}$ annihilation will be efficient enough to ensure that the contribution to $\Omega_{DM}$ of any surviving symmetric component remains subdominant, i.e. $\Omega_{\bar{\chi}} \ll \Omega_\chi \sim \Omega_{DM}$. Estimating the bounds on $m_\chi$ that follow from this argument is not a straightforward task, since for $m_\chi \gg M_W$ the annihilation cross section for $SU(2)_L$ multiplets is generically affected by non-perturbative Sommerfeld enhancements, which can result in a sizable suppression of the relic density. One of the most studied cases is that of an $SU(2)_L$ triplet with zero hypercharge, that is a wino-like DM, $\widetilde{W}$. With the tree level annihilation cross section, $\Omega_{\widetilde{W}} = \Omega_{DM}$ is obtained for $m_{\widetilde{W}} =$ 2.5 TeV [21]. More refined studies which include Sommerfeld and higher order corrections have found sizable enhancements of the annihilation rate, so that the condition $\Omega_{\widetilde{W}} = \Omega_{DM}$ is fulfilled for larger values of the mass. For example, Refs. [33–35] quote mass values in the range 2.7 TeV $\lesssim m_{\widetilde{W}} \lesssim$ 3.0 TeV, while more recent studies [36,37] give even higher values $m_{\widetilde{W}} \sim$ 3.1–3.2 TeV.[8] For a fermion triplet with $y = 1$ the tree level result quoted in [21] is $m_\chi \sim 1.9$ TeV, which is lower than in the $y = 0$ case because of the larger multiplicity of the complex multiplet. To our knowledge, no results have been reported in the literature for a $y = 1$ fermion triplet including Sommerfeld enhancements, however we would expect even larger effects than in the $y = 0$ case. This is because in the $T \gg T_{EW}$ limit the interaction range of the Sommerfeld potential is determined by the Debye screening length in the thermal plasma $\sim 1/(g_{1,2}T)$ (with $g_{1,2}$ the $U(1)_Y$ and $SU(2)_L$ couplings) rather than by the inverse gauge boson mass $1/M_W$. Although for $y \neq 0$ one expects that $SU(2)_L$ forces would result in non-perturbative corrections similar to the $y = 0$ case, the somewhat larger range of $U(1)_Y$ interactions can further enhance the effect. All in all, based on the results for the $y = 0$ case we make the educated guess that $\Omega_{DM}$ can be completely accounted for by a symmetric DM component in the mass range 2.7 TeV $\lesssim m_\chi \lesssim$ 2.8 TeV.

[7] We thank the authors of [32] for helping us in spotting a numerical error in the phase space factor for this case.

[8] It is worth remarking at this point that for $\chi\chi$ annihilation into $(\phi\phi)^{2y}$ the $U(1)_Y$ non-relativistic potential is repulsive, so that the rates for chemical equilibrating reactions will get suppressed rather than enhanced. This would raise the corresponding freeze-out temperature favoring the viability of the MADM scenario.

To the extent this is a reasonable estimate, then Fig. 1 shows that not much space is left for relevant contributions from the $\chi$–$\bar{\chi}$ asymmetry. For $y = 1$ scalar triplets similar arguments can be put forth, except that the lowest order result $m_\chi \sim 1.6$ TeV is a bit lower than in the fermion case, implying that the mass range in which an asymmetry could give relevant contributions to $\Omega_{DM}$ is accordingly reduced.

It was found in the previous section that for $y = 1/2$ multiplets the bounds from the two conditions (8) and (22) are in conflict: for fermions the lower bound on the cutoff scale ($\Lambda \gtrsim 4.1 \times 10^5 \left(\frac{T}{100\text{ GeV}}\right)^{1/2}$ TeV) is almost three times larger than the upper bound ($\Lambda \lesssim 1.5 \times 10^5$ TeV). For scalars the quantity $\frac{m_\chi}{\lambda}$ is bounded by the same lower limit, which is about five time larger than the upper limit ($\frac{m_\chi}{\lambda} \lesssim 8 \times 10^4$ TeV). The $y = 1/2$ cases of lowest dimension are, however, of particular interest since they correspond to a fermion doublet (similar to a pure Higgsino) and to a scalar doublet (similar to the scalar DM candidate of the inert doublet model [38]), and therefore it is worth checking if, in case the previous conflicts could be reconciled in some way, the $y = 1/2$ doublet MADM scenarios could become viable. For fermion doublets a tree level estimate of the $\chi$ mass that could account for $\Omega_{DM}$ via freeze-out of symmetric annihilation yields $m_\chi \sim 1.2$ TeV [21]. Non-perturbative corrections to this result have been found to be negligible [34]. Then the condition $m_\chi \gtrsim 2.5$ TeV that ensures that freeze-out of the relevant processes occur above $T_{EW}$ implies that the MADM relic density would largely overshoots the observed value of $\Omega_{DM}$. For a scalar doublet the value hinted by symmetric annihilation $m_\chi \sim 0.54$ TeV [21] is also not affected much by Sommerfeld corrections,[9] and the same conclusion holds. All in all, the results of the previous section together with considerations of the $m_\chi$ values needed to realize the condition $\Omega_\chi \approx \Omega_{DM}$ via symmetric annihilation, indicate that the MADM scenario cannot be relevant for fermions or scalars with $y = 1/2$.

The general conclusion is that among multiplets of minimal dimension, only $y = 1$ scalar/fermion triplets can marginally satisfy the condition of a sufficient suppression of the symmetric part of the relic density, so that the MADM scenario can become relevant. However, in the case of $y = 1$ multiplets of higher dimension (e.g. a quintuplet) the MADM scenario, subject to the constraint (4), becomes more easily viable. This is because the annihilation cross section for the symmetric component gets enhanced roughly as the fourth power of the multiplet dimension. This implies a strong suppression of the relic density, and correspondingly larger values of $m_\chi$ are required to saturate $\Omega_{DM}$ in the absence of an asymmetry. Moreover, for larger representations non-perturbative corrections to the annihilation processes become particularly important. As an example, it was found in Ref. [34] that for a fermion quintuplet with $y = 0$, $m_\chi \sim 4.4$ TeV obtained at tree level [21] gets boosted up to $m_\chi \sim 10$ TeV after the inclusion of Sommerfeld effects [34]. Therefore a thermally produced DM fermion quintuplet of mass $m_\chi \ll 10$ TeV could contribute the whole of DM only if its relic abundance is dominated by an initial asymmetry.

## 4. Other phenomenological implications

Let us finally discuss briefly other possible phenomenological implications of MADM candidates.

*Searches at colliders*

Searches at colliders of EW interacting new particles have been performed, but the current reach of LHC is only of a few hundred GeV [39–44] which does not constrain the interesting mass range for MADM. Future $e^+e^-$ colliders with an energy reach of $\sqrt{s} \sim 5$ TeV and $pp$ colliders with $\sqrt{s} \sim 100$ TeV will only marginally probe the multi TeV parameter space [45–47]. The chances that MADM particles will be ever produced at foreseeable colliders are thus rather feeble.

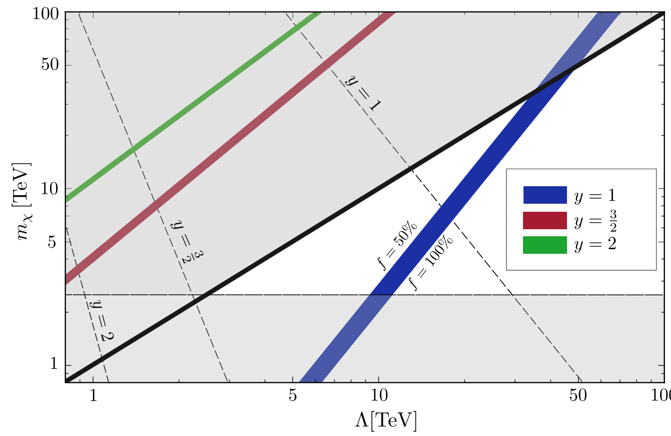


**Fig. 2.** Same as Fig. 1 for scalar DM.

*Direct detection experiments*

While MADM tree level $Z$ mediated interactions with nuclei are kinematically forbidden, non-vanishing DD cross sections appear at the loop level. However, an accidental cancellation among various contributions [48–51] results in suppressed cross sections $\sim \mathcal{O}(10^{-47})$ cm$^2$, which are by far below the current experimental bounds [52,53]. In the relevant mass range ($m_\chi \gtrsim$ TeV), the cross sections remain also below the reach of next generation DD experiments [54] and close to the neutrino scattering background.

*Indirect detection*

The possibilities to bound (or discover) MADM via indirect detection (ID) of signals from DM annihilation are more optimistic. While it is well known that any conclusion derived from searches of DM annihilation byproducts heavily depends on the DM halo model, large portions of the mass range remain ruled out also when adopting rather implausible profiles. The most relevant bounds come from cosmic-ray antiprotons measurements and from the absence of gamma-ray line features towards the galactic center. For example, for a $y = 0$ fermion triplet (wino-like DM) the corresponding bounds have been thoroughly studied, e.g. in [34,36,37], with the result that a mass range 1.8 TeV $\lesssim m_{\widetilde{W}} \lesssim 3.5$ TeV is excluded. There is a simple reason to expect that $y = 1$ MADM triplets could be also strongly disfavored by ID limits: the annihilation cross section gets a large enhancement from the formation of loose bound states when the range of the bounding interaction $\sim 1/M_W$ becomes of the order of the Bohr radius of the two particles state $\sim 1/(\alpha_2 m_\chi)$ [35], that is around $m_\chi \sim M_W/\alpha_2 \sim 2.4$ TeV, and we have seen that $y = 1$ triplets have the allowed values of $m_\chi$ rather close to this region. The same conclusion does not apply, however, to non-minimal $y = 1$ multiplets (e.g. quintuplets) for which the allowed mass range extends to $m_\chi \gg 2.5$ TeV. The issue of reliable MADM bounds from indirect detection for fermion and scalar quintuplets clearly deserves a specific study.

## 5. Conclusions

Our study started from the observation that any new $SU(2)_L$ multiplet carrying non-vanishing hypercharge and in chemical

[9] Larger values of $m_\chi$ are possible if annihilation into Higgs scalars largely dominates over annihilation mediated by gauge bosons.

equilibrium with the thermal bath at $T > T_{EW}$, will be unavoidably characterized by an asymmetry in its number density. We have assumed that some matter parity renders the lightest member of this multiplet stable, thus providing a candidate for DM. We have imposed a requirement of *minimality*, that is that no other new particle is introduced to help evading phenomenological constraints. We have also explored under which conditions the present-day relic abundance of such a DM candidate can be (mostly) determined by its initial asymmetry, which would justify denoting it as MADM.

A first set of constraints comes from limits from DD experiments, which exclude DM candidates interacting via (unsuppressed) tree-level $Z$ boson exchange. We have seen that this requirement can be satisfied by our MADM candidates in a minimal way: a single effective operator coupling DM to the Higgs field can in fact first be responsible (at $T > T_{EW}$) of enforcing chemical equilibrium between DM and the thermal bath, and next it can ensure that after EW symmetry breaking the lightest mass eigenstate corresponds to a real scalar or to a Majorana fermion, none of which couples (diagonally) to the $Z$ boson. Still, $Z$-mediated inelastic scatterings involving the next-to-lightest neutral state impose severe constraints on viable MADM scenarios. The requirement that reactions enforcing chemical equilibrium, as well as DM annihilation processes decouple before the EW phase transition, leaving the correct amount of DM, provides another set of constraints. Together with the former ones, these allow to exclude all MADM candidates except scalar and fermion multiplets with hypercharge $y = 1$, for which the lowest dimension representations containing a neutral member are triplets. However, even in this case, gauge annihilation could hardly erase sufficiently the symmetric component, which will eventually constitute most of the DM, and we have also seen that the mass range for which $y = 1$ triplets could constitute good MADM candidates is already strongly disfavored by ID limits. However, this conclusion does not necessarily apply for larger representations. For example, on the basis of the analysis presented in [34], we have argued that a (thermally produced) fermion quintuplet with $y = 1$ and a mass not much above the few TeV range, could account for the whole of DM only if its relic number density is sufficiently enhanced by an initial asymmetry. This case is thus interesting, and we think it deserves a dedicated study.

Finally, it should be mentioned that most of the constraints discussed in this paper can be evaded by departing from minimality. Perhaps the simplest possibility is to add a SM singlet to which the 'would be MADM' can decay, thus transferring its asymmetry-related relic density to a particle with no EW interactions. This would automatically bypass DD constraints and open up large portions of the parameter space. An ADM realization along this line involving an $SU(2)_L$ doublet fermion with $y = 1/2$ decaying into a SM singlet has been put forth for example in Ref. [23].

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