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Flavored dark matter and R-parity violation

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ABSTRACT: Minimal Flavor Violation offers an alternative symmetry rationale to R-parity conservation for the suppression of proton decay in supersymmetric extensions of the Standard Model. The naturalness of such theories is generically under less tension from LHC searches than R-parity conserving models. The flavor symmetry can also guarantee the stability of dark matter if it carries flavor quantum numbers. We outline general features of supersymmetric flavored dark matter (SFDM) models within the framework of MFV SUSY. A simple model of top flavored dark matter is presented. If the dark matter is a thermal relic, then nearly the entire parameter space of the model is testable by upcoming direct detection and LHC searches.

KEYWORDS: Cosmology of Theories beyond the SM, Supersymmetric Standard Model, Discrete and Finite Symmetries, Global Symmetries

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The hierarchy problem and the WIMP miracle independently motivate new dynamics at the weak scale. It is therefore a compelling possibility that both the naturalness and dark matter (DM) problems are resolved by the same new physics. Indeed, the paradigm for physics beyond the Standard Model (SM) for nearly three decades, weak-scale supersymmetry (SUSY) with R-parity conservation (RPC), accomplishes precisely this feat.

However, SUSY with RPC as a natural solution to the hierarchy problem is facing increasingly stringent constraints from a swath of searches at the LHC experiments. For example, searches for multi-jets and missing momentum have placed limits on gluinos and squarks in the TeV range [1, 2]. A characteristic feature of these searches involves the selection of events with large missing transverse momentum, as would inevitably occur from the decay of superpartners to the stable LSP. Thus, a simple way to evade these constraints is to allow interactions that violate R-parity [3–10]. With R-parity violation (RPV) the LSP is unstable, resulting in high multiplicity signatures with leptons and/or jets in the final states. In particular, final states with jets are generally very difficult to constrain due to the large QCD background. For recent studies see [11–18].

While RPV certainly helps to hide SUSY at the LHC, it is a step backwards both from theoretical and phenomenological perspectives. An understanding of the stability of the proton in terms of a protective symmetry is abandoned. Indeed, not all RPV couplings can be allowed simultaneously as then both baryon- and lepton-number-violating couplings will mediate rapid proton decay. This has motivated a number of theoretical proposals to explain the relative size of RPV couplings from a symmetry rationale [19–32].

A particularly interesting proposal to understand the stability of the proton in SUSY with RPV is to invoke the principle of Minimal Flavor Violation (MFV) [33–35], in which one assumes that the non-abelian flavor symmetry $G_F = \text{SU}(3)_Q \times \text{SU}(3)_u \times \text{SU}(3)_d \times \text{SU}(3)_L \times \text{SU}(3)_e$ is broken only by the Yukawa interactions [36–39]. Even in RPC theories, an MFV structure is generally imposed on the soft SUSY breaking interactions in order to suppress unwanted flavor changing neutral currents (FCNCs). It has been shown that

MFV by itself suppresses RPV couplings enough to explain the stability of the proton. Furthermore, the leading RPV superpotential operator relevant for collider phenomenology is $\bar{u}\bar{d}\bar{d}$, leading to signatures with multiple jets, bottom quarks and top quarks, which have much weaker direct LHC constraints compared to RPC SUSY.

While theoretical progress has been made towards understanding the size of RPV couplings, it would still seem we have forsaken our second compelling hint for physics at the weak scale — the WIMP miracle. As RPV renders the LSP unstable there is no viable DM candidate amongst the superpartners of the SM particles.

In this paper we demonstrate that the MFV hypothesis can also provide a symmetry rationale for WIMP DM. In ref. [40] it was shown using an operator analysis that MFV automatically implies exact stability for a large number of representations of the quark flavor group $G_q = \text{SU}(3)_Q \times \text{SU}(3)_u \times \text{SU}(3)_d$, leading to the scenario of flavored dark matter, where DM is charged under G_q . Here we demonstrate that this stability is the result of an underlying Z_3 symmetry, which we term *flavor triality*, that is a subgroup of $\text{SU}(3)_c \times \text{SU}(3)_Q \times \text{SU}(3)_u \times \text{SU}(3)_d$. Under this Z_3 symmetry the SM fields and Yukawa spurions transform trivially, while the FDM candidate is charged. In ref. [40], non-supersymmetric theories were investigated. Here we consider supersymmetric theories of flavored dark matter (SFDM). We will examine the general structure of SFDM models including the effects of SUSY breaking on the flavor splittings in the mass spectrum and couplings.

Finally, we will investigate in detail a model of top flavored dark matter. The DM candidate is taken to be a vector-like fermion contained in a gauge singlet, $\text{SU}(3)_{u_R}$ flavor triplet. A flavor singlet mediator field with SM gauge quantum numbers of right handed top allows the DM to interact with the SM. The DM is a thermal relic due to its efficient annihilation to $t\bar{t}$ pairs in the early universe. By virtue of its coupling to the top, the DM obtains a sizable one loop coupling to the Z boson, which mediates spin independent scattering with nuclei at rates that will be tested by LUX and future ton scale direct detection experiments. Furthermore, the mediator fields, being colored, can be produced at the LHC and decay to DM, leaving signatures with missing energy, jets and tops. For other studies of flavored DM, see refs. [40–46].

1 R-parity violation and MFV SUSY

In the minimal supersymmetric standard model (MSSM), the most general superpotential consistent with gauge invariance and renormalizability is given by

$$\begin{aligned}
 W = & Y_e L H_d \bar{e} + Y_u Q H_u \bar{u} + Y_d Q H_d \bar{d} + \mu H_u H_d \\
 & + \lambda L L \bar{e} + \lambda' L Q \bar{d} + \lambda'' \bar{u} \bar{d} \bar{d} + \mu' L H_u.
 \end{aligned}
 \tag{1.1}$$

It is useful to assign a matter parity to the superfields, $P_M = (-1)^{3(B-L)}$, with B , L the baryon and lepton number, respectively. Under matter parity the quark and lepton superfields have charge -1 and the Higgs superfields have charge $+1$. R-parity is then defined on the spin s component fields as $P_R = (-1)^{2s} P_M$, under which all of the SM fields are even and the superpartners are odd. The terms on the first line of eq. (1.1) conserve matter parity, while those on the second line do not.

The size of the RPV couplings, λ , λ' , λ'' , μ' is severely constrained by the non-observation of proton decay. For example, suppose one desired a small RPV $\bar{u}d\bar{d}$ coupling λ'' in order to facilitate the decay of the LSP to jets, with the aim of suppressing the amount of missing energy. Such a coupling must be at least larger than $\sim \mathcal{O}(10^{-8})$ to mediate a decay on detector scales. If there are also, e.g., $LQ\bar{d}$ couplings λ' present at some level, then squark exchange will mediate the decay of the proton with a lifetime

$$\tau_p \sim 10^{33} \text{ yr} \left(\frac{10^{-19}}{\lambda'} \right)^2 \left(\frac{10^{-8}}{\lambda''} \right)^2 \left(\frac{m_{\tilde{q}}}{\text{TeV}} \right)^4. \quad (1.2)$$

We see that to get interesting modifications to SUSY collider signatures from $\bar{u}d\bar{d}$, we must require $LQ\bar{d}$ to be extremely suppressed to avoid rapid proton decay. This is another way of saying that one can have B or L violation, but not both. Besides proton decay, there are a variety of additional strong constraints on RPV couplings; for a review see ref. [10].

The problem is even worse than this — one could imagine for example that sizable RPV couplings exist only amongst the second and third generation. However, as a result of electroweak symmetry breaking rotations from the gauge to the mass basis will generally induce sizable coupling amongst the first two generations. Clearly, it is therefore desirable to have a symmetry explanation for the suppression of dangerous RPV couplings.

Minimal flavor violation provides one such symmetry principle to explain the smallness of unwanted RPV couplings. The MFV hypothesis promotes the Yukawa couplings to spurion fields transforming under the nonabelian flavor symmetry $G_F = \text{SU}(3)_Q \times \text{SU}(3)_u \times \text{SU}(3)_d \times \text{SU}(3)_L \times \text{SU}(3)_e$. It is assumed that these spurions are the only source of G_F breaking. This assumption typically is imposed in any case of the soft breaking squark masses and trilinear scalar couplings to suppress FCNCs. Since the RPV couplings in the superpotential contain the quark and lepton fields, it is then natural to ask what constraints MFV places on the size of RPV couplings. This question has been addressed in refs. [33–35] which have shown that MFV suppresses RPV couplings enough to explain the stability of the proton, while generating neutrino masses and being consistent with $n - \bar{n}$ oscillation and dinucleon decay constraints.

The primary difference in approach between [33, 34] and [35] is that the latter imposes holomorphicity on the Yukawa spurions, which is required if one imagines the Yukawas to arise from the vacuum expectation values of chiral superfields in a UV completion. This assumption drastically reduces the number of allowed couplings in the superpotential, thus leading to a more predictive setup. For the remainder of this paper, we follow [35] and impose holomorphy on the Yukawa spurions.

In fact, in the limit of massless neutrinos, MFV allows only one operator in the superpotential [35]:

$$W = \frac{1}{2} w'' (Y_u \bar{u})(Y_d \bar{d})(Y_d \bar{d}) \equiv \frac{1}{2} \lambda'' \bar{u} d \bar{d}, \quad (1.3)$$

where the effective $\bar{u}d\bar{d}$ coupling is given by

$$\lambda''_{ijk} = w'' y_i^{(u)} y_j^{(d)} y_k^{(d)} \epsilon_{jkl} V_{il}^*. \quad (1.4)$$

One clearly observes the strong Yukawa and CKM suppression in the effective couplings λ'' amongst the first and second generation fields. At the same time, the coupling λ''_{tsb} , though also suppressed, is still large enough to cause the LSP to decay within the LHC detectors. For example, with an $\mathcal{O}(1)$ MFV coupling w'' , one obtains $\lambda''_{tsb} \sim 10^{-4}$.

With massless neutrinos, the baryon number violating operator in eq. (1.3) is the only superpotential coupling allowed. Additional lepton number violating couplings could in principle arise from the Kähler potential which has no constraints from holomorphicity. However, only $\Delta L = 3$ operators are in fact allowed due to an accidental $Z_3^L \subset \text{SU}(3)_L \times \text{SU}(3)_e$ symmetry present in the theory. Such operators can only occur at the nonrenormalizable level, and are suppressed. Thus the proton is safely stable in the limit of massless neutrinos. It is straightforward to incorporate neutrino masses into the theory, which allows for additional sources of lepton number violation, but MFV still can safely suppress proton decay and be consistent with various other RPV constraints [35].

The collider signatures of MFV SUSY depend primarily on the nature of the LSP. The third generation squarks, due to their large Yukawa couplings, can easily be split from those of the first and second generation. It is therefore possible that the stop or the sbottom is the LSP. A stop LSP will decay due to the interaction in eq. (1.3) via $\tilde{t} \rightarrow sb$, while a sbottom LSP can decay similarly via $\tilde{b} \rightarrow st$. These decays are prompt (except perhaps for the sbottom at low $\tan\beta$). Other possible LSPs, such as neutralinos, charginos, or gluinos can easily leave displaced vertices, as they will decay, e.g., through an offshell stop or sbottom. Such signatures are very distinctive, although presently the limits are not very strong. For further studies of the collider phenomenology in MFV SUSY, see refs. [35, 47, 48].

2 Dark matter stability from MFV

We have seen that MFV SUSY provides an attractive framework to explain the proton stability in supersymmetric theories with RPV. By design, the LSP will decay with a very short lifetime such that SUSY collider events have suppressed missing energy. Therefore, there is no WIMP DM candidate amongst the superpartners of the SM particles. This is of course true of any scenario with RPV couplings large enough to hide SUSY at the LHC. Therefore, one must look beyond the RPV MSSM if one takes the WIMP miracle seriously.¹

Despite the fact that MFV SUSY does not contain a WIMP candidate, MFV itself provides a motivation for DM, as we now discuss. In ref. [40], it was shown that the MFV hypothesis implies absolute stability for certain representations that transform nontrivially under the quark flavor group $G_q = \text{SU}(3)_Q \times \text{SU}(3)_u \times \text{SU}(3)_d$ and are singlets under $\text{SU}(3)_c$. If these states are also electrically neutral, they make excellent dark matter candidates.

In ref. [40] DM stability was demonstrated through an operator analysis. Conditions were derived for the existence of the most general operator composed of a single DM multiplet, SM fields and Yukawa spurions that would mediate the decay of the would-be DM particle if present. Here we wish to revisit this question from a symmetry perspective.

¹Another possible DM candidate outside the WIMP paradigm is a light gravitino, which we do not investigate here.

	$SU(3)_c$	$SU(3)_Q$	$SU(3)_u$	$SU(3)_d$
Q	$\mathbf{3}$	$\mathbf{3}$	$\mathbf{1}$	$\mathbf{1}$
\bar{u}	$\bar{\mathbf{3}}$	$\mathbf{1}$	$\bar{\mathbf{3}}$	$\mathbf{1}$
\bar{d}	$\bar{\mathbf{3}}$	$\mathbf{1}$	$\bar{\mathbf{3}}$	$\mathbf{1}$
Y_u	$\mathbf{1}$	$\bar{\mathbf{3}}$	$\mathbf{3}$	$\mathbf{1}$
Y_d	$\mathbf{1}$	$\bar{\mathbf{3}}$	$\mathbf{1}$	$\mathbf{3}$
G	$\mathbf{8}$	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$

Table 1. Representations of fields charged under $SU(3)_c \times SU(3)_Q \times SU(3)_u \times SU(3)_d$.

We will show that DM stability can be traced to the presence of an accidental Z_3 symmetry, which we call *flavor triality*, that is present under the assumption of MFV.

It is in fact very simple to see that an accidental symmetry is present that can stabilize DM. Consider the following discrete Z_3 transformation which is an element of $SU(3)_c \times SU(3)_Q \times SU(3)_u \times SU(3)_d$:

$$U = (\omega^2)_c \times (\omega)_Q \times (\omega)_u \times (\omega)_d, \quad (2.1)$$

where $\omega \equiv e^{2\pi i/3}$ and the subscript indicates the group which contains the corresponding Z_3 element. Using the representations listed in table 1, one can easily check that *all* of the SM fields and Yukawa spurions transform trivially under (2.1). For example, under (2.1) $Q \rightarrow \omega^3 Q = Q$, $\bar{u} \rightarrow \omega^{-3} \bar{u} = \bar{u}$, etc.

Now consider a new matter multiplet χ which is to be our DM candidate. This field is a color singlet (as DM is both color and electrically neutral) and transforms under G_q with irreducible representation

$$\chi \sim (n_Q, m_Q)_Q \times (n_u, m_u)_u \times (n_d, m_d)_d, \quad (2.2)$$

where we have used tensorial notation with n_Q, m_Q etc. taking possible values $0, 1, 2, \dots$. Under the Z_3 transformation (2.1), χ transforms as

$$\chi \rightarrow \omega^{n-m} \chi, \quad (2.3)$$

where we have defined $n \equiv n_Q + n_u + n_d$, $m \equiv m_Q + m_u + m_d$. Thus, χ will transform nontrivially under (2.1) provided the following condition is met:

$$(n - m) \bmod 3 \neq 0. \quad (2.4)$$

Since the SM fields and Yukawa spurions transform trivially under the Z_3 transformation, provided the condition (2.4) is met, χ transforms nontrivially and will therefore be absolutely stable. We will use the term *flavor triality* to refer to this Z_3 symmetry under which DM is charged and the SM is neutral. Examples of G_q representations for which (2.4) holds are flavor triplets, (e.g., $(\mathbf{3}, \mathbf{1}, \mathbf{1})$ etc.), sextets, (e.g., $(\mathbf{6}, \mathbf{1}, \mathbf{1})$ etc.) and certain mixed representations, (e.g., $(\mathbf{3}, \mathbf{3}, \mathbf{1})$ etc.). On the other hand flavor singlets, octets and certain

mixed representations, (e.g., $(\mathbf{3}, \bar{\mathbf{3}}, \mathbf{1})$ etc.) do not meet 2.4. We wish to emphasize that flavor triality does not require SUSY; it is simply a consequence of MFV.

Provided we strictly enforce MFV, the conclusion is that the dark matter candidate χ is exactly stable if (2.4) holds, even in the presence of arbitrary higher dimension operators. A natural question is whether stability holds in the presence of deviations from MFV. Without reference to a concrete UV completion, perhaps the most sensible way to parametrize deviations from the MFV hypothesis is to consider additional spurions that break flavor. Provided that these new spurions transform trivially under flavor triality, then DM stability will not be spoiled. For example, additional spurions transforming like the SM Yukawas or, e.g., as flavor octets will not cause the decay of the DM.

The SM flavor symmetries are anomalous, so one may worry that flavor triality is broken at the quantum level. An attractive possibility is that the SM flavor symmetries are gauged in the UV, which ultimately implies the presence of additional matter that would render the theory anomaly free. Flavor triality would then fundamentally be a discrete gauge symmetry and could not be spoiled by quantum gravitational effects [49]. This possibility is similar in spirit to obtaining R-parity from broken gauged $U(1)_{B-L}$ [50, 51].

Finally, we wish to comment on the possibility of lepton-flavored dark matter. Recall that the stability of quark-flavored dark matter considered above rests on the fact that the quark fields transform trivially under the element (2.1) of $SU(3)_c \times SU(3)_Q \times SU(3)_{u_R} \times SU(3)_{d_R}$. This is a consequence of the $SU(3)_c$ charge of the quark fields. However, since the leptons are not colored, there is no transformation analogous to (2.1) under which the lepton fields transform trivially. Therefore, MFV does not in general imply stability for lepton-flavored dark matter.

3 Flavored dark matter in MFV SUSY

With the stability condition (2.4) we are now in a position to consider explicit models of flavored DM within the framework of MFV SUSY. With only the addition of a DM multiplet, interactions between the DM and MSSM fields are not possible at the renormalizable level. One could consider effective theories of flavored DM, as was done in ref. [40], but in such theories one cannot address important questions relevant for supersymmetric theories such as gauge coupling unification. Renormalizable theories require an additional mediator field that couples DM to the MSSM fields. See refs. [44, 45] for renormalizable FDM models in non-supersymmetric setups.

The basic models of SFDM contain a vector-like DM multiplet X, \bar{X} and a vector-like mediator field Y, \bar{Y} . The superpotential is given by

$$W = \hat{M}_X X \bar{X} + \hat{M}_Y Y \bar{Y} + \hat{\lambda} X Y \Phi_{\text{SM}}, \tag{3.1}$$

where Φ_{SM} is one of the quark fields $Q, \bar{u},$ or \bar{d} . For simplicity we will restrict our discussion to models in which X is a flavor triplet and SM gauge singlet, while Y is a flavor singlet and carries SM gauge charges, but of course other possibilities exist.

One of the main differences between non-supersymmetric theories of FDM and SFDM is the constraint that holomorphy places on the superpotential mass and coupling param-

eters. No further Yukawa insertions are allowed in the superpotential, which naively leads one to expect that the masses and couplings of the individual flavors of X are not split. Obviously this constraint is not present in non-supersymmetric theories. However, as we will see below, non-holomorphic terms in the Kähler potential can lead to large flavor splittings, which will be important for phenomenology.

The mediator Y . The mediator Y will generally be charged under the gauge group of the SM, and will therefore typically ruin gauge coupling unification. To avoid this, we can embed Y, \bar{Y} in a complete $SU(5)$ multiplet. For example, if $X \sim (\mathbf{1}, \mathbf{1}, \mathbf{3})_{G_q}$, then one can embed Y into a $\mathbf{5} \supset (\mathbf{3}, \mathbf{1}, -\frac{1}{3})_{SM}$ which allows the coupling $X Y \bar{d}$. Alternatively, if $X \sim (\bar{\mathbf{3}}, \mathbf{1}, \mathbf{1})_{G_q}$ or $X \sim (\mathbf{1}, \mathbf{3}, \mathbf{1})_{G_q}$, we can embed Y in a $\bar{\mathbf{10}} \supset (\bar{\mathbf{3}}, \bar{\mathbf{2}}, -\frac{1}{6})_{SM} + (\mathbf{3}, \mathbf{1}, \frac{2}{3})_{SM}$, which allows couplings to Q or \bar{u} , respectively. Thus, our framework is compatible with gauge coupling unification.

Flavor splittings. As mentioned above, the masses and couplings of the individual flavors are degenerate at the level of the renormalizable superpotential due to holomorphy of the Yukawa spurions. However, there are several ways by which flavor splittings can be induced as we now discuss.

There is no holomorphy constraint on the Kähler potential, which may contain additional Yukawa insertions that are consistent with the flavor symmetry. The kinetic term need not be canonical, e.g.,

$$\int d^4\theta \left(X^\dagger \hat{k}_X X + \bar{X} \hat{\bar{k}}_X \bar{X}^\dagger \right), \tag{3.2}$$

where $\hat{k}_X, \hat{\bar{k}}_X$ are matrices in flavor space which contain Yukawa insertions. For example, if $X \sim (\mathbf{1}, \mathbf{3}, \mathbf{1})_{G_q}$, we have

$$\begin{aligned} \hat{k}_X &= 1 + k Y_u Y_u^\dagger + \dots, \\ \hat{\bar{k}}_X &= 1 + \bar{k} Y_u Y_u^\dagger + \dots, \end{aligned} \tag{3.3}$$

where we have written explicitly the leading Yukawa insertion. If the MFV couplings k, \bar{k} are $O(1)$, this can lead to a sizable splitting between the masses and physical couplings of the third and first two generations of flavors after canonical normalization of the kinetic term is carried out.

Another important effect comes from SUSY breaking terms in the Kähler potential, which can generate *effective* non-holomorphic mass terms in the superpotential, e.g.,

$$\int d^4\theta \left(\frac{S^\dagger}{M} \bar{X} \hat{\mu}_X X + \text{h.c.} \right) = \int d^2\theta \left(\frac{F}{M} \bar{X} \hat{\mu}_X X + \text{h.c.} \right), \tag{3.4}$$

where $S = \theta^2 F$ is a spurion parameterizing SUSY breaking, with \sqrt{F} the SUSY breaking scale and M the messenger scale. Since it originates in the Kähler potential, the coupling $\hat{\mu}_X$ need not be holomorphic, e.g., for $X \sim (\mathbf{1}, \mathbf{3}, \mathbf{1})_{G_q}$, we have

$$\hat{\mu}_X = \mu_0 + \mu_1 Y_u Y_u^\dagger + \dots, \tag{3.5}$$

where again we have written explicitly the leading Yukawa insertion and μ_0, μ_1 are $O(1)$ MFV couplings.

These two effects, eqs. (3.2) and (3.4), are enough to generate large splittings in the masses and couplings between the top flavor and up and charm flavors, in the example $X \sim (\mathbf{1}, \mathbf{3}, \mathbf{1})_{G_q}$. Typically, up-type Yukawa insertions will lead to a bigger splittings, although this is not necessarily true at large $\tan\beta$ since the bottom Yukawa coupling can be $O(1)$. For example, in the case of $X \sim (\mathbf{1}, \mathbf{1}, \mathbf{3})_{G_q}$, one can expect at large $\tan\beta$ sizable splittings between the bottom flavor and down and strange flavors of X .

Majorana mass term? It may be of phenomenological interest to consider the possible existence of a Majorana mass term for the DM, $W \supset \delta M_X X X + \delta \bar{M}_X \bar{X} \bar{X}$. For instance, such a term will split Dirac fermion DM into two Majorana states, which can dramatically affect the predictions for direct detection rates. In the context of flavored dark matter, this possibility was explored in ref. [45], where it was noted that such a Majorana mass violates MFV and can be regarded as an additional flavor-breaking spurion (e.g., if X is a flavor triplet then δM_X is a flavor anti-sextet). We would like to emphasize that such a spurion is charged under flavor triality (2.1) and will generically spoil the DM stability unless further symmetries are invoked.

The scale of SFDM. The superpotential masses \hat{M}_X, \hat{M}_Y in eq. (3.1) can naturally be tied to the weak scale through SUSY breaking, as can be seen from eq. (3.4) in which the effective mass term generated is of order $m_{\text{soft}} = F/M$. This is analogous to the Giudice-Masiero solution to the μ problem [52].

4 Example: top flavored dark matter

We now describe the phenomenology of a concrete model of SFDM in which the DM carries “top” flavor. The model contains the following additional chiral multiplets with quantum numbers,

$$\begin{aligned} X_i &\supset (\eta_i, \chi_i) \sim (\mathbf{1}, \mathbf{1}, 0)_{\text{SM}} \times (\mathbf{1}, \mathbf{3}, \mathbf{1})_{G_q}, \\ Y &\supset (\phi, \psi) \sim (\mathbf{3}, \mathbf{1}, \frac{2}{3})_{\text{SM}} \times (\mathbf{1}, \mathbf{1}, \mathbf{1})_{G_q}, \end{aligned} \tag{4.1}$$

along with conjugate fields \bar{X}^i, \bar{Y} . The index $i = u, c, t$ denotes the flavor of X , and η_i (ϕ) and χ_i (ψ) are the scalar and fermionic components of X_i (Y), respectively. The superpotential is given by

$$W = \hat{M}_X X_i \bar{X}^i + \hat{M}_Y Y \bar{Y} + \hat{\lambda} X_i Y \bar{u}^i. \tag{4.2}$$

The DM candidate can in principle be either a scalar or fermion component of X_i ; for concreteness we will focus on fermionic DM χ_i . As discussed in the previous section, the masses and couplings of the individual flavors in X_i are degenerate at the level of the renormalizable superpotential, which is a consequence of holomorphy of the Yukawa spurions. However, non-holomorphic and SUSY breaking terms in the Kähler potential

will generically induce flavor splittings which in this model can easily be $O(1)$ between the first two and third generations due to the large top Yukawa coupling, as we now discuss.

First, the non-canonical kinetic terms as in eqs. (3.2), (3.3) can be brought to canonical form through field redefinitions, $X \rightarrow Z_X X$. $\bar{X} \rightarrow \bar{Z}_X \bar{X}$, where

$$\begin{aligned} Z_X &\approx \text{diag}(1, 1, (1 + ky_t^2)^{-1/2}), \\ \bar{Z}_X &\approx \text{diag}(1, 1, (1 + \bar{k}y_t^2)^{-1/2}), \end{aligned} \tag{4.3}$$

with k, \bar{k} $O(1)$ MFV couplings defined in eq. (3.3). Furthermore, accounting for SUSY breaking terms in the Kähler potential which generate effective non-holomorphic superpotential mass terms as in eqs. (3.4), (3.5), the masses for the fermions $\chi_i^T = (\chi_u, \chi_c, \chi_t)$ can be written as

$$\begin{aligned} \mathcal{M}_\chi &= \bar{Z}_X \left(\hat{M}_X + \frac{F}{M} \hat{\mu}_X \right) Z_X \\ &\approx \text{diag} \left(m, m, \frac{m + (F/M)\mu_1 y_t^2}{\sqrt{(1 + ky_t^2)(1 + \bar{k}y_t^2)}} \right), \end{aligned} \tag{4.4}$$

where we have defined $m = \hat{M}_X + (F/M)\mu_0$, with μ_0, μ_1 $O(1)$ MFV couplings defined in eq. (3.5). We observe that the first two generation dark fermions χ_u, χ_c are degenerate up to fine splittings induced by the up and charm Yukawas, while the third generation top-flavored fermion χ_t can obtain a large mass splitting due to the $O(1)$ top Yukawa. In particular, χ_t can be the lightest state in the new sector in which case it will be stable and the DM candidate, and we specialize to this case for the remainder of the paper.

The main interactions governing the cosmology and phenomenology of the model come from the superpotential interaction with \bar{u} in eq. (4.2). In component form, the important terms are

$$- \mathcal{L} \supset \bar{u}_R^i \lambda_i^j \chi_j \phi + \tilde{u}_R^{\dagger i} \lambda_i^j \chi_j \psi + \bar{u}_R^i \lambda_i^j \eta_j \psi + \text{h.c.}, \tag{4.5}$$

where again $\eta_i^T = (\eta_u, \eta_c, \eta_t)$ are the scalar components of X_i , and ϕ (ψ) is the scalar (fermion) component of the mediator field Y . The couplings λ are split due to the canonical normalization of X (4.3):

$$\begin{aligned} \lambda &= \hat{\lambda} Z_X \\ &\approx \text{diag} \left(\hat{\lambda}, \hat{\lambda}, \frac{\hat{\lambda}}{\sqrt{(1 + ky_t^2)}} \right). \end{aligned} \tag{4.6}$$

This demonstrates that χ_u, χ_c have identical couplings up to the small y_u, y_c induced splittings, while the DM particle χ_t can have a larger or smaller coupling depending on the value of k . For later reference we write explicitly the terms involving the DM candidate χ_t ,

$$- \mathcal{L} \supset \lambda_t \bar{t}_R \chi_t \phi + \lambda_t \tilde{t}_R^* \chi_t \psi + \text{h.c.}, \tag{4.7}$$

where $\lambda_t \equiv \hat{\lambda} / \sqrt{1 + ky_t^2}$.

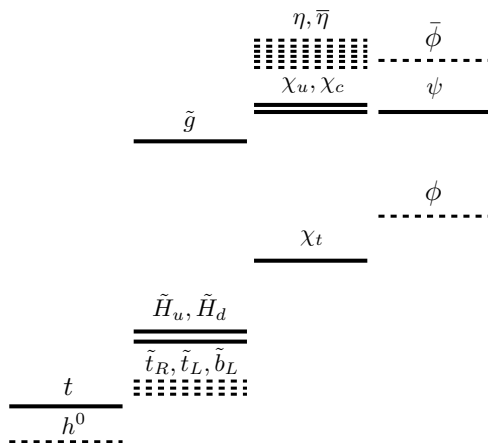


Figure 1. Example spectrum of the top-flavored DM model. The DM is χ_t .

Additional soft scalar squared mass terms will split the scalar components from their fermionic partners. In particular, the additional scalars $\eta_i, \bar{\eta}_i$ can be raised and will not be critical to our phenomenological discussion below, although they can lead to interesting flavor specific signatures if produced at the LHC. Of more direct importance is the mass of the mediators, $\phi, \bar{\phi}$ and ψ , which play a crucial role in χ_t annihilation processes in the early universe, mediate the scattering of χ_t with nuclei, and by virtue of their $SU(3)_c$ charge can be directly produced at the LHC. At the supersymmetric level, the scalar mediators ϕ and $\bar{\phi}$ are degenerate, but SUSY breaking can split and mass mix these states through soft scalar squared masses and b terms. We will make the simplifying assumptions below that ϕ - $\bar{\phi}$ mixing is negligible and that $m_\phi \ll m_{\bar{\phi}}, m_\psi$. In this regime ϕ will mediate the dominant interaction between the DM and the SM. We display in figure 1 one possible spectrum of the top-flavored scenario.

We note ref. [45] considered a non-supersymmetric scenario of top flavored dark matter as an explanation to the Tevatron top quark forward-backward asymmetry. They focused on light DM, $m_{\chi_t} \lesssim 100$ GeV, whereas we will be concerned with DM masses $m_{\chi_t} > m_t$.

Thermal relic abundance. In the regime $m_{\chi_t} > m_t$ and $m_\phi \ll m_\psi$, the dominant process governing the relic abundance of χ_t is

$$\chi_t \bar{\chi}_t \rightarrow t \bar{t}, \quad (4.8)$$

which is mediated by t -channel exchange of the scalar mediator ϕ . The thermally averaged annihilation cross section is

$$\langle \sigma v \rangle_{t\bar{t}} = \frac{N_c \lambda_t^4 m_{\chi_t}^2}{32\pi (m_{\chi_t}^2 + m_\phi^2 - m_t^2)^2} \left(1 - \frac{m_t^2}{m_{\chi_t}^2} \right)^{1/2}, \quad (4.9)$$

where $N_c = 3$ is the number of colors. The process $\chi_t \bar{\chi}_t \rightarrow \tilde{t}_i \tilde{t}_j^*$ can also be important if ψ is similar in mass to ϕ . While this will result in small numerical changes to the parameters

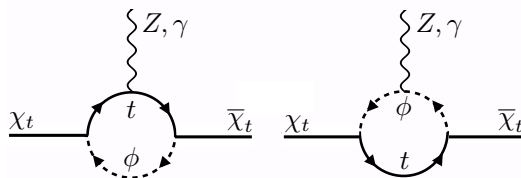


Figure 2. One-loop diagrams generating effective Z, γ couplings for top-flavored DM, with amplitudes given in eqs. (4.10), (4.13).

for which the correct relic abundance is obtained, it will not qualitatively change the conclusions presented below.

The observed cold DM relic abundance of $\Omega h^2 \approx 0.12$ [53] is obtained for an annihilation cross section of $\langle \sigma v \rangle \approx 1.5$ pb (for a Dirac fermion). A minimum coupling of $\lambda_t = 0.35$ is needed to achieve this cross section if the $t\bar{t}$ final state is considered. Exceptions to this statement are if annihilation to the other final states are present and substantial, or when $m_\phi - m_{\chi_t} < m_{\chi_t}/20$ and the effects of coannihilation can assist in achieving the correct relic density. We do not include the effects of coannihilation in our numerical results below, but see ref. [54] for a recent study.

The parameters needed for $\langle \sigma v \rangle = 1.5$ pb are shown in figure 3, given as the curve labeled σ_{th} . Note that in the parameter space between this contour and the line $m_\phi = m_{\chi_t}$, the annihilation cross section is larger and thus χ_t cannot be all of the DM. Outside of this curve the cross section is lower, and other channels are needed so that χ_t does not overclose the universe. In figure 3 we have included only the contribution from the $t\bar{t}$ mode (4.9), which is valid in the limit of heavy fermionic mediator $m_\psi \gg m_\phi$.

Direct detection. The strongest constraints on this model come from direct detection experiments. Although there is no tree-level coupling of DM with nucleons, a sizable interaction is generated at one-loop. Z exchange is the most important process, arising from an effective coupling of DM to the Z . The diagrams shown in figure 2 generate the operator

$$\mathcal{L} \supset g_Z Z_\mu \bar{\chi}_t \gamma^\mu P_L \chi_t \quad (4.10)$$

at one-loop. We include only the contributions from one-loop $t - \phi$ exchange in our computations, which is valid in the regime $m_\psi \gg m_\phi$. There are additional diagrams with \tilde{t}_i and ψ in the loop that can be numerically important if ϕ and ψ have comparable masses.

The full expression for g_Z can be found in the appendix. In the limit of $m_\phi \gg m_t, m_{\chi_t}$,

$$g_Z \simeq -\frac{g}{c_w} \frac{\lambda_t^2 N_c}{16\pi^2} \left(\frac{m_t}{m_\phi} \right)^2 \left(1 + \log \left[\frac{m_t^2}{m_\phi^2} \right] \right). \quad (4.11)$$

In general g_Z is suppressed by the mass of the fermion in the loop, but here the large top mass and $\mathcal{O}(1)$ DM - top coupling λ_t assist in generating a relatively large Z coupling.

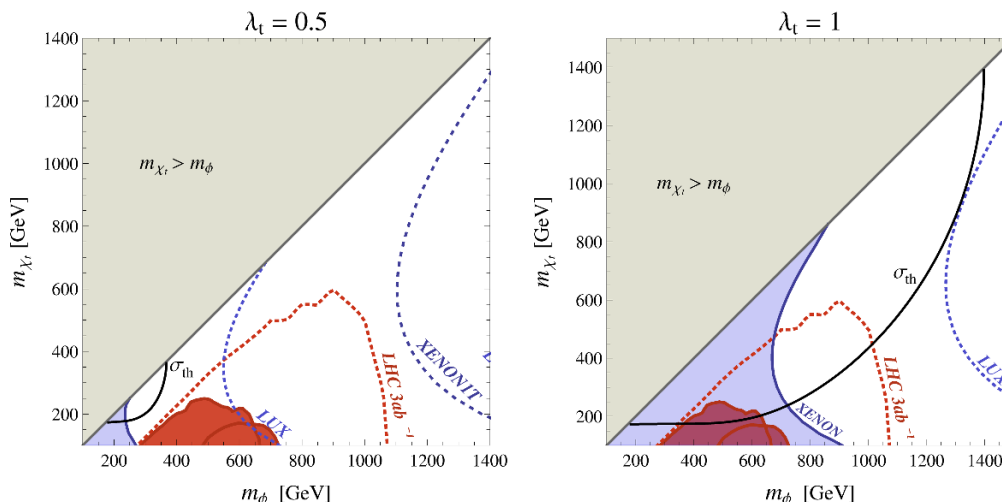


Figure 3. Limits on the parameter space of top-flavored DM from LHC stop searches with 8 TeV data (solid red area) and XENON100 (solid blue area) for $\lambda_t = 0.5$ and $\lambda_t = 1$. The red dashed line is a projection for the 95% CL exclusion limit with a 3000 fb^{-1} run at the 14 TeV LHC [65]. For the LUX (XENON1T) projection, we show the contour where 1 event is expected per 10000 (10^5) $\text{kg} \cdot \text{d}$. The direct detection limits assume that χ_t saturates the observed relic density. The solid black line labeled σ_{th} is where $\langle\sigma v\rangle = 1.5 \text{ pb}$, the annihilation cross section necessary to obtain a relic density of $\Omega h^2 = 0.12$. Note that below and to the right of the black line, $\langle\sigma v\rangle$ is lower and hence χ_t is too abundant, and additional annihilation channels are required.

The Z coupling mediates a spin-independent (SI) scattering of dark matter and nuclei. The differential cross section is

$$\frac{d\sigma_Z}{dE_R} = \frac{m_N (f_p Z + (A - Z)f_n)^2}{\pi v^2} F^2[E_R], \quad (4.12)$$

where the couplings $f_n = g_Z G_{Fc_w} / \sqrt{2}g$ and $f_p = -(1 - 4s_w^2)g_Z G_{Fc_w} / \sqrt{2}g$. Note that we focus on spin-independent (SI) interactions; although this operator give rise to spin-independent and spin-dependent scattering with similar cross sections, the experimental limits for SI interactions are much stronger.

It is also necessary to consider the same diagrams with an external photon instead of a Z boson, as shown in figure 2. These generate a magnetic dipole moment for the DM:

$$\mathcal{L} \supset \frac{\mu_\chi}{2} \bar{\chi}_t \sigma^{\mu\nu} \chi_t F_{\mu\nu}. \quad (4.13)$$

The dipole moment in the limit $m_\phi \gg m_t, m_{\chi_t}$ is

$$\mu_\chi \simeq \frac{e\lambda_t^2 m_{\chi_t}}{32\pi^2 m_\phi^2}, \quad (4.14)$$

with the full one-loop result given in the appendix.

We find that magnetic dipole interactions provide a non-negligible contribution to the rate for larger DM mass, where the dominant contribution is through the dipole-charge

interaction for Xenon. The scattering cross section is

$$\frac{d\sigma_{DZ}}{dE_R} = \frac{m_N e^2 Z^2}{\pi v^2} \frac{\mu_\chi^2}{4 m_N^2} \left(\frac{m_N v^2}{2E_R} - \frac{m_{\chi_t} + 2m_N}{2m_{\chi_t}} \right) F^2[E_R]. \quad (4.15)$$

For $m_{\chi_t} \lesssim m_\phi \sim 1.5$ TeV, the dipole contribution affects the rate by up to 40-50%, depending on the energy range considered.

Other interactions are also present at one-loop but negligible compared to Z exchange and dipole-interactions, as we discuss in the appendix.

Current experimental limits from XENON100 [55] constrain the parameter space for couplings $\lambda_t \gtrsim 0.5$. (Recall also for $\lambda_t < 0.35$, other annihilation channels must be present or the DM will overclose the universe if it is a thermal relic.) With increased couplings, larger masses m_{χ_t} and m_ϕ are required to satisfy these limits.

To calculate the limits for XENON100, we compute both Z and dipole mediated scattering rates over the stated nuclear recoil energy range of 6.6-30.5 keV_r. The published limits are stated in terms of a 90% CL exclusion limit on spin-independent nucleon scattering cross section σ_n , which apply for Z exchange. However, those limits are not directly applicable when dipole interactions are included because of the difference in energy dependence, as can be seen in eq. 4.15.

Instead, we match the onto the published limits by calculating the number of expected events from Z exchange only, finding that the published limit is well approximated by requiring ≤ 2 events with 2324 kg · d. We then include the dipole interactions in the event rate and re-evaluate the limit. The end result only changes the limits in m_{χ_t} and m_ϕ by less than 5%, so we do not try to model the energy-dependent acceptance of the experiment.

We also consider projected limits from LUX [56], for which results are expected in the near future. The expected energy range is 5-25 keV_r, and with 10000 kg · d exposure. The procedure for computing these limits is the same as described above. Again the effect of including dipole interactions only changes the limits by $\lesssim 5\%$. Finally, we include XENON1T [57] projections, assuming the same energy range as XENON100 and 10^5 kg · d exposure. For all of the above calculations we assume a Standard Halo Model with $\rho_\chi = 0.4$ GeV/cm³, $v_{esc} = 550$ km/s, $v_e = 240$ km/s.

We show these results in figure 3. The anticipated LUX reach can effectively test this model in the case that χ_t is a thermal relic annihilating primarily to tops. The limit and projection curves approximately follow lines of constant g_Z^2/m_{χ_t} , which keep the event rate constant. At fixed m_ϕ , the constraint becomes stronger with smaller DM mass primarily because the rate scales as $1/m_{\chi_t}$, while it again becomes stronger near $m_{\chi_t} \lesssim m_\phi$ due to the enhancement of g_Z .

LHC signatures. The collider signatures of this scenario at the LHC depend in detail on the spectrum of the SM superpartners and the SFDM sector. An example spectrum is illustrated in figure 1. We will focus here on the case of a stop LSP. Provided the stop is lighter than the DM χ_t (more specifically $m_{\tilde{t}_1} < m_{\chi_t} + m_\psi$), then it will decay via the baryon number violating vertex (1.3):

$$\tilde{t}_1 \rightarrow \bar{b}\bar{s}. \quad (4.16)$$

Stops decaying via (4.16) are not constrained by existing searches for paired dijet resonances [58–61].

The DM χ_t is of course stable and will lead to \cancel{E}_T if produced. There are several ways in which χ_t can be produced, as we now discuss. For example, the scalar mediator ϕ , being colored, can be pair produced through strong interactions. Once produced, it will decay via

$$\phi \rightarrow t \bar{\chi}, \tag{4.17}$$

resulting in a signature of $t\bar{t} + \cancel{E}_T$ for ϕ pair production. Remarkably, ϕ acts as a “fake stop”, in that the signature of ϕ pair production mimics exactly the signature of stop pair production in standard RPC scenarios when the stop decays to a top and a stable neutralino LSP. The limits (and projected reach) on direct stop pair production from CMS [62] and ATLAS [63, 64] can thus be directly applied to the case of ϕ production and are displayed in figure 3.

Another channel which is important is pair production the fermionic mediator ψ , which will decay via

$$\psi \rightarrow \tilde{t}_i \bar{\chi}. \tag{4.18}$$

If the LSP is mostly right handed stop, $\tilde{t}_1 \sim \tilde{t}_R$, then ψ will decay primarily to the stop LSP and the DM χ_t . The stop LSP will subsequently decay according to eq. (4.16), leading to a signature of $4j + \cancel{E}_T$ for ψ pair production. This signature is similar, though not identical, to gluino pair production followed by $\tilde{g} \rightarrow q\bar{q}\chi^0$ mediated by a heavy off-shell squark in standard RPC scenarios. The gluino searches using data from the 7 TeV run provide the strongest limits on ψ pair production. This is because the recent 8 TeV analyses are optimized for a high mass gluino, which, being a color octet, has a much larger production cross section than the color triplet ψ in our scenario. The ATLAS search [66] probes $m_\psi \lesssim 350$ GeV for $m_{\chi_t} \sim 200$ GeV, $m_{\tilde{t}_1} \lesssim 100$ GeV. If the DM χ_t is lighter, then the sensitivity can be extended up to higher masses (of order 800 GeV for massless χ_t), although in this range new annihilation channels are required to obtain a viable cosmology.

If the new states in the flavored dark sector lie below the MSSM superpartners, particularly the gluino and first and second generation squarks, additional signatures are possible. For example, if kinematically allowed the gluino can decay with a sizable branching ratio via

$$\tilde{g} \rightarrow \phi \bar{\psi}. \tag{4.19}$$

The mediators ϕ and ψ subsequently decay according to eqs. (4.17) and (4.18), leading to a multi-top, multi-jet + \cancel{E}_T final state signature of gluino pair production.

Through their coupling to the top and stops the new states χ_t , ϕ , and ψ provide a relevant contribution to the Higgs mass at two loops. In particular, for the case of a thermal relic χ_t considered here, the required $O(1)$ coupling λ_t implies that the new states χ_t , ϕ , and ψ cannot be too heavy without a significant tuning of the weak scale. The level of tuning induced by these states is similar to that of the gluino, which also contributes to the Higgs mass at two loops. However, in comparison to the gluino in the MSSM with RPC, the LHC limits on the ϕ and ψ are generically much weaker, as discussed above.

5 Outlook

Supersymmetry with R-parity violation is only weakly constrained by searches at the LHC, particularly if the LSP decays to jets through the $\bar{u}d\bar{d}$ operator. However, with RPV one abandons a symmetry rationale for the proton stability and the WIMP miracle. Minimal Flavor Violation can provide an explanation for the suppression of proton decay and, as we have demonstrated in this paper, the stability of WIMP dark matter. This framework provides a compelling explanation for the naturalness of the weak scale, dark matter, and the lack of evidence for new physics at the LHC.

In this work we have established the existence of a Z_3 discrete symmetry, *flavor triality*, which is a consequence of the MFV hypothesis. Flavor triality presents the prospect of flavored dark matter. We have investigated general aspects of theories of flavored dark matter in the MFV SUSY framework. Flavor splittings in the masses and couplings, which are relevant for cosmology and phenomenology, arise from non-holomorphic terms in the Kähler potential. The SFDM framework is compatible with gauge coupling unification and the mass scale of the dark sector states can naturally be tied to the weak scale through SUSY breaking.

We have studied in detail a specific scenario of top flavored dark matter. The dark matter is a thermal relic, annihilating to top quarks in the early universe. This model can be tested by LUX and future ton scale Xenon experiments. Furthermore, the scalar mediator ϕ can be directly produced at the LHC and can mimic a stop from standard RPC scenarios. Therefore, in this scenario it is conceivable that we will first discover the “fake stop” ϕ in searches for $t\bar{t}+\cancel{E}_T$, while the true stop responsible for canceling the Higgs mass quadratic divergences is buried in the QCD multi-jet background. Many other novel signatures are possible in this scenario, as we have discussed above.

We have explored just one possible incarnation of SFDM; there are many other models that would be worthwhile to explore. The condition in eq. (2.4) ensuring dark matter stability is satisfied for a variety of representations of G_q . There are additional possibilities in the choice of SM gauge quantum numbers for the dark matter multiplet X (the only requirement being that it contains a color and electrically neutral component), as well as the flavor and gauge representations of the mediator field. Within a given model, there are also several candidates for the dark matter particle. For instance in the model studied in this paper, one could also consider the scalar component as the dark matter, as well as, e.g., up-flavored dark matter, as has been studied recently in simplified DM models [54, 68–70]. A systematic investigation of these theories should be carried out.

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A Direct detection

This appendix summarizes one-loop amplitudes relevant to direct detection for the interactions in eq. (4.7). Useful related discussions can be found in refs. [44, 45]. As discussed in the text, there is a Z coupling, eq. (4.10). In the limit of zero momentum transfer ($q^2 \rightarrow 0$) the coupling is

$$g_Z = \frac{\lambda_t^2 N_c}{16\pi^2} \int_0^1 dy (1-y) I(y), \quad (\text{A.1})$$

$$I(y) = g_R \left[\log \left(\frac{\hat{\Delta}_S}{\hat{\Delta}_F} \right) - 1 - \frac{r_\chi y(1-y)}{\hat{\Delta}_S} \right] + g_L \left(\frac{r_t}{\hat{\Delta}_F} \right),$$

where $g_{L,R}$ are the couplings of the left- and right-handed top to Z . The other dimensionless quantities are

$$\hat{\Delta}_F \equiv y(1 + r_\chi(y-1)) - r_t(y-1), \quad (\text{A.2})$$

$$\hat{\Delta}_S \equiv y(r_t + r_\chi(y-1)) + (1-y), \quad (\text{A.3})$$

$$r_t \equiv \left(\frac{m_t}{m_\phi} \right)^2, \quad r_\chi \equiv \left(\frac{m_{\chi t}}{m_\phi} \right)^2, \quad (\text{A.4})$$

with $\hat{\Delta}_F$ ($\hat{\Delta}_S$) corresponding to the propagator factor for the diagram with emission of the Z from the fermion (scalar) in the loop.

For an external photon, the diagrams and the calculation are the same, but with $g_L = g_R = 2e/3$. The vector and axial-vector couplings to the photon vanish in the limit $q^2 \rightarrow 0$, as required by electromagnetic gauge invariance. There is however a nonzero magnetic dipole moment, $(\mu_\chi/2)\bar{\chi}\sigma^{\mu\nu}\chi F_{\mu\nu}$:

$$\mu_\chi = \frac{e\lambda_t^2 m_{\chi t}}{32\pi^2 m_\phi^2} \int_0^1 dy \frac{2y(1-y)^2(1+r_t+2r_\chi y(y-1))}{\hat{\Delta}_F \hat{\Delta}_S}. \quad (\text{A.5})$$

For reference we also give the amplitude for the (subdominant) charge-charge contribution to direct detection. Although the vector coupling to photons is zero in the $q^2 = 0$ limit (the DM is neutral), there are q^2 corrections that give rise to charge-charge interactions (as well as velocity-suppressed charge-dipole interactions). The amplitude has a contribution $b_\chi q^2 \bar{\chi}\gamma^\mu\chi$, where the coefficient is:

$$b_\chi = \frac{e\lambda_t^2 N_c}{48\pi^2 m_\phi^2} \int dy I(y), \quad (\text{A.6})$$

$$I(y) = \frac{(1-y)^3}{6} \left(\frac{1}{\hat{\Delta}_F} + \frac{r_t + r_\chi y^2}{\hat{\Delta}_F^2} + \frac{2r_\chi(y-1)y}{\hat{\Delta}_S^2} + \frac{(2y-1)(r_t-1)}{\hat{\Delta}_F \hat{\Delta}_S} \right).$$

The differential cross section is

$$\frac{d\sigma}{dE_R} = \frac{m_N}{2\pi v^2} Z^2 e^2 b_\chi^2 F^2[E_R]. \quad (\text{A.7})$$

Other interactions are present but negligible. The coupling of χ_t to the Higgs is not important for direct detection because the coupling of the Higgs to nucleons is small. There is a scalar coupling to gluons, but this is only generated at subleading order in $1/m_\phi^2$ and in the limit of large m_ϕ^2 the contribution to direct detection is parametrically suppressed relative to Z -exchange by $(m_{\chi_t} m_n / m_\phi^2)^2$ [67]. A box diagram with χ_u in the loop can also generate a 4-fermion operator coupling χ_t to u quarks, but the amplitude is suppressed by $(\lambda_u)^2$ and the mass of χ_u .

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