

RECEIVED: November 15, 2022

REVISED: May 10, 2023

ACCEPTED: June 18, 2023

PUBLISHED: July 3, 2023

Neutrino masses and self-interacting dark matter with mass mixing $Z - Z'$ gauge portal

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ABSTRACT: New light gauge bosons can affect several low-energy experiments, such as atomic parity violation or colliders. Here, we explore the possibility that a dark sector is charged under a new $U(1)$ gauge symmetry, and the portal to the Standard Model is through a $Z - Z'$ mass mixing. In our approach, breaking the new gauge symmetry is crucial to generate neutrino masses. We investigate the parameter space to reproduce neutrino masses, the correct dark matter relic abundance, and to produce the observed core-like DM distribution in galactic centers.

KEYWORDS: Models for Dark Matter, New Gauge Interactions, New Light Particles

ARXIV EPRINT: [2210.14863](https://arxiv.org/abs/2210.14863)

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1 Introduction

The existence of dark matter and the non-zero neutrino masses are not explained within the Standard Model (SM) of particle physics. Dark matter (DM) is a component of the universe that accounts for $\sim 27\%$ of its total matter-energy density [1]. No particle, fundamental or composite, in the SM can account for it. A possible way to incorporate dark matter into the SM framework is to extend the gauge group $G_{\text{SM}} = \text{SU}(3)_C \times \text{SU}(2)_L \times \text{U}(1)_Y$ to include a dark gauge sector, under which the DM candidate is charged. One of the simplest ways to extend the SM is with an extra abelian gauge symmetry $U(1)_D$, which enlarges the gauge boson content of the SM. The dark sector of such a theory may communicate with the SM particles via the kinetic mixing term of the abelian subgroups $U(1)_Y$ and $U(1)_D$, the mass mixing among the neutral gauge bosons, or the scalar sector. The connection through the kinetic mixing is a popular and widely explored paradigm known as the *dark photon* [2–7], where the dark gauge boson acquires vector couplings to the SM fermions. In the mass mixing case [8, 9], the dark gauge boson acquires both vector and axial vector couplings to the SM fermions, leading to possible signatures in parity-violation experiments. This new gauge boson can provide a viable communication channel between the DM and the SM, leading to the correct dark matter relic density through the freeze-out while avoiding direct-detection constraints [10–12].

| | L | N | N' | F | H_1 | H_2 | ϕ | χ_L | χ_R^c |
|-----------|------|-----|------|-----|-------|-------|--------|----------|------------|
| $SU(2)_L$ | 2 | 1 | 1 | 1 | 2 | 2 | 1 | 1 | 1 |
| $U(1)_Y$ | -1/2 | 0 | 0 | 0 | 1/2 | 1/2 | 0 | 0 | 0 |
| $U(1)_D$ | 0 | 1 | -1 | 0 | 0 | 1 | -1 | Q_D | $-Q_D$ |

Table 1. Matter content of the dark matter with mass mixed $U(1)_D$ gauge boson. Quarks and right handed charged leptons, not shown in this table, are not charged under the $U(1)_D$ gauge symmetry. The model contains two dark charges, we have chosen to absorb one of them into the definition of the dark gauge coupling, leaving the dark matter charge Q_D free.

On the other hand, many neutrino mass generation mechanisms explain the lightness of Majorana neutrinos, compared to the rest of SM fermions. Among them, a popular class of models is the so-called *seesaw* mechanism, where a large mass scale suppresses the electroweak scale in the neutrino masses, giving rise to small neutrino masses. A popular type of these models requires the introduction of right handed (RH) neutrinos, singlets of the SM. Different mass models may be obtained, depending on the mass terms present in the lagrangian after the breaking of G_{SM} and any other additional symmetries in the model, namely, the type-I seesaw [13–17], the linear seesaw [18–20] or the inverse seesaw [21, 22]. Each of these models results in different possible values for the heavy neutrino masses and active-sterile neutrino mixings. In this work, we study the SM extended with an anomaly-free $U(1)_D$ gauge symmetry. The fermions, charged under the new gauge symmetry, will be identified with the right-handed neutrinos and the dark matter candidate. An extra Higgs doublet, charged under the dark gauge symmetry, generates a mass mixing among the dark gauge boson and the electroweak neutral gauge boson. The right-handed neutrinos’ and Higgs fields’ charges shape the neutrino seesaw matrix [23]. Dark matter is connected to the SM matter fields through the neutral gauge boson mass mixing, opening up viable thermal freeze-out channels and signatures in direct detection experiments.

2 The model

We consider the model for the mass mixing of a new gauge boson with the Z boson, described by table 1.¹ The dark sector (stable after the SSB) consists of a vector-like pair of fermions χ_L and χ_R . Fermions charged under $U(1)_D$ transform trivially under the SM gauge symmetry, guaranteeing the cancellation of all mixed SM-dark anomalies. For each charged fermion under $U(1)_D$, there is a fermion with an opposite charge, such that the pure $U(1)_D$ and the $U(1)_D$ -gravity anomalies vanish. The right-handed neutrinos N , N' , F , participate in the seesaw mechanism, with their $U(1)_D$ charges shaping the seesaw mass matrix. The scalar sector will induce mass mixing among the electroweak and dark neutral

¹Note that the charge assignment under the new $U(1)_D$ for N and N' could take values different value as far as they remain vector-like, but it is fixed by convenience. To have this, we add to the Standard Model(SM) a new dark gauge symmetry $U(1)_D$, a scalar $SU(2)_L$ doublet H_2 charged under the new symmetry, and a scalar singlet ϕ to trigger the symmetry breaking. The right handed neutrinos are charged under $U(1)_D$. Therefore, at least two different sets of vector-like fermions N and N' are needed to cancel the anomalies. We also need the inclusion of an extra fermion singlet under the new symmetry, F , to break the lepton number.

bosons, linking the SM fermions with the dark sector. The χ fields can act as a dark matter candidate, interacting with SM fields through the mass-mixed dark gauge boson. In this way, we show that the $U(1)_D$ can drive the phenomenology of neutrino and dark matter. The neutral gauge boson mixing will impact the quark and lepton physics, such as parity violation in polarized electron-nucleon and electron-electron scattering.

2.1 Neutrino sector

The RH neutrinos are charged under the $U(1)_D$. To generate the Yukawa Lagrangian, their charges must match that of the H_2 Higgs doublet. To avoid extra Goldstone bosons, in the scalar sector we must have a term such as $H_2 H_1^\dagger \phi$ or $H_2 H_1^\dagger \phi^2$. From these two conditions, we conclude that the charge of ϕ equals one of the RH neutrinos charges. The two RH neutrinos N and N' have a Dirac mass term. In contrast, there is no way to generate a Majorana mass for any of those fields through the ϕ field. The only way to do so is to include an extra fermion with no $U(1)_D$ charge. In this way, the Lagrangian density of the neutrino sector is

$$\mathcal{L}_\nu = Y_1^\nu \bar{L} \tilde{H}_1 F + Y_2^\nu \bar{L} \tilde{H}_2 N + M_1 \bar{N}^c N' + Y^N \bar{N}^c F \phi + Y^{N'} \bar{N}'^c F \phi^* + M_F \bar{F}^c F + h.c. \quad (2.1)$$

After spontaneous symmetry breaking (SSB) the resulting neutrino mass matrix in the (ν_L, N, N', F) basis is

$$M = \begin{pmatrix} 0 & m_2^D & 0 & m_1^D \\ (m_2^D)^T & 0 & M_1 & Y^N v_\phi \\ 0 & (M_1)^T & 0 & Y^{N'} v_\phi \\ (m_1^D)^T & (Y^N)^T v_\phi & (Y^{N'})^T v_\phi & M_F \end{pmatrix}, \quad (2.2)$$

where $m_i^D = Y_i^\nu v_i$ are the Dirac mass matrices. The light neutrino mass matrix is given by

$$m_{\text{light}} = (m_1^D)^T \alpha m_1^D + (m_2^D)^T \beta m_2^D + (m_1^D)^T \delta m_2^D + (m_2^D)^T \epsilon m_2^D, \quad (2.3)$$

where the $\alpha, \beta, \delta, \epsilon$ matrices are defined as

$$\begin{aligned} \alpha &= \left(M_F + Y^N (M_1^T)^{-1} (Y^{N'})^T v_\phi^2 + Y^{N'} (M_1)^{-1} (Y^N)^T v_\phi^2 \right)^{-1}, \\ \beta &= \left((Y^N)^T v_\phi \right)^{-1} + \left((Y^N)^T v_\phi \right)^{-1} M_1 (M_1^T)^{-1} \\ &\quad \times \left[Id + M_1^T (Y^N v_\phi)^{-1} \left(Y^{N'} v_\phi M_1^{-1} (Y^N v_\phi)^T - M_F \right) \left((Y^{N'} v_\phi)^T \right)^{-1} \right]^{-1}, \\ \delta &= - (M_1^T)^{-1} (Y^{N'} v_\phi)^T \left[M_F + Y^N (M_1^T)^{-1} (Y^{N'})^T v_\phi^2 + Y^{N'} (M_1)^{-1} (Y^N)^T v_\phi^2 \right]^{-1}, \\ \epsilon &= - (Y^N v_\phi)^{-1} \left[Y^{N'} v_\phi (M_1^T)^{-1} \left[Id + M_1^T (Y^N v_\phi)^{-1} \left(Y^{N'} v_\phi M_1^{-1} (Y^N v_\phi)^{-1} - M_F \right) \right. \right. \\ &\quad \left. \left. \times \left((Y^{N'} v_\phi)^T \right)^{-1} \right]^{-1} + M_F \beta \right], \end{aligned} \quad (2.4)$$

where Id is the Identity Matrix. The minimal field content leading to two massive light neutrinos is $(Nr(N) = 2, Nr(N') = 2, Nr(F) = 2)$. There are several familiar limits to this framework:

1. The type-I seesaw limit can be obtained when $Y^N, Y^{N'} \rightarrow 0$ or $Y^N, m_2^D \rightarrow 0$ or $Y^{N'}, m_2^D \rightarrow 0$. The magnitude of light neutrino masses is

$$m_\nu \sim \frac{v_1^2}{M_F}. \tag{2.5}$$

2. When $M_F, m_1^D \rightarrow 0$, the light neutrino masses take the form of the inverse seesaw

$$m_\nu \sim \frac{(m_2^D)^2 Y_{N'}}{M_1 Y_N}. \tag{2.6}$$

3. When $M_F, m_2^D \rightarrow 0$, the light neutrino masses take an inverse seesaw form

$$m_\nu \sim \frac{(m_D^1)^2 M_1}{v_\phi^2 Y_N Y_{N'}}. \tag{2.7}$$

4. When $Y_N, m_1^D \rightarrow 0$, the light neutrino masses take the form

$$m_\nu \sim \frac{(m_2^D)^2 (Y_{N'})^2 v_\phi^2}{M_1^2 M_F}. \tag{2.8}$$

In this case, there is an extra suppression compared with the inverse seesaw from the light (heavy) scale v_ϕ (M_F).

We will examine the viability of each limit, depending on the scale of v_ϕ indicated by DM phenomenology in section 3.5.

2.2 Dark sector

We choose as dark matter candidate a Dirac fermion $\chi = \chi_L + \chi_R$ with mass term

$$\mathcal{L}_\chi^{\text{mass}} = \frac{1}{2} M_\chi \bar{\chi} \chi. \tag{2.9}$$

We choose Q_D so that Majorana mass terms are forbidden at any order in perturbation theory, with the scalar content of table 1. The condition to keep the Dirac character of χ is

$$Q_D \neq \frac{m}{2}, \quad m \in \mathbb{Z}. \tag{2.10}$$

Since $Q_D \neq 0$, χ couples to the dark gauge boson X ; once the gauge symmetry is broken this induces a coupling to both the physical Z boson and the dark photon Z' , see eq. (2.18). For definiteness we will choose $Q_D = 1/3$ that satisfies eq. (2.10). A similar dark matter model is described in [10].

2.3 Gauge sector

The $U(1)_D$ charges of the new fields, except for χ , are equal in magnitude. Therefore, we may redefine the gauge coupling and field charges such that $Q = \pm 1$ for these fields. With this in mind, the kinetic terms of the scalar fields in the model defined in table 1 are

$$\mathcal{L}_{SK} = \sum_{i=1}^2 \left[(D_\mu H_i)^\dagger (D^\mu H_i) \right] + \left[(D_\mu \phi)^\dagger (D^\mu \phi) \right], \tag{2.11}$$

where the covariant derivatives for the $SU(2)_L$ Higgs doublets are

$$D_\mu H_i = \left(\partial_\mu + \frac{ig}{2} \vec{\tau} \cdot \vec{W}_\mu + \frac{ig'}{2} B_\mu + ig_D Q_i X_\mu \right) H_i, \quad (2.12)$$

with $Q_1 = 0$ and $Q_2 = 1$. The corresponding covariant derivative for the $SU(2)_L$ scalar singlet is

$$D_\mu \phi = (\partial_\mu - ig_D X_\mu) \phi. \quad (2.13)$$

After electroweak and dark symmetry breaking, $H_{1,2}$ and ϕ acquire vacuum expectation values and we write

$$H_1 = \begin{pmatrix} H_1^+ \\ (v_1 + h_1 + ia_1)/\sqrt{2} \end{pmatrix}; \quad H_2 = \begin{pmatrix} H_2^+ \\ (v_2 + h_2 + ia_2)/\sqrt{2} \end{pmatrix}; \quad \phi = v_\phi + h_\phi + ia_\phi. \quad (2.14)$$

The SM vacuum expectation value is $v_{\text{SM}}^2 = v_1^2 + v_2^2$, and we define $\tan \beta = v_2/v_1$. There are five vector bosons, $W^\pm = (W_1 \mp iW_2)/\sqrt{2}$ correspond to the usual charged pair with mass $gv_{\text{SM}}/2$, one neutral gauge boson, $A = s_w W_3 + c_w B$ (where $s_w = \sin(\theta_w)$ and $\tan(\theta_w) = g'/g$) remains massless. For the remaining fields, let $\tilde{Z} = c_w W_3 - s_w B$, then the mass matrix for $\{\tilde{Z}, X\}$ becomes

$$m_{\tilde{Z}X}^2 = \frac{1}{4} \begin{pmatrix} g_z^2 v_{\text{SM}}^2 & -2v_2^2 g_D g_z \\ -2v_2^2 g_D g_z & 4g_D^2 (v_2^2 + v_\phi^2) \end{pmatrix}, \quad (2.15)$$

where $g_z = \sqrt{g^2 + g'^2}$. The \tilde{Z} - X mixing angle, θ_X , is given by

$$\tan 2\theta_X = \frac{g_z g_D v_2^2}{\frac{1}{4} g_z^2 v_{\text{SM}}^2 - g_D^2 (v_2^2 + v_\phi^2)}. \quad (2.16)$$

Denoting the mass eigenstates by Z and Z' , the corresponding masses are given by

$$M_{Z/Z'}^2 = \frac{1}{8} g_z^2 v_{\text{SM}}^2 + \frac{1}{2} g_D^2 (v_2^2 + v_\phi^2) \pm \frac{1}{8} \left\{ \left[g_z^2 v_{\text{SM}}^2 - 4g_D^2 (v_2^2 + v_\phi^2) \right]^2 + \left(4g_z g_D v_2^2 \right)^2 \right\}^{1/2}. \quad (2.17)$$

The \tilde{Z} - X mixing induces a coupling of the Z' to the electroweak neutral current J_{NC} , and a coupling of the Z to the dark current J_{DC} proportional to $\sin \theta_X$; explicitly,

$$\mathcal{L}_{\text{NC}} = -e J_{\text{EM}}^\mu A^\mu - Z_\mu (\cos \theta_X \frac{g_z}{2} J_{\text{NC}}^\mu + \sin \theta_X g_D J_{\text{DC}}^\mu) - Z'_\mu (-\sin \theta_X \frac{g_z}{2} J_{\text{NC}}^\mu + \cos \theta_X g_D J_{\text{DC}}^\mu), \quad (2.18)$$

where $e = gs_w$. The currents in eq. (2.18) are

$$\begin{aligned} J_{\text{EM}}^\mu &= \sum_r Q_r^{\text{EM}} \bar{f}_r \gamma^\mu f_r, & J_{\text{NC}}^\mu &= \sum_r t_L^3(r) \bar{f}_r \gamma^\mu (1 - \gamma^5) f_r - 2s_W^2 J_{\text{EM}}^\mu, \\ J_{\text{DC}}^\mu &= \frac{g_D}{3} \bar{\chi} \gamma^\mu \chi + g_D (\bar{N} \gamma^\mu N - \bar{N}' \gamma^\mu N'), \end{aligned} \quad (2.19)$$

where Q_r^{EM} is the EM charge of the f_r fermion and $t_L^3(r)$ its weak isospin.

2.4 Scalar sector

The scalar potential for the model in table 1 is given by

$$\begin{aligned}
 V_i = & \mu_1^2 H_1^\dagger H_1 + \mu_2^2 H_2^\dagger H_2 + \mu_\phi^2 \phi^* \phi + \kappa \phi^* H_1^\dagger H_2 + \lambda_1 (H_1^\dagger H_1)^2 + \lambda_2 (H_2^\dagger H_2)^2 + \lambda_3 (\phi^* \phi)^2 \\
 & + \lambda_4 (H_1^\dagger H_1)(H_2^\dagger H_2) + \lambda_5 (H_2^\dagger H_2)(\phi^* \phi) + \lambda_6 (H_1^\dagger H_1)(\phi^* \phi) + \lambda_7 (H_1^\dagger H_2)(H_2^\dagger H_1),
 \end{aligned}
 \tag{2.20}$$

where the only complex coupling is κ , however by a field redefinition it can be made real. From the 10 real scalar degrees of freedom, four goldstone bosons are absorbed in the vector boson masses; the remaining six correspond to a charged pair H^\pm , a pseudoscalar A , and three neutral scalars. Using the notation of eq. (2.14) the first three and their masses are given by

$$\begin{aligned}
 H^+ & \sim \sin \beta H_1^+ - \cos \beta H_2^+, & M_{H^+}^2 & = \frac{1}{2} s v_\phi^2 - \frac{1}{2} \lambda_7 v_{\text{SM}}^2, \\
 A & \sim v_\phi \sin \beta a_1 - v_\phi \cos \beta a_2 + \frac{1}{2} v_{\text{SM}} \sin(2\beta) a_\phi, & M_A^2 & = \frac{1}{2} s v_\phi^2 + \frac{\sin^2(2\beta)}{8} s v_{\text{SM}}^2,
 \end{aligned}
 \tag{2.21}$$

while in the $\{h_1, h_2, h_\phi\}$ basis the CP-even mass matrix is given by

$$M_E^2 = \begin{pmatrix} 2\lambda_1 v_1^2 - \frac{\kappa v_\phi}{\sqrt{2}} \tan \beta & v_1 v_2 (\lambda_4 + \lambda_7) + \frac{\kappa v_\phi}{\sqrt{2}} & v_1 v_\phi \lambda_6 + \frac{\kappa v_2}{\sqrt{2}} \\ v_1 v_2 (\lambda_4 + \lambda_7) + \frac{\kappa v_\phi}{\sqrt{2}} & 2\lambda_2 v_2^2 - \frac{\kappa v_\phi}{\sqrt{2}} \cot \beta & v_2 v_\phi \lambda_5 + \frac{\kappa v_1}{\sqrt{2}} \\ v_1 v_\phi \lambda_6 + \frac{\kappa v_2}{\sqrt{2}} & v_2 v_\phi \lambda_5 + \frac{\kappa v_1}{\sqrt{2}} & 2\lambda_3 v_\phi^2 - \frac{\kappa v_1 v_2}{\sqrt{2} v_\phi} \end{pmatrix}.
 \tag{2.22}$$

To simplify the expressions we defined

$$s = -\frac{\sqrt{8} \kappa}{v_\phi \sin(2\beta)},
 \tag{2.23}$$

which is positive since in our conventions $\kappa < 0$. The scalar potential must be bounded from below, leading to restrictions on the scalar couplings. We have collected these restrictions in appendix A. We note the existence of a decoupling limit, where the scalar masses become much heavier than the electroweak scale, save for the Higgs seen at LHC. This limit is achieved when $v_2 \rightarrow 0$ and $\kappa \rightarrow 0$, with μ_2^2 setting the heavy scalar scale. The decoupling limit drives the gauge boson mixing to zero, making the Z' and dark matter invisible.

3 Phenomenology

3.1 General constraints on light dark Z' bosons

The induced couplings of the dark Z' boson to the SM fermions can be probed by a variety of experiments [8, 9, 24–27]. The observables measured by those experiments, constrain the parameter space in the $\theta_{\mathbf{x}} - M_{Z'}$ plane. The most relevant experimental constraints are the following:

- Atomic Parity Violation. As noted above, the mass mixing among the SM and X bosons induces a Z' couplings to SM fermions (cf. eq. (2.19)), which inherit the parity

violating nature of the SM Z couplings. The Z' parity violating couplings may induce observable effects on low energy experiments, when the mass of the Z' is comparable to the energy scale of the experiment [28–30]. This parity-violating interaction of quarks and leptons mediated by the Z' has been probed in atomic transitions of Yb, Cs, Tl, Pb, and Bi. The measurement of the nuclear weak charge in these experiments can be used to constrain the Z' couplings to the SM fermions as a function of its mass [8]. The resulting constraint is approximately [26, 27, 31]

$$\sin\theta_x \lesssim 5 \times 10^{-5}, \text{ for } M_{Z'} < 40 \text{ MeV}; \quad \text{and} \quad \frac{\sin\theta_x}{M_{Z'}} \lesssim \frac{10^{-6}}{\text{MeV}}, \text{ for } 40 \text{ MeV} < M_{Z'} < 100 \text{ GeV}. \quad (3.1)$$

- Collider searches. As the Z' couples to the SM fermions, it can be produced in a multitude of collider experiments. The Z' -mediated Drell-Yan production of muons in hadron colliders yield some of the strongest constraints on the Z' couplings for masses below M_Z . Neutral gauge boson production and decay to leptons, in association with photon production has been searched for in e^+e^- collisions, at BaBar [32], CMS [33], LHCb [34], among others. The constraint from colliders in the mass region, $100 \text{ MeV} < M_{Z'} < 80 \text{ GeV}$, is roughly [26, 27]

$$\sin\theta_x \lesssim 5 \times 10^{-3}. \quad (3.2)$$

- Beam dump experiments. The production of neutral bosons in electron bremsstrahlung processes in beam dumps has been probed, for example at the NA64 [35], E141 [36], E137 [37] and E774 [38] experiments. The beam dump limits on $\sin\theta_x$ in the Z' gauge boson mass region, $1 \text{ MeV} < M_{Z'} < 500 \text{ MeV}$ are roughly [26, 27]

$$\sin\theta_x \lesssim 3 \times 10^{-8} \text{ or } \frac{\sin\theta_x}{(M_{Z'}/\text{GeV})^{-1.2}} \gtrsim 3.3 \times 10^{-7}. \quad (3.3)$$

Using the DarkCast code [26, 27], we derive the constraints from the APV, Collider and beam dump experiments mentioned. We show this results in figure 4.

3.2 Dark matter relic density

We consider the thermal freeze-out to determine the DM relic density, Ω_χ . We identify three scenarios which can result in a relic density of dark matter in accordance with cosmological measurements, $\Omega_\chi h^2 \leq 0.1198$:

1. When $M_\chi > M_{Z'}$, the annihilation t-channel, $\bar{\chi}\chi \rightarrow Z'Z'$, is kinematically allowed. This leads to a relic density which only depends on the Z' boson mass $M_{Z'}$, the dark gauge coupling g_D , and the dark matter mass M_χ . Numerically we find that for each value of $M_{Z'}$ there is a minimum value of g_D for which there is no dark matter overabundance. We show this behavior in figure 2. This scenario is well-studied and is known in the literature as Secluded WIMP Dark Matter [39–41].
2. When $M_\chi > M_f$, f being a SM fermion, the Z - Z' mixing allows the s-channel annihilation of dark matter into f , $\bar{\chi}\chi \rightarrow \bar{f}f$. In the resonant regions $M_\chi \sim M_{Z'}/2$

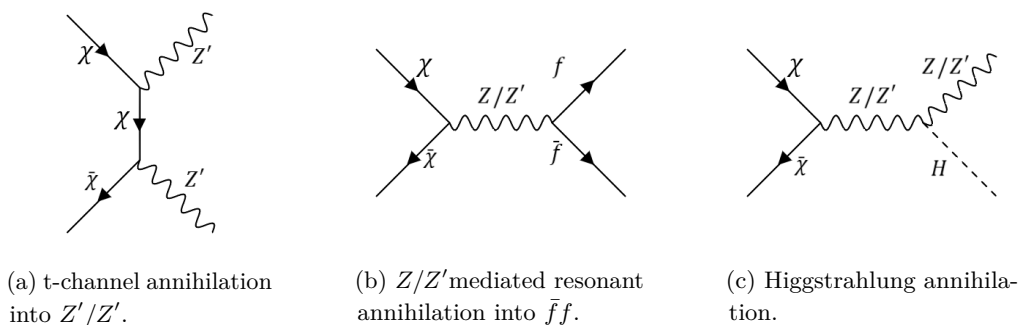


Figure 1. Dark matter annihilation channels for the determination of the freeze-out relic density.

the annihilation cross section can be enhanced enough to reach the required value to result in an allowed relic density, while keeping the value of θ_χ in the allowed region discussed in section 3.1 [42]. In this channel in addition to $M_{Z/Z'}$, M_χ and g_D , the Z - Z' mixing angle θ_χ is a crucial parameter.

3. When $2M_\chi > M_{Z(Z')} + M_S$, where S is one of the four neutral scalars in the model, the s-channel Higgsstrahlung channel ($\bar{\chi}\chi \rightarrow SZ(Z')$) is kinematically allowed. In this channel, the scalar masses and mixing angles become relevant to the relic density calculation. For dark matter masses above the W boson mass, the channel $\bar{\chi}\chi \rightarrow WW$ is open and can contribute significantly to the dark matter relic density.

We illustrate the tree-level Feynman diagrams of these processes in figure 1. To calculate the relic density for this model, we have implemented the model in SARAH [43] and micrOmegas [44, 45], scanning over a range of the free parameters. We study the first two cases, as they lead to an interesting interplay between the dark sector, the light gauge boson parameters and the neutrino sector.

3.3 Dark matter direct detection

After excluding the parameter space where the relic density of χ does not correspond to that of dark matter, we look at the Spin-Independent cross section of dark matter with nucleons. The $Z - Z'$ mass mixing leads to tree-level dark matter-nucleon elastic scattering, a process searched for in direct detection experiments. The scattering is mediated by Z and Z' exchange.

In the case where relic density is determined by t-channel $\bar{\chi}\chi \rightarrow Z'Z'$ annihilations, there is no correlation between dark matter relic density and the direct detection cross section, as the relic annihilation cross section is independent of the $Z - Z'$ mixing angle at leading order. However, a bound on the $Z - Z'$ mixing angle may be derived from the limits on this cross section. For the $M_{Z'}^2 \ll M_Z^2$ limit, the spin independent χ -nucleus elastic scattering cross section is approximately [49]

$$\sigma_{\text{SI}}^{Z,A} = \frac{\mu_{\chi N}^2 \sin^2 2\theta_\chi g_X^2 Q_\chi^2}{4\pi M_{Z'}^4} \left[Z(2g_{\text{SM}}^{Zu} + g_{\text{SM}}^{Zd}) + (A - Z)(g_{\text{SM}}^{Zu} + 2g_{\text{SM}}^{Zd}) \right]^2, \quad (3.4)$$

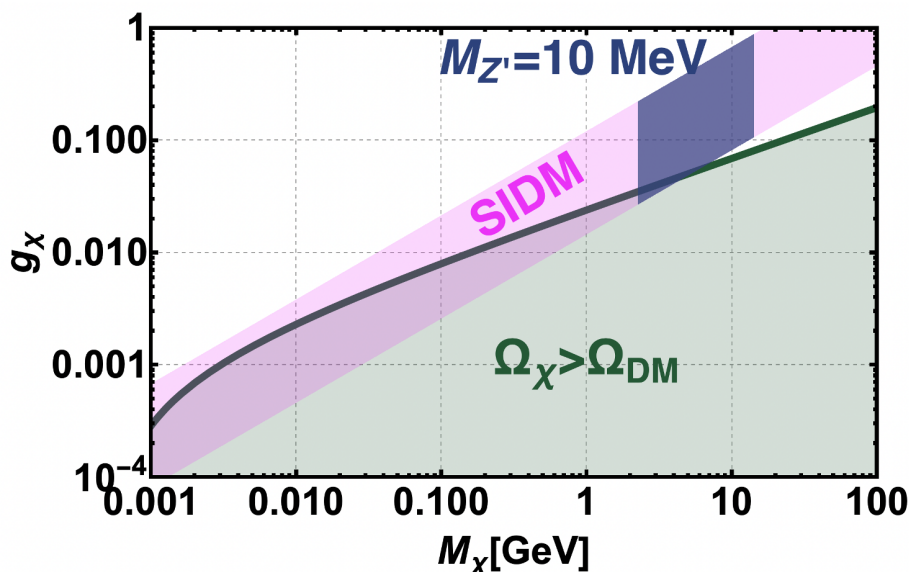


Figure 2. Case 1: t-channel annihilation into a Z' pair. Parameter space in the $M_\chi - g_\chi$ plane (where $g_\chi = g_D \cos \theta_\chi \approx g_D$) excluded by relic density overabundance. Area shaded in green results in relic density overabundance, while the green line corresponds to $\Omega_{\text{DM}} = 0.1195$. The area indicated in Magenta with the label “SIDM” corresponds to the region where dark matter self-interactions are consistent with astrophysical observations. Within this area, the blue region corresponds to a mediator mass $M_{Z'}$ of 10 MeV, considering the uncertainty as estimated in the text.

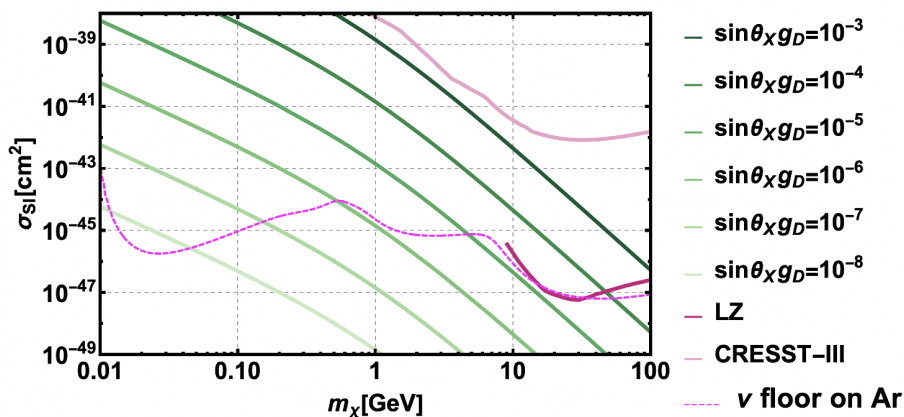


Figure 3. Spin Independent direct detection cross section (σ_{SI}) as a function of dark matter mass, for the Z' resonant case (s-channel). In the green lines the product $\sin \theta_\chi g_D$ is fixed to values between 10^{-8} and 10^{-3} as indicated in the caption. The purple lines show the experimental limit on σ_{SI} set by the LZ [46] and CRESST-III [47] experiments. The dashed magenta line shows the neutrino floor on Argon [48]. Note that the LZ experiment target nucleus is Xenon, and it has not yet reached the Xenon neutrino floor.

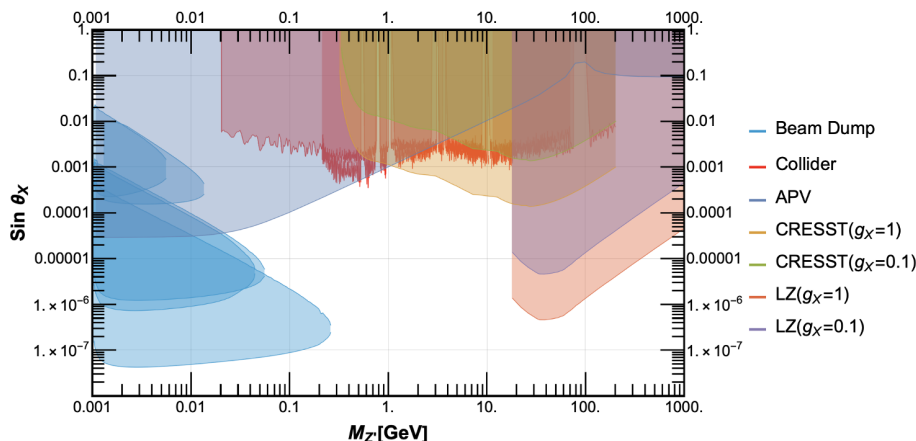


Figure 4. Low energy constraints and resonant dark matter direct detection constraints. In this figure we show the low energy constraints from Atomic Parity Violation [28], collider (BaBar [32], CMS [33] and LHCb [34]) and beam dump (NA64 [35], E141 [36], E137 [37], E774 [38], KEK [50] and Orsay [51]) experiments on the $M_{Z'}$ – $\sin \theta_X$ parameter space. These constraints are obtained as discussed in section 3.1. We also show in the same space the constraints derived from direct detection experiments (LZ [46] and CRESST-III [47]) in the resonant dark matter models, as discussed in section 3.3.

where $\mu_{\chi N}^2$ is the reduced χ -nucleus mass, g_{SM}^{Zq} is the vector coupling of $q = u, d$ to the Z boson in the SM, and Z and A are the electric charge and atomic mass number of the nucleon respectively. Using the most stringent limits on dark matter direct detection we can list the following constraints

- $\frac{1}{2} |g_X Q_\chi \sin 2\theta_X| \lesssim 10^{-7}$ at $M_{Z'} = 10$ GeV, and $M_\chi = 30$ GeV from LZ [46].
- $\frac{1}{2} |g_X Q_\chi \sin 2\theta_X| \lesssim 10^{-6}$ at $M_{Z'} = 100$ MeV, and $M_\chi = 1$ GeV from CRESST-III [47].

For the resonant Z' channel case we observe a clear correlation between the parameter product $g_D \sin \theta_X$ and the SI cross section in figure 3. For dark matter masses above 1 GeV and below 10 GeV, CRESST-III results exclude models where $g_D \sin \theta_X > 10^{-4}$. For dark matter masses above 10 GeV LZ results exclude models with $g_D \sin \theta_X > 10^{-6}$ up to 30 GeV, after which the constraint relaxes. In figure 4 we show direct detection constraints alongside Z' searches constraints, for a choice of $g_D = 1, 0.1$. We see that for Z' masses above 20 GeV, direct detection constraints are stronger than low energy constraints for these choices of g_D . For lighter Z' masses CRESST-III results are stronger than Z' searches, down to 1 GeV, with $g_D = 1$. For $g_D = 0.1$ CRESST-III results are weaker than collider or APV constraints.

3.4 Dark matter self-interactions

Estimations of the DM distribution in dwarf galaxies indicate that the DM density at the core does not exhibit a spike, as would be expected if it behaved as an ideal gas; this is known as the “core vs. cusp” problem [52–54]. This problem can be alleviated [55, 56]

by including self-interactions within the dark sector;² such interactions must be relatively strong and velocity-dependent. Models of this type are often referred to as self-interacting dark matter (SIDM) models. In this section we determine the extent to which the present model can address the core-vs-cusp problem.

Existing data constraints the SIDM cross section for galaxy clusters and for dwarf and low-surface-brightness galaxies; since the typical velocity in each environment is different, the cross section must have an appropriate velocity-dependence.

Dark matter self-interactions in this model consist of $\chi\chi \rightarrow \chi\chi$ scatterings mediated by the Z' boson. Defining

$$\beta_\chi = \sqrt{1 - \frac{4M_\chi^2}{s}}, \quad g_\chi = g_D \cos \theta_\chi, \quad (3.5)$$

we obtain the self-interaction cross section σ_{SIDM} [59]

$$\frac{\sigma_{\text{SIDM}}}{M_\chi} = \frac{g_\chi^4}{4\pi s M_\chi} \left\{ \frac{(2s + 3M_{Z'}^2)s\beta_\chi^2 + 2(M_{Z'}^2 + 2M_\chi^2)^2}{2M_{Z'}^2(M_{Z'}^2 + s\beta_\chi^2)} - \frac{(s\beta_\chi^2 + 2M_{Z'}^2)(3M_{Z'}^2 + 4M_\chi^2) + 2(M_{Z'}^2 + 2M_\chi^2)^2 - 4M_\chi^4 \ln\left(1 + \frac{s\beta_\chi^2}{M_{Z'}^2}\right)}{s\beta_\chi^2(2M_{Z'}^2 + s\beta_\chi^2)} \right\}; \quad (3.6)$$

The requirements for SIDM are met in this model when this cross section is enhanced by the kinematic condition $M_\chi \gg M_{Z'}$, in the small relative velocity β_χ regime. The SIDM cross section magnitude is determined by astrophysical data, namely dwarf and low surface brightness galaxies. A velocity-dependence of the cross section is obtained by the different typical velocities of dark matter in each environment. The central values of the cross sections and velocities are [60]

$$\frac{\sigma_{\text{SIDM}}}{M_\chi} \Big|_{\text{galaxy}} = 1.9 \frac{\text{cm}^2}{\text{gr}}, \quad \frac{\sigma_{\text{SIDM}}}{M_\chi} \Big|_{\text{cluster}} = 0.1 \frac{\text{cm}^2}{\text{gr}}; \quad \beta_\chi|_{\text{galaxy}} = 3.3 \times 10^{-4}, \quad \beta_\chi|_{\text{cluster}} = 5.4 \times 10^{-3}. \quad (3.7)$$

The galaxy data were derived from dwarf and low surface brightness galaxies; the estimates of the average collision velocity β_χ were obtained using a generic halo model and a Maxwellian velocity distribution. Fitting to these values we find

$$M_{Z'} = \frac{M_\chi}{566}, \quad g_\chi = \left(\frac{M_\chi}{75 \text{ GeV}} \right)^{3/4}. \quad (3.8)$$

The errors in the data (eq. (3.7)) are large; to take these into account we allow a factor of 2.5 in the numerical coefficients in eq. (3.8), (e.g., that the first coefficient ranges from 566/2.5 to 2.5 * 566). We find that these conditions can be met when dark matter relic density is obtained through annihilations to Z' pairs. In figure 2 we show the band where

²There are other puzzles associated with the standard cold-DM scenario, such as the “too big to fail” puzzle [57, 58]. In this paper we will not consider the extent to which the present model can address such issues.

SIDM is viable. We note that there is an overlap between the SIDM band and the line where χ accounts for dark matter completely. In this scenario, the mass of the Z' is of order ~ 10 MeV, which is where low energy experiments are most sensitive.

3.5 Neutrino masses and $U(1)_D$ breaking scale

From the constraints on the dark sector parameters from the dark matter phenomenology, we can infer the following possibilities for the neutrino mass mechanism.

For all dark matter relic density channels, the inverse seesaw-like scenario

$$m_\nu \sim \frac{(m_D^2)^2 Y_{N'}}{M_1 Y_N}, \tag{3.9}$$

and the type-I seesaw scenario

$$m_l \sim \frac{v_1^2}{M_F}, \tag{3.10}$$

are viable with heavy neutrinos of canonical seesaw scale.

Of special interest are the cases where the neutrino masses are linked to the $U(1)_D$ breaking scale, namely the limits considered in eqs. (2.7) and (2.8). In the limit considered in eq. (2.8)

$$m_\nu \sim \frac{(m_D^2)^2 Y_{N'} v_\phi^2}{M_1^2 M_F}, \tag{3.11}$$

a light neutrino masses of the order $\mathcal{O}(\text{eV})$ can be achieved with $\mathcal{O}(\text{TeV})$ heavy neutrinos, as in the canonical inverse seesaw. In this way the smallness of neutrino masses is linked to the smallness of the Z' boson mass.

Lets consider now the limit in eq. (2.7). For the t-channel annihilation case, the correct relic density is determined for values of g_D larger than 3×10^{-4} for a dark matter mass of ~ 1 MeV, or g_D larger than 2×10^{-1} for a dark matter mass of ~ 100 GeV (see figure 2). In the light Z' paradigm, with small $Z - Z'$ mixing we have

$$M_{Z'} > g_D v_\phi, \tag{3.12}$$

and the t-channel dominated annihilation has the kinematic condition

$$M_\chi > M_{Z'}. \tag{3.13}$$

These two equations rule out the inverse seesaw-like limit of neutrino masses

$$m_\nu \sim \frac{(m_D^1)^2 M_1}{v_\phi^2 Y_N Y_{N'}}, \tag{3.14}$$

as the low scale of v_ϕ needed to obtain a light $M_{Z'}$ would result in either ~ 1 GeV scale sterile neutrinos with large mixings with the active neutrinos, or neutrino Yukawa couplings much smaller than the electron Yukawa coupling of the SM, calling into question the necessity of the seesaw scheme.

4 Conclusions

Abelian gauge extensions of the SM with light gauge mediators have become popular recently because they lead to interesting implications in low energies experiments and observables. In this work we have studied a scenario where the coupling of such a light gauge boson with the SM is generated through a mass mixing with the Z boson. We extended the model by including a DM fermion charged under the new $U(1)_D$ symmetry which is automatically stable without the inclusion of extra symmetries. In this way the DM phenomenology further constrains the new gauge sector. The DM self-interactions are also discussed and we present the parameter space where the cusp-core problem is resolved. Focusing on the mass region for the Z' gauge boson where low energy parity violation experiments are more sensitive, $M_Z \sim 10$ MeV we have found the parameter space for the dark gauge coupling, g_D and mass mixing parameter to reproduce the DM relic abundance and search for the region where the DM self-interaction explain the approximately constant dark matter density in the inner parts of galaxies. We have found an overlap to explain correctly the DM relic abundance and the DM self-interaction in the DM mass region $\sim 2 - 7$ GeV with g_D above ~ 0.01 . We also identify a scenario where there is no significant DM self-interactions, but DM can be seen in direct detection experiments. The neutrino masses are generated through the seesaw mechanism. In some cases, the $U(1)_D$ breaking scale plays a crucial role in the neutrino mass generation. In these scenarios it is possible to generate neutrino masses through a low-energy seesaw.

A Tree-level stability of the potential

Let

$$x_1 = \sqrt{2}|H_1|^2, \quad x_2 = \sqrt{2}|H_2|^2, \quad x_3 = \sqrt{2}|\phi|^2, \quad (\text{A.1})$$

then the quartic part of the potential takes the form

$$V_4 = \frac{1}{2} \sum_{i=1}^3 \lambda_i x_i + \eta_3 x_1 x_2 + \eta_1 x_2 x_3 + \eta_2 x_3 x_1, \quad (\text{A.2})$$

where $\eta_1 = \lambda_5/2$, $\eta_2 = \lambda_6/2$, $\eta_3 = (\lambda_4 + \zeta^2 \lambda_7)/2$ with $\zeta = |H_1^\dagger H_2|/(|H_1||H_2|)$ (note that $0 \leq \zeta \leq 1$). The tree level stability conditions are then: $\lambda_i > 0$ and $\eta_i > -\sqrt{\lambda_j \lambda_k}$ with $\{i, j, k\}$ a cyclic permutations of $\{1, 2, 3\}$; in addition,

- If $\eta_i > 0$, $\eta_j, \eta_k < 0$, then $\lambda_i \eta_i > \eta_j \eta_k - \sqrt{(\lambda_i \lambda_j - \eta_k^2)(\lambda_i \lambda_k - \eta_j^2)}$.
- If $\eta_{i,j,k} < 0$ then $\lambda_1 \lambda_2 \lambda_3 + 2\eta_1 \eta_2 \eta_3 > \lambda_1 \eta_1^2 + \lambda_2 \eta_2^2 + \lambda_3 \eta_3^2$.

Acknowledgments

The authors thank Luis J. Flores for kindly providing the neutrino floor data and R. Ferro-Hernandez for useful discussions. LMGDLV thanks CONACYT for the funding of his PhD studies. This work has been supported by the University of California Institute for Mexico and the United States (UC MEXUS) (CN 18-128) and the Consejo Nacional de Ciencia

y Tecnología (CONACYT) (CN 18-128), by the German-Mexican research collaboration grant SP 778/4-1 (DFG) and 278017 (CONACYT), the CONACYT CB-2017-2018/A1-S-13051, and DGAPA-PAPIIT IN107621. EP is grateful to funding from ‘Cátedras Marcos Moshinsky’ (Fundación Marcos Moshinsky).

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