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(In)visible Z' and dark matter

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Abstract: We study the consequences of an extension of the standard model containing an invisible extra gauge group under which the SM particles are neutral. We show that effective operators, generated by loops of heavy chiral fermions charged under both gauge groups and connecting the new gauge sector to the Standard Model, can give rise to a viable dark matter candidate. Its annihilations produce clean visible signals through a gamma-ray line. This would be a smoking gun signature of such models observable by actual experiments.

Keywords: Beyond Standard Model, Cosmology of Theories beyond the SM, Anomalies in Field and String Theories

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Contents

1 Introduction

Extensions of the SM with additional Z' symmetries were widely discussed in the literature [1]. There are several ways of defining semi-invisible Z' theories, under which SM fermions are neutral, with effective operators connecting directly Z' to the SM sector. The simplest one can imagine is connecting the two sectors by a kinetic mixing operator

$$\delta \cdot \; F_X^{\mu\nu} \; F_{\mu\nu}^Y \,, \tag{1.1}$$

with $F_{\rho\sigma}^I = \partial_\rho B_\sigma^I - \partial_\sigma B_\rho^I$, which is naturally generated at the loop-level by heavy fermions charged both under SM and the Z' gauge symmetry. This scenario is very interesting from a dark matter perspective and was studied intensively over the last years [2–4]. The second and less studied is the effect of the vertex connecting Z' to two SM gauge fields. The simplest examples are the generalized Chern-Simons (GCS) terms [5–7], for example

$$\epsilon^{\mu\nu\rho\sigma} Z'_\mu B_\nu F_{\rho\sigma}^Y \,, \tag{1.2}$$

which were argued to be generated by heavy fermions in [5].

In this letter we study from an effective operator viewpoint the GCS terms of the type (1.2) and their decoupling properties at low-energy. We argue that the $Z'Z\gamma$ and $Z'WW$ vertices necessarily have a heavy-fermion mass suppression $1/M^2$. This is due to the fact that, if the heavy fermion masses M which are decoupled ($M \to \infty$) are SM symmetric, coming only from the Higgs mechanism breaking the Z' gauge symmetry, the operator (1.2) should be invariant under the non-linearly (in the broken-phase) realized Z' symmetry, whereas it should be well-defined (non-singular) in the unbroken SM phase. The gauge invariant version of (1.2) is then

$$\frac{i}{M^2} \, \epsilon^{\mu\nu\rho\sigma} \mathcal{D}_\mu \theta_X (H^+ D_\nu H - (D_\nu H)^+ H) \; F_{\rho\sigma}^Y \tag{1.3}$$

where D_ν is the generic covariant derivative, $\mathcal{D}_\mu$ defined as $\mathcal{D}_\mu \theta_X = \partial_\mu \theta_X - g_X Z'_\mu$ is the Stueckelberg gauge-invariant combination with θ_X being the Stueckelberg axion, g_X is the Z' gauge coupling and M is related to the heavy fermion masses. After electroweak symmetry breaking, we get a GCS term (1.2) with a coeff. prop. to v^2/M^2, where v is the electroweak vev.

In the case of two Z' gauge bosons, we show that there is a genuine non-decoupling effect, i.e. independent of the heavy-fermion masses.[1] Indeed, for two extra gauge symmetries $\mathrm{U}(1)_X$ and $\mathrm{U}(1)'_X$ with Stueckelberg realization of gauge symmetries , the dimension-four operator

$$\epsilon^{\mu\nu\rho\sigma} \mathcal{D}_\mu \theta_X \mathcal{D}_\nu \theta'_X F^Y_{\rho\sigma} \,, \tag{1.4}$$

where θ_X and θ'_X are the axions of the two Z' respectively, is gauge invariant and can be generated by a heavy chiral but anomaly-free fermion spectrum charged under the two U(1)'s and the SM. We will check this explicitly in section 4 by using the formulae of ref. [5]. The masses of the heavy fermions that are taken to infinity $M \to \infty$ have to come from the Higgs breaking of the two U(1)'s and has to be SM invariant. In both cases, of one or two Z', we allow possible SM-like mass elements $m \sim v$ for the heavy fermions, which we keep fixed in the decoupling limit $M \to \infty$. The term (1.4) provides an interesting counter-example of the decoupling theorem [14]. This is different from the non-decoupling effects studied in [15] for two reasons. First, the heavy fermionic spectrum we will consider, albeit chiral, is free of any gauge anomalies. Secondly, whereas in [15] the fermions which were decoupled had a SM-like mass, in the case studied here the masses which decouple are SM invariant.

The GCS terms were already discussed from the viewpoint of anomalous three-gauge boson vertices, notably $Z'Z\gamma$ in various papers [5, 8, 10]. The consequences for the LHC were subsequently analyzed in [8–10] and [11]. In ref. [12] the supersymmetric partner of the axion, the axino, was discussed as a dark matter candidate. The main point of the present paper is that, whereas the mass suppression in (1.3) make the LHC signatures of such an (in)visible Z' difficult to detect, the interactions described by the dimension-six operators (1.3) and others similar to it, discussed in more detail in section 2, make the lightest fermion in the Z' sector a viable dark matter candidate. Indeed, due to the couplings (1.2) such a fermion can annihilate into a Z and a photon, via the s-channel Z' virtual exchange, with an appropriate relic abundance. An interesting signature of this channel of dark-matter annihilation is the gamma ray in the final state, which is monochromatic and can be tested with the FERMI/GLAST [13] experiment in the near future. We would like to emphasize that the reason this (loop-suppressed) coupling can produce a visible gamma ray signal is that in our case it is the same diagram which describes the main annihilation channel for the dark matter and simultaneously generate the monochromatic gamma ray. In contrast, in other BSM models (e.g supersymmetry), the $Z'Z\gamma$ vertex is loop suppressed, whereas the dark matter annihilation occurs at tree-level, making the gamma ray signal highly suppressed. We argue that if the vertex (1.2)

[1]see [11] for a recent similar example.

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dominates[2] or gives similar effects compared to the kinetic mixing (1.1), the monochromatic gamma ray remains visible, providing an astrophysical window towards high-energy physics. For the two Z' case, the unsuppressed coupling (1.4) leads to an enhancement of the diagram producing the relic density and the monochromatic gamma ray, which now is generated by the anomalous coupling $Z'Z''\gamma$, provided the kinematic constraint $M_{Z''} < 2M_{\rm DM}$ is fulfilled. For TeV scale DM mass, both Z' and Z" should then have TeV masses, in order to get good relic density via the DM annihilation into $Z''\gamma$.

2 (In)visible Z', effective operators and decoupling

We consider here an effective model where a left and right dark matter fermion $\psi_L^{\rm DM}$, $\psi_R^{\rm DM}$, charged under a spontaneously broken extra $\mathrm{U}(1)_X$, with charges $X_L^{\rm DM}$ and $X_R^{\rm DM}$ respectively, is added to the Standard Model sector. If the $\mathrm{U}(1)_X$ is invisible to the SM, i.e if quarks and leptons are neutral with respect to the extra gauge symmetry, the only way Z' can contribute to the low-energy physics, is through effective interactions obtained after integrating out the UV physics. Typical examples of this kind of effects are given by considering a heavy sector of fermions charged under $\mathrm{U}(1)_X$ and the SM gauge group, that we briefly discuss later on. When the heavy fermions decouple, loop effects give rise to the general effective Lagrangian:

$$\begin{aligned}\mathcal{L} = &\mathcal{L}_{\rm SM} + \bar{\psi}_L^{\rm DM}\left(i\gamma^\mu\partial_\mu + g_X X_L^{\rm DM}\gamma^\mu Z'_\mu\right)\psi_L^{\rm DM} + \bar{\psi}_R^{\rm DM}\left(i\gamma^\mu\partial_\mu + g_X X_R^{\rm DM} Z'_\mu\right)\psi_R^{\rm DM} \\ &- \left(\bar{\psi}_L^{\rm DM} M_{\rm DM}\psi_R^{\rm DM} + \mathrm{h.c.}\right) + \frac{1}{2}(\partial_\mu a_X - M_{Z'}Z'_\mu)^2 - \frac{1}{4}F^X_{\mu\nu}F^{X\,\mu\nu} \\ &+\mathcal{L}_1(Z'_\mu) + \mathcal{L}_2(B_\mu, W^a_\mu) + \mathcal{L}_{\rm mix}(Z'_\mu, B_\mu, W^a_\mu)\,, \end{aligned} \tag{2.1}$$

where $\mathcal{L}_{\rm SM}$ is the Standard Model Lagrangian, $\mathcal{L}_1$ and $\mathcal{L}_2$ represent the new effective operators generated separately in the SM gauge sector and Z' one, whereas in $\mathcal{L}_{\rm mix}$ we collect all the induced terms mixing them. The aim of this paper is to show how these terms give the possibility to detect the presene of an "invisible" Z'. Z' gauge symmetry is spontaneously broken by the vev of a Higgs field S. The Stueckelberg axion a_X assures the gauge invariance of the effective action, and g_X and F^X are the Z' gauge coupling and gauge field strength. The Stueckelberg mechanism can be understood as a heavy Higgs mechanism, where the extra Higgs field S takes the form $S = (V+s)\ \exp[i\frac{a_X}{V}]$, where V is the heavy Higgs vev, and the axion transforms non-linearly under $\mathrm{U}(1)_X$ gauge transformations

$$\delta A_X^\mu \ = \ \partial^\mu\alpha \quad , \quad \delta a_X \ = \ \alpha\, g_X V\,, \tag{2.2}$$

for a Higgs field S of X-charge equal to 1.

The lightest fermion charged (only) under Z' will be our dark matter candidate. The dark matter mass $M_{\rm DM}$ can be of two types:

- If $\Psi_L^{\rm DM}$ and $\Psi_R^{\rm DM}$ have equal charges $X_L^{\rm DM} = X_R^{\rm DM}$, then $\Psi^{\rm DM}$ is vector-like and therefore we can write the Dirac mass $M\bar{\Psi}^{\rm DM}\Psi^{\rm DM}$. The magnitude of the DM mass in this case is completely unrelated to the Z' mass.

[2]This can be realized if the heavy fermions generating these operators come in complete SU(5) representations.

- If they are chiral, i.e. the left and right $\mathrm{U}(1)_X$ charges are different $X_L^{\mathrm{DM}} = X_R^{\mathrm{DM}} \pm 1$, then we can write down Higgs-type masses $\lambda_{\mathrm{DM}} S \bar{\Psi}_L^{\mathrm{DM}} \Psi_R^{\mathrm{DM}} +$ h.c. or $\lambda_{\mathrm{DM}} S^+ \bar{\Psi}_L^{\mathrm{DM}} \Psi_R^{\mathrm{DM}}$ +h.c. In this case, for DM Yukawa couplings of the order of the Z' gauge couplings, naturally $M_{\mathrm{DM}} \sim M_{Z'}$.

Since we will be interested in electroweak values for DM mass, in both cases we will consider a standard range $100\ GeV \lesssim M_{\mathrm{DM}} \lesssim 1\ TeV$.

The Higgs S can be also invoked to provide a mass for the heavy fermions.

Let us enter more into the details of the effective interactions. For notations convenience we define:

$$\theta_X \equiv \frac{a_X}{V} \quad , \qquad \mathcal{D}_\mu \theta_X \equiv \partial_\mu \theta_X - g_X Z'_\mu ,$$
$$\widetilde{F}_{\mu\nu} \equiv \epsilon_{\mu\nu\rho\sigma} F^{\rho\sigma} \quad , \qquad (FG) \equiv \mathrm{Tr}[F_{\mu\nu} G^{\mu\nu}] \quad , \qquad \mathcal{T}r(EFG) \equiv \mathrm{Tr}[E_\mu{}^\lambda F_{\lambda\nu} G^{\nu\mu}] , \tag{2.3}$$

where Tr takes into account a possible trace over non-abelian indices. Of crucial importance in what follows are the symmetries of the high-energy theory, which includes the heavy fermions $\Psi_{L,R}^{(H)}$ transforming under both $\mathrm{U}(1)_X$ and the SM gauge group. There are two cases:

- If they are vector-like, i.e. the left and right $\mathrm{U}(1)_X$ charges are equal $X_L = X_R$ and therefore they have Dirac masses $M \bar{\Psi}_L^{(H)} \Psi_R^{(H)}$ + h.c., the effective operators obtained after integrating them out have to respect the charge conjugation symmetry C. In this case, a straightforward generalization of the Furry's theorem applies and the first effective operator constructed out of gauge fields mixing the two sectors is of the Euler-Heisenberg type $(1/M^4) F^4$. Due to the big M^4 mass suppression, this is not the case of interest for us.

- If they are chiral, i.e. the left and right $\mathrm{U}(1)_X$ charges are different $X_L = X_R \pm 1$ and therefore they have Higgs-type masses $\lambda S \bar{\Psi}_L^{(H)} \Psi_R^{(H)}$ + h.c. or $\lambda S^+ \bar{\Psi}_L^{(H)} \Psi_R^{(H)}$ + h.c. , the effective operators obtained after integrating them out violate C. They respect the CP symmetry if all couplings are real and can violate CP for complex couplings.

We use gauge invariance and CP symmetry of the lagrangian in order to classify the effective interaction terms invariant under $\mathrm{SU}(2) \times U_Y(1) \times \mathrm{U}(1)_X$ at low-energy. An important point in what follows is that, while the $\mathrm{U}(1)_X$ gauge symmetry is necessarily realized in the broken (Stueckelberg) phase, if they are generated by heavy states respecting the SM gauge symmetry, the effective operators have to be invariant under the unbroken SM gauge group. CP symmetry is a useful tool in classifying the effective operators since non-decoupling (mass-independent) effects have to respect it. We restrict in what follows for simplicity to CP-even operators.[3] Restricting then to CP-invariant operators mixing the two sectors, we then find:

[3]The CP-odd operators have the form:

Dimension-four, CP-odd operators: $\mathcal{D}_\mu \theta_X (H^\dagger D^\mu H + c.c.)$, $\quad \partial^\mu \mathcal{D}_\mu \theta_X H^\dagger H$.

- Dimension-four operators:

$$\delta\; F^{Y}_{\mu\nu} F^{X\,\mu\nu} \qquad , \qquad i\eta\; \mathcal{D}_\mu \theta_X H^\dagger D_\mu H + c.c. \tag{2.4}$$

- Dimension-six operators:

$$\begin{aligned}
\mathcal{L}_{\rm mix} = \; & \frac{1}{M^2}\Big\{ b_1 \mathcal{T}r(F^X F^Y \widetilde{F}^Y) + 2b_2 \mathcal{T}r(F^X F^W \widetilde{F}^W) + b_3 \mathcal{T}r(F^Y F^X \widetilde{F}^X) \\
& + \mathcal{D}^\mu \theta_X \left[i (D^\nu H)^\dagger (c_1 \widetilde{F}^Y_{\mu\nu} + c_2 \widetilde{F}^W_{\mu\nu} + c_3 \widetilde{F}^X_{\mu\nu}) H + c.c. \right] \\
& + \partial^\mu \mathcal{D}_\mu \theta_X \left[d_1 (F^Y \widetilde{F}^Y) + 2 d_2 (F^W \widetilde{F}^W) + d_3 (F^Y \widetilde{F}^X) \right] \\
& + \mathcal{D}_\mu \theta_X \mathcal{D}^\mu \theta_X \left[d_4 (F^Y F^Y) + 2 d_5 (F^W F^W) \right] \Big\} \, .
\end{aligned} \tag{2.5}$$

For our aim we are interested in the gauge invariant terms which couple to the Higgs field, in order to reproduce the coupling to the axion a_X and the SM neutral Golstone boson a_H, which in the SM broken phase is:

$$\epsilon^{\mu\nu\rho\sigma} \mathcal{D}_\mu \theta_X \; \mathcal{D}_\nu \theta_H \; F^Y_{\rho\sigma} \, , \tag{2.6}$$

where $\theta_H = a_H / v$.

In (2.4), the coefficient δ parameterizes the kinetic mixing term of Z' with the hypercharge gauge field. The parameter η generates, after electroweak symmetry breaking, a mass mixing between Z and Z' that can be estimated to be small in such effective theories and we ignore it in our analysis. On the other hand, b_i, c_i and d_i in (2.5) contain the possible low-energy three gauge-boson interaction terms. Their values are fixed by the properties of the more fundamental theory, and in particular many of them can vanish depending on the spectrum of the sector integrated out. Some remarks are in order: the kinetic mixing term is a general feature of all the heavy fermion spectra with coupling to both $\mathrm{U}(1)_X$ and $\mathrm{U}(1)_Y$. However, it follows from the Furry's theorem that the cubic gauge boson interaction terms appear only when chiral fermions are integrated out. This is then consistent with the heavy fermions getting masses from Yukawa couplings to the heavy Higgs S breaking $\mathrm{U}(1)_X$. Effective operators like the ones in the second line in (2.5) contain,after electroweak symmetry breaking, the generalized Chern-Simons (GCS) dimension-four operator (1.3)

$$\frac{v^2}{M^2}\; \epsilon^{\mu\nu\rho\sigma} Z'_\mu B_\nu F^Y_{\rho\sigma} \; . \tag{2.7}$$

Since it originates from a dimension six operator, (2.7) is suppressed by the heavy mass scale M^2. As we will see in section 4, this can also be explicitly checked by using the

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Dimension-six, CP-odd:

$$\frac{1}{M^2} \mathcal{D}^\mu \theta_X \left[i (D^\nu H)^\dagger F^V_{\mu\nu} H + c.c. \right] , \qquad \frac{1}{M^2} \left[(D^\nu H)^\dagger \widetilde{F}^V_{\mu\nu} H + c.c. \right] ,$$
$$\frac{1}{M^2} \partial^\mu \mathcal{D}_\mu \theta_X (F^V F^V) , \qquad \frac{1}{M^2} \mathcal{D}_\mu \theta_X \mathcal{D}^\mu \theta_X (F^V \widetilde{F}^V)$$

.

We remind the reader that the axion a_X is CP odd. Some of these operators could be of some interest for CP violation in the Higgs sector, but this is beyond the goals of the present paper.

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formulae for axionic couplings and GCS terms in [5] in the decoupling limit $M \to \infty$. As a general result, for one Z' the mass-independent three gauge-boson interaction terms indeed vanish after imposing the anomaly cancelations for the heavy spectrum, and the leading non vanishing contributions come from the dimension-six operators listed previously. It is interesting to notice that, with the exception of a possible kinetic mixing in the first line in (2.5), all the operators mixing the (in)visible Z' to the SM are mass-suppressed and therefore decouple at low-energy, in agreement to the decoupling theorem [14]. The only ingredients we need in order to prove this are the SM gauge invariance in the unbroken phase and CP symmetry of the effective operators in the decoupling limit.

Finally, notice that our analysis concerning the effective operators mixing Z' to the SM gauge bosons, remains valid when the heavy sector integrated out is free only from mixed gauge anomalies (and for example with a non-vanishing $\mathrm{U}(1)_X^3$ anomaly).

Due to the mass suppression, effects of the operators discussed in the previous paragraph at low-energy are generically suppressed by E^2/M^2 or v^2/M^2 and can have important effects only for energies not far below the heavy fermion masses. For vector-like heavy fermions the suppression is more severe E^4/M^4 due to the charge conjugation invariance constraints. An obvious question is then if it is possible at all to generate genuine non-decoupling effects by integrating-out heavy fermions which are chiral with respect to extra U(1)'s but cancel all triangle gauge anomalies between themselves. For one massive $\mathrm{U}(1)_X$, as we proved above, this is impossible. For two extra symmetries, $\mathrm{U}(1)_X$ and $\mathrm{U}(1)'_X$ with Stueckelberg realization of gauge symmetries, this is however possible [11]. Indeed, in this case the dimension-four operator

$$\epsilon^{\mu\nu\rho\sigma} \; \mathcal{D}_\mu \theta_X \; \mathcal{D}_\nu \theta'_X \; F^Y_{\rho\sigma} \tag{2.8}$$

is gauge invariant and can be generated by a heavy chiral but anomaly-free fermion spectrum charged under the two U(1)'s and simultaneously under the SM. We explicitly verify this in section 4 by using the formulae of ref. [5]. The term (2.8) provides an interesting counter-example of the decoupling theorem [14], different from the non-decoupling effects studied in [15], since the heavy fermionic spectrum, albeit chiral, is free of any gauge anomalies.

3 (In)visible Z' as a mediator of dark matter annihilation

Several Dark Matter candidates have been proposed and widely discussed in the literature. The distinct phenomenological feature of the present model is a clear dark matter annihilation signature in the galactic halo. We are interested in particular in the trilinear couplings of the form $Z'Z\gamma$ and $Z'ZZ$. These terms can provide a clear signature for the indirect detection of dark matter. The main idea is the following: if the dark matter candidate is lighter than the fermionic sector which we integrated out, the unique tree level annihilation diagram is given by the exchange of Z'. Then, Z' can couple to the visible sector only via the couplings to the SM gauge bosons. We also stress that, as shown in figure 2, this could also give one of the very few available signatures of such (in)visible Z'.

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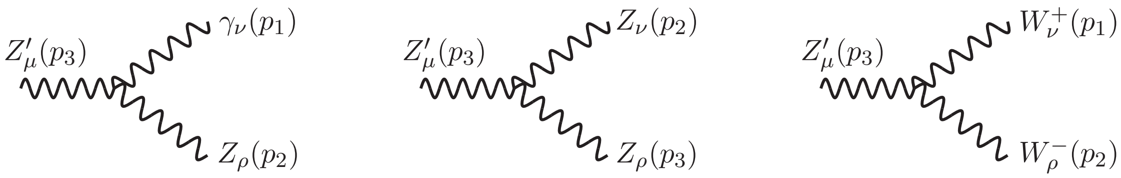

Figure 1. Three vertices of interest generated by (2.5).

The relevant information for the corresponding analysis is contained in the operators $\mathcal{L}_{\rm mix}$ in (2.5). Indeeed, we can easily extract the $Z'VV$ interaction vertices generated by (2.5):

$$\begin{aligned}
\Gamma^{Z'\gamma Z}_{\mu\nu\rho}(p_3;p_1,p_2) &= -8\,\frac{(d_1-d_2)}{M^2} g_X \sin\theta_W \cos\theta_W (p_1+p_2)^\mu \epsilon_{\nu\rho\sigma\tau} p_2^\sigma p_1^\tau \\
&\quad -2\,\frac{e\; g_X}{\cos\theta_W \sin\theta_W}\frac{v^2}{M^2}\left[c_1\cos\theta_W + c_2 \sin\theta_W\right]\epsilon_{\mu\nu\rho\sigma} p_1^\sigma \\
\Gamma^{Z'ZZ}_{\mu\nu\rho}(p_3;p_1,p_2) &= -4\,\frac{(d_1 \sin^2\theta_W + d_2\cos^2\theta_W)}{M^2} g_X (p_1+p_2)^\mu \epsilon_{\nu\rho\sigma\tau} p_2^\sigma p_1^\tau \\
&\quad -\frac{e\; g_X}{\cos\theta_W \sin\theta_W}\frac{v^2}{M^2}\left[c_2\cos\theta_W - c_1 \sin\theta_W\right]\epsilon_{\mu\nu\rho\sigma} (p_2^\sigma - p_1^\sigma) \\
\Gamma^{Z'W^+W^-}_{\mu\nu\rho}(p_3;p_1,p_2) &= -4\,\frac{d_2}{M^2} g_X (p_1+p_2)^\mu \epsilon_{\nu\rho\sigma\tau} p_2^\sigma p_1^\tau \\
&\quad -\frac{e\; g_X}{\cos\theta_W \sin\theta_W}\frac{v^2}{M^2} c_2 \epsilon_{\mu\nu\rho\sigma} (p_2^\sigma - p_1^\sigma)
\end{aligned} \tag{3.1}$$

where θ_W is the Weinberg angle, e is the electric charge and M is the typical scale in the massive fermionic sector ($\sim$ TeV). As we will see in UV completions discussed in section 4, the coefficients (c_i, d_i) are combinations of gauge charges and are suppressed by a loop factor ($\sim 10^{-2}$). Before analyzing in more details such a model, it is interesting to observe the dependence of the vertices on each coupling. For instance, $\Gamma^{Z'VV}_{\mu\nu\rho}$ are independent of the coefficients b_i after symmetrization under the exchange $(p_1, \nu \leftrightarrow p_2, \rho)$ and the dependence of $\Gamma^{Z'Z\gamma}_{\mu\nu\rho}$ on d_i is proportional to the total momentum $(p_1+p_2)^\mu$. These facts have significant consequences on the gamma-ray spectrum. Indeed, it follows from (3.1) that the Z γ final state can naturally compete with the ZZ one in a large part of the parameter space.

We also notice that all these trilinear coupling can be written symbolically as $\frac{(\partial)^3}{M^2}VVV$, with V a generic gauge boson, except for the terms depending explicitly on the Higgs field, which after electroweak symmetry breaking are of the form $\frac{v^2}{M^2}\partial VVV$. The ratio between the two contributions in the process we are interested in, will then be roughly depending on the ratio $\frac{v^2}{M^2_{\rm DM}}$. Therefore, the larger the mass of the dark matter candidate, the smaller the contributions of the terms related to the operators parameterized by the coefficients c_i, with respect to the other ones.

Concerning the kinetic mixing δ, its presence induces a redefinition of the gauge boson mass matrix eigenvectors and eigenvalues, as extensively studied in the literature [1]. The

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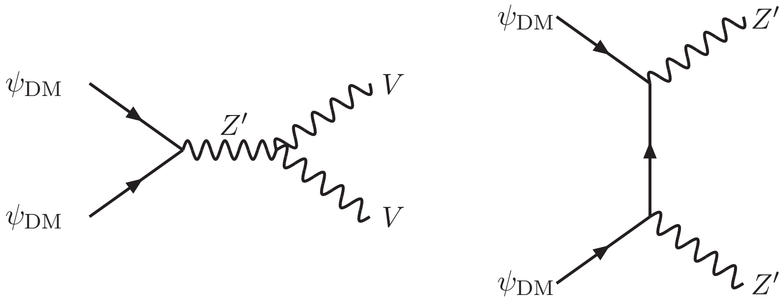

Figure 2. Feynman diagrams contributing to the dark matter annihilation.

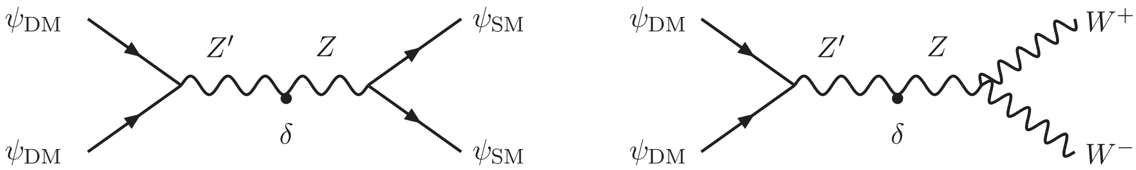

Figure 3. Rôle of the mixing parameter δ.

kinetic mixing also contributes to the DM annihilation, as shown in figure 3. A general comment is that the smaller the kinetic mixing compared to the $Z'Z\gamma$ vertex, the cleaner is the monochromatic gamma ray.

We should also parameterize the $\psi^{\rm DM}\psi^{\rm DM}Z'$ coupling. For a generic DM fermion, the vertex can be written as

$$\Gamma_{\psi^{\rm DM}\psi^{\rm DM}Z'} = i\frac{g_X}{4}\gamma^\mu(V_{\rm DM} - A_{\rm DM}\gamma^5) \tag{3.2}$$

where $V_{\rm DM}$ and $A_{\rm DM}$ are the vectorial and axial couplings, related to the $\mathrm{U}(1)_X$ charges of the Weyl components of $\psi^{\rm DM}$. For a DM candidate with a pure vectorial coupling, the operators with coeff. d_i in the third line of eq. (2.5) do not contribute to the DM annihilation amplitudes, due to the conservation of the DM vector current. On the other hand, the same operators, in the case of a DM with a pure axial coupling, contribute to the amplitudes proportionally to the divergence of the DM axial-vector current and therefore proportionally to the DM mass $M_{\rm DM}$. In our study we consider the general case with both vector and axial-vector couplings, with $(V_{\rm DM}, A_{\rm DM}) = (1, 0)$ for vectorial coupling and $(V_{\rm DM}, A_{\rm DM}) = (0, 1)$ for axial couplings. The specific values of $(V_{\rm DM}, A_{\rm DM})$ are not relevant in our study as they can be absorbed by a redefinition of the $\mathrm{U}(1)'$ coupling g_X. As noticed above, the important point in the formula (3.1) is that, depending on the relative values of the coefficients c_i and d_i, we obtain different final states in the annihilation process $\psi^{\rm DM}\psi^{\rm DM} \to Z' \to VV$. The most interesting for us is the $Z\gamma$ final state. Indeed, we can show easily that for such a final state, because the annihilation occurs for DM particles at

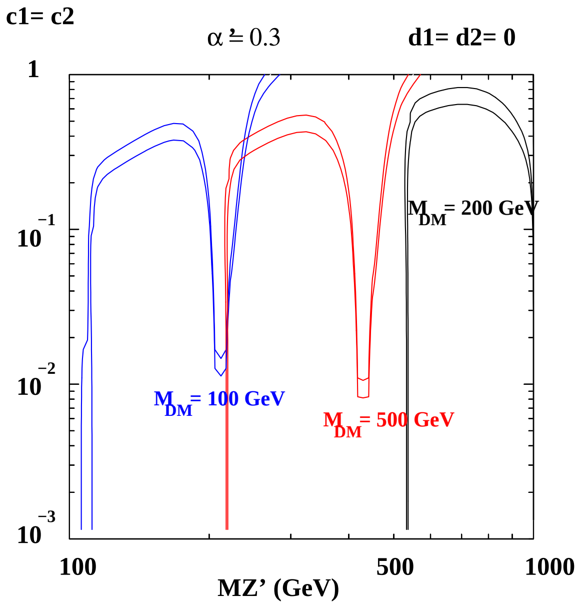

Figure 4. Scan on the mass of Z' (in logarithmic scale) versus the couplings $c_1 = c_2$ for $d_1 = d_2 = 0$ and $M = 1\,\text{TeV}$. We also defined $\alpha' = g_X^2/4\pi$. Colored lines represent the WMAP limits on the dark matter relic density for different values of the dark matter mass. Notice that the results are invariant under the rescaling $M \to \alpha M$, $(c_i, d_i) \to (\alpha^2 c_i, \alpha^2 d_i)$.

rest, the energy of the photon is monochromatic[4] and equal to

$$E_\gamma \;=\; M_{\text{DM}} \left[1 - (\frac{M_Z}{2M_{\text{DM}}})^2\right] \,. \tag{3.3}$$

We computed the relic density Ωh^2 using the last released version of the Micromegas code [19], modified to include the (in)visible Z’ and its couplings to the SM. We show the results in figure 4 as a scan on ($M_{Z'}$, $c_1 = c_2$) for $d_1 = d_2 = 0$. The region between the two lines of the same color corresponds to the 5σ region of WMAP [16]. The different dominant annihilation channels contributing to Ωh^2 are depicted in figures 2. We clearly see the pôle regions when $M_{Z'} \sim 2M_{DM}$. In the Z'-pôle region, the main annihilation channel for $c_2 = d_2 = 0$ is the γZ final state at more than 80 %. This comes mainly from the reduction of the $Z'ZZ$ coupling and a smaller phase space in the final state. For $c_2, d_2 \neq 0$, the annihilation channel $\psi_{\text{DM}}\psi_{\text{DM}} \to Z' \to W^+W^-$ becomes stronger and Ωh^2 decreases, leading to slight changes in the figure 4. When the W^+W^- channel is open ($d_2, c_2 > d_1, c_1$), this final state is the dominant one at 60%. We plot in figure 5 the branching fraction of Z' into the $Z\gamma$ (red full line), ZZ (green dotted line) and W^+W^- (blue dashed line) as a function of c_2/c_1 for $M_{\text{DM}} = 200\,\text{GeV}$ and $M_{Z'} = 215\,\text{GeV}$. W^+W^- final state begins to be dominant for $c_2 \sim 2c_1$. Another interesting feature in regions where the Z' is light ($M_{Z'} < M_{DM}$) is that we observe a new zone, respecting WMAP, almost

[4]One should note that models with DM annihilations producing enhanced gamma ray lines have also been recently suggested in a completely different framework in [18].

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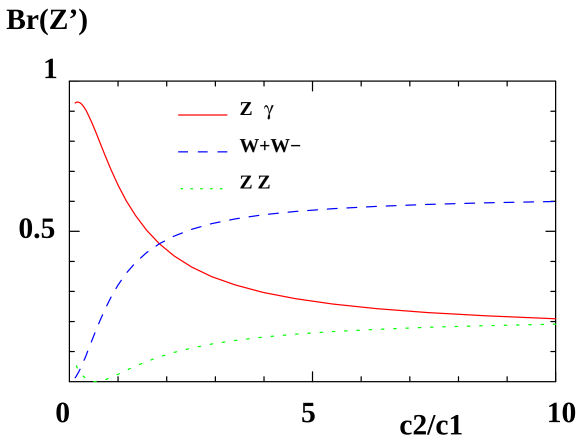

Figure 5. Branching fractions of Z' boson decays into $Z\gamma$ (red full line), ZZ (green dotted line) and W^+W^- (blue dashed line) as a function of c_2/c_1 for $M_{\mathrm{DM}} = 200\,\mathrm{GeV}$, $M_{Z'} = 215\,\mathrm{GeV}$ and $d_1 = d_2 = 0$.

independent of the value of d_1. This region comes from the t-channel annihilation into $Z'Z'$ final state, depicted in figure 2. Depending on the energy of the Z' in the final state, we can observe its decay modes to SM gauge bosons.

Concerning the direct detection prospects, it will be hard to see any spin-dependent or spin-independent signal. Indeed, the idea of direct detection experiment is based upon the measurement of the recoil energy of a target nucleus hit by a dark matter particle. As no effective coupling exists between the Z' and the constituent quarks of the proton, the cross section of such process is simply suppressed.[5] However the indirect detection possibilities seem more promising as the dark matter annihilate only into WW, ZZ or $Z\gamma$ final states. The monochromatic photon present in the $Z\gamma$ final state could be a smoking gun signal of such models. Indeed, monochromatic processes exist also in supersymmetric or KK-like models at one loop order, but they are invisible after including the total amount of diffuse flux coming from the tree level processes. In the case of an (in)visible Z', all the final states come with an amplitude of the same order of magnitude: the γ-monochromatic line can easily be disentangled from the diffuse background. We plotted in figures 6 the diffuse gamma fluxes from the galactic center that we expect in different scenarios for a DM mass of 250 GeV and 700 GeV. For heavy axial dark matter, (figure 6b) more than 90% of the signal is coming from the term proportional to d_i. We used the Pythia Monte Carlo to simulate the gamma-ray spectrum using an analysis similar to that performed in [20].

For a Z' mass close to pole $M_{Z'} \sim 2M_{\mathrm{DM}}$, the main process contributing to the relic abundance is the s-channel exchange of a Z': the monochromatic line (3.3) is clearly visible (figure 6a and b). However, if the Z' is lighter than ψ_{DM}, the t-channel annihilation $\psi_{\mathrm{DM}}\psi_{\mathrm{DM}} \to Z'Z'$ dominate the relic density annihilation processes. The Z' decays finally

[5] This is one of the main reason that the (in)visible Z' will be hardly visible at LHC. Indeed, the relevant Feynman diagrams are the same, whereas the main production channel of the Z' is by vector-vector fusion, as described in [10].

into ZZ or $Z\gamma$, but not exactly at rest: we measure the "would-be" monochromatic γ-ray line, deformed by the kinetic component of the Z' (figure 6c). In some cases, the result can be spectacular and could be seen by the satellite GLAST/FERMI-LAT [21] after 5 years of data taking.[6] Finally, we also checked how the kinetic mixing between Z and Z' could play an important role in such analysis. Even with large mixing (figure 6d), we are still able to observe a (slightly reduced) gamma-ray line of the same order of magnitude as the continuous signal. The reduction of the amplitude comes from the fact that the diagrams depicted in figure 3 also contribute to the annihilation of the DM particle . The chosen value of the mixing δ is the maximal value consistent with the direct detection limit coming from the proton-dark matter elastic scattering cross section ($\sigma^{\rm SI}_{\Psi^{\rm DM}-p} \sim 10^{-8}$pb) [22]. Let us mention that recent works ([3] and references therein) studied Stueckelberg Z' extensions with Z-Z' kinetic mixing, but without the Chern-Simons terms.

In the case of two Z', as discussed in section 2 and further discussed in section 4, there is a genuine non-decoupling operator (2.8). If its coefficient is comparable or dominant over the kinetic mixings ZZ' and ZZ'', the dark-matter annihilation proceeds via the process $\psi^{\rm DM}\psi^{\rm DM} \rightarrow$ virtual $Z' \rightarrow Z''\,\gamma$, which is unsuppressed by the heavy mass and generates a clean gamma ray signal at an energy $E_\gamma = M_{\rm DM}\left[1-(\frac{M_{Z''}}{2M_{\rm DM}})^2\right]$.

4 UV renormalizable theories

Finally, we would like to discuss possible UV completions that give the earlier discussed effective operators. The question if such patterns emerge in renormalizable quantum field theories was addressed in [5]. The framework is a consistent (i.e. anomaly-free) and renormalizable gauge theory with spontaneously-broken gauge symmetry via the Brout-Englert-Higgs mechanism. Through appropriate Yukawa couplings, some large masses can be given to a subset of the fermions . We consider the general case of several spontaneously broken U(1)'s. We denote by $\psi^{(h)}_{L,R}$ such massive chiral fermions. Their $\mathrm{U}(1)_i$ charges are $X^{(h)i}_{L,R}$. In the sequel, we briefly review and adapt the results of the calculation [5] of the effective axion and GCS couplings at low-energy , generated by the loops of the heavy chiral fermions. In what follows, we write explicitly only the massive gauge fields A_i and we consider for simplicity a number of Higgs fields S_i equal to the number of massive U(1)'s. This allows to simplify the formulae, but the results can be easily extended to a more general case. Therefore we can assume that the Higgs field S_i has charge 1 under the gauge transformation related to the vector field A_i (more complicated charge assignments can be reduced to this one after rotations and redefinitions).

The relevant terms in the effective action of the heavy fermion sector of the theory are

$$\begin{aligned} L_h \;=\; & \bar{\psi}^{(h)}_L\left(i\gamma^\mu\partial_\mu + g^i X^{(h)i}_L\gamma^\mu A^i_\mu\right)\psi^{(h)}_L + \bar{\psi}^{(h)}_R\left(i\gamma^\mu\partial_\mu + g^i X^{(h)i}_R\gamma^\mu A^i_\mu\right)\psi^{(h)}_R \\ & -\left(\bar{\psi}^{(h)}_L M^{(h)}\psi^{(h)}_R + \text{h.c.}\right), \end{aligned} \tag{4.1}$$

[6] Work in progress.

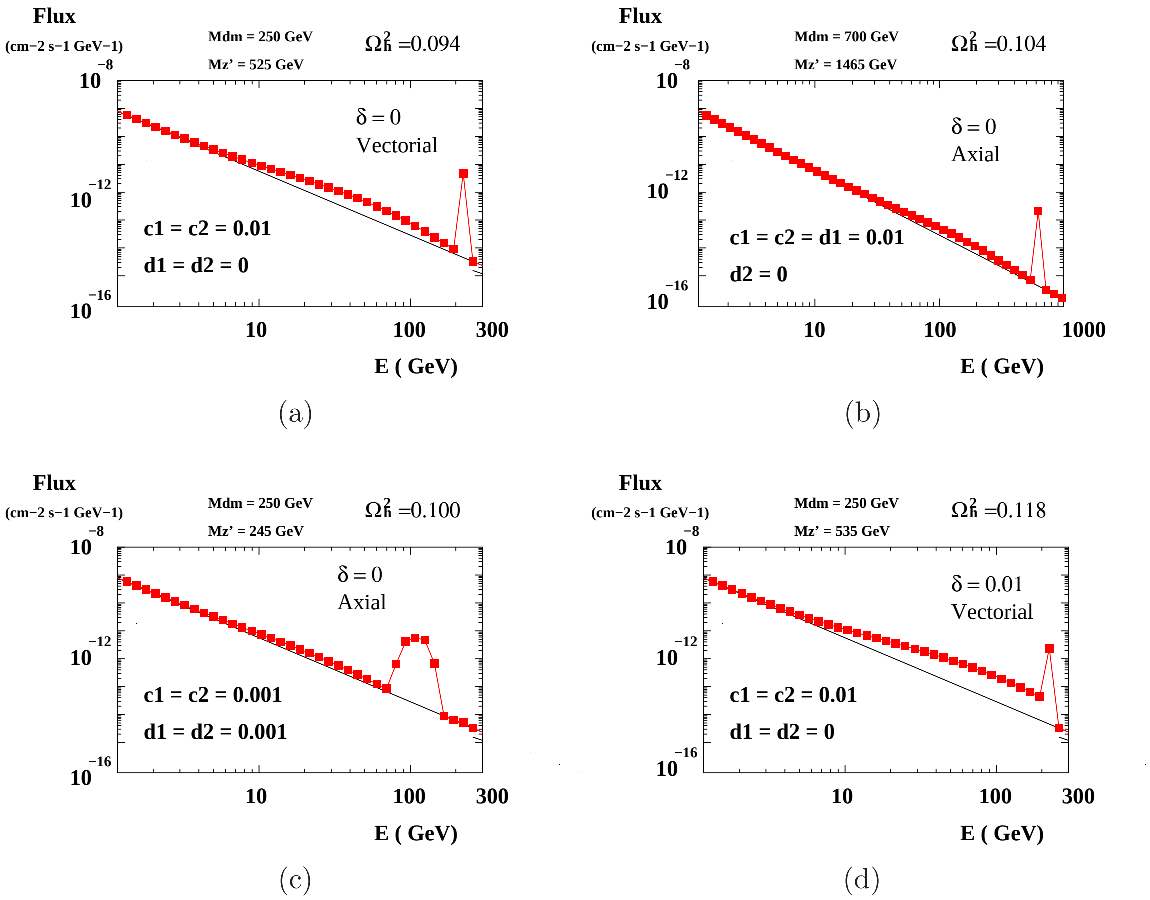

Figure 6. Typical example of a gamma-ray differential spectrum for different masses of dark matter and Z' and $Z - Z'$ mixing angle, compared with the background (black line [20]). All fluxes are calculated for a classical NFW halo profile and $M = 1\,\mathrm{TeV}$. "Axial" and "vectorial" stands for the nature of the $\psi_{\rm DM}\psi_{\rm DM}Z'$ coupling.

where M^h is the mass matrix of heavy fermions, with matrix elements

$$\begin{aligned} M^{(h)}_{ab} &= \lambda^h_{ab}\, S_i \qquad \text{case (a)} \quad \text{or} \\ M^{(h)}_{ab} &= \lambda^h_{ab}\, \bar{S}_i \qquad \text{case (b)}\,, \end{aligned} \tag{4.2}$$

where S_i is the Higgs field of charge $+1$ under the gauge group $\mathrm{U}(1)_i$ and singlet with respect to the other gauge groups. The Higgses spontaneously break the abelian gauge symmetries via their vevs, $\langle S_i \rangle = V_i$. For every gauge group $\mathrm{U}(1)_i$, there is a set h_i of fermions We charges of the fermions satisfy the relations

$$X_L^{(h_i)i} - X_R^{(h_i)i} = \pm 1 \equiv \epsilon^{(h_i)i} \tag{4.3}$$

in order to couple with the Higgs field S_i. The fermions we are considering are charged both under the SM gauge group and the additional U(1)'s. Whereas they are *chiral* wrt the Z'-type symmetries, they are *vector-like* wrt the SM gauge group, as required by the existence of the mass terms in (4.2). If the associated Yukawa coupling eigenvalues are large, $\lambda_{ab}^h \gg g_i$, spontaneous symmetry breaking generates large Dirac fermion masses. We consider the heavy fermion decoupling limit, with fixed Higgs vev's and fixed gauge boson masses, whereas $M^{(h)} \to \infty$. Since we are interested in the (in)visible Z' where the SM fermions are neutral under the massive U(1)'s, the heavy fermion sector is by itself anomaly-free

$$\sum_h (X_L^i X_L^j X_L^k - X_R^i X_R^j X_R^k)^{(h)} \;=\; 0\;. \tag{4.4}$$

We are interested in low-energy couplings generated by the loops of the heavy fermions. After gauge symmetry breaking, we parameterize the scalar fields by

$$S_i \;=\; (V_i + s_i)\; e^{\frac{i\;a_i}{V_i}}\;, \tag{4.5}$$

where s_i are massive Higgs-like fields and a_i are axions. The gauge transformations of gauge fields and axions are

$$\delta A_\mu^i \;=\; \partial_\mu \alpha^i \quad , \quad \delta a_i \;=\; V_i \;\; \alpha^i\;. \tag{4.6}$$

The GCS terms and axionic couplings can be computed by performing a diagrammatic computation with the action (4.1), by starting from the corresponding three gauge boson amplitude induced by triangle diagram loops of heavy fermions and expanding in powers of external momenta $k/M^{(h)}$. We define the effective action after integrating out the heavy states by

$$\begin{aligned}\mathcal{S} = &-\sum_i \int \frac{1}{4} F_{i,\mu\nu} F_i^{\mu\nu} + \frac{1}{2}\int \sum_i (\partial_\mu a^i - g_i V_i A_\mu^i)^2\,,\\ &+\frac{1}{96\pi^2} C_{ij}^i\;\epsilon^{\mu\nu\rho\sigma}\int a^i F_{\mu\nu}^i F_{\rho\sigma}^j + \frac{1}{48\pi^2} E_{ij,k}\;\epsilon^{\mu\nu\rho\sigma}\int A_\mu^i A_\nu^j F_{\rho\sigma}^k\,,\end{aligned} \tag{4.7}$$

where the coefficients $E_{ij,k}$ satisfy the cyclic relation

$$E_{ij,k} \;+\; E_{jk,i} \;+\; E_{ki,j} \;=\; 0 \tag{4.8}$$

and the gauge invariance conditions, in the presence of an anomaly free spectrum, read

$$\begin{aligned} C_{jk}^i g_i V_i \;-\; E_{ij,k} \;-\; E_{ik,j} &= 0\,,\\ C_{jk}^i g_i V_i \;+\; C_{ki}^j g_j V_j \;+\; C_{ij}^k g_k V_k &= 0\;.\end{aligned} \tag{4.9}$$

One can easily find the solution of (4.9)

$$E_{ij,k} \;=\; \frac{1}{3}\left(g_i V_i C_{jk}^i \;-\; g_j V_j C_{ik}^j\right)\;. \tag{4.10}$$

The result obtained in [5], in the decoupling limit $M^{(h)} \to \infty$ with finite Higgs vev's V_i is

$$\begin{aligned} E_{ij,k} &= \frac{1}{4}\sum_h (X_L^i X_R^j - X_R^i X_L^j)^{(h)} (X_R^k + X_L^k)^{(h)}\,,\\ C_{ij}^I &= \frac{1}{4 g_I V_I}\;\sum_{h_I} \epsilon^{(h_I)I} [2(X_L^i X_L^j + X_R^i X_R^j) + X_L^i X_R^j + X_R^i X_L^j]^{(h_I)}\,,\end{aligned} \tag{4.11}$$

where the index h_I in (4.11) refers to the heavy fermionic spectrum coupling to the axion a_I. Notice that, while within the more general framework of [5] the GCS terms $E_{ij,k}$ had a freedom in their definition related to the different possible distribution of the gauge anomalies in the low-energy theory among different U(1) currents, in the context of the present paper, they are uniquely fixed since the low-energy spectrum is neutral under the massive abelian gauge fields.

Actually, due to gauge invariance of the low-energy effective action, the unique gauge invariance combination of GCS and axionic terms is [5, 7, 11]

$$\frac{1}{48\pi^2} d_{ij,k}\, \epsilon^{\mu\nu\rho\sigma}\, (\partial a^i - g^i V^i A^i)_\mu\, (\partial a^j - g^j V^j A^j)_\nu\, F^k_{\rho\sigma}\,, \tag{4.12}$$

which justifies the form of the operators used in section 2, in terms of the Stueckelberg gauge-invariant combinations. They lead to the couplings

$$E_{ij,k} \;=\; d_{ij,k} g^i g^j V^i V^j \quad , \quad C^i_{jk} \;=\; d_{ij,k}\, g^j V^j + d_{ik,j} g_k V_k \;. \tag{4.13}$$

We now analyze in a model-independent way the resulting low-energy couplings in the cases of one and two extra U(1)'s, using (4.11) and the anomaly cancelation constraints, and compare the results with the operatorial analysis performed in section 2. We don't need to assume that the whole mass matrix comes from the breaking of the additional U(1)'s. For appropriate quantum numbers, some SM like mass entries $\lambda'_{ab} H \bar{\psi}^{(h,a)}_L \psi^{(h,b)}_R$, where H is the SM Higgs, can also exist without changing our conclusions, as long as we keep $\lambda'_{ab}\langle H\rangle$ fixed in the decoupling limit $M^{(h)} \to \infty$. If the DM is chiral under Z' and a singlet under the SM , we need of course to assume that there exists a heavier partner singlet under the SM which cancel its $U(1)^3_{Z'}$ anomaly.

4.1 One Z'

We are interested in the XYY GCS term, where $X \equiv Z'$ is the massive gauge boson and Y is the hypercharge one. The relevant information in the high-energy spectrum is encoded in the Y and X charges

$$\begin{array}{ccc} & Y & X \\ \Psi^a_L & y_a & x_a \\ \Psi^a_R & y_a & x_a - \epsilon_a \end{array} \tag{4.14}$$

We denote by $l_a = dim R_a$ the dimension of each fermion representation. The mixed anomaly $Tr(Y^2X)$, the GCS coeff. $E_{XY,Y}$ and the axionic[7] coupling C_{YY} are computed to be

$$Tr\;(Y^2X) = \sum_a l_a \epsilon_a y_a^2 \quad , \quad E_{XY,Y} \;=\; \frac{1}{2}\sum_a l_a \epsilon_a y_a^2\,,$$
$$C_{YY} = \frac{3}{2gV}\; \sum_a l_a \epsilon_a y_a^2 \;. \tag{4.15}$$

[7] In this case, there is only one axion.

Since they are all proportional to each other and we consider anomaly-free spectra $Tr\ (Y^2 X) = 0$, the GCS and axionic couplings vanish in the decoupling limit, in agreement with the effective operator analysis performed in section 2. On the other hand, the dimension-six operators in (2.5) certainly do exist and are the ones we took into account in the phenomenological analysis performed in section 3.

4.2 Two Z'

In contrast with the case of one Z', the effective operator analysis in section 2 revealed the possible existence of the non-decoupling dimension-four operator (2.8). In order to check its existence in the UV theory with heavy chiral fermions, we consider the following table of charge assignements

$$\begin{array}{cccc} & Y & X_1 & X_2 \\ \Psi_L^a & y_a & x_a & z_a \\ \Psi_R^a & y_a & x_a - \epsilon_a & z_a \\ \chi_L^m & y_m & x_m & z_m \\ \chi_R^m & y_m & x_m & z_m - \epsilon_m \end{array} \tag{4.16}$$

As transparent from the table, the first group of fermions Ψ acquire masses from the first Higgs field breaking X_1, whereas the second group of fermions χ acquire masses from the second Higgs field breaking X_2. Analogously to the previous example, we define $l_m = dim R_m$. In this case we are interested in $E_{X_1 X_2, Y}$ and the two axionic couplings $C^1_{X_2 Y}$ and $C^2_{X_1 Y}$. Using again (4.11), we find

$$\begin{aligned} Tr\ (X_1 X_2 Y) &= \sum_a l_a \epsilon_a y_a z_a + \sum_m l_m \epsilon_m x_m y_m\,, \\ E_{X_1 X_2, Y} &= \frac{1}{2}\left(\sum_a l_a \epsilon_a y_a z_a - \sum_m l_m \epsilon_m x_m y_m\right)\,, \\ C^{X_1}_{X_2 Y} &= \frac{3}{2 g_1 V_1} \sum_a l_a \epsilon_a y_a z_a \quad , \quad C^{X_2}_{X_1 Y} = \frac{3}{2 g_2 V_2} \sum_m l_m \epsilon_a x_m y_m\,. \end{aligned} \tag{4.17}$$

In this case, by imposing cancelation of the mixed anomaly $Tr\ (X_1 X_2 Y) = 0$, we find that the GCS and the two axionic couplings precisely fit into the gauge invariant term

$$d_{X_1 X_2, Y}\ \epsilon^{\mu\nu\rho\sigma}\ (\partial a_1 - g_1 V_1 X_1)_\mu (\partial a_2 - g_2 V_2 X_2)_\nu\ F^Y_{\rho\sigma}\,, \tag{4.18}$$

with $d_{X_1 X_2, Y} = \frac{1}{g_1 g_2 V_1 V_2} E_{X_1 X_2, Y}$, which is the non-decoupling operator (2.8) we were searching for. It is also easy to compute $E_{Y X_1, X_2}$ and $E_{X_2 Y, X_1}$ and check (4.8) and (4.9). There are other gauge invariant operators purely within the Z' sector that are however irrelevant for our purposes.

5 Conclusions

In this paper we studied the consequences of an extension of the standard model containing an invisible extra gauge group under which the SM particles are neutral. We showed that effective operators mixing the two sectors are generated by loops of heavy fermions, which are chiral wrt U(1) and are vector-like wrt to the SM gauge group. This implies in particular that the decoupling limit is taken by considering large SM invariant masses $M \gg v$ of the heavy fermions, whereas keeping fix SM like masses $m \sim v$. The induced operators mix SM with Z' via a $Z'VV$ vertices. If the lightest fermion in the Z' sector is stable, the induced operators allow for its annihilations that can give rise to a viable dark matter candidate. Its annihilations produce clean visible signals through a gamma-ray line which seems to be quite an universal clear feature of this type of constructions. The fact that only one gamma-ray line is produced and the fact that no signal is expected from direct detection experiments can be a distinctive signature of the model. Indeed a supersymmetric neutralino or inert Higgs scalar for instance would instead annihilate into $Z\gamma$ *and* $\gamma\gamma$ final states with similar ratios and interact non-trivially with the nuclei. Such smoking gun signatures could be observable by near future experiments like FERMI/GLAST or, after more data taking, by the HESS/MAGIC telescopes.

From a theoretical side, we showed that heavy, chiral wrt U(1) but vector-like wrt the SM gauge group, fermions can generate the effective operators (2.5), even if they cancel among themselves all gauge anomalies. Since we discussed only the case where the heavy fermion masses which decouple are SM invariant, in the decoupling limit the effective operators do respect the SM gauge symmetry. This is to be contrasted with the case of decoupling a SM-like mass [15], where the low-energy, non-decoupling effective operators arise only in the broken phase, after electroweak symmetry breaking. For the same reason, the effective operators realize the Z' gauge symmetry in the broken phase.

As a result, in the case of one Z', in addition to the well-studied $Z' - Z$ kinetic mixing and $Z' - Z$ mass mixing described in eq. (2.4), the other operators mixing SM with Z' are the dimension six operators, eqs. (2.5). Although they contain a heavy-mass suppression, they can induce the viable dark-matter annihilation channel discussed in detail in section 3. On the other hand, for two Z', as discussed briefly in section 2 and further in section 4, there is a genuine non-decoupling producing an unsuppressed process $\Psi^{\rm DM}\Psi^{\rm DM} \rightarrow$ virtual $Z' \rightarrow Z''\,\gamma$, if both additional gauge bosons are light, in particular if $M_{Z''} < 2M_{\rm DM}$. The class of (in)visible Z' discussed in our paper is another example of an "abelian hidden sector" [23], with potential collider signatures awaiting a dedicated analysis. It would be interesting to perform a systematic study of the effects of the effective operators (2.5) at low-energy from a decoupling perspective. It is interesting to note that a recent study [24] in the context of the Green-Schwarz mechanism gave similar phenomenological features.

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