Physics of the Dark Universe 18 (2017) 127–135

Contents lists available at ScienceDirect

Physics of the Dark Universe

journal homepage: www.elsevier.com/locate/dark

Perturbative unitarity constraints on gauge portals

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ARTICLE INFO

Article history:
Received 29 January 2017
Received in revised form 3 September 2017
Accepted 13 September 2017

Keywords:
Dark matter models

ABSTRACT

Dark matter that was once in thermal equilibrium with the Standard Model is generally prohibited from obtaining all of its mass from the electroweak phase transition. This implies a new scale of physics and mediator particles to facilitate dark matter annihilation. In this work, we focus on dark matter that annihilates through a generic gauge boson portal. We show how partial wave unitarity places upper bounds on the dark gauge boson, dark Higgs and dark matter masses. Outside of well-defined fine-tuned regions, we find an upper bound of 9 TeV for the dark matter mass when the dark Higgs and dark gauge bosons both facilitate the dark matter annihilations. In this scenario, the upper bound on the dark Higgs and dark gauge boson masses are 10 TeV and 16 TeV, respectively. When only the dark gauge boson facilitates dark matter annihilations, we find an upper bound of 3 TeV and 6 TeV for the dark matter and dark gauge boson, respectively. Overall, using the gauge portal as a template, we describe a method to not only place upper bounds on the dark matter mass but also on the new particles with Standard Model quantum numbers. We briefly discuss the reach of future accelerator, direct and indirect detection experiments for this class of models.

1. Introduction

Understanding the nature of dark matter (DM) is one of the most pressing unresolved problems in particle physics. Dark matter is needed to understand structure formation, the observed galactic rotation curves [1–3] and the acoustic peaks in the cosmic microwave background [4]. Moreover, the dark matter relic abundance is measured to be [4]

$$h^2\,\Omega_c = 0.1199 \pm 0.0027. \quad (1)$$

A compelling argument for the origin of this abundance is to assume dark matter was once in thermal contact with the baryon–photon plasma during the early universe. Since all known forms of matter in the universe were once in thermal equilibrium, this type of dark matter is theoretically persuasive. In this scenario, the measured relic abundance is controlled by dark matter annihilations into Standard Model (SM) particles. Because of constraints from the observed large scale structure in the universe, dark matter must be stable and non-relativistic (cold) when departing thermal equilibrium [2].

The Standard Model (SM) alone cannot account for the missing matter in the universe [5]. Current experimental constraints, however, provide some guidance on the structure of the underlying theory. For example, the lack of large missing energy signatures at the Large Hadron Collider (LHC) [6–16] and other colliders [17–24] suggest that dark matter is either heavy or has very small couplings with the SM so that it is not produced in high-energy collisions. Additionally, direct detection experiments [25–27], updated precision electroweak constraints, and precision Z-pole experiments [28–30] all severely constrain the direct coupling of dark matter to the SM Higgs and/or Z bosons. These constraints all imply dark matter cannot obtain all of its mass from the SM Higgs alone [28]. Thus, if dark matter is a weakly interacting massive particle (WIMP), we are led to scenarios where new mediators facilitate dark matter interactions with the SM. Moreover, a new fundamental scale of physics is needed that is (at least partly) responsible for the dark matter mass. Mediator-facilitated interactions help to evade current experimental constraints by partially decoupling the dark matter from the SM. Should these scenarios be realized in nature, the discovery of the mediator particles would be an important step in understanding the nature of dark matter. It is therefore crucial to place bounds on the masses and couplings of these mediators. The most popular ways for dark matter to annihilate via a mediator particle are through the Higgs [31] boson, through scalars that are coloured or charged, or via a new neutral gauge boson. Some of us considered the perturbative unitarity constraints on the Higgs portal in [32–34]. In this work, we focus

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https://doi.org/10.1016/j.dark.2017.09.006

on placing constraints on a scenario where fermionic dark matter is charged under a new, dark gauge group, $U(1)_D$. This gauge group is spontaneously broken by a dark Higgs, Φ, generating a massive, dark Z' boson. This boson is also known in the literature as a dark photon. The dark Z' mixes kinetically as well as through mixed mass terms with the SM Z boson. Thus, the mixing between the hidden sector and the SM allows dark matter (DM) to annihilate via the Higgses, Z and Z' bosons.

We apply unitarity constraints in a manner reminiscent of Griest and Kamiokowski [35]. However, there are important differences: Here we focus on perturbative unitarity constraints which determine, in particular, when the dark matter couplings become strong. WIMP dark matter and perturbativity have always had an important conceptual association. Dark matter masses that violate the perturbative unitarity bounds imply the dark matter is efficiently forming bound states as well as annihilating as the temperature decreases toward the thermal decoupling temperature. Because the dark matter annihilates into lighter states, the annihilation diagrams can be altered (and sometimes dressed with these lighter states) to produce diagrams in which the bound states decay. The dark matter decays have a lifetime well shorter than the age of the universe. Elementary dark matter with a mass beyond the perturbative unitarity bounds may not be an asymptotic state. The resulting bound state is not a viable dark matter candidate. Thus within the gauge portal, the bounds presented in this paper are a first step in understanding the nature of thermally produced dark matter when a coupling necessary for dark matter annihilation becomes strong. Note, it is well known that viable dark matter candidates exist that are the result of strongly coupled or confining hidden sectors. However, in these models the dark matter annihilation processes are still perturbative [36]. We show our perturbative unitarity constraints are improved in comparison to the updated Griest and Kamiokowski bounds [37]. Of central importance is the fact that our methodology places constraints on any particle associated with the dark matter annihilation. For this paper, our bounds on the masses and couplings of the new Higgs and dark gauge boson are novel.

Our basic perturbative unitarity arguments are straightforward. The DM annihilation cross section depends on the masses of the dark matter, the dark Higgs, and the Z', as well as the dark matter couplings to the dark Higgs and dark Z'. As the dark matter gets heavier, its annihilation cross section decreases. In order for heavy dark matter to satisfy the relic abundance constraints, it must annihilate more efficiently and therefore have sizeable couplings to the SM and hidden sectors. Eventually, the couplings required to obtain the correct relic abundance are so large that perturbative unitarity is violated. Perturbative unitarity arguments therefore set an upper bound on the dark matter mass. If the dark Higgs and gauge boson masses are raised to be larger than the dark matter mass, fewer (and more suppressed) annihilation channels are available. The annihilation cross section in these regimes of parameter space is thus diminished. Therefore these arguments yield bounds on the dark matter mass as well as on the mass of any other particle involved with the dark matter annihilation.

In the next section, we introduce a generic $U(1)_D$ model on which to place our unitary bounds and introduce the parameters that need to be constrained as well as the constraints from electroweak precision tests (EWPT). In Section 3, we show how to apply unitarity constraints on the various sectors of the model. Section 4 details how relic abundance and direct detection constraints on the DM sector impact the masses and couplings of the theory. Section 5 gives our results by detailing the bounds on the particle masses obtained by applying the EWPT, unitarity and relic abundance constraints. Conclusion and Appendices follow.

2. A representative model

We extend the SM with an additional $U(1)_D$ gauge group with coupling g_D that is spontaneously broken at a high scale. This dark group is associated to a dark gauge boson that mixes kinetically and via mass terms with SM hypercharge. The $U(1)_D$ gauge group is broken by a new, dark Higgs, that gets a vev u. The model then includes two Higgs fields

$$H = \frac{1}{\sqrt{2}}\begin{pmatrix} \sqrt{2}\,G^{\pm} \\ v + h + i\,G^0 \end{pmatrix} \qquad \Phi = \frac{1}{\sqrt{2}}\left(u + \rho + i\,G^0_\rho\right) \tag{2}$$

where H is the SM Higgs. We also introduce a DM candidate χ, which is a chiral fermion, neutral under the SM gauge groups but charged under $U(1)_D$. All SM particles are taken to be neutral under $U(1)_D$. The dark charge assignments for the DM and the dark Higgs are

$$Q_\Phi = -2 \qquad Q_{\chi_L} = -1 \qquad Q_{\chi_R} = 1. \tag{3}$$

Anomaly cancellation mandates the introduction of the second chiral fermion. In this work, we take this additional fermion to be much heavier than the other particles so that it does not have any influence on the final results.

We adopt the notations and conventions from [38]. The relevant parts of the lagrangian associated to new physics is

$$\mathcal{L} = \mathcal{L}_{\text{gauge}} + \mathcal{L}_{\text{DM}} + \mathcal{L}_{\text{Higgs}} \tag{4}$$

where the dark matter and gauge sectors are

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4}\hat{B}_{\mu\nu}\hat{B}^{\mu\nu} - \frac{1}{4}\hat{Z'}_{\mu\nu}\hat{Z'}^{\mu\nu} - \frac{\sin\delta}{2}\hat{B}_{\mu\nu}\hat{Z'}^{\mu\nu} \tag{5}$$

$$\mathcal{L}_{\text{DM}} = \bar{\chi}_L \not{D}_\mu \chi_L + \bar{\chi}_R \not{D}_\mu \chi_R - \lambda_\chi \bar{\chi}_L \Phi \chi_R + h.c. \tag{6}$$

and

$$\mathcal{L}_{\text{Higgs}} = \left|D_\mu H\right|^2 + \left|D_\mu \Phi\right|^2 - V(H, \Phi) \tag{7}$$

$$V(H,\Phi) = \lambda_1\left(H^\dagger H - \frac{v^2}{2}\right)^2 + \lambda_2\left(\Phi^\dagger\Phi - \frac{u^2}{2}\right)^2 + \lambda_3\left(H^\dagger H - \frac{v^2}{2}\right)\left(\Phi^\dagger\Phi - \frac{u^2}{2}\right) \tag{8}$$

is the Higgs sector. The kinetic mixing is parameterized by the mixing angle δ. The kinetic terms can be diagonalized by defining new fields B_μ and Z'_μ such that [38]

$$\begin{pmatrix} \hat{B}_\mu \\ \hat{Z'}_\mu \end{pmatrix} = \begin{pmatrix} 1 & -\tan\delta \\ 0 & \sec\delta \end{pmatrix}\begin{pmatrix} B_\mu \\ Z'_\mu \end{pmatrix} \tag{9}$$

where the hatted fields are the fields before diagonalizing kinetic mixing. Denoting g_1 and g_2 as the SM hypercharge and weak couplings respectively, the covariant derivatives for the Higgs and DM fields then become

$$DH = \partial H - ig_2 W^a \sigma^a H - \frac{ig_1}{2}BH + \frac{ig_1}{2}\tan\delta\, Z'H \tag{10}$$

$$D\Phi = \partial\Phi - \frac{2ig_D}{\cos\delta}Z'\Phi \tag{11}$$

$$D\chi_R = \partial\chi_R - \frac{ig_D}{\cos\delta}Z'\chi_R \tag{12}$$

$$D\chi_L = \partial\chi_L + \frac{ig_D}{\cos\delta}Z'\chi_L. \tag{13}$$

Any SM particle with non-zero hypercharge will then acquire a dark charge. We now have the effective dark gauge coupling

$$g' = \frac{g_D}{\cos\delta}. \tag{14}$$