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Anomalous $U(1)$ Gauge Bosons as Light Dark Matter in String Theory

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Present experiments are sensitive to very weakly coupled extra gauge symmetries which motivates further investigation of their appearance in string theory compactifications and subsequent properties. We consider extensions of the standard model based on open strings ending on D-branes, with gauge bosons due to strings attached to stacks of D-branes and chiral matter due to strings stretching between intersecting D-branes. Assuming that the fundamental string mass scale saturates the current LHC limit and that the theory is weakly coupled, we show that (anomalous) $U(1)$ gauge bosons which propagate into the bulk are compelling light dark matter candidates. We comment on the possible relevance of the $U(1)$ gauge bosons, which are universal in intersecting D-brane models, to the observed 3σ excess in XENON1T.

INTRODUCTION

The primary objective of the High Energy Physics (HEP) program is to find and understand what physics may lie beyond the Standard $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ Model (SM), as well as its connections to gravity and to the hidden sector of particle dark matter (DM). This objective is pursued in several distinct ways. In this Letter, we explore one possible pathway to join the vertices of the HEP triangle using string compactifications with large extra dimensions [1], where sets of D-branes lead to chiral gauge sectors close to the SM [2, 3].

D-branes provide a nice and simple realization of non-abelian gauge symmetry in string theory. A stack of N identical parallel D-branes eventually generates a $U(N)$ theory with the associated $U(N)$ gauge group where the corresponding gauge bosons emerge as excitations of open strings ending on the D-branes. Chiral matter is either due to strings stretching between intersecting D-branes, or to appropriate projections on strings in the same stack. Gravitational interactions are described by closed strings that can propagate in all dimensions; these comprise parallel dimensions extended along the D-branes and transverse ones.

String compactifications could leave characteristic footprints at particle colliders:

- the emergence of Regge recurrences at parton collision energies $\sqrt{s} \sim$ string mass scale $\equiv M_s = 1/\sqrt{\alpha'}$ [4–6];
- the presence of one or more additional $U(1)$ gauge symmetries, beyond the $U(1)_Y$ of the SM [7–9].

Herein we argue that the (anomalous) $U(1)$ gauge bosons that do not partake in the hypercharge combination could become compelling dark matter candidates. Indeed, as noted elsewhere [10] these gauge fields could live in the bulk and the four-dimensional $U(1)$ gauge coupling would become infinitesimally small in low string scale models, $g \sim M_s/M_{\text{Pl}}$, where M_{Pl} is the Planck mass (for previous investigations in different regions of parameters and different string scenarios,

see for example [11–13]). Note that for typical energies E of the order of the electron mass, the value of g is still bigger than the gravitational coupling $\sim E/M_{\text{Pl}}$, and the strength of the new force would be about 10^7 times stronger than gravity, where we have taken $M_s \sim 8$ TeV, saturating the LHC bound [14].

To develop some sense for the orders of magnitude involved, we now make contact with the experiment. The XENON1T Collaboration has recently reported a surplus of events in $1 \lesssim$ electronic recoils/keV $\lesssim 7$, peaked around 2.8 keV [15]. The total number of events recorded within this energy window is 285, whereas the expected background is 232 ± 15 . Taken at face value this corresponds to a significance of roughly 3σ , but unknown backgrounds from tritium decay cannot be reliably ruled out [15]. Although the excess is not statistically significant, it is tempting to imagine that it corresponds to a real signal of new physics. A plethora of models have already been proposed to explain the excess, in which the DM particle could be either the main component of the abundance in the solar neighborhood, $n_{\text{DM}} \sim 10^5 (m_{\text{DM}}/2.8 \text{ keV})^{-1} \text{ cm}^{-3}$, or else a sub-component of the DM population. Absorption of a ~ 2.8 keV mass dark vector boson that saturates the local DM mass density provides a good fit to the excess for a $U(1)_X$ gauge coupling to electrons of $g_{X,\text{eff}} \sim 2 \times 10^{-16} - 8 \times 10^{-16}$ [15–20]. For such small masses and couplings, the cosmological production should be non-thermal [17], avoiding constraints from structure formation [21, 22]. Leaving aside attempts to fit the XENON1T excess, we might consider a wider range of dark photon masses and couplings. For light and very weakly coupled dark photons, the cooling of red giants and horizontal branch stars give stronger or similar bounds on $g_{X,\text{eff}}$ than direct detec-

tion experiments [23, 24].¹ For instance, rescaling the bounds quoted in [19] leads to an upper bound $g_{X,\text{eff}} \gtrsim 10^{-16} - 10^{-14}$ for m_X varying from 10 to 100 keV. As an example, if we take $m_X \sim 15$ keV in agreement with the bound of ~ 5 keV [21, 22], the upper bound is about $g_{X,\text{eff}} \lesssim 5 \times 10^{-16}$. Obtaining such small values of masses and couplings for the dark photon are challenging as we will show.

GENERATING THE SMALL $U(1)_X$ COUPLINGS

Open string models

We start from ten-dimensional type I string theory compactified on a six-dimensional space of volume $V_6 M_s^6$. The relation between the Planck mass, the string scale, the string coupling g_s , and the total volume of the bulk $V_6 M_s^6$ reads:

$$M_{\text{Pl}}^2 = \frac{8}{g_s^2} M_s^8 \frac{V_6}{(2\pi)^6}. \quad (1)$$

A hierarchy between the Planck and string scales can be due to either a large volume $V_6 M_s^6 \gg 1$ or a very small string coupling. We discuss these two possibilities successively.

From now on, we denote by d the total number of dimensions that are large. For simplicity, we assume that they have a common radius R while the other $6 - d$ dimensions have a radius M_s^{-1} . Obviously, the latter have no more a classical geometry and supergravity description as space dimensions; they represent new degrees of freedom that have a stringy description, for example through the corresponding world-sheet conformal field theories. The couplings of all light states are under control and our formulae still hold. The $U(1)_X$ gauge fields live on a $D(3 + \delta_X)$ -brane that wraps a δ_X -cycle of volume V_X , while its remaining four dimensions extend into the uncompactified space-time. The corresponding gauge coupling is given by:

$$g_X^2 = \frac{(2\pi)^{\delta_X+1} g_s}{V_X M_s^{\delta_X}}. \quad (2)$$

Assuming all the δ_X -cycles are sub-spaces of internal d large dimensions with the same radius, the substitution of (1) into (2) leads to:

$$g_X^2 = 2\pi g_s \left(\frac{8}{g_s^2}\right)^{\delta_X/d} \left(\frac{M_s}{M_{\text{Pl}}}\right)^{2\delta_X/d}. \quad (3)$$

It is straightforward to see that to realize the weakest gauge interaction the volume seen by $U(1)_X$ must exhaust the total

large internal volume suppressing the strength of gravitational interactions $\delta_X = d$ (as in [16]), yielding

$$g_X = \sqrt{\frac{16\pi}{g_s} \frac{M_s}{M_{\text{Pl}}}} \sim 4 \times 10^{-14} \left(\frac{0.2}{g_s}\right)^{1/2} \left(\frac{M_s}{10 \text{ TeV}}\right), \quad (4)$$

where we have taken as reference values $g_s = 0.2$ and $M_s \gtrsim 10$ TeV. The latter is a conservative bound from non-observation of stringy excitations at colliders [14] while a slightly stronger bound of order, but model dependent, can be obtained from limits on dimension-six four-fermion operators [26–29]. As for g_s we will consider that it is in the range $0.01 - 0.2$, and could be fixed after a careful study of running of the gauge couplings. In the case of toroidal compactifications, the internal six-dimensional volume is expressed in terms of the parallel and transversal radii as

$$V_6 = (2\pi)^6 \prod_{i=1}^{d_{\parallel}} R_i^{\parallel} \prod_{j=1}^{d_{\perp}} R_j^{\perp}, \quad (5)$$

where now for each stack of D_p -branes we identify the corresponding $d_{\parallel} = \delta$. For instance, if the SM arise from D3-branes and the $U(1)_X$ from D7-branes with an internal space having four large dimensions all parallel to the D7-brane world-volume ($\delta_X = d = 4$), we get for g_X the result in (4).

Little strings models

Another possibility for engineering extremely weak extra gauge symmetries is to consider a scenario which allows very small value of g_s . Such a possibility is provided by small instantons [30, 31] or Little String Theory (LST) [32, 33] where we localize the SM gauge group on Neveu-Schwarz (NS) branes (dual to the D-branes).

In the case of LST [32, 33], we start with a compactification on a six-dimensional space of volume V_6 with the Planck mass given by (1) (up to a factor 2 in the absence of an orientifold). The internal space is taken as a product of a two-dimensional space, of volume V_2 , times a four-dimensional compact space, of volume V_4 . However, instead of D-brane discussed above, we assume that the SM degrees of freedom emerge on a stack of NS5-branes wrapping the two-cycle of volume V_2 . We take for simplicity this to be a torus made of two orthogonal circles with radii R_1 and R_2 . The corresponding (tree-level) gauge coupling is given by:

$$g_{\text{SM}}^2 = \frac{R_1}{R_2} \quad (\text{Type IIA}) \quad \text{and} \quad g_{\text{SM}}^2 = \frac{1}{R_1 R_2 M_s^2} \quad (\text{Type IIB}); \quad (6)$$

thus, an order one SM coupling imposes $R_1 \simeq R_2 \simeq M_s^{-1}$. On the other hand, the $U(1)_X$ is supposed to appear in the bulk and has a coupling given by (2). If $U(1)_X$ arises from a D9-brane then:

$$M_{\text{Pl}}^2 = \frac{8}{g_s^2} M_s^8 \frac{V_2 V_4}{(2\pi)^6}, \quad \text{and} \quad g_X^2 = \frac{(2\pi)^7 g_s}{V_2 V_4 M_s^6}. \quad (7)$$

¹ A point worth noting at this juncture, however, is that there are several stellar systems that exhibit a mild preference for an over-efficient cooling mechanism when compared to theoretical models [25]. Thus, the argument can be turned around and the anomalous cooling could be interpreted as evidence for $U(1)_X$ production in dense star cores.

Now, taking all the internal space radii to be of the order of the string length, $M_s^6 V_2 V_4 \simeq (2\pi)^6$, leads to:

$$g_X \simeq \sqrt{32\pi} \sqrt{\frac{M_s}{M_{\text{Pl}}}} \sim 5 \times 10^{-7} \left(\frac{M_s}{10 \text{ TeV}} \right)^{1/2}. \quad (8)$$

Note however that the $U(1)_X$ from a D-brane does not interact directly with the electrons of the SM on the NS5-brane. Such interaction could arise via a closed string exchange which is likely to be suppressed by two powers of the string coupling, leading to an effective interaction of the order of 10^{-14} .

Small instantons models

In heterotic strings compactified on $K3$, of volume V_{K3} fibered over a two-dimensional base P^1 of volume V_{P^1} with integrated volume $\langle V_{K3} V_{P^1} \rangle$, the Planck mass reads:

$$M_{\text{Pl}}^2 = \frac{64\pi}{g_s^2} M_s^8 \langle V_{K3} V_{P^1} \rangle. \quad (9)$$

Taking the limit of instanton small size leads to emergence of a gauge group, identified with the SM one, supported at particular points on $K3$. The corresponding gauge coupling reads:

$$g_{\text{SM}}^2 = \frac{2\pi^2}{M_s^2 \langle V_{K3} V_{P^1} \rangle}, \quad (10)$$

implying that to give phenomenologically acceptable values, the compactification radius should remain of order of the string scale. The $U(1)_X$ is identified within the bulk theory descending from the ten-dimensional gauge symmetry:

$$g_X = \frac{g_s}{2} \frac{1}{M_s^3 \sqrt{\langle V_{K3} V_{P^1} \rangle}} = 4 \sqrt{\pi} \frac{M_s}{M_{\text{Pl}}} \sim 6 \times 10^{-14} \frac{M_s}{10 \text{ TeV}}. \quad (11)$$

Taking $\langle V_{K3} V_{P^1} \rangle \simeq \langle V_{K3} \rangle \langle V_{P^1} \rangle$, we see that the weakness of gravitational interactions, and a consequence of the $U(1)_X$ coupling, can be due either to a large volume of the $K3$ internal space or to a small string coupling: $\langle V_{K3} \rangle^{1/4} \sim \text{GeV}^{-1}$ or $g_s \sim 10^{-13}$ for $M_s \sim 10 \text{ TeV}$.

DARK PHOTON MASS GENERATION

We turn now to the generation of a mass for the dark photon.

Higgs mechanism

Let's denote by v_X the vacuum expectation value for the Higgs h_X that breaks the $U(1)_X$ symmetry. The simplest quartic potential $-\mu_X^2 h_X^2 + \lambda_X h_X^4$ leads to $v_X = \mu_X / \sqrt{2\lambda_X}$, a Higgs mass of order μ_X and a mass for the dark photon

$$m_X = \frac{g_X \mu_X}{\sqrt{2\lambda_X}} = \sqrt{2\pi} g_s \left(\frac{8}{g_s^2} \right)^{\delta/2d} \left(\frac{M_s}{M_{\text{Pl}}} \right)^{\delta/d} v_X. \quad (12)$$

This gives for $d = \delta = 6$:

$$m_X \sim \left(\frac{0.2}{g_s} \right)^{1/6} \left(\frac{M_s}{1000 \text{ TeV}} \right)^2 \left(\frac{v_X}{M_s} \right) \text{keV}. \quad (13)$$

Taking $v_X \simeq M_s$, this leads to a mass of order 0.1 to $1.4 \times 10^3 \text{ eV}$ when varying M_s from 10 to 1000 TeV, and g_s from 0.2 to 0.02. For this region of the parameter space, the gauge coupling varies in the range $4 \times 10^{-14} \lesssim g_X \lesssim 2 \times 10^{-11}$. Higher photon masses are of course easier to obtain with smaller number of d_{\parallel} dimensions. For example, an $M_s \sim 10 \text{ TeV}$, and $M_s \sim 100 \text{ TeV}$ lead respectively to $m_X \sim 6 \text{ keV}$, $g_X \sim 6 \times 10^{-10}$, and $m_X \sim 270 \text{ keV}$, $g_X \sim 6 \times 10^{-9}$ for $\delta_X = 4$, $d = 6$ and $g_s \sim 0.2$.

Stückelberg mechanism

Another possibility is that the abelian gauge field $U(1)_X$ becomes massive via a Stückelberg mechanism as a consequence of a Green-Schwarz (GS) anomaly cancellation [34, 35], which is achieved through the coupling of twisted Ramond-Ramond axions [36, 37]. The mass of the anomalous $U(1)_X$ can be unambiguously calculated by a direct one-loop string computation. Assuming the $U(1)_X$ arises from a brane wrapping δ_X dimensions among the d large dimensions, it is given by

$$m_X = \kappa \sqrt{\frac{V_a M_s^2}{V_X M_s^{\delta_X}}} M_s = \frac{\kappa}{(2\pi)^{\frac{\delta_X-2}{2}}} \left(\frac{\sqrt{8} M_s}{g_s M_{\text{Pl}}} \right)^{\frac{\delta_X-\delta_a}{d}} M_s, \quad (14)$$

where κ is the anomaly coefficient (which is in general an ordinary loop suppressed factor), V_a is the two-dimensional internal volume corresponding to the propagation of the axion field [10] and δ_a is the number of large dimensions in V_a . For $\delta_a = 0$, it leads to:

$$m_X = \frac{\kappa}{(2\pi)^{\frac{\delta_X-2}{2}}} \left(\frac{\sqrt{8} M_s}{g_s M_{\text{Pl}}} \right)^{\delta_X/d} M_s, \quad (15)$$

which gives $\sim 0.8\kappa \text{ keV}$ and $\sim 38\kappa \text{ keV}$ and for $M_s \sim 10 \text{ TeV}$ and $M_s \sim 100 \text{ TeV}$, respectively ($\delta_X = 4$, $d = 6$ and $g_s = 0.2$). For this region of the parameter space, the gauge coupling varies in the range $6 \times 10^{-10} \lesssim g_X \lesssim 6 \times 10^{-9}$. The case $\delta_a = 2$ and $\delta_X = d = 4$ leads to:

$$m_X = \frac{\kappa}{(2\pi)} \left(\frac{\sqrt{8} M_s}{g_s M_{\text{Pl}}} \right)^{\frac{1}{2}} M_s \sim 172 \kappa \left(\frac{0.2}{g_s} \right)^{\frac{1}{2}} \left(\frac{M_s}{10 \text{ TeV}} \right)^{\frac{3}{2}} \text{keV}. \quad (16)$$

For a concrete example of such case, consider 2 D7-branes intersecting in two common directions; namely, $D7_1 : 1234$

² Note that the $U(1)$ is not necessarily anomalous in four dimensions. A mass can be generated for a non-anomalous $U(1)$ by a six-dimensional (6d) GS term associated to a 6d anomaly cancellation in a sector of the theory.

and $D7_2 : 1256$, where 123456 denote the internal six directions. Take now 1234 large and 56 small (order the string scale) compact dimensions. The gauge fields of $D7_1$ have a suppression of their coupling by the 4-dimensional internal volume V_X while the states in the intersection of the two $D7$ branes see only the 12 large dimensions and give 6 dimensional anomalies, cancelled by an axion living in the same intersection, so V_a is the volume of 12 only.

$U(1)$ kinetic mixing

We have seen that the tiny couplings are not trivial to obtain and lead often to too small dark photon masses. This issue can be alleviated by resorting to the case where effective smaller couplings of $U(1)_X$ to SM states are obtained when the dark photons do not couple directly to the visible sector, but do it through kinetic mixing with ordinary photons. It can be generated by non-renormalisable operators, but it is natural to assume that it is generated by loops of states carrying charges $(q^{(i)}, q_X^{(i)})$ under the two $U(1)$'s and having masses m_i :

$$\epsilon_{\gamma X} = \frac{eg_X}{16\pi^2} \sum_i q^{(i)} q_X^{(i)} \ln \frac{m_i^2}{\mu^2} \equiv \frac{eg_X}{16\pi^2} C_{\text{Log}} \quad (17)$$

where μ^2 denotes the renormalization scale³, where we absorbed also the constant contribution. The effective coupling to SM is then:

$$g_{X,\text{eff}} = e\epsilon_{\gamma X} = \frac{\alpha_{\text{em}} g_X}{4\pi} C_{\text{Log}} \sim 6 \times 10^{-4} g_X C_{\text{Log}}. \quad (18)$$

We can try to fit both desired values of $g_{X,\text{eff}}$ and m_X . For a mass of the dark photon arising from a Higgs mechanism, we determine $g_X \sim m_X/M_s$, with $v_X \sim M_s$, this constrains:

$$\begin{aligned} C_{\text{Log}} &\simeq 1.7 \times 10^3 g_{X,\text{eff}} \frac{M_s}{m_X} \\ &\simeq 0.05 \left(\frac{g_{X,\text{eff}}}{8 \times 10^{-16}} \right) \left(\frac{M_s}{100 \text{ TeV}} \right) \left(\frac{m_X}{2.8 \text{ keV}} \right)^{-1}. \end{aligned} \quad (19)$$

A cancellation in the logarithm can be total, and the contribution appears at higher loops [38], or partial, for instance between particles with (order one) charges $(q^{(i)}, q_X^{(i)})$ and $(q^{(j)}, q_X^{(j)} = -q_X^{(i)})$ and masses $m_j = m_i + \Delta m_{ij}$. For $\Delta m_{ij} \ll m_i$, we have an approximation:

$$C_{\text{Log}} \sim \sum_{i,j} \frac{\Delta m_{ij}}{m_i}. \quad (20)$$

TOWARDS EXPLICIT MODELS

We shall now discuss more explicitly the emergence of such extra abelian gauge groups in D -brane models. The minimal

embedding of the SM particle spectrum requires at least three brane stacks [39] leading to three distinct models of the type $U(3) \otimes U(2) \otimes U(1)$ that were classified in [39, 40]. Only one of them, model (C) of [40], has baryon number as a gauge symmetry that guarantees proton stability (in perturbation theory), and can be used in the framework of low mass scale string compactifications. In addition, because the charge associated to the $U(1)$ of $U(2)$ does not participate in the hypercharge combination, $U(2)$ can be replaced by the symplectic $Sp(1)$ representation of Weinberg-Salam $SU(2)_L$, leading to a model with one extra $U(1)$ added to the hypercharge [41]. Note that the abelian factor associated to the $U(2)$ stack of D-branes couples to the lepton doublet, and consequently this anomalous $U(1)$ cannot be a good dark matter candidate, because the left-handed neutrinos make it unstable. One can add to these three stacks another D9-brane which will provide the $U(1)_X$ which will mix with the photon through loops of states living in the intersections of the D9 and the $U(3)$ and $U(1)$ stacks. The dark $U(1)_X$ is of course unstable as it decays to three ordinary photons. However, the partial decay width is found to be [42]

$$\Gamma_{X \rightarrow 3\gamma} \sim 10^{-28} \left(\frac{m_X}{2.8 \text{ keV}} \right)^9 \left(\frac{g_{X,\text{eff}}}{5 \times 10^{-16}} \right)^2 \text{ Gyr}^{-1}, \quad (21)$$

and so for the range of small gauge coupling considered here, the life-time is big enough to allow it be a viable candidate for dark matter.

Actually, the SM embedding in four D-brane stacks leads to many more models that have been classified in [43, 44]. The total gauge group of interest here,

$$\begin{aligned} G &= U(3)_C \otimes U(2)_L \otimes U(1)_1 \otimes U(1)_X \\ &= SU(3)_C \otimes U(1)_C \otimes SU(2)_L \otimes U(1)_L \otimes U(1)_1 \otimes U(1)_X, \end{aligned} \quad (22)$$

contains four abelian factors. The non-abelian structure determines the assignments of the SM particles. The quark doublet Q corresponds to an open string with one end on the color stack of D-branes and the other on the weak stack. The anti-quarks u^c and d^c must have one of their ends attached to the color branes. The lepton doublet and possible Higgs doublets must have one end on the weak set of branes. Per contra, the abelian structure is not fixed because the $U(1)_Y$ boson, which gauges the usual electroweak hypercharge symmetry, could be a linear combination of all four abelian factors. However, herein we restrict ourselves to models in which the bulk $U(1)_X$ does not contribute to the hypercharge, in order to avoid an unrealistically small gauge coupling. Of particular interest here are models (3) and (5) of reference [43]. The general properties of their chiral spectra are summarized in Table I and II. One can check by inspection that for both models the hypercharge,

$$\begin{aligned} q_Y &= -\frac{1}{3}q_C + \frac{1}{2}q_L + q_1 && \text{for model (3)} \\ q_Y &= \frac{2}{3}q_C + \frac{1}{2}q_L + q_1 && \text{for model (5)} \end{aligned} \quad (23)$$

³ In string theory, it is replaced by the string scale M_s .

TABLE I: Chiral fermion spectrum of the D-brane model (3).

Fields	Representation	q_C	q_L	q_1	q_X	q_Y
Q	$(\mathbf{3}, \mathbf{2})$	1	1	0	0	$\frac{1}{6}$
u^c	$(\bar{\mathbf{3}}, \mathbf{1})$	-1	0	-1	0	$-\frac{2}{3}$
d^c	$(\bar{\mathbf{3}}, \mathbf{1})$	-1	0	0	-1	$\frac{1}{3}$
L	$(\mathbf{1}, \mathbf{2})$	0	1	-1	0	$-\frac{1}{2}$
e^c	$(\mathbf{1}, \mathbf{1})$	0	0	1	1	1

TABLE II: Chiral fermion spectrum of the D-brane model (5).

Fields	Representation	q_C	q_L	q_1	q_X	q_Y
Q	$(\mathbf{3}, \mathbf{2})$	1	-1	0	0	$\frac{1}{6}$
u^c	$(\bar{\mathbf{3}}, \mathbf{1})$	-1	0	0	1	$-\frac{2}{3}$
d^c	$(\bar{\mathbf{3}}, \mathbf{1})$	-1	0	1	0	$\frac{1}{3}$
L	$(\mathbf{1}, \mathbf{2})$	0	1	-1	0	$-\frac{1}{2}$
e^c	$(\mathbf{1}, \mathbf{1})$	0	0	1	1	1

is anomaly free. In addition, the $U(1)_X$ is long-lived (because it only couples to the e^c and to either u^c or d^c) and therefore a viable DM candidate.

SUMMARY

We have investigated the possibility of identification of the light dark photon with one of the ubiquitous $U(1)$ gauge bosons of D-brane string theory constructions. We have first investigated how small the gauge coupling can be made. We found that open strings allow values as weak as $g_X \mathcal{O}(10^{-14})$ for a string scale of order $M_s \sim \mathcal{O}(10)$ TeV. For the case of six and four large extra dimensions, the Kaluza-Klein excitations appear above the GeV and MeV scales, respectively. They are very weakly coupled and decay quickly, but one could hope to observe their cumulative effect at TeV scales energies in future collider searches. The case of small instantons leads to similar conclusions with an interesting additional possibility: no large extra dimension below the TeV but a tiny string coupling. This possibility is also realized in Little String Theory but with a stronger gauge coupling $g_X \mathcal{O}(10^{-7})$. We have then looked at possible realization of models with dark photon masses in the range of keV. We have found two possibilities. A stringy Stückelberg mechanism generates masses in the desirable range in models with four large extra dimensions. If the dark photon mass is generated instead by a (low energy field theoretical) Higgs mechanism, then a kinetic mixing between the visible and dark $U(1)$ allows to get simultaneously the desired mass and coupling strength.

Here, we have left aside a couple of issues whose investigations are beyond the scope of this work. String models typically exhibit a large number of moduli fields that need to be fixed. In particular, the size of extra dimensions and couplings discussed here are vacuum expectation of such fields and all the hierarchies should be dynamically generated (see

[45] for example). Also, obtaining the correct relic density for such weakly coupled particles to form the main component of dark matter is challenging. The dark photons can not be thermally produced and one needs to resort to different mechanisms. A promising possibility is the proposal of [46], the desired abundance of dark photons can be generated by the quantum fluctuations during inflation for an appropriate set of parameters: the scale of inflation, reheating temperature and coupling to the curvature scalar.

Finally, we have discussed possible D-brane model can accommodate the excess of events with 3σ significance over background recently observed at XENON1T. The model is fully predictive, and can be confronted with future data from dark matter direct-detection experiments, LHC Run 3 searches, and astrophysical observations.

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