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.

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, $g \sim M_s/M_{Pl}$, 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 $\sim E/M_{Pl}$, 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 $\pm$ 15. Taken at face value this corresponds to a significance of roughly $3\sigma$, but unknown backgrounds from tritium decay cannot be reliably ruled out [15]. Although the excess is not statistical 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, $\rho_{DM} \sim 10^7 (n_{DM}/2.8$ keV)$^{-1}$ cm$^{-3}$, or else a sub-component of the DM population. Absorption of a $\sim 2.8$ 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 $g_{X,\text{eff}} \sim 2 \times 10^{-6} - 8 \times 10^{-16}$ [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 detection experiments [23,24]. For instance, rescaling the bounds $\lesssim 3 \times 10^{-10}$.

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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 [23]. Thus, the argument

$\lesssim 3 \times 10^{-10}$.
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 $\sim 5$ 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.

We start from ten-dimensional type I string theory compactified on a six-dimensional space of volume $V_6M_6^6$. The relation between the Planck mass, the string scale, the string coupling $g_s$, and the total volume of the bulk $V_6M_6^6$ reads:

$$M_{\text{Pl}}^2 = \frac{8}{g_s^2} M_6^8 \frac{V_6}{(2\pi)^6}.$$  \hspace{1cm} (1)

A hierarchy between the Planck and string scales can be due to either a large volume $V_6M_6^6 \gg 1$ 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 $M_6^{-1}$. The $U(1)_X$ gauge fields live on a D$(3+\delta_X)$-brane that wraps a $\delta_X$-cycle of volume $V_X$, while its remaining four dimensions extend into the uncompactified space-time. The corresponding gauge coupling is given by:

$$g_X^2 = \frac{(2\pi)^{d_X+1}}{V_X} \frac{g_s^2}{M_6^{d_X}}.$$  \hspace{1cm} (2)

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

$$g_X^2 = 2\pi g_s \left( \frac{8}{g_s^2} \right)^{\delta_X/d} \left( \frac{M_s}{M_\text{Pl}} \right)^{2\delta_X/d}.$$  \hspace{1cm} (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

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

where we have taken as reference values $g_s = 0.2$ and $M_s \sim 10$ 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][27]. 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

$$V_6 = (2\pi)^6 \prod_{i=1}^{d_1} R_i^6 \prod_{j=1}^{d_2} R_j^4,$$  \hspace{1cm} (5)

where now for each stack of D$p$-branes we identify the corresponding $d_i = \delta$. For instance, if the SM degrees of freedom emerge on a stack of D7-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).

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:

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

thus, an order one SM coupling imposes $R_1 \approx R_2 \approx M_s^{-1}$. 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:

$$M_{\text{Pl}}^2 = \frac{8}{g_s^2} M_6^8 \frac{V_2 V_4}{(2\pi)^6}, \quad \text{and} \quad g_X^2 = \frac{(2\pi)^3 g_s}{V_2 V_4 M_6^6}.$$  \hspace{1cm} (7)

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

$$g_X \approx \sqrt{32\pi} \frac{M_s}{M_\text{Pl}} \sim 5 \times 10^{-7} \left( \frac{M_s}{10 \text{ TeV}} \right)^{1/2}.$$  \hspace{1cm} (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 $10^{-14}$.

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 $< V_{K3} V_{P^1} >$, the Planck mass reads:

$$M_{\text{Pl}}^2 = \frac{64 \pi}{g_s^2} M_s^8 < V_{K3} V_{P^1} >.$$  \hspace{1cm} (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 $K^3$. The corresponding gauge coupling reads:

$$g_{\text{SM}}^2 = \frac{2\pi^2}{M_{\text{Pl}}^2 < V_{\text{pl}}>,}$$

(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:

$$g_X = \frac{g_s}{2 M_{\text{pl}}^2 < V_{K3} V_{pl}>,} = 4 \sqrt{\frac{\delta}{M_{\pi}}} \sim 6 \times 10^{-14} \frac{M_{\text{pl}}}{10 \text{ TeV}}.$$ 

(11)

Taking $< V_{K3} V_{pl}>, \ll < V_{K3}, > < V_{pl},>$ 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: $< V_{K3}, > \sim 1/4 \sim$ GeV$^{-1}$ or $g_s \sim 10^{-13}$ for $M_X \sim 10 \text{ TeV}$.

We turn now to the generation of a mass for the dark photon. 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 $\mu_X$ and a mass for the dark photon

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

(12)

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

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

(13)

Taking $v_X \approx M_{\pi}$, this leads to a mass of order 0.1 to $1.4 \times 10^3 \text{ eV}$ when varying $M_{\pi}$ 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 $4 \times 10^{-14} \leq g_X \leq 2 \times 10^{-11}$. Higher photon masses are of course easier to obtain with smaller number of brane dimensions. For example, an $M_{\pi} \sim 10 \text{ TeV}$, and $M_{\pi} \sim 100 \text{ TeV}$ 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 $d_X = 4, d = 6$ and $g_s \approx 0.2$.

Another possibility is that the abelian gauge field $U(1)_X$ becomes massive via a Stuckelberg 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 $^3U(1)_X$ can be unambiguously calculated by a direct one-loop string computation. Assuming the $U(1)_X$ arises from a brane wrapping $\delta_X$ dimensions among the $d$ large dimensions, it is given by

$$m_X = \sqrt{\frac{V_X S_{\text{area}}^2}{V_{M_{\text{pl}}^2}}} M_{\pi} = \frac{\sqrt{8}}{(2\pi)^{\delta/d}} \left( \frac{\sqrt{8}}{g_s M_{\pi}} \right)^{\delta/d} g_s M_{\pi},$$

(14)

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

$$m_X = \frac{\kappa}{(2\pi)^{\delta/d}} \left( \frac{\sqrt{8}}{g_s M_{\pi}} \right)^{\delta/d} M_{\pi}.$$

(15)

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

$$m_X = \frac{\kappa}{(2\pi)^{\delta/d}} \left( \frac{\sqrt{8}}{g_s M_{\pi}} \right)^{\delta/d} M_{\pi} \sim 172 \kappa \left( \frac{0.2}{g_s} \right)^{\delta/d} \left( \frac{M_{\pi}}{10 \text{ TeV}} \right)^{\delta/d} \text{ keV}.$$ 

(16)

For a concrete example of such case, consider 2 D7-branes intersecting in two common directions; namely, $D7_1 : 1234$ 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 $V_X$ is the volume of 12 only.

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 $(q_i, q_i^X)$ under the two $U(1)$’s and having masses $m_i$:

$$\epsilon_{X} = \frac{e g_X}{16 \pi^2} \sum_i q_i q_i^X \frac{m_i^2}{\mu^2} \equiv \frac{e g_X}{16 \pi^2} C_{\log}.$$ 

(17)

where $\mu^2$ denotes the renormalization scale$^3$, where we absorbed also the constant contribution. The effective coupling to SM is then:

$$g_{X,\text{eff}} = \frac{g_{X}}{\epsilon_{X}} = \frac{g_{X}}{4\pi} C_{\log} \sim 6 \times 10^{-4} g_X C_{\log}.$$ 

(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 $g_s \sim m_X/M_{\pi}$, with $v_X \sim M_{\pi}$, this constrains:

$$C_{\log} \approx 1.7 \times 10^3 g_{X,\text{eff}} \frac{M_{\pi}}{m_X} \approx 0.05 \left( \frac{g_{X,\text{eff}}}{8 \times 10^{-10}} \right) \left( \frac{M_{\pi}}{100 \text{ TeV}} \right) \left( \frac{m_X}{2.8 \text{ keV}} \right)^{-1}.$$ 

(19)

$^3$ In string theory, it is replaced by the string scale $M_{\text{pl}}$. 

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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.
A cancellation in the logarithm can be total, and the contribution appears at higher loops [33], or partial, for instance between particles with (order one) charges \((q_i^{(i)}, q_i^{(j)})\) and \((q_i^{(i)}, q_X^{(j)})\) and masses \(m_i = m_i + \Delta m_i\). For \(\Delta m_i \ll m_i\), we have an approximation:

\[
C_{\text{Log}} \sim \sum_{i,j} \frac{\Delta m_{ij}}{m_i}.
\]  

(20)

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 at least over three D-brane stacks [39] leading to three distinct models of the type \(U(3) \otimes U(2) \otimes U(1)\) 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]

\[
\Gamma_{\chi \rightarrow 3\gamma} \sim 10^{-28} \left( \frac{m_X}{2.8 \text{ keV}} \right)^9 \left( \frac{\xi_{\text{eff}}}{5 \times 10^{-16}} \right)^2 \text{ Gyr}^{-1},
\]  

(21)

and so for the range of small gauge coupling considered here, the life-time is big enough to allow it to 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 [33, 44]. The total gauge group of interest here,

\[
G = U(3)_C \otimes U(2)_L \otimes U(1)_L \otimes U(1)_X
\]  

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\]  

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,

\[
q_Y = -\frac{1}{3} q_c + \frac{1}{2} q_L + q_1
\]

for model \((3)\)

\[
q_Y = \frac{2}{3} q_c + \frac{1}{2} q_L + q_1
\]

for model \((5)\)

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.

In summary, we have investigated a model of light dark matter based on ubiquitous \(U(1)\) gauge bosons of D-brane string compactifications. We have shown that this model can accommodate the excess of events with \(3\sigma\) 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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