Illuminating the Hidden Sector of String Theory by Shining Light through a Magnetic Field

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Many models of physics beyond the Standard Model predict minicharged particles to which current and near future low-energy experiments are highly sensitive. Such minicharges arise generically from kinetic-mixing in theories containing at least two $U(1)$ gauge factors. Here, we point out that the required multiple $U(1)$ factors, the size of kinetic-mixing, and suitable matter representations to allow for a detection in the near future occur naturally in the context of string theory embeddings of the Standard Model. A detection of minicharged particles in a low energy experiment would likely be a signal of an underlying string theory and may provide a means of testing it.

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The absorption probability and the propagation speed of polarized light propagating in a magnetic field may depend on the relative orientation of the polarization and the magnetic field. These effects are known as vacuum magnetic dichroism and birefringence, respectively.

In 2006, the PVLAS collaboration 1 reported an anomalously large rotation of the polarization plane of light after its passage through a transverse magnetic field in vacuum. In Ref. 2, it was shown that such a signal may be originating from the dichroism caused by pair production of minicharged fermions of sub-eV mass and fractional electric charge. More recent measurements by the PVLAS collaboration with an improved apparatus 3 did not confirm this signal. Accordingly these new measurements provide a bound of roughly \( \epsilon \leq \frac{Q_f}{e} \lesssim \text{few } \times 10^{-7} \), for \( m_f \lesssim 0.1 \text{ eV} \). \( \text{(1)} \)

This is the best known laboratory bound on the existence of light minicharged particles demonstrating that optical experiments are a powerful tool to search for such particles.

Moreover, motivated by the initial PVLAS result it has been demonstrated that minicharged particles and hidden-sector $U(1)$ gauge bosons can also be searched for in a variety of other low-energy laboratory experiments and significant improvements in the sensitivity are expected in the near future (see discussion at the end of the paper).

In this letter, we argue that models with minicharged fermions can naturally and generically arise in string theory. Detection of minicharged particles would therefore not only address the fundamental question of charge quantization, but also provide insight into the underlying theory of nature.

Particles with a small, unquantized charge arise very naturally in so-called paraphoton 4 models, containing, beyond the usual electromagnetic $U(1)$ gauge factor, at least one additional hidden-sector $U(1)$ factor. The basic observation is that particles with paracharge get an induced electric charge proportional to some small mixing angle between the kinetic terms of photons and paraphotons 5. Moreover, in models containing more than one paraphoton with at least one paraphoton being exactly massless and one light, \( keV \gg m_{\gamma'} \neq 0 \), the prohibitively strong astrophysical bounds on the fractional charge, \( \epsilon \lesssim 2 \times 10^{-14} \), for \( m_f \lesssim \text{few keV} \), arising from energy loss considerations of stars 6, can be relaxed considerably.

In a simple model analysed in 10, there are two paraphotons: one massless and one light, and the fermion transforms in the bifundamental representation of these two $U(1)$ factors. In vacuum, the fermion acquires an electric charge \( \epsilon \) due to a kinetic-mixing between the photon and the two paraphotons. Importantly, however, this electric charge is reduced in the stellar plasma by a multiplicative factor \( \frac{m_{\gamma'}^2}{\omega_p^2} \), where \( \omega_p \sim \text{few keV} \) is the plasma frequency. This charge screening mechanism is caused by a partial cancellation between two paraphotons interacting with the bifundamental fermion 10. The vacuum value 10 is therefore perfectly compatible with astrophysical bounds (as well as cosmological bounds based on big bang nucleosynthesis) as long as

\[ m_{\gamma'} \lesssim 0.1 \text{ eV}. \] \( \text{(2)} \)

This minimal model can be supplemented by an axion-like spin-zero particle, coupled to the minicharged fermions 10. A triangle diagram then leads to a coupling of the axion-like particle to two photons. The resulting production of axion-like particles gives an additional (to the one from minicharged fermions) contribution to the vacuum magnetic dichroism and birefringence. One can then expect to have observable effects with even smaller values of \( \epsilon \), while still not being in conflict with astrophysics.

The purpose of this letter is to point out that the required multiple $U(1)$ factors, the size of kinetic-mixing, and suitable matter representations to allow for detection in near future experiments occur very naturally within the context of realistic extensions of the Standard Model (SM) based on string theory. It is a feature of our approach that we do not construct a model specifically for
FIG. 1: (a) One-loop diagram which contributes to kinetic-mixing in field theory, and (b) its equivalent in open string theory (from Ref. 14).
FIG. 2: Kinetic-mixing in open string models with SUSY breaking on “hidden” branes. The visible sector consists of a phenomenologically well determined supersymmetric configuration of D3-branes at a fixed point in the 6 dimensional compact manifold, possibly with D7-branes passing through to cancel local tadpoles. Global absence of tadpoles is assumed to require additional branes and/or anti-branes in the bulk. Closed string interactions are mediated from hidden to visible sector by cylinder diagrams, and are equivalent to Fig. 1(b).

This phenomenon generally arises in these models because the hidden sector of D-branes in the bulk carry additional hidden U(1) factors (possibly emerging from U(N)), which interact with the visible sector MSSM branes by exchanging closed string modes through the bulk (cf. Fig. 2). Such a closed string exchange (a cylinder diagram) can also be interpreted in the “open string channel” as a kinetic-mixing diagram as shown in Fig. 1(b); the one-loop open string diagram has a heavy string in the loop stretched between brane and anti-brane. The masses of the modes in the loop are given by their stretching energy proportional to their length (i.e. the distance between brane and anti-brane). The reason, the one-loop contributions do not cancel, is that the presence of anti-branes breaks supersymmetry, which is in fact integral to this particular scenario. Consequently there is a residual contribution to kinetic-mixing and hence $\chi$, which is again given by the amount of supersymmetry breaking.

Consider non-degenerate radii, with our three infinite dimensions, $d + p - 3$ large dimensions of radius $R_{4=p+d} = R$, and $9 - d - p$ small space dimensions of radius $R_{4=p+d+1..10} = r$, with the visible and hidden sectors living on stacks of Dp-branes wrapping the small dimensions (where $p \geq 3$), and with the distance between hidden and visible branes being generally of size the compact dimension. The result for the mixing parameter when the hidden $U(1)_b$ is unbroken can be written as

\[
\chi \sim \frac{\pi \alpha_p}{N X_a X_b} \left( \frac{2(8-p)/2}{\alpha_p} \frac{M_N}{M_P} \right)^{2(8-p)/6-p} \left( \frac{R}{r} \right) \frac{d-p+3}{d-p+1}. \tag{6}
\]

(For the present discussion the spatial dimensions of the branes is assumed to be $p$ for both visible and hidden sectors; for the more general set-up, see Ref. 13.) The integer $N$ is a factor corresponding to the $Z_N$ point-group symmetry of the fixed point, a typical value being $N = 3$. $X_a, b$ represent factors coming from the traces of Chan-Paton matrices in the vertex operators for photon emission off the open string loop, corresponding to the crosses in Figs. 1(b, 2) typically $X_{a, b} \sim 1$, but we will comment more precisely on the meaning of these factors below. Finally, $\alpha_p$ is the value of the gauge coupling, for which one may take $\alpha_p \sim 1/24$ (the MSSM unification value). For example if $p = d = 3$, then

\[
\chi \approx 5 \times 10^{-10} \left( \frac{R}{r} \right) \left( \frac{M_s}{10^{11}\text{GeV}} \right)^{4/3}, \tag{7}
\]

so that an $R/r \sim 10^4$ would give the right value. Varying $p \geq 3$ and $d$, there is considerable freedom to reduce the ratio $R/r$, whilst still having $\chi \sim 10^{-7}$ (cf. Fig. 3). (Note that the strong bounds on $M_s$ obtained in Ref. 13 from the astrophysical bounds on $\epsilon$ do not apply in general due to the charge screening mechanism proposed in 10.)

Before continuing with the generic consideration of volume and scale dependence, we should briefly digress, to consider a subtle and somewhat technical issue, namely the question of whether such $U(1)$s can remain massless (or at least very much lighter than the string scale) but still kinetically mix\(^1\). This is a delicate question because the kinetic-mixing diagram is also the diagram for a mass term mixing visible and hidden photons, and the one-loop open string diagrams one evaluates in fact correspond to two terms in the Lagrangian of the form

\[
m_{ab}^2 A^a_{\mu} A^b_{\mu} + \chi_{ab} F_{\mu \nu}^{(a)} F^{(b) \mu \nu}. \tag{8}
\]

The mass term $m_{ab}$ is a St"uckelberg mass mixing, which is associated with mixed anomalies and their cancellation via the Green-Schwarz mechanism 21, 22. Anomaly free $U(1)$s must have $m_{ab} = 0$. (Also note that $U(1)$s that are anomaly-free in 4d may still get St"uckelberg masses due to 6d anomalies.) In the second term, $\chi_{ab}$ is the kinetic mixing parameter. Since both of the terms in Eq. 8 arise from the same diagram, how can $\chi_{ab}$ be non-vanishing in an anomaly free theory where $m_{ab} = 0$?

The answer is that in order to get a contribution to the St"uckelberg mass one has to extract a $1/k^2$ pole from the appropriate one-loop integral. From the closed string point of view this corresponds to the St"uckelberg mass only getting contributions from massless closed string modes. Such contributions are blind to the location in the compact dimensions of the different sources. The non-pole contributions in this integral gives rise to $\chi_{ab}$. Importantly these contributions to $\chi_{ab}$ are from both massless and massive Kaluza Klein modes. The latter certainly do care about the location of the sources in the compact dimensions, and so contributions to $\chi_{ab}$ do not

\(^1\) We thank Mark Goodsell for extensive input and collaboration on these and related issues which are discussed in detail in Ref. 20.

\[\pi\]
generally cancel even though the contributions to $m_{ab}$ must. Therefore $\chi$ arises for $U(1)$s that are anomaly free provided that the anomaly-free combination is from branes that are located at different points or wrapping different cycles.

This is in fact generic in large volume compactifications. For example for branes parallel to orientifold planes, the anomaly-free $U(1)$ comes from the original brane plus its displaced orientifold image, and the two contributions to kinetic mixing do not cancel. In a future publication [20], this will be confirmed by showing explicit constructions where kinetic mixing occurs between anomaly-free $U(1)$s even if they come purely from D-branes in completely supersymmetric and tadpole-free configurations.

Let us now return to the generic implications for the volume and scale dependence that can be derived from Eq. (6), and comment on the fermion sector. A crucial ingredient for the electric charge screening mechanism in the stellar plasma is that there have to be two hidden sector $U(1)$s, and the hidden sector fermions have to have charge $(0, e, -e)$ under the visible and the two hidden sector $U(1)$ factors, respectively [11]. Note that this is generic in open string models with hidden D-brane sectors, since hidden sector fields arising from open strings stretched between hidden sector branes naturally fall into the bifundamental representation of the two hidden sector $U(1)$s.

The remaining question is whether there is any reason to expect the mass of the minicharged particles to be $\sim 0.1$ eV or smaller, and indeed there is. Since one of the $U(1)$s is by assumption unbroken, it is natural to expect some fermions on the hidden brane to be initially massless. However, as we have seen, the $U(1)$ mixes with the visible sector symmetries, and of course here the MSSM requires mass terms, namely the $\mu$-term for the Higgs. Generally, these induce two-loop mass terms in the hidden sector, as shown in Fig. 4. These contributions are diluted by the same volume factors ($V_{\|}, V_{\perp}$) that cause the dilution of gravity, given by [23]

$$M_{\text{st}}^2/M_P^2 \sim \alpha_p^2 V_{\|}/V_{\perp};$$

at one-loop the stretched states get a mass-splitting (the inner loop of Fig. 4) of order $(V_{\|}/V_{\perp})M_P^2/\mu$, where $\mu \sim 1$ TeV, and at two-loops the diagram receives another volume factor $V_{\|}/V_{\perp}$. In total, therefore, the mass induced in the hidden sector is

$$m_{\text{hidden}} = \alpha_p^{-4} (M_s^6/M_P^4 \mu) \sim \alpha_p^{-4} (M_W^2/M_P),$$

which is roughly of the right order of magnitude for $\alpha_p \sim 1/24$. Note, that no new scales beyond what is assumed for the MSSM have been introduced. The conclusion is quite general: sub eV masses are induced for hidden sector fermions if there are fermions with mass $\sim 1$ TeV in the visible sector.

Therefore, both closed and open string models not only predict the necessary extra $U(1)$ factors and correct fermion representation, but can also accommodate values of the kinetic-mixing parameter and fermion masses that may allow for detection in the near future.

As far as the gauge boson mass is concerned, the good news is that, naturally, a Higgs appears in the bifundamental representation of the hidden $U(1)$s, leaving automatically one mixed $U(1)$ massless, as required in the charge screening mechanism of [10]. It remains to be seen whether one can come up with a mechanism, perhaps based on accidental symmetries, to stabilize its small sub-eV scale (cf. Eq. (2)). Finally, there may still be room for an additional light spin-zero particle coupled to the hidden sector fermions which could then play the role of an axion-like particle [10, 23, 24].

There are a number of exciting possibilities to test such a scenario in laboratory experiments, allowing for experimental insights into string theory with less model dependence than astrophysical or cosmological considerations. The existence of minicharged particles can be tested [2] by improving the sensitivity of instruments for the detection of vacuum magnetic birefringence and dichroism [1, 23, 24, 27, 28, 29, 30]. Another sensitive tool
is Schwinger pair production in strong electric fields, as they exist, for example, in accelerator cavities [31]. A classical probe is the search for invisible orthopositronium decays [32, 33]. We expect that all these laboratory experiments will probe into the range $\epsilon \sim 10^{-9} - 10^{-6}$. Hidden-sector $U(1)_{\text{gauge}}$ bosons [6] and additional axion-like particles [34, 35], coupled to the minicharged fermions, may be observed in photon regeneration experiments [28, 29, 36, 37, 38, 39, 40, 41, 42, 43, 44] some of which have recently published data [45, 46]. These searches are complementary to and presently more sensitive than collider techniques based on the effect of kinetic mixing on precision electroweak observables [47, 48, 49].

In conclusion, many string theory models with intermediate string scales $M_s \sim 10^{11}$ GeV and/or large volumes predict the existence of minicharged particles with $\epsilon \gtrsim 10^{-9}$, testable with near future laboratory experiments.

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