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Illuminating the hidden sector of string theory by shining light through a magnetic field

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ABSTRACT

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.

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 [4–6]

$$\epsilon \equiv \frac{Q_f}{e} \lesssim \text{few} \times 10^{-7}, \quad \text{for } m_f \lesssim 0.1 \text{ eV}. \tag{1}$$

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 Letter).

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 [7] 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 [8]. 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 $m_f \lesssim$ few keV, arising from energy loss considerations of stars [9], can be relaxed considerably. In a simple model analyzed 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 ϵ 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 $m_{\gamma'}^2/\omega_p^2$, where $\omega_p \sim$ 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 (1) is therefore

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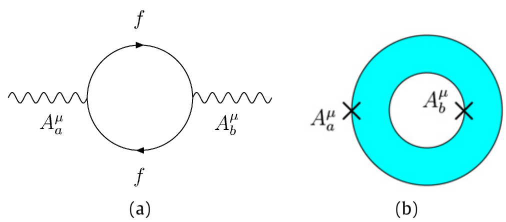

Fig. 1. (a) One-loop diagram which contributes to kinetic mixing in field theory, and (b) its equivalent in open string theory (from Ref. [14]).

perfectly compatible with astrophysical bounds (as well as cosmological bounds based on big bang nucleosynthesis) as long as

$$m_{\gamma'} \lesssim 0.1 \text{ eV}. \tag{2}$$

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 ϵ, 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 the purpose of producing minicharged particles but instead argue that the required minicharged particles are a generic, but also testable prediction of a large class of string theory models.

Let us begin by recalling the essentials of gauge kinetic mixing. It arises generally in theories that have, in addition to some visible $U(1)_a$, at least one other additional $U(1)_b$ factor in a hidden sector. In the basis in which the interaction terms have the canonical form, the pure gauge part of the Lagrangian for an arbitrary $U(1)_a \times U(1)_b$ theory can be written as (generalization to more than one additional $U(1)$ factor is straightforward)

$$\mathcal{L}_{\text{gauge}} = -\frac{1}{4}F^{\mu\nu}_{(a)}F_{(a)\mu\nu} - \frac{1}{4}F^{\mu\nu}_{(b)}F_{(b)\mu\nu} + \frac{\chi}{2}F^{\mu\nu}_{(a)}F_{(b)\mu\nu}, \tag{3}$$

where χ parametrizes the mixing. This parameter is directly linked to charge shifts. Notably, starting for example with two fermion species f_a and f_b with charges $(e, 0)$ and $(0, e)$, respectively, under $U(1)_a \times U(1)_b$, one finds that their charges, after diagonalization of the gauge kinetic term, are shifted by the amount

$$\epsilon \simeq \chi, \tag{4}$$

to leading order in $\chi \ll 1$ [8]. In other words, the hidden sector fermion f_b picks up a small charge $\epsilon \simeq \chi$ under the visible sector $U(1)_a$.

In view of the current experimental sensitivity, the following question immediately arises in this context:

Are there theoretically appealing and phenomenologically viable theories which naturally lead to a value of $\chi \simeq \epsilon \sim 10^{-7}$ for the kinetic mixing parameter?

We will argue that the answer is yes. Many SM extensions, and indeed most extensions coming from string theory, predict additional hidden $U(1)$ factors which can give rise to the kinetic mixing phenomenon (e.g., [11–16]).

In the context of field theory, a non-zero value of χ can be induced quite generally at the one-loop level when there are states which are simultaneously charged under both the visible $U(1)$ and the hidden $U(1)$ factors [8] (cf. the field theory diagram in Fig. 1(a)). The expected magnitude of χ in the field theory setting may be estimated [8,11] by considering the contribution of two fermions with charges (e_a, e_b) and $(e_a, -e_b)$ and masses m and m', respectively. Their joint contribution to χ is given by [8]

$$\chi \simeq (e_a e_b/(6\pi^2)) \log(m'/m), \tag{5}$$

and can be of order 10^{-7}, for natural values of $e_a e_b/(4\pi) \sim \alpha$, with $\alpha = e^2/(4\pi)$ being the electromagnetic fine-structure constant, and nearly degenerate masses, $m'/m \sim 1.0002$. Kinetic mixing could be avoided in such theories if the particle spectrum has some particular properties, as discussed in Ref. [11]. For example, if one or both of the $U(1)$ gauge factors sits within an unbroken non-abelian gauge symmetry, then kinetic mixing is not possible simply because of the tracelessness of the generators.

In the string theory context (e.g., [17]), the story is more subtle and more varied, but quite generally there are clear hints that string theory naturally leads to a generation of $\chi \neq 0$ [11–14,16], and can naturally give values of the right order of magnitude. The expected size of χ in a particular string theory setting can be investigated by performing the equivalent one-loop string theory calculation (cf. Fig. 1(b)).

Weakly coupled heterotic closed string models (in which the string scale, $M_s \approx 5 \times 10^{17}$ GeV, is close to the Planck scale, $M_P = 1.2 \times 10^{19}$ GeV) were treated this way in Ref. [11]. In these models, even if the particular string vacuum does not have the $U(1)$ explicitly embedded in a non-abelian group, there remains a "memory" of the underlying non-abelian structure, and the contributions vanish at leading order. Nevertheless, various effects below the string scale reintroduce kinetic mixing. In particular, low energy supersymmetry (SUSY) breaking will *always* split the matter multiplets contributing to the kinetic mixing by an amount of order the supersymmetry breaking scale in the hidden sector. If supersymmetry breaking is mediated by gravity, then the latter is roughly $M_{\text{SUSY}} \sim 10^{11}$ GeV and $\chi \sim 10^{-7}$ can naturally be obtained [11].

Other interesting classes of string models are mostly based on configurations of Dirichlet (D)-branes and/or fluxes. (Dp branes are membrane-like objects with p spatial dimensions, where open strings can be attached, thereby becoming matter and gauge fields.) Particularly interesting are the intermediate scale models with $M_s \sim \sqrt{M_W M_P} \sim 10^{11}$ GeV. The so-called bottom-up approach [18,19] is a good way to ascertain general properties of these models; the idea is to fix the gauge properties of the SM or minimal supersymmetric SM (MSSM) *locally* in the compactified space using D-branes, without fixing the *global* details of the compactification. The latter have only a minor affect on the visible sector phenomenology and may vary. However, the *global* set-up almost certainly requires anti-D-branes in the bulk which absorb certain types of unwanted tadpoles but leave supersymmetry broken. Inevitably, these set-ups predict kinetic mixing and therefore the existence of minicharged particles [13,14]. This phenomenon generally arises in these models because the hidden sector of $\overline{\text{D}}$-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 χ, which is again given by the amount of supersymmetry breaking.

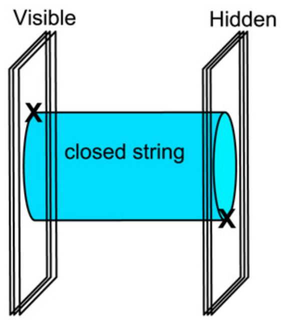

Fig. 2. Kinetic mixing in open string models with SUSY breaking on “hidden” branes. The visible sector consists of a phenomenologically well determined supersymmetric configuration of D3-branes at a fixed point in the 6-dimensional compact manifold, possibly with D7-branes passing through to cancel local 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).

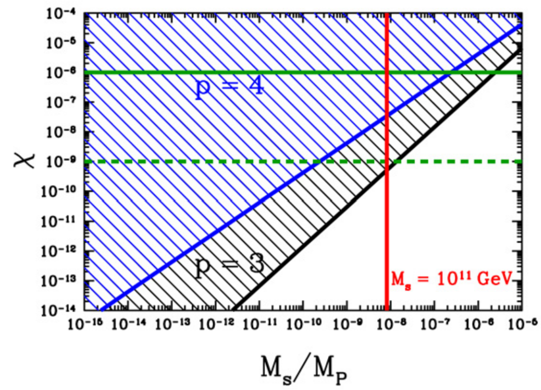

Fig. 3. Possible values for the kinetic mixing parameter χ, as a function of the string scale, in the bottom-up approach discussed in the text. Values in the shaded region above the black (lower) and blue (upper) lines are predicted in models with D3- and D4-branes, respectively, for $R/r \geqslant 1$. The red (vertical) line gives the largest string scale allowed by phenomenology in these models. The area above the green solid (horizontal) line is excluded by current experiments searching for minicharged particles while the green dashed line gives an idea of the expected sensitivity in the near future. (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this Letter.)

Consider non-degenerate radii, with our three infinite dimensions, $d + p - 3$ large dimensions of radius $R_{i=4,\dots,p+d} = R$, and $9 - d - p$ small space dimensions of radius $R_{i=p+d+1..10} = r$, with the visible and hidden sectors living on stacks of Dp-branes wrapping the small dimensions (where $p \geqslant 3$), and with the distance between hidden and visible branes being generically of size the compact dimension. The result for the mixing parameter when the hidden $U(1)_b$ is unbroken can be written as [13]

$$\chi \sim \pi \frac{\alpha_p}{N} X_a X_b \left(\frac{2^{(8-p)/2}}{\alpha_p} \frac{M_s}{M_P} \right)^{\frac{2(5-p)}{6-p}} \left(\frac{R}{r} \right)^{\frac{d-p+3}{6-p}}. \tag{6}$$

(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, α_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

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

so that an $R/r \sim 10^4$ would give the right value. Varying $p \geqslant 3$ 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 [13] from the astrophysical bounds on ϵ 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

$$m_{ab}^2 A_\mu^a A_b^\mu + \chi_{ab} F_{\mu\nu}^{(a)} F^{(b)\mu\nu}. \tag{8}$$

The mass term m_{ab} is a Stückelberg 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 Stückelberg masses due to 6d anomalies.) In the second term, χ_{ab} is the kinetic mixing parameter. Since both of the terms in Eq. (8) arise from the same diagram, how can χ_{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 Stückelberg 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 Stückelberg 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 χ_{ab}. Importantly these contributions to χ_{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 χ_{ab} do not generally cancel even though the contributions to m_{ab} must. Therefore χ 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 [10]. 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.

[1] We thank Mark Goodsell for extensive input and collaboration on these and related issues which are discussed in detail in Ref. [20].

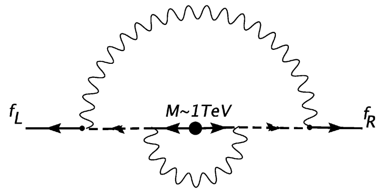

Fig. 4. Two-loop induced fermion masses in the hidden sector are directly related to the supersymmetry breaking. The dashed lines and gauge bosons are all stretched between visible and hidden branes, while the internal propagator is a fermion mass propagator of order 1 TeV.

The remaining question is whether there is any reason to expect the mass of the minicharged particles to be ~ 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 $U(1)_b$ mixes with the visible sector symmetries, and of course here the MSSM requires mass terms, namely the μ-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]

$$M_s^2/M_P^2 \sim \alpha_p^2 V_{||}/V_{\perp}; \tag{9}$$

at one-loop the stretched states get a mass-splitting (the inner loop of Fig. 4) of order $(V_{||}/V_{\perp})M_s^2/\mu$, 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

$$m_{\text{hidden}} = \alpha_p^{-4}\left(M_s^6/M_P^4\mu\right) \sim \alpha_p^{-4}\left(M_W^2/M_P\right), \tag{10}$$

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 ~ 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 [10,24].

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 [1,25–30]. 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 $U(1)$ gauge bosons [6] and additional axion-like particles [34,35], coupled to the minicharged fermions, may be observed in photon regeneration experiments [28,29,36–44] 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 [47–49].

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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