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Various theoretical and experimental considerations motivate models with high scale supersymmetry breaking. While such models may be difficult to test in colliders, we propose looking for signatures at much lower energies. We show that a keV line in the X-ray spectrum of galaxy clusters (such as the recently disputed 3.5 keV observation) can have its origin in a universal string axion coupled to a hidden supersymmetry breaking sector. A linear combination of the string axion and an additional axion in the hidden sector remains light, obtaining a mass of order 10 keV through supersymmetry breaking dynamics. In order to explain the X-ray line, the scale of supersymmetry breaking must be about 10^{11-12} GeV. This motivates high scale supersymmetry as in pure gravity mediation or minimal split supersymmetry and is consistent with all current limits. Since the axion mass is controlled by a dynamical mass scale, this mass can be much higher during inflation, avoiding isocurvature (and domain wall) problems associated with high scale inflation. In an appendix we present a mechanism for dilaton stabilization that additionally leads to $\mathcal{O}(1)$ modifications of the gaugino mass from anomaly mediation.

I. INTRODUCTION

The hierarchy problem has dominated much of the discussion on physics beyond the Standard Model (SM) in the past three decades, and supersymmetry emerged as the leading contender to solve this problem. In order to solve the problem fully, there was much anticipation that supersymmetry should be discovered very soon after the LHC began operating. Unfortunately, the LHC Run-I at 7–8 TeV placed a very strong lower limit, typically above a TeV, on superparticle masses [1], even though the quantitative limits are quite sensitive to the assumptions on the mass spectrum as well as the decay modes.

In addition, the discovered mass of the Higgs boson at 125 GeV [2] is higher than what was expected in the Minimal Supersymmetric Standard Model (MSSM). If we rely on the radiative corrections [3] from superparticles to raise the mass of the Higgs boson, we need to place scalar top quarks above a TeV. Finally, there have been long standing issues with supersymmetry, such as the absence of effects from large flavor-changing neutral currents, cosmological problems with the gravitino, and string moduli, which all prefer a supersymmetry spectrum with scalars around $m_{\mathrm{SUSY}} \approx$ 100–1000 TeV. If we take these hints seriously, direct searches for supersymmetry at collider experiments will be very difficult in the foreseeable future.

It is important to ask the question of whether there are alternative ways to find an experimental hint for su-persymmetry. We argue in this letter that the energy scale $m_{\mathrm{SUSY}}^2/M_{\mathrm{pl}} \approx$ keV may provide us with an indirect window to supersymmetry beyond the reach of accelerator experiments. Here M_{pl} is the reduced Plank scale, $M_{\mathrm{pl}} \simeq 2.4 \times 10^{18}$ GeV.

The recent observation of an unidentified line at about 3.5 keV in the X-ray spectrum of galaxy clusters [4, 5] hints at new particles at the keV energy scale. Although it has since been disputed by several other (non-) observations [6], it is interesting to consider that it (or a line observed in the future) could be a signal of dark matter decaying into photons. Even if we attribute this particular line to astrophysical processes, looking for new lines in X-rays is a continuing prospect. The possibility of linking such a low energy signal to physics at very high scales is an intriguing new avenue that we will explore in this work, using the 3.5 keV line as our guiding example. However, the types of models we will consider are rather generic and are not tied only to this specific experimental result.

Inspired by this observation, we investigate how supersymmetry may be relevant to the observed excess in X-rays from clusters of galaxies. Given the monochromatic line feature, it is tempting to consider a dark matter particle decaying into two photons. Note the Landau–Yang theorem that a vector cannot decay into two photons. Thus we consider a scalar particle decaying into two photons.[1] Then we need to understand the radiative stability of the keV energy scale, in addition to the origin of the keV scale itself.

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[1] However, there may be an alternative possibility that a fermion, such as a sterile neutrino, decays into a light active neutrino and a photon, through a suppressed mixing between the sterile and active neutrinos (see e.g. the review [7]).

The minute we assume that $m_{\rm SUSY}$ may be around 1000 TeV, there is a possible derived energy scale of $m_{\rm SUSY}^2/M_{\rm pl} \approx$ keV. One immediate possibility that comes to mind is that $m_{\rm SUSY}$ may be the scale of supersymmetry breaking itself, such as in low-energy gauge mediation [8], and keV is the mass scale of the gravitino or moduli. This possibility was examined already in the literature. For example, the gravitino [9] or moduli [10–12] may be dark matter. The decay of the moduli in this context may produce an X-ray signal from the clusters of galaxies [13, 14]. However, there are several non-trivial problems in gauge mediation, such as the μ-problem, the overproduction of gravitinos, and producing the correct Higgs boson mass (there are consistent models evading such difficulties, though, as in [8]).

We point out in this letter that there is an alternative possibility. $m_{\rm SUSY} \approx 100$–1000 TeV may be the gravitino mass. This possibility has attracted quite a bit of interest in the literature recently, starting from anomaly mediation [15] and leading up to pure gravity mediation [16] or minimal split supersymmetry [17]. In this case the scale of supersymmetry breaking is $\Lambda_{\rm SUSY} \approx (m_{\rm SUSY} M_{\rm pl})^{1/2} \approx 10^{12}$ GeV. The keV scale emerges parametrically as $\Lambda_{\rm SUSY}^4/M_{\rm pl}^3$.

If the new particle is a scalar, the keV mass scale must be protected against radiative corrections. The most effective mechanism is if the particle is a pseudo-Nambu–Goldstone-boson (pNGB). We call it generically an *axion* even though it may not have anything to do with the solution to the strong CP problem of QCD. The possibility that a pNGB may explain the origin of the 3.5 keV photon line has been also been considered elsewhere [18]. In this letter, we point out how such a pNGB can have a natural origin in the context of high scale supersymmetry.

The scalar decay proceeds through a dimension five operator suppressed by scale M with a rate $\Gamma \approx m^3/8\pi M^2$. For a 7 keV particle, the observed decay rate[2] is well-described by the energy scale $M \approx 0.1 M_{\rm pl}$. If interpreted as the axion decay constant $M = 32\pi^2 f$, $f \approx 10^{15}$ GeV. Therefore, discussing only two important scales, Λ_{SUSY} and $M_{\rm pl}$, seems well-warranted.

Given the large scale M and the coupling to electromagnetism, a well motivated possibility is to consider the scalar to be a modulus or axion field from string theory, where such properties can occur naturally, e.g. [19]. We will consider the "universal" or model-independent string axion, the defining properties of which we take to be the high scale decay constant and a universal coupling to all $F\widetilde{F}$. Thus the string axion couples to the hidden sector responsible for dynamical SUSY breaking and it may have a mass $m \sim \Lambda_{\rm SUSY}^2/f$. However, *this is not the case if the hidden sector contains an anomalous, global $U(1)$ symmetry that is spontaneously broken.* In this case, a second axion emerges which mixes with the string axion and leaves a massless eigenstate. Note that a spontaneously broken, anomalous $U(1)$ is a common feature of dynamical SUSY breaking models; the necessity of lifting flat directions in order to break supersymmetry typically induces non-zero vacuum expectation values, thus breaking global symmetries.

In the above scenario, where the SUSY breaking sector contains both a string axion and a hidden sector axion, instead of an exactly massless axion *we actually expect a non-zero, suppressed mass for the axion.* Gravity is believed to not respect global symmetries (see, e.g., [20]) and these violations may show up in a low-energy effective theory as higher dimension operators that explicitly break a global symmetry. Such explicit violations of the hidden sector $U(1)$ give a small, non-zero mass to the light axion. An axion with a keV scale mass and $f \approx 10^{15}$ GeV together with a high supersymmetry breaking scale suggest an explicit $U(1)$ violating mass-squared operator suppressed by $1/M_{\rm pl}^2 f$ leading to an axion mass $m \approx \Lambda^4/M_{\rm pl}^2 f$.

The rest of this letter explores an explicit example of the general scenario outlined above. We consider a string axion coupled to the IYIT model of dynamical supersymmetry breaking [21]. When the scale of supersymmetry breaking is large, $\Lambda_{\rm SUSY} \approx 10^{11-12}$ GeV, this model contains an axion that can produce the 3.5 keV X-ray line seen in [4, 5]. As discussed above, we believe the phenomena seen in our explicit example to be common. For example, we note that it occurs in other models of dynamical SUSY breaking such as the 4-1 model [22].

Following the demonstration of the 7 keV axion dark matter candidate, we address potential cosmological issues that arise in our explicit example. Some of these issues, such as isocurvature fluctuations, are common to setups based on our general mechanism. However, we believe the mechanisms employed to overcome certain cosmological issues in our explicit example can be applied in more general scenarios.

We also include two appendices. In the first, we give in detail the calculation of the axion spectrum for our explicit example. While the techniques there can be found throughout the literature, we include the derivation to keep our results self-contained. The second appendix presents a new mechanism for dilaton stabilization. As a result of this mechanism the axion develops an F-term; interestingly, this effects the gaugino masses at the $\mathcal{O}(1)$ level compared to their values from anomaly mediation.

II. AN EXPLICIT MODEL

As an explicit realization of our setup, we consider the minimal IYIT model [21]. The model consists of four quark superfields Q^i, $i = 1,\ldots,4$, charged under a $Sp(1) \simeq SU(2)$ gauge symmetry together with gauge singlets Z_{ij} in the **6** of the $SU(4)$ flavor symmetry. Supersymmetric $SU(2)$ gauge dynamics lead to a quantum

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[2] Of course, by changing slightly the scale M the rate can be below current experimental bounds.

modified moduli space with Pf $(QQ) = \Lambda^4$, where Λ is the dynamical scale of the theory. The gauge singlets couple to the quarks via a tree-level superpotential $W = \lambda Z_{ij} Q^i Q^j$. Supersymmetry is broken by the F-term for Z, which cannot be simultaneously satisfied with the quantum constraint.

The model contains a non-anomalous R-symmetry and an anomalous $U(1)_h$ symmetry under which Q and Z have charges $(0, 1)$ and $(2, -2)$, respectively. The $U(1)_h$ symmetry, which has a non-anomalous $\mathbb{Z}_4$ subgroup, is spontaneously broken by the quantum constraint. Therefore the phase of Q is the hidden sector axion a_h with decay constant $f_h \sim \Lambda$.

In addition to the fields Q and Z, we consider a string axion coupled to the gauge dynamics with strength $1/(32\pi^2 f_s)$. The $SU(2)$ dynamical scale then contains the string axion a_s,

$$\Lambda = \mu e^{-\frac{8\pi^2}{b_0 g^2}} e^{\frac{i a_s}{b_0 f_s}} = |\Lambda| e^{\frac{i a_s}{b_0 f_s}}, \tag{1}$$

where g is the gauge coupling and b_0 is the coefficient of the one-loop beta function. For $Sp(N_c)$ gauge theories $b_0 = 2(N_c + 1)$, so for the scenario at hand $b_0 = 4$. Presently, we consider the dilaton and fermion partners of the string axion to be stabilized and therefore non-dynamical. Otherwise, we would replace $i a_s$ by the chiral multiplet A_s in Eq. (1). In Appendix B we present a possible mechanism of stabilization.

The superpotential (with all indices suppressed) is

$$W = \lambda Z Q Q + \frac{\mathcal{A}}{\Lambda^2} \left(\text{Pf } (QQ) - \Lambda^4 \right), \tag{2}$$

where the quantum constraint is enforced by the Lagrange multiplier $\mathcal{A}$. In Appendix A we work out the effective theory and axion spectrum in detail while keeping track of factors of 4π using naïve dimensional analysis (NDA) [23, 24]. However, it is simple to see the basic results. Schematically, taking $QQ \sim \Lambda^2 e^{2i a_h/f_h}$ and replacing $\Lambda^4 \to \Lambda^4 e^{i a_s/f_s}$ in the quantum constraint, it is easy to see that the F-term for $\mathcal{A}$ produces a potential for the axions a_h and a_s,

$$V(a_h, a_s) \sim \Lambda^4 \left[1 - \cos \left(\frac{4 a_h}{f_h} - \frac{a_s}{f_s} \right) \right]. \tag{3}$$

The above potential[3] makes it clear that one linear combination of axions gains a mass of order Λ^2/f_h (for $f_s \gg f_h$) while the orthogonal combination is massless.[4]

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[3] The axion can be used as an inflaton in a chaotic inflation scenario since the Kähler potential for the axion has a shift symmetry [25]. In fact, the superpotential in Eq. (A13) may generate a potential that is assumed in a chaotic natural inflation in supergravity proposed in [26]. The light axion has an effective decay constant $b_0 f_s$ so we need $b_0 f_s > M_{\text{pl}}$ for successful natural inflation [27].

[4] Once a constant is added to the superpotential to cancel the cosmological constant, the field Z has a small expectation value $\langle Z \rangle \sim m_{3/2}$. Therefore, the massless axion is, in fact, a linear combination of a_h, a_s, and the R axion (the phase of Z).

As discussed previously, we generically expect quantum gravity to violate the $U(1)_h$ symmetry. Such explicit violations give the massless axion from above a small, non-zero mass. To this end, we consider the leading operator that violates the $U(1)_h$ symmetry while respecting the R-symmetry and the non-anomalous discrete $\mathbb{Z}_4 \subset U(1)_h$. With this criteria, the leading operator is a deformation of the superpotential of the form[5]

$$W \supset \lambda' \frac{1}{M_{\text{pl}}^4} Z(QQ)^3. \tag{4}$$

We note that since the vacuum is located at $\langle Z \rangle = 0$ [28] (see also Appendix A), the above is the leading order term to the superpotential containing Z.

There are lower dimension operators that explicitly violate $U(1)_h$, e.g. $\delta W = c\text{Pf } (QQ)/M_{\text{pl}}$, and therefore lead to different parametrics for the axion mass. These operators violate the R-symmetry and it is conceivable that this leads to their suppression, e.g. $c \sim m_{3/2}/M_{\text{pl}} \sim \Lambda^2/M_{\text{pl}}^2$ which gives a parametrically similar axion mass as the operator in Eq. (4). Thus we will consider only the operator of Eq. (4) in the following analysis.

The explicit violation of the $U(1)_h$ symmetry in Eq. (4) gives a mass to the light axion through the F-term for Z and is worked out in detail in the first appendix. To leading order, the mass of the light axion is

$$m_a^2 \approx \frac{2\lambda\lambda'}{(4\pi)^4} \frac{\Lambda^8}{M_{\text{pl}}^4 f_s^2} = 2 \frac{\lambda\lambda'}{(\lambda/4\pi)^4} \frac{F^4}{M_{\text{pl}}^4 f_s^2}. \tag{5}$$

where $F = \lambda \Lambda^2/(4\pi)^2$ is the scale of SUSY breaking [23] (see also Appendix A, Eq. (A12)). As emphasized previously in a more general context, here we explicitly see that the spectrum contains an axion with a suppressed mass $m_a \approx \Lambda^4/M_{\text{pl}}^2 f_s$.

Through its string axion component, the light axion couples directly to Standard Model photon operator $F\widetilde{F}$ with strength $1/f_s$. We can express the dynamical scale Λ in terms of the decay rate,

$$\Gamma = \frac{\alpha_{\text{EM}}^2}{64\pi^3} \frac{m_a^3}{f_s^2}, \tag{6}$$

as

$$\Lambda = \left(\frac{2\pi \alpha_{\text{EM}}^2}{\lambda\lambda'} \frac{m_a^5 M_{\text{pl}}^4}{\Gamma} \right)^{1/8}. \tag{7}$$

Experimental results [4, 5] determine $m_a \approx 7$ keV and $\Gamma \approx 5.7 \times 10^{-53}$ GeV. In the strongly-coupled vacuum the coupling λ becomes non-perturbative and $\mathcal{O}(4\pi)$. Taking $\lambda' \sim 1$, the supersymmetry breaking scale is

$$\sqrt{F} \sim 10^{11.5} \text{ GeV}, \tag{8}$$

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[5] A Kähler operator of the form $Z^* Z(QQ)^2/M_{\text{pl}}^4$ also respects the same symmetries and leads to the same order mass term as the operator in Eq. (4).

with a gravitino mass

$$m_{3/2} = \frac{F}{M_{\rm pl}} \sim \mathcal{O}(10)\text{-}\mathcal{O}(100) \text{ TeV}. \qquad (9)$$

We see that we have constructed an explicit model for the string axion coupled to a hidden supersymmetry breaking sector where the scale of supersymmetry breaking must be high to match the experimental X-ray line.

We also know more about the spectrum of this model. The field Z has charge 2 under the $\mathbb{Z}_4$ symmetry; it cannot couple to $W^\alpha W_\alpha$ to give the gauginos mass. Thus we are lead to anomaly mediation (although see Appendix B for modifications to the gaugino mass), which also fits nicely with the gravitino mass above and the known Higgs mass. We can easily incorporate this model in pure gravity mediation [16] or minimal split supersymmetry [17] models, to complete the extension of the SM.

III. COSMOLOGY

Although the recent observation of B-modes in the CMB by the BICEP2 collaboration [29] has been shown to be consistent with expectations from dust [30, 31], it has nevertheless renewed interest in models with a high inflationary scale, $H_{\rm inf} \sim 10^{14}$ GeV, at the upper end of the currently-allowed range [32]. Large values of $H_{\rm inf}$ present several cosmological challenges to any realistic model; for instance, isocurvature fluctuations of the nearly-massless axions must be suppressed. Furthermore, given a SUSY breaking scale of $\Lambda \sim 10^{11-12}$ GeV $< H_{\rm inf}$, domain walls are a potential problem due to the spontaneous breaking of the $\mathbb{Z}_4$ symmetry after inflation.

It should be noted that the domain wall issue is a model-specific one, which may be avoided by altering the dynamical SUSY breaking sector. For example, the $\mathbb{Z}_4$ symmetry may be gauged, or a model without a residual discrete symmetry may be chosen. Of course, it may be that $H_{\rm inf} < \Lambda$, although in this case isocurvature fluctuations may still pose a problem. For the purposes of this section, we will focus on the model presented in the previous section and present a consistent cosmological history that addresses the aforementioned issues in the presence of a high inflationary scale.

For $\Lambda \sim 10^{12}$ GeV, as required by the analysis in the previous section, the dynamical sector is weakly coupled during inflation. Thus domain walls would be formed after reheating, once the temperature fell below Λ. However, since Λ is a dynamical scale, we will show that it may be temporarily increased during inflation. Taking $H_{\rm inf} \lesssim \Lambda \lesssim \rho_{\rm inf}^{1/4}$ will ensure that the $\mathbb{Z}_4$ symmetry is broken during inflation without the IYIT vacuum energy dominating that of the inflaton.

Consider the gauge coupling of $SU(2)$ set by a gauge kinetic function $f = \{(1/g_0^2) + c(\phi^2/M_{\rm pl}^2)\}W^\alpha W_\alpha$ with W^α the hidden sector gauge field strength, ϕ a singlet,

and g_0 a coupling set by string theory and compactification. A superpotential of the form $\kappa Y(\phi^2 - M_{\rm pl}^2)$, for a coupling κ and a superfield Y, gives the singlet a large vev. Generically, ϕ has a Hubble induced soft mass during inflation; if κ is sufficiently small, $\kappa \lesssim H_{\rm inf}/M_{\rm pl}$, then $\langle\phi\rangle \sim 0$ and the effective coupling $1/g^2 \approx 1/g_0^2$ is strong. The dynamical scale, Λ, is easily of order 10^{15} GeV during inflation. Domain walls are thus avoided as the $\mathbb{Z}_4$ symmetry is broken during inflation by the IYIT meson condensate, so long as the reheating temperature is sufficiently low, $T_R \lesssim \Lambda$.

Even with $\Lambda \sim 10^{15}$ GeV, the light axion is still essentially massless compared to $H_{\rm inf}$, leading to a potential isocurvature problem. This may be avoided by further increasing the light axion mass during inflation, so that $m_a \sim H_{\rm inf}$. [6] This may be accomplished by including a $U(1)_h$-violating coupling of the IYIT quarks to the inflaton sector, giving an inflaton-field-dependent mass to the light axion. For concreteness, we assume a chaotic inflation scenario as described in [25]. Here the Kähler potential respects a shift symmetry on the inflaton chiral multiplet, H, which is broken by a mass term in the super potential, $W \supset mHX$. We couple the IYIT quarks to the X chiral multiplet in the Kähler potential, $K \supset X^\dagger X \mathrm{Pf}\,(QQ)/M^4 + \text{h.c.}$ Once the dynamical SUSY breaking sector becomes strongly coupled, this gives a mass term for a_h:

$$V \supset m^2 h^\dagger h \left[1 - \frac{\Lambda^4}{M^4}\cos\left(\frac{4a_h}{f_h}\right)\right]. \qquad (10)$$

Here, h, the psuedo-scalar component of H, is the inflaton. Taking $m^2 \langle h^\dagger h\rangle \sim H_{\rm inf}^2 M_{\rm pl}^2$ and $f_h \sim \Lambda$, we have $m_a \sim 4H_{\rm inf}\left(\Lambda M_P/M^2\right)$. Thus there will be no isocurvature problem for $M \lesssim 10^{17}$ GeV.

Having given the light axion a large mass during inflation, it remains to show that the appropriate misalignment can still be generated. If the coupling, c, of the singlet to the gauge kinetic term is complex, then the imaginary part may be removed by a shift in the string axion, $a_s \to a_s + \mathrm{Im}(c)\langle\phi\rangle/M_{\rm pl}$. When $\langle\phi\rangle = 0$, the axions will find their minimum at the origin. After inflation, however, we have $\langle\phi\rangle \sim M_{\rm pl}$, and the heavy axion will relax to its new minimum at $(4a_h/f_h - a_s/f_s) \sim \mathrm{Im}(c)$. The effective potential for the light axion is then approximately

$$V \sim \frac{2\lambda\lambda'}{(4\pi)^4}\frac{\Lambda^8}{M_{\rm pl}^4}\cos\left(2\,\mathrm{Im}(c) + \frac{2a_s}{f_s}\right). \qquad (11)$$

Since both axions were pinned to the origin during inflation, this generates a misalignment of $a \sim \mathrm{Im}(c)f_s$. In order that this misalignment is sufficient to reproduce

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[6] Note that the heavy axion mass, $m_{a'} \sim \Lambda$, is already heavy relative to $H_{\rm inf}$ due to the impact of the ϕ singlet on the dynamical scale.

the observed abundance of dark matter, we must have $\text{Im}(c) \sim 10^{-4}$.

Since ϕ was trapped at the origin during inflation, one may worry that coherent oscillations of ϕ about its new minimum would come to dominate the energy density of the Universe. Subsequent ϕ decays could lead to overproduction of winos. However, if the Hubble induced mass remains significant as ϕ relaxes, then ϕ will adiabatically roll to its minimum and such oscillations do not occur. It is straightforward to check that, at the time the inflaton decays, ϕ is displaced from its minimum at M_{pl} by an amount $\Delta\phi \sim H^2/(\kappa^2 M_{\text{pl}})$, where the Hubble constant H is evaluated at the inflaton decay time. For a reheating temperature $T_R \sim 10^9$ GeV and $\kappa \sim 10^{-5}$, $\Delta\phi \sim 10^{-7}$ GeV. This very small misalignment does not produce any appreciable amount of ϕ oscillations.

One may further worry that domain walls are formed from the spontaneous breaking of the $\mathbb{Z}_2$ symmetry on ϕ by $\langle\phi\rangle \neq 0$. These domain walls are an artifact of our choice of the function of ϕ in front of the gauge kinetic term and in the superpotential. Other functions of ϕ will do just as well. In particular, we can add an explicit $\mathbb{Z}_2$ breaking term $\epsilon\phi$ into the superpotential to collapse the domain walls without changing our main results.

The dilaton superpartner to the string axion may similarly become misaligned after inflation, depending on the exact form of its Kähler potential. A mechanism such as that proposed in [33] allows for the dilaton to adiabatically roll to its minimum, and therefore its misalignment is not dangerous.

IV. DISCUSSION

The 3.5 keV X-ray line observed in [4, 5] allows the exciting possibility that it may originate from dark matter. Tests of this dark matter hypothesis are predominately limited to indirect detection. Since the initial observation followup studies [6] have called into doubt this X-ray line. As there are now several conflicting observations and the possible implications are far reaching, it behooves us to continue to observe with future X-ray experiments. As briefly studied in [4], future X-ray observations of galaxy clusters by the Astro-H Observatory can distinguish the dark matter hypothesis from other astrophysical sources that may be masquerading as a dark matter signal. We also note that a promising place to look for a clean signal is from dwarf spheroidal galaxies. These smaller galaxies can be dominated by non-baryonic matter and yield a signal with smaller backgrounds. Even if this 3.5 keV line is a red herring, the types of models we have presented give an exciting link from a future X-ray signal to physics at very high scales. This motivates further dedicated time to observe these dwarf galaxies and galaxy clusters.

We briefly comment on other phenomenological consequences. For example, the wino is stable and will generically be a sub-dominant component of the dark matter in the Universe. It is in principle possible to indirectly detect the wino dark matter utilizing gamma ray observations of dwarf spheroidal galaxies or the galactic center [34]. However, it may be very challenging since the wino indirect cross-section scales as the density squared. As for the direct detection of the wino at the LHC, see [35].

Embedding our mechanism into pure gravity mediation gives generic predictions about the gluino mass. For example, for $m_{3/2} \approx 50-100$ TeV then the gluino mass is in a detectable range at the LHC. For a slightly larger gravitino mass, $m_{3/2} \approx 200$ TeV, we would expect $m_{\tilde{g}} \approx 4-5$ TeV, and the gluino would be hard to detect at the LHC. However, if the $\mathcal{O}(1)$ mass cancellation from the dilaton stabilization mechanism in Appendix B takes place, then the gluino mass can be $m_{\tilde{g}} \approx 2-3$ TeV even for $m_{3/2} \approx 200$ TeV, and is therefore detectable.

Motivated by recent experimental results in the X-ray spectrum of galaxy clusters and the current situation in particle physics beyond the SM, we have explored the possibility of linking a keV signal to supersymmetry breaking at a much higher scale, around 10^{11-12} GeV. This exciting experimental link between such different energy scales is possible through a light axion, a mixture of a string theory and a hidden (supersymmetry-breaking) sector axion, which gains only a small mass from the supersymmetry breaking sector.

As an example demonstrating this possibility, we constructed a model with dynamical supersymmetry breaking from the IYIT model coupled to a string axion. One sees explicitly that there is a linear combination of axions which does not gain a large mass, but only a small mass once the superpotential includes higher dimensional operators. This mass is directly related to the scale of the hidden sector, and thus supersymmetry breaking. Using the X-ray results as input, we find the scale of supersymmetry breaking to be $\mathcal{O}(10^{11.5})$ GeV. This scale fits nicely with models like pure gravity mediation or minimal split supersymmetry.

Rather than producing a light axion to explain an X-ray signal, one can instead construct similar models for the QCD axion. In this case one needs to suppress operators to even higher dimension to produce a lighter axion. Instead of using an $Sp(1)$ gauge group, a larger group such as $Sp(5)$ should be used. Then we have a model for high-scale supersymmetry breaking with a light QCD axion solving the strong CP problem.

Given the lack of experimental evidence for supersymmetry thus far, together with theoretical arguments for considering models which may be difficult to see at colliders in the immediate future, it can be fruitful to pursue new avenues for signals of supersymmetry. In this letter we have shown a model which is well-motivated theoretically and experimentally, and suggests a hint for supersymmetry in the keV sky.

Appendix A: Axion potential

In this appendix we derive the low-energy effective axion potential for the IYIT model discussed in the main text. In order to properly capture the axion dynamics—in particular, the mixing of the hidden sector axion and string axion—the Lagrange multiplier that enforces the $SU(2)$ quantum constraint must be kept in the spectrum. This is because, in the absence of the string axion, the hidden sector axion and Lagrange multiplier pair up to become heavy together. Therefore, this analysis differs from the usual situation where the Lagrange multiplier is immediately integrated out of the spectrum, and we feel it is worthwhile, especially for non-experts of supersymmetric dynamics, to carefully lay out the steps of the calculation. In order to elucidate the physics, we first describe the simpler case of the model with no tree-level superpotential and no string axion and then add these terms to find the axion spectrum quoted in the text.

1. $W_{\mathrm{tree}} = 0$

We begin by considering $SU(2)$ supersymmetric gauge theory with four quark superfields and no tree-level super-potential. Mesons $M^{ij} = \epsilon^{\alpha\beta} Q^i_\alpha Q^j_\beta$—in the **6** of the $SU(4)$ flavor symmetry—parameterize the moduli space of the low-energy supersymmetric vacua. In the quantum theory, instantons deform the moduli space and the mesons are subject to the constraint $\mathrm{Pf}\, M = \Lambda^4$ [36], where Λ is the $SU(2)$ dynamical scale and can always be made real by an anomalous $U(1)_h$ rotation under which Q has unit charge. This constraint may be enforced in the low energy theory through a Lagrange multiplier, $W = \mathcal{A}(\mathrm{Pf}\, M - \Lambda^4)$. In the following, we make use of the local Lie group isomorphism $SU(4) \simeq SO(6)$ to describe the flavor symmetry; in this language, the mesons are in the vector representation of $SO(6)$ and the quantum constraint is $M_i^2 = \Lambda^4$.

Let us describe the qualitative features of the low-energy vacua. The quantum constraint spontaneously breaks the $SO(6)$ flavor symmetry. At points of enhanced symmetry, the flavor symmetry is broken from $SO(6) \to SO(5)$ by the vacuum expectation value

$$\langle M_6 \rangle = \Lambda^2, \tag{A1}$$

giving rise to five massless Nambu-Goldstone bosons. The quantum constraint also spontaneously breaks the anomalous $U(1)_h$ symmetry under which Q has unit charge. The would-be NGB associated with the spontaneous breaking of $U(1)_h$ gets a mass of order Λ through the anomaly and should be integrated out of the low energy theory. This would-be NGB, analogous to the η' meson of QCD, is what we refer to as the axion a_h.

In summary, the quantum constraint, when satisfied as in Eq. (A1), breaks the flavor symmetry as $SO(6) \times U(1)_h \to SO(5)$ and five of the mesons are massless while the sixth one, the axion, gains a mass of order the dynamical scale. In the rest of this subsection, we show how this qualitative picture works out quantitatively in the effective theory. We then demonstrate how introducing a string-like axion leaves a massless axion in the low-energy theory.

In terms of the canonically normalized mesons $\hat{M} \equiv M/\Lambda$, the effective superpotential and Kähler potential are given by

$$W_{\mathrm{eff}} = \frac{1}{16\pi^2} \mathcal{A}\big(\hat{M}_i^2 - \Lambda^2\big) \tag{A2a}$$

$$K_{\mathrm{eff}} = \frac{1}{16\pi^2} K_{\mathrm{dyn}}(\hat{M}, \mathcal{A}), \tag{A2b}$$

where $\mathcal{A}$ is a Lagrange multiplier enforcing the quantum constraint. The factors of 4π are included to ensure that the effective theory becomes strongly coupled at the scale Λ and are counted using naïve dimensional analysis [23, 24]. We note that estimates using naïve dimensional analysis have an uncertainty factor of a few; we take this uncertainty to be implicit in our results and do not explicitly keep track of it. Since the quantum moduli space is smoothly described

by the meson fields [36], we can take a canonical kinetic term for the meson fields as an apporiximation. Our results are not sensitive to this approximation. Further, a kinetic term for $\mathcal{A}$ is generated at one-loop via the interaction $\mathcal{A}\hat{M}^2$ in the effective superpotential. Therefore, at leading order the dynamical Kähler potential is given by

$$K_{\text{dyn}}(\hat{M}, \mathcal{A}) \approx \hat{M}^\dagger \hat{M} + \kappa \mathcal{A}^\dagger \mathcal{A}$$

where $\kappa \approx 5$ since there are five light mesons—the Nambu-Goldstone bosons—running in the loop that generates the kinetic term for $\mathcal{A}$.

To study the $SO(6) \times U(1)_h/SO(5)$ vacuum we parameterize the mesons as

$$\hat{M} = e^{2A_h/f_h}\left(\hat{M}_a, \sqrt{\Lambda^2 - \hat{M}_a^2}\right), \ a = 1, \ldots, 5. \tag{A3}$$

The $\hat{M}_a$ are the five NGB supermultiplets associated with the breaking $SO(6) \to SO(5)$, while the axion supermultiplet A_h, with scalar component $s_h + ia_h$, is associated with the breaking of $U(1)_h$. Inserting this parameterization into Eq. (A2b), expanding around small field values, and requiring a canonical kinetic term for the axion, we find the axion decay constant is given by $f_h = \Lambda/(\sqrt{2}\pi)$. The superpotential is

$$W = \frac{\Lambda^2}{(4\pi)^2}\mathcal{A}\left(e^{4A_h/f_h} - 1\right), \tag{A4}$$

and the F-term for $\mathcal{A}$ gives the axion potential. In components, the vacuum lies at $\langle s_h \rangle = 0$ and the axion potential is

$$V(a_h) = \frac{\Lambda^4}{8\pi^2\kappa}\left(1 - \cos\frac{4a_h}{f_h}\right). \tag{A5}$$

The axion mass is easily seen to be

$$m_{a_h}^2 = \frac{2\Lambda^4}{\kappa\pi^2 f_h^2} = \frac{4}{\kappa}\Lambda^2.$$

Now we consider an additional string axion coupled to the $SU(2)$ gauge dynamics,

$$\mathcal{L} \supset -\int d^2\theta \frac{1}{16\pi^2}\frac{A_s}{f_s}\text{Tr}(W_\alpha^2) + \text{h.c.}, \ A_s = s_s + ia_s + \sqrt{2}\theta\psi_s + \ldots \tag{A6}$$

This coupling means that the $SU(2)$ dynamical scale now depends on A_s,

$$\Lambda^4 \to \Lambda^4 e^{A_s/f_s}.$$

It is a simple matter to find the effective potential including the string axion; we proceed exactly as above and find that the superpotential is

$$W = \frac{\Lambda^2}{(4\pi)^2}\mathcal{A}\left(e^{4A_h/f_h} - e^{A_s/f_s}\right) \tag{A7}$$

Here we will assume the dilaton in A_s to be heavy and decoupled by some unspecified dynamics, while in Appendix B we present a new method to fix the dilaton. Then, the low energy axion potential is given by

$$V(a_s, a_h) = \frac{\Lambda^4}{8\pi^2\kappa}\left[1 - \cos\left(\frac{4a_h}{f_h} - \frac{a_s}{f_s}\right)\right]. \tag{A8}$$

The above potential makes it clear that one linear combination of axions gets a mass of order Λ while the orthogonal direction is massless. It is a simple procedure to find the mass eigenstates; in the limit of $f_s \gg f_h$ the heavy and light modes are given by:

$$
\begin{aligned}
\text{Light } a: && m_a^2 &= 0, & a &\approx a_s + \frac{f_h}{4f_s}a_h + \mathcal{O}\left(\frac{f_h^2}{f_s^2}\right), & f_a &= \sqrt{16f_s^2 + f_h^2} \approx 4f_s \\
\text{Heavy } a': \ m_{a'}^2 &= \frac{1}{8\pi^2\kappa}\frac{\Lambda^4}{f_{a'}^2} \approx \frac{2\Lambda^4}{\kappa\pi^2 f_h^2}, & a' &\approx -a_h + \frac{f_h}{4f_s}a_s + \mathcal{O}\left(\frac{f_h^2}{f_s^2}\right), & f_{a'} &= \frac{f_s f_h}{\sqrt{16f_s^2 + f_h^2}} \approx \frac{f_h}{4}
\end{aligned}
\tag{A9}
$$

Note that the light axion picks up the larger decay constant, $f_a \approx b_0 f_s$.

2. $W_{\text{tree}} \neq 0$

Now we consider the theory with the tree level superpotential considered in the text,

$$W_{\text{tree}} = W_{\text{IYIT}} + W_{\cancel{U(1)}}, \tag{A10}$$

where $W_{\text{IYIT}} = \lambda ZQQ$ spontaneously breaks supersymmetry and $W_{\cancel{U(1)}}$ is a term which explicitly breaks the $U(1)_h$ symmetry.

We first consider $W_{\cancel{U(1)}} = 0$ and briefly review how the IYIT superpotential spontaneously breaks SUSY [21] and the location of the vacuum [28]. For small Z the low-energy theory is still described by mesons and the effective superpotential is

$$W_{\text{eff}} = \frac{1}{(4\pi)^2}\Big[\lambda\Lambda Z_i \hat{M}_i + \mathcal{A}\big(\hat{M}_i^2 - \Lambda^2\big)\Big] \tag{A11}$$

where the singlets Z_i are in the $\mathbf{6}$ of the $SO(6)$ flavor symmetry. SUSY is broken through the F-term for Z, which is incompatible with the quantum constraint $\hat{M}_i^2 = \Lambda^2$. The vacuum lies in the direction where the $SO(6)$ symmetry is broken to $SO(5)$; this gives rise to five massless NGBs while their supersymmetric scalar partners gain a SUSY breaking mass of order $\lambda\Lambda$. Note that the singlets have canonical kinetic terms, $K_{\text{eff}} \supset Z^\dagger Z$ (with no factors of 4π). Therefore the SUSY breaking scale is suppressed from the dynamical scale Λ by extra factors of 4π [23]

$$F = \frac{\lambda}{(4\pi)^2}\Lambda^2. \tag{A12}$$

The superpotential in Eq. (A11) is an O'Raifeartaigh model of SUSY breaking and therefore comes with a classically flat direction; namely, in the vacuum $\langle \hat{M}_6 \rangle = \Lambda$ the singlet Z_6 is massless at tree-level and its value is undetermined. For perturbative values of the coupling λ, the theory is calculable near the origin and one finds that there is a stable, local minimum located at $\langle Z_6 \rangle = 0$ [28].[7]

Turning on the explicit $U(1)_h$ breaking term, $W_{\cancel{U(1)}}$, the light axion in Eq. (A9) will gain a small, non-zero mass. It is simple to estimate the size of this mass; the explicit breaking term is of the form

$$W_{\cancel{U(1)}} = \frac{\lambda'}{M_{\text{pl}}^n}\mathcal{O}_{\cancel{U(1)}}$$

with $\mathcal{O}_{\cancel{U(1)}}$ a dimension $n+3$ operator that explicitly breaks $U(1)_h$. The light axion carries the string decay constant f_s (see Eq. (A9)) and the only other dimensionful scales in the problem are M_{pl} and Λ (since we are in the vacuum where $\langle Z_6 \rangle = 0$). Therefore, up to numerical factors, the light axion mass is

$$m_a^2 \sim \lambda' \frac{\Lambda^{n+4}}{M_{\text{pl}}^n f_s^2}.$$

In the text we consider the leading operator that breaks $U(1)_h$ while preserving $U(1)_R$, $\mathcal{O}_{\cancel{U(1)}} \sim \lambda' Z(QQ)^3 \sim \Lambda^3 \lambda'^{ijkl} Z_i \hat{M}_j \hat{M}_k \hat{M}_l$. For simplicity, we may take the $SO(6)$ preserving interaction with $\lambda'^{ijkl} = \lambda'\delta^{ij}\delta^{kl}$—flavor violating effects can be considered as perturbations around this and do not significantly change the results. Thus, we examine the effective superpotential

$$W_{\text{eff}} = \frac{1}{(4\pi)^2}\Big[\lambda\Lambda Z_i \hat{M}_i + \lambda'\frac{\Lambda^3}{M_{\text{pl}}^4} Z_i \hat{M}_i \hat{M}_j^2 + \mathcal{A}\big(\hat{M}_i^2 - \Lambda^2 e^{A_s/f_s}\big)\Big]. \tag{A13}$$

It is straightforward to compute the axion potential; there is a contribution from $\mathcal{A}$'s F-term (Eq. (A8)) while the explicit breaking manifests itself in the F-term for Z giving

$$\delta V(a_h, a_s) = \frac{\Lambda^4}{(4\pi)^4}\left[\lambda^2 + \lambda'^2\frac{\Lambda^8}{M_{\text{pl}}^8} + 2\lambda\lambda'\frac{\Lambda^4}{M_{\text{pl}}^4}\cos\left(\frac{4a_h}{f_h}\right)\right] \tag{A14}$$

---

[7] If we explicitly keep the heavy axion and Lagrange multiplier $\mathcal{A}$ in the effective theory, as in the meson parameterization in Eq. (A3), then the classically flat direction is not Z_6 but instead it is a linear combination of Z_6 and $\mathcal{A}$. The results of the preceding paragraph and reference [28] remain the same for this classically flat direction.

The full axion potential is given by Eq. (A8) plus the above contribution. To leading order in f_h/f_s the mass of the heavy axion is unchanged from Eq. (A9) while the light axion gains a mass of

$$m_a^2 \approx \frac{2\lambda\lambda'}{(4\pi)^4} \frac{\Lambda^8}{M_{\rm pl}^4 f_s^2} = \frac{2\lambda\lambda'}{(\lambda/4\pi)^4} \frac{F^4}{M_{\rm pl}^4 f_s^2}, \tag{A15}$$

where in the last equality we expressed the mass in terms of the SUSY breaking scale in Eq. (A12).

Finally, we comment on the value of the yukawa coupling λ. In the strongly coupled vacuum, λ quickly becomes non-perturbative as the wave-function for the quarks shrinks to zero. For $\lambda = 4\pi$ we cannot prove that $\langle Z \rangle = 0$ is a stable minimum [28]. Thus we must assume it is the case.

Appendix B: Dilaton Stabilization

Dilaton stabilization, or more generally, moduli stabilization, is a notorious problem in string theory. When the string axion couples to the sector responsible for dynamical SUSY breaking there are several typical issues. First, there is a runaway direction in which SUSY is preserved. More concretely, the vacuum energy is order $V \sim \Lambda^4$ where Λ is the dynamical scale of the hidden sector. With the string axion, this is modified to $V \sim \Lambda^4 e^{(\phi+\phi^*)/f_s}$, where ϕ is the scalar modulus of the string axion supermultiplet so that $s \equiv \phi + \phi^*$ is the dilaton. Clearly, the potential is minimized for $s \to -\infty$ with $V \to 0$ and SUSY is restored.

There are ways to stabilize moduli, such as KKLT [37] or racetrack scenarios [38]. However, these are supersymmetric preserving mechanisms, so both the dilaton and the string axion are fixed. One may want, as in this work, a mechanism which stabilizes the dilaton but leaves the string axion free. This is clearly a non-supersymmetric request, and therefore any such mechanism that achieves this must make use of the dynamical supersymmetry breaking sector or some other source of supersymmetry breaking. In these setups, the dilaton typically has a gravitino sized mass, $m_{3/2}$. Such scenarios are not easily constructed, and they typically have other issues, such as so-called dilaton domination of SUSY breaking [39]. Here, due to the dilaton's coupling to the Standard Model gauge sector, gauginos get very large masses, $m_{1/2} \gtrsim m_{3/2}$. In anomaly mediated SUSY breaking $m_{3/2} \sim \mathcal{O}(100)$ TeV, making the phenomenology uniteresting.

In this appendix, we suggest a novel mechanism that both stabilizes the dilaton and does not introduce a large gaugino mass despite a large F term for the dilaton. As we will see, the gaugino mass coming from the dilaton is comparable to its mass from anomaly mediation. This means that the anomaly mediated mass may be $\mathcal{O}(1)$ changed, which is exciting in its own right. The key observation in this setup is to take both the modulus and the IYIT singlet fields Z to live in a strongly coupled sector [40]. By analogy with composite models, we make the crucial assumption that a form of naïve dimensional analysis (NDA) also holds for these fields. Properly counting the factors of 4π coming from NDA then give the results outlined above.

A comment on notation: in this appendix ϕ refers to a modulus field like the string axion supermultiplet. In relation to string axion multiplet A_s and the hidden sector multiplet A_h considered in the text, ϕ is the linear combination of them that is light.[8] For clarity, we ignore $\mathcal{O}(1)$ constants such as the one-loop beta function coefficient b_0; these are easily restored and do not alter our results.

We assume that there is a Kähler coupling between the string axion multiplet and the singlet fields,

$$K \supset g(\phi + \phi^*) + h(\phi + \phi^*)Z^\dagger Z \tag{B1a}$$

$$\approx M(\phi + \phi^*) + \frac{1}{2}(\phi + \phi^*)^2 + \cdots + Z^\dagger Z \left(1 + \frac{\phi + \phi^*}{M} + \dots\right), \tag{B1b}$$

where the scale M corresponds to the Plank scale, $M \simeq 2.4 \times 10^{18}$ GeV. In this case, the vacuum energy is given by

$$V = K^{ZZ^\dagger} \left|\frac{\partial W}{\partial Z}\right|^2 \sim \frac{\lambda^2 \Lambda^4 e^{s/f_s}}{1 + \frac{s}{M} + \dots}. \tag{B2}$$

The potential is minimized for $s = -M + f_s \sim -M$ with vacuum energy $V \sim \lambda^2 \Lambda^4 e^{-M/f_s}/(f_s/M)$. While the dilaton is technically stabilized and SUSY is still broken, the vacuum energy is tiny for the typical string axion parameters we are considering as it is suppressed by $e^{-M/f_s} \sim e^{-10^3}$.

---

[8] In Eq. (A7) and the following discussion it is easy to see that an entire chiral superfield is left massless if the dilaton is not fixed. Note that this linear combination generically picks up the larger decay constant (e.g. section VI.F.4 of [41]), i.e. the string decay constant (see also Eq. (A9)). Therefore this ϕ really does behave like a string modulus and the results of this section apply more generally.

Moreover, the dilaton acquires a large F-term, $F_\phi \sim M m_{3/2}$:

$$F_\phi \approx K^{\phi\phi^*}\left(W_\phi + K_\phi W/M_{\rm pl}^2\right) \approx \frac{K_\phi W}{M_{\rm pl}^2} \approx M m_{3/2}. \tag{B3}$$

where we have taken the superpotential to contain a constant so as to cancel the cosmological constant, $W \sim Z\Lambda^2 e^{\phi/f_s} + m_{3/2}M_{\rm pl}^2$, and evaluated F_ϕ in the vacuum $\langle Z \rangle = 0$. As the string axion couples to the Standard Model sector, this F-term leads to a large gaugino mass,

$$\int d^2\theta \frac{\phi}{32\pi^2 f_s} W_{{\rm SM},\alpha} W_{\rm SM}^\alpha \Rightarrow m_{1/2} \sim \frac{M}{32\pi^2 f_s} m_{3/2} \gg m_{3/2}, \tag{B4}$$

with $32\pi^2 f_s \sim 10^{17}$ GeV for $f_s \sim 10^{15}$ GeV.

These results change if dilaton sector is strongly coupled. We then imagine that the proper low-energy effective theory becomes strongly coupled at the compactification scale. We are led to consider the string scale as a composite scale, which we label M_c, and apply the rules of naïve dimensional analysis.

Let us briefly review the rules of NDA [23, 24, 42]: multiply the effective action by an overall factor of $1/16\pi^2$, replace the composite fields as $\Phi \to 4\pi\Phi$, and relabel the cutoff $M \to M_c$. For example, the coupling of the dilaton to SM gauge fields in Eq. (B4) becomes

$$\frac{\phi}{M} W_{{\rm SM},\alpha} W_{\rm SM}^\alpha \to \frac{4\pi\phi}{16\pi^2 M_c} W_{{\rm SM},\alpha} W_{\rm SM}^\alpha. \tag{B5}$$

As the above operator is responsible for the string axion decay to photons considered in the text, it sets the scale M_c:

$$\int d^2\theta \frac{\phi}{4\pi M_c} W_{{\rm EM},\alpha} W_{\rm EM}^\alpha \supset \frac{a}{32\pi^2 f_s} F_{\rm EM} \tilde{F}_{\rm EM}, \tag{B6}$$

so that for the observed value $f_s \sim 10^{15}$ GeV we have $M_c \sim 10^{16}$ GeV.

Let us assume that the singlet fields Z also live in the strongly coupled sector; then we can view them as "composite" particles just like the string axion. Using NDA, the relevant terms in the effective Kähler potential are

$$K \sim \frac{1}{16\pi^2}\left[4\pi M_c(\phi + \phi^*) + \frac{(4\pi)^3}{M_c} Z^\dagger Z (\phi + \phi^*)\right]. \tag{B7}$$

The vacuum energy as a function of the dilaton is then of the same form as Eq. (B2) with $M = M_c/4\pi$ and $f_s = M_c/8\pi$. Since $M \sim f_s \sim 0.1 M_c$, we immediately see that the vacuum energy is no longer tiny: at the minimum $V'(s) = 0$ we have $s = -M + f_s \sim -f_s \sim -M_c/8\pi$ so that the vacuum energy is $V \sim \lambda^2\Lambda^4 e^{-1}/\mathcal{O}(1)$. The dilaton F-term is $F_\phi = M_c m_{3/2}/4\pi \approx f_s m_{3/2}$; this gives the gauginos a mass of order

$$m_{1/2} \approx \frac{F_\phi}{32\pi^2 f_s} m_{3/2} \approx \frac{1}{32\pi^2} m_{3/2}. \tag{B8}$$

This contribution to the gaugino mass from the dilaton is comparable in size to the gaugino mass coming from anomaly mediation.

In summary, we have outlined a mechanism to stabilize the dilaton—while leaving the string axion free—that is phenomenoligically viable with supersymmetry breaking and a string axion decay constant that could explain the 3.5 keV line, as described in the main text. Moreover, the stabilization mechanism may allow for the anomaly mediation relations for gaugino masses to be changed by an $\mathcal{O}(1)$ amount, which could prove useful for model building. The crucial assumption in achieving these results is that the effective action of the dilaton should become strongly coupled at the compactification scale. This led us to applying naïve dimensional analysis to ensure this behavior of the effective action. By also considering the singlets involved in SUSY breaking to live in the strongly coupled sector we obtain the stated results.

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[1] Georges Aad et al. (ATLAS Collaboration), "Search for squarks and gluinos with the ATLAS detector in final states with jets and missing transverse momentum us-
ing $\sqrt{s} = 8$ TeV proton–proton collision data," (2014), arXiv:1405.7875 [hep-ex]; CMS Collaboration (CMS Collaboration), "Search for supersymmetry in hadronic

final states using M_{T2} with the CMS detector at $\sqrt{s} = 8$ TeV," (2014).

[2] Georges Aad <u>et al.</u> (ATLAS Collaboration), "Observation of a new particle in the search for the Standard Model Higgs boson with the ATLAS detector at the LHC," Phys.Lett. **B716**, 1–29 (2012), arXiv:1207.7214 [hep-ex]; Serguei Chatrchyan <u>et al.</u> (CMS Collaboration), "Observation of a new boson at a mass of 125 GeV with the CMS experiment at the LHC," Phys.Lett. **B716**, 30–61 (2012), arXiv:1207.7235 [hep-ex].

[3] Yasuhiro Okada, Masahiro Yamaguchi, and Tsutomu Yanagida, "Upper bound of the lightest Higgs boson mass in the minimal supersymmetric standard model," Prog.Theor.Phys. **85**, 1–6 (1991); Y. Okada, Masahiro Yamaguchi, and T. Yanagida, "Renormalization group analysis on the Higgs mass in the softly broken supersymmetric standard model," Phys.Lett. **B262**, 54–58 (1991); John R. Ellis, Giovanni Ridolfi, and Fabio Zwirner, "On radiative corrections to supersymmetric Higgs boson masses and their implications for LEP searches," Phys.Lett. **B262**, 477–484 (1991); Howard E. Haber and Ralf Hempfling, "Can the mass of the lightest Higgs boson of the minimal supersymmetric model be larger than m(Z)?" Phys.Rev.Lett. **66**, 1815–1818 (1991).

[4] Esra Bulbul, Maxim Markevitch, Adam Foster, Randall K. Smith, Michael Loewenstein, <u>et al.</u>, "Detection of An Unidentified Emission Line in the Stacked X-ray spectrum of Galaxy Clusters," Astrophys.J. **789**, 13 (2014), arXiv:1402.2301 [astro-ph.CO].

[5] Alexey Boyarsky, Oleg Ruchayskiy, Dmytro Iakubovskyi, and Jeroen Franse, "An unidentified line in X-ray spectra of the Andromeda galaxy and Perseus galaxy cluster," (2014), arXiv:1402.4119 [astro-ph.CO].

[6] D. Malyshev, A. Neronov, and D. Eckert, "Constraints on 3.55 keV line emission from stacked observations of dwarf spheroidal galaxies," Phys.Rev. **D90**, 103506 (2014), arXiv:1408.3531 [astro-ph.HE]; Michael E. Anderson, Eugene Churazov, and Joel N. Bregman, "Non-Detection of X-Ray Emission From Sterile Neutrinos in Stacked Galaxy Spectra," (2014), arXiv:1408.4115 [astro-ph.HE]; Takayuki Tamura, Ryo Iizuka, Yoshitomo Maeda, Kazuhisa Mitsuda, and Noriko Y. Yamasaki, "An X-ray Spectroscopic Search for Dark Matter in the Perseus Cluster with Suzaku," (2014), arXiv:1412.1869 [astro-ph.HE].

[7] Alexander Kusenko, "Sterile neutrinos: The Dark side of the light fermions," Phys.Rept. **481**, 1–28 (2009), arXiv:0906.2968 [hep-ph].

[8] Koichi Hamaguchi, Masahiro Ibe, Tsutomu T. Yanagida, and Norimi Yokozaki, "Testing the Minimal Direct Gauge Mediation at the LHC," (2014), arXiv:1403.1398 [hep-ph].

[9] T. Moroi, H. Murayama, and Masahiro Yamaguchi, "Cosmological constraints on the light stable gravitino," Phys.Lett. **B303**, 289–294 (1993).

[10] Andre de Gouvea, Takeo Moroi, and Hitoshi Murayama, "Cosmology of supersymmetric models with low-energy gauge mediation," Phys.Rev. **D56**, 1281–1299 (1997), arXiv:hep-ph/9701244 [hep-ph].

[11] J. Hashiba, M. Kawasaki, and T. Yanagida, "Cosmological constraint on the string dila-

[12] T. Asaka, J. Hashiba, M. Kawasaki, and T. Yanagida, "Cosmological moduli problem in gauge mediated supersymmetry breaking theories," Phys.Rev. **D58**, 083509 (1998), arXiv:hep-ph/9711501 [hep-ph].

[13] M. Kawasaki and T. Yanagida, "Constraint on cosmic density of the string moduli field in gauge mediated supersymmetry breaking theories," Phys.Lett. **B399**, 45–48 (1997), arXiv:hep-ph/9701346 [hep-ph].

[14] Alexander Kusenko, Michael Loewenstein, and Tsutomu T. Yanagida, "Moduli dark matter and the search for its decay line using Suzaku X-ray telescope," Phys.Rev. **D87**, 043508 (2013), arXiv:1209.6403 [hep-ph].

[15] Gian F. Giudice, Markus A. Luty, Hitoshi Murayama, and Riccardo Rattazzi, "Gaugino mass without singlets," JHEP **9812**, 027 (1998), arXiv:hep-ph/9810442 [hep-ph]; James D. Wells, "Implications of supersymmetry breaking with a little hierarchy between gauginos and scalars," (2003), arXiv:hep-ph/0306127 [hep-ph]; "PeV-scale supersymmetry," Phys.Rev. **D71**, 015013 (2005), arXiv:hep-ph/0411041 [hep-ph].

[16] M. Ibe, Takeo Moroi, and T.T. Yanagida, "Possible Signals of Wino LSP at the Large Hadron Collider," Phys.Lett. **B644**, 355–360 (2007), arXiv:hep-ph/0610277 [hep-ph]; Masahiro Ibe and Tsutomu T. Yanagida, "The Lightest Higgs Boson Mass in Pure Gravity Mediation Model," Phys.Lett. **B709**, 374–380 (2012), arXiv:1112.2462 [hep-ph]; Masahiro Ibe, Shigeki Matsumoto, and Tsutomu T. Yanagida, "Pure Gravity Mediation with $m_{3/2} = 10$-100 TeV," Phys.Rev. **D85**, 095011 (2012), arXiv:1202.2253 [hep-ph].

[17] Nima Arkani-Hamed, Arpit Gupta, David E. Kaplan, Neal Weiner, and Tom Zorawski, "Simply Unnatural Supersymmetry," (2012), arXiv:1212.6971 [hep-ph].

[18] Tetsutaro Higaki, Kwang Sik Jeong, and Fuminobu Takahashi, "The 7 keV axion dark matter and the X-ray line signal," (2014), arXiv:1402.6965 [hep-ph]; Joerg Jaeckel, Javier Redondo, and Andreas Ringwald, "A 3.55 keV hint for decaying axion-like particle dark matter," (2014), arXiv:1402.7335 [hep-ph]; Hyun Min Lee, Seong Chan Park, and Wan-Il Park, "Cluster X-ray line at 3.5 keV from axion-like dark matter," (2014), arXiv:1403.0865 [astro-ph.CO]; Farinaldo S. Queiroz and Kuver Sinha, "The Poker Face of the Majoron Dark Matter Model: LUX to keV Line," Phys.Lett. **B735**, 69–74 (2014), arXiv:1404.1400 [hep-ph].

[19] Peter Svrcek and Edward Witten, "Axions In String Theory," JHEP **0606**, 051 (2006), arXiv:hep-th/0605206 [hep-th].

[20] Tom Banks and Nathan Seiberg, "Symmetries and Strings in Field Theory and Gravity," Phys.Rev. **D83**, 084019 (2011), arXiv:1011.5120 [hep-th].

[21] Ken-Iti Izawa and Tsutomu Yanagida, "Dynamical supersymmetry breaking in vector - like gauge

theories," Prog.Theor.Phys. **95**, 829–830 (1996), arXiv:hep-th/9602180 [hep-th]; Kenneth A. Intriligator and Scott D. Thomas, "Dynamical supersymmetry breaking on quantum moduli spaces," Nucl.Phys. **B473**, 121–142 (1996), arXiv:hep-th/9603158 [hep-th].

[22] Michael Dine, Ann E. Nelson, Yosef Nir, and Yuri Shirman, "New tools for low-energy dynamical supersymmetry breaking," Phys.Rev. **D53**, 2658–2669 (1996), arXiv:hep-ph/9507378 [hep-ph].

[23] Markus A. Luty, "Naive dimensional analysis and supersymmetry," Phys.Rev. **D57**, 1531–1538 (1998), arXiv:hep-ph/9706235 [hep-ph].

[24] Andrew G. Cohen, David B. Kaplan, and Ann E. Nelson, "Counting 4 pis in strongly coupled supersymmetry," Phys.Lett. **B412**, 301–308 (1997), arXiv:hep-ph/9706275 [hep-ph].

[25] M. Kawasaki, Masahide Yamaguchi, and T. Yanagida, "Natural chaotic inflation in supergravity," Phys.Rev.Lett. **85**, 3572–3575 (2000), arXiv:hep-ph/0004243 [hep-ph].

[26] Renata Kallosh, Andrei Linde, and Bert Vercnocke, "Natural Inflation in Supergravity and Beyond," (2014), arXiv:1404.6244 [hep-th].

[27] Kazuya Yonekura, "Notes on natural inflation," (2014), arXiv:1405.0734 [hep-th]; Michael Dine, Patrick Draper, and Angelo Monteux, "Monodromy Inflation in SUSY QCD," (2014), arXiv:1405.0068 [hep-th].

[28] Z. Chacko, Markus A. Luty, and Eduardo Ponton, "Calculable dynamical supersymmetry breaking on deformed moduli spaces," JHEP **9812**, 016 (1998), arXiv:hep-th/9810253 [hep-th].

[29] P.A.R. Ade et al. (BICEP2 Collaboration), "Detection of B-Mode Polarization at Degree Angular Scales by BICEP2," Phys.Rev.Lett. **112**, 241101 (2014), arXiv:1403.3985 [astro-ph.CO].

[30] Raphael Flauger, J. Colin Hill, and David N. Spergel, "Toward an Understanding of Foreground Emission in the BICEP2 Region," JCAP **1408**, 039 (2014), arXiv:1405.7351 [astro-ph.CO].

[31] R. Adam et al. (Planck Collaboration), "Planck intermediate results. XXX. The angular power spectrum of polarized dust emission at intermediate and high Galactic latitudes," (2014), arXiv:1409.5738 [astro-ph.CO].

[32] P.A.R. Ade et al. (Planck Collaboration), "Planck 2013 results. XXII. Constraints on inflation," Astron.Astrophys. **571**, A22 (2014), arXiv:1303.5082 [astro-ph.CO].

[33] Andrei D. Linde, "Relaxing the cosmological moduli problem," Phys.Rev. **D53**, 4129–4132 (1996), arXiv:hep-th/9601083 [hep-th]; Kazunori Nakayama, Fuminobu Takahashi, and Tsu-tomu T. Yanagida, "Gravity mediation without a Polonyi problem," Phys.Lett. **B714**, 256–261 (2012), arXiv:1203.2085 [hep-ph].

[34] Biplob Bhattacherjee, Masahiro Ibe, Koji Ichikawa, Shigeki Matsumoto, and Kohei Nishiyama, "Wino Dark Matter and Future dSph Observations," (2014), 10.1007/JHEP07(2014)080, arXiv:1405.4914 [hep-ph]; Timothy Cohen, Mariangela Lisanti, Aaron Pierce, and Tracy R. Slatyer, "Wino Dark Matter Under Siege," JCAP **1310**, 061 (2013), arXiv:1307.4082; JiJi Fan and Matthew Reece, "In Wino Veritas? Indirect Searches Shed Light on Neutralino Dark Matter," JHEP **1310**, 124 (2013), arXiv:1307.4400 [hep-ph].

[35] Georges Aad et al. (ATLAS Collaboration), "Search for charginos nearly mass degenerate with the lightest neutralino based on a disappearing-track signature in pp collisions at $\sqrt{(s)}$=8TeV with the ATLAS detector," Phys.Rev. **D88**, 112006 (2013), arXiv:1310.3675 [hep-ex].

[36] Nathan Seiberg, "Exact results on the space of vacua of four-dimensional SUSY gauge theories," Phys.Rev. **D49**, 6857–6863 (1994), arXiv:hep-th/9402044 [hep-th].

[37] Shamit Kachru, Renata Kallosh, Andrei D. Linde, and Sandip P. Trivedi, "De Sitter vacua in string theory," Phys.Rev. **D68**, 046005 (2003), arXiv:hep-th/0301240 [hep-th].

[38] N.V. Krasnikov, "On Supersymmetry Breaking in Superstring Theories," Phys.Lett. **B193**, 37–40 (1987); Lance J. Dixon, "Supersymmetry Breaking in String Theory," (1990); J.A. Casas, Z. Lalak, C. Munoz, and Graham G. Ross, "Hierarchical Supersymmetry Breaking and Dynamical Determination of Compactification Parameters by Nonperturbative Effects," Nucl.Phys. **B347**, 243–269 (1990).

[39] Masha Baryakhtar, Edward Hardy, and John March-Russell, "Axion Mediation," JHEP **1307**, 096 (2013), arXiv:1301.0829 [hep-ph].

[40] Emilian Dudas, Andrei Linde, Yann Mambrini, Azar Mustafayev, and Keith A. Olive, "Strong moduli stabilization and phenomenology," Eur.Phys.J. **C73**, 2268 (2013), arXiv:1209.0499 [hep-ph].

[41] Jihn E. Kim and Gianpaolo Carosi, "Axions and the Strong CP Problem," Rev.Mod.Phys. **82**, 557–602 (2010), arXiv:0807.3125 [hep-ph].

[42] Aneesh Manohar and Howard Georgi, "Chiral Quarks and the Nonrelativistic Quark Model," Nucl.Phys. **B234**, 189 (1984); Howard Georgi and Lisa Randall, "Flavor Conserving CP Violation in Invisible Axion Models," Nucl.Phys. **B276**, 241 (1986).