Axiverse - the many axions of string theory

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Contents

1 Axions in string theory

We mainly follow [1] and [2].

1.1 Recap: Strong CP problem and the QCD axion

Most general gauge invariant Lagrangian of QCD includes topological P and CP violating Θ-term [1]

$$\mathcal{L} = -\Theta \frac{\alpha_s}{8\pi} G^a_{\mu\nu} \tilde{G}^{a,\mu\nu} \,, \tag{1}$$

where G is the field strength of the gluon field. Eq. (1) can be written as a total derivative and hence does not contribute to perturbation theory but plays a non-perturbative role (instantons). Non-observation of electric dipole moment of the neutron implies $|\bar{\Theta}| \lesssim 10^{-10}$.[2]

Strong CP problem [3]: *Why is* $\bar{\Theta} \in [-\pi, \pi]$ *so small?*

One solution: *Dynamical relaxation* [4, 5]. ⇒ Introduce a boson - the axion - a with shift symmetry

$$a \to a + const \,, \tag{2}$$

which is only violated by axionic couplings to gauge fields

$$\mathcal{L} = \frac{1}{2}\partial_\mu a \partial^\mu a - \frac{\alpha_s}{8\pi}\left(\bar{\Theta} + \frac{a}{f_a}\right) G^a_{\mu\nu}\tilde{G}^{a,\mu\nu} - \frac{e^2}{32\pi^2} C_{a\gamma}\frac{a}{f_a} F^{em}_{\mu\nu}\tilde{F}^{\mu\nu}_{em} + \frac{C_{ae}}{2f_a}\bar{e}\gamma^\mu\gamma_5\, e\,\partial_\mu a \,, \tag{3}$$

[1] see talk by Andreas Ringwald.
[2] The actual physical parameter is $\bar{\Theta} = \Theta + \arg\det M_q$ where M_q is the quark mass matrix.

with axionic decay constant f_a, electromagnetic field strength $F^{em}_{\mu\nu}$ and dimensionless couplings to photons $C_{a\gamma}$ and electrons C_{ae}. This effectively promotes $\bar{\Theta}$ to a field

$$\bar{\Theta}_{eff} = \bar{\Theta} + \frac{a}{f_a} \,. \tag{4}$$

The topological charge density $\propto \langle G^a_{\mu\nu}\tilde{G}^{a,\mu\nu}\rangle$ induced by QCD instantons at low energies provides a non-trivial potential for the axion with minimum at $\langle a \rangle = 0$ with parametrically small mass

$$m_a = \frac{m_\pi f_\pi}{f_a}\frac{\sqrt{m_u m_d}}{m_u + m_d} \simeq 0.6\,\text{meV} \times \left(\frac{10^{10}\text{GeV}}{f_a}\right) , \tag{5}$$

where m_π and f_π are the mass and decay constant of the pion and m_u and m_d are the masses of the u and d quark.

Astrophysical and cosmological constraints on f_a

For large axion decay constant f_a the coupling of the axion to photons $C_{a\gamma}/f_a$ and electrons C_{ae}/f_a is strongly suppressed. There are hard lower limits on f_a/C_{ai} from laboratory experiments as well as astrophysics, in particular non-observation of drastic energy losses in horizontal branch stars or white dwarfs:

$$\frac{f_a}{C_{a\gamma}} > 10^6\,\text{GeV}\,, \qquad \frac{f_a}{C_{ae}} > 10^9\,\text{GeV}\,. \tag{6}$$

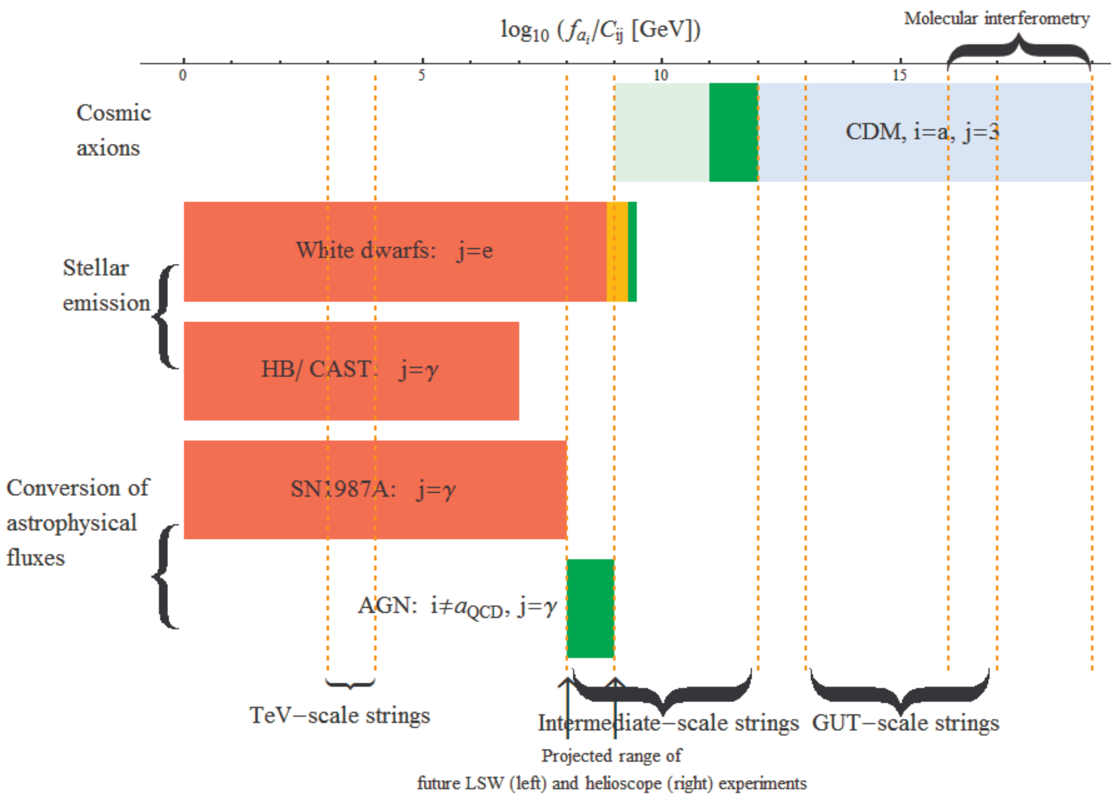

Figure 1: Constraints and hints for the coupling f_{a_i}/C_{ij}. Figure taken from [1]

There is also a softer cosmological upper bound. The axion can contribute significantly to the cold dark matter density for $f_a \gtrsim 10^{11}$ GeV, being non-thermally produced for $T_{RH} > f_a$ via a vacuum realignment mechanism to a coherent state of non-relativistic particles:

$$\Omega_a h^2 \simeq 0.71 \times \left(\frac{f_a}{10^{12}\,\mathrm{GeV}}\right)^{7/6} \left(\frac{\Theta_a}{\pi}\right)^2 , \tag{7}$$

where Θ_a is the initial misalignment angle. For $f_a > 10^{12}$ GeV and $\Theta_a \sim \mathcal{O}(1)$ there is the danger of the axions overclosing the universe. The observational constraints on f_a point to very high scale physics and hence it is natural to look for the axion and axion like particles (ALP's) in string theory. As we will see one can often even find an 'axiverse' [6] in string theory.

Furthermore, in string compactifications one should be aware of the cosmological moduli problem (CMP) induced by the scalar partners of the axions, the 'saxions' ϕ. The generically high number of moduli in string compactifications tend to cause problems with big bang nucleosynthesis (BBN), reheating and overclosure of the universe. For reheating temperature $T \gtrsim \mathcal{O}(1)$ MeV, necessary for BBN, this problem is solved if $m_\phi \gtrsim 30$ TeV.

1.2 Where do the string axions come from?

String theory is formulated in 10 dimensions such that 6 D have to be compactified to make contact with 4D real world physics. Starting with the 10D effective supergravity action one can derive a 4D low energy effective action that depends on the 10D effective field theory to start with and the compactification manifold and potential brane and orientifold configurations on the latter. Here we focus on type IIB compactified on Calabi-Yau manifolds X ($\Rightarrow \mathcal{N} = 2$ SUSY) with D7 branes and orientifold planes ($\Rightarrow \mathcal{N} = 1$ SUSY). The (extension of) the standard model is realized on a stack of D7 branes.

The bosonic field content of the 10D effective supergravity action of type IIB is given by the dilaton ϕ, the 2-form B_2, the metric G_{MN} and the n-forms C_0, C_2 and C_4. After compactification, there is a large number of complex scalar fields the moduli [7]:

- Size / Kähler moduli: $T_i = t_i + i\tau_i$ with $t_i = \mathrm{Vol}(D_i)$ and $\tau_i = \int_{D_i} C_4$ where D_i are the four-cycles of X for $i = 1, .., h^{1,1}$.[3]
- Axio-dilaton: $S = e^{-\phi} + iC_0$ with string coupling $g_s = e^{\phi}$.
- Shape / complex structure moduli $U_i = u_i + i\nu_i$ for $i = 1, .., h^{2,1}_-$

The imaginary parts of these scalar fields are dubbed axions because the corresponding couplings to $U(1)$ gauge field strengths $F_i\tilde{F}_i \sim \mathrm{tr}(F_i \wedge F_i)$ can be found in the 4D effective action of D7 branes wrapping 4-cycles D_i [8]:

$$\begin{aligned}
\mathcal{L} \supset &- \left(d\tau_\alpha + \frac{M_P}{\pi} A_i q_{i\alpha}\right) \frac{K_{\alpha\beta}}{8} \wedge * \left(d\tau_\beta + \frac{M_P}{\pi} A_j q_{j\beta}\right) \\
&+ \frac{1}{4\pi M_P} r^{i\alpha} \tau_\alpha \mathrm{tr}(F_i \wedge F_i) - \frac{r^{i\alpha} t_\alpha}{4\pi M_P} \mathrm{tr}(F_i \wedge *F_i) , \\
= &- \frac{K_{\alpha\beta}}{8} d\tau_\alpha \wedge *d\tau_\beta - \frac{K_{\alpha\beta} M_P}{8\pi} \left(d\tau_\alpha \wedge *A_j q_{j\beta} + A_i q_{i\alpha} \wedge *d\tau_\beta\right) + \frac{M_P^2}{2(2\pi)^2} A_i A_j q_{i\alpha} K_{\alpha\beta} q_{j\beta} \\
&+ \frac{1}{4\pi M_P} r^{i\alpha} \tau_\alpha \mathrm{tr}(F_i \wedge F_i) - \frac{r^{i\alpha} t_\alpha}{4\pi M_P} \mathrm{tr}(F_i \wedge *F_i) ,
\end{aligned} \tag{8}$$

[3] For simplicity we assume $h^{1,1} = h^{1,1}_+$, i.e. $h^{1,1}_- = 0$.

where $M_P = 2.4 \cdot 10^{18} GeV$ is the reduced Planck mass, $K_{\alpha\beta} = \frac{\partial^2 K}{\partial t_\alpha \partial t_\beta}$ is the Kähler metric and the A_i are the gauge fields with $F_i = d\,A_i$. Furthermore, $r^{i\alpha}$ are the expansion coefficients of the 2-forms $\hat{D}_i = r^{i\alpha}\omega_\alpha$ that are Poincaré dual to the 4-cycles D_i, i.e.

$$r^{i\alpha} = l_s^{-4} \int_{D_i} \tilde{\omega}^\alpha = l_s^{-4} \int_X \hat{D}_i \wedge \tilde{\omega}^\alpha \quad , \alpha = 1, .., h^{1,1} \,, \tag{9}$$

for a basis of 2-forms ω_α and a basis of 4-forms $\tilde{\omega}^\alpha$ on X with $l_s^{-4} \int \omega_\beta \wedge \tilde{\omega}^\alpha = \delta^\alpha_\beta$. The gauge coupling is given by the volume of the 4-cycle, i.e.

$$\frac{1}{g_i^2} \sim \frac{r^{i\alpha} t_\alpha}{\pi M_P} \,. \tag{10}$$

The couplings $q_{i\alpha} \neq 0$ give mass to the $U(1)$ gauge bosons A_i via the Stückelberg mechanism and are given as

$$q_{i\alpha} = l_s^{-2} \int_{D_i} \omega_\alpha \wedge \frac{\mathcal{F}}{2\pi} \,, \tag{11}$$

where $\mathcal{F}$ is the gauge flux on D_i. In the case $q_{i\alpha} \neq 0$, the axion τ_α is 'eaten' by the now massive $U(1)$ gauge boson and obtains a mass of order of the string scale

$$M_S = \frac{M_P}{\sqrt{4\pi \mathcal{V}}} \,. \tag{12}$$

Hence, it disappears from the 4D low energy effective field theory and cannot function as the QCD axion. The $\tau_\alpha \mathrm{tr}(F_i \wedge F_i)$ term in eq. (8) that gives τ_α axionic properties originates from dimensional reduction of a 8D Chern-Simons term of the D7 brane that is responsible for anomaly cancellation in the Green-Schwarz mechanism.

1.3 Axion decay constants in string theory

There seems to be an upper bound $f_a < M_P$ in effective theories from string theory due to moduli space dualities and new corrections that become important when one naively tries to send $f_a > M_P$, rendering again $f_a < M_P$. More precisely the authors of [9] e.g. found:

- For the weakly coupled heterotic string compactified on a six torus with radius R (16 supersymmetries $\Leftrightarrow \mathcal{N} = 4$ in 4D) with axion moduli from antisymmetric tensor field B_2

$$\frac{f_a^2}{M_P^2} = \frac{1}{(RM_S)^4} \,. \tag{13}$$

  So for small R it seems that one can realize trans-Planckian decay constants. However, T duality $R \leftrightarrow 1/R$ relates this region in moduli space to a physically equivalent region with large R where $f_a < M_P$.

- For theories in 4D with 32 supersymmetries (type II compactified on a 6-torus) one finds e.g. for the C_1 axion in type IIA

$$\frac{f_a^2}{M_P^2} = \frac{g^2}{(RM_S)^2} \,. \tag{14}$$

  It was shown [10] that the space of all compactifications is a connected space and every direction can be mapped by duality transformations into type II with weak coupling and radii larger than the string scale and hence $f_a < M_P$.

- For theories with less than 16 supersymmetries it is not possible to make general statements regarding $f_a < M_P$. But the non-successful search for $f_a > M_P$ in [9] might be considered indicative of the general situation.
  To circumvent the simplest equivalence $R \leftrightarrow 1/R$ that forbids $f_a > M_P$ one could consider for instance type IIB near a conifold point in moduli space, which is a singular region in the geometry which is topologically $S^2 \times S^3$. For the B_2 axion one finds

$$\frac{f_a^2}{M_P^2} = \frac{1}{(V_{S^2} M_S^2)^2} \,. \tag{15}$$

  One could hope to obtain $f_a > M_P$ for $V_{S^2} \to 0$ since there is no simple dual description of the theory for small V_{S^2}. However, worldsheet instantons violate the axionic shift symmetry and become sizable for $V_{S^2} \to 0$, effectively rendering $f_a > M_P$. Also often there appear additional light states in the effective theory as $V_{S^2} \to 0$ that couple to the respective axion and again yield $f_a < M_P$ in the end.

Accomplishing $f_a > M_P$ would be interesting from the point of view of axion inflation as we will see in Robert Richter's talk.

2 Axions and moduli stabilization in type IIB flux compactifications

In type IIB flux compactifications, the axio-dilaton S and the $h^{2,1}$ complex structure moduli U_i are fixed by fluxes at tree-level.[4] The Kähler moduli are massless at tree-level due to the no-scale structure of the effective potential but can be stabilized by taking into account perturbative α' and g_s corrections to the Kähler potential and non-perturbative corrections to the superpotential. These non-perturbative corrections break the shift symmetry of the axions. In the following, we discuss the axion masses in the moduli stabilization scenarios of KKLT [11] and the LARGE Volume scenario (LVS) [12].

If one is looking for an axion to solve the strong CP problem, i.e. the QCD axion, there is the following obstruction: There are many stringy effects that can generate a potential for the axions, e.g. worldsheet instantons and gaugino condensation, that can dominate over the in terms of Planck scale physics very weak QCD instantons. If this is the case the axions cannot solve the strong CP problem because they will not lead to $\langle \Theta \rangle = 0$. Hence, we will be looking for at least one axion, whose mass is zero up to QCD instantons after moduli stabilization.

2.1 KKLT

Consider non-perturbative corrections

$$W_{np} = \sum_{i=1}^{h^{1,1}} A_i(S, U_j) e^{-a_i T_i} \,, \tag{16}$$

with $a_i = 6\pi/b_0$ with b_0 the coefficient of the one-loop beta function and A_i considered to be $\mathcal{O}(1)$ constants after tree level stabilization of S and U_j. This can lead to stable AdS vacua that are supersymmetric, i.e.

$$D_{T_i} W = D_S W = D_{U_j} W = 0 \,. \tag{17}$$

As was realized in [13], this always leaves all saxions t_i and axions τ_i with a mass of the order of

$$m_{t_i} \sim m_{\tau_i} \sim a_i W_0 M_P / \mathcal{V} \,, \tag{18}$$

[4] See Francesco Pedro's talk.

such that there remains no candidate for a QCD axion or any ALP.

However, in a variation of the KKLT scenario [14] a single non-perturbative effect on an ample divisor

$$D_{am} = \sum_{i=1}^{h^{1,1}} \lambda_i D_i \,, \tag{19}$$

with $\lambda_i > 0 \; \forall i = 1, .., h^{1,1}$, i.e.

$$W_{np} = A e^{-aT_{am}} = A e^{-a \sum_{i=1}^{h^{1,1}} \lambda_i T_i} \,, \tag{20}$$

can lead to a stable supersymmetric AdS vacuum with one massive axion $\tau_{am} = \mathrm{Im}(T_{am})$ and $h^{1,1} - 1$ at leading order massless axions. These receive masses by very small higher-order non-perturbative corrections

$$W_{np} = \sum_{i=1}^{h^{1,1}-1} A_i e^{n_i a_i \tilde{T}_i} \,, \tag{21}$$

where $\tilde{T}_i$ is a combination of moduli orthogonal to T_{am}, leaving $h^{1,1} - 1$ axions with a logarithmically hierarchic mass spectrum.

In both cases, the volume scales as $\mathcal{V} \sim -\ln W_0$, such that large volumes are difficult to achieve and the string scale is typically of the order of the Planck or GUT scale which generically leaves one with axion decay constants f_a in conflict with the (standard) cosmological upper bounds.

2.2 LVS

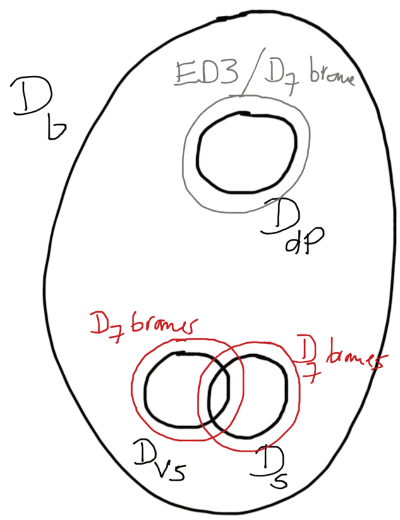

Figure 2: Minimal setup for a Calabi-Yau with $h^{1,1} = 4$ in the LVS with an MSSM / GUT like sector. Figure taken from [1]

In the LVS exponentially large volumes can be realized for generic $\mathcal{O}(1)$ values of the flux superpotential W_0. For the purpose of moduli stabilization one needs at least two Kähler moduli: one small cycle T_{dP} and one big cycle T_b whose volume dominates the overall 'swiss-cheese' volume form

$$\mathcal{V} = \gamma_b t_b^{3/2} - \sum_{i=1}^{h^{1,1}-1} \gamma_i t_i^{3/2} , \tag{22}$$

with triple self-intersections γ_b and γ_i and $t_{dP} \in \{t_i\}$. T_{dP} carries the non-perturbative effect

$$W_{np} = A e^{-aT_{dP}} , \tag{23}$$

which stabilizes the overall volume and T_{dP}.

If one also demands an MSSM or GUT like sector on a stack of D7 branes wrapping a 4-cycle, one has to add 2 more Kähler moduli corresponding to small cycles [15]: a 4-cycle corresponding to T_{vs} which supports the MSSM / GUT like sector and whose gauge coupling is determined by its volume t_{vs} and a 4-cycle which intersects this 4-cycle stabilizing $t_{vs} > 0$ by D-terms that originate from gauge flux on this cycle. T_{dP} corresponds to a del Pezzo divisor, i.e. the divisor is rigid which guarantees $A \neq 0$ and it does not intersect any other divisor. The latter secures that there are no chiral modes at the intersection with the brane stack of the visible sector that force $A = 0$.

Regarding the stabilization of the axions one finds [2]:

- The del Pezzo axion τ_{dP} becomes heavy because its shift symmetry is broken by $Ae^{-aT_{dP}}$.
- The non-zero gauge flux on the N D7 stack on T_s induces the breaking $U(N) \to SU(N) \times U(1)$ with an anomalous $U(1)$. The anomaly is canceled by the Green-Schwarz mechanism which induces a Stückelberg mass term for the $U(1)$ gauge boson by 'eating' the axion τ_s. This generically induces a high mass for τ_s. Note that this 'eating' of axions happens for every cycle where one switches on non-vanishing gauge flux to stabilize the cycle modulus by D-term constraints.
- The axions τ_b and τ_{vs} remain light because we have not switched on neither gauge flux nor are there large non-perturbative effects on these cycles. The perturbative effects that are used to stabilize the saxions do not break the axionic shift symmetry and hence do not imply a potential for the axions. Tiny higher oder corrections W_{np} and K_{np} in the end give a mass to these axions.

Hence there are at least two light axions in the LVS with a MSSM / GUT like sector. For $h^{1,1} > 4$ there will be an even larger number of ALP's corresponding to axions of 4-cycles that are fixed by perturbative effects.

For typical values $g_s \sim 0.1$, $W_0 \sim 1$ and $\mathcal{V} = 10^{14}$ the main scales in the model are

$$\begin{aligned}
M_s &= \frac{M_P}{\sqrt{4\pi\mathcal{V}}} \sim 10^{10}\,\mathrm{GeV} , \\
M_{KK} &\sim \frac{M_P}{\mathcal{V}^{2/3}} \sim 10^{9}\,\mathrm{GeV} , \\
m_{t_s} &\sim \frac{M_P}{\mathcal{V}^{1/2}} \sim 10^{10}\,\mathrm{GeV} , \\
m_{t_{dP}} &\sim m_{\tau_{dP}} \sim \sqrt{g_s} W_0 \frac{M_P}{\mathcal{V}} \ln \mathcal{V} \sim 30\,\mathrm{TeV} , \\
m_{\mathrm{soft}} &\sim m_{3/2} \sim \sqrt{g_s} W_0 \frac{M_P}{\mathcal{V}} \sim 1\,\mathrm{TeV} , \\
m_{t_{vs}} &\sim \alpha_{vs} m_{3/2} \sim 40\,\mathrm{GeV} , \\
m_{t_b} &\sim \frac{M_P}{\mathcal{V}^{3/2}} \sim 0.1\,\mathrm{MeV} .
\end{aligned} \tag{24}$$

The light modulus m_{t_b} might cause problems in terms of the CMP.

In Kähler uplifting scenarios the possible emergence of an axiverse has not been studied.

Axion decay constants and couplings to the visible sector

Generically, the axion decay constants f_a is of the order of the string scale in many models in string theory [9, 16]. For models with GUT-like phenomenology this is often too large to be in agreement with standard cosmological bounds.

To calculate the axion decay constant in our above LVS example, we need to transform the two axion fields τ_{vs} and τ_b to a basis a_{vs} and a_b where the fields are canonically normalized. One finds

$$\begin{aligned} \frac{\tau_b}{t_b} &\simeq \mathcal{O}(1)a_b + \mathcal{O}(t_{vs}^{3/4}\mathcal{V}^{-1/2})a_{vs}\,, \\ \frac{\tau_{vs}}{t_{vs}} &\simeq \mathcal{O}(1)a_b + \mathcal{O}(t_{vs}^{-3/4}\mathcal{V}^{1/2})a_{vs}\,. \end{aligned} \tag{25}$$

The decay constants and couplings to gauge bosons and matter fields be obtained by comparing eq. (8) and eq. (3) to be

$$\begin{gathered} f_{a_b} \simeq \frac{M_{KK}}{4\pi}\,, \qquad f_{a_{vs}} \simeq \frac{M_S}{\sqrt{4\pi}t_{vs}^{1/4}}\,, \\ C_{bb} \simeq \mathcal{O}(1)\,, \quad C_{vs\,b} \simeq \mathcal{O}(\mathcal{V}^{-1/3})\,, \quad C_{b\,vs} \simeq \mathcal{O}(\mathcal{V}^{-2/3})\,, \quad C_{vs\,vs} \simeq \mathcal{O}(1)\,. \end{gathered} \tag{26}$$

For an intermediate string scale $M_S \sim 10^{9-12}$ GeV we are in the phenomenologically interesting window

$$f_a \equiv \frac{f_{a_{vs}}}{C_{vs\,vs}} \sim M_S \sim 10^{9-12}\,\text{GeV}\,. \tag{27}$$

The decay constant of a_b is smaller $f_{a_b} \sim M_{KK}$ but its coupling to the standard model gauge bosons can be completely neglected since $C_{b\,vs} \simeq \mathcal{O}(\mathcal{V}^{-2/3})$.

a_{vs} is an excellent candidate for the QCD axion since it does not obtain a mass by stronger effects than QCD instantons and it actually couples to QCD via

$$\mathcal{L} \supset \frac{g_3^2}{32\pi^2}\left(\Theta - \frac{a_{vs}C_{vs\,vs}}{f_{a_{vs}}}\right)\text{tr}(G_3 \wedge G_3)\,. \tag{28}$$

Note that if we would not have found $C_{vs\,vs} \simeq \mathcal{O}(1)$ but some very small value $C_{vs\,vs} \ll 1$, a vacuum angle $\Theta > 2\pi C_{vs\,vs}$ could not have been absorbed into the VEV of a_{vs} since a_{vs} only varies through $[0, 2\pi f_a\}$, i.e. the axionic potential is $2\pi f_a$ periodic.

One can obtain more ALP's with

$$f_{ALP_i} \sim f_{a_{vs}} \sim M_S\,, \tag{29}$$

and

$$C_{ALP_i vs} \sim C_{vs\,vs} \simeq \mathcal{O}(1)\,, \tag{30}$$

if there are more small cycles that intersect with the visible sector cycle T_{vs} and there is no gauge flux on these cycles such that the axions are not eaten by anomalous $U(1)$'s. These ALPs are expected to be even lighter than the QCD axion with a logarithmic mass hierarchy

$$m_{ALP_i} \sim e^{-n\pi t_{ALP_i}} \times \begin{cases} M_P\,, & \text{for } \delta W_{np} \text{ terms or QCD corrections}\,, \\ m_{3/2} & \text{for } \delta K_{np} \text{ terms}\,. \end{cases} \tag{31}$$

In [2], globally consistent orientifolded Calabi-Yau models with semi-realistic D7-brane and gauge flux setup were found with the above features explicitly realized.

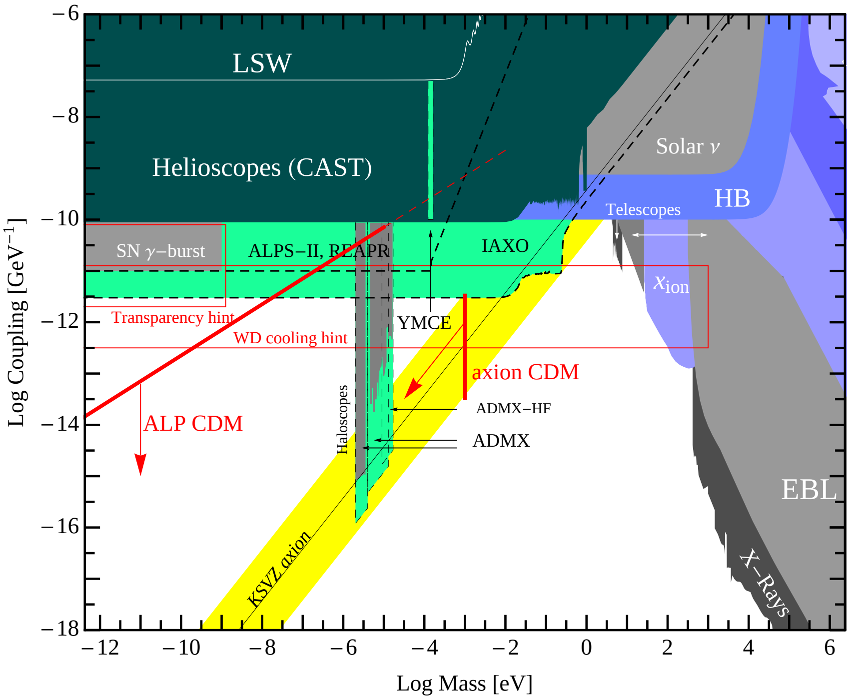

Figure 3: Axion and ALP coupling to photons $g_{i\gamma} \equiv \alpha C_{i\gamma}/(2\pi f_{a_i})$ vs. its mass. The yellow band is the generic prediction for the QCD axion. Figure taken from [1]

3 Prospects to discover the axiverse

See [6, 1, 17, 18, 19].

- *Haloscope searches:* If the QCD axion contributes significantly to the cold dark matter density, i.e. $f_a \gtrsim 10^{11-12}$ GeV, one can try to exploit the axion-photon coupling by searching for the signal of dark matter axion to photon conversions in a narrow bandwidth microwave cavity sitting in a strong magnetic field.

- *Helioscope searches:* If the QCD axion decay constant is at the lower end of the phenomenological window, i.e. $f_a \gtrsim 10^{9-10}$ GeV, one can try to detect solar axions by their conversion into photons inside of a strong magnet pointed towards the sun. Since helioscope searches do not loose sensitivity towards small masses, one can also use these experiments to look for ALPs with approximately the same coupling to photons as the QCD axion.
  There are astrophysical hints such as the anomalous transparency of the universe for TeV photons and the anomalous cooling of white dwarfs which might be explained by an ALP with
  $$\frac{f_{a_i}}{C_{i\,e}} \sim 10^9\,\mathrm{GeV}\,, \quad \frac{f_{a_i}}{C_{i\,\gamma}} \sim 10^8\,\mathrm{GeV}\,, \quad m_{a_i} \lesssim 10^{-9} - 10^{-10}\,\mathrm{eV}\,, \tag{32}$$
  which could be tested by helioscopes.

- *Light-shining-through-walls searches:* Laser photons are send along a strong magnetic field which allows them to convert to ALPs. These transfer through a wall that is opaque for photons and a strong magnetic field behind the wall allows them to convert into photons again that can be detected.

- *Rotation of the CMB polarization:* ALPs with mass $10^{-33} - 10^{-28}$ eV cause a rotation in the polarization of the CMB if they start oscillating around their VEV anytime between recombination and today, violating parity due to $\langle a \rangle \neq 0$ at recombination. Planck and / or other CMB experiments could test this scenario in the future.
- *Steps in the power spectrum of small scale density perturbations:* If axions make up a significant part of the dark matter density, they can suppress power in small scale (< 1 Mpc) density perturbations. This is due to a quantum pressure originating from the uncertainty principle forbidding gravitational collapse below a scale proportional to $1/\sqrt{m_{a_i}}$. This might be observable for masses up to 10^{-18} eV e.g. by 21 cm line observations. For different axions contributing to the dark matter density, there could be multiple steps in the power spectrum.
- *Black holes:* For masses $10^{-22} - 10^{-10}$ eV, ALP's can affect the dynamics of black holes by forming gravitational bound states with a rotating black hole. This bound state can be deexcited through graviton emission that carries away the black hole's angular momentum. For heavy black holes, if the angular momentum of the black hole is supported by accretion one might be able to detect these gravity wave pulsars by future gravity wave experiments. If the black hole is light and the axion heavy, i.e. the QCD axion, the loss of angular momentum leads to a spindown of the black hole, resulting in gaps in the mass spectrum of rapidly rotating black holes.
- *Dark Radiation:* After inflation the energy of the universe has to be transferred predominantly into thermal relativistic SM degrees of freedom (reheating). Constraints on this are measured by N_{eff}, the effective number of neutrino species which is $N_{eff,BBN} = 3$ and $N_{eff,CMB} = 3.04$ in the SM. Observations show a mild but consistent preference for $N_{eff} - N_{eff,SM} > 0$ from BBN and(?) CMB measurements. [18, 19] show that dark radiation from ALPs is unavoidable in sequestered LVS scenarios which strongly constrains the couplings of the ALPs.

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