Cosmological implications of a light dilaton

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Abstract: Supersymmetric Peccei–Quinn symmetry and string theory predict a complex scalar field comprising a dilaton and an axion. These fields are massless at high energies, but it is known since long that the axion is stabilized in an instanton dominated vacuum. Instantons and axions together also provide a mechanism to stabilize a dilaton, thus accounting for a dilaton as a possible cold dark matter component accompanying the axion. We briefly review the prospects of this scenario and point out further implications.

Based on a talk at COSMO–97, First International Workshop on Particle Physics and the Early Universe, Ambleside (England) 15–19 September 1997.

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Cosmology and physics of the early universe play an increasing role in physics, nurtured by seminal experimental and theoretical interactions with particle physics and astrophysics. A particularly exciting field of interaction between particle physics, astrophysics and cosmology concerns the nature of the dark matter in the universe: While there is consensus that an appreciable fraction has to consist of non-baryonic components, it is unclear yet whether we are talking about LSPs or axions eventually coming with a dilaton or something else, and what the contribution from massive neutrinos is. It is also unclear, whether cold dark matter comes as WIMPs or through coherent oscillations of very light (pseudo-)scalars. A light axion/dilaton pair motivated from supersymmetric Peccei–Quinn symmetry or string theory is an interesting candidate for the latter scenario.

The axion [1, 2] is widely recognized as a leading competitor for the role of cold dark matter in the universe [3, 4], and intense searches are under way [5]. However, in models which originate from a supersymmetric theory or realize a duality symmetry on an abelian subgroup, the axion $a$ is accompanied by a dilaton $\phi$, which would also contribute an appreciable amount to the energy density of the universe. The relevant couplings before taking into account non-perturbative and curvature effects are contained in a Lagrangian

$$\mathcal{L} = -\frac{1}{2} \partial_\mu \phi \cdot \partial^\mu \phi - \frac{1}{2} \exp\left(-\frac{2\phi}{f_\phi}\right) \partial_\mu a \cdot \partial^\mu a - \frac{1}{4} \exp\left(\frac{\phi}{f_\phi}\right) F_{\mu\nu}^j F^\mu\nu_j + \frac{a}{4 f_\phi} \tilde{F}_{\mu\nu}^j F^{\mu\nu}_j,$$

where we specialized to the supersymmetric/S–dual case [6] in fixing the axion scale $f_a = (\alpha_s/2\pi)f_\phi$.

The observation that a CDM axion should come with a dilaton is already implicit in [4], since it was pointed out that a dilaton with a large coupling scale $f_\phi$ should also meet cosmological constraints from $\Omega \leq 1$. It was shown in [8] that dilatons with $f_\phi$ of the order of the Planck mass should acquire a mass of the order of the gravitino mass if supersymmetry is softly broken at low energies. Another mass generating mechanism was proposed in [9], where it was pointed out that instantons and axions together may generate a dilaton mass due to the specific couplings of the dilaton to the kinetic terms of the axion and the gluons. This implies a scenario where a dilaton is stabilized at a very late stage together with an axion, whence both particles would arise through coherent oscillations at the chiral phase transition. On the other hand, dilatons which are stabilized at a high temperature generically acquire higher masses.
and may appear as WIMPs rather than through coherent oscillations. In this sense dilatons were mentioned as dark matter already in \cite{10, 11}, and mechanisms for dilaton stabilization at or above the SUSY scale are discussed in \cite{12, 8, 13, 14}.

In the present contribution we are mainly concerned with the prospects of a dilaton as a CDM component, neglecting interesting implications of possible dilaton–curvature couplings, for which we refer to \cite{10, 13, 15} and references there. However, before focusing on axions and dilatons let us make a general remark on cold dark matter from particle physics beyond the standard model:

The characteristic features of cold dark matter are very weak coupling and negligible pressure, whence the energy $\rho_\phi$ in the CDM degrees of freedom decays with the third power of the scale factor $R^{-3}$ and catches up with the radiation dominated heat reservoir according to $\rho_\phi/\rho_\gamma \sim R \sim \sqrt{t}$. Due to its weak coupling CDM is also decoupled from the heat bath, and this has simple but interesting implications: Both WIMPs and light axions and dilatons begin to behave like dust when mass terms are created during a phase transition or become comparable to the temperature. However, before the onset of mass terms or for temperatures well above the mass scale of a WIMP, CDM components can be considered as factually or effectively massless, and the question arises about evolution of their energy density during this early phase: If the CDM degrees of freedom, in spite of their weak coupling, have a sufficiently strong internal coupling for internal thermalization, then their energy decays in parallel to the heat bath $\rho_\phi \sim R^{-4}$ and $\rho_\phi/\rho_\gamma$ remains constant during the pre-CDM phase. However, if all the couplings of the CDM components are very weak, then internal thermalization cannot be assumed, and before onset of mass terms CDM degrees of freedom behave like a stiff fluid $p_\phi = \rho_\phi$, resulting in a sharp decline of their energy density $\rho_\phi/\rho_\gamma \sim R^{-2}$ above the temperature $T_c$ where mass terms become relevant (see Fig. 1).

Superficially, we refer to the degrees of freedom in the stiff fluid as moduli. In Fig. 1 $T_s$ labels a scale where both the moduli and the radiation originate (this might be a string scale, but this we do not know). The scale $T_c$ indicates, that a transition from the stiff fluid behaviour to CDM–like behaviour emerges upon creation or increasing importance of mass terms, and $T_r$ is the scale of photon recombination, when most energy in the heat reservoir is converted into pressureless matter. Here, all temperatures refer to the radiation dominated heat bath, whence $T \sim R^{-1}$. 
The scale $T_c$ is readily calculated in terms of $T_s, T_r$ and the ratios $\eta = \rho_\gamma(T_s)/\rho_\phi(T_s)$ and $\xi = \rho_\phi(T_r)/\rho_\gamma(T_r)$:

$$T_c = \left(\eta \xi T_s^2 \right)^{1/3}. \quad (1)$$

The corresponding time scales can be inferred from the Friedmann equation under the proviso of negligible coupling between the radiation and the moduli: For $T > T_c$ the scale factor evolves according to

$$\frac{2}{\sqrt{3}m_{Pl}} \rho_\gamma(t - t_0) = x \sqrt{x^2 \rho_\gamma + \rho_\phi} - \sqrt{\rho_\gamma + \rho_\phi}$$

$$- \frac{\rho_\phi}{\sqrt{\rho_\gamma}} \ln \left( \frac{x \sqrt{\rho_\gamma} + \sqrt{x^2 \rho_\gamma + \rho_\phi}}{\sqrt{\rho_\gamma} + \sqrt{\rho_\gamma + \rho_\phi}} \right),$$

while for lower temperatures the evolution follows

$$\frac{\sqrt{3}}{2m_{Pl}} \rho_\phi^2(t - t_0) = (x \rho_\phi - 2 \rho_\gamma) \sqrt{x \rho_\phi + \rho_\gamma} - (\rho_\phi - 2 \rho_\gamma) \sqrt{\rho_\phi + \rho_\gamma}.$$

Here $x = R/R_0$, $m_{Pl} = 2.4 \times 10^{18}$ GeV, and the subscript 0 in each equation indicates an arbitrary fiducial time during the respective epoch.
Consider the heterotic string [16], e.g.: This predicts an effective number of massless states $g_* = 7560$ below the string scale, among which 120 belong to the very weakly coupled gravitational sector. Assuming thermalization of the other degrees of freedom and equipartition at the string scale gives an estimate $\eta = 62$. Furthermore, it was pointed out in [17] that the heterotic string scale should be as low as $T_s \simeq 4 \times 10^{16}$ GeV, because for an earlier transition the product of comoving time and string scale would spoil the uncertainty relation. Inserting $T_r \simeq 0.3$ eV and $\xi \simeq 10$ in (1) then gives $T_c \sim 10^9$ GeV, while the same calculation with $T_s = m_{Pl}$ yields $T_c \sim 10^{10}$ GeV. These values comply nicely with bounds on an axion scale (see e.g. [18]), and with Pati’s proposal of a grand fiesta of new physics around $10^{11}$ GeV [19].

On the other hand, if the moduli would be thermalized already for temperatures above $T_c$, eq. (1) would be replaced by $T_c = \eta \xi T_r$, and large values of $T_c$ would require a huge ratio $\eta$.

Eq. (1) generically indicates a high temperature phase transition if we insist on continuous evolution of the energy density in the moduli. Conversely, it means that the energy in an axion and a dilaton stabilized through instanton effects at a scale $T \sim 1$ GeV arises from conversion of a tiny fraction $\sim 10^{-8} \rho_\gamma$ of radiation during the chiral phase transition.

The instanton–induced potential [9]

$$V(a, \phi) = \frac{1}{6} m_\phi f_\phi^2 \left( 2 \exp \left( \frac{\phi}{f_\phi} \right) + \exp \left( -2 \frac{\phi}{f_\phi} \right) \right) + m_a f_a^2 \left( 1 - \cos \left( \frac{a}{f_a} \right) \right) \quad (2)$$

yields parameters $m_\phi f_\phi \simeq m_a f_a \simeq m_\pi f_\pi$. However, a dilaton mass raises the same issues as an axion mass: A large–$f_\phi$ dilaton in an open or flat universe would prevent the chiral phase transition in the same way as an axion [3], and the well-known upper bound on decay constants arising from chiral symmetry breaking and $\Omega \leq 1$ is $f \leq 10^{12}$ GeV. Possibilities to relax this bound include large entropy production [20] or a subsequent phase of accelerated expansion [21]. Recently, Banks and Dine also proposed two very different scenarios, supposing either moduli domination until a reheating temperature between the QCD scale and the scale of nucleosynthesis is reached, or an axion driven to $\sqrt{\langle a^2 \rangle} \leq 10^{-4} f_a$ by an effective potential during inflation [14]. However, overall the bound $f \leq 10^{12}$ GeV is very robust, and even the assumption of strong coupling for temperatures well above the phase transition turned out not to ameliorate the problem [22].
Here we would like to briefly comment on another proposal: Above the phase transition both the axion and the dilaton are massless and one should expect strong fluctuations in both fields. Therefore, one might expect chiral symmetry breaking also to occur in presence of a large–$f$ axion, but not everywhere at the same time, i.e. the transition hypersurfaces would not be flat (for $k = 0$) or hyperbolic (for $k = -1$). As a consequence, a large–$f$ axidilaton would induce an extra time scale and generate extra perturbations during the transition. Of course, there still would exist an upper bound, since very large $f$ would imply too long a time scale for the transition and generate too large perturbations at the QCD scale. Stated differently, the average displacement of the axidilaton from the vacua of the potential (2) prevents emergence of the chiral symmetry breaking low temperature phase, since it costs too much energy to create a mass term. But the scenario just mentioned presumes, that this would not apply to those regions where the axidilaton field is close to the minimum of (2). In those regions the instanton gas could dominate the partition function and induce both a quark condensate and axidilaton mass terms, thus confining the axion and the dilaton in that region to absolute values $\ll f$. Outside of those instanton dominated bubbles the axidilaton continues to fluctuate (almost) freely, until it comes close enough to zero, whence it also gets trapped through the onset of the phase transition.

The caveat with this proposal concerns the emergence of fluctuations required to relax the bound on $f$: The quantum fluctuations of the axidilaton in the radiation dominated background are only vacuum fluctuations, which are neglected e.g. in calculating the seeds for structure formation in de Sitter space \cite{23}, and would also be subtracted if one calculates the two-point function according to Bunch and Davies or Vilenkin and Ford \cite{24}. Therefore, any appreciable effect should result from thermal fluctuations, and this touches upon the problem of decoherence in subluminally expanding spacetimes. We have nothing to say about the latter problem at this stage, and therefore we conclude with the assertion that one can think of plausible scenarios for accomodating chiral symmetry breaking in $\Omega \leq 1$ universes with large–$f$ axions and dilatons, but overall the bound $f \leq 10^{12}$ GeV seems very robust.

**Acknowledgement:** R.D. would like to thank the organizers of COSMO–97 for hospitality during a very interesting and stimulating meeting. Support by the DFG through SFB 375–95 is gratefully acknowledged.
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