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Scalar field dark matter with spontaneous symmetry breaking and the 3.5 keV line

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ABSTRACT

We show that the present dark matter abundance can be accounted for by an oscillating scalar field that acquires both mass and a non-zero expectation value from interactions with the Higgs field. The dark matter scalar field can be sufficiently heavy during inflation, due to a non-minimal coupling to gravity, so as to avoid the generation of large isocurvature modes in the CMB anisotropies spectrum. The field begins oscillating after reheating, behaving as radiation until the electroweak phase transition and afterwards as non-relativistic matter. The scalar field becomes unstable, although sufficiently long-lived to account for dark matter, due to mass mixing with the Higgs boson, decaying mainly into photon pairs for masses below the MeV scale. In particular, for a mass of ~ 7 keV, which is effectively the only free parameter, the model predicts a dark matter lifetime compatible with the recent galactic and extragalactic observations of a 3.5 keV X-ray line.

One of the most important open problems in modern cosmology is the nature of dark matter (DM), an invisible form of matter that can explain the observed structure of the Universe on large scales, the galaxy rotation curves and the anisotropies in the Cosmic Microwave Background (CMB). However, despite the large number of candidates, there are still no definite answers concerning its origin [1]. An interesting possibility is an interaction between DM and the Higgs field, widely known as "Higgs-portal DM". This has been extensively studied in the literature, namely in the context of thermal production [2–14]. However, the lack of evidence for WIMP-like particles in the various ongoing experiments [15] suggests looking for alternative candidates, such as oscillating scalar fields, as considered e.g. in Refs. [16–18].

In this Letter, we show for the first time that a scalar field dark matter coupled to the Higgs field can naturally explain the 3.5 keV X-ray line detected by the XMM-Newton observatory. Our model considers a complex scalar field, Φ, interacting with the Higgs doublet, $\mathcal{H}$, only through scale-invariant interactions given by the Lagrangian density:

$$\mathcal{L}_{int} = \pm\, g^2\, |\Phi|^2\, |\mathcal{H}|^2 + \lambda_\phi\, |\Phi|^4 + V\,(\mathcal{H}) + \xi\, R\, |\Phi|^2\ , \tag{1}$$

where the Higgs potential $V(\mathcal{H})$ has the standard "mexican hat" shape. We assume that the scale invariance of the Φ interactions is a consequence of an underlying scale invariance of the full theory, that is spontaneously broken in the Higgs and gravitational sectors by some mechanism that has no influence on the effective dynamics of the dark matter scalar field (see also Ref. [19]). This allows for the Higgs-dark scalar interaction with coupling, g, the dark scalar field quartic self-interactions with coupling, λ_ϕ, and for a non-minimal coupling, ξ, of the DM to the Ricci scalar, R.

The interaction Lagrangian (1) also exhibits a U(1) symmetry and we may consider two cases. When the Higgs-dark scalar interaction has a positive sign, the U(1) symmetry remains unbroken and the DM field is stable. For a negative coupling, the U(1) symmetry can be spontaneously broken in the vacuum and the DM field may decay, allowing for astrophysical signatures, as we will see below. In this Letter, we focus on the latter case, leaving the discussion of the former to a longer companion paper.

The background dynamics of the homogeneous dark scalar field mode is determined by the equation of motion:

$$\ddot{\phi} + 3H\dot{\phi} + V'\,(\phi) + 2\xi\, R\phi = 0\ , \tag{2}$$

where $\Phi = \phi/\sqrt{2}$ since the complex phase has a trivial dynamics. From the associated energy-momentum tensor, we obtain the

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effective energy density and pressure of the field, which are, respectively, given by

$$\rho_\phi = \frac{\dot\phi^2}{2} + V(\phi) + 12\xi H\phi\dot\phi + 6\xi\phi^2 H^2 , \tag{3}$$

$$p_\phi = \frac{1}{2}\,(1-8\xi)\,\dot\phi^2 - V(\phi) + 4\xi\phi V'(\phi) + 4\xi\phi\dot\phi H + \xi\phi^2\left[(8\xi-1)\,R + 2\,\frac{\ddot a}{a} + 4\,H^2\right] , \tag{4}$$

where a is the scale factor. We will see below that the introduction of a non-minimal coupling does not significantly change the usual form of ρ_ϕ and p_ϕ for an oscillating scalar field. As pointed out in Ref. [18], the initial conditions for the scalar field oscillations are set by the inflationary dynamics. In the parametric regime where $\xi \gg g, \lambda_\phi$, which will henceforth be the focus of our discussion, the field's mass during inflation is dominated by the non-minimal coupling to the curvature scalar, $R \simeq 12\,H_{inf}^2$, where $H_{inf} \simeq 2.5\times 10^{13}\,(r/0.01)^{1/2}$ GeV is the Hubble parameter during inflation and r is the tensor-to-scalar ratio. This yields $m_\phi \simeq \sqrt{12\xi}\,H_{inf} \gtrsim H_{inf}$ for $\xi \gtrsim 0.1$. As pointed out in Ref. [18], this super-Hubble mass prevents the field from acquiring large fluctuations during inflation that would give rise to observable isocurvature modes in the CMB spectrum, which are now significantly constrained [20]. For $m_\phi/H_{inf} > 3/2$, quantum fluctuations in the field get stretched and amplified during inflation, yielding a spectrum [21]:

$$|\delta\phi_k|^2 \simeq \left(\frac{H_{inf}}{2\pi}\right)^2 \left(\frac{H_{inf}}{m_\phi}\right) \frac{2\pi^2}{(a\,H_{inf})^3} . \tag{5}$$

Integrating over the comoving momentum k on super-horizon scales, we can obtain the field variance at the end of inflation, which sets the typical homogeneous field amplitude at the onset of oscillations in the post-inflationary era, ϕ_{inf}:

$$\phi_{inf} = \sqrt{\langle\phi^2\rangle} \simeq \alpha\,H_{inf} , \qquad \alpha \simeq 0.05\,\xi^{-1/4} . \tag{6}$$

Note that, during inflation, all terms in Eq. (3) are $\sim H_{inf}^4$ and therefore the dark scalar plays a negligible role in the inflationary dynamics.

We should briefly mention that, during the (p)reheating period, the Ricci scalar oscillates with the inflaton field, χ, since $R = (3p_\chi - \rho_\chi)/M_{Pl}^2 \sim m_\chi^2\chi^2/M_{Pl}^2$, inducing an effective bi-quadratic coupling between the dark scalar and the inflaton, $g_{\phi\chi}^2 \sim \xi m_\chi^2/M_{Pl}^2 \ll 1$. This interaction will lead to ϕ-particle production during reheating but, since $q_\phi = g_{\phi\chi}^2\chi^2/4m_\chi^2 \sim \xi\chi^2/M_{Pl}^2 \lesssim 1$ with $\chi \lesssim M_{Pl}$ during reheating, this should not be very efficient. In particular, it is natural to assume that the inflaton couples more strongly to other fields, which will thus be produced more efficiently and consequently reduce the amplitude of the inflaton's oscillations before any significant ϕ-particle production occurs. In addition, such particles remain relativistic until $T < m_\phi \ll T_{EW}$, and as we will see this implies that their density is much more diluted than the density of the homogeneous dark scalar condensate. We therefore expect ϕ-particle production during reheating to yield a negligible contribution to the present dark matter abundance.

After inflation and the reheating period, the Universe becomes dominated by radiation, and $R \simeq 0$. For temperatures above the electroweak scale, thermal effects keep the Higgs close to the origin (see e.g. [22]), such that the dark scalar field potential is dominated by the quartic term, $V(\phi) \simeq \lambda_\phi\phi^4/4$. Once the effective field mass $m_\phi = \sqrt{3\lambda_\phi}\phi$ exceeds the Hubble parameter in this era, the field starts oscillating about the origin with an amplitude that decays as $a^{-1} \propto T$.

It is easy to check that, in the oscillating phase, the last two terms in Eqs. (3) and (4) become subdominant since $m_\phi \gg H$. In addition, the remaining terms in Eq. (4) proportional to ξ cancel out upon averaging over the field oscillations, since $\langle\dot\phi^2\rangle = \langle\phi\,V'(\phi)\rangle$. This implies that the field's energy density and pressure are approximately given by the corresponding $\xi = 0$ expressions once it begins oscillating, such that $\rho_\phi \propto a^{-4}$ as long as the quartic potential term is dominant. During this period, the field thus behaves as *dark radiation*.

Equating the Hubble parameter in the radiation era with the effective field mass, we obtain for the cosmic temperature at the onset of field oscillations:

$$T_{rad} = \left(\sqrt{3\lambda_\phi}\,\phi_{inf}\,M_{Pl}\sqrt{\frac{90}{\pi^2 g_*}}\right)^{1/2} , \tag{7}$$

where g_* is the number of relativistic degrees of freedom. This is below the reheating temperature if the inflaton decays sufficiently fast after inflation, with $T_R \sim \sqrt{H_{inf}M_P}$ for instantaneous reheating.

Once the temperature drops below the electroweak scale, the Higgs field acquires a vacuum expectation value (vev) and the relevant Lagrangian density for the real ϕ and Higgs components is:

$$\mathcal{L}_{int} = -\frac{g^2}{4}\,\phi^2 h^2 + \frac{\lambda_\phi}{4}\,\phi^4 + \frac{\lambda_h}{4}\left(h^2 - \tilde v^2\right)^2 , \tag{8}$$

where $\lambda_h \simeq 0.13$ is the Higgs self-coupling. The Higgs and dark scalar vevs are, respectively:

$$h_0 = \left(1 - \frac{g^4}{4\lambda_\phi\lambda_h}\right)^{-1/2} \tilde v \equiv \mathrm{v}, \quad \phi_0 = \frac{g}{\sqrt{2\lambda_\phi}}\,\mathrm{v} , \tag{9}$$

where $\mathrm{v} = 246\,\mathrm{GeV}$. Note that a non-vanishing vev for ϕ implies $g^4 < 4\lambda_\phi\lambda_h$, which we assume to hold.

The interaction Lagrangian above is valid once the leading thermal contributions to the Higgs potential become Boltzmann-suppressed, which should occur below $T_{EW} \sim m_W$, where m_W is the W boson's mass. At this point, the field starts oscillating about ϕ_0 rather than about the origin. To determine the amplitude of oscillations at this stage, note that at T_{EW} the amplitude of field oscillations about the origin has been redshifted to:

$$\phi_{EW} \simeq \left(\frac{4\pi^2 g_*}{270}\right)^{1/4} \left(\frac{\phi_{inf}}{M_{Pl}}\right)^{1/2} \frac{T_{EW}}{\mathrm{v}}\,\frac{\lambda_\phi^{1/4}}{g}\,\phi_0$$
$$\simeq 10^{-4}\,g_*^{1/4}\,\xi^{-1/8}\left(\frac{T_{EW}}{m_W}\right)\left(\frac{r}{0.01}\right)^{1/4}\frac{\lambda_\phi^{1/4}}{g}\,\phi_0. \tag{10}$$

We thus see that $\phi_{EW} \lesssim \phi_0$ for $g \gtrsim 10^{-4}\lambda_\phi^{1/4}$ for $\xi \sim \mathcal{O}(1)$, with a larger non-minimal coupling to curvature localizing the field even closer to the origin at the electroweak phase transition (EWPT). This implies that, in these parametric regimes, the field will start oscillating about the non-zero vev below T_{EW}, with an amplitude $\phi_{DM} \equiv x_{DM}\,\phi_0$ with $x_{DM} \lesssim 1$ [23]. The field's equation of state then smoothly changes from a dark radiation to a cold dark matter behavior as the potential becomes quadratic about the minimum.

Therefore, the field amplitude evolves with the temperature as $\phi(T) = \phi_{DM}(T/T_{EW})^{3/2}$ and the number of particles per comoving volume becomes constant:

$$\frac{n_\phi}{s} = \frac{45}{4\pi^2 g_{*S}}\,\frac{m_\phi\phi_{DM}^2}{T_{EW}^3} , \tag{11}$$