Thermal Friction as a Solution to the Hubble Tension

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The phenomenological early dark energy (EDE) provides a promising solution to the Hubble tension in the form of an extra beyond-ΛCDM component that acts like a cosmological constant at early times ($z \geq 3000$) and then dilutes away as radiation or faster. We show that a rolling axion coupled to a non-Abelian gauge group, which we call the ‘dissipative axion’ (DA), mimics this phenomenological EDE at the background level and presents a particle-physics model solution to the Hubble tension, while also eliminating fine-tuning in the choice of scalar-field potential. We compare the DA model to the EDE fluid approximation at the background level and comment on their similarities and differences. We determine that CMB observables sensitive only to the background evolution of the Universe are expected to be similar in the two models, strengthening the case for exploring the perturbations of the DA as well as for this model to provide a viable solution to the Hubble tension.

I. INTRODUCTION

The flat ΛCDM model of cosmology has been tremendously successful at correctly describing numerous observables, including the complex cosmic microwave background (CMB) spectra \([1, 2]\). However, its predictions for servables, including the complex cosmic microwave background level, mimics the evolution of early dark energy. This component behaves like a cosmological constant at early times (\(z \geq 3000\)) and then dilutes away as fast or faster than radiation at some critical redshift \(z_c\), localising its influence on cosmology around \(z_c\). It can increase the pre-recombination expansion rate, decreasing the size \(r_s\) of the sound horizon. The CMB inference of \(H_0\) is based on \(r_s\) and its angular size \(\theta_s\) on the surface of last scatter. Precise observations of \(\theta_s\) combined with a ΛCDM-based deduction of \(r_s\) lead to a determination of \(H_0\) as \(\theta_s \sim r_s H_0\). Hence, a theory that predicts a smaller \(r_s\) also infers a greater \(H_0\) to preserve the precisely measured \(\theta_s\), alleviating the Hubble tension. Remarkably, the EDE solution improves the CMB \(\chi^2\) relative to ΛCDM. It was proposed as a phenomenological solution, the dynamics of which could emerge from various particle-physics models \([13–15]\).

In this letter, we present a dynamical particle-physics model that could solve the Hubble tension, which at the background level, mimics the evolution of early dark energy. This dissipative axion (DA) model is presented in Section II. Although we leave the details of the perturbations of this model to future work, in Section III, we argue why the background dynamics of this model are promising and indicate that the DA can form the required EDE and resolve the Hubble tension. Finally, we conclude in Section IV, and discuss the broader implications of this model as well as the way forward.

II. MODEL

We add a pure dark non-Abelian gauge group (SU(2)) and an axion \(\phi\) to the standard model particle content. The dark gauge bosons interact with \(\phi\) via a CP-odd coupling:

\[
\mathcal{L}_{\text{int}} = \frac{\alpha}{16\pi} \phi \tilde{A}_\mu A^{a}_{\mu} \tag{1}
\]

where \(A^{a}_{\mu} = \epsilon^{\mu\nu\alpha\beta} A^{\alpha\beta}_a\) is the field strength of the dark gauge bosons and \(\alpha = \frac{g^2}{4\pi}\), where \(g\) is the gauge coupling of the dark group. The dark sector is decoupled from the standard model. We give the axion, which is displaced from its minimum, a simple UV-potential \(1\):

\[
V(\phi) = \frac{1}{2} m^2 \phi^2 \tag{2}
\]

This potential intuitively illustrates the dynamics of our model, as the axion is essentially a harmonic oscillator. It mimics a cosmological constant while it slow-rolls down this potential. In the absence of any interaction, Hubble friction alone would govern the equation of motion of \(\phi\). However, the interaction term \(\mathcal{L}_{\text{int}}\) adds an additional

\[1\] The IR potential from the confining group is rapidly suppressed at temperatures above the confining scale and we have checked that its contribution is sub-dominant for our parameters.
friction $\Upsilon(T_{dr})$ to the equation of motion, dissipating energy through the production of dark radiation $\rho_{dr}$ which is comprised of dark gauge bosons, where $T_{dr}$ is the temperature of the dark radiation. In the small coupling limit ($\alpha \ll 1$) and $m \ll \alpha^2 T_{dr}$, this friction can be inferred from the sphaleron rate for a pure non-Abelian gauge group [18–20] and scales as:

$$\Upsilon(T_{dr}) = \kappa \alpha^3 \frac{T_{dr}^3}{f^2} \tag{3}$$

where $\kappa$ is an O(10) number with weak dependence on $\alpha$ and $f > T_{dr}$. The following equations of motion then describe the homogeneous evolution of the axion-radiation system:

$$\ddot{\phi} + (3H + \Upsilon(T_{dr})) \dot{\phi} + m^2 \phi = 0$$

$$\rho_{dr} + 4H \rho_{dr} = \Upsilon(T_{dr}) \dot{\phi}^2 \tag{4}$$

where $\rho_{dr} = \frac{\alpha^2}{30} g_* T_{dr}^4$ and $g_*$ denotes the relativistic degrees of freedom in the dark sector.

In the original EDE work, an oscillating scalar field subject only to Hubble friction had been proposed, whose energy must dilute like radiation or faster after the field becomes dynamical in order to satisfy precision cosmology constraints. This requirement places rigid demands on the scalar-field potential $V \propto \left(1 - \cos \frac{\phi}{T}\right)^n$ considered by [15] (or $V \propto \phi^{2n}$ as in [13]) with $n \geq 2$. These potentials do not easily emerge from a UV-complete theory without extreme fine-tuning, as quantum corrections naively should give rise to a dominant $\propto \phi^2$ term, which is absent in these potentials. Other proposed phenomenological EDE candidates [14] have similar fine-tuning issues.

In our DA model, the particle production friction $\Upsilon \gg m, 3H$, overdamps the motion of the scalar field. Thus, because the field is not oscillating, its dynamics are not sensitive to the potential $V(\phi)$. Instead, the friction $\Upsilon$ extracts energy from the scalar field into the dark radiation, which automatically dilutes away as $a^{-4}$.

We approximate the solution to the equation of motion Eq. (4) as

$$\phi(z) \approx \phi_0 e^{-\frac{4\alpha^2}{30} g_* T_{dr}^4 z^2}, \tag{5}$$

which is the solution of an overdamped oscillator where we approximated $t \simeq H(z)^{-1}$. Equation (5) illustrates that the DA begins to roll faster when $\frac{\Upsilon(z_d)}{m^4} \equiv H(z_d)$, where $z_d$ denotes the redshift at which the axion field becomes dynamical. At high redshifts ($z \gg z_d$) the axion is slowly rolling, building up to a steady-state temperature on time scales of order $\Upsilon^{-1}$ in the dark sector:

$$T_{dr}(z) \approx \left(\frac{m^4 f^2 \phi^2(z)}{2 \frac{4\alpha^2}{30} g_* \kappa \alpha^3 H(z)} \right)^{\frac{1}{2}} \tag{6}$$

by continuously extracting energy from the rolling field [21]. The generation of a steady-state temperature is independent of the presence of an initial dark temperature, as even starting with temperature fluctuations of the order of Hubble is sufficient to rapidly build up to the temperature in equation (6) [21].

As the field begins to roll faster, the temperature $T_{dr}$ in the dark sector rises steadily and the field continuously dumps its energy into the dark radiation bath. However, due to the weak dependence of the temperature on the background quantities, this change is $O(1)$. Therefore, approximating the friction $\Upsilon(z)$ as roughly constant does not change the qualitative behavior of our model at the background level, as we discuss in more detail in Section III. Eventually, as the axion energy is depleting, the source term $\Upsilon \dot{\phi}^2$ becomes smaller than $4H \rho_{dr}$, which leads to a decrease in temperature $T_{dr}$ until $\Upsilon \dot{\phi}^2$ becomes negligible and the dark radiation dilutes away as $a^{-4}$.

The main features discussed in this section are universal in the presence of a large friction coefficient ($\Upsilon \gg H(z)$) with temperature dependence $\Upsilon \propto T^p$ with $p < 4$, independent of the particular particle model that gives rise to the friction. On the micro-physical level however, it is difficult to construct a model that gives rise to a large friction, as couplings to light degrees of freedom generically also give rise to an even larger thermal back-reaction [22]. This back reaction is a finite contribution to the two-point function of the scalar field due to the thermal background, which modifies its effective mass. In our DA model, there is no large thermal back-reaction due to the shift symmetry of the axion $^2 [23–27]$. Hence, we are able to generate a large friction coefficient in the equation of motion of the axion without other unwanted side effects, making this an ideal candidate for converting a constant early dark energy into radiation. Moreover, the proposed model is minimal - it requires only an non-Abelian gauge group that couples to an axion in a well-understood manner, while placing no specific requirement on the shape of the axion’s potential. There exists another class of Abelian dissipative rolling axion-field models that convert vacuum energy to radiation through tachyonic instabilities [23–26, 28], which may be a promising alternative to the particle candidate presented in this paper.

### III. BACKGROUND DYNAMICS

Having laid the groundwork for the background evolution of the DA, we turn to its ability to form the required EDE and draw comparisons with the best-fit parameters of Ref. [12, henceforth labelled P18]. The particle setup described in Sec. II results in a rolling scalar field that behaves like a cosmological constant at early times plus a dark radiation component. The total contribution $\rho_{DA}$

\[2\] We have checked that the back-reaction due to softly breaking the shift-symmetry with the UV-potential is negligible.
fraction reaches a maximum at $\rho$ and as $\Omega$ fractional DA energy density peaks at redshift $z$ of the total energy density of the Universe, where $\rho$ to EDE as illustrated in Fig. 1. This constitutes a total $z < z_c$ which is only a function of the axion potential and its dilution component is subdominant and $\phi$ where $\rho$ the sum of the two, $\rho$ to an EDE-like component from the DA is then given by $z$ is roughly constant and the dark radiation component $\Omega$ in the scalar field (blue) is constantly and the dark radiation component $\Omega$ (yellow) is subdominant. At intermediate times ($z_{\text{peak}} < z < z_d$), the dark radiation $\Omega$ transitions to become dominant as $\Omega$ drops. Shortly after $T_{\text{dr}}$ reaches a maximum, the total fractional DA energy density peaks at redshift $z_{\text{peak}}$.

to an EDE-like component from the DA is then given by the sum of the two,

$$\rho_{\text{DA}}(z) = \rho_\phi(z) + \rho_{\text{dr}}(z), \quad (7)$$

where $\rho_\phi(z) \approx \frac{1}{2} m^2 \phi^2(z)^3$. At very early times, the radiation component is subdominant and $\phi$ is essentially frozen, acting like a cosmological constant with energy density

$$\rho_{\text{DA}}(z \gg z_d) \approx \frac{1}{2} m^2 \phi_0^2, \quad (8)$$

which is only a function of the axion potential and its initial conditions. Sometime after the axion thaws ($z < z_d$), the dark radiation becomes the dominant contributor to EDE as illustrated in Fig. 1. This constitutes a total fraction

$$f_{\text{DA}}(z) = \frac{\rho_{\text{DA}}(z)}{\rho_m(z) + \rho_r(z) + \rho_{\text{DA}}(z)} \quad (9)$$

of the total energy density of the Universe, where $\rho_m$ and $\rho_r$ denote the matter and radiation densities. This fraction reaches a maximum at $z_{\text{peak}}$. Relating this to the ‘critical redshift’ $z_c$ of the EDE as defined in P18, their best fit $z_c = 5345^4$ for the EDE that dilutes as radiation, which corresponds to $z_{\text{peak}} = 3322$. Roughly at this time, the source term $T\phi^2$ in Eq. (4) becomes negligible and the dark radiation dilutes away as $a^{-4}$ as shown in Fig. 1.

By approximating the friction $\Upsilon(z_{\text{peak}}) = \Upsilon_0$ as a constant, we illustrate how to estimate $z_{\text{peak}}$ analytically. In this limit, the approximation for the temperature of the dark radiation simplifies to

$$T_{\text{dr}}(z > z_{\text{peak}}) \simeq \left( \frac{m^2 \phi(z)}{2 \sqrt{\frac{\pi^2 g_*(H(z)\Upsilon_0)}}} \right)^{\frac{1}{2}}, \quad (10)$$

which together with Eqs. (5) and (7), allows us to approximate $f_{\text{DA}}$ as an analytical function in $z$:

$$f_{\text{DA}}(z \geq z_{\text{peak}}) \simeq \frac{1}{\rho_m(z) + \rho_r(z)} \left( 1 + \frac{m^2 \phi_0^2}{2 H(z)\Upsilon_0} \right). \quad (11)$$

Solving $\frac{df_{\text{DA}}}{dz}|_{z_{\text{peak}}} = 0$, and assuming that the peak lies close to matter-radiation equality, we can approximate $z_{\text{peak}}$ as

$$z_{\text{peak}} \simeq \left( \frac{1}{2 \sqrt{\frac{\pi^2 g_*(H_0)\Upsilon_0}}^2} \right)^{\frac{1}{2}} m^2 \phi_0^2, \quad (12)$$

where $\Omega_m$ is the fractional matter density today and $z_{\text{peak}}$ is now dependent only on $\frac{H_0}{m^2}$.

Equations (10)- (12) demonstrate how the physical observables depend exclusively on $\frac{H_0}{m^2}$, which sets the time scale at which the axion becomes dynamical, and $\frac{1}{2} m^2 \phi_0^2$ which scales the total amount of early dark energy. Therefore, at the background level, we effectively introduce only two new parameters beyond $\Lambda$CDM, but expect the perturbations to depend on more than just these two parameters. Including the full temperature dependence of the friction at the background level requires solving the coupled differential Eq. (5) numerically by specifying an initial condition $\Upsilon(z_1)$ at some $z_1$, increasing the effective number of background parameters to three. While this does not have a significant impact on the qualitative behavior of the DA system, it does change $\Upsilon(z_1)$, and $\frac{1}{2} m^2 \phi_0^2$ by $O(1)$ when keeping $z_{\text{peak}}$ and $f_{\text{DA}}(z_{\text{peak}})$ fixed.

For redshifts smaller than $z_{\text{peak}}$, the early dark energy is dominated by the radiation component which dilutes as:

$$f_{\text{DA}}(z < z_{\text{peak}}) \simeq \frac{1}{1 + \frac{z_{\text{peak}}}{z}} {\frac{z_{\text{peak}}}{1 + z_{\text{peak}}}}^4. \quad (13)$$

The fractional energy density $f_{\text{DA}}$ is then peaked at $z_{\text{peak}}$, as shown in Fig. 2. Our proposed model hence mimics

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4 The posteriors for EDE parameters in P18 are non-Gaussian. The best-fit parameters quoted here therefore do not correspond to their mean values and we hence do not include errors on the quoted parameters here.
the EDE proposed in P18 with \( n = 2 \), which resolves the Hubble tension.

The primary difference between the two models at the background level is a narrower peak for the DA (the effect being more pronounced for the constant friction approximation), as seen in Fig. 2. Based on this, we explore the expected differences between the background observables of the two models. In particular, we discuss the impact on CMB observables that capture the important features of the full CMB spectrum, but depend only on the background evolution of the Universe [11, 12, 29]. These are the size \( r_s \) of the sound horizon, the ratio \( r_{damp}/r_s \) of the damping scale to the sound horizon, the height of the first peak and the horizon size at matter-radiation equality.

As our model adds more radiation to the Universe, we naively expect the redshift of matter-radiation equality to shift. Quantifying this shift correctly requires a full MCMC to allow other cosmological parameters, in particular the physical density \( \omega_{cdm} \) of cold dark matter to compensate for some or all of the shift. We expect that the results of the MCMC will pull our posteriors in a direction that minimizes change to \( z_{eq} \). We hence leave further discussion of changes to \( z_{eq} \) for future work. We expect an increase in \( \omega_{cdm} \) to similarly compensate for a change to the height of the first CMB peak. Such an increase was observed by P18 for EDE - the best-fit \( \omega_{cdm} \) increases by \( \sim 9\% \) in the \( n = 2 \) EDE cosmology relative to \( \Lambda \)CDM. To compare, their maximum \( f_{EDE} \leq 7\% \). Moreover, the dark radiation peaks during matter-domination, further minimising the effect of adding dark radiation to the Universe. Consequently, in this letter, we limit our comparisons of the two models to investigating the effects of the sharper peak in \( f_{DA} \).

To do so, we first note that a slight narrowing of the peak of \( f_{DA} \) relative to \( f_{EDE} \) has minimal impact on the recombination redshift \( z_s \). This was verified using a modified version of the equation of state parameterization of the EDE of P18, similar to Ref. [14], sharpening the peak in \( f_{EDE} \) and calculating \( z_s \) with the CLASS [30, 31] cosmology code. As \( z_s \) is a background quantity, and \( f_{DA} \) is nearly identical to a narrower \( f_{EDE} \), we expect \( z_s \) for the DA to be similar to the EDE scenario. Then, the main change to \( r_s \) comes not from the limits of its integral, but the integrand, specifically, the expansion rate. Knowing how the expansion rate for the DA differs from EDE, we can calculate \( r_s \) by fixing the background cosmology to the best fit of the \( n = 2 \) EDE of P18, and the DA parameters such that the temperature dependent (independent) \( z_{peak} \) and \( f_{DA} \) (\( f_{EDE} \)) match the best-fit EDE (values specified in the caption of Fig. 2), giving

\[
\frac{1}{2} m^2 \varphi_0^2 = 0.55 \text{ eV}^4 \quad \text{and} \quad \frac{1}{2} m^2 \varphi_0^2 = 0.21 \text{ eV}^4 \quad \text{for the temperature dependent (independent) DA model.}
\]

We can see that the the approximations in Eqs. (11) and (13) qualitatively capture the behavior of \( f_{DA} \) as a function of \( \frac{1}{2} m^2 \varphi_0^2 \) and \( \frac{1}{2} m^2 \varphi_0^2 \); however, the best fit numerical values change by O(1) when solving the full coupled temperature dependent system.

\[
\frac{\sigma}{m} \frac{\varphi_0^2}{r_{peak}} = 1.3 \times 10^{36} \text{ GeV}^{-1} \left( \frac{\varphi_0}{m} \right) = 5.7 \times 10^{36} \text{ GeV}^{-1}
\]

\[
\frac{1}{2} m^2 \varphi_0^2 = 0.55 \text{ eV}^4 \quad \text{and} \quad \frac{1}{2} m^2 \varphi_0^2 = 0.21 \text{ eV}^4 \quad \text{for the temperature dependent (independent) DA model.}
\]

\[
\frac{1}{2} m^2 \varphi_0^2 \approx 0
\]

\[
H(z) = 140.0 (140.1) \text{ Mpc},
\]

compared to \( r_s = 139.8 \text{ Mpc} \) in P18. Here, \( c_s(z) \) is the speed of sound in plasma and the DA enters into the expansion rate \( H(z) \). This is well within 1σ of the \( r_s \) in the best-fit EDE scenario of P18 for \( n = 2 \), for which the best-fit Hubble constant increases to \( H_0 = 71.1 \text{ km/s/Mpc} \). This along with a larger error on \( H_0 \) resolves the tension in the EDE case. As the CMB inferences of \( r_s \) and \( H_0 \) are degenerate, with the expectation of a reduced \( r_s \) that matches P18 in the DA model, we similarly expect a high \( H_0 \) that will significantly ease the Hubble tension, if not resolve it.

For \( r_{damp} \), we expect a smaller change still, as the integral for \( r_{damp} \) is sharply peaked close to recombination and less sensitive to the expansion rate \( \sim z_{eq} \). The small difference in \( r_s \) can be absorbed by \( H_0 \), thereby diminishing the Hubble tension, while changes to \( r_{damp}/r_s \) can be absorbed by the tilt \( n_s \) of the primordial power spectrum as noted by Ref. [11, 12].

Another requirement of EDE models that succeed in resolving this discrepancy is an effective sound speed \( c_s^2 \) < 1 of perturbations in the new component [13–15]. This in part led to the success of Refs. [12, 15]. The DA model consists of a scalar field and dark radiation. The effective sound speed of a rolling scalar field is 1, while \( c_s^2 = 1/3 \) for dark radiation [32]. The coupling between the two components complicates matters, however as \( \rho_\phi < 20\% \) at \( z_{peak} \), the rest of the energy density being made up of

\[
\frac{1}{2} m^2 \varphi_0^2 \approx 0
\]
dark radiation, naively, we expect $c_s^2$ for the DA to be between $1/3 < c_s^2 < 1$.

We reiterate, that to fully understand the effect of the DA on observables beyond what is discussed here would require detailed knowledge of how the perturbations of this model evolve and MCMC simulations with all available cosmological data. The equations of evolution of the perturbations in the DA are complicated by the coupling between the dark radiation and the scalar $\phi$, and the evolving dark temperature $T_{\text{dr}}$. This derivation is further complicated by the initial conditions of the model. Depending on whether the field is present as a spectator field during inflation, its initial conditions may be isocurvature. Here, we simply seek to motivate the relevance of this model as a particle theory solution to the Hubble tension, and leave the exploration of perturbations to subsequent work.

Nonetheless, as the DA model produces a value for $r_s$ extremely close the EDE value, and little to no difference is expected in $r_{\text{damp}}$ between these two models, these expectations coupled with the predicted increase in $\omega_{\text{cdm}}$ makes the DA a promising theoretical model to deliver the EDE required to resolve the Hubble tension.

IV. DISCUSSION

In this letter, we proposed the DA as a particle-model solution to the Hubble tension. The axion couples to a dark non-Abelian gauge group, which adds an additional friction to the equation of motion of the axion, and sources a dark radiation bath as the field rolls down its potential. This overdamped system has a well understood UV-completion and has no fine-tuning concerns which are present for the phenomenological scalar field EDE solutions. From a phenomenological perspective, the injection time and total amount of EDE is quantified by two linear combinations of parameters: $\frac{\zeta_{\text{eff}}}{m}$ and $\frac{1}{2} m^2 \phi_0^2$. The full theory has additional parameters, as the friction is determined by: $\Upsilon = \kappa \alpha^5 \frac{T_{\text{dr}}}{f}$. Here, $\kappa$ is an O(10) number, $\alpha < 0.1$, $T_{\text{dr}} < f$, and $m \ll \alpha^2 T_{\text{dr}}$. Taking the sample values specified in the caption of Fig. 2, we find that these conditions are easily satisfied for many different combinations of viable parameters, for example: $m = 4 \times 10^{-25}$ eV, $T_{\text{dr}}(z_{\text{peak}}) = 0.4$ eV, $f = 0.3$ GeV, $\alpha = 0.1$, $\phi_0 = 10^{-3} M_{\text{Pl}}$, where $M_{\text{Pl}}$ is the reduced Planck scale. As discussed, only two linear combinations of degenerate parameters are constrained by the background dynamics. We expect the full perturbative analysis to lift some of the degeneracy in the parameters of the full theory and also in the potential choice for the DA.

We have solely investigated the overdamped DA regime here. Particle-sourcing friction could also play a role in an underdamped regime. Moreover, the DA can be theorised to have a UV-completion that ties its friction to the dark matter abundance. The symmetry breaking scale $f$ can, for example, be linked to the presence of heavy quarks charged under the dark SU(N). Thus, the dark matter abundance could be determined by $f$, which also controls the friction $\Upsilon$, potentially allowing a dynamical explanation for why the DA begins to roll close to matter-radiation equality. We leave a detailed exploration of this to future work.

We note that $N_{\text{eff}}$ constraints will not restrict this model. While the CMB was emitted at the redshift of recombination, the peaks of the CMB spectra in fact encode information from redshifts $z \lesssim 10^6$ [17, 33]. The DA adds dark radiation to the Universe only after $\sim z_{\text{eq}}$, unlike $N_{\text{eff}}$ which adds radiation to the Universe at all times. Their imprints on the CMB peaks are hence different - the DA is expected to cause its largest change to the CMB close to the first peak in the TT spectrum based on Refs. [17, 33], while $N_{\text{eff}}$ is not only constrained by matter-radiation equality, but also through its effect on the higher peaks in the CMB TT spectrum [34]. These distinct effects on the CMB imply that the DA model cannot be quantified by $N_{\text{eff}}$, nor be restricted by $N_{\text{eff}}$ constraints.

Numerous works [12–15] have demonstrated that the evolution of perturbations of EDE greatly impacts its effect on observables and thus its ability to solve the Hubble tension. At the background level, the DA dynamics are degenerate with the specifics of the axion potential. However, perturbations of this model and their impact on CMB data may be able to lift that degeneracy. In future work, we will derive the perturbation equations of the model and test it with an MCMC and cosmological data.

Lastly, we have invoked the DA model here as an explanation of early dark energy that resolves the Hubble tension, but this model has applications far beyond this tension. It has already been shown to be a viable candidate for cosmic inflation [21], and could similarly drive the current cosmic acceleration (for example, [35]). A family of scalar fields have often been theorised to cause the two known eras of cosmic expansion [36, 37]. We add the DA to this list.

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