

New Source for QCD Axion Dark Matter Production: Curvature Induced

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We discuss a novel mechanism for generating dark matter from a fast-rolling scalar field, relevant for both inflation and rotating axion models, and apply it specifically to the (QCD) axion. Dark matter comes from scalar field fluctuations generated by the product of the curvature perturbation and the fast-rolling background field. These fluctuations can explain the totality of dark matter in a vast axion parameter space, particularly for the QCD axion, which will be targeted by upcoming experiments. We review the constraints on this mechanism and potential gravitational-wave signatures.

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Introduction—Axion-like particles (ALP) are leading candidates to explain the dark matter (DM) of the Universe [1–3]. Several mechanisms to produce ALPs have been advocated: thermal production in the case of large axion couplings to the standard model as is the case for the QCD axion with a low decay constant f_a [4], standard misalignment [1–3], large misalignment [5], kinetic misalignment [6–10]. In the latter case, it was shown that in most of the parameter space, the axion field is fragmented and the axion energy density is not confined in the zero mode [11–14]. Axions can also be produced from the emission of the cosmic string-domain wall network [15–18]. Axions from standard misalignment and from cosmic strings are the most studied cases, although they correspond to a region of parameter space in the (axion mass m_a , axion coupling $\propto f_a^{-1}$) plane that is difficult to probe experimentally. In contrast, kinetic misalignment followed by axion fragmentation, as arising in models of

rotating axions, opens a large region of parameter space for ALP DM that will be tested in upcoming experiments.

This Letter presents a completely novel mechanism for scalar particle production in the early Universe, that can be applied generally to scalar DM, and specifically to rotating axion models. Rather than coming from the energy initially stored in the zero-mode background field, the energy density comes from the fluctuations of the scalar field that are induced by primordial (inflationary) curvature perturbations, see also Refs. [19–22]. Such a source term is effective if the background scalar field has a large kinetic energy.

The generation of these fluctuations was pointed out in [12,23], and some important implications, when applied to massless fields, were shown in [24]. In this Letter, we study how they can behave as DM, as suggested in [14]. Depending on the scalar mass value as well as other UV parameters, their contribution may be the dominant one over the abundance from the zero mode, as shown in Fig. 1. This general mechanism could have other applications. For instance, if the scalar field is the inflaton itself and inflation is ended by a fast-rolling stage (like in quintessential inflation [25,26] and string cosmologies [23]), fluctuations can be efficiently produced and could explain DM. The next two sections are rather general and model independent, while the last section applies to a model of rotating axions.

Curvature-induced fluctuations of fast-rolling field—In the absence of anisotropic stress, the space-time metric in the conformal Newtonian gauge is

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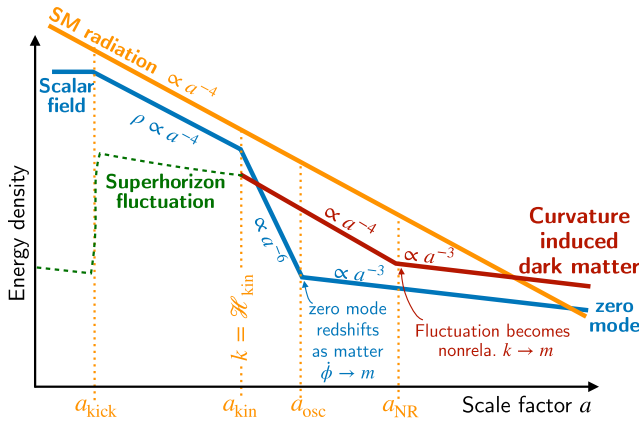


FIG. 1. Cartoon of the energy-density evolution of a fast-moving scalar field condensate in blue, which behaves as radiation $\sim a^{-4}$, kination $\sim a^{-6}$, and matter $\sim a^{-3}$, consecutively, during the radiation era of the Universe. The second stage starts at a_{kin} when the field dynamics is kinetic-energy dominated. The last stage happens when the field gets trapped and oscillates inside the potential barrier at a_{osc} , leading to DM via kinetic misalignment or fragmentation. The red line shows the energy-density evolution of scalar-field fluctuations ρ_{fluct} induced by primordial curvature perturbations. Initially produced relativistically, they become nonrelativistic at a_{NR} , and can serve as DM. The orange line shows the standard model radiation. Another possible scenario is if the scalar field is the inflaton itself and fast rolls after inflation, that would correspond to setting $a_{\text{kick}} = a_{\text{kin}}$. Plot is schematic and not to scale.

$$ds^2 = a^2 [-(1 + 2\Phi)d\eta^2 + (1 - 2\Phi)\delta_{ij}dx^i dx^j], \quad (1)$$

where Φ is the scalar metric fluctuation whose Fourier modes are defined by $\Phi_{\mathbf{k}} = \int d^3\mathbf{x} e^{i\mathbf{k}\cdot\mathbf{x}} \Phi(\mathbf{x})$, and η is conformal time. On superhorizon modes, Φ is frozen and related to the gauge-invariant comoving curvature perturbation $\mathcal{R} \equiv -\Phi - 2(\mathcal{H}^{-1}\Phi' + \Phi)/3(1 + \omega) \simeq -3\Phi/2$, where $(\cdot)' \equiv d(\cdot)/d\eta$, $\mathcal{H} = aH$ is the comoving Hubble rate and, ω is the Universe equation of state. We assume radiation domination $\omega = 1/3$ (unlike [24]) and treat the case of kination domination in Fig. 3.

Consider a fast-rolling scalar field, $\phi = \bar{\phi} + \delta\phi$, with $\bar{\phi}$ its homogeneous part and $\delta\phi$ its fluctuations, as already discussed in [24]. During fast rolling—when kinetic and gradient energies dominate and potential energy is negligible—and under the conditions $\Phi, \delta\phi \ll 1$, the scalar field equation of motion becomes

$$\delta\phi_{\mathbf{k}}'' + 2\mathcal{H}\delta\phi_{\mathbf{k}}' + k^2\delta\phi_{\mathbf{k}} = 4\Phi_{\mathbf{k}}'\bar{\phi}'. \quad (2)$$

The right-hand side of this equation is the key point. The nonzero scalar speed $\bar{\phi}'$ combined with the curvature perturbation $\Phi_{\mathbf{k}}'$ sources the scalar fluctuation with momentum k , leading to the subhorizon solution ($k \gg \mathcal{H}$)

$$\phi_{\mathbf{k}}(\eta) \simeq \alpha \mathcal{R}_{\mathbf{k}}(0) \frac{\bar{\phi}'}{\mathcal{H}} \cos(k\eta), \quad (3)$$

for modes entering during the fast-rolling stage $k < \mathcal{H}_{\text{kin}}$ and $\alpha \simeq 0.52$ (see Ref. [12], Appendix B.2, and Sec. A of Supplemental Material [27] for a rederivation). At all orders in $\delta\phi$, the scalar field energy density is $\rho_{\phi} = [(\phi')^2 + (\nabla\phi)^2]/(2a^2) = \bar{\rho}_{\phi}(1 + \delta_{\phi} + \delta_{\phi}^2/2)$, with $\bar{\rho}_{\phi} = (\bar{\phi}')^2/2$ and the linear component of the density contrast $\delta_{\phi} \equiv (\rho_{\phi}/\bar{\rho}_{\phi} - 1)_{\text{linear}} = 2\delta\phi'/\bar{\phi}'$. We deduce the volume average of the energy density fluctuations [24]

$$\rho_{\text{fluct}}(\eta) \equiv \langle \rho_{\phi} - \bar{\rho}_{\phi} \rangle = \frac{\bar{\rho}_{\phi}(\eta)}{2} \int d\log k \mathcal{P}_{\delta_{\phi}}(k, \eta), \quad (4)$$

where $\mathcal{P}_{\delta_{\phi}}(k) \equiv \int d^3\mathbf{x} e^{i\mathbf{k}\cdot\mathbf{x}} \langle \delta_{\phi}(0)\delta_{\phi}(\mathbf{x}) \rangle$ is the dimensionless power spectrum of the linear density contrast. We used the important fact that $\langle \delta_{\phi} \rangle \propto \langle \mathcal{R}_{\mathbf{k}} \rangle = 0$ for adiabatic primordial curvature perturbations. Using Eq. (3) and averaging over oscillations, we get $\mathcal{P}_{\delta_{\phi}} \simeq 2\alpha^2 \mathcal{P}_{\mathcal{R}}(k)(k/\mathcal{H})^2$ and Eq. (4) leads to the fractional energy density $\Omega_{\text{fluct}} \equiv \rho_{\text{fluct}}/\rho_{\text{tot}}$ of the fluctuations of the fast-rolling field

$$\frac{d\Omega_{\text{fluct}}(\eta)}{d\log k} = \alpha^2 \Omega_{\phi}(\eta_{\text{kin}}) \left(\frac{k}{\mathcal{H}_{\text{kin}}} \right)^2 \mathcal{P}_{\mathcal{R}}(k), \quad (5)$$

where $\mathcal{H}_{\text{kin}} = a_{\text{kin}} H_{\text{kin}}$ is the value of the Hubble rate when the scalar field starts rolling at a_{kin} . We used $\mathcal{H} \propto a^{-1}$ and introduced $\Omega_{\phi}(\eta) \equiv \bar{\rho}_{\phi}/\rho_{\text{tot}}$ which redshifts like $\propto a^{-2}$ during radiation domination. We now study the novel implications for DM.

Dark matter from curvature-induced scalar fluctuations—We consider a scalar field ϕ with mass m . Below a scale factor $a_{\text{NR}}(k) = k/m$, the mode k turns nonrelativistic and its energy density redshifts like matter. From integrating Eq. (5) over all modes, we deduce the total fractional energy abundance of scalar fluctuation today

$$\frac{\Omega_{\text{fluct}}^0}{\Omega_{\text{rad}}^0} = \alpha^2 \Omega_{\phi}(\eta_{\text{kin}}) \left(\frac{a_0 m}{\mathcal{H}_{\text{kin}}} \right) \int_{k_{\text{osc}}}^{k_{\text{kin}}} \frac{dk}{k} \left(\frac{k}{\mathcal{H}_{\text{kin}}} \right) \mathcal{P}_{\mathcal{R}}(k), \quad (6)$$

where $h^2 \Omega_{\text{rad}}^0 \simeq 4.21 \times 10^{-5}$ [30] with $h \simeq 0.68$ the reduced Hubble factor and a_0 is today's scale factor. The expression at any time is provided in Eq. (A25) of Supplemental Material [27]. We model the curvature power spectrum as a power law

$$\mathcal{P}_{\mathcal{R}}(k) = \mathcal{P}_{\mathcal{R}}(k_{\star})(k/k_{\star})^{n_s(k)-1}, \quad (7)$$

where the amplitude and spectral index are measured by Planck at the pivot scale $k_{\star} \equiv 0.05 \text{ Mpc}^{-1}$ [31], $\mathcal{P}_{\mathcal{R}}(k_{\star}) \simeq 2.099(29) \times 10^{-9}$, $n_s(k_{\star}) \simeq 0.965(4)$. For k away from $k_{\star}$, we use the Taylor expansion $n_s(k) = n_s(k_{\star}) - (3.65 \times 10^{-4}) \ln(k/k_{\star})$ obtained after analyzing

300 single-field inflation models [32]. Plugging Eq. (7) into Eq. (6), the total abundance for $\mathcal{H}_{\text{kin}} \gg k_{\text{osc}}$ is dominated by the mode k_{kin} that reenters at a_{kin}

$$\frac{\Omega_{\text{fluct}}^0}{\Omega_{\text{rad}}^0} = \alpha^2 \Omega_{\phi}(\eta_{\text{kin}}) \left(\frac{a_0 m}{\mathcal{H}_{\text{kin}}} \right) \frac{\mathcal{P}_{\mathcal{R}}(k_{\text{kin}})}{n_s}. \quad (8)$$

Note that $\mathcal{P}_{\mathcal{R}}(k_{\text{kin}})$ is poorly constrained and could be in principle many orders of magnitude larger than $\mathcal{P}_{\mathcal{R}}(k_*)$, a fact exploited in inflationary models that predict primordial black holes [33]. The abundance of fluctuations $\delta\phi_{\mathbf{k}}$ should be compared with the one of the zero mode $\bar{\phi}$

$$\Omega_{\text{zero}}^0 \simeq \Omega_{\text{rad}}^0 \Omega_{\phi}(\eta_{\text{kin}}) \left(\frac{a_0}{a_{\text{osc}}} \right) \left(\frac{a_{\text{kin}}}{a_{\text{osc}}} \right)^2, \quad (9)$$

where a_{osc} is the scale factor when the homogeneous scalar field stops fast rolling and starts oscillating, redshifting like matter [6,7]. This occurs when the scalar field energy density drops below $\bar{\rho}_{\phi}(\eta_{\text{osc}})$. Denoting $\bar{\rho}_{\phi}(\eta_{\text{osc}}) \simeq m_a^2 f_a^2$ in view of applying this to axions, with mass m denoted by m_a , and decay constant f_a , this leads to $a_{\text{kin}}/a_{\text{osc}} \simeq (m_a f_a)^{1/3} / \bar{\rho}_{\phi}^{1/6}(\eta_{\text{kin}})$. We deduce

$$\frac{\Omega_{\text{fluct}}^0}{\Omega_{\text{zero}}^0} \simeq 0.9 \frac{\alpha^2}{n_s} \Omega_{\phi}^{1/2}(\eta_{\text{kin}}) \left(\frac{10^{10} \text{ GeV}}{f_a} \right) \left(\frac{\mathcal{P}_{\mathcal{R}}(k_{\text{kin}})}{2.1 \times 10^{-9}} \right). \quad (10)$$

The fluctuation energy density can be further enhanced by up to one order of magnitude when the scalar field dominates the Universe, inducing a kination era, see Sec. E of Supplemental Material [27]. Such a situation can be natural in rotating axion models and is associated with large backgrounds of gravitational waves (GW) [34–36]. We illustrate the observability prospects at future GW observatories in Fig. 3.

Fluctuations are produced with $k \gg m$ and initially behave as radiation, so they are constrained by structure formation. The Lyman- α limit on thermal DM, $m_{\text{WDM}} \gtrsim 2 \text{ keV}$ [37], translates into an upper bound on the DM equation of state at matter-radiation equality, $\omega_{\text{DM,eq}} \lesssim 10^{-8}$ [38]. In the nonrelativistic limit ($k \ll m$), the equation of state for a mode is $\omega_{\text{fluct}} \simeq k/(3am)$, see Sec. B of Supplemental Material [27]. The dominant contribution arises from $k \simeq k_{\text{kin}}$, which leads to the bound

$$T_{\text{kin}} \lesssim 1.1 \times 10^{15} \text{ GeV} \left(\frac{m}{1 \text{ eV}} \right) \frac{g_{*s}^{1/3}(T_{\text{kin}})}{g_*^{1/2}(T_{\text{kin}})}, \quad (11)$$

where T_{kin} is the Universe temperature at a_{kin} . This condition is satisfied in the model parameter space of Fig. 2, but is modified if the scalar field dominates the Universe and leads to a constraint in Fig. 3.

Equation (5) suggests that the energy density fraction in radiationlike fluctuations is fixed over time and is

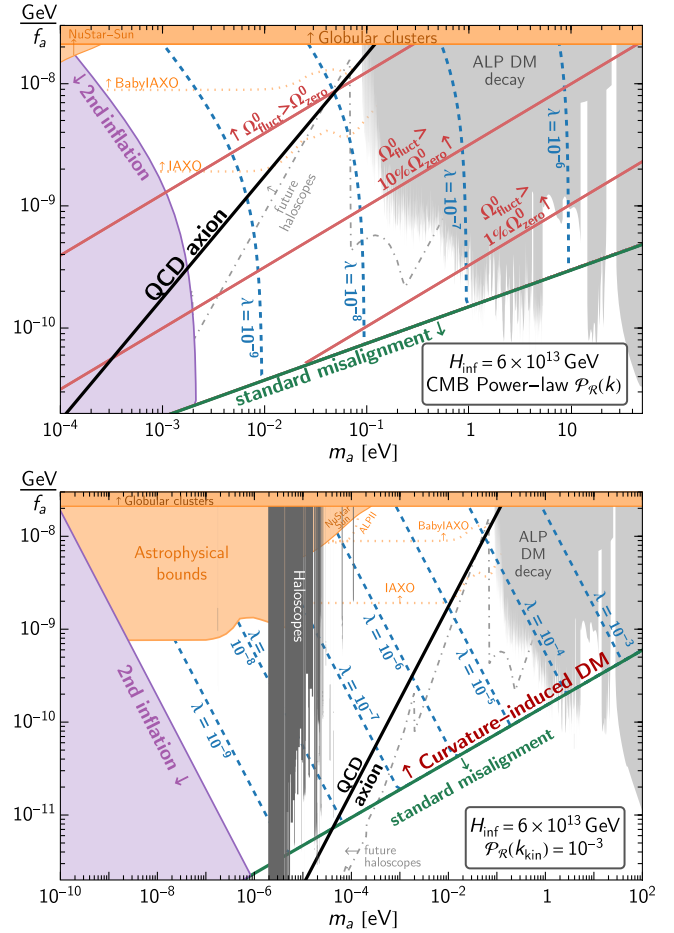


FIG. 2. The rotating axion with quartic potential (12) produces the observed DM abundance in the region above the green line, via curvature-induced production and kinetic misalignment, i.e., $\Omega_{\text{DM}}^0 = \Omega_{\text{fluct}}^0 + \Omega_{\text{zero}}^0$. For each λ , DM is overproduced on the right of each dashed blue line. The top panel assumes $\mathcal{P}_{\mathcal{R}}(k)$ using Eq. (7). Red lines show the fraction of curvature-induced DM to zero-mode DM. The bottom panel assumes $\mathcal{P}_{\mathcal{R}}(k_{\text{kin}}) = 10^{-3}$. The purple region is ruled out by the second inflation bound (13), while the warmness constraint (11) does not appear. The gray regions are the combined bounds for ALP DM from haloscopes and ALP-decay searches [39–49], while the orange constraints apply to any ALP. The bounds from ALP decay would become weaker if axion does not form DM [50]. Regions above gray and orange dotted and dot-dashed will be probed by future experiments. All constraints are compiled from [51], assuming $g_{\phi\gamma\gamma} \simeq 1.02\alpha_{\text{em}}/(2\pi f_a)$ as motivated by the KSVZ model [52,53].

proportional to $\Omega_{\phi} \mathcal{P}_{\mathcal{R}}(k_{\text{kin}}) \ll 1$. Therefore, the curvature-induced axion DM is exempted from the dark radiation bound imposed by big-bang nucleosynthesis [54] and cosmic microwave background [31].

In the above analysis, we neglected the interactions between the curvature-and-fast-roll-induced fluctuations and the zero mode. We provide an estimate of this effect in Sec. C of Supplemental Material [27] for the

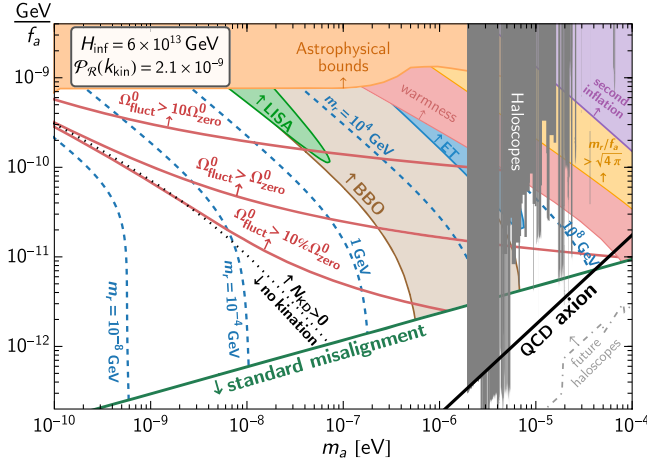


FIG. 3. Detectability of the kination peak of inflationary gravitational waves at LISA (green), BBO (brown), and ET (blue), using the scalar potential (16) with $\{\kappa, l\} = \{10^{-4}, 13\}$. The yellow region violates perturbativity ($m_a/f_a > \sqrt{4\pi}$). The purple region is excluded by the second inflation bound [Eq. (8.22) of [35]]. The pink region has DM too warm [$\omega_{\text{fuct}}^{\text{eq}} \simeq k_{\text{kin}}/(3a_{\text{eq}}m_a) > 10^{-8}$]. Red lines represent the ratio $\Omega_{\text{fuct}}^0/\Omega_{\text{zero}}^0$. Below the black dotted line, there is no kination era; the curvature-induced DM is produced in radiation era and becomes less abundant due to $\Omega_\phi(\eta_{\text{kin}}) < 1$ in Eq. (10).

rotating-axion model, which indicates that the zero-mode contribution could be further suppressed and the dominant relic would come from the fluctuations. Nonetheless, a dedicated numerical simulation would be needed for estimating the size of such corrections.

Application: DM from rotating axions—We now apply our findings to a well-motivated framework realizing fast-moving axions: rotating complex fields, as discussed extensively in Refs. [6,8,9,12,14,34–36,55]. They rely on a complex scalar field $\chi = S e^{i\phi/f_a}$, corresponding to the Peccei-Quinn field in the case of the QCD axion [56,57], where S and ϕ describe the radial and angular field directions. The latter is identified with the axion. The key evolution is that the dynamics start at a large- S value rather than at the origin of the potential. Such a scenario requires four ingredients: (i) a scalar potential invariant under a global $U(1)$ symmetry with the spontaneous breaking minimum at $S = f_a$ that allows a motion along the axion direction at late times, (ii) an explicit $U(1)$ -breaking term that operates only at early times and kicks the complex field into an elliptic orbit, (iii) a large radial initial field value S_{ini} that enables the angular kick, and (iv) a radial-motion damping mechanism to bring the field orbit into a circle, which later settles down to the potential minimum and starts the kination behavior (i.e., $\rho_\phi \propto a^{-6}$). Here, we remain agnostic on the origins of some ingredients as they depend on the complete realization of such models, which were discussed extensively in, e.g., Refs. [14,35].

Let us consider for ingredient (i) a $U(1)$ -invariant quartic potential and for ingredient (iii) a Hubble-induced mass term:

$$V(\chi) = \lambda(|\chi|^2 - f_a^2)^2 - \begin{cases} \lambda_H H^2 |\chi|^2 & \text{during inflation} \\ 0 & \text{after inflation,} \end{cases} \quad (12)$$

which also suppresses the axion isocurvature perturbation.

During inflation, the field is driven to the potential minimum $S_{\text{ini}} = f_a \sqrt{1 + \beta^2}$ with $\beta^2 \equiv \lambda_H H_{\text{inf}}^2 / (2\lambda f_a^2)$, where H_{inf} is the Hubble scale during inflation. After inflation, we assume for simplicity that the Universe reheated instantaneously [58]. The scalar field energy density is set by the potential, $\rho_{\chi,\text{ini}} = V(S_{\text{ini}}) = \lambda f_a^4 \beta^4 = \lambda_H^2 H_{\text{inf}}^4 / (4\lambda)$, and is subdominant with respect to the total radiation energy density of the Universe. The field remains frozen due to Hubble friction.

The complex field starts moving when $3H_{\text{kick}} = m_{S,\text{kick}} \equiv \sqrt{V'(S_{\text{ini}})/S_{\text{ini}}}$. This happens at the Hubble scale, $H_{\text{kick}} = \sqrt{2\lambda_H} H_{\text{inf}}/3$, which is slightly after inflation ends. To prevent the “second inflation” (i.e., the scalar field dominates the Universe before the kick occurs), we require that $\rho_{\chi,\text{ini}} < \rho_{\text{kick}}^{\text{rad}} = 3M_{\text{Pl}}^2 H_{\text{kick}}^2$ with $M_{\text{Pl}} \simeq 2.44 \times 10^{18}$ GeV, which translates into the lower bound of λ ,

$$\lambda > \frac{3\lambda_H}{8} \left(\frac{H_{\text{inf}}}{M_{\text{Pl}}} \right)^2 \simeq 2.3 \times 10^{-10} \left(\frac{\sqrt{\lambda_H} H_{\text{inf}}}{6 \times 10^{13} \text{ GeV}} \right)^2. \quad (13)$$

As the field orbits inside a quartic potential, the total energy density of the complex field redshifts like radiation as $\sim a^{-4}$. The last remaining ingredient (iv) “the radial damping” must occur before the field reaches the potential minimum. We refer to, e.g., Refs. [9,14,35,36,55] for the realizations of radial damping using the interactions between the complex field and other particles in the thermal plasma. Once the orbit reaches the minimum $S = f_a$, the axion can be considered as a freely rolling field, whose fluctuations can be efficiently induced by the curvature perturbation.

From the time of the kick to the time when the complex field reaches the bottom of its potential, the Universe has expanded by $a_{\text{kin}}/a_{\text{kick}} = S_{\text{ini}}/f_a = \sqrt{1 + \beta^2}$. The energy densities in the rotating homogeneous axion $\vec{\phi}$ and the total energy density of the Universe at the start of the kination behavior read,

$$\bar{\rho}_{\phi,\text{kin}} = \epsilon \rho_{\chi,\text{ini}} \left(\frac{a_{\text{kick}}}{a_{\text{kin}}} \right)^4 = \epsilon \left(\frac{\lambda_H^2}{4\lambda} \right) \frac{H_{\text{inf}}^4}{(1 + \beta^2)^2}, \quad (14)$$

$$\rho_{\text{tot,kin}} = \rho_{\text{kick}}^{\text{rad}} \left(\frac{a_{\text{kick}}}{a_{\text{kin}}} \right)^4 = \left(\frac{2\lambda_H}{3} \right) \frac{M_{\text{Pl}}^2 H_{\text{inf}}^2}{(1 + \beta^2)^2}, \quad (15)$$

where we use that both quantities scale as a^{-4} , and ϵ parametrizes the energy fraction in the axion compared to that of the radial mode. ϵ encodes the orbit ellipticity and can be determined by the exact form of the explicit $U(1)$ -breaking term, see, e.g., Refs. [8,35]. Since $\Omega_{\text{fluct}}^0/\Omega_{\text{zero}}^0 \propto \sqrt{\epsilon} \mathcal{P}_{\mathcal{R}}(k_{\text{kin}})$, a smaller ϵ can be compensated by a larger $\mathcal{P}_{\mathcal{R}}$. We assume $\epsilon \sim 1$, as for $\epsilon \ll 1$, the complex field could also experience parametric resonance [8]. Note that $\epsilon \sim 1$ can be obtained in supersymmetric realizations where the potential is nearly quadratic [14,35,55].

Using Eqs. (14) and (15) in Eq. (10), we show in Fig. 2 the region where DM comes from the curvature-induced fluctuation and the kinetic misalignment, $\Omega_{\text{DM}}^0 = \Omega_{\text{fluct}}^0 + \Omega_{\text{zero}}^0$ where $\Omega_{\text{DM}}^0 = 0.2657$ [30] is the fractional energy density in DM today and we assume no interaction between the fluctuation and zero modes. For $\mathcal{P}_{\mathcal{R}}(k)$ from Eq. (7), the curvature-induced axion contributes to DM with at least more than 1% of the zero-mode abundance, which goes up to more than 10% for the QCD axion. The purple region is ruled out by the second inflation constraint (13). The curvature-induced fluctuation opens up the model parameter space. Note also the dependence on the inflationary scale H_{inf} and the Hubble-mass coefficient λ_H in Eqs. (14) and (15) such that [Eq. (10)] scales as $\Omega_{\text{fluct}}^0/\Omega_{\text{zero}}^0 \propto \sqrt{\lambda_H} H_{\text{inf}}$. Figure 2 uses the largest value suggested by Planck data [59] and a fixed $\lambda_H = 1$. In the bottom panel, we show predictions for $\mathcal{P}_{\mathcal{R}}(k_{\text{kin}}) = 10^{-3}$, close to the primordial black holes overproduction bound [60], that leads to more enhanced axion fluctuations and larger parameter space. We have checked that, for $m_a < 100$ eV, the curvature-induced fluctuation dominates the DM abundance if $\mathcal{P}_{\mathcal{R}}(k_{\text{kin}}) \gtrsim 3.5 \times 10^{-7}$. We show the parameter spaces for other $\mathcal{P}_{\mathcal{R}}(k_{\text{kin}})$ values in Supplemental Material [27]. For the parameter space in Fig. 2, the temperature at which the fluctuation (of k_{kin} mode) becomes nonrelativistic is $T_{\text{NR}} \gtrsim \text{MeV}$.

Experimental tests—In our example model, the DM from kinetic misalignment alone cannot explain DM for $m_a \lesssim 2.1$ meV due to the second inflation bound (13). As shown in Fig. 2, the new mechanism opens up this range of m_a , which is a target region for the ALPS II experiment [61], for helioscopes BabyIAXO and IAXO [62,63], and for future haloscopes [51]. Regarding gravitational signatures, it was shown in [13] that kinetic fragmentation of the axion condensate can lead to compact minihalos. A new analysis will be required to study the formation of compact structures in the new scenario studied in this Letter. Next, we also stress the potential signatures in gravitational waves, which will arise in the regime where the field temporarily dominates the energy density of the Universe over radiation, see Fig. 3.

From radiation-domination to kination—If the fast-rolling axion dominates the Universe, it can lead to a kination era if the quartic potential (12) is replaced by a nearly quadratic one [9,34–36,55]:

$$V(\chi) = m_r^2 |\chi|^2 \left(\log \frac{|\chi|^2}{f_a^2} - 1 \right) + m_r^2 f_a^2 + \frac{\kappa^2}{M_{\text{Pl}}^{2l-6}} |\chi|^{2l-2}, \quad (16)$$

together with the Hubble-mass term in Eq. (12). Its cosmological history is reviewed exhaustively in [35]. With the same ingredients as for the quartic model, the fluctuations and the zero mode can produce axion DM over a large region of parameter space shown in Fig. 3 (see Supplemental Material [27] for the slightly different relic abundance expressions). Strikingly, the kination era induced by this DM axion imprints a triangular smoking-gun signature [34–36] in the inflationary gravitational-wave background that could be observed at Laser Interferometer Space Antenna (LISA) [64], Einstein Telescope (ET) [65], and Big Bang Observer (BBO) [66]. Note that our results rely on the fluctuations produced during the kination era only. The effect of the matter era preceding kination could, in fact, enhance the relic abundance, which we leave for a future analysis.

Conclusion—We presented a novel source for scalar DM in models where the scalar field experienced a short stage of fast rolling in its early history. We studied the implications in the framework of rotating axion models. Assumptions and constraints on parameters are similar to models of kinetic misalignment, in particular those conditions related to radial mode damping as studied extensively in [35] and [14]. What we learn in this letter is that axion DM can be dominantly produced from primordial curvature perturbations, as fluctuations rather than as the well-known oscillations of the zero-mode condensate, without demanding any specific assumption on the realization of inflation.

Our findings are highly relevant to a large number of upcoming axion experiments that will probe the specific region of parameter space (where the axion cannot be DM from the standard misalignment mechanism [1–3]) where we predict our novel source of axions to be the dominant one and sufficient to explain DM. Precise predictions will need to be checked with numerical simulations because of backreaction effects. While fluctuations from the condensate fragmentation in [12] do not impact much the axion relic abundance because they are produced when mostly nonrelativistic, in this Letter, the fluctuations are highly relativistic at the time of production and interactions between the fluctuations and the background could modify the predictions presented in our analysis. The corresponding lattice simulations will be challenging as they involve a long evolution and a large hierarchy of scales.

Note added—We recently became aware of an upcoming related paper by A. Bodas, A. Ghalsasi, K. Harigaya, R. Co, and L.-T. Wang [67].

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Data availability—The data that support the findings of this article are not publicly available upon publication because it is not technically feasible and/or the cost of preparing, depositing, and hosting the data would be prohibitive within the terms of this research project. The data are available from the authors upon reasonable request.

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