

Coherent Freeze-Out of Dark Matter

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We propose a novel coherent freeze-out mechanism where a weakly interacting massive particle (WIMP) is quadratically coupled to a light axionlike particle (ALP). Although the coupling is too feeble to thermalize the ALP, coherent forward scattering induces medium-dependent mass shifts that significantly modify both WIMP freeze-out and ALP misalignment dynamics. The symmetry of the ALP potential is broken at high temperatures and restored through either a first-order transition or a crossover. In the former, WIMPs alone compose dark matter with annihilation cross sections enhanced by up to 3 orders of magnitude relative to the standard scenario; in the latter, a Planck-suppressed coupling naturally yields an ALP abundance of the order of the observed dark matter density, largely independent of its initial displacement and mass.

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Introduction—The particle nature of dark matter (DM) remains unknown. Two well-motivated candidates beyond the standard model (SM) are weakly interacting massive particles (WIMPs) and axionlike particles (ALPs) (see Refs. [1,2] for recent reviews). They inhabit very different mass ranges and arise from distinct theoretical frameworks, and their cosmological production mechanisms likewise differ sharply. WIMPs begin in thermal equilibrium and obtain their relic abundance through freeze-out, a process largely insensitive to ultraviolet (UV) physics. ALPs, by contrast, never thermalize due to their extremely weak couplings; their relic abundance is instead determined by the initial field displacement at the onset of coherent oscillation, typically set by inflationary fluctuations or other high-scale dynamics.

In this Letter, we raise a generic question: If both a WIMP and an ALP are present in the theory, what are the cosmological consequences of an interaction between them? Since the coupling is typically too weak to thermalize the ALP, one might expect their dynamics to remain independent, as no significant momentum exchange occurs. However, this intuition fails. The ALP behaves as a coherent, homogeneous field, while the WIMP population forms an approximately uniform thermal bath. Coherent forward scattering between the two backgrounds induces

mean-field corrections to the dispersion relations of both species: The ALP potential is deformed by the thermal WIMP bath, and the WIMP effective mass is shifted by the ALP background. These medium effects lead to coupled, nontrivial dynamics in their cosmological evolution. The goal of this Letter is to show that even an extremely weak (Planck-suppressed) coupling between the two sectors leads to striking cosmological consequences.

To illustrate the mechanism, we consider a simple model in which WIMP and ALP fields interact through an effective quadratic coupling [see Eq. (1)]. The thermal bath of WIMPs induces a temperature-dependent correction to the ALP potential, spontaneously breaking its discrete symmetry at high temperatures and restoring it at low temperatures. Such an “inverse phase transition” occurs not only in many particle-physics models [3–8], but also in condensed-matter systems [9]. In the broken phase, the ALP field is displaced to a new vacuum, whose backreaction reduces the effective WIMP mass and, therefore, affects its freeze-out. At the same time, ALP coherent oscillations begin only after the phase transition, with the initial value of the field fixed by the high-temperature dynamics, in contrast with the usual misalignment mechanism.

Modifications to the dynamics of a light scalar in the presence of a thermal bath have been discussed in several contexts (see, e.g., Refs. [10–18]). The crucial distinction in our analysis is that we fully incorporate the backreaction of the scalar on the thermal bath—an effect that plays a central role in our mechanism.

Dynamics—We begin with the following effective coupling:

$$\mathcal{L}_{\text{eff}} \supset \frac{1}{\Lambda} \bar{\chi} \chi \frac{\phi^2}{2}, \quad (1)$$

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where Λ is the cutoff scale, χ is a fermionic WIMP (either Dirac or Majorana), and ϕ is a light real scalar. For convenience, we refer to ϕ as the ALP hereafter, although it may be CP -even or CP -odd. The model contains three hierarchical scales, $\Lambda \gg m_\chi \gg m_\phi$, with m_χ and m_ϕ denoting the physical zero-temperature masses of χ and ϕ , respectively. For simplicity, we do not include higher-order terms in ϕ in the zero-temperature potential. The coupling in (1) is generic and can arise in a wide range of UV completions. We take a definite sign for the interaction, which is essential for our analysis. Scalar-fermion interactions with the same structure as (1), including the sign, appear in familiar contexts. For example, this is the structure of the leading nonderivative interaction between nucleons and pions in the SM or between nucleons and the QCD axion in minimal extensions, both of which follow directly from chiral effective theory [19–24]. Throughout, we assume the coupling is small enough that ϕ never thermalizes (neither chemically nor kinematically), which requires $\Lambda \gg \sqrt{T_{\text{RH}} M_{\text{Pl}}}$, where T_{RH} is the reheating temperature of the Universe and $M_{\text{Pl}} \approx 1.2 \times 10^{19}$ GeV is the Planck mass.

Even without momentum exchange, coherent forward scattering modifies the propagation of both species. In the region of interest, χ is in thermal equilibrium with the SM. The thermal χ background shifts the ALP mass, and the classical ALP field shifts the WIMP mass:

$$m_{\phi,\text{eff}}^2 = m_\phi^2 - \frac{\langle \bar{\chi}\chi \rangle_T}{\Lambda}, \quad m_{\chi,\text{eff}} = \left| m_\chi - \frac{\phi^2}{2\Lambda} \right|, \quad (2)$$

where

$$\langle \bar{\chi}\chi \rangle_T = g_\chi \int \frac{d^3\mathbf{k}}{(2\pi)^3} \frac{m_{\chi,\text{eff}}}{E_{\mathbf{k}}} f_\chi(\mathbf{k}) \quad (3)$$

is the thermally averaged number density, with $E_{\mathbf{k}} \equiv \sqrt{\mathbf{k}^2 + m_{\chi,\text{eff}}^2}$ and $g_\chi = 4(2)$ if χ is Dirac (Majorana). Crucially, the correction to m_ϕ depends on $m_{\chi,\text{eff}}$, which itself depends on ϕ ; the dynamics of the two sectors are, therefore, *coupled*.

For analytic control, we adopt a Maxwell-Boltzmann distribution, $f_\chi = e^{-E_{\mathbf{k}}/T}$. (A more complete treatment using the full one-loop thermal potential is provided in Supplementary Material Sec. S4 [25]; however, it does not affect the main results discussed below.) This leads to the ALP potential $V(\phi, T) \equiv (g_\chi/2\pi^2)m_\chi^4 \mathcal{V}(\phi, x)$ with the normalized potential

$$\mathcal{V}(\phi, x) = -\frac{\phi^2 \gamma^2 K_1(\gamma x)}{2x} + \kappa \frac{\phi^2}{2}, \quad (4)$$

where K_1 is the modified Bessel function. Here, we have defined $x \equiv m_\chi/T$, $\phi \equiv \phi/\sqrt{m_\chi \Lambda}$, $\gamma \equiv |1 - \phi^2/2|$, and

$$\kappa \equiv \frac{2\pi^2 m_\phi^2 \Lambda}{g_\chi m_\chi^3}. \quad (5)$$

The first term in (4) encodes the thermal effect of the WIMP bath, while the second is the bare ALP mass. The resulting potential, shown in the left panel in Fig. 1, contains a term quartic in ϕ due to the backreaction on $m_{\chi,\text{eff}}$, ensuring it is bounded from below. At early times ($\gamma x \ll 1$), the thermal term acts like a negative T^2 mass and spontaneously breaks the Z_2 symmetry ($\phi \rightarrow -\phi$); at late times, the bare mass dominates and restores it. The order of this transition is set by a critical coupling $\kappa_c \approx 0.27$, which can be calculated in Ginzburg-Landau theory (see Supplementary Material Sec. S1 [25]). For $\kappa < \kappa_c$, it corresponds to a first-order phase transition (FOPT), and

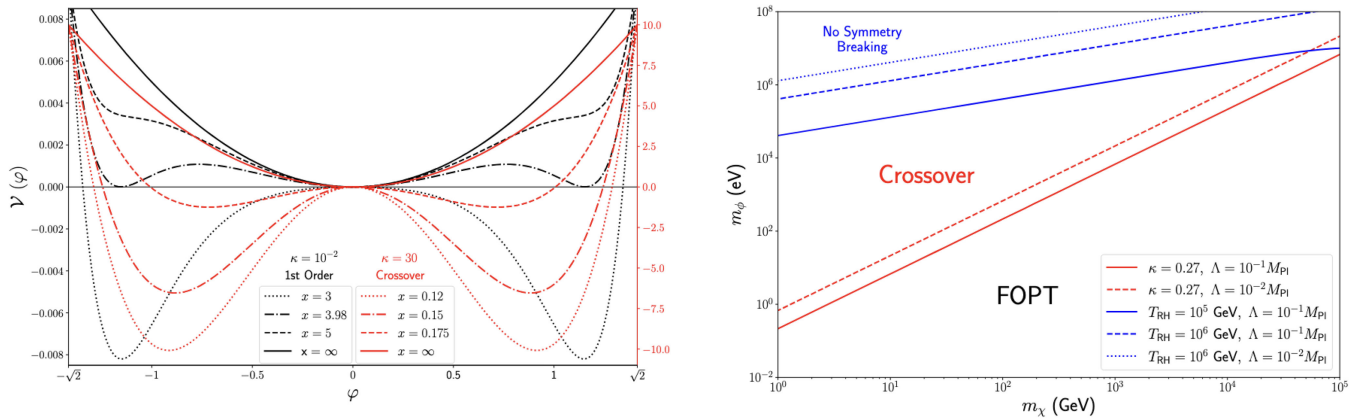


FIG. 1. Left: normalized ALP potential (4) at various temperatures. For $\kappa < \kappa_c \approx 0.27$, the transition is first order (black curves); for $\kappa > \kappa_c$, it is a crossover (red curves). Right: phase diagram in the (m_χ, m_ϕ) plane. The $\kappa = \kappa_c$ contour (red lines) separates first-order and crossover regimes. The condition $\kappa = K_1(x_{\text{RH}})/x_{\text{RH}}$ (blue lines), with $x_{\text{RH}} = m_\chi/T_{\text{RH}}$, marks where symmetry is unbroken at reheating temperature.

for $\kappa > \kappa_c$, it is a crossover. We truncate the potential at $|\varphi| = \sqrt{2}$, where $m_{\chi,\text{eff}} \rightarrow 0$ and higher-order terms become important. Since the field starts with negligible velocity after inflation, $|\dot{\varphi}|$ never exceeds this bound.

In the radiation era, the ALP field obeys the following equation of motion (EOM):

$$\varphi'' + \frac{2}{x}\varphi' - \eta x \left(1 - \frac{\varphi^2}{2}\right) \times \left[(1 - \varphi^2)K_1(\gamma x) + \frac{\varphi^2}{2}\gamma x K_0(\gamma x) \right] \varphi + \eta \kappa x^2 \varphi = 0, \quad (6)$$

where prime denotes the derivative with respect to x and

$$\eta \equiv \frac{g_\chi}{2\pi^2} \frac{1}{1.66^2 g_*} \frac{M_{\text{Pl}}^2}{m_\chi \Lambda}, \quad (7)$$

with g_* the number of relativistic degrees of freedom. The EOM (6) contains only two free parameters (κ and η), and its solutions are studied in detail in Supplementary Material Sec. S2 [25]. Depending on κ , the resulting parameter space falls into three classes (see right panel of Fig. 1). (1) *No symmetry breaking*— $\kappa > K_1(x_{\text{RH}})/x_{\text{RH}}$ with $x_{\text{RH}} \equiv m_\chi/T_{\text{RH}}$. Symmetry is never broken, and the two sectors decouple. (2) *Crossover*— $\kappa_c < \kappa < K_1(x_{\text{RH}})/x_{\text{RH}}$. Symmetry is broken at high T and restored smoothly at low T . (3) *FOPT*— $\kappa < \kappa_c$ and $\kappa < K_1(x_{\text{RH}})/x_{\text{RH}}$. Symmetry is broken at high T and restored at low T through a first-order transition. Below, we focus on the latter two scenarios, where the WIMP and ALP exhibit genuinely coupled dynamics.

First-order phase transition regime—At high temperatures, the ALP potential develops a symmetry-breaking minimum $\varphi_* \neq 0$ (see Fig. 1). We define $\tilde{x}_{\text{cfo}}$ as the time when this local minimum disappears:

$$\left. \frac{\partial \mathcal{V}(\varphi, \tilde{x}_{\text{cfo}})}{\partial \varphi} \right|_{\varphi=\varphi_*} = \left. \frac{\partial^2 \mathcal{V}(\varphi, \tilde{x}_{\text{cfo}})}{\partial \varphi^2} \right|_{\varphi=\varphi_*} = 0, \quad (8)$$

which determines both φ_* and $\tilde{x}_{\text{cfo}}$ as functions of κ . In the limit of $\kappa \ll 1$ and $x \gg 1$, the symmetry-breaking vacuum evolves as

$$\varphi_* = \sqrt{2 \left(1 - \frac{c_1}{x}\right)}, \quad (9)$$

where $c_1 \approx 1.33$ is the root of $c_1 K_0(c_1) = K_1(c_1)$ and $\tilde{x}_{\text{cfo}}$ follows the simple scaling

$$\tilde{x}_{\text{cfo}} = c_2 \kappa^{-1/3}, \quad c_2 \approx 0.83. \quad (10)$$

Here, we use the full thermal potential presented in Supplementary Material Sec. S4 [25] to obtain (10), providing a more accurate quantitative description. For

$x < \tilde{x}_{\text{cfo}}$, the tunneling rate is highly suppressed until very close to $\tilde{x}_{\text{cfo}}$, so the field remains trapped at φ_* . It returns to $\varphi = 0$ only when $x > \max\{\tilde{x}_{\text{cfo}}, x_{\text{osc}}\}$, where $x_{\text{osc}} \sim m_\chi/\sqrt{m_\phi M_{\text{Pl}}}$ marks the onset of ALP oscillations.

We now show how the above dynamics affect WIMP freeze-out. The Boltzmann equation governing the WIMP number density is [26]

$$\dot{n}_\chi + 3Hn_\chi = -\langle \sigma v \rangle_{\text{eff}}(n_\chi^2 - n_{\chi,\text{eq}}^2). \quad (11)$$

The ALP background reduces the WIMP effective mass, modifying both $\langle \sigma v \rangle_{\text{eff}}$ and the equilibrium distribution:

$$n_{\chi,\text{eq}} = \frac{g_\chi}{2\pi^2} m_{\chi,\text{eff}}^2 T K_2(m_{\chi,\text{eff}}/T). \quad (12)$$

When $\varphi = \varphi_*$, Eq. (9) implies

$$m_{\chi,\text{eff}}/T = (1 - \varphi_*^2/2)x = c_1 \approx 1.33. \quad (13)$$

So the argument of the Bessel function remains of the order of unity even at $x \gg 25$. This allows the WIMP to stay in equilibrium far longer than in the standard freeze-out scenario.

Freeze-out is initiated when symmetry is restored. In the parameter space of interest, $\tilde{x}_{\text{cfo}} > x_{\text{osc}}$, so after $\tilde{x}_{\text{cfo}}$, the ALP energy density redshifts as matter, $\varphi(x > \tilde{x}_{\text{cfo}}) = \varphi_*(\tilde{x}_{\text{cfo}}) \times (\tilde{x}_{\text{cfo}}/x)^{3/2}$. This increases the WIMP effective mass:

$$\frac{m_{\chi,\text{eff}}}{m_\chi} = 1 - \frac{\varphi_*^2(\tilde{x}_{\text{cfo}})}{2} \frac{\tilde{x}_{\text{cfo}}^3}{x^3} = 1 - \frac{\tilde{x}_{\text{cfo}}^3}{x^3} + c_1 \frac{\tilde{x}_{\text{cfo}}^2}{x^3}. \quad (14)$$

The rapid approach to $m_{\chi,\text{eff}} = m_\chi$ after $\tilde{x}_{\text{cfo}}$ drives χ out of equilibrium. This indicates the beginning of *coherent freeze-out*.

Using this expression in the Boltzmann equation and expanding $\sigma v = \sigma_0 + \sigma_1 v^2 + \dots$, the coherent freeze-out temperature x_{cfo} is found to be

$$x_{\text{cfo}} \approx \tilde{x}_{\text{cfo}} + \frac{1}{3} \log c + \frac{7}{6} \log \left(\frac{1}{3} \log c \right), \quad (15)$$

with

$$c \equiv \frac{27g_\chi}{2\pi^3} \sqrt{\frac{15}{2g_*}} m_\chi M_{\text{Pl}} (\sigma_0 + 6\sigma_1/\tilde{x}_{\text{cfo}}) \tilde{x}_{\text{cfo}}^{-4}. \quad (16)$$

The resulting asymptotic yield is

$$Y_{\chi,\infty}^{\text{cfo}} \approx \frac{14}{9} \sqrt{\frac{45}{\pi g_*}} \frac{1}{m_\chi M_{\text{Pl}}} \frac{x_{\text{cfo}}}{\sigma_0 + 21\sigma_1/(10x_{\text{cfo}})}. \quad (17)$$

The derivations of (15)–(17) are provided in Supplementary Material Sec. S3 [25]. These analytical

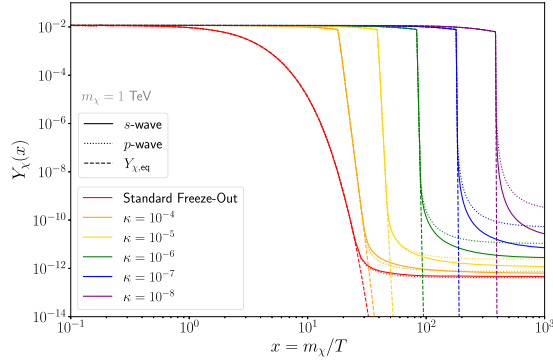


FIG. 2. WIMP yield $Y_\chi = n_\chi/s$ versus inverse temperature in the coherent freeze-out scenario. Solid (dotted) curves show numerical solutions of (11) for s -wave (p -wave) annihilation; dashed curves show the equilibrium yield (12). The standard freeze-out result is shown in red. Benchmark parameters used are $g_\chi = 2$, $m_\chi = 1$ TeV, $\sigma_0 = 2.2 \times 10^{-26}$ cm³/s, and $\sigma_1 = 1.9 \times 10^{-25}$ cm³/s.

expressions are highly accurate in the small- κ regime, where $x_{\text{cfo}} \approx \tilde{x}_{\text{cfo}}$; deviations from the exact numerical results remain below $\mathcal{O}(1)\%$ for $\tilde{x}_{\text{cfo}} \gtrsim 50$ (see Table I in Supplementary Material Sec. S3 [25]). The evolution of Y_χ with temperature is shown in Fig. 2.

At this point, it is interesting to compare the yield predicted in our mechanism with that in the standard freeze-out case:

$$\frac{Y_{\chi,\infty}^{\text{cfo}}}{Y_{\chi,\infty}^{\text{fo}}} = \frac{14}{9} \frac{x_{\text{cfo}}}{x_{\text{fo}}} \frac{\sigma_0^{\text{fo}} + 3\sigma_1^{\text{fo}}/x_{\text{fo}}}{\sigma_0^{\text{fo}} + 21\sigma_1^{\text{fo}}/(10x_{\text{cfo}})}, \quad (18)$$

where σ_i^{cfo} and σ_i^{fo} denote the annihilation cross sections in the coherent and standard freeze-out cases, respectively, and $x_{\text{fo}} \approx 25$ is the standard freeze-out temperature. For $\sigma_i^{\text{cfo}} = \sigma_i^{\text{fo}}$, we find an enhancement of the yield of the order of $x_{\text{cfo}}/x_{\text{fo}}$ for s -wave and $x_{\text{cfo}}^2/x_{\text{fo}}^2$ for p -wave annihilation. Since $x_{\text{cfo}} \sim \kappa^{-1/3} \gg 25$, the enhancement—especially for p -wave—is substantial.

If κ is extremely small, freeze-out may occur when χ is still relativistic. For instance, if WIMPs annihilate via exchanging weak gauge bosons, for temperatures below the weak scale, we have $\langle \sigma v \rangle_{\text{eff}} \sim T^2 G_F^2/x^\ell$, with $\ell = 0$ ($\ell = 1$) for s -wave (p -wave), and G_F is the Fermi constant. Requiring $n_\chi \langle \sigma v \rangle_{\text{eff}} \sim H$ with $n_\chi \sim T^3$ gives the relativistic freeze-out temperatures:

$$T_{\text{rfo}}^s \sim G_F^{-2/3} M_{\text{Pl}}^{-1/3} \sim \text{MeV}, \quad (19)$$

$$T_{\text{rfo}}^p \sim G_F^{-1/2} \left(\frac{m_\chi}{M_{\text{Pl}}} \right)^{1/4} \sim 10 \text{ MeV} \left(\frac{m_\chi}{10 \text{ GeV}} \right)^{1/4}. \quad (20)$$

Relativistic freeze-out yields too large χ relic abundance, so viable models require $x_{\text{cfo}} < m_\chi/T_{\text{rfo}}$.

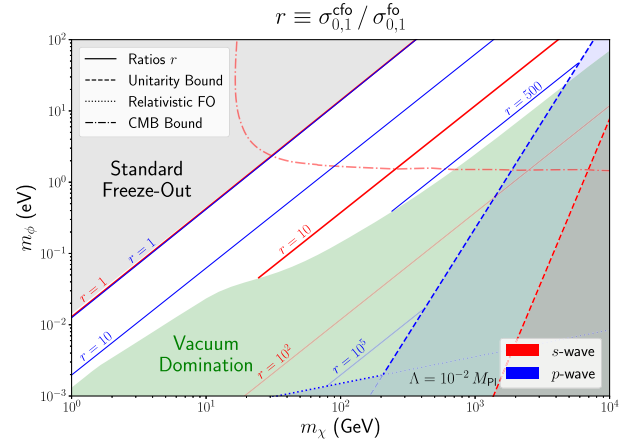


FIG. 3. Coherent freeze-out in the FOPT regime, where DM consists solely of the WIMP. Contours show the ratio of WIMP annihilation cross sections in the coherent and standard freeze-out scenarios for s -wave (red solid) and p -wave (blue solid) processes, assuming both reproduce the observed relic abundance. The green shaded region denotes the parameter space where the vacuum energy dominates before tunneling occurs. Regions below the dotted curves are excluded because of relativistic freeze-out. Regions to the right of the red (blue) dashed curves violate unitarity for s -wave (p -wave). The gray shaded region marks $r = 1$, where coherent and standard freeze-out coincide. The red dash-dotted line shows the CMB constraint on the s -wave cross section, excluding the region below it; no analogous CMB bound applies to p -wave due to velocity suppression.

The evolution of ϕ during freeze-out puts another upper bound on x_{cfo} . During the phase transition, the field is first trapped in the symmetry-breaking vacuum and then undergoes nonadiabatic evolution, tunneling under the potential barrier. For sufficiently small κ , vacuum energy can dominate before tunneling occurs, leading to a new inflationary phase. Avoiding vacuum domination places a conservative bound on the latest possible onset of coherent freeze-out, $x_{\text{cfo}} \lesssim \mathcal{O}(500)$ (see Supplementary Material Sec. S3 [25] for further discussion).

In the coherent freeze-out scenario, the WIMP annihilation cross section can exceed the standard value while still yielding the correct relic abundance. Figure 3 shows the ratio of the required cross sections in the two scenarios, assuming both reproduce the observed DM density, where we fix $g_\chi = 2$. The $r = 1$ contour corresponds to $\kappa \approx 10^{-4}$. For $\kappa \gtrsim 10^{-4}$, the ALP symmetry is restored before $x_{\text{fo}} \approx 25$, reducing to the standard freeze-out case (gray region), whereas for $\kappa \lesssim 10^{-4}$ the coherent effect becomes significant.

Cosmic microwave background (CMB) anisotropies place stringent limits on the WIMP s -wave annihilation cross section σ_0 [27,28]. For $m_\chi \lesssim 30$ GeV, the thermal value $\sigma_0^{\text{fo}} \approx 2.2 \times 10^{-26}$ cm³/s is excluded [29], with the bound weakening at higher masses. In the coherent

freeze-out scenario, the required s -wave cross section can exceed the standard value, extending the CMB constraint to heavier WIMPs. This bound appears as the red dash-dotted line in Fig. 3, and the region below it is excluded for s -wave-dominated annihilation. Even under this constraint, the s -wave cross section can still be enhanced by up to 2 orders of magnitude before reaching the unitarity limit [30] (red dashed line) and by a factor of about 30 before reaching the vacuum-domination limit (green shaded region).

In contrast, the enhancement needed for p -wave annihilation is much larger—scaling as $\propto x_{\text{cfo}}^2/x_{\text{fo}}^2$ rather than $\propto x_{\text{cfo}}/x_{\text{fo}}$ —and can reach values of the order of 10^5 in Fig. 3 before saturating unitarity and about 10^3 before reaching the vacuum-domination limit. Consequently, p -wave-dominated WIMPs (e.g., Majorana WIMPs lighter than the W boson [31]) can have annihilation cross sections up to 3 orders of magnitude above the standard thermal value under the conservative assumption that vacuum domination does not occur while still yielding the observed relic abundance. This greatly enlarges the viable parameter space and may improve the prospects for the indirect detection of p -wave DM.

Finally, we comment on the ALP evolution after WIMP freeze-out. For $x > \tilde{x}_{\text{cfo}}$, the field oscillates around $\phi = 0$, with the ALP energy density redshifting as matter. An explicit estimate gives the ALP density after freeze-out: $\rho_\phi(T) \sim m_\chi T^3$, which, if stable, would exceed the observed DM density and overclose the Universe. Thus, ϕ must be unstable and should also decay rapidly into radiation after the onset of coherent oscillations to avoid significant dilution of the WIMP yield. The decay of ϕ is generally model dependent and does not affect the parameter space shown in Fig. 3; for completeness, an explicit example realizing the rapid decay of ϕ is provided in Supplementary Material Sec. S3 [25]. We therefore conclude that in the FOPT regime, ϕ behaves as a spectator field, and DM is composed solely of the WIMP.

Crossover regime—We now turn to the crossover regime, where the ALP evolution is nearly adiabatic. In this regime, the symmetry of the ALP potential is restored before $x_{\text{fo}} \approx 25$, so the ALP has little influence on the WIMP freeze-out. Instead, the thermal effect from the WIMP bath significantly alters the ALP evolution, producing a relic abundance that departs from the standard misalignment result. Consequently, the ALP can naturally serve as a second DM component in this regime.

Since the two sectors decouple while χ is still relativistic, we may take the limit $\gamma\chi \ll 1$ in the ALP EOM (6), which reduces to

$$\phi'' + \frac{2}{x}\phi' - \eta(1 - \kappa x^2)\phi + \eta\phi^3 = 0. \quad (21)$$

The symmetry is broken for $x < x_c \equiv 1/\sqrt{\kappa}$ and restored for $x > x_c$. As shown by the numerical solution in Supplementary Material Sec. S2 [25], the ALP is initially

frozen by Hubble friction for $x \ll 1/\sqrt{\eta}$; for $1/\sqrt{\eta} \lesssim x \lesssim x_c$, it evolves toward the symmetry-breaking minimum at $\phi_* = 1$ and stabilizes there, erasing any dependence on its initial value; and for $x \gtrsim x_c$, the symmetry is restored and the field relaxes back toward $\phi = 0$.

The adiabatic approximation holds when

$$\eta|1 - \kappa x^2| \gg 1/x^2, \quad (22)$$

i.e., when the effective ALP mass dominates over Hubble friction. This condition fails near x_c , where the thermal and bare mass terms cancel. We define $x_1 = x_c + \delta x$ by requiring $|\eta(1 - \kappa x_1^2)| \equiv 1/x_1^2$ so that adiabaticity is restored for $x > x_1$. In the limit $\eta \gg \kappa$ (which is always satisfied for the relevant parameter space), one finds $\delta x/x_c \approx \kappa/(2\eta)$, implying that the loss of adiabaticity is extremely brief.

For $x > x_1$, adiabatic evolution gives

$$m_{\phi,\text{eff}}(x)\phi^2(x)a^3(x) = m_{\phi,\text{eff}}(x_1)\phi^2(x_1)a^3(x_1), \quad (23)$$

where $a(x)$ is the cosmic scale factor. The quantity in (23) is the adiabatic invariant, corresponding to the conservation of the comoving ALP number density.

At present, $x_0 \equiv m_\chi/T_0$, the thermal mass is negligible, $m_{\phi,\text{eff}}(x_0) = m_\phi$, so we obtain

$$\frac{\phi^2(x_0)a^3(x_0)}{\phi^2(x_1)a^3(x_1)} = \frac{H(x_1)}{m_\phi} \approx \frac{H(x_c)}{m_\phi} = \left(\frac{\kappa}{\eta}\right)^{1/2} \ll 1. \quad (24)$$

This corresponds to a rapid drop in the ALP field around x_c during symmetry restoration, which agrees well with the numerical solution in Supplementary Material Sec. S2 [25]. The present-day energy density is, therefore,

$$\rho_\phi(x_0) \approx m_\phi^2 \phi^2(x_c) \left(\frac{\kappa}{\eta}\right)^{1/2} \left(\frac{x_c}{x_0}\right)^3 \frac{g_{*S}(x_0)}{g_{*S}(x_c)}, \quad (25)$$

where we used $x_1 \approx x_c$ for $\eta \gg \kappa$ and g_{*S} is the entropy-weighted number of relativistic degrees of freedom. Substituting $\phi(x_c) = \sqrt{m_\chi/\Lambda}$ into (25) and using (5) and (7), we arrive at the relic abundance:

$$\begin{aligned} \Omega_\phi &= \frac{1.66 \cdot 4g_{*S}(x_0)\sqrt{2g_\chi g_*(x_c)} m_\chi^{3/2} \Lambda^{1/2} T_0^3}{3g_{*S}(x_c) H_0^2 M_{\text{Pl}}^3} \\ &\approx 0.3 \left(\frac{m_\chi}{10 \text{ GeV}}\right)^{3/2} \left(\frac{\Lambda}{0.1 M_{\text{Pl}}}\right)^{1/2}, \end{aligned} \quad (26)$$

where, to obtain the numerical value, we used $g_\chi = 4$, $g_{*S}(x_c) = g_*(x_c) = 106.75$, and $g_{*S}(x_0) = 3.93$.

We thus arrive at an “ALP miracle”: An ALP that couples quadratically to a weak-scale fermion with a Planck-suppressed interaction can naturally obtain the correct relic abundance. Importantly, this prediction is insensitive to both the initial ALP field value and the

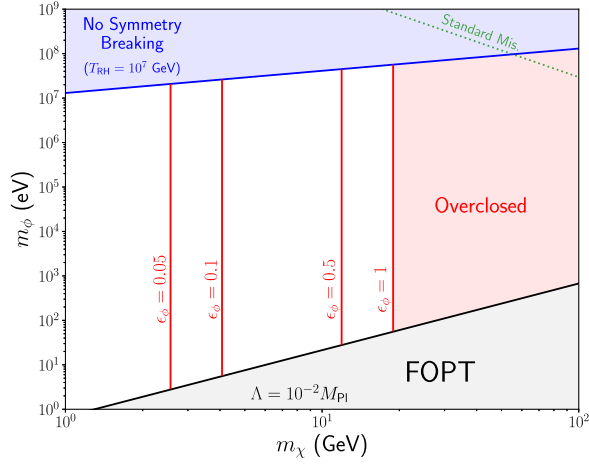


FIG. 4. Coherent freeze-out in the crossover regime, where DM may include both WIMPs and ALPs. Red contours show the ALP fraction $\epsilon_\phi \equiv \Omega_\phi/\Omega_{\text{DM}}$; regions with $\epsilon_\phi > 1$ overproduce DM. The area below the black line ($\kappa = \kappa_c$) corresponds to the FOPT regime, while regions above the blue line exhibit no symmetry breaking (relaxable with a higher reheating temperature). For comparison, the green dotted line shows the standard misalignment prediction of $\epsilon_\phi = 1$, using $\phi_i = \sqrt{m_\chi}\Lambda$ as the initial field displacement.

ALP mass, provided the latter lies in the crossover regime, $\text{eV} \lesssim m_\phi \lesssim \text{MeV}$ (see the right panel in Fig. 1).

In Fig. 4, we show contour lines of the ALP fraction $\epsilon_\phi \equiv \Omega_\phi/\Omega_{\text{DM}}$, where $\Omega_{\text{DM}} \simeq 0.26$ is the observed DM density [29]. For comparison, the dotted line shows the standard misalignment prediction of $\epsilon_\phi = 1$, obtained by taking the initial displacement $\phi_i = \sqrt{m_\chi}\Lambda$.

Conclusions—In this Letter, we considered the cosmological dynamics of a WIMP-like fermion and an ALP-like scalar field. We found that even a tiny effective coupling between the WIMP and ALP fields can have a dramatic effect on their evolution due to coherent forward scattering. At high temperatures, spontaneous symmetry breaking by the ALP field may be induced. The dynamics of the subsequent phase transition determine the present-day densities of both WIMPs and ALPs, which may deviate by many orders of magnitude from the predictions of the standard WIMP freeze-out and ALP misalignment production mechanisms. Many of the conventional experimental targets for WIMP and ALP searches need to be reconsidered in this scenario. For example, the indirect detection of a p -wave WIMP could become more promising. In addition, part of the parameter space of our model predicts a strongly first-order phase transition associated with the WIMP freeze-out, potentially sourcing an observable gravitational-wave signal. The mass suppression during the coherent freeze-out phase also implies efficient momentum exchange between WIMPs and the thermal bath, suggesting that kinetic decoupling occurs no earlier than the phase transition, although its precise timing and the

resulting impact on small-scale structure formation are model dependent. These possibilities will be explored in future work [32].

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Data availability—The data are not publicly available. The data are available from the authors upon reasonable request.

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